An Introductory Course on Causal Fermion Systems

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Comments are welcome!
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Preface

This book is intended as an easily accessible introduction to the theory of causal fermion systems. After giving the physical and mathematical background (Part 1), the theory of causal fermion systems is introduced (Part 2). We proceed by providing mathematical methods which can be regarded as a toolbox for analyzing causal fermion systems (Part 3). We conclude with an outlook on the applications (Part 4).

In order to address as large an audience as possible, the book contains extensive preliminaries which cover both physical and mathematical aspects. We have two typical audiences in mind when writing these preliminaries: physicists with only basic knowledge of mathematics and mathematicians without physical background.

The book is based on three main resources: First, the lecture notes of the spring school “Relativistic Fermion Systems” held in Regensburg in April 2013, adapted for the spring school “Causal Fermion Systems” held in Regensburg in March 2016. Second, the lecture “Causal Variational Principles” given at the University of Regensburg in the summer semester 2017. Finally, the online course “An Introduction to Causal Fermion Systems” held in the summer semester 2021.

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Felix Finster, Sebastian Kindermann and Jan-Hendrik Treude
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Part 1

Physical and Mathematical Background
CHAPTER 1

Physical Preliminaries

In this chapter we give the necessary physical background. Clearly, our presentation cannot replace the standard physics textbooks, which will be cited along the way. We also fix our conventions and introduce the notation which will be used consistently throughout this book.

1.1. The Schrödinger Equation

We begin by recalling a few basics on non-relativistic quantum mechanics. For a self-contained introduction we refer to good standard textbooks like [128, 130, 105].

A quantum mechanical particle is described by its wave function \( \psi(t, \vec{x}) \), where \( t \in \mathbb{R} \) is time and \( \vec{x} \in \mathbb{R}^3 \) is the position vector. We begin with nonrelativistic quantum mechanics without spin. In this case, the wave function \( \psi \) is complex-valued. Its absolute square \( |\psi(t, \vec{x})|^2 \) has the interpretation as the probability density of the particle to be located at the position \( \vec{x} \). For this interpretation to be sensible, the integral over the probability density must be equal to one,

\[
\int_{\mathbb{R}^3} |\psi(t, \vec{x})|^2 \, d^3x = 1.
\]  

This equation must hold for all times. This entails that the dynamical equations must preserve the integral (1.1.1).

A basic ingredient to quantum mechanics is the superposition principle. It states that for any wave functions \( \psi \) and \( \phi \), also their complex linear combination

\[
\tilde{\psi} = \alpha \psi + \beta \phi \quad \text{with } \alpha, \beta \in \mathbb{C}
\]  

(defined pointwise by \( \tilde{\psi}(t, \vec{x}) = \alpha \psi(t, \vec{x}) + \beta \phi(t, \vec{x}) \)) is a physically admissible wave function. In more mathematical terms, the physical wave functions form a complex vector space. Evaluating (1.1.1) for the wave function \( \tilde{\psi} \) and using that the probability integral must be preserved in time for all \( \alpha \) and \( \beta \), one concludes that the integral

\[
\int_{\mathbb{R}^3} \overline{\phi(t, \vec{x})} \psi(t, \vec{x}) \, d^3x
\]  

(1.1.3)

must be time independent for any wave functions \( \psi \) and \( \phi \). The procedure to deduce (1.1.3) from (1.1.1) is sometimes referred to as polarization. The integral (1.1.3) defines a scalar product on the wave functions, which we denote by

\[
\langle \phi | \psi \rangle_{\mathcal{H}} := \int_{\mathbb{R}^3} \overline{\phi(t, \vec{x})} \psi(t, \vec{x}) \, d^3x.
\]  

(1.1.4)

The complex vector space of wave functions endowed with this scalar product forms the Hilbert space \( L^2(\mathbb{R}^3, \mathbb{C}) \) of square-integrable functions, which we also denote by \( (\mathcal{H}, \langle \cdot | \cdot \rangle_{\mathcal{H}}) \) (for basics on Hilbert spaces see Section 2.2 below).
The dynamics of the wave function is described by a linear evolution equation on $\mathcal{H}$, the Schrödinger equation, which we write as

$$i\partial_t \psi = H \psi,$$

where $H$, the so-called Hamiltonian, is a linear operator acting on the Hilbert space $\mathcal{H}$. The Schrödinger equation is linear; this is essential in order to ensure that the time evolution is compatible with the superposition principle. The requirement that the scalar (1.1.4) must be time independent implies that

$$0 = \partial_t \langle \phi \mid \psi \rangle_{\mathcal{H}} = -i \left( \langle H \phi \mid \psi \rangle_{\mathcal{H}} - \langle \phi \mid H \psi \rangle_{\mathcal{H}} \right)$$

(1.1.6)

for all wave functions $\psi, \phi$. In other words, the Hamiltonian must be a symmetric operator on the Hilbert space $\mathcal{H}$ (mathematical issues like domains and the distinction between symmetric and selfadjoint operators are postponed to Section 3.2 below).

In the simplest setting without spin, the Hamiltonian is chosen as

$$H = -\frac{1}{2m} \Delta + V,$$

where we have set $\hbar = c = 1$ (we will do so throughout this book). Here, $\Delta$ is the Laplacian on $\mathbb{R}^3$, and $V(t, \vec{x})$ is a real-valued potential. The parameter $m > 0$ is the rest mass of the particle.

The Schrödinger equation can be analyzed with various methods. If the Hamiltonian is time independent, the Schrödinger equation can be solved by exponentiating (for details on the spectral theorem see Section 3.2 below),

$$\psi(t) = e^{-iHt} \psi(0).$$

Then the dynamics of $\psi$ can be related to spectral properties of the Hamiltonian. Another method, which has the advantage that it also applies if the Hamiltonian depends on time, is to make use of the fact that the time evolution forms a strongly continuous semigroup of operators (see for example [110], Section 34). Alternatively, one can analyze the Schrödinger equation as a parabolic partial differential equation. Since our focus are the relativistic equations, we do not cover these methods in this course, but refer the interested reader to the textbooks [132, Chapter 6] or [31, Section II.7.1].

### 1.2. Special Relativity and Minkowski Space

We now give a brief introduction to special relativity. For more details and the physical background we recommend the textbooks [124, 112, 106].

In special relativity, space and time are combined to a four-dimensional spacetime. Spacetime is described mathematically by Minkowski space $(\mathcal{M}, \langle \cdot, \cdot \rangle)$, a real four-dimensional vector space endowed with an inner product $\langle \cdot, \cdot \rangle$ of signature $(+ - - -)$. Choosing a pseudo-orthonormal basis $(e_i)_{i=0,\ldots,3}$ and representing the vectors of $\mathcal{M}$ in this basis, $\xi = \sum_{i=0}^{3} \xi^i e_i$, the inner product takes the form

$$\langle \xi, \eta \rangle = \sum_{j,k=0}^{3} g_{jk} \xi^j \eta^k,$$

(1.2.1)

where $g_{ij}$, the Minkowski metric, is the diagonal matrix $g = \text{diag}(1, -1, -1, -1)$. In what follows we usually omit the sums using Einstein’s summation convention (i.e. we sum over all indices which appear twice, once as an upper and once as a lower index). Also, we sometimes abbreviate the Minkowski inner product by writing $\xi \eta := \langle \xi, \eta \rangle$ and $\xi^2 := \langle \xi, \xi \rangle$. A pseudo-orthonormal basis $(e_i)_{i=0,\ldots,3}$ is also referred to as a reference
Special Relativity and Minkowski Space

Consider a frame, because the corresponding coordinate system \((x^i)\) of Minkowski space gives the time and space coordinates for an observer in a system of inertia. We also refer to \(t := x^0\) as time and denote the spatial coordinates by \(\vec{x} = (x^1, x^2, x^3)\).

The sign of the Minkowski metric encodes the causal structure of spacetime. Namely, a vector \(\xi \in \mathcal{M}\) is said to be

\[
\begin{align*}
\text{timelike} & \quad \text{if } \langle \xi, \xi \rangle > 0 \\
\text{spacelike} & \quad \text{if } \langle \xi, \xi \rangle < 0 \\
\text{lightlike} & \quad \text{if } \langle \xi, \xi \rangle = 0 .
\end{align*}
\]

Lightlike vectors are also referred to as null vectors. Moreover, the term non-spacelike refers to timelike or lightlike vectors. The null vectors form the double cone

\[
L := \{ \xi \in \mathcal{M} \mid \langle \xi, \xi \rangle = 0 \} ,
\]

referred to as the light cone. Physically, the light cone is formed of all light rays through the origin of \(\mathcal{M}\). Similarly, the timelike vectors correspond to velocities slower than the speed of light; they form the

\[
\text{interior light cone } \quad I := \{ \xi \in \mathcal{M} \mid \langle \xi, \xi \rangle > 0 \} .
\]

Likewise, the non-spacelike vectors form the

\[
\text{closed light cone } \quad J := \{ \xi \in \mathcal{M} \mid \langle \xi, \xi \rangle \geq 0 \} = I \cup L .
\]

We denote the future and past light cones by superscripts \(\lor\) and \(\land\), i.e.

\[
J^\lor := \{ \xi \in \mathcal{M} \mid \langle \xi, \xi \rangle \geq 0, \xi^0 \geq 0 \} \\
J^\land := \{ \xi \in \mathcal{M} \mid \langle \xi, \xi \rangle \geq 0, \xi^0 \leq 0 \} ,
\]

and similarly for \(I\). These notions are illustrated in Figure 1.1.

The spacetime trajectory of a moving object is described by a curve \(q(\tau)\) in Minkowski space (with \(\tau\) an arbitrary parameter). We say that the spacetime curve \(q\) is timelike if the tangent vector to \(q\) is everywhere timelike. Spacelike, null, and non-spacelike curves are defined analogously. The usual statement of causality that no information can travel faster than the speed of light can then be expressed as follows:

**Causality:** Information can be transmitted only along non-spacelike curves.

The set of all points which can be joined with a given spacetime point \(x\) by a non-spacelike curve is precisely the closed light cone centered at \(x\), denoted by \(J_x := J - x\). It is the
union of the two single cones
\begin{align*}
J'_x & = \{ y \in \mathcal{M} \mid (y-x)^2 \geq 0, \ (y^0 - x^0) \geq 0 \} \\
J'^{\omega}_x & = \{ y \in \mathcal{M} \mid (y-x)^2 \geq 0, \ (y^0 - x^0) \leq 0 \},
\end{align*}
interpreted as the points in the causal future and past of \( x \), respectively. Therefore, we refer to \( J'_x \) and \( J'^{\omega}_x \) as the closed future and past light cones centered at \( x \), respectively.

The sets \( I'_x \), \( I'^{\omega}_x \) and \( L'_x \), \( L'^{\omega}_x \) are introduced similarly.

Special relativity demands that physical equations be Lorentz invariant, which means that they must be formulated in Minkowski space, independent of the choice of reference frame. The simplest relativistic wave equation is the Klein-Gordon equation
\[ (-\Box - m^2) \psi = 0, \] (1.2.3)
where \( \Box \equiv \partial_j \partial^j \) is the wave operator. This equation describes a scalar particle (i.e. a particle without spin) of mass \( m \). If the particle has electric charge, one needs to suitably insert the electromagnetic potential \( A \) into the Klein-Gordon equation. One finds empirically that the equation
\[ - (\partial_k - ieA_k)(\partial^k - ieA^k) \psi = m^2 \psi \] (1.2.4)
describes a scalar particle of mass \( m \) and charge \( e \) in the presence of an electromagnetic field.

### 1.3. The Dirac Equation

We now give a brief introduction to the Dirac equation, the relativistic generalization of the Schrödinger equation. The physical background is introduced in more detail in the classic textbooks [14, 106, 118].

In order to describe a relativistic particle with spin, Dirac had the idea to work with a first order differential operator whose square is the wave operator. This operator is constructed by introducing the Dirac matrices \( \gamma^j \) as 4 \times 4-matrices which satisfy the anti-commutation relations
\[ 2 g^{jk} \mathbf{1} = \{ \gamma^j, \gamma^k \} \equiv \gamma^j \gamma^k + \gamma^k \gamma^j. \] (1.3.1)
Then the square of the operator \( \gamma^j \partial_j \) indeed gives the wave operator,
\[ (\gamma^j \partial_j)^2 = \gamma^j \gamma^k \partial_j \partial_k = \frac{1}{2} \{ \gamma^j, \gamma^k \} \partial_{jk} = \Box. \] (1.3.2)
For convenience, we shall always work in the Dirac representation, i.e. we set
\[ \gamma^0 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \gamma^i = \begin{pmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{pmatrix}, \] (1.3.3)
where \( \sigma^i \) are the three Pauli matrices
\[ \sigma^1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \sigma^2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \sigma^3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \] (1.3.4)
Including the mass \( m \geq 0 \), the Dirac equation in the vacuum (i.e. without any interaction) reads
\[ \left( i \gamma^k \frac{\partial}{\partial x^k} - m \right) \psi(x) = 0, \] (1.3.5)
where \( \psi(x) \), the Dirac spinor, has four complex components. If we multiply (1.3.5) by the operator \( (i \gamma^j \partial_j + m) \) and use (1.3.2), we find that each component of \( \psi \) satisfies the Klein-Gordon equation (1.2.3). Following the standard conventions in physics, we also
1.3. THE DIRAC EQUATION

denote the contraction with Dirac matrices by a slash, i.e. \( \mathsl{\gamma}^j u_j \) for \( u \) a vector of Minkowski space and \( \mathsl{\phi} = \mathsl{\gamma}^j \partial_j \). The action of the matrix \( \mathsl{\gamma} \) on a spinor \( \psi \), i.e. \( \mathsl{\gamma} \psi \), is often referred to as Clifford multiplication by the vector \( u \). The leptons and quarks in the standard model are Dirac particles. Thus, on the fundamental level, all matter is described by the Dirac equation.

In the presence of an electromagnetic field, the Dirac equation is modified to

\[
i \mathsl{\gamma}^k (\partial_k - iA_k) \psi = m \psi. \tag{1.3.6}
\]

Multiplying by the operator \( (i \mathsl{\gamma}^j (\partial_j - iA_j) + m) \) and using again the anti-commutation relations, we obtain the equation

\[
- (\partial_k - iA_k)(\partial^k - iA^k) + \frac{i}{2} F_{jk} \mathsl{\gamma}^j \mathsl{\gamma}^k - m^2 \psi = 0,
\]

where \( F_{jk} = \partial_j A_k - \partial_k A_j \) (see Exercise 1.1). This differs from the Klein-Gordon equation (1.2.4) by the extra term \( \frac{i}{2} F_{jk} \mathsl{\gamma}^j \mathsl{\gamma}^k \), which describes the coupling of the spin to the electromagnetic field.

The Dirac spinors at every spacetime point are endowed with an indefinite inner product of signature (2,2), which we call spin inner product and denote by

\[
\langle \psi | \phi \rangle(x) := \sum_{\alpha=1}^{4} s_{\alpha} \psi^{\alpha}(x)^{\dagger} \phi^{\alpha}(x), \quad s_1 = s_2 = 1, \ s_3 = s_4 = -1, \tag{1.3.7}
\]

where \( \psi^\dagger \) is the complex conjugate wave function (this inner product is often written as \( \overline{\psi} \phi \) with the so-called adjoint spinor \( \overline{\psi} = \psi^\dagger \mathsl{\gamma}^0 \)). By the adjoint \( A^* \) of an operator acting on spinors we always mean the adjoint with respect to the spin inner product. Thus it is defined by the relation

\[
\langle A^* \psi | \phi \rangle = \langle \psi | A \phi \rangle \quad \text{for all } \psi, \phi.
\]

In an obvious way, this definition of the adjoint gives rise to the notions selfadjoint, anti-selfadjoint and unitary. With these notions, the Dirac matrices are selfadjoint, meaning that

\[
\langle \gamma^l \psi | \phi \rangle = \langle \psi | \gamma^l \phi \rangle \quad \text{for all } \psi, \phi. \tag{1.3.8}
\]

To every solution \( \psi \) of the Dirac equation we can associate a vector field \( J \) by

\[
J^k = \langle \psi | \gamma^k \psi \rangle, \tag{1.3.9}
\]

referred to as the Dirac current. It is either timelike or lightlike (see Exercise 1.2). Moreover, it is divergence-free, as the following computation shows,

\[
\partial_k J^k = \partial_k \langle \psi | \gamma^k \psi \rangle = \langle \partial_k \psi | \gamma^k \psi \rangle + \langle \psi | \gamma^k \partial_k \psi \rangle
\]

\[
= i (\langle i \partial \psi | \psi \rangle - \langle \psi | i \partial \psi \rangle)
\]

\[
= i (\langle i (\partial + A - m) \psi | \psi \rangle - \langle \psi | (i \partial + A - m) \psi \rangle) = 0. \tag{1.3.10}
\]

This is referred to as current conservation.

Current conservation is closely related to the probabilistic interpretation of the Dirac wave function, as we now explain. Suppose that \( \psi \) is a smooth solution of the Dirac equation with suitable decay at spatial infinity (for example of spatially compact support;
Then current conservation allows us to apply the Gauß divergence theorem to obtain

\[ 0 = \int_{t_1}^{t_2} dt \int_{\mathbb{R}^3} d^3 x \partial_t \langle \gamma^k \psi \rangle(t, \vec{x}) \]

\[ = \int_{\mathbb{R}^3} \langle \psi \gamma^0 \psi \rangle(t_2, \vec{x}) d^3 x - \int_{\mathbb{R}^3} \langle \psi \gamma^0 \psi \rangle(t_1, \vec{x}) d^3 x \]

(1.3.11)

(this argument works similarly on a region \( \Omega \subset M \) whose boundary consists of two spacelike hypersurfaces). Polarizing (similar as explained after (1.1.2)), we conclude that for any two solutions \( \phi, \psi \) of the Dirac equation, the spatial integral

\[ (\phi | \psi) := 2\pi \int_{\mathbb{R}^3} \langle \phi \gamma^0 \psi \rangle(t, \vec{x}) d^3 x \]

(1.3.12)

is time independent. Since the inner product \( \langle . | \gamma^0 . \rangle \) is positive definite, the integral (1.3.12) defines a scalar product. We denote the Hilbert space corresponding to this scalar product by \( \mathcal{H} = L^2(\mathbb{R}^3)^4 \). In analogy to the integrand in (1.1.4) in non-relativistic quantum mechanics, the quantity \( \langle \psi | \gamma^0 \psi \rangle \) can be interpreted as the probability density. Current conservation (1.3.11) ensures that the probability integral is time independent.

The previous considerations generalize immediately to the situation in the presence of an external potential. To this end, we replace the operator \( \hat{A} \) in the Dirac equation (1.3.6) by a multiplication operator \( \mathcal{B}(x) \), which we assume to be smooth and symmetric with respect to the spin inner product, i.e.

\[ \langle \mathcal{B} \psi | \phi \rangle = \langle \psi | \mathcal{B} \phi \rangle \quad \text{for all } \psi, \phi. \]

(1.3.13)

We write the Dirac equation with a Dirac operator \( \mathcal{D} \) as

\[ (\mathcal{D} - m) \psi = 0 \quad \text{where} \quad \mathcal{D} := i\partial \phi + \mathcal{B}. \]

(1.3.14)

The symmetry assumption (1.3.13) is needed for current conservation to hold (as one sees immediately if in (1.3.10) one replaces \( \hat{A} \) by \( \mathcal{B} \)).

Similar to (1.1.5), we can again rewrite the dynamics with a symmetric operator \( H \). To this end, we multiply the Dirac equation (1.3.14) by \( \gamma^0 \) and bring the \( t \)-derivative on a separate side of the equation,

\[ i\partial_t \psi = H \psi \quad \text{where} \quad H := -\gamma^0(i\vec{\gamma}\nabla + \mathcal{B} - m) \]

(1.3.15)

(note that \( \gamma^j \partial_j = \gamma^0 \partial_0 + \vec{\gamma} \nabla \)). We refer to (1.3.15) as the Dirac equation in the Hamiltonian form. Now we can again apply (1.1.6) to conclude that the Hamiltonian is a symmetric operator on \( \mathcal{H} \).

We remark that in the Hamiltonian formulation, one often combines the prefactor \( \gamma^0 \) in (1.3.15) with the other Dirac matrices by working with the new matrices

\[ \beta := \gamma^0 \quad \text{and} \quad \vec{\alpha} := \gamma^0 \vec{\gamma}. \]

This is convenient because these new matrices are Hermitian with respect to the standard scalar product on \( \mathbb{C}^4 \). In this book, we shall not work with \( \alpha \) and \( \vec{\beta} \). We prefer the notation (1.3.15), because it becomes clearer which parts of the operators are Lorentz invariant. For calculations using \( \beta \) and \( \vec{\alpha} \) we refer for example to the monograph (134).
In addition to integrating over space (1.3.12), one can also introduce an inner product on spinorial wave functions by integrating the spin inner product over all of spacetime,

\[ <\psi|\phi> = \int_{\mathcal{M}} <\psi|\phi>_x d\mu_\mathcal{M} . \] (1.3.16)

This inner product will in general not be well-defined on solutions of the Dirac equation, because the time integral may diverge. But it can be considered for example on spinorial wave functions which are compactly supported in spacetime (but are no solutions of the Dirac equation). This \textit{spacetime inner product} will be important for the constructions in Chapter 15. In this context, it is very useful that the Dirac operator is symmetric with respect to the spacetime inner product, meaning that that

\[ <\mathcal{D}\psi|\phi> = <\psi|\mathcal{D}\phi> \] (1.3.17)

for all spinorial wave functions which decay sufficiently fast at spatial infinity and for large times. Indeed, the symmetry property (1.3.17) holds in curved spacetime as well (see the explanation after (4.2.31) below).

So far, Dirac spinors were introduced in a given reference frame. Let us verify that our definitions are independent of the choice of reference frame. A linear transformation of Minkowski space which leaves the form of the Minkowski metric (1.2.1) invariant is called a \textit{Lorentz transformation}. All Lorentz transformations form a group, the so-called \textit{Lorentz group}. The Lorentz transformations which preserve both the time direction and the spatial orientation form a subgroup of the Lorentz group, the \textit{orthochronous proper Lorentz group}. We consider two reference frames \((x^j)\) and \(\tilde{(x}^l)\) with the same orientation of time and space. Then the reference frames are related to each other by an orthochronous proper Lorentz transformation \(\Lambda\), i.e. in components

\[ \tilde{x}^l = \Lambda_j^l x^j , \quad \frac{\partial}{\partial \tilde{x}^l} = \frac{\partial \tilde{x}^l}{\partial x^j} \frac{\partial}{\partial x^j} = \Lambda_j^l \frac{\partial}{\partial x^j} , \]

and \(\Lambda\) leaves the Minkowski metric invariant,

\[ \Lambda_j^l \Lambda_k^m g_{lm} = g_{jk} . \] (1.3.18)

Under this change of spacetime coordinates, the Dirac operator \(i\gamma^l \left( \frac{\partial}{\partial \tilde{x}^l} - iA_l \right)\) transforms to

\[ i\tilde{\gamma}^l \left( \frac{\partial}{\partial \tilde{x}^l} - i\tilde{A}_l \right) \quad \text{with} \quad \tilde{\gamma}^l = \Lambda_j^l \gamma^j . \] (1.3.19)

This transformed Dirac operator does not coincide with the Dirac operator \(i\gamma^l \left( \frac{\partial}{\partial x^l} - iA_l \right)\) as defined in the reference frame \((\tilde{x}^l)\) because the new Dirac matrices have a different form. However, the next lemma shows that the two Dirac operators do coincide after a suitable unitary transformation of the spinors.

\textbf{Lemma 1.3.1.} \textit{For any orthochronous proper Lorentz transformation} \(\Lambda\) \textit{there is a unitary matrix} \(U(\Lambda)\) \textit{unitary with respect to the spin inner product (1.3.7)} \textit{such that}

\[ U(\Lambda) \Lambda_j^l \gamma^j U(\Lambda)^{-1} = \gamma^l . \]

\textbf{Proof.} Since \(\Lambda\) is orthochronous and proper, we can write it in the form \(\Lambda = \exp(\lambda)\), where \(\lambda\) is a suitable generator of a rotation and/or a Lorentz boost. Then \(\Lambda(t) := \exp(t\lambda), \ t \in \mathbb{R},\) is a family of Lorentz transformations, and differentiating (1.3.18) with respect to \(t\) at \(t = 0\), we find that

\[ \Lambda_j^l g_{lk} = -g_{jm} \Lambda_k^m \]
(note that $\Lambda(t)^l_j = \delta^l_j + t \lambda^l_j + \cdots$). Using this identity together with the fact that the Dirac matrices are selfadjoint, it is straightforward to verify that the matrix

$$u := \frac{1}{4} \lambda_k \gamma^l \gamma^k$$

is anti-selfadjoint. As a consequence, the family of matrices

$$U(t) := \exp (tu)$$

is unitary. We now consider for a fixed index $l$ the family of matrices

$$A(t) := U(t) \Lambda(t)^l_j \gamma^j U(t)^{-1}.$$

Clearly, $A(0) = \gamma^l$. Furthermore, differentiating with respect to $t$ gives

$$\frac{d}{dt} A(t) = U \Lambda^l_j \left\{ u \gamma^j - \gamma^j u + \lambda^l_k \gamma^k \right\} U^{-1},$$

and a short calculation using the commutation relations

$$[\gamma^l \gamma^k, \gamma^j] = 2 (\gamma^l \gamma^{kj} - \delta^l_j \gamma^k)$$

shows that the curly brackets vanish. We conclude that $A(1) = A(0)$, proving the lemma.

Applying this lemma to the Dirac operator in (1.3.19), one sees that the Dirac operator is invariant under the joint transformation of the spacetime coordinates and the spinors

$$x^j \rightarrow \Lambda^l_k x^k, \quad \psi \rightarrow U(\Lambda) \psi.$$

Moreover, since the matrix $U(\Lambda)$ is unitary, the representation of the spin inner product (1.3.7) is valid in any reference frame. We conclude that our definition of spinors is indeed Lorentz invariant.

For what follows, it is important to keep in mind that, in contrast to the spin inner product, the combination $\psi^\dagger \phi = \langle \psi | \gamma^0 \phi \rangle$ is not Lorentz invariant. Instead, it is the zero component of a Minkowski vector. Consequently, the integrand in (1.3.12) is not a scalar, but a density. Its integral, on the other hand, is again Lorentz invariant due to current conservation.

Out of the Dirac matrices one can form the pseudo-scalar matrix $\Gamma$ by

$$\Gamma = \frac{i}{4!} \epsilon_{jklm} \gamma^j \gamma^k \gamma^l \gamma^m = i \gamma^0 \gamma^1 \gamma^2 \gamma^3$$

(in the physics literature, this matrix is usually denoted by $\gamma^5$). Here $\epsilon_{jklm}$ is the totally antisymmetric symbol (i.e. $\epsilon_{jklm}$ is equal to $\pm 1$ if $(j,k,l,m)$ is an even and odd permutation of $(0,1,2,3)$, respectively, and vanishes otherwise). A short calculation shows that the pseudo-scalar matrix is anti-selfadjoint and $\Gamma^2 = 1$. As a consequence, the matrices

$$\chi_L = \frac{1}{2} (1 - \Gamma), \quad \chi_R = \frac{1}{2} (1 + \Gamma)$$

satisfy the relations

$$\chi^2_{L/R} = \chi_{L/R}, \quad \Gamma \chi_L = -\chi_L, \quad \Gamma \chi_R = \chi_R, \quad \chi^*_L = \chi_R, \quad \chi_L + \chi_R = 1.$$

They can be regarded as the spectral projectors of the matrix $\Gamma$ and are called the chiral projectors. The projections $\chi_L \psi$ and $\chi_R \psi$ are referred to as the left- and right-handed components of the spinor, respectively. A matrix is said to be even and odd if it
commutes or anti-commutes with $\Gamma$, respectively. It is straightforward to verify that the Dirac matrices are odd, and therefore
\[ \gamma^j x_{L|R} = x_{R|L} \gamma^j. \]

By multiplying the Dirac equation (1.3.6) from the left by $x_{L|R}$, one can rewrite it as a system of equations for the left- and right-handed components of $\psi$,
\[ i\gamma^k (\partial_k - iA_k) x_L \psi = m x_R \psi, \quad i\gamma^k (\partial_k - iA_k) x_R \psi = m x_L \psi. \]

If $m = 0$, these two equations decouple, and we get separate equations for the left- and right-handed components of $\psi$. This observation is the starting point of the 2-component Weyl spinor formalism. Here we shall not use this formalism. Instead, we will describe chiral massless particles (like massless neutrinos) by the left- or right-handed component of a Dirac spinor.

1.4. The Hilbert Space of Dirac Solutions

We now formulate the structures of Dirac theory with a convenient notation, which has the main advantage of generalizing to curved spacetime. The Dirac wave functions are four-component complex wave functions in Minkowski space. More generally, one can consider them as sections of a vector bundle. In view of these general concepts (which will be introduced in Section 2.5 below), we denote the Cartesian product
\[ S_M := M \times \mathbb{C}^4, \]

as the spinor bundle. For every $x \in M$, the corresponding subset $S_x := \{x\} \times \mathbb{C}^4$ is referred to as the spinor space at the spacetime point $x$. Clearly, the spinor bundle is the disjoint union of all the spinor spaces,
\[ S_M = \bigcup_{x \in M} S_x. \]

(Endowed with a topology to be introduced in Section 2.5). The wave functions take values in the corresponding spinor spaces, i.e.
\[ \psi(x) \in S_x \sim \mathbb{C}^4. \]

A mapping $\psi : M \to S_M$, which to every $x \in M$ associates a point in the corresponding spinor space $S_x$, is referred to as a section of the spinor bundle. The spinor space $S_x$ is the fiber of the spinor bundle at the spacetime point $x$. The spin inner product (1.3.7) is an inner product on the fiber, as we often clarify by an additional subscript $x$,
\[ \langle \cdot, \cdot \rangle_x : S_x \times S_x \to \mathbb{C}. \]

We next highlight the obtained analytic structures, again anticipating concepts and results to be introduced later in this book. The scalar product (1.3.12) on the Dirac solutions gives rise to a Hilbert space structure (for a mathematical introduction to Hilbert spaces see Section 2.2 below). In order to construct the right class of solutions, one can begin by solving the Cauchy problem for smooth initial data $\psi_0$ of compact support given for example on the hypersurface $\{t = 0\}$. Rewriting the Dirac equation as a linear symmetric hyperbolic system (see Chapter 13 below), one sees that this Cauchy problem has a unique global solution in Minkowski space. Moreover, this solution is smooth and, due to finite propagation speed, has compact support on any other hypersurface $\{t = \text{const}\}$ (see Figure 1.2). One says that the solution has spatially compact support. The smooth and spatially compact sections of the spinor bundle (not necessarily solutions of
the Dirac equation) are denoted by $C_c^\infty (\mathcal{M}, S\mathcal{M})$. Clearly, for spatially compact solutions, the scalar product (1.3.12) is well-defined and finite. Taking the completion, one obtains a Hilbert space denoted by $(\mathcal{H}_m, \langle ., . \rangle)$ (for details on the completion and Hilbert spaces see Section 2.2 and Exercise 2.2). By construction, we know that

$$C_c^\infty (\mathcal{M}, S\mathcal{M}) \cap \mathcal{H}_m \quad \text{is dense in } \mathcal{H}_m.$$

For the reader familiar with the theory of partial differential equations, we finally remark that the solutions in $\mathcal{H}_m$ can also be characterized in terms of Sobolev spaces (see for example [31], Section II.5). More precisely, the solutions in $\mathcal{H}_m$ are in $H^{1,2}_\text{loc}(\mathcal{M}, S\mathcal{M})$. By the trace theorem (see for example [31], Section II.5.2), their restriction to a hypersurface $\{ t = \text{const} \}$ is in $L^2_\text{loc}(\mathbb{R}^3, \mathbb{C}^4)$. As a consequence, the integrand of the spatial integral in (1.3.12) is locally integrable. The solutions in $\mathcal{H}_m$ have the additional property that their restriction to the hypersurfaces are even in $L^2(\mathbb{R}^3, \mathbb{C}^4)$, so that the integral in (1.3.12) exists and is finite.

### 1.5. Dirac’s Hole Theory and the Dirac Sea

The Dirac theory gives rise to anti-matter and pair creation, as we now briefly explain. For the present purposes, it suffices to consider the Dirac equation in the vacuum (1.3.5). It can be solved by the plane-wave ansatz (for more details see for example [14], Section 3.1, [118], Section 3.3 or [134], Section 1.4.1)

$$\psi(x) = \chi(k) e^{-ikx},$$

where $k \in \mathcal{M}$ is the four-momentum, and $kx = \langle k, x \rangle$ is the Minkowski inner product (for a mathematically precise treatment of the Fourier transform see Section 24.1 below).

Using this ansatz in (1.3.5) gives the matrix equation

$$\begin{pmatrix} \hat{k} & m \end{pmatrix} \chi(k) = 0. \quad (1.5.1)$$

Multiplying by $\hat{k} + m$ and using the anti-commutation relations (1.3.1) gives the necessary condition

$$k^2 = m^2. \quad (1.5.2)$$

If this condition is satisfied, the matrix $\hat{k} - m$ has a two-dimensional kernel, which coincides with the image of the matrix $\hat{k} + m$. Thus the general solution of (1.5.1) can be written as

$$\chi(k) = (\hat{k} + m) \phi \quad \text{with} \quad \phi \in \mathbb{C}^4.$$

The zero component $\omega := k^0$ of four-momentum has the interpretation as $2\pi$ times the frequency of the wave. Equation (1.5.2), also referred to as the *dispersion relation,*
can be written as
\[ \omega^2 = |\vec{k}|^2 + m^2 \quad \text{or} \quad \omega = \pm \sqrt{|\vec{k}|^2 + m^2}. \]

Here the plus and the minus sign correspond to positive and negative frequency, respectively. The corresponding solutions are said to be on the upper and lower mass shell (see Figure 1.3). Using Planck’s relation \( E = \hbar \omega \), the frequency can be related to the energy of the solution. Thus the solutions on the upper and lower mass shell have positive respectively negative energy.

At first sight, the occurrence of solutions of negative energy is problematic, because particles of negative energy have never been observed. Moreover, at least in a naive consideration, solutions of negative energy should make the physical system unstable, because by bringing a particle into a state of larger and larger negative energy, using energy conservation one could extract more and more positive energy from the system.

This problem was resolved by Dirac in 1930 \[28\] and led to the prediction of particle creation and anti-matter, as we now outline (an excellent more detailed explanation can be found in [14, Section 5.1]). We work in the setting of non-interacting many-particle quantum mechanics, where the many-particle wave function is described by a product of one-particle wave functions. In other words, the quantum state is described by occupying many one-particle states. Dirac’s concept is that in the vacuum, all the states of negative energy should be occupied, forming the so-called Dirac sea. According to the Pauli exclusion principle, each state may be occupied by at most one electron. Therefore, adding particles to the system, the additional particles must occupy states of positive energy, giving rise to electrons. By convention, the electrons have negative electric charge.

Moreover, one can create “holes” in the Dirac sea. The resulting “hole in a sea of negative energy” appears as a particle of again positive energy, but with the opposite and thus positive electric charge. These “holes” can be observed as positrons. Furthermore, starting from the completely filled Dirac sea, one can “excite” a particle of the sea by a transition from a state of negative energy to a state of positive energy. As a result, one gets a particle (=electron) plus a hole (=positron). This explains why matter and anti-matter can be created in pairs in a process called pair creation.

The above intuitive picture of the Dirac sea has important observable consequences, because it explains fundamental physical phenomena like anti-matter and pair creation. Nevertheless, the naive picture suffers from the problems that the Dirac sea has an infinite negative charge density and an infinite energy density. In modern quantum field theory, these problems are bypassed by introducing a suitable vacuum state and working “relative” to this vacuum state. Here we shall not enter these constructions. Instead, we
shall take Dirac’s concept of a “sea of interacting particles” seriously, as will be explained in more detail in Section 5.8.

Exercises

Exercise 1.1. Show that, multiplying the Dirac equation (1.3.6) by the operator $(iγ^j(∂_j−iA_j)+m)$ and using the anti-commutation relations, we obtain the equation

$$−(∂_k−iA_k)(∂^k−iA^k)+\frac{i}{2}F_{jk}\gamma^j\gamma^k−m^2)\psi=0.$$ 

This differs from the Klein-Gordon equation (1.2.4) by the extra term $\frac{i}{2}F_{jk}\gamma^j\gamma^k$, which describes the coupling of the spin to the electromagnetic field.

Exercise 1.2. In this exercise, we shall verify that for any non-zero spinor $ψ$, the corresponding Dirac current vector $J^k=⟨ψ|γ^kψ⟩$ is non-spacelike.

(a) Show that the matrix $γ^0γ^1$ is Hermitian and has eigenvalues $±1$. Deduce that

$$⟨ψ,γ^0γ^1ψ⟩_C^4≤∥ψ∥^2_{C^4}.$$ 

(b) Show that the last inequality implies that $|J^1|≤|J^0|$.

(c) Use the rotational symmetry of the Dirac equation to conclude that $J^0≥|J|$ (where $J=(J^1,J^2,J^3)∈\mathbb{R}^3$).

Exercise 1.3. (frequency splitting) Let us define for every $k∈\mathbb{R}^3$ the energy $ω(k):=\sqrt{k^2+m^2}$ and the matrices

$$p_±(k):=\frac{k+m}{2k^0}\gamma^0|k^0=±ω(k)∈\text{Mat}(4,\mathbb{C})$$

with $k:=k^0γ^0−k\cdotγ$.

(a) Referring to the standard scalar product of $\mathbb{C}^4$, show that the matrices $p_±(k)$ are symmetric, idempotent, add up to the identity and have orthogonal images. Conclude that the spinor space $\mathbb{C}^4$ can be decomposed into the orthogonal direct sum

$$\mathbb{C}^4=W_k^+⊕W_k^−, \quad \text{with} \quad W_k^±:=\text{Im}p_±(k).$$

(b) Let $φ∈C_∞^∞(\mathbb{R}^3,\mathbb{C}^4)$ be a smooth solution of the Dirac equation with spatially compact support, i.e.

$$(iγ^0∂_0+iγ^α∂_α−m)φ=0, \quad \text{with} \quad φ(t,⋅)∈C_0^∞(\mathbb{R}^3,\mathbb{C}^4) \text{ for all } t∈\mathbb{R}.$$ 

Let $φ$ be the smooth function on $\mathbb{R}^4$ defined by taking the Fourier transform of $φ$ in the spatial variables only. Find $h∈C_∞(\mathbb{R}^3,\text{Mat}(4,\mathbb{C}))$ such that

$$i∂tφ(t,k)=h(k)⋅φ(t,k) \quad \text{for all } t∈\mathbb{R}, \quad k∈\mathbb{R}^3.$$ 

Show that $h(k)$ is also symmetric with respect to the standard scalar product of $\mathbb{C}^4$ and satisfies

$$h(k)p_±(k)=±ω(k)p_±(k).$$

In particular, $±ω(k)$ form the spectrum of $h(k)$.

(c) Referring to point (b), conclude that

$$φ(t,x)=\int_{\mathbb{R}^3}\frac{d^3k}{(2π)^{3/2}}(p_−(k)φ(0,k)e^{iω(k)t}+p_+(k)φ(0,k)e^{−iω(k)t})e^{ik⋅x}$$

*Hint:* Remember that the Cauchy problem admits unique smooth solutions.
From a mathematical point of view, the Dirac sea is described by the Hilbert space generated by all the smooth solutions with spatially compact support and the property that \( p_-(k) \hat{\psi}(0, k) = \hat{\psi}(0, k) \).

**Exercise 1.4. (Left- and right-handed spinors)** Out of the Dirac matrices \( \gamma^i \) one can form the **pseudo-scalar matrix** \( \gamma^5 \) by

\[
\gamma^5 := \frac{i}{4!} \epsilon_{jklm} \gamma^j \gamma^k \gamma^l \gamma^m = i \gamma^0 \gamma^1 \gamma^2 \gamma^3,
\]

where \( \epsilon_{jklm} \) is the totally antisymmetric symbol. Using the commutation relations of the Dirac matrices, shows that \((\gamma^5)^\dagger = \gamma^5 = - (\gamma^5)^*\), \((\gamma^5)^2 = 1\) and \{\(\gamma^5, \gamma^i\)\} := \(\gamma^5 \gamma^i + \gamma^i \gamma^5 = 0\). As a corollary, show that the matrices

\[
\chi_L := \frac{1}{2}(1 + \gamma^5), \quad \chi_R := \frac{1}{2}(1 - \gamma^5),
\]

called the **chiral projectors**, satisfy the following identities

\[
(\chi_{L/R})^2 = \chi_{L/R}, \quad \gamma^5 \chi_L = - \chi_L, \quad \gamma^5 \chi_R = \chi_R, \quad (\chi_{L/R})^* = \chi_{L/R},
\]

\[(\chi_{L/R})^\dagger = \chi_{L/R}, \quad \chi_L + \chi_R = 1.
\]

Given a spinor \( \psi \in \mathbb{C}^4 \), the projections \( \chi_L \psi \) and \( \chi_R \psi \) are referred to as the **left- and right-handed components** of the spinor, respectively. Show that

\[
\gamma^j \chi_{L/R} = \chi_{R/L} \gamma^j \quad \text{for all } j \in \{0, 1, 2, 3\}.
\]

Using these relations, show that the Dirac equation with an external electromagnetic field \( A_j \) can be rewritten as a system of equations for the left- and right-handed components of \( \psi \),

\[
i \gamma^k (\partial_k - i A_k) \chi_L \psi = m \chi_R \psi, \quad i \gamma^k (\partial_k - i A_k) \chi_R \psi = m \chi_L \psi
\]

What happens in the limiting case \( m = 0 \)?

**Exercise 1.5. (Invariance of the scalar product under Lorentz transformations)** Let \((\mathbb{R}^4, \eta)\) be Minkowski Space, where \( \eta = \text{diag}(-1, 1, 1, 1) \). The matrix Lie group

\[
\mathcal{L}_+ := \{ \Lambda \in \text{Mat}(4, \mathbb{R}) \mid \eta_{ij} \Lambda^i_l \Lambda^j_m = \eta_{lm}, \ \Lambda^0_0 > 0, \ \det \Lambda = +1 \}
\]

is known as the **proper orthochronous Lorentz group**. A **Cauchy hyperplane** is a 3-dimensional hyperplane \( \mathcal{N} \subset \mathbb{R}^4 \) with the property that its tangent vectors are spacelike. A vector field \( \nu \in T\mathbb{R}^4 \) is said to be **normal to \( \mathcal{N} \)** and **future directed** if, for all \( p \in \mathcal{N} \),

\[
\eta_{ij} \nu^i \nu^j |_p = 1, \ \nu(p) > 0 \quad \text{and} \quad \eta_{ij} \nu^i u^j |_p = 0 \quad \text{for all } u \in T\mathcal{N}.
\]

On the manifold \( \mathcal{N} \) we define a **volume form**

\[
d \text{vol}_\mathcal{N} := \nu_\iota(d \text{vol}_{\mathbb{R}^4}) = \nu^0 \epsilon_{i_0 i_1 i_2 i_3} dx^{i_1} \otimes dx^{i_2} \otimes dx^{i_3},
\]

where \( \epsilon_{ijkl} \) is the totally antisymmetric symbol and \( d \text{vol}_{\mathbb{R}^4} = \epsilon_{i_0 i_1 i_2 i_3} dx^{i_0} \otimes dx^{i_1} \otimes dx^{i_2} \otimes dx^{i_3} \) is the canonical volume form on \( \mathbb{R}^4 \).

Finally, consider the set of smooth spatially compact solutions of the Dirac equation

\[
\mathcal{D} := \{ \psi \in C^\infty_{sc}(\mathbb{R}^4, \mathbb{C}^4) \mid (i \gamma^j \partial_j - m) \psi(x) = 0 \}.
\]

Prove the following statements.

(a) The normal vector field \( \nu \) is time-like, i.e. \( \eta_{ij} \nu^i \nu^j > 0 \).

(b) **Current conservation**: for all \( \psi \in \mathcal{D} \),

\[
\partial_k J^k = 0, \quad \text{with} \quad J^k(x) := \langle \psi(x), \gamma^k \psi(x) \rangle.
\]
(c) The following integral is independent of $\mathcal{N}$, with $\nu_i := \eta_{ij} \nu^j$,
\[
\langle \psi, \phi \rangle := \int_{\mathcal{N}} \langle \psi(x), \gamma^i \phi(x) \rangle \nu_i(x) \, d\text{vol}_N(x), \quad \psi, \phi \in \mathcal{D}. \tag{1.5.4}
\]

Hint: Can you apply the Gauss divergence theorem?

(d) The volume form $d\text{vol}_N$ is invariant under Lorentz transformations, i.e.
\[
f^*_\Lambda (d\text{vol}_N) = d\text{vol}_N, \quad \hat{\mathcal{N}} = f_\Lambda(\mathcal{N}),
\]
where $f_\Lambda : \mathbb{R}^4 \ni x \mapsto \Lambda x \in \mathbb{R}^4$ and $\phi^*$ denotes the pull-back.

Hint: What is the relation between $\epsilon_{jklm}$ and the determinant of a $4 \times 4$ Matrix?

(e) The following linear mapping is well-defined,
\[
\mathcal{D} \ni \psi \mapsto \hat{\psi} := U(\Lambda^{-1}) \psi(\Lambda \cdot ) \in \mathcal{D}, \tag{1.5.5}
\]
where $U : \mathcal{L}^+_4 \ni \Lambda \mapsto U(\Lambda) \in \text{Mat}(4, \mathbb{C})$ satisfies
\[
U(\Lambda^{-1}) = U(\Lambda)^* = U(\Lambda)^{-1}, \quad U(\Lambda)^* \gamma^i U(\Lambda) = \Lambda^i_j \gamma^j.
\]
Moreover, the product (1.5.4) is invariant under (1.5.5).
CHAPTER 2

Mathematical Preliminaries

2.1. Basics on Topology

We here recall a few basic concepts from topology. A more systematic treatment can be found in many good elementary textbooks like for example [129] or [101]. Since topological spaces are a rather abstract concepts, we prefer to begin with metric spaces.

**Definition 2.1.1.** Let $E$ be a set. A mapping

$$d : E \times E \to \mathbb{R}_0^+$$

is a **metric** if it has the following properties:

(i) **Positivity:** For all $x, y \in E$,

$$d(x, y) \geq 0 \quad \text{and} \quad d(x, y) = 0 \iff x = y.$$

(ii) **Symmetry:** For all $x, y \in E$,

$$d(x, y) = d(y, x).$$

(iii) **Triangle inequality:** For all $x, y, z \in E$,

$$d(x, y) \leq d(x, z) + d(y, z).$$

The pair $(E, d)$ is a **metric space**.

A simple example is $E = \mathbb{R}^3$ with the Euclidean distance function

$$d(x, y) = \|x - y\| := \left(\sum_{\alpha=1}^{3} |x^\alpha - y^\alpha|^2 \right)^{\frac{1}{2}}.$$ 

In view of this example, the metric $d(x, y)$ is sometimes also referred to as the **distance** between $x$ and $y$.

A metric gives rise to a corresponding topology:

**Definition 2.1.2.** For any $x \in E$ and $r > 0$, the set

$$B_r(x) := \{ y \in E \mid d(x, y) < r \}$$

is referred to as the **open ball** of radius $r$ centered at $x$. A subset $\Omega \subset E$ is **open** if for every $x \in \Omega$ there is a radius $r > 0$ such that $B_r(x) \subset \Omega$ (see Figure 2.1). The **topology** $\mathcal{O}$ of $E$ is the family of all open subsets,

$$\mathcal{O} := \{ \Omega \subset E \text{ open} \} \subset \mathcal{P}(E)$$

(where $\mathcal{P}(E)$ is the power set, i.e. the set of all subsets of $E$).

We next collect the properties of the topology (we omit the proof, which can be found in many textbooks).

**Lemma 2.1.3.** Given a metric space $(E, d)$, the corresponding topology $\mathcal{O}$ has the following properties:
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Figure 2.1. An open set $\Omega \subset E$.

(i) $\emptyset, E \in \mathcal{O}$

(ii) Closedness under finite intersections, i.e.
$$\Omega_1, \ldots, \Omega_n \in \mathcal{O} \implies \Omega_1 \cap \cdots \cap \Omega_n \in \mathcal{O}.$$ 

(iii) Closedness under arbitrary unions: For any family $(\Omega_\lambda)_{\lambda \in \Lambda}$ of subsets of $E$ (with an arbitrary, possibly infinite index set $\Lambda$), the following implication holds,
$$\Omega_\lambda \in \mathcal{O} \quad \forall \lambda \in \Lambda \implies \bigcup_{\lambda \in \Lambda} \Omega_\lambda \in \mathcal{O}.$$

A topological space is defined by taking the topology its just-stated properties as the starting point:

**Definition 2.1.4.** A set $E$ together with a distinguished family of subsets $\mathcal{O} \subset \mathcal{P}(E)$ having the properties (i)-(iii) in Lemma 2.1.3 is referred to as a topological space. The family of subsets $\mathcal{O}$ is the topology of $E$. The sets in $\mathcal{O}$ are called open.

Clearly, topological spaces are a general and abstract concept. In particular, the topology of a topological space does not necessarily need to come from an underlying metric. In most of our applications, the topology has additional properties. In particular, we shall always assume that the topology is Hausdorff, meaning that for all distinct point $x, y \in E$, there are disjoint open set $U, V \in \mathcal{O}$ with $x \in U$ and $y \in V$ (see Figure 2.2). Note that the topology coming from a metric is always Hausdorff, as one sees immediately by choosing $U = B_r(x)$ and $V = B_r(y)$ with $r = d(x, y)/3$.

The significance of the above definitions lies in the fact that many notions from analysis can be formulated purely in topological terms. We conclude by recalling a few of such topological definitions. A set $A \subset E$ is closed if its complement $E \setminus A$ is open. The properties in Lemma 2.1.3 can be restated for closed sets by saying that the empty set and $E$ are closed, and that finite unions as well as arbitrary intersections of closed sets are again closed (see Exercise 2.1). The **closure** $\overline{A}$ of a subset $A \subset E$ is defined by
$$\overline{A} := \bigcap \{B \supset A \text{ closed} \}.$$ 

It is by definition the smallest closed set containing $A$. Similarly, the **interior** $\mathring{A}$ of a set $A \subset E$ is defined as the largest open set contained in $A$, i.e.
$$\mathring{A} := \bigcup \{B \subset A \text{ open} \}.$$ 

Next, a subset $K \subset E$ is compact if every open covering of $K$ has an open subcovering. Finally, a mapping $f : E \to F$ between two topological spaces $(E, \mathcal{O}_E)$ and $(F, \mathcal{O}_F)$ is continuous if the pre-image of any open set is open, i.e. if
$$\Omega \in \mathcal{O}_F \implies f^{-1}(\Omega) \in \mathcal{O}_E.$$
2.2. Banach Spaces, Hilbert Spaces and Linear Operators

In this section, we consider complex vector spaces with additional structures like a norm and a scalar product. We also recall the notion of completeness and introduce linear operators. For more details and further reading we recommend the textbooks [126, 121, 110].

We begin with normed spaces.

**Definition 2.2.1.** Let $V$ be a complex vector space. A **norm** $\| \cdot \|$ is a mapping $\| \cdot \| : V \to \mathbb{R}_0^+$ with the following properties:

(i) **Homogeneity:** For all $x \in V$ and $\lambda \in \mathbb{C}$,
    $$\| \lambda x \| = |\lambda| \| x \| .$$

(ii) **Definiteness:** For all $x \in V$,
    $$\| x \| = 0 \iff x = 0 .$$

(iii) **Triangle inequality:** For all $x, y, z \in V$,
    $$\| x + y \| \leq \| x \| + \| y \| .$$

The pair $(V, \| \cdot \|)$ is a **normed space**.

Every normed space is a metric space (see Definition 2.1.1) with the metric $d(x, y) := \| x - y \|$.

We next recall the notion of completeness. A sequence $(x_n)_{n \in \mathbb{N}}$ of a metric space $(E, d)$ is a **Cauchy sequence** if for all $\varepsilon > 0$ there is $N \in \mathbb{N}$ with $d(x_n, x_m) < \varepsilon$ for all $m, n \geq N$. The sequence converges to $x$ if for all $\varepsilon > 0$ there is $N \in \mathbb{N}$ with $d(x, x_n) < \varepsilon$ for all $n \geq N$. The metric space is **complete** if every Cauchy sequence converges. A complete normed space is referred to as a **Banach space**.

In what follows, we usually assume that all metric spaces are complete. This is no major restriction because a non-complete metric space can made complete by going over to its abstract completion (for details see Exercise 2.2). To this end, one considers the set of Cauchy sequences in $E$ endowed with the distance function

$$d\left((x_n)_{n \in \mathbb{N}}, (y_n)_{n \in \mathbb{N}}\right) := \lim_{n \to \infty} d(x_n, y_n) . \quad (2.2.1)$$

This distance function is not necessarily a metric because the distance of two distinct sequences may be zero. Therefore, we introduce an equivalence relation on the Cauchy sequences by

$$(x_n)_{n \in \mathbb{N}} \simeq (y_n)_{n \in \mathbb{N}} \quad \text{if} \quad d\left((x_n)_{n \in \mathbb{N}}, (y_n)_{n \in \mathbb{N}}\right) = 0 . \quad (2.2.2)$$
Then the equivalence classes form a complete metric space. The constant sequences give an isometric embedding of $E$ into the completion.

We next specialize the setting by going over from a norm to a scalar product.

**Definition 2.2.2.** Let $V$ be a complex vector space. A **scalar product** $\langle \cdot, \cdot \rangle$ is a mapping

$$\langle \cdot, \cdot \rangle : V \times V \to \mathbb{C}$$

with the following properties:

(i) **Linearity the second argument:** For all $u, v, w \in V$ and $\alpha, \beta \in \mathbb{C}$,

$$\langle u | \alpha v + \beta w \rangle = \alpha \langle u | v \rangle + \beta \langle u | w \rangle .$$

(ii) **Hermitian symmetry:** For all $u, v \in V$,

$$\langle u | v \rangle = \overline{\langle v | u \rangle} .$$

(iii) **Positive definiteness:** For all $u \in V$,

$$\langle u | u \rangle \geq 0 \quad \text{and} \quad \langle u | u \rangle = 0 \iff u = 0 .$$

The pair $(V, \| \cdot \|)$ is a **scalar product space**.

Every scalar product space is a normed space with the norm $\| u \| := \sqrt{\langle u | u \rangle}$ (see Exercise [2.3]). A complete scalar product space is a **Hilbert space**.

Throughout this book, all Hilbert spaces will be **separable**, meaning that there is a countable subset $D \subset \mathcal{H}$ which is **dense** in the sense that its closure is the whole Hilbert space, $\overline{D} = \mathcal{H}$. In a separable Hilbert space, one can choose an **orthonormal Hilbert space basis** $(e_i)_{i \in I}$ characterized by the following properties:

(i) The index set $I$ is at most countable.

(ii) The system $(e_i)_{i \in I}$ is **orthonormal**, i.e.

$$\langle e_i | e_j \rangle = \delta_{ij}$$

(where $\delta_{ij}$ is the Kronecker delta).

(iii) The system $(e_i)_{i \in I}$ is **complete**, meaning that every vector $u \in \mathcal{H}$ has the representation as a (possibly infinite) linear combination

$$u = \sum_{i \in I} c_i e_i \quad \text{with} \quad c_i \in \mathbb{C} ,$$

where the series converges in $\mathcal{H}$.

Combining the properties (i) and (ii), every vector can be written as

$$u = \sum_{i \in I} \langle e_i | u \rangle e_i .$$

Moreover, the cardinality of the index set $I$ does not depend on the choice of the basis, making it possible to define the **dimension** of the Hilbert space

$$\dim \mathcal{H} := \# I \in \mathbb{N}_0 \cup \{ \infty \} .$$

We next introduce bounded linear operators.

**Definition 2.2.3.** Let $(V, \| \cdot \|_V)$ and $(W, \| \cdot \|_W)$ be normed spaces. A mapping

$$A : V \to W$$

is a **bounded linear operator** from $V$ to $W$ if it has the following properties:
(i) **Linearity:** For all \( u, v \in V \) and \( \alpha, \beta \in \mathbb{C} \),
\[
A(\alpha u + \beta v) = \alpha A(u) + \beta A(v).
\]

(ii) **Boundedness:** There is a constant \( c > 0 \) such that
\[
\|A(u)\|_W \leq c \|u\|_V \quad \text{for all } u \in V.
\]

Usually, one also writes \( A(u) \) simply as \( Au \).

The bounded linear operators form a complex vector space with the vector operations defined pointwise, i.e.
\[
(\alpha A + \beta B)(u) := \alpha Au + \beta Bu.
\]

This vector space is denoted by \( \mathcal{L}(V, W) \). It is endowed with the norm
\[
\|A\| := \sup_{u \in V, \|u\|_V = 1} \|Au\|_W,
\]
referred to as the sup-norm or the **operator norm**. If \( W \) is complete, then so is \( \mathcal{L}(V, W) \) (for details see Exercise 2.4). We remark that the linear operators from a normed space \( V \) to another normed space \( W \) are introduced similarly. They form a vector space denoted by \( \mathcal{L}(V, W) \), which is complete if and only if \( W \) is. We again refer for details to Exercise 2.4.

We will be concerned mainly with bounded linear operators on a Hilbert space \( (\mathcal{H}, \langle \cdot, \cdot \rangle) \). Two cases are of specific interest: mappings to the complex numbers, and mappings back to the Hilbert space \( \mathcal{H} \). The first case, the resulting operator is also referred to as a **bounded linear form**. These operators form the so-called **dual space** \( \mathcal{H}^* \),
\[
\mathcal{H}^* := \mathcal{L}(\mathcal{H}, \mathbb{C}).
\]

The dual space of a Hilbert space can be canonically identified with the Hilbert space using the following result.

**Theorem 2.2.4. (Fréchet-Riesz)** Let \( \mathcal{H} \) be a Hilbert space. Then for any bounded linear functional \( \phi \in \mathcal{H}^* \) there is a unique vector \( v \in \mathcal{H} \) such that
\[
\phi(u) = \langle v | u \rangle \quad \text{for all } u \in \mathcal{H}. \tag{2.2.3}
\]

In the case of a separable Hilbert space of interest here, this theorem can be understood in simple terms by writing \( v \) explicitly in an orthonormal Hilbert space basis \( (e_i)_{i \in I} \) as
\[
v = \sum_{i \in I} \phi(e_i) \ e_i.
\]

However, the proof of convergence of this series is not quite straightforward but makes it necessary to employ the orthogonal projection to closed subspaces of a Hilbert space. We refer the interested reader to [121], Section II.2, [110], Sections 6.2 and 6.3 or to Exercises 2.5 and 2.6.

We finally consider the space \( \mathcal{L}(\mathcal{H}, \mathcal{H}) \) of bounded linear endomorphisms. For brevity, this space is also denoted by \( \mathcal{L}(\mathcal{H}) \). In the context of such linear endomorphisms, the following additional notions are important.

**Definition 2.2.5.** A bounded linear operator \( A \in \mathcal{L}(\mathcal{H}) \) is **symmetric** if
\[
\langle Au \mid v \rangle = \langle u \mid Av \rangle \quad \text{for all } u, v \in \mathcal{H}.
\]

An operator \( U \in \mathcal{L}(\mathcal{H}) \) is **unitary** if it has a bounded inverse and if
\[
\langle Uu \mid Uv \rangle = \langle u \mid v \rangle \quad \text{for all } u, v \in \mathcal{H}.
\]

It has **finite rank** if its image \( A(\mathcal{H}) \) is a finite-dimensional.
For clarity, we remark that symmetric operators are also referred to as being selfadjoint. For bounded linear operators, these two notions are equivalent. However, for unbounded operators, these notions no longer coincide, and selfadjointness is the property which is more difficult to obtain, but which is needed for obtaining a spectral decomposition (see Section 3.2). With this in mind, when talking about bounded operators, in this book we usually prefer the notion of a symmetric operator.

If the Hilbert space $\mathcal{H}$ is finite-dimensional, symmetric and unitary operators can be diagonalized by choosing an orthonormal basis of eigenvectors (for details see standard textbooks on linear algebra like [96, 131]). The eigenvalues of a symmetric operator are all real, whereas the eigenvalues of a unitary operator have modulus one. The generalization of this result to the infinite-dimensional setting is provided by the spectral theorem. In most parts of this book, the operators have finite rank. With this in mind, we now explain how the results of linear algebra carry over to operators of finite rank (on possibly infinite-dimensional Hilbert spaces). The general spectral theorem for bounded symmetric operators, which will be needed mainly in Section 15, will be treated in Section 3.2 below.

Thus let $A \in \mathcal{L}(\mathcal{H})$ be a symmetric bounded operator of finite rank. Then its image $I := A(\mathcal{H})$ is finite-dimensional. Its orthogonal complement
\[ I^\perp := \{v \in \mathcal{H} \mid \langle v, u \rangle = 0 \text{ for all } u \in I\} \] (2.2.4)

is a closed subspace of $\mathcal{H}$, and every vector $u \in \mathcal{H}$ can be decomposed uniquely as
\[ u = u^\parallel + u^\perp \quad \text{with} \quad u^\parallel \in I, u^\perp \in I^\perp \] (2.2.5)
(see Exercise 2.7). Moreover, for any $u \in I^\perp$ the computation
\[ 0 = \langle A^2 u \mid u \rangle = \langle Au \mid Au \rangle \]
shows that $Au = 0$. Therefore, $A$ vanishes identically on $I^\perp$. Therefore, it suffices to consider the restriction $A|I : I \to I$. Being an operator on a finite-dimensional vector space, it can be diagonalized as in linear algebra.

When taking products of symmetric operators $A_1, \ldots, A_n$ of finite rank, one chooses $I$ as the finite-dimensional vector space spanned by the images of all the operators. Then the restrictions $A_k|I$ with $k = 1, \ldots, n$ all map $I$ to itself, making it possible to work again in a finite-dimensional subspace of $\mathcal{H}$.

### 2.3. Basics on Abstract Measure Theory

We now outline a few basics on abstract measure theory. For more details we refer to good standard textbooks like [126, 97, 15].

Let $\mathcal{F}$ be a set. A measure is a mapping which to subsets of $\mathcal{F}$ associates the corresponding measure, which can be thought of as the “volume” of the set. In order to get into a mathematically sensible setting, the measure must not be defined on any subset, but only on a distinguished family of subsets of $\mathcal{F}$. This subset must form a $\sigma$-algebra, which we now define.

**Definition 2.3.1.** A system of subsets $\mathcal{M}$ of $\mathcal{F}$ is a $\sigma$-algebra if it has the following properties:

(i) $\emptyset \in \mathcal{M}$

(ii) $\mathcal{M}$ is closed under taking complements, i.e.
\[ A \in \mathcal{M} \implies \mathcal{F} \setminus A \in \mathcal{M} \]
The sets in $\mathcal{M}$ are also referred to as the measurable sets.

In other words, a $\sigma$-algebra is closed under at most countable set operations.

We next introduce a measure as a mapping which to every measurable set associates its “volume.” This mapping is compatible with at most countable set operations, as is made precise by the notion of $\sigma$-additivity.

**Definition 2.3.2.** A measure $\rho$ is a mapping from a $\sigma$-algebra to the non-negative numbers or infinity, $\rho : \mathcal{M} \to \mathbb{R}^+_0 \cup \{\infty\}$, which has the following properties:

(i) $\rho(\emptyset) = 0$

(ii) $\sigma$-additivity holds, i.e. for every sequence $(A_n)_{n \in \mathbb{N}}$ of pairwise disjoint measurable sets,

$$\rho\left(\bigcup_{n \in \mathbb{N}} A_n\right) = \sum_{n=1}^{\infty} \rho(A_n).$$

The structure $(\mathcal{F}, \mathcal{M}, \rho)$ is a measure space.

On a measure space $(\mathcal{F}, \mathcal{M}, \rho)$ a notion of integration is introduced as follows. We begin with complex-valued functions. A function $f : \mathcal{F} \to \mathbb{C}$ is measurable if the pre-image of any open set is measurable. For a measurable function, the integral of its absolute value is well-defined, but could be infinite,

$$\int_{\mathcal{F}} |f(x)| \, d\rho(x) \in \mathbb{R}^+_0 \cup \{\infty\}.$$  \hspace{1cm} (2.3.1)

If this integral is finite, then the function $f$ is called integrable. In this case, the integral is well-defined and finite even without the absolute value,

$$\int_{\mathcal{F}} f(x) \, d\rho(x) \in \mathbb{C}.$$  

This notion of integrability can be understood immediately from the requirement that in integrals one must always avoid expressions of the form “$\infty-\infty$.” The integrable functions form a vector space denoted by $L^1(\mathcal{F}, d\rho)$. Similarly, the measurable functions $f$ whose power $|f|^p$ with $p \in (1, \infty)$ is integrable, form the vector space $L^p(\mathcal{F}, d\rho)$. Finally, the space $L^\infty(\mathcal{F}, d\rho)$ is defined as the functions which are essentially bounded in the sense that there is a number $c > 0$ such that the pre-image $|f|^{-1}((c, \infty])$ has $\rho$-measure zero. The spaces $L^p(\mathcal{F}, d\rho)$ with $p \in [1, \infty]$ are all Banach spaces if endowed with the corresponding norms

$$\|f\|_p := \left(\int_{\mathcal{F}} |f|^p \, d\rho\right)^{\frac{1}{p}} \quad \text{if} \quad p \in [0, \infty)$$

$$\|f\|_\infty := \inf \left\{c \geq 0 \left| |f|^{-1}((c, \infty]) = 0\right.\right\}.$$  

The space $L^2(\mathcal{F}, d\rho)$ is even a Hilbert space, endowed with the scalar product

$$\langle f \mid g \rangle_{L^2(\mathcal{F}, d\rho)} := \int_{\mathcal{F}} f(x) \overline{g(x)} \, d\rho(x).$$  \hspace{1cm} (2.3.2)
We remark that these constructions generalize immediately to functions taking values in a Banach or Hilbert spaces, if one simply replaces the absolute value in (2.3.1) by the norm on the Banach space and the inner product $fg$ in (2.3.2) by the Hilbert space scalar product.

We now specialize the setting by considering a class of measures which will be of major importance in this book, namely Borel measures. To this end, we assume that $(\mathcal{F}, \mathcal{O})$ is a topological space. Then the Borel algebra is defined as the smallest $\sigma$-algebra which contains all the open sets (see Exercise 2.8). A measure on the Borel algebra is referred to as a Borel measure. The Borel measures of relevance to us will typically harmonize with the topology in the following sense.

**Definition 2.3.3.** A Borel measure $\rho$ on $\mathcal{F}$ is called regular if for any measurable set $A$,

$$\rho(A) = \sup_{K \subset A \text{ compact}} \rho(K) = \inf_{\Omega \supset A \text{ open}} \rho(\Omega).$$

It is locally finite if every point of $\mathcal{F}$ has an open neighborhood $\Omega$ with $\rho(\Omega) < \infty$. Regular, locally finite Borel measures are also referred to as Radon measures.

We conclude this section by introducing a few other notions which will be needed later on. The first notion is the support of a measure.

**Definition 2.3.4.** Let $(\rho, \mathcal{M})$ be a measure on the topological space $(\mathcal{F}, \mathcal{O})$. The **support** of $\rho$ is defined as the complement of the largest open set of measure zero, i.e.

$$\text{supp } \rho := \mathcal{F} \setminus \bigcup \{ \Omega \subset \mathcal{F} \mid \Omega \text{ is open and } \rho(\Omega) = 0 \}.$$  \quad (2.3.3)

Note that the support is by definition a closed subset of $\mathcal{F}$.

Suppose that we want to compare two Radon measures $\rho$ and $\tilde{\rho}$ on $\mathcal{F}$. A natural idea is to take the difference $\rho - \tilde{\rho}$ and to look at its size. The difficulty is that for a measurable set $A \subset \mathcal{F}$, its measures $\rho(A)$ and $\tilde{\rho}(A)$ could both take the value $+\infty$, in which case their difference would be ill-defined. In order to avoid this problem, we use the regularity and exhaust by compact sets. Assuming that $\mathcal{F}$ is locally compact, the fact that the measures are locally finite implies that the measures of compact sets are always finite. This leads us to the following definition:

**Definition 2.3.5.** Given Radon measures $\rho$ and $\tilde{\rho}$ on a locally compact topological space $\mathcal{F}$, we define the measures $\mu^\pm$ by

$$\mu^+(A) = \sup_{K \subset A \text{ compact}} (\rho(K) - \tilde{\rho}(K))$$

and

$$\mu^-(A) = \sup_{K \subset A \text{ compact}} (\tilde{\rho}(K) - \rho(K)).$$

The difference of measures $\rho - \tilde{\rho}$ is said to have **bounded total variation** the measures $\mu^\pm$ are bounded, i.e. if

$$\mu^+(\mathcal{F}), \mu^-(\mathcal{F}) < \infty.$$  

If this is the case, the **total variation measure** $|\rho - \tilde{\rho}|$ is defined by

$$|\rho - \tilde{\rho}| = \mu^+ + \mu^-.$$  

We remark that $\rho - \tilde{\rho}$ can also be defined in the context of signed measures; we refer the interested reader to [97, §28] or [126, Section 6.1].

Another notion of measure theory which we will use frequently is the push-forward measure, which we now define (for more details see for example [15, Section 3.6] or
Exercise \[2.9\]. To this end, let \((\mathcal{F}, \mathcal{M}, \rho)\) be a measure space. Moreover, we consider a mapping 

\[ f : \mathcal{F} \rightarrow X, \]

where \(X\) is a point set. Then \(f\) induces a measure on \(X\). Indeed, we let \(\mathcal{M}_X\) be the set of all subsets \(\Omega \subset X\) whose pre-image \(f^{-1}(\Omega)\) is \(\rho\)-measurable. Using the elementary identities

\[ f^{-1}\left(\bigcup_{n \in \mathbb{N}} \Omega_n\right) = \bigcup_{n \in \mathbb{N}} f^{-1}(\Omega_n) \quad \text{and} \quad f^{-1}(X \setminus \Omega) = \mathcal{F} \setminus f^{-1}(\Omega), \]

one verifies that \(\mathcal{M}_X\) is indeed a \(\sigma\)-algebra on \(X\). The push-forward measure, denoted by \(f_*\rho\) is a measure on \(\mathcal{M}_X\) defined by

\[ (f_*\rho)(\Omega) := \rho(f^{-1}(\Omega)). \]

### 2.4. Distributions and Fourier Transform

We now recall a few basics on distribution theory and the Fourier transform. For more details, we recommend the textbook \[84, 121\] Sections V.3 and IX or \[120\] §2.1, §2.2 and Appendix A.

**Example 2.4.1. (The \(\delta\)-distribution)** The prime example of a distribution is Dirac’s \(\delta\)-distribution which in physics textbook is introduced as a “function” \(\delta(x)\) which is zero except at the origin, where it takes the value \(\infty\). The infinite contribution at the origin is “normalized” by demanding that its integral is equal to one. These properties can be summarized by saying that

\[ \int_{-\infty}^{\infty} f(x) \delta(x) \, dx = f(0). \tag{2.4.1} \]

There are various ways to make mathematical sense of this equation. One method is to regard the combination \(\delta(x) \, dx\) as a measure \(\delta_0\) supported at the origin of total volume one, i.e. \(\delta_0(\mathbb{R}) = 1\). In this way, the equation \eqref{eq:delta} makes sense if \(f\) is any continuous function. An alternative method is to take \eqref{eq:delta} as the definition of a linear functional on a space of suitable test functions \(f\). The latter method has the advantage that it makes it possible to even define the derivative of the \(\delta\)-distribution by

\[ \int_{-\infty}^{\infty} f(x) \delta'(x) \, dx = -f(0). \tag{2.4.2} \]

In order to allow for distributions to include an arbitrary number of derivatives, we choose a space of smooth test functions. Distributions will be defined as linear functionals on this space of test functions. Derivatives of distributions can be defined similar as in \eqref{eq:delta} by “formally integrating by parts.” Since a function \(g\) defines a linear functional \(T_g\) by integrating,

\[ T_g(f) := \int_{-\infty}^{\infty} g(x) f(x) \, dx, \tag{2.4.3} \]

every function gives rise to a corresponding distribution. In this sense, distributions are generalized functions. In order to make sure that the integral in \eqref{eq:delta} exists and is finite, it is a good idea to assume that the test function \(f\) has suitable decay properties at infinity.

In order to make these ideas mathematically precise, we need to specify the space of test functions. Moreover, we need to endow this space of test function with a topology. Then we can introduce distributions as the space of linear functionals on the test
functions. As we shall see, working with the right space of test functions, one can make mathematical sense of the Fourier transform for distributions.

In preparation of our constructions, we need to introduce the multi-index notation in \( \mathbb{R}^n \). A multi-index is an \( n \)-tuple \( \alpha = (p_1, \ldots, p_n) \in \mathbb{N}_0^n \) of non-negative integers. We define a monomial \( x^\alpha \) and a corresponding combination of partial derivatives \( \partial^\alpha \) by

\[
x^\alpha := (x_1)^{p_1} \cdots (x^n)^{p_n} \quad \text{and} \quad \partial^\alpha := \left( \frac{\partial}{\partial x^1} \right)^{p_1} \cdots \left( \frac{\partial}{\partial x^n} \right)^{p_n}.
\]

The order \( |\alpha| \) of the multi-index \( \alpha \) is defined by

\[
|\alpha| := p_1 + \cdots + p_n.
\]

On smooth functions in \( \mathbb{R}^n \) we introduce the Schwartz norms \( \| \cdot \|_{p,q} \) with \( p, q \in \mathbb{N}_0 \) by

\[
\| f \|_{p,q} := \max_{\alpha \text{ with } |\alpha| \leq p} \max_{\beta \text{ with } |\beta| \leq q} \sup_{x \in \mathbb{R}^n} |x^\alpha \partial^\beta f(x)|.
\]

The Schwartz space \( S(\mathbb{R}^n) \) is formed of all functions for which all the Schwartz norms are finite, i.e.

\[
S(\mathbb{R}^n) := \left\{ f \in C^\infty(\mathbb{R}^n, \mathbb{C}) \mid \| f \|_{p,q} < \infty \text{ for all } p, q \in \mathbb{N}_0 \right\}.
\]

We always consider complex-valued functions, but the constructions work similarly for real-valued functions. Defining the vector operators pointwise, \( S(\mathbb{R}^n) \) is a complex vector space. The functions in \( S(\mathbb{R}^n) \) are also referred to as the Schwartz functions. These functions as well as all their partial derivatives have rapid decay in the sense that multiplying them by a polynomial of arbitrary order still gives a bounded function.

The Schwartz norms induce the following topology on \( S(\mathbb{R}^n) \). We say that a set \( \Omega \subset S(\mathbb{R}^n) \) is open if for every \( f \in \Omega \) there exists \( p, q \in \mathbb{N}_0 \) and \( r > 0 \) such that the open \( r \)-ball corresponding to the norm \( \| \cdot \|_{p,q} \) is contained in \( \Omega \), i.e.

\[
\left\{ g \in S(\mathbb{R}^n) \mid \| f - g \|_{p,q} < r \right\} \subset \Omega.
\]

For our purposes, it is most convenient to characterize this topology by the corresponding notion of convergence. Then we can say that (see Exercise 2.10)

\[
f_n \to f \text{ in } S(\mathbb{R}^n) \quad \text{if} \quad \| f_n - f \|_{p,q} \to 0 \text{ for all } p, q \in \mathbb{N}_0.
\]

We point out that for a set to be open, it suffices to satisfy the condition (2.4.4) for some \( p, q \in \mathbb{N}_0 \). Alternatively, we could also have defined the open sets of \( S(\mathbb{R}^n) \) by the stronger condition that \( \Omega \) must be open with respect to all Schwartz norms. The latter definition would have given rise to a coarser topology, meaning that there would be fewer open sets. In contrast, demanding (2.4.4) merely for some \( p, q \in \mathbb{N}_0 \) gives a finer topology. Working with a fine topology has the following purpose: For a finer topology, fewer sequences converge (as is obvious in (2.4.5), where the sequence must converge for all Schwartz norms). As a consequence, there are more continuous linear functionals, simply because sequential continuity must be verified for fewer sequences. In other words, choosing a finer topology on a vector space has the effect that its dual space becomes larger (where the dual space is defined as the space of all continuous linear functionals). This is what we want, because defining distributions as this dual space will give us a sufficiently rich and general class of objects.

**Definition 2.4.2.** The space of tempered distributions, denoted by \( S'(\mathbb{R}^n) \), is defined as the dual space of the Schwartz space,

\[
S'(\mathbb{R}^n) := S^*(\mathbb{R}^n) = L(S(\mathbb{R}^n), \mathbb{C}).
\]
Let $T \in \mathcal{S}'(\mathbb{R}^n)$. The continuity of the functional $T : \mathcal{S}(\mathbb{R}^n) \to \mathbb{R}$ means that there are $p, q \in \mathbb{N}_0$ and a constant $c > 0$ such that (see again Exercise 2.10)

$$|T(f)| \leq c \|f\|_{p,q} \quad \text{for all } f \in \mathcal{S}(\mathbb{R}^n).$$

(2.4.6)

Tempered distributions can be regarded as generalized functions, as we now explain.

**Example 2.4.3. (regular distributions)** Let $g \in L^\infty(\mathbb{R}^n)$ be a bounded, measurable function. We define a linear functional $T_g$ on the Schwartz function by

$$T_g(f) := \int_{\mathbb{R}^n} g(x) f(x) \, dx.$$

This functional is continuous because

$$|T_g(f)| \leq \int_{\mathbb{R}^n} g(x) f(x) \, dx = \int_{\mathbb{R}^n} \frac{|g(x)|}{\left(1 + |x|^2\right)^{\frac{p+1}{2}}} |f(x)| \left(1 + |x|^2\right)^{-\frac{p+1}{2}} \, dx \leq C(n) \|g\|_{L^\infty(\mathbb{R}^2)} \|f\|_{n+1,0} \int_{\mathbb{R}^n} \frac{d^n x}{\left(1 + |x|^2\right)^{\frac{n+1}{2}}},$$

(2.4.7)

and the fact that the last integral is finite implies that the inequality holds for a suitable constant $c(n) > 0$ if we choose $p = n + 1$ and $q = 0$. Therefore, $T_g \in \mathcal{S}'(\mathbb{R}^n)$ is a tempered distribution.

In this way, every function $g \in L^\infty(\mathbb{R}^n)$ gives rise to a corresponding tempered distribution $T_g$. Let us verify that the corresponding linear mapping

$$T : L^\infty(\mathbb{R}^n) \to \mathcal{S}'(\mathbb{R}^n), \quad g \mapsto T_g$$

is injective.

(2.4.8)

To this end, we use the following denseness argument: Let $g \in L^\infty(\mathbb{R}^n)$ be non-zero. Then there is a bounded open set $\Omega \subset \mathbb{R}^n$ with $g|_{\Omega} \neq 0$. Choosing a compactly supported non-negative test function $\eta \in C_0^\infty(\mathbb{R}^n)$ with $\eta|_{\Omega} \equiv 1$, we find that the product $\eta g$ has compact support and is non-zero. Moreover, being bounded and compactly supported, this function is square integrable, $\eta g \in L^2(\mathbb{R}^n)$. Using that $C_0^\infty(\mathbb{R}^n)$ is dense in $L^2(\mathbb{R}^n)$, we conclude that there is a function $f \in C_0^\infty(\mathbb{R}^n)$ with

$$\int_{\mathbb{R}^n} f(x) (\eta g)(x) \neq 0.$$

Hence $T_g(\eta f) \neq 0$, showing that $T_g \neq 0$. We conclude that the mapping $T$ in (2.4.8) is indeed injective.

The fact that the mapping (2.4.8) is an embedding means that distributions can be regarded as “generalized functions.” Distributions which can be represented in the form $T_g$ with $g \in L^\infty(\mathbb{R}^n)$ are referred to as regular distributions.

We finally remark that $T_g$ can be defined more generally for functions $g$ which increase at most polynomially at infinity. But we do not need this generalization here. $\diamond$

We next define the *derivative* of distributions. In generalization of the procedure used in Example 2.4.1 for the $\delta$-distribution, the method is to “formally integrate by parts” and to use the resulting formula as the definition:

**Definition 2.4.4.** For any tempered distribution $T$ and multi-index $\alpha$, we define the distributional derivative $\partial^\alpha T$ by

$$(\partial^\alpha T)(f) = (-1)^{|\alpha|} T(\partial^\alpha f) \quad \text{for all } f \in \mathcal{S}(\mathbb{R}^n).$$
The estimate using (2.4.6)

\[ |(\partial^\alpha T)(f)| \leq |T(\partial^\alpha f)| \leq c \|\partial^\alpha f\|_{p,q} \leq c \|f\|_{p,q+|\alpha|} \]

shows that \(\partial^\alpha T\) is again a continuous linear functional.

We now come to the Fourier transformations. We first introduce it for Schwartz functions.

**Definition 2.4.5.** For \(f \in S(\mathbb{R}^n)\), we define the Fourier transform \((\mathcal{F}f)\) and the adjoint Fourier transform \((\mathcal{F}^* f)\) by

\[
(\mathcal{F}f)(p) = \int_{\mathbb{R}^n} f(x) e^{ipx} \, dx \tag{2.4.9}
\]

\[
(\mathcal{F}^* f)(x) = \int_{\mathbb{R}^n} f(p) e^{-ipx} \frac{dp}{(2\pi)^n}, \tag{2.4.10}
\]

where \(x,p \in \mathcal{M}\), and \(px = \langle p, x \rangle\) denotes the Minkowski inner product.

**Lemma 2.4.6.** The Fourier transform and its adjoint map Schwartz functions to Schwartz functions, \(\mathcal{F}, \mathcal{F}^* : S(\mathbb{R}^n) \to S(\mathbb{R}^n)\).

**Proof.** In order to prove (2.4.11), we differentiate (2.4.9) to obtain

\[
p^\alpha \partial^\beta_p (\mathcal{F}f)(p) = p^\alpha \int_{\mathbb{R}^n} f(x) x^\beta e^{ipx} \, dx = (-i)^{|\alpha|} \int_{\mathbb{R}^n} f(x) x^\beta (\partial^\alpha_x e^{ipx}) \, dx
\]

where in the last step we integrated by parts. Taking the absolute value, we obtain the estimate

\[
|p^\alpha \partial^\beta_p (\mathcal{F}f)(p)| \leq \int_{\mathbb{R}^n} |(\partial^\alpha f(x))| x^\beta | \, dx \leq \|\partial^\alpha f\|_{|\beta|+n+1,0} \leq \|f\|_{|\beta|+n+1,|\alpha|},
\]

where in (*) the integral was estimated similar as in (2.4.7). This estimate shows that the Fourier transform of a Schwartz function is again a Schwartz function. The estimate for the inverse Fourier transform is the same. \(\square\)

**Lemma 2.4.7. (Fourier inversion formula)** The Fourier transform and its adjoint are inverses of each other,

\[
\mathcal{F} \circ \mathcal{F}^* = \mathcal{F}^* \circ \mathcal{F} = 1_{S(\mathbb{R}^n)}. \tag{2.4.12}
\]

A detailed proof of this lemma can be found in [120, Theorem 2.2.4] or [84, Theorem 8.2.2]. With this in mind, we only give a sketch of the proof in one dimension. Computing the Fourier transforms \(\mathcal{F}^* \mathcal{F} f\) iteratively, one obtains \(f\) in a formal computation if one interchanges the integrals and uses the distributional identity

\[
\int_{-\infty}^{\infty} e^{ip(x-y)} \, dp = 2\pi \delta(x-y). \tag{2.4.13}
\]
One method of proving Lemma 2.4.7 is to make this computation mathematically sound by inserting a convergence-generating factor $e^{-\varepsilon p^2}$. More precisely,

$$
\mathcal{F}^*(\mathcal{F} f) = \lim_{\varepsilon \downarrow 0} \int_{-\infty}^{\infty} (\mathcal{F} f)(p) e^{-\varepsilon p^2} e^{-ipx} \frac{dp}{2\pi} \\
= \lim_{\varepsilon \downarrow 0} \int_{-\infty}^{\infty} \left( \int_{-\infty}^{\infty} f(y) e^{ipy} dy \right) e^{-\varepsilon p^2} e^{-ipx} \frac{dp}{2\pi} \\
= \lim_{\varepsilon \downarrow 0} \int_{-\infty}^{\infty} f(y) \left\{ \int_{-\infty}^{\infty} e^{i(p-y)x} e^{-\varepsilon p^2} \frac{dp}{2\pi} \right\} dy.
$$

(2.4.14)

The integral inside the curly brackets is Gaussian and can be computed explicitly. The resulting family of Gaussians tends to the $\delta$ distribution (see Exercise 2.12), making it possible to use standard estimates to obtain the result.

Having the Fourier transform for Schwartz functions at our disposal, we can now introduce the Fourier transform of tempered distributions. Similar as in the definition of the distributional derivative, the idea is to let the Fourier transform act on the test function. In order to make this idea more concrete, it is helpful to consider a regular tempered distribution $T_{\mathcal{F}g}$ corresponding to the Fourier transform of a Schwartz function $g$. Then

$$
T_{\mathcal{F}g}(f) = \int_{\mathbb{R}^n} (\mathcal{F}g)(p) f(p) \, dp = \int_{\mathbb{R}^n} \left( \int_{\mathbb{R}^n} g(x) e^{ipx} \, dx \right) f(p) \, dp \\
= \int_{\mathbb{R}^n} g(x) \left( \int_{\mathbb{R}^n} e^{ipx} f(p) \, dp \right) \, dx = T_g(\mathcal{F} f).
$$

The right side can be used to define the Fourier transform of a tempered distribution.

**Definition 2.4.8.** The Fourier transform and the adjoint Fourier transform of a tempered distribution $T \in \mathcal{S}'(\mathbb{R}^n)$ are defined by

$$
\mathcal{F}, \mathcal{F}^* : \mathcal{S}'(\mathbb{R}^n) \to \mathcal{S}'(\mathbb{R}^n), \quad (\mathcal{F} T)(f) = T(\mathcal{F} f), \quad (\mathcal{F}^* T)(f) = T(\mathcal{F}^* f).
$$

A direct computation shows that $\mathcal{F}^*$ is the inverse of $\mathcal{F}$ (see Exercise 2.11). Examples for how to compute Fourier transforms of distributions can be found in the exercises (see Exercise 2.13).

We close with two remarks. First, it is often very useful to consider the Fourier transform on $L^2$-functions, where the Fourier transform is unitary:

**Theorem 2.4.9.** (Plancherel) For any $f \in L^2(\mathbb{R}^n)$,

$$
\|\mathcal{F} f\|_{L^2(\mathbb{R}^n)} = (2\pi)^{\frac{n}{2}} \|f\|_{L^2(\mathbb{R}^n)}.
$$

We again omit the proof, which can be found in [120, Theorem 2.3.4] or [84, Theorem 9.2.2]. On a formal level, Plancherel’s formula is obtained by a direct computation using again (2.4.13). Similar to (2.4.14), this computation can be made mathematically sound by introducing a convergence-generating factor.

Finally, we remark that in some applications (for example when working in local charts on a manifold), it is not feasible to work with functions defined in all of $\mathbb{R}^n$ with suitable decay properties. In this case, instead of the Schwartz functions, one considers the space of test functions denoted by

$$
\mathcal{D}(\mathbb{R}^n) = C_0^\infty(\mathbb{R}^n),
$$

endowed with the topology induced by the family of norms $\|\cdot\|_q$ with $q \in \mathbb{N}_0$ given by

$$
\|f\|_q := \max_{\beta \text{ with } |\beta| \leq q} \sup_{x \in \mathbb{R}^n} \|\partial^\beta f(x)\|.
$$
2.5. Manifolds and Vector Bundles

We now recall the basic definitions of a manifold and vector bundle. For more details we refer to good textbooks like [107] or to the basic definitions in [111] §1 and §2.

**Definition 2.5.1.** A *topological manifold* of dimension \( n \in \mathbb{N} \) is a Hausdorff topological space \( \mathcal{M} \) which is \( \sigma \)-compact (i.e., can be written as an at most countable union of compact subsets) and has the property that every point in \( \mathcal{M} \) has an open neighborhood which is homeomorphic to an open subset of \( \mathbb{R}^n \).

More specifically, for every \( x \in \mathcal{M} \) there is an open neighborhood \( U \subset \mathcal{M} \) and a mapping \( \phi : U \to \mathbb{R}^n \) such that the image \( \phi(U) \) is open in \( \mathbb{R}^n \) and that the mapping \( \phi : U \to \phi(U) \) is a homeomorphism onto an open subset of \( \mathbb{R}^n \). Therefore, we may define the transition map (see Figure 2.4)

\[
\phi\big|_{U \cap U'} \circ (\phi'\big|_{U \cap U'})^{-1} : \phi\big(U \cap U'\big) \to \phi\big(U \cap U'\big).
\]

Being a mapping between two open subsets of \( \mathbb{R}^n \), it is clear what “differentiability” of this mapping means. A *differentiable manifold* is a topological manifold together with a complete atlas with the property that all transition maps are differentiable. Likewise, a *smooth manifold* is defined by the requirement that all transition maps are smooth. A differentiable manifold is equipped with a *differentiable structure*. In particular, we can define a differentiable mapping \( f : \mathcal{M} \to \mathcal{M}' \) between two differentiable manifolds \( \mathcal{M} \) and \( \mathcal{M}' \) by the requirement that it be differentiable in all charts.
We next introduce the notion of a *vector bundle*. This notion is helpful for the understanding of causal fermion systems because, under suitable regularity assumptions, a causal fermion system will give rise to a vector bundle over spacetime with the spin spaces as fibers (see Section 11.1). We begin with the notion of a topological vector bundle and explain the differentiable structure afterward.

**Definition 2.5.2.** Let $\mathcal{B}$ and $\mathcal{M}$ be topological spaces and $\pi : \mathcal{B} \to \mathcal{M}$ a continuous surjective map. Moreover, let $Y$ be a (real or complex) vector space and $G \subseteq \text{GL}(Y)$ a group acting on $Y$. Then $\mathcal{B}$ is a topological vector bundle with fiber $Y$ and structure group $G$ if every point $x \in \mathcal{M}$ has an open neighborhood $U$ equipped with a homeomorphism $\phi_U : \pi^{-1}(U) \to U \times Y$, called a local trivialization or bundle chart, such that the diagram

$$
\pi^{-1}(U) \xrightarrow{\phi_U} U \times Y \xrightarrow{\pi} U
$$

commutes, where the projection maps are $\pi$ and the projection onto the first factor, respectively. Furthermore, on overlaps $U \cap V$, we have

$$
\phi_U \circ \phi_V^{-1} \big|_{\{x\} \times Y} = g_{UV}(x) \quad \text{for all } x \in U \cap V,
$$

where $g_{UV} : U \cap V \to G$ is a continuous transition function.

A simple example of a vector bundle is the Cartesian product $\mathcal{B} = \mathcal{M} \times Y$. According to (2.5.1), a vector bundle has this product structure “locally” in $\pi^{-1}(U)$, but in general not globally.

A differentiable (or smooth) vector bundle is a topological vector bundle where the basis $\mathcal{M}$ is a differentiable (or smooth) manifold together with an atlas of bundle charts such that all transition maps are differentiable (or smooth, respectively).

**Exercises**

**Exercise 2.1. (Closed sets)** Show that the closed sets (defined as the complements of the open sets) have the following properties:

(i) The sets $\emptyset$ and $E$ are closed.

(ii) Closedness under finite unions:

$$
A_1, \ldots, A_n \text{ closed} \implies A_1 \cup \cdots \cup A_n \text{ closed}.
$$
(iii) Closedness under arbitrary intersections:
\[ A_\lambda \text{ closed } \quad \forall \lambda \in \Lambda \implies \bigcap_{\lambda \in \Lambda} A_\lambda \text{ closed}. \]

**Exercise 2.2. (Completion of metric spaces)**

(a) Show that the limit in (2.2.1) exists. Show that the so-defined distance function has the properties (ii) and (iii) in Definition 2.1.1.
(b) Verify that (2.2.2) is an equivalence relation. Why is the distance function well-defined on the equivalence classes. Verify property (i) in Definition 2.1.1.
(c) Modify the construction in order to form the completion of a scalar product space.

**Exercise 2.3. (Norm of a scalar product space)** Given a scalar product space \((V, \langle \cdot, \cdot \rangle)\), show that \[ \| u \| := \sqrt{\langle u | u \rangle} \] defines a norm (see Definitions 2.2.2 and 2.2.1).

**Exercise 2.4. (Completeness of \(L(V,W)\))**

(a) Show that the operator norm on \(L(V,W)\) is indeed a norm, i.e. verify that it has all the properties in Definition 2.2.1.
(b) Show that \(L(V,W)\) is complete if and only if \(W\) is complete.

**Exercise 2.5. (Orthogonal projection to closed subspaces of a Hilbert space)**

Let \((\mathcal{H}, \langle \cdot, \cdot \rangle)\) be a Hilbert space and \(V \subset \mathcal{H}\) a closed subspace.

(a) Show that parallelogram law: For all \(u, v \in \mathcal{H}\),
\[ \| u + v \|^2 + \| u - v \|^2 = 2 \| u \|^2 + 2 \| v \|^2. \]
(b) Given \(u \in \mathcal{H}\), let \((v_n)_{n \in \mathbb{N}}\) be a sequence in \(V\) which is a minimizing sequence of the distance to \(u\), i.e.
\[ \| u - v_n \| \to \inf_{v \in V} \| u - v \|. \]
Prove that the sequence \((v_n)_{n \in \mathbb{N}}\) converges. *Hint:* Apply the parallelogram law to show that the sequence is Cauchy. Then make use of the completeness of the Hilbert space.
(c) Show that the limit vector \(v := \lim_{n \to \infty} v_n\) has the property
\[ \langle u - v, w \rangle = 0 \quad \text{for all } w \in V. \]
In view of this equation, the vector \(v\) is also referred to as the orthogonal projection of \(u\) to \(V\).

In the finite-dimensional setting, the orthogonal projection can be given more explicitly as explained in Exercise 2.7.

**Exercise 2.6. (Proof of the Fréchet-Riesz theorem)** Let \(\phi \in \mathcal{H}^*\) be non-zero.

(a) Show that the kernel of \(\phi\) is a closed subspace of \(\mathcal{H}\).
(b) Apply the result of Exercise 2.5 to construct a nonzero vector \(v\) which is orthogonal to \(\ker \phi\). Show that this vector is unique up to scaling.
(c) Show that, after a suitable scaling, the vector \(v\) satisfies (2.2.3).
(d) Show that the vector \(v\) satisfying (2.2.3) is unique.

**Exercise 2.7. (Orthogonal complement of a finite-dimensional subspace)**

(a) Let \(I\) be a finite-dimensional subspace of the Hilbert space \((\mathcal{H}, \langle \cdot, \cdot \rangle)\). Show that its orthogonal complement \(I^\perp\) defined by (2.2.4) is again a complex vector space.
(b) Show that restricting the scalar product to \(I\), one gets again a Hilbert space. In particular, why is it again complete?
(c) Show that every vector \( u \in \mathcal{H} \) has a unique decomposition of the form \( (2.2.5) \).

*Hint:* Choosing an orthonormal basis \( e_1, \ldots, e_n \) of \( I \), a good ansatz for \( u \) is

\[
u = \sum_{k=1}^{n} (e_k|u) e_k.
\]

**Exercise 2.8. (Borel algebra)**

(a) Show that the power set of \( \mathcal{F} \) (i.e. the set of all subsets) forms a \( \sigma \)-algebra.

(b) Show that the intersection of \( \sigma \)-algebras is again a \( \sigma \)-algebra.

(c) Combine (a) and (b) to conclude that there is a smallest \( \sigma \)-algebra which contains all open subsets of \( \mathcal{F} \).

**Exercise 2.9. (Understanding the push-forward measure)** The purpose of this exercise is to introduce the so-called push-forward measure, which will be used later for the construction of causal fermion systems. Let \( \mathcal{M} \subset \mathbb{R}^3 \) be a smooth surface described by a parametrization \( \Phi \). More precisely, given an open subset \( \Omega \subset \mathbb{R}^2 \), we consider a smooth injective map

\[
\Phi : \Omega \rightarrow \mathbb{R}^3
\]

with the property that \( D\Phi |_p : \mathbb{R}^2 \rightarrow \mathbb{R}^3 \) has rank two for all \( p \in \Omega \). Then the surface \( \mathcal{M} \) is defined as the image \( \Phi(\Omega) \subset \mathbb{R}^3 \). We now introduce a measure \( \rho \) on \( \mathbb{R}^3 \) as the push-forward measure of the Lebesgue measure on \( \mathbb{R}^2 \) through \( \Phi \): Let \( \mu \) be the Lebesgue measure on \( \mathbb{R}^2 \). We define a set \( U \subset \mathbb{R}^3 \) to be \( \rho \)-measurable if and only if its pre-image \( \Phi^{-1}(U) \subset \mathbb{R}^2 \) is \( \mu \)-measurable. On the \( \rho \)-measurable sets we define the measure \( \rho \) by

\[
\rho(U) = \mu(\Phi^{-1}(U)).
\]

Verify that the \( \rho \)-measurable sets form a \( \sigma \)-algebra, and that \( \rho \) is a measure. What are the sets of \( \rho \)-measure zero? What is the support of the measure \( \rho \)?

Suppose that \( \Phi \) is no longer assumed to be injective. Is \( \rho \) still a well-defined measure? Is \( \rho \) well-defined if \( \Phi \) is only assumed to be continuous? What are the minimal regularity assumptions on \( \Phi \) needed for the push-forward measure to be well-defined? What is the support of \( \rho \) in this general setting?

**Exercise 2.10. (Topology on the Schwartz space)**

(a) Show that the topology on \( S(\mathbb{R}^n) \) defined by \( (2.4.4) \) gives rise to the notion of convergence \( (2.4.5) \).

(b) Show that a linear functional \( T : S(\mathbb{R}^n) \rightarrow \mathbb{R} \) is continuous if and only if there are \( p, q \in \mathbb{N}_0 \) and \( c > 0 \) such that the inequality \( (2.4.6) \) holds.

**Exercise 2.11. Prove the Fourier inversion formula for tempered distributions**

\[
\mathcal{F} \circ \mathcal{F}^* = \mathcal{F}^* \circ \mathcal{F} = 1_{S'(\mathbb{R}^n)} : S'(\mathbb{R}^n) \rightarrow S'(\mathbb{R}^n).
\]

*Hint:* Use Lemma \( 2.4.7 \) together with the definition of the Fourier transform of a tempered distribution.

**Exercise 2.12. (Dirac sequence)** Given \( \varepsilon > 0 \), consider the Gaussian

\[
\eta_\varepsilon(x) := \frac{1}{\sqrt{4\pi\varepsilon}} e^{-\frac{x^2}{4\varepsilon}}.
\]

(a) Show that \( \varepsilon \) is a Schwartz function.
(b) Show that the corresponding regular distribution converges to the δ distribution in the sense that for all $f \in \mathcal{S}(\mathbb{R})$,

$$\lim_{\varepsilon \searrow 0} T_{\eta_\varepsilon}(f) = \delta(f)$$

(the δ-distribution as introduced informally in Example 2.4.1 is defined by $\delta(f) = f(0)$ for all $f \in \mathcal{S}(\mathbb{R})$).

Exercise 2.13. (Fourier transform of δ distribution) Let $T \in \mathcal{S}'(\mathbb{R})$ be the regular distribution corresponding to the constant function, i.e.

$$T(f) := \int_{-\infty}^{\infty} f(x) \, dx.$$  

(a) Show that $T(f) = 2\pi (\mathcal{F}^*f)(0)$.

(b) Use the Fourier inversion theorem for tempered distributions to conclude that $\mathcal{F}T = 2\pi \delta$ (where $\delta$ is again the δ-distribution).

(c) Apply the Fourier inversion theorem again to compute the Fourier transform of the δ-distribution.

(d) Alternatively, one can compute these Fourier transforms directly working with convergence-generating factors in the style of (2.4.14). Do this carefully step by step, making sure that every computation step is mathematically well-defined.

Exercise 2.14. (Schwartz space)

(a) Schwartz functions decay faster than polynomially at infinity together with all their partial derivatives. More precisely, show that for every multi-index $\alpha$ and for every $N \in \mathbb{N}$ there exists $C_{\alpha,N} \geq 0$ such that

$$|D^\alpha f(x)| \leq \frac{C_{\alpha,N}}{1 + |x|^N} \quad \text{for all } x \in \mathbb{R}^n.$$  

(b) Decide which of the following functions belong to $\mathcal{S}(\mathbb{R})$ and which do not.

$$f_1(x) := e^{-x^2}, \quad f_2(x) := \frac{1}{1 + x^4}, \quad f_3(x) := \frac{e^{-x^2}}{2 + \sin(e^{x^2})}.$$  

Motivate your answers!

Exercise 2.15. (The principal value integral) For every $f \in \mathcal{S}(\mathbb{R})$ we define

$$A(f) := \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}\setminus(-\varepsilon,\varepsilon)} \frac{f(x)}{x} \, dx.$$  

Does the limit exist? In fact, $A$ is a tempered distribution. Show it.

Exercise 2.16. (Multiplication operators) Let $n \in \mathbb{N} \setminus \{0\}$.

(a) Let $g$ be a (not necessarily continuous) function $g : \mathbb{R}^n \to \mathbb{R}$ such that there is an $r \in (0, \infty)$ with $x \mapsto g(x)/(1 + |x|^r)$ in $L^1(\mathbb{R}^n)$. Show that the assignment $f \mapsto T_g(f) := \int_{\mathbb{R}^n} f(x) \cdot g(x) \, dx$ defines a tempered distribution $T_g \in \mathcal{S}(\mathbb{R}^n)\ast$.

(b) Construct a function $g \in C^\infty(\mathbb{R}^n) \cap L^1(\mathbb{R}^n)$ (such that $T_g \in \mathcal{S}(\mathbb{R}^n)\ast$) not pointwise bounded by any polynomial, i.e., such that there is no polynomial function $P : \mathbb{R} \to \mathbb{R}$ with $|g(x)| \leq P(|x|) \ \forall x \in \mathbb{R}^n$. 
EXERCISE 2.17. (Tempered distributions as integrals) Find a smooth function $g \in L^1(\mathbb{R}^n)$ which is not pointwise bounded by any polynomial, i.e. there is no polynomial $p$ such that
$$|g(x)| \leq p(|x|) \quad \text{for all } x \in \mathbb{R}^n.$$ Show that, nevertheless, the corresponding functional
$$T_g : \mathcal{S}(\mathbb{R}^n) \ni f \mapsto \int_{\mathbb{R}^n} g(x) f(x) \, d^n x$$ (2.5.3)
is a well-defined tempered distribution.

EXERCISE 2.18. (Approximating the Dirac delta distribution) Not all distributions can be written as in (2.5.3). Nevertheless, it can be shown that every tempered distribution can be approximated by such functionals. Let us see this in the concrete example of the Dirac delta distribution. Let $\varphi \in C^\infty_0(\mathbb{R}^n)$ fulfill $\varphi \geq 0$, supp $\varphi \subset B_1(0)$ and $\|\varphi\|_{L^1} = 1$. Define, for every $\varepsilon > 0$,
$$\varphi_\varepsilon(x) := \frac{1}{\varepsilon^n} \varphi \left( \frac{x}{\varepsilon} \right) \in C^\infty_0(\mathbb{R}^n) \quad \text{and} \quad \delta_\varepsilon := T_{\varphi_\varepsilon} \quad \text{(defined as in (2.5.3))}$$
Each functional $\delta_\varepsilon$ is a tempered distribution (why?). Show that, for every $f \in \mathcal{S}(\mathbb{R}^n)$,
$$\delta_\varepsilon(f) \to \delta(f) = f(0) \quad \text{as } \varepsilon \searrow 0.$$

EXERCISE 2.19. (Another regular tempered distribution) Let $n, k \in \mathbb{N}$ with $n > k$. Show that the mapping
$$\mathcal{S}(\mathbb{R}^n) \ni f \mapsto \int_{\mathbb{R}^n} \frac{f(x)}{|x|^k} \, d^n x \in \mathbb{C}$$is a well-defined tempered distribution.

EXERCISE 2.20. (Convolution) Let $n \geq 1$, $\alpha \in \mathbb{N}^n$, $f \in \mathcal{S}(\mathbb{R}^n)$ and $T \in \mathcal{S}^*(\mathbb{R}^n)$. Prove the following statements.
(a) $D^\alpha (f * T) = f * (D^\alpha T)$,
(b) $f * \delta = T_f$ where $\delta$ is the Dirac delta distribution.

EXERCISE 2.21. (Fourier transform on $\mathcal{S}(\mathbb{R}^n)$ and $\mathcal{S}^*(\mathbb{R}^n)$) Compute the Fourier transform of the following functions and tempered distributions.
(a) $f \in \mathcal{S}(\mathbb{R}^n)$ defined by $f(x) := e^{-\lambda x^2}$ for $\lambda \geq 0$;
(b) $T_g \in \mathcal{S}^*(\mathbb{R}^n)$ for the functions $g \in L^1(\mathbb{R})$ defined by
$$g(x) = e^{-|x|} \quad \text{and} \quad g(x) = \frac{1}{1 + x^2}.$$ 

EXERCISE 2.22. (Fourier transform on $L^1(\mathbb{R}^n)$) The functions in $L^1(\mathbb{R}^n)$ define tempered distributions by means of the identification $g \mapsto T_g$. As distributions they admit Fourier transform. However, for these functions the Fourier transform can also be defined directly via the usual integral form. The goal of this exercise is to show that this is indeed true and that the two definitions coincide.
(a) Show that the Fourier transform
$$(\mathcal{F}g)(p) := \frac{1}{(2\pi)^{n/2}} \int_{\mathbb{R}^n} e^{-ipx} g(x) \, d^n x$$
defines a map
$$\mathcal{F} : L^1(\mathbb{R}^n) \to C^0(\mathbb{R}^n) \cap L^\infty(\mathbb{R}^n).$$
Moreover, show that there exists a constant \( C_n \) such that
\[
\| \mathcal{F}g \|_\infty \leq C_n \| g \|_{L^1} \quad \text{for all } g \in L^1(\mathbb{R}^n).
\]
(b) Let \( g \in L^1(\mathbb{R}^n) \). Show that the Fourier transform of the distribution \( T_g \in \mathcal{S}^*(\mathbb{R}^n) \) satisfies
\[
\mathcal{F}(T_g) = T_{\mathcal{F}g}.
\]
Hint: You may use that \( \mathcal{S}(\mathbb{R}^n) \) is dense in \( L^1(\mathbb{R}^n) \).

Exercise 2.23. (On the topology of \( \mathcal{S}(\mathbb{R}^n) \)) Consider the Schwartz space \( \mathcal{S}(\mathbb{R}^n) \) equipped with the family of norms \( \| \cdot \|_{p,q} \), where \( p, q \in \mathbb{N}_0 \). Show the following statements.
(a) For any \( f, g \in \mathcal{S}(\mathbb{R}^n) \) the following series converges:
\[
d(f, g) := \sum_{p,q \in \mathbb{N}_0} \frac{1}{2^{p+q}} \frac{\| f - g \|_{p,q}}{1 + \| f - g \|_{p,q}}.
\]
Moreover it defines a metric on \( \mathcal{S}(\mathbb{R}^n) \).
(b) The metric space \( (\mathcal{S}(\mathbb{R}^n), d) \) is complete.
(c) Show that if a topology on a vector space is induced by a norm, then there is a neighborhood \( U \) of 0 such that for every other neighborhood \( V \) of 0, there is a positive number \( r \) such that \( U \subset r \cdot V \). Can the topology of \( \mathcal{S}(\mathbb{R}^n) \) be induced by a norm?
CHAPTER 3

Elements of Operator Theory

In this chapter we introduce some material from functional analysis which will be needed later in this book. More precisely, in Section 3.1 we explain the concept that linear operators with certain properties form a submanifold of the space of all linear operators. This concept will be useful when generalizing the causal action principle to causal variational principles in Chapter 6. In Section 3.2 we recall the spectral calculus for selfadjoint operators. Although this material is covered in most functional analysis lectures, in this book we do not expect that the reader is already familiar with this topic. The spectral theorem will be used only when developing functional analytic methods in spacetime in Chapter 15.

3.1. Manifolds of Operators

In this book, it is sometimes useful to observe that certain sets of operators on a Hilbert space form a smooth manifold. For the purposes in this book, it suffices to work out this concept in the case of a finite-dimensional Hilbert space (the generalization to the infinite-dimensional case is a bit more technical; for details see [73]).

We begin with a simple example which illustrates the basic concept.

Example 3.1.1. (Grassmann manifold) Given $f \in \mathbb{N}$, we consider $\mathbb{C}^n$ with the canonical scalar product $\langle . , . \rangle_{\mathbb{C}^f}$. Let $\mathcal{G}$ be the set of all orthogonal projection operators to one-dimensional subspaces of $\mathbb{C}^f$,

$$\mathcal{G} := \{ \pi_V \text{ orthogonal projection to a one-dimensional subspace } V \subset \mathbb{C}^f \}.$$

Let us verify that $\mathcal{G}$ is a smooth manifold of dimension

$$\dim \mathcal{G} = 2f - 2.$$

To this end, let $\pi_V$ be such a projection operator. We choose a unit vector $v$ which spans $V$. Next we let $V^\perp$ be the orthogonal complement of $V$ and $W = B_1(0) \cap V^\perp$ its unit ball. We consider the mapping

$$F : W \to \mathcal{G}, \quad u \mapsto \pi_{\text{span}(v+u)}.$$

It is verified by direct computation that $F$ is injective and that its image is an open neighborhood of $\pi_V$ in $\mathcal{G}$. Also, one easily verifies that it is a local homeomorphism and thus defines a chart

$$\phi = F^{-1}_{|F(W)} : F(W) \to W \subset V^\perp \simeq \mathbb{C}^{n-1} \simeq \mathbb{R}^{2n-2}.$$

Moreover, one verifies directly that all the charts obtained in this way form a smooth atlas. This manifold is called Grassmann manifold.

This concept can be generalized to so-called flag manifolds (see for example [99], page 142):
DEFINITION 3.1.2. (flag manifold) Given \( r \in \mathbb{N} \), we choose \( r \) integers
\[ 0 < d_1 < \cdots < d_r < f. \]
Consider a sequence \( (L_1, \ldots, L_r) \) of nested subspaces
\[ L_1 \subset \cdots \subset L_r \subset \mathbb{C}^f \]
with \( \dim L_i = d_i \) for \( i = 1, \ldots, r \). Then the set of such sequences \( (L_1, \ldots, L_r) \) is referred to as the flag manifold \( \mathcal{F}_{d_1, \ldots, d_r} \).

Here we do not need to verify that a flag manifold is indeed a manifold. Instead, it suffices to consider a specific set of operators which is related to a flag manifold. Namely, choosing again \( \mathcal{H} = \mathbb{C}^f \) as well as integers \( p, q \) with \( p + q \leq f \), we let \( \mathcal{F}^{p,q} \) be the set
\[ \mathcal{F}^{p,q} = \{ A \in \mathcal{L}(\mathcal{H}) \mid A \text{ is symmetric and has } p \text{ positive and } q \text{ negative eigenvalues} \}, \]
where we count the eigenvalues with multiplicities. Taking \( L_1 \) as the subspace spanned by the positive eigenvalues and \( L_2 \) as the image of \( A \), one gets a corresponding flag manifold with \( r = 2 \) and \( d_1 = p, d_2 = p + q \). But the operators in \( \mathcal{F}^{p,q} \) contain more information, namely the eigenvalues and the corresponding eigenspaces. Therefore the set \( \mathcal{F}^{p,q} \) can be regarded as a flag manifold with additional structures. We now prove that this set is again a smooth manifold, following the method in [66, Section 3].

PROPOSITION 3.1.3. The set \( \mathcal{F}^{p,q} \) is a smooth manifold of dimension
\[ \dim \mathcal{F}^{p,q} = 2f(p + q) - (p + q)^2. \]

PROOF. Let \( x \in \mathcal{F}^{p,q} \). We denote its image by \( I \) and set \( J = I^\perp \). Thus, using a block matrix representation in \( \mathbb{C}^f = I \oplus J \), the matrix \( x \) has the representation
\[ x = \begin{pmatrix} X & 0 \\ 0 & 0 \end{pmatrix}. \]
Now we let \( A \) be a symmetric linear operator on \( I \). By choosing its norm sufficiently small, we can arrange that the matrix \( (X + A) \) has again \( p \) positive and \( q \) negative eigenvalues. In particular, this operator is invertible. Next, we choose a linear operator \( B : J \to I \). We now form the matrix
\[ M := \begin{pmatrix} 1 & 0 \\ B^* (X + A)^{-1} & 1 \end{pmatrix} \begin{pmatrix} X + A & 0 \\ 0 & 0 \end{pmatrix} \begin{pmatrix} 1 & (X + A)^{-1}B \\ 0 & 1 \end{pmatrix} \]
\[ = \begin{pmatrix} X + A & B \\ B^* & B^*(X + A)^{-1}B \end{pmatrix}. \]
This matrix has again \( p \) positive and \( q \) negative eigenvalues (for details see Exercise 3.1). Thus, for sufficiently small \( \varepsilon \), we obtain the mapping
\[ \Lambda : (\text{Symm}(I) \oplus L(I, J)) \cap B_\varepsilon(0) \to \mathcal{F}^{p,q}, \quad (A, B) \mapsto M \]
(where Symm(I) denotes the symmetric linear operators). Let us verify that (again for sufficiently small \( \varepsilon \)) this mapping is a homeomorphism to an open neighborhood of \( x \in \mathcal{F}^{p,q} \). It is obvious from (3.1.2) that \( \Lambda \) is injective. In order to verify that it maps to an open neighborhood of \( x \), we let \( y \in \mathcal{F}^{p,q} \) with \( \|x - y\| < \delta \) (with \( \delta > 0 \) to be specified below). Diagonalizing \( y \) with a unitary matrix \( U \), we obtain the block matrix representation
\[ y = \begin{pmatrix} U_{11} & 0 \\ U_{21} & U_{22} \end{pmatrix} \begin{pmatrix} X + C & 0 \\ 0 & 0 \end{pmatrix} \begin{pmatrix} U_{11}^* & U_{21}^* \\ U_{12}^* & U_{22}^* \end{pmatrix}, \]
where $C$ is a symmetric linear operator on $I$. In the limit $y \to x$, the image of $y$ converges to the image of $x$, implying that the matrix $U_{11}$ becomes unitary. Therefore, for sufficiently small $\delta > 0$, the matrix $U_{11}$ is invertible, giving rise to the representation

$$y = \begin{pmatrix} 1 & 0 \\ U_{21} U_{11}^{-1} & 1 \end{pmatrix} \begin{pmatrix} U_{11} (X + C) U_{11}^* & 0 \\ 0 & 0 \end{pmatrix} \begin{pmatrix} 1 & (U_{11}^*)^{-1} U_{21}^* \\ 0 & 1 \end{pmatrix} .$$

This is of the form (3.1.1), and one can even read off $A$ and $B$,

$$A = U_{11} (X + C) U_{11}^* - x$$

$$B = (U_{11} (X + C) U_{11}^*) (U_{11}^*)^{-1} U_{21}^* .$$

We conclude that $\Lambda$ is a bijection to an open neighborhood of $x \in \mathcal{F}_{p,q}$. The continuity of $\Lambda$ and of its inverse are obvious. We have thus constructed a chart around $x$.

Performing the above construction around every point $x \in \mathcal{F}_{p,q}$ gives an atlas. By direct computation one verifies that the transition maps are smooth. We conclude that, with the above atlas, $\mathcal{F}_{p,q}$ is indeed a smooth manifold.

We finally determine its dimension. The linear operator $B$ is represented by a $(p+q) \times (f-p-q)$-matrix, giving rise to $2(p+q)(f-p-q)$ real degrees of freedom. The symmetric linear operator $A$, on the other hand, is represented by a Hermitian $(p+q) \times (p+q)$-matrix, described by $(p+q)^2$ real parameters. Adding these dimensions gives the result. □

### 3.2. The Spectral Theorem for Selfadjoint Operators

**Exercises**

Exercise 3.1. Let $x$ be a Hermitian $f \times f$-matrix of rank $p + q$ which (counting multiplicities) has $p$ positive and $q$ negative eigenvalues. Let $A$ be another $f \times f$-matrix (not necessarily Hermitian). Prove the following statements:

(a) The matrix $A^* x A$ has at most $p$ positive and at most $q$ negative eigenvalues.

*Hint:* Consider the maximal positive and negative definite subspaces of the bilinear forms $\langle.,x.\rangle_C$ and $\langle.,A^* x A.\rangle_C$. Use that

$$\langle u, A^* x A u \rangle_C = \langle (A u), x (A u) \rangle_C .$$

(b) If $A$ is invertible, then the matrix $A^* x A$ has again $p$ positive and $q$ negative eigenvalues.
CHAPTER 4

Spinors in Curved Spacetime

This chapter provides a brief introduction to spinors in curved spacetime. In order to make this book accessible to a broader readership, we mainly restrict attention to systems in Minkowski space. Nevertheless, many constructions and results carry over to curved spacetime in a straightforward way. The present section is intended for providing the necessary background for these generalizations. The reader not interested in gravity and Lorentzian geometry may skip this section. More specifically, the results of this chapter will be used only in Sections 11.3, 13.6 and 19.3. We follow the approach in [36].

4.1. Curved Spacetime and Lorentzian Manifolds

The starting point for general relativity is the observation that a physical process involving gravity can be understood in different ways. Consider for example an observer at rest on earth looking at a freely falling person (e.g. a diver who just jumped from a diving board). The observer at rest argues that the earth’s gravitational force, which he can feel himself, also acts on the freely falling person and accelerates him. The person in free fall, on the other hand, does not feel gravity. He can take the point of view that he himself is at rest, whereas the earth is accelerated towards him. He concludes that there are no gravitational fields, and that the observer on earth merely feels the force of inertia corresponding to his acceleration. Einstein postulated that these two points of view should be equivalent descriptions of the physical process. More generally, it depends on the observer whether one has a gravitational force or an inertial force. In other words,

\textit{equivalence principle: no physical experiment can distinguish between gravitational and inertial forces.}

In mathematical language, observers correspond to coordinate systems, and so the equivalence principle states that the physical equations should be formulated in general (i.e. “curvilinear”) coordinate systems, and should in all these coordinate systems have the same mathematical structure. This means that the physical equations should be invariant under diffeomorphisms, and thus spacetime is to be modeled by a \textit{Lorentzian manifold} $(\mathcal{M}, g)$.

A Lorentzian manifold is “locally Minkowski space” in the sense that at every spacetime point $p \in \mathcal{M}$, the corresponding \textit{tangent space} $T_p\mathcal{M}$ is a vector space endowed with a scalar product $\langle \cdot, \cdot \rangle_p$ of signature $(+−−−)$. Therefore, we can distinguish between spacelike, timelike and null tangent vectors. Defining a non-spacelike curve $q(\tau)$ by the condition that its tangent vector $u(\tau) \in T_{q(\tau)}\mathcal{M}$ be everywhere non-spacelike, our above definition of light cones and the notion of causality immediately carry over to a Lorentzian manifold. In a coordinate chart, the scalar product $\langle \cdot, \cdot \rangle_p$ can be represented in the form (1.2.1) where $g_{jk}$ is the so-called \textit{metric tensor}. In contrast to Minkowski space, the metric tensor is not a constant matrix but depends on the spacetime point, $g_{jk} = g_{jk}(p)$. Its ten components can be regarded as the relativistic analogue of Newton’s
gravitational potential. For every \( p \in \mathcal{M} \) there are coordinate systems in which the metric tensor coincides with the Minkowski metric up to second order,

\[
g_{jk}(p) = \text{diag}(1, -1, -1, -1), \quad \partial_j g_{kl}(p) = 0.
\]

Such Gaussian normal coordinates correspond to the reference frame of a “freely falling observer” who feels no gravitational forces. However, it is in general impossible to arrange that also \( \partial_j g_{kl}(p) = 0 \). This means that by going into a suitable reference frame, the gravitational field can be transformed away locally (=in one point), but not globally. With this in mind, a reference frame corresponding to Gaussian normal coordinates is also called a local inertial frame.

The physical equations can be carried over to a Lorentzian manifold by the prescription that they should in a local inertial frame have the same form as in Minkowski space; this is referred to as the strong equivalence principle. It amounts to replacing all partial derivatives by the corresponding covariant derivatives \( \nabla \) of the Levi-Civita connection; we write symbolically

\[
\partial \longrightarrow \nabla.
\]

The gravitational field is described via the curvature of spacetime. More precisely, the Riemannian curvature tensor is defined by the relations

\[
R^i_{jk} u^l = \nabla_j \nabla_k u^l - \nabla_k \nabla_j u^l.
\]

Contracting indices, one obtains the Ricci tensor \( R_{jk} = R^i_{jk} \) and scalar curvature \( R = R^j_j \).

The relativistic generalization of Newton’s gravitational law are the Einstein equations

\[
R_{jk} - \frac{1}{2} R g_{jk} = 8 \pi \kappa T_{jk},
\]

where \( \kappa \) is the gravitational constant. Here the energy-momentum tensor \( T_{jk} \) gives the distribution of matter and energy in spacetime.

### 4.2. The Dirac Equation in Curved Spacetime

Dirac spinors are often formulated on a manifold using frame bundles, either an orthonormal frame \([9, 85]\) or a Newman-Penrose null frame \([117, 20]\). We here outline an equivalent formulation of spinors in curved spacetime in the framework of a \( U(2, 2) \) gauge theory (for details see \([36]\)). We begin with constructions in local coordinates, whereas global issues like topological obstructions to the existence of spin structures will be discussed in Section 4.4 below. We let \( \mathcal{M} \) be a 4-dimensional manifold (without Lorentz metric) and define the spinor bundle \( S\mathcal{M} \) as a vector bundle over \( \mathcal{M} \) with fiber \( \mathbb{C}^4 \). The fibers are endowed with a scalar product \( \langle \cdot, \cdot \rangle \) of signature \((2, 2)\), referred to as the spin scalar product. Sections in the spinor bundle are called spinors or wave functions. In local coordinates, a spinor is represented by a 4-component complex function on spacetime, usually denoted by \( \psi(x) \). Choosing at every spacetime point a pseudo-orthonormal basis \((e_\alpha)_{\alpha=1,\ldots,4}\) in the fibers,

\[
\langle e_\alpha | e_\beta \rangle = s_\alpha \delta_{\alpha\beta}, \quad s_1 = s_2 = 1, \ s_3 = s_4 = -1
\]

and representing the spinors in this basis, \( \psi = \psi^\alpha e_\alpha \), the spin scalar product takes again the form \((1.3.7)\). Clearly, the basis \((e_\alpha)\) is not unique, but at every space-point can be transformed according to

\[
e_\alpha \longrightarrow (U^{-1})^\beta_\alpha e_\beta,
\]
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where \( U \) is an isometry of the spin scalar product, \( U \in \text{U}(2,2) \). Under this basis transformation the spinors behave as follows,

\[
\psi^\alpha(x) \longrightarrow U^\beta_\alpha(x) \psi^\beta(x).
\]  

(4.2.2)

In view of the analogy to gauge theories, we interpret this transformation of the wave functions as a local gauge transformation with gauge group \( \mathcal{G} = \text{U}(2,2) \). We refer to a choice of the spinor basis \( (e_\alpha) \) as a gauge.

Our goal is to formulate classical Dirac theory in such a way that the above \( \text{U}(2,2) \) gauge transformations correspond to a physical symmetry, the \( \text{U}(2,2) \) gauge symmetry. To this end, we shall introduce a Dirac-type operator as the basic object on \( \mathcal{M} \), from which we will then deduce the Lorentz metric and the gauge potentials. We define a differential operator \( \mathcal{D} \) of first order on the wave functions by the requirement that in a chart and gauge it should be of the form

\[
\mathcal{D} = iG^j(x) \frac{\partial}{\partial x^j} + B(x)
\]  

(4.2.3)

with suitable \((4 \times 4)\)-matrices \( G^j \) and \( B \). This definition does not depend on coordinates and gauge, although the form of the matrices \( G^j \) and \( B \) clearly does. More precisely, under a change of coordinates \( x^i \rightarrow \tilde{x}^i \) the operator (4.2.3) transforms into

\[
i \left( G^k(x) \frac{\partial}{\partial \tilde{x}^k} \right) \frac{\partial}{\partial \tilde{x}^j} + B(\tilde{x}),
\]  

(4.2.4)

whereas a gauge transformation \( \psi \rightarrow U \psi \) yields the operator

\[
UDU^{-1} = i \left( UG^jU^{-1} \right) \frac{\partial}{\partial x^j} + \left( UBU^{-1} + iUG^j(\partial_jU^{-1}) \right).
\]  

(4.2.5)

We define an operator of Dirac type by the requirement that by choosing suitable coordinates and gauge, one can arrange that the coefficient matrices \( G^j \) of the partial derivatives “coincide locally” with the Dirac matrices of Minkowski space.

**Definition 4.2.1.** A differential operator \( \mathcal{D} \) of first order is called operator of Dirac type (or Dirac-type operator) if for every \( p \in \mathcal{M} \) there is a chart \((x^i, U)\) around \( p \) and a gauge \((e_\alpha)_{\alpha=1,...,4}\) such that \( \mathcal{D} \) is of the form (4.2.3) with

\[
G^j(p) = \gamma^j,
\]  

(4.2.6)

where the \( \gamma^j \) are the Dirac matrices of Minkowski space in the Dirac representation (1.3.3).

It may seem unconventional that we defined Dirac-type operators without having a spinor bundle nor a connection. We shall now construct from the Dirac-type operator a gauge-covariant derivative \( D \), also referred to as the spin derivative. To this end, we first write the transformation law (4.2.2) in the shorter form

\[
\psi(x) \longrightarrow U(x) \psi(x)
\]  

(4.2.7)

with \( U \in \text{U}(2,2) \). Clearly, partial derivatives of \( \psi \) do not have a nice behavior under gauge transformations because we pick up derivatives of \( U \). This problem disappears if instead of partial derivatives we consider gauge-covariant derivatives

\[
D_j = \partial_j - iA_j,
\]  

(4.2.8)

provided that the gauge potentials transform according to

\[
A_j \longrightarrow U A_j U^{-1} + iU (\partial_j U^{-1}).
\]  

(4.2.9)
Namely, a short calculation shows that the gauge-covariant derivative behaves under gauge transformations according to

\[ D_j \rightarrow U D_j U^{-1}, \tag{4.2.10} \]

and thus the gauge-covariant derivatives of \( \psi \) obey the simple transformation rule

\[ D_j \psi \rightarrow U D_j \psi. \]

Our task is to find matrices \( A_j \) which transform under local gauge transformations according to (4.2.9). This construction will also reveal the structure of the matrix \( B \), and this will finally lead us to the definition of the Dirac operator, which involves precisely the gravitational and electromagnetic fields.

In the chart and gauge where (4.2.6) holds, it is obvious from (1.3.1) that the anti-commutator of the matrices \( G_j(p) \) gives the Minkowski metric. Using the transformation rules (4.2.4) and (4.2.5), one sees that in a general coordinate system and gauge, their anti-commutator defines a Lorentz metric,

\[ g_{jk}(x) = \frac{1}{4!} \epsilon_{jklm} G^l(x) G^m(x). \tag{4.2.11} \]

Moreover, using that the Dirac matrices in Minkowski space are symmetric w.r.t. the spin scalar product (see (1.3.8)), one sees that the same is true for the matrices \( G^j \), i.e.

\[ \langle G^i \psi | \phi \rangle = -\langle \psi | G^i \phi \rangle \quad \text{for all } \psi, \phi. \tag{4.2.12} \]

Via (4.2.11), the Dirac-type operator induces on the manifold a Lorentzian structure. We refer to the matrices \( G_j \) as the Dirac matrices in curved spacetime. Since we can arrange that these matrices coincide locally with the Dirac matrices of Minkowski space, all manipulations of Dirac matrices can be performed at any given spacetime point in an obvious way. In particular, the pseudo-scalar operator (1.3.20) now takes the more general form

\[ \Gamma(x) = \frac{1}{4!} \epsilon_{jklm} G^j(x) G^k(x) G^l(x) G^m(x). \tag{4.2.13} \]

where the anti-symmetric tensor \( \epsilon_{jklm} = \sqrt{|\det g|} \epsilon_{jklm} \) differs from the anti-symmetric symbol \( \epsilon_{jklm} \) by the volume density. The pseudo-scalar operator gives us again the notion of even and odd matrices and of chirality (1.3.21). Furthermore, we introduce the bilinear matrices \( \sigma^{jk} \) by

\[ \sigma^{jk}(x) = \frac{i}{2} [G^j, G^k]. \tag{4.2.14} \]

As in Minkowski space, the matrices

\[ G^j, \quad \Gamma G^j, \quad 1, \quad i \Gamma, \quad \sigma^{jk} \tag{4.2.14} \]

form a basis of the 16-dimensional (real) vector space of symmetric matrices (symmetric with respect to the spin scalar product \( \langle \cdot | \cdot \rangle \)). The matrices \( G^j \) and \( \Gamma G^j \) are odd, whereas \( 1, i \Gamma \) and \( \sigma^{jk} \) are even.

For the construction of the spin connection we must clearly consider derivatives. The Lorentzian metric (4.2.11) induces the Levi-Civita connection \( \nabla \), which defines the covariant derivative of tensor fields. Taking covariant derivatives of the Dirac matrices, \( \nabla_k G^j = \partial_k G^j + \Gamma^j_{kl} G^l \), we obtain an expression which behaves under coordinate transformations like a tensor. However, it is not gauge covariant, because a gauge transformation (4.2.7) yields contributions involving first derivatives of \( U \). More precisely, according
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4.2.5. \( \nabla_k G^j \rightarrow \nabla_k (UG^j U^{-1}) = (\nabla_k U)G^j U^{-1} + UG^j (\partial_k U)U^{-1} \)

\[ = U(\nabla_k G^j)U^{-1} - [U(\partial_k U), UG^j U^{-1}] . \]

(4.2.15)

We can use the second summand in (4.2.15) to partially fix the gauge.

**Lemma 4.2.2.** For every spacetime point \( p \in \mathcal{M} \) there is a gauge such that

\( \nabla_k G^j(p) = 0 \) for all indices \( j, k \).

(4.2.16)

**Proof.** We start with an arbitrary gauge and construct the desired gauge with two subsequent gauge transformations:

1. The matrix \( \partial_j \Gamma \) is odd, because

\[ 0 = \partial_j 1 = \partial_j (\Gamma \Gamma) = (\partial_j \Gamma) \Gamma + \Gamma (\partial_j \Gamma) . \]

As a consequence, the matrix \( i \Gamma (\partial_j \Gamma) \) is selfadjoint. We can thus perform a gauge transformation \( U \) with \( U(p) = 1, \partial_j U(p) = \frac{1}{2} \Gamma (\partial_j \Gamma) \). In the new gauge the matrix \( \partial_j \Gamma(p) \) vanishes,

\[ \partial_j \Gamma |p \rightarrow \partial_j (U \Gamma U^{-1}) |p = \partial_j \Gamma |p + \frac{1}{2} \left[ \Gamma (\partial_j \Gamma), \Gamma |p \right] = \partial_j \Gamma |p - \Gamma^2 (\partial_j \Gamma) |p = 0 . \]

Differentiating the relation \( \{ \Gamma, G^j \} = 0 \), one sees that the matrix \( \nabla_k G^j |p \) is odd. We can thus represent it in the form

\[ \nabla_k G^j |p = \Lambda^j_{km} G^m |p + \Theta^j_{km} \Gamma G^m \]

(4.2.17)

with suitable coefficients \( \Lambda^j_{km} \) and \( \Theta^j_{km} \).

This representation can be further simplified: According to Ricci’s Lemma, \( \nabla_n g^{jk} = 0 \). Expressing the metric via the anti-commutation relations and differentiating through with the Leibniz rule, we obtain

\[ 0 = \{ \nabla_n G^j, G^k \} + \{ G^j, \nabla_n G^k \} \]

\[ = 2 \Lambda^j_{nm} g^{mk} \Theta^j_{nm} 2i \Gamma \sigma^{mk} + 2 \Lambda^k_{nm} g^{mj} \Theta^j_{nm} 2i \Gamma \sigma^{mj} \]

(4.2.18)

and thus

\[ \Lambda^j_{nm} g^{mk} = -\Lambda^k_{nm} g^{mj} . \]

(4.2.19)

In the case \( j = k \neq m \), (4.2.18) yields that \( \Theta^j_{nm} = 0 \). For \( j \neq k \), we obtain \( \Theta^j_{nj} \sigma^{lk} + \Theta^k_{nk} \sigma^{kj} = 0 \) and thus \( \Theta^j_{nj} = \Theta^k_{nk} \) (\( j \) and \( k \) denote fixed indices, no summation is performed). We conclude that there are coefficients \( \Theta_k \) with

\[ \Theta^j_{km} = \Theta_k \delta^j_m . \]

(4.2.20)

2. We perform a gauge transformation \( U \) with \( U(p) = 1 \) and

\[ \partial_k U = -\frac{1}{2} \Theta_k \Gamma - \frac{i}{4} \Lambda^m_{kn} g^{nl} \sigma_{ml} . \]
Using the representation (4.2.17) together with (4.2.19) and (4.2.20), the matrix \( \nabla_k G^j \) transforms into

\[
\nabla_k G^j \rightarrow \nabla_k G^j + [\partial_k U, G^j] = \Lambda^j_{km} G^m + \Theta_k \Gamma G^j - \Theta_k \Gamma G^j - \Lambda^j_{km} g^{nl} [\sigma_{ml}, G^j] = \Lambda^j_{km} G^m - \frac{1}{2} \Lambda^j_{km} g^{nj} G^m = 0 .
\]

We call a gauge satisfying condition (4.2.16) a normal gauge around \( p \). In order to analyze the remaining gauge freedom, we let \( U \) be a transformation between two normal gauges. Then according to (4.2.15) and (4.2.16), the commutator \( [U(\partial_k U^{-1}), UG^j U^{-1}] \) vanishes at \( p \) or, equivalently,

\[
[i(\partial_k U^{-1}) U, G^j]|_p = 0 .
\]

As is easily verified in the basis (4.2.14) using the anti-commutation relations, a matrix which commutes with all Dirac matrices is a multiple of the identity matrix. Moreover, the matrix \( i(\partial_j U^{-1}) U \) is selfadjoint because \( (i(\partial_j U^{-1}) U)^* = -iU^{-1} (\partial_j U) = -i \partial_j(U^{-1} U) + i(\partial_j U^{-1}) U = i(\partial_j U^{-1}) U \). We conclude that the matrix \( i(\partial_j U^{-1}) U \) is a real multiple of the identity matrix. Transforming it unitarily with \( U \), we see that it also coincides with the matrix \( iU (\partial_j U^{-1}) \). Restricting attention to normal gauges, it is easy to find expressions with the required behavior (4.2.9) under gauge transformations. Namely, setting

\[
a_j = \frac{1}{4} \text{Re} \text{Tr} \left( G_j B \right) \mathbb{1} ,
\]

where "Tr" denotes the trace of a \( 4 \times 4 \)-matrix, one sees from (4.2.5) that

\[
a_j \rightarrow a_j + \frac{1}{4} \text{Re} \text{Tr} \left( G_j G^k i(\partial_k U^{-1}) U \right) \mathbb{1} = a_j + iU(\partial_j U^{-1}) .
\]

We can identify the \( a_j \) with the gauge potentials \( A_j \) and use (4.2.8) as the definition of the spin connection.

**Definition 4.2.3.** The spin derivative \( D \) is defined by the condition that it behaves under gauge transformations (4.2.7) according to (4.2.10), and that in normal gauges around \( p \) it has the form

\[
D_j(p) = \frac{\partial}{\partial x^j} - ia_j
\]

with the potentials \( a_j \) according to (4.2.21).

In general gauges, the spin derivative can be written as

\[
D_j = \frac{\partial}{\partial x^j} - iE_j - ia_j
\]

with additional matrices \( E_j(x) \), which involve the Dirac matrices and their first derivatives. The components of \( E_j \) are sometimes referred to as spin coefficients. A short calculation shows that the trace of the matrix \( E_j \) does not change under gauge transformations, and since it vanishes in normal gauges, we conclude that the matrices \( E_j \) are
trace-free. It is straightforward to verify that they are given explicitly by (for details see Appendix A)

\[ E_j = \frac{i}{2} \Gamma (\partial_j \Gamma) - \frac{i}{16} \mathrm{Tr}(G^m \nabla_j G^m) G_m G_n + \frac{i}{8} \mathrm{Tr}(\Gamma G_j \nabla_m G^m) \Gamma . \]  

(4.2.24)

In the next two theorems we collect the basic properties of the spin connection.

**Theorem 4.2.4.** The spin derivative satisfies for all wave functions \( \psi, \phi \) the equations

\[ [D_k, G^j] + \Gamma^j_{kl} G^l = 0 \]  

(4.2.25)

\[ \partial_j <\psi|\phi> = <D_j \psi|\phi> + <\psi|D_j \phi>. \]  

(4.2.26)

**Proof.** The left side of (4.2.25) behaves under gauge transformations according to the adjoint representation \( U \to U . U^{-1} \) of the gauge group. Therefore, it suffices to check (4.2.25) in a normal gauge, where

\[ [D_k, G^j] + \Gamma^j_{kl} G^l = \nabla_k G^j - \frac{i}{4} \text{Re} \text{Tr} (G_j B) [1, G^j] = 0 . \]

Since both sides of (4.2.26) are gauge invariant, it again suffices to consider a normal gauge. The statement is then an immediate consequence of the Leibniz rule for partial derivatives and the fact that the spin derivative differs from the partial derivative by an imaginary multiple of the identity matrix (4.2.22).

The identity (4.2.25) means that the coordinate and gauge invariant derivative of the Dirac matrices vanishes. The relation (4.2.26) shows that the spin connection is compatible with the spin scalar product.

We define **torsion** \( T \) and **curvature** \( R \) of the spin connection as the following 2-forms,

\[ T_{jk} = \frac{i}{2} ([D_j, G_k] - [D_k, G_j]) , \quad R_{jk} = \frac{i}{2} [D_j, D_k] . \]

**Theorem 4.2.5.** The spin connection is torsion-free. Curvature has the form

\[ R_{jk} = \frac{1}{8} R_{mnjk} \sigma^{mn} + \frac{1}{2} (\partial_j a_k - \partial_k a_j) , \]

where \( R_{mnjk} \) is the the Riemannian curvature tensor and the \( a_j \) are given by (4.2.21).

**Proof.** The identity (4.2.25) yields that

\[ [D_j, G_k] = [D_j, g_{kl} G^l] = (\partial_j g_{kl}) G^l - g_{kl} \Gamma^l_{jm} G^m = \Gamma^m_{jk} G_m . \]

Thus, using that the Levi-Civita connection is torsion-free, we obtain

\[ T_{jk} = \frac{i}{2} (\Gamma^m_{jk} - \Gamma^m_{kj}) G_m = 0 . \]

Next, again using (4.2.25), we can rewrite the covariant derivative as a spin derivative,

\[ G_l \nabla_k u^l = [D_k, G_l u^l] . \]

Iterating this relation, we can express the Riemann tensor (4.2.2) by

\[ G_i \nabla_{jk} u^i = [D_j, [D_k, G_l u^l]] - [D_k, [D_j, G_l u^l]] \]

\[ = [[D_j, D_k], G_l u^l] = -2i \left[ R_{jk}, G_l u^l \right] . \]

This equation determines curvature up to a multiple of the identity matrix,

\[ R_{jk}(x) = \frac{1}{8} R_{mnjk} \sigma^{mn} + \lambda_{jk} 1 . \]
Thus it remains to compute the trace of curvature,
\[
\frac{1}{4} \text{Tr}(\mathcal{R}_{jk}) 1 = \frac{1}{8} \text{Tr} \left( \partial_j A_k - \partial_k A_j \right) 1 = \frac{1}{2} \left( \partial_j a_k - \partial_k a_j \right),
\]
where we used (4.2.23) and the fact that the matrices $E_j$ are trace-free.

We come to the physical interpretation of the above construction. According to Lemma 4.2.2 we can choose a gauge around $p$ such that the covariant derivatives of the Dirac matrices vanish at $p$. Moreover, choosing normal coordinates and making a global (=constant) gauge transformation, we can arrange that $G(p) = \gamma^j$ and $\partial_j g_{kl}(p) = 0$. Then the covariant derivatives at $p$ reduce to partial derivatives, and we conclude that
\[
G^j(p) = \gamma^j, \quad \partial_k G^j(p) = 0.
\]
These equations resemble the conditions for normal coordinates (4.1.1), except that the role of the metric is now played by the Dirac matrices. Indeed, by differentiating (4.2.11) one sees that (4.2.27) implies (4.1.1). Therefore, (4.2.27) is a stronger condition which not only poses a condition for the coordinates, but also for the gauge. We call a coordinate system and gauge where (4.2.27) is satisfied a normal reference frame around $p$.

In a normal reference frame around $p$, the Dirac matrices, and via (4.2.11) also the metric, are the same as in Minkowski space up to the order $\sim (x-p)^2$. Combining the equivalence principle with the usual minimal coupling procedure in physics, it seems a sensible physical assumption that the Dirac equation at $p$ should coincide with that in Minkowski space. This implies that there should be a normal gauge such that all gauge potentials vanish at $p$, and thus the Dirac operator at $p$ should coincide with the vacuum Dirac operator $i\partial$. This physical argument makes it possible to specify the zero order term in (4.2.3).

**Definition 4.2.6.** A Dirac-type operator $D$ is called Dirac operator if for any $p \in \mathcal{M}$ there is a normal reference frame around $p$ such that $B(p) = 0$.

Equivalently, the Dirac operator could be defined as a differential operator of first order (4.2.3) with the additional structure that for any $p \in \mathcal{M}$ there is a coordinate chart and gauge such that the following three conditions are satisfied,
\[
G^j(p) = \gamma^j, \quad \partial_k G^j(p) = 0, \quad B(p) = 0.
\]
This alternative definition has the disadvantage that it is a-priori not clear whether the second condition $\partial_k G^j(p) = 0$ can be satisfied for a general metric. This is the reason why we preferred to begin with only the first condition (Definition 4.2.1), then showed that the second condition can be arranged by choosing suitable coordinates and gauge, and satisfied the third condition at the end (Definition 4.2.6).

In general coordinates and gauge, the Dirac operator can be written as
\[
D = iG^j D_j = iG^j \left( \partial_j - iE_j - ia_j \right),
\]
where $D$ is the spin connection of Definition 4.2.3. The matrices $E_j$ take into account the gravitational field and are called spin coefficients, whereas the $a_j$ can be identified with the electromagnetic potential (compare (1.3.6)). We point out that the gravitational field cannot be introduced into the Dirac equation by the simple replacement rule $\partial \rightarrow D$, because gravity has an effect on both the Dirac matrices and the spin coefficients. But factorizing the gauge group as $\text{U}(2,2) = \text{U}(1) \times \text{SU}(2,2)$, the $\text{SU}(2,2)$-gauge transformations are linked to the gravitational field because they influence $G^j$ and $E_j$, whereas the $\text{U}(1)$ can be identified with the gauge group of electrodynamics. In this sense, we
4.3. COMPUTATION OF THE DIRAC OPERATOR

obtain a unified description of electrodynamics and general relativity as a U(2, 2) gauge theory. The Dirac equation

\[(\mathcal{D} - m) \psi = 0\]  

(4.2.29)
describes a Dirac particle in the presence of a gravitational and electromagnetic field. According to Theorem 4.2.5, the curvature of the spin connection involves both the Riemann tensor and the electromagnetic field tensor. One can express the classical action in terms of these tensor fields, so that the corresponding Euler-Lagrange equations give rise to the classical Einstein-Dirac-Maxwell equations.

For the probabilistic interpretation of the Dirac equation in curved spacetime, we choose a spacelike hypersurface \(\mathcal{N}\) (corresponding to “space” for an observer) and consider in generalization of (1.3.12) on solutions of the Dirac equation the scalar product

\[
(\psi | \phi) = \hat{N} \prec \psi | G^j \nu_j \phi \succ d\mu_N,
\]

(4.2.30)
where \(\nu\) is the future-directed normal on \(\mathcal{N}\) and \(d\mu_N\) is the invariant measure on the Riemannian manifold \(\mathcal{N}\). Then \((\psi | \psi)\) is the normalization integral, which we again normalize to one. Its integrand has the interpretation as the probability density. In analogy to (1.3.9) the Dirac current is introduced by \(J^k = \prec \psi | G^k \psi \succ\). Using Theorem 4.2.4 one sees similar as in Minkowski space that the Dirac current is divergence-free, \(\nabla_k J^k = 0\). From the Gauß divergence theorem one obtains that the scalar product (4.2.30) does not depend on the choice of the hypersurface \(\mathcal{N}\).

In analogy to (1.3.16), we can introduce the inner product

\[
\langle \psi | \phi \rangle := \int_{\mathcal{M}} \langle \psi | \phi \rangle_x \, d\mu_M.
\]

(4.2.31)
in which the wave functions (which need not satisfy the Dirac equation but must have a suitable decay at infinity) are integrated over all of spacetime. We finally remark that, using Theorem 4.2.4 together with Gauß’ divergence theorem, one easily verifies that the Dirac operator is symmetric with respect to this inner product.

### 4.3. Computation of the Dirac Operator

We now explain how the Dirac operator can be computed in an efficient way in a given spacetime. Thus suppose that the Lorentzian metric \(g_{ij}\) is given in a chosen chart. The general procedure is to first choose matrices \(G^j(x)\) which are symmetric w.r.t. the spin scalar product (4.2.12) (where in our formulation, the spin scalar product is always given by (4.2.1)) and which satisfy the anti-commutation relations (4.2.11). Then the spin coefficients as given by (4.2.24) are obtained by a straightforward computation. Then the spin derivative is given by (4.2.23) (where \(a_j\) are the components of the electromagnetic potential; they are set to zero if no electromagnetic field is present). The Dirac operator is given by (4.2.28), i.e.

\[
\mathcal{D} = iG^j \partial_j = iG^j \partial_j + G^j E_j + G^j a_j.
\]

(4.3.1)
In this construction, one has a lot of freedom to choose the Dirac matrices \(G^j(x)\) (as described systematically by the U(2, 2)-gauge transformations (4.2.7) and (4.2.10)). It is a promising strategy to use this gauge freedom such as to choose Dirac matrices for which the formulas for the spin coefficients (4.2.24) become as simple as possible. Moreover, one should keep in mind that for the computation of the Dirac operator, one does not need to know all the matrices \(E_j\), but it suffices to compute the combination \(G^j E_j\) in (4.3.1).
Indeed, in many spacetimes of physical interest, making use of the gauge freedom, the combination $G^j E_j$ can be computed easily (for details see the computations in black hole geometries in \[80, 63\] or various examples in \[58, Section 9\]). We here illustrate this method by the example of a diagonal metric, in which case it is even unnecessary to compute the Christoffel symbols:

**Proposition 4.3.1.** Assume that there is a local chart $(x^i)$ in which the metric is diagonal, i.e.

$$ds^2 = \sum_{i=0}^{3} g_{ii}(x) \, dx_i^2 .$$  \hfill (4.3.2)

Then there is a gauge in which the Dirac operator (without electromagnetic field) takes the form

$$\mathcal{D} = i G^j \left( \frac{\partial}{\partial x^j} \right) + B ,$$  \hfill (4.3.3)

where

$$G^j(x) = g_{jj}(x)^{-\frac{1}{2}} \gamma^j$$  \hfill (4.3.4)

$$B(x) = \frac{i}{2\sqrt{|\det g|}} \frac{\partial_j}{\partial_j} \left( \sqrt{\det g} \right) G^j .$$  \hfill (4.3.5)

(Here $\gamma^j$ are again the Dirac matrices in Minkowski space).

**Proof.** With (4.3.4) we have satisfied the anti-commutation relations

$$\{G^j, G^k\} = 2 \, g^{jk} \, \mathbb{1} .$$

Moreover, the choice (4.3.4) ensures that the pseudo-scalar operator is constant, and that all derivatives of the $G^j$ are in the span of $\gamma^0, \ldots, \sigma^k$. Therefore, the formula for the zero order term in the Dirac operator (4.3.1) simplifies to

$$B = -\frac{i}{16} \text{Tr} \left( G^m \left( \nabla_j G_n \right) \right) G^j G^m G^n ,$$  \hfill (4.3.6)

where $\nabla_j G_n \equiv \partial_j G_n - \Gamma^k_{jn} G_k$ is the covariant derivative acting on the components of the spinorial matrix. Using the algebra of the Dirac matrices, one finds that (4.3.6) has a vectorial component (obtained by using the anti-commutation relations), and an axial component which is totally antisymmetric in the indices $j, m, n$. This totally antisymmetric term vanishes for the following reasons: First, since the Levi-Civita connection is torsion-free, we may replace the covariant derivative by a partial derivative. Second, it follows from (4.3.4) that the matrix $\partial_j G_n$ is a multiple of $G_n$, implying that the trace $\text{Tr}(G_m(\partial_j G_n))$ is symmetric in the indices $m$ and $n$.

It remains to compute the vectorial component of (4.3.6). A short computation yields

$$B = \frac{i}{2} \nabla_j G^j ,$$

and the usual formula for the covariant divergence of a vector field gives the result. $\Box$
4.4. Formulation with Vector Bundles, the Spinor Bundle

So far, the Dirac operator was introduced in a local chart. We intentionally left a large local gauge freedom, having the advantage that this freedom can be used to simplify the form of the Dirac operator. The remaining question is whether our constructions in local charts can be made global to obtain a Dirac operator \( D \) acting on the sections of the so-called spinor bundle \( S_M \). To this end, we shall consider the Dirac operator in different charts and path the Dirac operators in the overlapping regions.

In preparation, we recall the structures introduced so far, using a more notation which clarifies the dependence on gauge and coordinates. In our local construction at the beginning of Section 4.2, the spinor space at a point \( x \in M \) is simply \( \mathbb{C}^4 \) with the inner product (4.2.1). Using the same notation as in Section 1.4 in Minkowski space, we now denote the spinor space by \( (S_x M, \langle \cdot, \cdot \rangle_x) \). Moreover, we denote the linear operators on \( S_x M \) which are symmetric with respect to the spin scalar product by \( \text{Symm}(S_x M) \).

It is a 16-dimensional real vector space spanned by the operators in (4.2.14). Given a Dirac-type operator \( D \), the Dirac matrices \( G^j(x) \) span a four-dimensional subspace \( K_x \) of \( \text{Symm}(S_x M) \),

\[
K_x := \text{span}\{G^0(x), \ldots, G^3(x)\} \subset \text{Symm}(S_x M),
\]

referred to as a Clifford subspace at \( x \). Contracting a tangent vector \( u \) with the Dirac matrices gives rise to a mapping

\[
\gamma : T_x M \to K_x, \quad u \mapsto u^j G_j.
\]

Multiplying a spinor by \( \gamma(u) \) is referred to as Clifford multiplication. The anti-commutation relations (4.2.11) can be written as

\[
\frac{1}{2} \{\gamma(u), \gamma(v)\} = g_x(u, v) \mathbf{1}_{S_x M},
\]

showing that Clifford multiplication encodes the Lorentzian metric.

In view of the transformation law (4.2.4), the Clifford subspace does not depend on the choice of coordinates. But it clearly depends on the gauge. Indeed, in view of (4.2.5), it transforms according to

\[
K_x \to U K_x U^{-1} \quad \text{with} \; U \in U(S_x).
\]

In order to simplify our problem, it is a good idea to arrange by a gauge transformation that the Clifford subspace agrees at every spacetime point with the standard Clifford subspace:

**Lemma 4.4.1.** By a gauge transformation (4.3.2) we can arrange that

\[
K_x = \text{span}\{\gamma^0, \ldots, \gamma^3\}
\]

(where \( \gamma^j \) are again the Dirac matrices in the Dirac representation).

**Proof.** We consider the pseudo-scalar operator \( \Gamma(x) \) as defined by (4.2.13). By direct computation, one verifies that it satisfies the same relations as in Minkowski space

\[
\Gamma(x)^* = -\Gamma(x) \quad \text{and} \quad \Gamma(x)^2 = 1.
\]

The first relation implies that \( \Gamma(x) \) maps positive definite spinors to negative definite spinors and vice versa. Therefore, there is a pseudo-orthonormal basis of the spinor
space in which \( \Gamma(x) \) takes the same form as in Minkowski space,
\[
\Gamma(x) = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}.
\] (4.4.3)

Rewriting the change of basis as a gauge transformation, we have arranged by a transformation of the form (4.4.2) that the pseudo-scalar operator has the same form as in Minkowski space.

It follows from (4.2.13) and the anti-commutation relations that every vector in \( K \) anti-commutes with \( \Gamma \). Therefore,
\[
K \subset \text{span}\{\gamma^0, \ldots, \gamma^3, \gamma^0 \Gamma, \ldots, \gamma^3 \Gamma\}.
\]

We next show that the vector space \( K \cap \text{span}\{\gamma^0, \gamma^0 \Gamma\} \) is one-dimensional. To this end, let \( u, v \in T_x\mathcal{M} \) with \( \gamma(u) = (a + b\Gamma) \gamma^0 \) and \( \gamma(v) = (c + d\Gamma) \gamma^0 \) with real coefficients \( a, b, c, d \). Then their anti-commutator is computed by
\[
\{\gamma(u), \gamma(v)\} = 2(ac - bd) \mathbb{1} + 2(bc - ad) \Gamma,
\]
implying that \( bc - ad = 0 \). This implies that \( \gamma(u) \) and \( \gamma(v) \) are linearly dependent, giving the claim.

Repeating the last argument for \( \gamma^1, \ldots, \gamma^3 \), we conclude that there is a (not necessarily pseudo-orthonormal) basis \( u_0, \ldots, u_3 \) of \( T_x\mathcal{M} \) such that
\[
\gamma(u_j) = (a_j + b_j \Gamma) \gamma_j.
\]
Then for any \( j \neq k \),
\[
\{\gamma(u_j), \gamma(u_k)\} = (b_j a_k - a_j b_k) \Gamma [\gamma_j, \gamma_k],
\]
implying that the four vectors \( (a_j, b_j) \in \mathbb{R}^2 \) with \( j = 0, \ldots, 3 \) are all linearly independent. Therefore, by rescaling the basis vectors \( u_j \) we can arrange that
\[
\gamma(u_j) = (a + b\Gamma) \gamma_j
\]
for real parameters \( a \) and \( b \).

The signature \((1, 3)\) of the Lorentzian metric implies that \(|a| > |b|\). Moreover, by flipping the sign of the vectors \( u_j \) if necessary we can arrange that \( a > 0 \). Therefore, we may represent \( K \) as
\[
K = \text{span}\{e^{\alpha \Gamma} \gamma^0, \ldots, e^{\alpha \Gamma} \gamma^3\}
\]
for some \( \alpha \in \mathbb{R} \) (note that \( e^{\alpha \Gamma} = \cosh \alpha + \Gamma \sinh \alpha \)). Performing the gauge transformation (4.4.2) with \( U \) according to
\[
U = \exp\left(-\frac{\alpha}{2} \Gamma\right)
\]
gives the result. \(\square\)

After these preparations, we are ready to enter the patching construction. Thus let \( (x, U) \) and \( (\tilde{x}, \tilde{U}) \) be two local charts on \( (\mathcal{M}, g) \) with non-empty overlap \( U \cap \tilde{U} \). For technical simplicity, we only consider the case that the spacetime is time-oriented. We choose the charts such that \( x^0 \) and \( \tilde{x}^0 \) are time functions which increase to the future. Then we can write the Dirac operator in each chart according to (4.2.28), where for clarity we denote the objects in the chart \( \tilde{x} \) with an additional tilde. We first consider the case without electromagnetic field where the potentials \( a_j \) vanish. According to Lemma (4.4.1)
there is no loss of generality to restrict attention to gauges where the Dirac matrices are linear combinations of the Dirac matrices in Minkowski space, i.e.,
\[ G^j(x) = h^j_k(x) \gamma^k \quad \text{and} \quad \tilde{G}^j(\tilde{x}) = \tilde{h}^j_k(\tilde{x}) \gamma^k. \]

Since \( x^0 \) is a time coordinate, the bilinear form \( \langle . | G^0(x) | . \rangle_x \) is definite at very spacetime point \( x \), and similar for the tilde coordinates. We choose the signs of the Dirac matrices such that the bilinear forms \( \langle . | G^0(x) | . \rangle_x \) and \( \langle . | \tilde{G}^0(\tilde{x}) | . \rangle_y \) are all positive definite. Moreover, as explained in Lemma 4.4.1, we choose the gauge such that the pseudo-scalar operator (4.2.13) has the same form as in Minkowski space (4.4.3). This implies that the spatial orientations of the two charts are compatible.

The transformation from the chart \((x, U)\) to \((\tilde{x}, \tilde{U})\) involves the coordinate transformation as described by (4.2.4). After this transformation, the Dirac matrices
\[ \tilde{G}^j(\tilde{x}) \quad \text{and} \quad G^k(x) \frac{\partial \tilde{x}^j}{\partial x^k} \]
will in general not coincide. But since the matrices are all formed as linear combinations of the Dirac matrices in Minkowski space, satisfy the same anti-commutation relations, and have the same time and spatial orientations, they can be obtained from each other by an orthochronous and proper Lorentz transformation, i.e.
\[ \tilde{G}^j(\tilde{x}) = \Lambda^j_l G^k(x) \frac{\partial \tilde{x}^l}{\partial x^k}. \]

Now we can proceed just as in the proof of Lorentz invariance of the Dirac equation in Minkowski space (see Lemma 1.3.1) to conclude that there is a unitary transformation \( U(x) \in \mathbb{U}(S_x) \) of the form
\[ U := \exp \left( \frac{1}{4} \lambda_{lk} \gamma^l \gamma^k \right) \quad (4.4.4) \]
(with an anti-symmetric tensor \( \lambda_{lk} \)) such that the Dirac matrices agree after the gauge transformation, i.e.
\[ U(\tilde{x}) \tilde{G}^j(\tilde{x}) U(\tilde{x})^{-1} = G^j(\tilde{x}) \frac{\partial \tilde{x}^l}{\partial x^k}. \]

Since the spin coefficients \( E_j \) in (4.2.23) are given explicitly in terms of the Dirac matrices and their derivatives (see (4.2.24)), the lower order terms in the resulting Dirac operators (4.2.28) also agree. Moreover, using that the only matrices which commute with all Dirac matrices are multiples of the identity, one sees that the gauge transformation \( U(\tilde{x}) \) of the form (4.4.4) is uniquely determined up to a sign. In this way, to every coordinate transformation, we have found a gauge transformation, unique up to a sign, such that the Dirac operators agree.

With the above construction, we have found a procedure for matching the Dirac operators in two overlapping charts. The involved gauge transformations of the form (4.4.4) are unique up to signs. Therefore, once we have decided on the signs, there is a unique way of identifying the Dirac wave functions in the overlapping region of two charts, such as to obtain Dirac wave functions in a larger patch of the manifold \( \mathcal{M} \). Proceeding inductively, one can hope to obtain Dirac wave functions on all of \( \mathcal{M} \). The subtle point is whether the signs of the transformations can be chosen in a compatible way for all charts. In more mathematical terms, one must satisfy the so-called cocycle conditions, and it turns out that these conditions can be fulfilled if and only if \( M \) satisfies a topological condition, the so-called spin condition (for details see for example [109 § II.1 and § II.2]). If the
spin condition is satisfied, one can identify the spinor spaces via the mappings which patch the charts together. In this way, one obtains a vector bundle over $\mathcal{M}$, referred to as the spinor bundle $S\mathcal{M}$. The fibers of the spinor bundle are the spinor spaces $S_x\mathcal{M}$, which are four-dimensional complex vector spaces endowed with an inner product $\langle \cdot , \cdot \rangle_x$ of signature $(2,2)$. The transformations of the form (4.4.4) generate a group, the so-called spin group denoted by

$$\text{Spin}_x^+ \subset \text{U}(S_x)$$

(4.4.5)

(the reason why we write “generated by” is that the operators of the form (4.4.4) do not form a group; see Exercise 4.2). Elements of the spin group act on vectors of the Clifford subspace by the adjoint representation,

$$\gamma(v) \rightarrow U \gamma(v) U^{-1},$$

we obtain another vector of the Clifford subspace, i.e.

$$U \gamma(v) U^{-1} = \gamma(\mathcal{O}_U(u)).$$

(4.4.6)

Since the anti-commutation relations remain unchanged, the resulting transformation of the tangent space is an isometry. Indeed, by Lemma 1.3.1 it is a proper orthochronal Lorentz transformation,

$$\mathcal{O}_U \in \text{SO}^+(T_x\mathcal{M}).$$

(4.4.7)

The indices $\uparrow$ in (4.4.5) and (4.4.7) indicate that we restrict attention to orthochronous transformations. We thus obtain the usual commutative diagram

$$\mathbb{Z}_2 \longrightarrow \text{Spin}_x^+ \longrightarrow \text{SO}^+(T_x\mathcal{M}) \longrightarrow 0.$$ 

The connection to the usual spin group is obtained as follows. We say that a tangent vector $u \in T_x\mathcal{M}$ is a unit vector if $\langle u, u \rangle = \pm 1$. The spin group is defined by (see for example [9, 109], the concise summary in [5, Section 2] or similarly [85] in the Riemannian setting)

$$\text{Spin}_x := \{ \text{group generated by } \gamma(u) \gamma(v) \text{ with unit vectors } u, v \in T_x\mathcal{M} \}. \quad (4.4.8)$$

By expanding the exponential in (4.4.4), one sees that this matrix is generated by even products of Dirac matrices, showing that the group $\text{Spin}_x^+$ in (4.4.5) is a subgroup of $\text{Spin}_x$. The group $\text{Spin}_x$ also includes operators which are not unitary but satisfy instead the relation $U^*U = -1$. These transformations describe reversals of the time orientation. Working with the general spin group (4.4.8) is of advantage in general dimension or signature. In four-dimensional time-oriented and orientable spacetimes, however, we can just as well restrict attention to orthochronous proper Lorentz transformations and the gauge transformations in (4.4.5).

We finally mention how to treat an electromagnetic field. Then the starting point is a time-oriented Lorentzian spin manifold $(\mathcal{M}, g)$ together with an anti-symmetric two-tensor $F$ (the field tensor). In this situation, after the above coordinate and gauge transformations, the electromagnetic potentials $a_j$ and $\tilde{a}_j$ in the two charts will in general not coincide. But, since the field tensor is prescribed, they coincide after a local $U(1)$-gauge transformation. Identifying the spinor spaces after this gauge transformation defines the Dirac operator as acting on the spinor bundle $S\mathcal{M}$. The resulting effective gauge group is $\text{U}(1) \times \text{Spin}_x^+$. We point out that this effective gauge group is obtained under the condition that the Clifford subspace is fixed at each spacetime point according to Lemma 4.4.1. Dropping this condition gives rise to the larger local gauge group $\text{U}(2,2)$.
4.5. The Dirac Solution Space in Globally Hyperbolic Spacetimes

We now turn attention to solutions of the Dirac equation. In Minkowski space, a convenient method for constructing solutions is the Fourier transformation (see Section 24.1). However, this method can be used only for PDEs with constant coefficients, and therefore it does not apply to the Dirac equation in curved spacetime. Instead, a general method is to solve the Cauchy problem for given initial data, making it possible to identify the solution space of the Dirac equation with the space of suitable initial data. Clearly, for this method to be applicable, one needs to decompose spacetime into “space” and “time,” because otherwise it would not be clear how to prescribe initial data at some initial time. In order to describe the Dirac solutions in all of spacetime by initial data, this splitting of spacetime into space and time must be performed globally. This splitting will not be unique, as is already clear in Minkowski space because different reference frames give rise to different splittings. In curved spacetime, such a splitting does not even need to exist. For example, in spacetimes containing closed timelike curves, we cannot expect that the Cauchy problem is well-posed.

The necessary assumptions on spacetime needed for a sensible formulation of the Cauchy problem are subsumed in the mathematical notion of **global hyperbolicity**. We first give the formal definition and then explain its consequences. Let \((M, g)\) be a Lorentzian manifold. We assume that \(M\) is time-oriented. Then a parametrized piecewise \(C^1\)-curve \(\gamma(\tau)\) in \(M\) is said to be **causal** if its tangent vector \(\dot{\gamma}(\tau)\) is causal (i.e. timelike or lightlike) for all \(\tau\) where \(\gamma\) is differentiable. Moreover, it is **future-directed** and **past-directed** if its tangent vectors are future- and past-directed, respectively. The curve \(\gamma(\tau)\) is said to be **inextendible** if it cannot be extended as a piecewise \(C^1\)-curve. The manifold \(\mathcal{M}\) is said to satisfy the **strong causality condition** if there are no almost closed causal curves in the sense that for all \(x \in M\) and for each open neighborhood \(U\) of \(x\) there is an open neighborhood \(V \subset U\) of \(x\) such that every causal curve in \(M\) which is starting and ending in \(V\) is entirely contained in \(U\). Moreover, in straightforward generalization of the corresponding notions in Minkowski space as introduced after (1.2.2), we let \(J_\omega^\lor\) \((\text{and } J_\omega^\land)\) be the set of all points \(y \in M\) which can be joined from \(x\) by a future-directed (respectively past-directed) causal curve. The manifold \(\mathcal{M}\) is said to be **globally hyperbolic** if the strong causality condition holds and if the set \(J_\omega^\lor \cap J_\omega^\land\) is compact for all \(x, y \in \mathcal{M}\).

For more details on the abstract definitions we refer to [98, Section 6.6], [6, Section 1.3], [11, Section 3.2] or [115, Chapter 14]). A globally hyperbolic Lorentzian manifold \((\mathcal{M}, g)\) has remarkable properties, as we now explain. First, global hyperbolicity implies that \((\mathcal{M}, g)\) is diffeomorphic to a product,

\[
\mathcal{M} \simeq \mathbb{R} \times \mathcal{N},
\]

where \(\simeq\) means that there is a smooth diffeomorphism from \(\mathcal{M}\) to \(\mathbb{R} \times \mathcal{N}\). Thus every point \(p \in \mathcal{M}\) can be written as \(p = (t, x)\) with \(t \in \mathbb{R}\) and \(x \in \mathcal{N}\). One also refers to the above property that \(\mathcal{M}\) admits a **smooth foliation** \(\mathcal{M} = (\mathcal{N}_t)_{t \in \mathbb{R}}\), where \(\mathcal{N}_t := \{t\} \times \mathcal{N}\). Moreover, the foliation can be chosen such as to have the following properties:

(i) **Every surface** \(\mathcal{N}_t\) is spacelike (meaning that the metric induced by \(g\) on \(\mathcal{N}_t\) is negative definite).

(ii) **Every surface** \(\mathcal{N}_t\) is a Cauchy surface, meaning that every inextendible timelike curve in \(\mathcal{M}\) intersects \(\mathcal{N}_t\) exactly once.
The function $t$ is also referred to as a global time function. These above properties of globally hyperbolic manifolds were proven in [12] (for more details and more references see again [6, Section 1.3]).

The property of $N_t$ of being a Cauchy surface implies that the Cauchy problem for the Dirac equation is well-posed, as we now explain. To this end, let $(\mathcal{M}, g)$ be a four-dimensional globally hyperbolic spacetime. Then the topological splitting (4.5.1) implies that the spin condition mentioned before (4.4.5) is satisfied. Therefore, there is a spinor bundle $(\mathcal{S}M, \prec, |, \succ)$, being a vector bundle with fibers $S_x\mathcal{M} \cong \mathbb{C}^4$ (there may be different spin structures, but we shall not go into this here). Moreover, the Dirac operator $D$ is well-defined; in local coordinates and local spinor bases it takes the form (4.2.28). In the Cauchy problem, one seeks for solutions of the Dirac equation of mass $m$ for prescribed initial data at time $t_0$, i.e.

$$(D - \imath m)\psi = \phi \quad \text{with} \quad \psi|_{N_t} = \psi_0.$$  (4.5.2)

The following result holds:

**Theorem 4.5.1.** For smooth initial data $\psi_0 \in C^\infty(N_{t_0}, \mathcal{S}M)$ and a smooth inhomogeneity $\phi \in C^\infty(\mathcal{M}, \mathcal{S}M)$ the Cauchy problem (4.5.2) has a unique global solution $\psi \in C^\infty(\mathcal{M}, \mathcal{S}M)$.

The proof of this theorem uses methods of hyperbolic partial differential equations and will be given in Section 13.6 later in this book.

Having a Cauchy surface is also very useful because we can then a scalar product on the solution space as the spatial integral (4.2.30), where $N$ is chosen as a Cauchy surface. However, for the integral in (4.2.30) to be well-defined, working with smooth solutions is not suitable. Instead, similar as explained in Minkowski space in Section 1.4, we better assume that the solution has compact support on the Cauchy surface. Due to finite propagation speed for solutions of hyperbolic partial differential equations (as will be made precise in Section 13.2), the following result holds:

**Theorem 4.5.2.** If the initial data and the inhomogeneity have compact support,

$$\psi_0 \in C^\infty_0(N_{t_0}, \mathcal{S}M) \quad \text{and} \quad \phi \in C^\infty_0(\mathcal{M}, \mathcal{S}M),$$

then the solution $\psi$ of the Cauchy problem (4.5.2) also has compact support on any other Cauchy surface $N_t$.

The proof of this theorem will again be given in Section 13.6 below.

Using the same notion as in Section 1.4 we refer to smooth solutions as in the above theorem as having spatially compact support. Smooth and spatially compact sections of the spinor bundle are again denoted by $C^\infty_0(\mathcal{M}, \mathcal{S}M)$. For Dirac solutions in this class, the scalar product (4.2.30) is well-defined. Moreover, due to current conservation, this scalar product does not depend on the choice of the Cauchy surface (as explained after (4.2.30)). Exactly as explained in Section 1.4, taking the completion gives the Hilbert space $\mathcal{H}_m, (\cdot, \cdot))$ of weak solutions of the Dirac equation with the property that their restriction to any Cauchy surface is square integrable (where “square integrable” is defined via (4.2.30)).

### 4.6. Hamiltonian Formulation in Stationary Spacetimes

Given a foliation $(N_t)_{t \in \mathbb{R}}$ by Cauchy surfaces of the globally hyperbolic spacetime $(\mathcal{M}, g)$, one can rewrite the Dirac equation in the Hamiltonian form,

$$i\partial_t \psi = H \psi$$  (4.6.1)
with a Hamiltonian $H$. In order to compute $H$ in a local chart, one chooses a coordinate system $(x^i)$ such that $x^0 = t$ coincides with the time function. Then, writing the Dirac operator in (4.2.29) in the form (4.2.28) and solving for the time derivatives, one obtains in generalization of (1.3.15)

$$H = -(G^0)^{-1} \left( \sum_{\alpha=1}^{3} iG^\alpha \left( \partial_\alpha - iE_\alpha - ia_\alpha \right) - m \right) - E_0 - a_0.$$ 

When analyzing the Dirac equation in the Hamiltonian form, one must be careful because the Hamiltonian in general is not symmetric with respect to the Hilbert space scalar product (4.2.30). This can be seen as follows. For the Dirac equation in Minkowski space, the symmetry of the Hamiltonian is obtained just as for the Schrödinger equation by using that the scalar product is conserved in time (1.1.6). In curved spacetime, the scalar product is still conserved (due to current conservation). But when taking the time derivative, one must take into account that the scalar product itself is time-dependent. More precisely, assuming for notational simplicity that the Cauchy surfaces admit global charts,

$$0 = \partial_t \langle \phi \mid \psi \rangle = \frac{\partial}{\partial t} \int_{\mathcal{N}_t} <\phi \mid G^j \nu_j \psi> \, d\mu_{\mathcal{N}_t}$$

$$= (\partial_t \phi \mid \psi) + (\phi \mid \partial_t \psi) + \int_{\mathbb{R}^3} <\psi \mid \left( \partial_t (G^j \nu_j \sqrt{\det g_{\mathcal{N}_t}}) \right) \phi> \, dx$$

$$= -i \left( \langle H \phi \mid \psi \rangle - \langle \phi \mid H \psi \rangle \right) + \int_{\mathbb{R}^3} <\psi \mid \left( \partial_t (G^j \nu_j \sqrt{\det g_{\mathcal{N}_t}}) \right) \phi> \, dx$$

(4.6.2)

(where $g_{\mathcal{N}_t}$ denotes the induced Riemannian metric on the Cauchy surface $\mathcal{N}_t$). The integral in the last line is in general non-zero. In this case, the Hamiltonian is obviously not symmetric. At first sight, this might seem surprising because it seems to contradict the axioms of quantum mechanics (for a detailed account on this issue see [2, 3]). However, one should keep in mind that the non-symmetric contributions to the Hamiltonian are needed in order to compensate for the fact that the scalar product itself is time-dependent.

Our interpretation of the above problem is that the Hamiltonian formulation of the Dirac equation is useful only in situations when the integral in (4.6.2) vanishes. This can be arranged if all the coefficients of the metric are time-independent. In other words, spacetime should be stationary with corresponding Killing field given by $\partial_t$. Under these assumptions, the Hamiltonian $H$ is also time-independent. Moreover, the computation (4.6.2) shows that the operator is symmetric. Using that the time evolution maps smooth and compactly supported initial data on the Cauchy surface at time $t_0$ to a smooth and compactly supported solution at an arbitrary time $t$, one can use abstract methods to construct a selfadjoint extension of $H$ (see for example [78] for a general situation involving additional boundary conditions). Then the Cauchy problem can be solved immediately using the spectral theorem for selfadjoint operators,

$$\psi(t, x) = (e^{-itH} \psi_0)(x).$$

This formulation is particularly useful for analyzing the long-time behavior of the solutions (see for example the analysis in the Kerr geometry in [65, 64]).
Exercises

Exercise 4.1. Verify by elementary integration by parts in a chart that for a diagonal metric \((4.3.2)\), the Dirac operator \((4.3.3)\) is symmetric with respect to the inner product \((4.2.31)\).

Exercise 4.2. The goal of this exercise is to show that the unitary operators of the form \((4.4.4)\) do not form a group (in more mathematical language, the spin group is not exponential; for details see \([29]\) and the references therein). We proceed in several steps:

(a) Let \(\lambda_{jk}\) be an anti-symmetric tensor. Show using the anti-commutation relations that
\[
\left(\frac{1}{4} \lambda_{lk} \gamma^l \gamma^k\right)^2 = -\frac{1}{16} \lambda_{lk} \lambda^{lk} 1 + \frac{i}{16 \cdot 4!} \Gamma \varepsilon^{ijkl} \lambda_{ij} \lambda_{kl}.
\]
(b) Deduce from (a) that the corresponding unitary transformation \((4.4.4)\) is a linear combination of the matrices
\[
1, \quad \Gamma, \quad \lambda_{lk} \gamma^l \gamma^k \quad \text{and} \quad \lambda_{lk} \Gamma \gamma^l \gamma^k.
\]
(c) Show under the additional assumption \(\varepsilon^{ijkl} \lambda_{ij} \lambda_{kl} = 0\) that
\[
\exp \left(\frac{1}{4} \lambda_{lk} \gamma^l \gamma^k\right) = \begin{cases} 
1 \cos \alpha + \frac{1}{4} \lambda_{lk} \gamma^l \gamma^k \frac{\sin \alpha}{\alpha} & \text{if } \lambda_{lk} \lambda^{lk} > 0 \\
1 \cosh \alpha + \frac{1}{4} \lambda_{lk} \gamma^l \gamma^k \frac{\sinh \alpha}{\alpha} & \text{if } \lambda_{lk} \lambda^{lk} < 0,
\end{cases}
\]
where \(\alpha := \sqrt{|\lambda_{lk} \lambda^{lk}|}/4\). \textit{Hint:} Use \((4.6.3)\) in the power series of the exponential.

(d) Choose a specific tensor \(\lambda_{lk}\) for which the matrix in \((4.6.4)\) is equal to minus the identity.

(e) Show that the last expression involves no contribution \(\sim \Gamma\) only if either \(\alpha = 0\) or \(\beta = 0\) (\textit{Hint:} It might be convenient to work in an eigenvector basis of \(i\Gamma\)). Infer that in the case \(\alpha = 0\), this expression is a linear combination of the matrices \(1\) and \(\gamma^0\gamma^1\). Conclude that
\[
\frac{1}{4} \lambda_{lk} \gamma^l \gamma^k = (\alpha + i\Gamma \beta) \gamma^0 \gamma^1.
\]
(f) Deduce from (c) and (e) that the matrices of the form \((4.4.4)\) do not form a group.
Part 2

Causal Fermion Systems: Fundamental Structures
A Brief Introduction to Causal Fermion Systems

In this chapter we introduce and explain the basic objects and structures of a causal fermion system. Starting from a simple example (Section 5.1), we explain how to get to the basic objects of a causal fermion system (Section 5.2). After these preparations, the general definition of a causal fermion system will be given (Section 5.3). Next, as a further example, we explain how the Minkowski vacuum can be described by a causal fermion system (Section 5.4). In order to formulate equations describing the dynamics of a causal fermion system, we introducing a variational principle, the so-called causal action principle (Section 5.5). We conclude by discussing the form of the causal action principle (Section 5.7) and by explaining the underlying physical concepts (Section 5.8).

5.1. Motivating Example: A Two-Dimensional Lattice System

In order to motivate causal fermion systems, we begin with the familiar example of a cubic spacetime lattice. For simplicity, we consider a two-dimensional lattice (one space and one time dimension), but higher-dimensional lattices could be described similarly. Thus let $\mathcal{M} \subset \mathbb{R}^{1,1}$ be a cubic lattice in two-dimensional Minkowski space. We denote the spacing in time direction by $\Delta t$ and in spatial direction by $\Delta x$ (see Figure 5.1). The usual procedure for setting up equations on a lattice is to replace derivatives by difference quotients, giving rise to an evolution equation which can be solved time step by time step according to deterministic rules.

As a concrete example, let us consider a discretization of the two-dimensional wave equation on the lattice,

$$
0 = \Box \phi(t,x) := \frac{1}{(\Delta t)^2} \left( \phi(t + \Delta t, x) - 2\phi(t, x) + \phi(t - \Delta t, x) \right) - \frac{1}{(\Delta x)^2} \left( \phi(t, x + \Delta x) - 2\phi(t, x) + \phi(t, x - \Delta x) \right) .
$$

(5.1.1)

Solving this equation for $\phi(t + \Delta t, x)$ gives a rule for computing $\phi(t + \Delta t, x)$ from the values of $\phi$ at earlier times $t$ and $t - \Delta t$ (see again Figure 5.1).

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig5_1.png}
\caption{Time evolution of a lattice system $\mathcal{M} \subset \mathbb{R}^{1,1}$.}
\end{figure}
While this method is very simple and gives well-defined evolution equations, it also has several drawbacks:

- The above method of discretizing the continuum equations is very *ad hoc*. Why do we choose a regular lattice, why do we work with difference quotients? There are many other ways of discretizing the wave equation.
- The method is *not background-free*. In order to speak of the “lattice spacing,” the lattice must be thought of as being embedded in a two-dimensional ambient space-time.
- The concept of a spacetime lattice is not invariant under general coordinate transformations. In other words, the assumption of a spacetime lattice is *not compatible with the equivalence principle*.

In view of these shortcomings, the following basic question arises:

Can one formulate equations without referring to the nearest neighbor relation and the lattice spacing?

The answer to this question is yes, and we will now see how this can be done in the example of our two-dimensional lattice system. Although our example is somewhat oversimplified, this consideration will lead us quite naturally to the setting of causal fermion systems.

We next consider complex-valued wave functions \( \psi_1, \ldots, \psi_f \) on the lattice \( \mathcal{M} \) (for simplicity a finite number of them, i.e. \( f < \infty \)). A-priori, these wave functions are not assumed to satisfy any wave equation. On the complex vector space spanned by these wave functions we introduce a scalar product \( \langle \cdot | \cdot \rangle_{\mathcal{H}} \) and assume that the wave functions are orthonormal, i.e.

\[
\langle \psi_k | \psi_l \rangle_{\mathcal{H}} = \delta_{kl} .
\]

We thus obtain an \( f \)-dimensional Hilbert space \( (\mathcal{H}, \langle \cdot | \cdot \rangle_{\mathcal{H}}) \). Note that the scalar product is given abstractly (meaning that it has no representation in terms of the wave functions as a sum over lattice points). Next, for any lattice point \((t, x) \in \mathcal{M}\) we introduce the so-called *local correlation operator* \( F(t, x) : \mathcal{H} \to \mathcal{H} \) as a matrix with elements given by

\[
(F(t, x))_{jk} = \psi_j(t, x) \overline{\psi_k(t, x)} .
\]

The diagonal elements of this matrix are the absolute squares \( |\psi_k(t, x)|^2 \) of the corresponding wave functions. The off-diagonal elements, on the other hand, tell us about the correlation of the \( j \)th and \( k \)th wave function at the lattice point \((t, x)\). This is the reason for the name “local correlation operator.” This operator can also be characterized in a basis-invariant way by the relations

\[
\langle \psi, F(t, x) \phi \rangle_{\mathcal{H}} = \psi(t, x) \overline{\phi(t, x)} ,
\]

to be satisfied for all \( \psi, \phi \in \mathcal{H} \). Taking the complex conjugate, one sees immediately that the matrix defined by (5.1.3) is Hermitian. Stated equivalently independent of bases, the local correlation operator is a *symmetric* linear operator on \( \mathcal{H} \). Moreover, a local correlation operator has *rank at most one* and is *positive semi-definite*. This can be seen by writing it as

\[
F(t, x) = e^*e \quad \text{with} \quad e : \mathcal{H} \to \mathbb{C} , \quad \psi \mapsto \psi(x) .
\]

It is useful to denote the set of all operators with the above properties by \( \mathcal{F} \):

\[
\mathcal{F} := \{ F \in \text{L}(\mathcal{H}) \mid F \text{ is symmetric, positive semi-definite and has rank at most one} \} .
\]

Varying the lattice point, we obtain a mapping (see Figure 5.2)

\[
F : \mathcal{M} \to \mathcal{F} , \quad (t, x) \mapsto F(t, x) .
\]
For clarity, we note that the set $\mathcal{F}$ is not a vector space, because the linear combination of operators in $\mathcal{F}$ will in general have rank bigger than one. But it is a conical set in the sense that a positive multiple of any operator in $\mathcal{F}$ is again in $\mathcal{F}$ (this is why in Figure 5.2 the set $\mathcal{F}$ is depicted as a cone).

We point out that the local correlation operators do not involve the lattice spacing or the nearest neighbor relation; instead they contain information only on the local correlations of the wave functions at each lattice point. With this in mind, our strategy for formulating equations which do not involve the structures of the lattice is to work exclusively with the local correlation operators, i.e. with the subset $F(\mathcal{M}) \subset \mathcal{F}$. In other words, in Figure 5.2 we want to disregard the lattice on the left and work only with the objects on the right.

How can one set up equations purely in terms of the local correlation operators? In order to explain the general procedure, we consider a finite number of operators $F_1, \ldots, F_L$ of $\mathcal{F}$. Each of these operators can be thought of as giving information on the local correlations of the wave functions at a corresponding spacetime point. However, this “spacetime point” is no longer a lattice point, but at the moment it is merely a point without additional structure. In order to obtain a “spacetime” in the usual sense, one needs additional structures and relations between the spacetime points. Such relations can be obtained by multiplying the operators. Indeed, the operator product $F_i F_j$ tells us about correlations of the wave functions at different spacetime points. Taking the trace of this operator product gives a real number. Our method for formulating physical equations is to set up a variational principle. This variational formulation has the advantage that symmetries give rise to conservation laws by Noether’s theorem (as will be explained in Chapter 9). Therefore, we want to minimize an action $\mathcal{S}$. A simple example is to

\[
\text{minimize} \quad \mathcal{S}(F_1, \ldots, F_n) := \sum_{i,j=1}^{L} \text{Tr}(F_i F_j)^2
\]

under variations of the points $F_1, \ldots, F_L \in \mathcal{F}$. In order to obtain a mathematically sensible variational principle, one needs to impose certain constraints. Here we do not enter the details, because the present example is a bit too simple. Instead, we merely use it as a motivation for the general setting of causal fermion systems, which we now introduce.

5.2. Towards the General Definition

In order to get from our example to the general setting of causal fermion systems, we extend the above constructions in several steps:
(a) The previous example works similarly in higher dimensions, in particular for a lattice $\mathcal{M} \subset \mathbb{R}^{1,3}$ in four-dimensional Minkowski space. This has no effect on the resulting structure of a finite number of distinguished operators $F_1, \ldots, F_L \in \mathcal{F}$.

(b) Suppose that on the lattice we consider multi-component wave functions $\psi(t, x) \in \mathbb{C}^N$. Then the pointwise product on the right side of (5.1.3) must be replaced by a complex inner product, which we denote by $\langle \cdot | \cdot \rangle$ (in mathematical terms, this inner product is a non-degenerate sesquilinear form; we always use the convention that the wave function in the first argument is complex conjugated). Accordingly, the definition of the local correlation operator (5.1.3) is to be replaced by

$$ (F(t, x))^i_j = -\langle \psi_j(t, x) | \psi_k(t, x) \rangle $$

(the minus sign compared to (5.1.3) merely is a useful convention). The resulting local correlation operator is no longer an operator of rank at most one, but it has rank at most $N$ (as can be seen for example by writing it similar to (5.1.4) in the form $F(t, x) = -e^* e$ with $e : \mathcal{H} \to \mathbb{C}^N$). If the inner product $\langle \cdot | \cdot \rangle$ on $\mathbb{C}^N$ is positive definite, then the operator $F(t, x)$ is negative semi-definite. However, in the physical applications in mind, this inner product will not be positive definite. Indeed, a typical example in mind is that of four-component Dirac spinors. The Lorentz invariant inner product $\bar{\psi}\phi$ on Dirac spinors in Minkowski space (with the usual adjoint spinor $\bar{\psi} := \psi^\dagger \gamma^0$) is indefinite of signature $(2, 2)$. In order to describe systems involving leptons and quarks, one must take direct sums of Dirac spinors, giving the signature $(n, n)$ with $n \in 2\mathbb{N}$. With this in mind, we assume more generally that

$$ \langle \cdot | \cdot \rangle \text{ has signature } (n, n) \text{ with } n \in \mathbb{N}. $$

Then the resulting local correlation operators are symmetric operators of rank at most $2n$, which (counting multiplicities) have at most $n$ positive and at most $n$ negative eigenvalues.

(c) Finally, it is useful to generalize the setting such as to allow for continuous spacetimes and for spacetimes which may have both continuous and discrete components. In preparation, we note that the sums over the operators $F_1, \ldots, F_L$ in (5.1.5) can be written as integrals,

$$ S(\rho) = \int_\mathcal{F} d\rho(x) \int_\mathcal{F} d\rho(y) \text{ Tr}(F_i F_j)^2, \quad (5.2.1) $$

if $\rho$ is a measure on $\mathcal{F}$ chosen as the sum of Dirac measures supported at these operators,

$$ \rho = \sum_{i=1}^L \delta_{F_i}. \quad (5.2.2) $$

In this formulation, the measure plays a double role: First, it distinguishes the points $F_1, \ldots, F_L$ as those points where the measure is non-zero, as is made mathematically precise by the notion of the support of the measure, i.e.

$$ \text{supp } \rho = \{F_1, \ldots, F_L\}. \quad (5.2.3) $$

Second, a measure makes it possible to integrate over its support, an operation which for the measure (5.2.2) reduces to the sum over $F_1, \ldots, F_L$.

Now one can extend the setting simply by considering (5.2.1) for more general measures on $\mathcal{F}$ (like for example regular Borel measures). The main advantage
of working with measures is that we get into a mathematical framework in which variational principles like (5.1.5) can be studied with powerful analytic methods.

5.3. General Definition of a Causal Fermion System

Motivated by the previous considerations we now give the basic definition of a causal fermion system.

**Definition 5.3.1. (causal fermion system)** Given a separable complex Hilbert space \( \mathcal{H} \) with scalar product \( \langle .,| . \rangle_{\mathcal{H}} \) and a parameter \( n \in \mathbb{N} \) (the “spin dimension”), we let \( \mathcal{F} \subset L(\mathcal{H}) \) be the set of all symmetric operators on \( \mathcal{H} \) of finite rank, which (counting multiplicities) have at most \( n \) positive and at most \( n \) negative eigenvalues. On \( \mathcal{F} \) we are given a positive measure \( \rho \) (defined on a \( \sigma \)-algebra of subsets of \( \mathcal{F} \)). We refer to \((\mathcal{H}, \mathcal{F}, \rho)\) as a causal fermion system.

This definition evolved over several years. Based on preparations in [39], the present formulation was first given in [57].

This definition is illustrated in Figure 5.3. Now the set \( \mathcal{F} \) is invariant in addition under the transformation where an operator is multiplied by a negative number, as is indicated in the figure by the double cones. The support of the measure, denoted by

\[
M := \text{supp } \rho, \tag{5.3.1}
\]

is referred to as spacetime. In generalization of the example of the lattice system, where spacetime consisted of discrete points (5.2.3), now the measure \( \rho \) can also have continuous components. The measure \( \rho(\Omega) \) of a measurable subset \( \Omega \subset M \) can be regarded as the four-dimensional volume of the spacetime region \( \Omega \). It is a specific feature of a causal fermion system that a spacetime point \( x \in M \) is a linear operator on the Hilbert space \( \mathcal{H} \). This endows spacetime with a lot of additional structure. In particular, as will be explained in Section 5.6, the spacetime point operators give rise to an ensemble of spinorial wave functions and to causal and geometric structures. The general idea is that a causal fermion system describes a spacetime together with all structures therein. Before entering these structures in more detail, we illustrate the general definition by the simple and concrete example of Dirac wave functions in Minkowski space.

5.4. Example: Dirac Wave Functions in Minkowski Space

As a further example, we now explain how to construct a causal fermion system in Minkowski space. Recall that in Section 1.4 (and similarly in curved spacetime in Section 4.5) we introduced the Hilbert space \((\mathcal{H}_m, \langle | . \rangle)\) of all solutions of the Dirac equation. We now choose a closed subspace \( \mathcal{F}_m \) of this Hilbert space and denote the scalar product \( \langle | . \rangle \) restricted to this subspace by \( \langle | . \rangle_{\mathcal{F}_m} \). We thus obtain the Hilbert space \((\mathcal{H}, \langle | . \rangle_{\mathcal{F}_m})\).
By construction, the vectors in this Hilbert space are solutions of the Dirac equation. They can be thought of as the “occupied states” of the system. We prefer the notion of physical wave functions, where “physical” means intuitively that these wave functions are realized in our physical system (whatever this means; we shall not enter philosophical issues here). At this stage, we do not need to specify \( \mathcal{H} \). More explicit formulas and computations can be found in [48, Section 1.2] or in the later chapters of this book (Chapters 15-19).

We point out that the functions in \( \mathcal{H} \) do not need to be continuous (instead, as mentioned at the end of Section 1.4, their restriction to any Cauchy surface merely is an \( L^2 \)-function). Therefore, we cannot evaluate the wave functions pointwise at a spacetime point \( x \in \mathcal{M} \). However, for the following constructions it is crucial to do so. The way out is to introduce so-called regularization operators \( (\mathcal{R}_\varepsilon) \) with \( 0 < \varepsilon < \varepsilon_{\text{max}} \) which map \( \mathcal{H} \) to the continuous wave functions,\[ \mathcal{R}_\varepsilon : \mathcal{H} \to C^0(\mathcal{M}, S\mathcal{M}) . \tag{5.4.1} \]

In the limit \( \varepsilon \downarrow 0 \), these operators should go over to the identity (in a suitable sense which we do not need to specify here). The physical picture is that on a small length scale, which can be thought of as the Planck length scale \( \varepsilon \approx 10^{-35} \) meters, the structure of spacetime must be modified. The regularization operators specify this microscopic structure of spacetime. A simple example of a regularization operator is obtained by mollifying with a test function. Thus we let \( h \in C^\infty_0(\mathcal{M}, \mathbb{R}) \) be a non-negative test function with \( \hat{h}(x) = 1 \). We define the operators \( \mathcal{R}_\varepsilon \) for \( \varepsilon > 0 \) as the convolution operators\[ (\mathcal{R}_\varepsilon u)(x) := \frac{1}{\varepsilon^4} \int_{\mathcal{M}} h\left(\frac{x-y}{\varepsilon}\right) u(y) \, d^4 y . \]

Another method is to work in Fourier space (for preliminaries see Sections 1.5 and 2.4), i.e.\[ \psi(x) = \int \frac{d^4 k}{(2\pi)^4} \hat{\psi}(k) e^{-ikx} , \] and to regularize by multiplication with an exponentially decaying cutoff function, i.e.\[ (\mathcal{R}_\varepsilon \psi)(x) = \int \frac{d^4 k}{(2\pi)^4} \hat{\psi}(k) e^{-\varepsilon |\omega|} e^{-ikx} \quad \text{with} \quad \omega = k^0 . \tag{5.4.2} \]

This so-called \( \varepsilon \)-regularization is most convenient for explicit computations (for more details see [48 §2.4.1]). Clearly, these methods of regularizing Dirac solutions are very special and should be thought of merely as a mathematical tool for constructing simple and explicit examples of causal fermion systems.

Next, for any \( x \in \mathcal{M} \) we consider the bilinear form \[ b_x : \mathcal{H} \times \mathcal{H} \to \mathbb{C} , \quad b_x(u, v) = -\langle (\mathcal{R}_\varepsilon u)(x) \mid (\mathcal{R}_\varepsilon v)(x) \rangle . \tag{5.4.3} \]

This bilinear form is well-defined and bounded because \( \mathcal{R}_\varepsilon \) is defined pointwise and because evaluation at \( x \) gives a linear operator of finite rank. Thus for any \( v \in \mathcal{H} \), the anti-linear form \( b_x(\cdot, v) : \mathcal{H} \to \mathbb{C} \) is continuous. By the Fréchet-Riesz theorem (Theorem 2.2.4), there is a unique vector \( w \in \mathcal{H} \) such that \( b_x(u, v) = \langle u|w\rangle_{\mathcal{H}} \) for all \( u \in \mathcal{H} \). The
5.5. The Causal Action Principle

Having given the general definition of a causal fermion system (see Definition 5.3.1), the question arises how physical equations can be formulated in this setting. To this end, we now formulate a variational principle, the so-called causal action principle. The mathematical structure of the causal action is similar to the action (5.1.5) given in our example of the lattice system. Its detailed form, however, is the result of many computations and longer considerations, as will be outlined in Section 5.7 below.

For any \( x, y \in \mathcal{F} \), the product \( xy \) is an operator of rank at most \( 2n \). However, in general it is no longer a symmetric operator because \( (xy)^* = yx \), and this is different from \( xy \) unless \( x \) and \( y \) commute. As a consequence, the eigenvalues of the operator \( xy \) are in general complex. We denote these eigenvalues counting algebraic multiplicities by \( \lambda_1^{xy}, \ldots, \lambda_k^{xy} \in \mathbb{C} \) (more specifically, denoting the rank of \( xy \) by \( k \leq 2n \), we choose \( \lambda_1^{xy}, \ldots, \lambda_k^{xy} \) as all the non-zero eigenvalues and set \( \lambda_{k+1}^{xy}, \ldots, \lambda_{2n}^{xy} = 0 \)). We introduce the Lagrangian and the causal action by

\[
\text{Lagrangian: } \quad \mathcal{L}(x, y) = \frac{1}{4n} \sum_{i,j=1}^{2n} \left( |\lambda_i^{xy}| - |\lambda_j^{xy}| \right)^2 \\
\text{causal action: } \quad S(\rho) = \int_{\mathcal{F} \times \mathcal{F}} \mathcal{L}(x, y) \, d\rho(x) \, d\rho(y) .
\]
The causal action principle is to minimize $S$ by varying the measure $\rho$ under the following constraints:

- **volume constraint:** $\rho(\mathcal{F}) = \text{const}$ \hfill (5.5.3)
- **trace constraint:** $\int_{\mathcal{F}} \text{tr}(x) \, d\rho(x) = \text{const}$ \hfill (5.5.4)
- **boundedness constraint:** $\mathcal{T}(\rho) := \int_{\mathcal{F} \times \mathcal{F}} \left( \sum_{i=1}^{2n} |\lambda_{ij}^{xy}| \right)^2 \, d\rho(x) \, d\rho(y) \leq C$, \hfill (5.5.5)

where $C$ is a given parameter (and $\text{tr}$ denotes the trace of a linear operator on $\mathcal{H}$).

In order to make the causal action principle mathematically well-defined, one needs to specify the class of measures in which to vary $\rho$. To this end, on $\mathcal{F}$ we consider the topology induced by the operator norm $\| A \| := \sup \{ \| Au \|_{\mathcal{H}} \text{ with } \| u \|_{\mathcal{H}} = 1 \}$ \hfill (5.5.6)

(for basics see the preliminaries in Sections 2.1 and 2.2). In this topology, the Lagrangian as well as the integrands in (5.5.4) and (5.5.5) are continuous. The $\sigma$-algebra generated by the open sets of $\mathcal{F}$ consists of the so-called Borel sets. A regular Borel measure is a measure on the Borel sets with the property that it is continuous under approximations by compact sets from inside and by open sets from outside (for basics see the preliminaries in Section 2.3). The right prescription is to vary $\rho$ within the class of regular Borel measures on $\mathcal{F}$. There are two settings one can consider:

(a) The finite-dimensional setting: $\dim \mathcal{H} < \infty$ and $\rho(\mathcal{F}) < \infty$.

In this case, we will prove the existence of minimizers in Chapter 12. This will also clarify the significance of the constraints.

(b) The infinite-dimensional setting: $\dim \mathcal{H} = \infty$ and $\rho(\mathcal{F}) = \infty$.

An obvious complication in this setting is that the volume constraint (5.5.3) is infinite. Likewise, the other constraints as well as the causal action may diverge. These divergences can be avoided by restricting attention to variations which change the measure only on a set of finite volume. By doing so, the differences of the action and the constraints are well-defined and finite (this method will be explained in Section 12.8).

With this in mind, the main difficulty is to deal with infinite-dimensional Hilbert spaces. This case also seems to make mathematical and physical sense. However, the existence theory has not yet been developed. At least, it is known that the Euler-Lagrange equations corresponding to the causal action principle still have a mathematical meaning (for details see [48]). One way of getting along without an existence theory in the infinite-dimensional setting is to take the point of view that on a fundamental physical level, the total volume is finite and the Hilbert space $\mathcal{H}$ is finite-dimensional, whereas the infinite-dimensional setting merely is a mathematical idealization needed in order to describe systems in infinite volume involving an infinite number of quantum particles.

We now explain what spacetime and the underlying causal structure is. Given a minimizing measure $\rho$, we again define spacetime $\mathcal{M}$ as the support of the $\rho$ \hfill (5.3.1) (this is illustrated in Exercise 6.1). Thus the spacetime points are symmetric linear operators on $\mathcal{H}$. On $\mathcal{M}$ we consider the topology induced by $\mathcal{F}$ (generated by the sup-norm (5.5.6) on $\text{L}(\mathcal{H})$). Moreover, the measure $\rho|_{\mathcal{M}}$ restricted to $\mathcal{M}$ can be regarded as...
a volume measure on spacetime. This makes spacetime to a topological measure space. Furthermore, one has the following notion of causality:

**Definition 5.5.1. (causal structure)** For any \( x, y \in F \), the product \( xy \) is an operator of rank at most \( 2n \). We denote its non-trivial eigenvalues (counting algebraic multiplicities) by \( \lambda_{xy}^1, \ldots, \lambda_{xy}^n \). The points \( x \) and \( y \) are called *spacelike* separated if all the \( \lambda_{xy}^j \) have the same absolute value. They are said to be *timelike* separated if the \( \lambda_{xy}^j \) are all real and do not all have the same absolute value. In all other cases (i.e. if the \( \lambda_{xy}^j \) are not all real and do not all have the same absolute value), the points \( x \) and \( y \) are said to be *lightlike* separated.

Restricting the causal structure of \( F \) to \( M \), we get causal relations in spacetime.

The Lagrangian (5.5.1) is compatible with the above notion of causality in the following sense. Suppose that two points \( x, y \in F \) are spacelike separated. Then the eigenvalues \( \lambda_{xy}^j \) all have the same absolute value. As a consequence, the Lagrangian (5.5.1) vanishes. Thus pairs of points with spacelike separation do not enter the action. This can be seen in analogy to the usual notion of causality where points with spacelike separation cannot influence each other. This analogy is the reason for the notion “causal” in “causal fermion system” and “causal action principle.”

Moreover, a causal fermion system distinguishes a direction of time. To this end, we let \( \pi_x \) be the orthogonal projection in \( H \) on the subspace \( x(H) \subset H \) and introduce the functional

\[
\mathcal{C} : M \times M \to \mathbb{R}, \quad \mathcal{C}(x, y) := i \, \text{tr} \left( yy^* \pi_x - xy^* \pi_y \right).
\]

Obviously, this functional is anti-symmetric in its two arguments, making it possible to introduce the notions

\[
\begin{cases} 
& y \text{ lies in the future of } x \quad \text{if } \mathcal{C}(x, y) > 0 \\ & y \text{ lies in the past of } x \quad \text{if } \mathcal{C}(x, y) < 0.
\end{cases}
\]

By distinguishing a direction of time, we get a structure similar to a causal set (see for example [17]). However, in contrast to a causal set, our notion of “lies in the future of” is not necessarily transitive.

### 5.6. Basic Inherent Structures

It is the general concept that a causal fermion system describes spacetime as well as all structures therein (like the causal and metric structures, particles, fields, etc.). In order to recover all these structures, we merely give structures of the causal fermion system suitable names. We point out that we do not introduce additional structures, but we merely work with structures which are already encoded in the causal fermion system. We call these structures as being inherent in the causal fermion system. We now introduce and explain the most important of these structures: the *spin spaces*, the *physical wave functions* and the *kernel of the fermionic projector*. Other inherent structures will be introduced later in this book (see Chapters 9–11); for a more complete account we also refer to [48 Chapter 1].

The causal action principle depends crucially on the eigenvalues of the operator product \( xy \) with \( x, y \in F \). For computing these eigenvalues, it is convenient not to consider this operator product on the (possibly infinite-dimensional) Hilbert space \( H \), but instead to restrict attention to a finite-dimensional subspace of \( H \), chosen such that the operator product vanishes on the orthogonal complement of this subspace. This construction
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Hilbert space $\mathcal{H}$

Figure 5.4. The spin spaces

$$S_x \subset \mathcal{H}$$

Figure 5.5. The kernel of the fermionic projector

leads us to the spin spaces and to the kernel of the fermionic projector, which we now introduce. For every $x \in \mathcal{F}$ we define the spin space $S_x$ by $S_x = x(\mathcal{H})$; it is a subspace of $\mathcal{H}$ of dimension at most $2n$ (see Figure 5.4). Moreover, we let

$$\pi_x : \mathcal{H} \to S_x$$

be the orthogonal projection in $\mathcal{H}$ on the subspace $S_x \subset \mathcal{H}$. For any $x, y \in M$ we define the kernel of the fermionic projector $P(x, y)$ by (see Figure 5.5).

$$P(x, y) = \pi_x y|_{S_y} : S_y \to S_x$$

where $\pi_x$ is again the orthogonal projection on the subspace $x(\mathcal{H}) \subset \mathcal{H}$). Taking the trace of $(5.6.2)$ in the case $x = y$, one finds that $\text{tr}(x) = \text{Tr}_{S_x}(P(x, x))$, making it possible to express the integrand of the trace constraint (5.5.4) in terms of the kernel of the fermionic projector. In order to also express the eigenvalues of the operator $xy$ in terms of the kernel of the fermionic projector, we introduce the closed chain $A_{xy}$ as the product

$$A_{xy} = P(x, y) P(y, x) : S_x \to S_x.$$

Computing powers of the closed chain, one obtains

$$A_{xy} = (\pi_x y) (\pi_y x)|_{S_x} = \pi_x y x|_{S_x}, \quad (A_{xy})^p = \pi_x (y x)^p|_{S_x}.$$

Taking the trace, one sees in particular that

$$\text{Tr}_{S_x}(A_{xy}^p) = \text{tr} ((y x)^p) = \text{tr} ((x y)^p)$$

(5.6.4)

(5.6.2)

(5.6.1)

(5.6.3)

(5.6.4)

Since all our operators have finite rank, there is a finite-dimensional subspace $I$ of $\mathcal{H}$ such that $xy$ maps $I$ to itself and vanishes on the orthogonal complement of $I$. Then the non-trivial eigenvalues of the operator product $xy$ are given as the zeros of the characteristic polynomial of the restriction $xy|_I : I \to I$. The coefficients of this characteristic polynomial (like the trace, the determinant, etc.) are symmetric polynomials in the eigenvalues and can therefore be expressed in terms of traces of powers of $A_{xy}$. We conclude that the eigenvalues of the closed chain coincide with the non-trivial eigenvalues $\lambda_1^{xy}, \ldots, \lambda_n^{xy}$ of the
5.6. BASIC INHERENT STRUCTURES

The physical wave function

Figure 5.6. The physical wave function

operator $xy$ in Definition 5.5.1. This argument also shows that the operator products $xy$ and $yx$ are isospectral. In particular, one sees that kernel of the fermionic projector encodes the causal structure of $M$. Moreover, we see that the causal structure is symmetric in $x$ and $y$. The main advantage of working with the kernel of the fermionic projector is that the closed chain (5.6.3) is a linear operator on a vector space of dimension at most $2^n$, making it possible to compute the $\lambda_{xy}^1, \ldots, \lambda_{xy}^{2^n}$ as the eigenvalues of a finite matrix.

Next, it is very convenient to arrange that the kernel of the fermionic projector is symmetric in the sense that

$$P(x,y)^* = P(y,x).$$

(5.6.5)

To this end, one chooses on the spin space $S_x$ the spin inner product $\langle \cdot \mid \cdot \rangle_x$ by

$$\langle u \mid v \rangle_x = -\langle u \mid xv \rangle_H \quad \text{(for all $u, v \in S_x$)}.$$  

(5.6.6)

Due to the factor $x$ on the right, this definition really makes the kernel of the fermionic projector symmetric, as is verified by the computation

$$\langle u \mid P(x,y) v \rangle_x = -\langle u \mid P(x,y) v \rangle_H = -\langle u \mid xy v \rangle_H$$

$$= -\langle \pi_y x u \mid y v \rangle_H = -\langle P(y,x) u \mid v \rangle_y$$

(where $u \in S_x$ and $v \in S_y$). The spin space $(S_x, \langle \cdot \mid \cdot \rangle_x)$ is an indefinite inner product of signature $(p,q)$ with $p,q \leq n$ (for textbooks on indefinite inner product spaces see [16, 89]). In this way, indefinite inner product spaces arise naturally when analyzing the mathematical structure of the causal action principle.

The kernel of the fermionic projector plays a central role in the analysis for several reasons:

- The Lagrangian can be expressed in terms of $P(x,y)$ (via the closed chain (5.6.3) and its eigenvalues).
- Being a mapping from one spin space to another, $P(x,y)$ gives relations between different spacetime points. In this way, it carries geometric information. This will be explained in Chapter 11 (see also [55] or the introductory survey paper [50]).
- The kernel of the fermionic projector also encodes all the wave functions of the system. In order to see the connection, for a vector $u \in \mathcal{H}$ one introduces the corresponding physical wave function $\psi^u$ as (see Figure 5.6)

$$\psi^u : M \to \mathcal{H}, \quad \psi^u(x) = \pi_x u \in S_x.$$  

(5.6.7)

Then, choosing an orthonormal basis $(e_i)$ of $\mathcal{H}$ and using the completeness relation as well as (5.6.6), one obtains for any $\phi \in S_y$

$$P(x,y) \phi = \pi_x y |S_y \phi \phi = \sum_i \pi_x e_i \langle e_i \mid y \phi \rangle_H = -\sum_i \psi^{e_i}(x) \langle \psi^{e_i}(y) \mid \phi \rangle_y,$$
showing that $P(x, y)$ is indeed composed of all the physical wave functions, i.e. in a bra/ket notation

$$P(x, y) = - \sum_i \langle \psi^{\ell_i}(x) \rangle \langle \psi^{\ell_i}(y) \rangle.$$  \hfill (5.6.8)

### 5.7. How Did the Causal Action Principle Come About?

Causal fermion systems and the causal action principle came to light as a result of many considerations and computations carried out over several years. We now give an outline of these developments, also explaining the specific form of the causal action principle.

The starting point for the considerations leading to causal fermion systems was the belief that in order to overcome the conceptual problems of quantum field theory, the structure of spacetime should be modified. Moreover, instead of starting from differential equations in a spacetime continuum, one should formulate the physical equations using the new structures of spacetime, which might be non-smooth or discrete. A more concrete idea in this direction was that the spacetime structures should be encoded in the ensemble of wave functions which is usually associated to the Dirac sea (for basics see Section 1.5). Thus, instead of disregarding the sea states, one should take all these wave functions into account. The mutual interaction of all these wave functions should give rise to the structures of spacetime as we experience them.

The first attempts towards making this idea more precise go back to the early 1990s. The method was to consider families of Dirac solutions (the formalism of quantum fields was avoided in order to keep the setting as simple and non-technical as possible). In order to describe such a family mathematically, the corresponding two-point kernel $P(x, y)$ was formed

$$P(x, y) := - \sum_{l=1}^f \langle \psi_l(x) \rangle \langle \psi_l(y) \rangle$$

(5.7.1)

(5.7.2)

where $\psi_1, \ldots, \psi_f$ are suitably normalized solutions of the Dirac equation; for preliminaries see Section 1.3. The kernel $P(x, y)$ is also referred to as the kernel of the fermionic projector. In the Minkowski vacuum, this kernel is formed of all the states of the Dirac sea. Then the sum goes over to an integral over the lower mass shell

$$P^{\text{vac}}(x, y) = \int \frac{d^4k}{(2\pi)^4} (\not{k} + m) \delta(k^2 - m^2) \theta(-k_0) e^{-ik(x-y)}$$

this integral is well-defined as the Fourier transform of a tempered distribution; see the preliminaries in Section 2.4. Likewise, a system involving particles and anti-particles is described by “occupying additional states of positive energy” and by “creating holes in the Dirac sea,” respectively. Thus, more technically, one sets

$$P(x, y) = P^{\text{vac}}(x, y) - \sum_a \langle \psi_a(x) \rangle \langle \psi_a(y) \rangle + \sum_b \langle \phi_b(x) \rangle \langle \phi_b(y) \rangle,$$

where $\psi_a$ and $\phi_b$ are suitably normalized Dirac solutions of positive and negative energy, respectively. In case a bosonic interaction is present, the kernel of the fermionic projector should no longer satisfy the vacuum Dirac equation, but the Dirac equation in the presence of a, say, external potential $B$. Clarifying the dependence on the bosonic potential with an additional tilde, we write the resulting Dirac equation as

$$(i\not{\partial} + B - m) \tilde{P}(x, y) = 0.$$  \hfill (5.7.3)
Analyzing the distribution $\tilde{P}(x,y)$ in Minkowski space reveals the following facts:

(a) The kernel $\tilde{P}(x,y)$ contains all the information on the wave functions of the particles and anti-particles of the system. This can be understood from the representation (5.7.2) in which all these wave functions appear.

(b) The kernel $\tilde{P}(x,y)$ has singularities on the light cone. The detailed form of the singularities involves integrals of the potential $\mathcal{B}$ and its derivatives along the light cone. In particular, knowing the kernel $\tilde{P}(x,y)$ makes it possible to reconstruct the potential $\mathcal{B}$ at every spacetime point.

(c) The singularity structure of $\tilde{P}(x,y)$ encodes the causal structure of Minkowski space.

These findings show that, at least for Dirac systems in the presence of classical bosonic potentials, the kernel $\tilde{P}(x,y)$ contains all the information on the physical system. This led to the concept to regard $\tilde{P}(x,y)$ as the basic physical object in spacetime. The more familiar structures and objects like Minkowski space with its causal structure, the Dirac equation, the classical field equations for the bosonic fields (like the Maxwell or Einstein equations), however, should no longer be considered as being fundamental. Consequently, the physical equations should be formulated directly in terms of the kernel of the fermionic projector.

Formalizing this idea in a clean way also makes it necessary to disregard or to prescind from the usual spacetime structures. This led to the principle of the fermionic projector as formulated around 1990 (see the unpublished preprint [33] and the monograph [39]). We here a slightly different but equivalent formulation which is somewhat closer to the setting of causal fermion systems. Let $M$ be a discrete set (i.e. a point set without additional structures), the discrete spacetime. Moreover, to every $x \in M$ we associate an indefinite inner product space $(S_x, \prec \cdot, \prec \cdot)_x$, referred to as the spin space at $x$. Next, we consider wave functions $\psi_a$, being mappings which to every discrete spacetime point $x \in M$ associates a vector $\psi_a(x) \in S_x$ of the corresponding spin space. Out of these wave functions, one can form the kernel of the fermionic projector

$$P(x,y) := -\sum_a |\psi(x)\prec \psi(y)| : S_y \to S_x.$$  

The principle of the fermionic projector asserts that the physical equations should be formulated purely in terms of the kernel of the fermionic projector in discrete spacetime.

The next question was how precisely these physical equations should look like. This was a difficult question which took many years to be answered. Apart from the structural requirements coming from the principle of the fermionic projector, the following considerations served as guiding principles:

(i) In analogy to classical field theory, a variational approach should be used. One main advantage is the resulting connection between symmetries and conservation laws (corresponding to the classical Noether theorem), which seems of central importance in physical applications.

(ii) Classical field theory should be obtained in a certain limiting case. More specifically, the Euler-Lagrange equations coming from our variational principle should reproduce the Maxwell and Einstein equations.

(iii) Also the Dirac equation should be recovered in a certain limiting case.

\footnote{Of course, it is also an important requirement that our variational principle should give agreement with quantum field theory. But this connection was not used for finding the causal action principle. It was worked out more recently; for more details see Chapter 22.}
A brief introduction to causal fermion systems

More mathematically, the strategy was to form composite expressions of the kernel of the fermionic projector. More precisely, choosing \( n \) points \( x_1, \ldots, x_n \in M \), one can form the closed chain
\[
A_{x_1, \ldots, x_n} := P(x_1, x_2) P(x_2, x_3) \cdots P(x_{n-1}, x_n) P(x_n, x_1) : S_{x_1} \to S_{x_1}.
\] (5.7.4)

Being an endomorphism of the spin space, one can compute the eigenvalues of the closed chain and form a Lagrangian \( \mathcal{L}[A_{x_1, \ldots, x_n}] \) as a symmetric function of these eigenvalues. Summing over the spacetime points gives an ansatz for the \( n \)-point action
\[
S = \sum_{x_1, \ldots, x_n \in M} \mathcal{L}[A_{x_1, \ldots, x_n}].
\]

This ansatz can be made more specific by considering gauge phases. This consideration was motivated by the fact that the kernel of the fermionic projector \( \tilde{P}(x, y) \) formed of Dirac solutions involves gauge phases. More specifically, choosing the potential in the Dirac equation (5.7.3) as an electromagnetic potential, i.e. \( \mathcal{B} = \mathcal{A} \), then the leading contribution to kernel are gauge phases described by line integrals over the electromagnetic potential,
\[
\tilde{P}(x, y) = e^{-i \int_x^y A_j \xi^j} P^{\text{vac}}(x, y) + \cdots,
\] (5.7.5)

where
\[
\int_x^y A_j \xi^j = \int_0^1 A_j(\alpha y + (1 - \alpha)x) (y - x)^j \, d\alpha.
\]

Here \( \cdots \) stands for many other contributions to \( \tilde{P}(x, y) \) which involve derivatives of the potential (like the field tensor, the Maxwell current, etc.). All these additional contributions are small in the sense that they are less singular on the light cone. These findings will be made precise by the Hadamard and light-cone expansions of the kernel of the fermionic projector in Chapter 21 of this book. At this stage, we do not need to be specific. All we need is that gauge phases come into play, which involves integrals of the potential along the line segment joining the points \( x \) and \( y \).

Let us analyze the effect of the gauge phases on the closed chain (5.7.4). First of all, the closed chain is gauge invariant. Indeed, if one considers a pure gauge potential \( A_j = \partial_j \Lambda \), then the gauge phases in (5.7.5) simplify to
\[
\tilde{P}(x, y) = e^{-i \Lambda(y) + i \Lambda(x)} P^{\text{vac}}(x, y),
\] (5.7.6)

and the phase factors of neighboring factors cancel in (5.7.4). This consideration also explains why the arguments of the adjacent factors in (5.7.4) coincide. In the case \( n = 1 \), the kernel of the fermionic projector is evaluated only on the diagonal \( P(x, x) \). This turns out to be too simple, as can be understood from the fact that relations between spacetime points are not taken into account. If \( n \geq 3 \), the gauge phases in (5.7.4) can be rewritten using Stokes’ theorem as flux integrals of the electromagnetic field through the two-dimensional polygon with vertices \( x_1, \ldots, x_n \). Analyzing the situation in some more detail, one finds that the resulting Euler-Lagrange equations will be satisfied only if all fluxes vanish. This implies that the electromagnetic potential must be a pure gauge potential. In other words, the case \( n \geq 3 \) does not allow for an interaction via gauge potentials. This is the reason why this case was disregarded (for some more details on this argument see [39, Remark 6.2.5]).

After these considerations, we are left with the two-point action
\[
S = \sum_{x, y \in M} \mathcal{L}[A_{xy}],
\] (5.7.7)
where $A_{xy}$ is the closed chain formed of two points,

$$A_{xy} := P(x, y) P(y, x).$$

(5.7.8)

In this case, the polygon with vertices $x$ and $y$ degenerates to a straight line, implying that the flux through this polygon vanishes as desired. The starting point for a more quantitative analysis was to choose the Lagrangian formed by taking products and sums of traces or powers the closed chain. A typical example is the Lagrangian

$$\mathcal{L}_{A_{xy}} := \text{Tr}_{S_x} (A_{xy}^2) - c (\text{Tr}_{S_x} (A_{xy}))^2$$

(5.7.9)

with a real parameter $c$. In such examples, the Lagrangian is a symmetric polynomial in the eigenvalues of the closed chain. The methods and results of this early analysis can be found in the unpublished preprints [33, 34].

Generally speaking, working with such polynomial Lagrangians seemed a promising approach. However, the more detailed analysis revealed the basic problem that chiral gauge phases come into play. As just explained after (5.7.6), the closed chain and therefore also the Lagrangian are gauge invariant for the electromagnetic potential. However, the situation changes if chiral gauge potentials (in particular the left-handed gauge potentials of electroweak theory) are considered. In this case, the left- and right-handed components of $P(x, y)$ involve phase transformations by the left- and right-handed gauge potentials, respectively. When forming the closed chain (5.7.8), the left- and right-handed components of $P(x, y)$ are multiplied together. As a consequence, the closed chain involves relative phases of the left- and right-handed gauge potentials, i.e. phase factors of the form

$$e^{\pm i \int_x^y (A_L - A_R) \gamma \xi}$$

where $A_L$ and $A_R$ are the left- and right-handed gauge potentials (here for simplicity again Abelian). As a consequence, also the eigenvalues of the closed chain are multiplied by these relative phases. The traces of powers of the closed chain as in (5.7.9) are still real-valued (this is because the phase factors always come as complex conjugate pairs), but they do not have fixed signs. Working out the Euler-Lagrange equations, one sees that they also involve the relative gauge phases, making it difficult to allow for chiral gauge fields. In order to bypass these difficulties, from around 1999 on Lagrangians were considered which involved absolute values of the eigenvalues of the closed chain. This had two major advantages:

(a) The chiral gauge phases drop out of the Lagrangian.

(b) It became natural to formulate non-negative Lagrangians. As a consequence, in the variational principle one minimize the action instead of merely looking for critical points.

(c) A connection to causality was obtained. In order to see how this comes about, we give a simple computation in the Minkowski vacuum. Suppose that the points $x$ and $y$ are either timelike or spacelike separated. Then $P(x, y)$ is well-defined and finite even without regularization and, due to Lorentz symmetry, it has the form

$$P(x, y) = \alpha \xi \gamma^j + \beta \mathbb{1}$$

with two complex-valued functions $\alpha$ and $\beta$ (where again $\xi = y - x$). Taking the adjoint with respect to the spin scalar product, we see that

$$P(y, x) = \overline{\alpha} \xi \gamma^j + \overline{\beta} \mathbb{1}.$$
As a consequence,
\[ A_{xy} = P(x, y) P(y, x) = a \xi_j \gamma^j + b \mathbf{1} \]
with two real parameters \(a\) and \(b\) given by
\[ a = \alpha \beta + \beta \alpha, \quad b = |\alpha|^2 \xi^2 + |\beta|^2. \]
Applying the formula \((A_{xy} - b \mathbf{1})^2 = a^2 \xi^2 \mathbf{1}\), the roots of the characteristic polynomial of \(A_{xy}\) are computed by
\[ b \pm \sqrt{a^2 \xi^2}. \]
Therefore, the eigenvalues of the closed chain are either real, or else they form a complex conjugate pair. Spacelike separation is characterized by the fact that all eigenvalues of \(A_{xy}\) have the same absolute value. A Lagrangian which only depends on differences of absolute values of the eigenvalues vanishes for spacelike separation.

The further analysis led to the class of Lagrangians
\[ \mathcal{L} = \sum_{i,j} \left( |\lambda_i^{xy}|^p - |\lambda_j^{xy}|^p \right)^2. \]
with a parameter \(p \in \mathbb{N}\), where the \(\lambda_i^{xy}\) are the eigenvalues of \(A_{xy}\) (again counted with algebraic multiplicities). The case \(p = 1\) gives the causal Lagrangian \((5.5.1)\) (albeit with the difference of working instead of the local correlation operators with the kernel of the fermionic projector; the connection will be explained below). The decision for \(p = 1\) was taken based on the so-called state stability analysis, which revealed that the vacuum Dirac sea configuration \((5.7.1)\) is a local minimizer of the causal action only if \(p = 1\) (for details see [39, Section 5.5]). Now that the form of the causal action was fixed, the monograph [39] was completed and published. The causal action principle is given in this book as an example of a variational principle in discrete spacetime (see [39, Section 3.5]). The boundedness constraint \((5.5.5)\) already appears, and the causal Lagrangian \((5.5.1)\) arises when combining the Lagrangian with the Lagrange multiplier term corresponding to the boundedness constraint. The volume constraint \((5.5.3)\) is also implemented, however in discrete spacetime simply as the condition that the number of spacetime points be fixed (and \(\rho\)-integrals are replaced by sums over the spacetime points). The trace constraint, however, was not yet recognized as being necessary and important.

After the publication of the monograph [39], the causal action principle was analyzed in more detail and more systematically, starting from simple systems and proceeding to more realistic physical models, concluding with systems showing all the interactions of the standard model and gravity (see [48, Chapters 3-5]). This detailed study also led to the causal action principle in the form given in Section 5.5 above. The path from the monograph [39] to the present formulation in [48] is outlined in [39, Preface to second online edition]. We now mention a few points needed for the basic understanding.

One major conceptual change compared to the setting in indefinite inner product spaces was to recognize that an underlying Hilbert space structure is needed in order for the causal variational principle to be mathematically well-defined. This became clear when working on the existence theory in discrete spacetime [41]. This Hilbert space structure is built in most conveniently by working instead of the kernel of the fermionic projector with the local correlation operators which relate the Hilbert space scalar product to the spin inner product by
\[ \langle \psi | F(x) \phi \rangle_{\mathcal{F}} = -\langle \psi(x) | \phi(x) \rangle_x. \]
Using that the operator product $F(x)F(y)$ has the same non-trivial eigenvalues as the closed chain $A_{xy}$ given by \((5.7.8)\) (as we already observed in Section 5.6 after \((5.6.3)\)), the causal action principle can also be formulated in terms of the local correlation operators $F(x)$ with $x \in M$. Moreover, it turned out that measure-theoretic methods can be used to generalize the setting such as to allow for the description of not only discrete, but also continuous spacetimes. In this formulation, the sums over the discrete spacetime points are replaced by integrals with respect to a measure $\mu$ on $M$. This setting was first introduced in [43] when working out the existence theory. In this formulation, the only a-priori structure of spacetime is that of a measure space $(M, \mu)$. The local correlation operators give rise to a mapping

$$F : M \to \mathcal{F}, \quad x \mapsto F(x),$$

where $\mathcal{F}$ is the subset of finite rank operators on $\mathcal{H}$ which are symmetric and (counting multiplicities) have at most $n$ positive and at most $n$ negative eigenvalues (where $n$ is introduced via the signature $(n, n)$ of the indefinite inner product in \((5.7.10)\)). This analysis also revealed the significance of the trace constraint. As the final step, instead of working with the measure $\mu$, the causal action can be expressed in terms of the push-forward measure $\rho = F_*\mu$, being a measure on $\mathcal{F}$ (defined by $\rho(\Omega) = \mu(F^{-1}(\Omega))$). Therefore, it seems natural to leave out the measure space $(M, \mu)$ and to work instead directly with the measure $\rho$ on $\mathcal{F}$.

These considerations led to the general definition of causal fermion systems in Section 5.3 where the physical system is described by a Hilbert space $(\mathcal{H}, \langle \cdot, \cdot \rangle_{\mathcal{H}})$ and the measure $\rho$ on $\mathcal{F}$. The causal action principle takes the form as stated in Section 5.5.

5.8. Underlying Physical Concepts

We now briefly explain a few physical concepts behind causal fermion systems and the causal action principle. The aim is to convey the reader the correct physical picture in a non-technical way. Doing so already here makes it necessary to anticipate some ideas on a qualitative level which will be introduced more systematically and thoroughly later in this book.

It is a general feature of causal fermion systems that the usual distinction between the structure of spacetime itself (being modelled by Minkowski space or a Lorentzian manifold) and structures in spacetime (like wave functions and matter fields) ceases to exist. Instead, all these structures are described as a whole by a single object: the measure $\rho$ on $\mathcal{F}$. Spacetime and all structures therein are different manifestations of this one object. The dynamics of spacetime and of all objects in spacetime are described in a unified and holistic manner by the causal action principle. Clearly, in order to get a connection to the conventional description of physics, one still needs to construct the familiar physical objects from a causal fermion system. Also, one needs to rewrite the dynamics as described by the causal action principle in terms of these familiar physical objects. This study is a main objective of this book. As already exemplified in Section 5.6 by the spin spaces and physical wave functions, the strategy is to give structures present in a causal fermion system suitable names. This must be done carefully in such a way that these conveys the correct physical picture. Ultimately, the inherent structures serve the purpose of getting a better understanding of the causal action principle. As we shall see, this will be achieved by reformulating the Euler-Lagrange equations of the causal action principle in terms of the inherent structures. When this is done, also the physical names
of the inherent structures will be justified by showing that they agree with the familiar physical objects in specific limiting cases and generalize these objects in a sensible way.

In view of this unified description of all physical structures by a single mathematical object, it is difficult to describe the essence of causal fermion systems using the familiar notions from physics. One simple way of understanding the causal action principle is to focus on the structure of the physical wave functions and the kernel of the fermionic projector. Clearly, the resulting picture is a bit oversimplified, because it only captures part of the structures encoded in a causal fermion system. Nevertheless, it conveys a good and the correct intuition of what the causal action principle is about. We saw in Section 5.6 that a causal fermion system gives rise to the ensemble of physical wave functions $\psi^u \in H$ (see (5.6.7)). The kernel of the fermionic projector (5.6.8) is built up of all the physical wave functions and thus describes the whole ensemble. It gives rise to the closed chain (5.6.3), which in turn determines the causal action and the constraints. In this way, the causal action principle becomes a variational principle for the ensemble of physical wave functions. Thus the interaction described by the causal action principle can be understood as a direct mutual interaction of all the physical wave functions. In simple terms, the causal action principle aims at bringing the ensemble of wave functions into an “optimal” configuration. For such optimal configurations, the ensemble of wave functions gives rise to the spacetime structures as we know them: the causal and metric structure, the bosonic fields, and all that.

The last step can be understood more concretely starting from Dirac’s hole theory and the picture of the Dirac sea (for basics see again Section 1.5). In our approach, the Dirac sea is taken literally. Thus all the states of the Dirac sea correspond to physical wave functions. All the information contained in these wave functions induces spacetime with the familiar structures. As a simple example, the bosonic potentials $B$ are determined via the Dirac equation (5.7.3) from the ensemble of wave functions as described by $\tilde{P}(x,y)$. Clearly, on order to make this picture precise, one needs to verify that, in a certain limiting case, the kernel of the fermionic projector corresponding to a minimizer of the causal action principle indeed satisfies a Dirac equation of the form (5.7.3) and thus gives rise to a potential $B$. This will be one of the objectives of the later chapters in this book.

We now discuss which physical principles enter the approach, and how they were incorporated. Causal fermion systems evolved from an attempt to combine several physical principles in a coherent mathematical framework. As a result, these principles appear in the framework in a specific way:

- **The principle of causality**: A causal fermion system gives rise to a causal structure (see Definition 5.5.1). The causal action principle is compatible with this notion of causality in the sense that the pairs of points with spacelike separation do not enter the Euler-Lagrange equations. In simple terms, points with spacelike separation do not interact.

- **The local gauge principle**: Already in the above discussion of how the causal action principle came about, we mentioned that the Lagrangian is gauge invariant in the sense that gauge phases drop out of the Lagrangian (see the explanation after (5.7.6) in Section 5.7). When starting from a general causal fermion system, local gauge principle becomes apparent when representing the physical wave functions in bases of the spin spaces. More precisely, choosing a pseudo-orthonormal basis $(\epsilon_\alpha(x))_{\alpha=1,\ldots,\dim S_x}$
5.8. UNDERLYING PHYSICAL CONCEPTS

of $S_x$, a physical wave function can be represented as

$$\psi(x) = \sum_{\alpha=1}^{\dim S_x} \psi^\alpha(x) \epsilon_\alpha(x)$$

(5.8.1)

with component functions $\psi^1, \ldots, \psi^{\dim S_x}$. The freedom in choosing the basis $(\epsilon_\alpha)$ is described by the group of unitary transformations with respect to the indefinite spin inner product. This gives rise to the transformations

$$\epsilon_\alpha(x) \to \sum_\beta U^{-1}(x)_\alpha^\beta \epsilon_\beta(x) \quad \text{and} \quad \psi^\alpha(x) \to \sum_\beta U(x)_\alpha^\beta \psi^\beta(x)$$

(5.8.2)

with $U \in U(p, q)$. As the basis $(\epsilon_\alpha)$ can be chosen independently at each spacetime point, one obtains local gauge transformations of the wave functions, where the gauge group is determined to be the isometry group of the spin inner product. The causal action is gauge invariant in the sense that it does not depend on the choice of spinor bases.

- The **Pauli exclusion principle** is incorporated in a causal fermion system, as can be seen in various ways. One formulation of the Pauli exclusion principle states that every fermionic one-particle state can be occupied by at most one particle. In this formulation, the Pauli exclusion principle is respected because every wave function can either be represented in the form $\psi^u$ (the state is occupied) with $u \in \mathcal{F}$ or it cannot be represented as a physical wave function (the state is not occupied). Via these two conditions, the fermionic projector encodes for every state the occupation numbers 1 and 0, respectively, but it is impossible to describe higher occupation numbers.

  More technically, one may obtain the connection to the fermionic Fock space formalism by choosing an orthonormal basis $u_1, \ldots, u_f$ of $\mathcal{F}$ and forming the $f$-particle Hartree-Fock state

$$\Psi := \psi^{u_1} \wedge \cdots \wedge \psi^{u_f}.$$

Clearly, the choice of the orthonormal basis is unique only up to the unitary transformations

$$u_i \to \tilde{u}_i = \sum_{j=1}^f U_{ij} u_j \quad \text{with} \quad U \in U(f).$$

Due to the anti-symmetrization, this transformation changes the corresponding Hartree-Fock state only by an irrelevant phase factor,

$$\tilde{\psi}^{\tilde{u}_1} \wedge \cdots \wedge \tilde{\psi}^{\tilde{u}_f} = \det U \, \psi^{u_1} \wedge \cdots \wedge \psi^{u_f}.$$

Thus the configuration of the physical wave functions can be described by a fermionic multi-particle wave function. The Pauli exclusion principle becomes apparent in the total anti-symmetrization of this wave function.

Clearly, the above Hartree-Fock state does not account for quantum entanglement. Indeed, the description of entanglement requires more general Fock space constructions (this will be described in more detail in Chapter 22).

- The **equivalence principle**: Starting from a causal fermion system $(\mathcal{F}, \mathcal{F}, \rho)$, spacetime $M$ is given as the support of the measure $\rho$. Thus spacetime is a topological space (with the topology on $M$ induced by the operator norm on $L(\mathcal{F})$). In situations when spacetime has a smooth manifold structure, one can describe spacetime by choosing coordinates. However, there is no distinguished coordinate systems, giving rise to the freedom of performing general coordinate transformations. The causal action as well
as all the constraints are invariant under such transformations. In this sense, the equivalence principle is implemented in the setting of causal fermion systems. However, other physical principles are missing. For example, the principle of locality is not included. Indeed, the causal action principle is non-local, and locality is recovered only in the continuum limit. Moreover, our concept of causality is quite different from causation (in the sense that the past determines the future) or microlocality (stating that the observables of spacelike separated regions must commute).

Exercises

Exercise 5.1. (A causal fermion system on $\ell_2$) Let $\mathcal{H} = \ell_2$ the Hilbert space of square-summable complex-valued sequences, equipped with the scalar product

$$\langle u | v \rangle = \sum_{i=1}^{\infty} \bar{u}_i v_i, \quad u = (u_i)_{i \in \mathbb{N}}, \quad v = (v_i)_{i \in \mathbb{N}}.$$ 

For any $k \in \mathbb{N}$, let $x_k \in L(\mathcal{H})$ the operator defined by

$$(x_k u)_i := u_{k+1}, \quad (x_k u)_{k+1} := u_k, \quad (x_k u)_i = 0 \text{ for } i \notin \{k, k+1\}.$$ 

In other words,

$$x_k u = (0, \ldots, 0, u_{k+1}, u_k, 0, \ldots)$$

$k - 1$ entries

Finally, let $\mu$ the counting measure on $\mathbb{N}$ (i.e. $\mu(X) = |X|$ equals the cardinality of $X \subset \mathbb{N}$.)

(a) Show that every operator $x_k$ has rank two, is symmetric, and has one positive and one negative eigenvalue. Make yourself familiar with the concept that every operator is a point in $\mathcal{F}$ for spin dimension $n = 1$.

(b) Let $F : \mathbb{N} \to \mathcal{F}$ be the mapping which to every $k$ associates the corresponding operator $x_k$. Show that the push-forward measure $\rho = F_* \mu$ defined by $\rho(\Omega) := \mu(F^{-1}(\Omega))$ defines a measure on $\mathcal{F}$. Show that this measure can also be characterized by

$$\rho(\Omega) = |\{k \in \mathbb{N} \mid x_k \in \Omega\}|.$$ 

(c) Show that $(\mathcal{H}, \mathcal{F}, \rho)$ is a causal fermion system of spin dimension one.

(c) Show that the support of $\rho$ consists precisely of all the operators $x_k$. What is spacetime $M$? What is the causal structure on $M$? What is the resulting causal action?

Exercise 5.2. (On the trace constraint) Let $(\mathcal{H}, \mathcal{F}, \rho)$ be a causal fermion system of spin dimension $n$.

(a) Let $x \in \mathcal{F}$. Show that $\mathcal{L}(x, x) > 0$ whenever $\text{tr}(x) \neq 0$.

(b) Assume that $\int_{\mathcal{F}} \text{tr}(x) \, d\rho(x) \neq 0$. Show that $S(\rho) > 0$.

Exercise 5.3. (Well-posedness of the causal action principle) This exercise explains why the causal action principle is ill-posed if $\dim \mathcal{H} = \infty$ and $\rho(\mathcal{F}) < \infty$.

(a) Let $\mathcal{H}_0$ be a finite-dimensional Hilbert space of dimension $n$ and $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$ be a causal fermion system of finite total volume $\rho_0(\mathcal{F}_0) < \infty$. Let $\iota : \mathcal{H}_0 \to \mathcal{H}$ be an isometric embedding of Hilbert spaces. Construct a causal fermion system $(\mathcal{H}, \mathcal{F}, \rho)$ which has the same action, the same total volume and the same values for the trace and boundedness constraints as the causal fermion system $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$. 


(b) Let $\mathcal{H}_1 := \mathcal{H}_0 \oplus \mathcal{H}_0$. Construct a causal fermion system $(\mathcal{H}_1, \mathcal{F}_1, \rho_1)$ which has the same total volume and the same value of the trace constraint as $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$ but a smaller action and a smaller value of the boundedness constraint.

Hint: Let $F_{1/2} : L(\mathcal{H}_0) \to L(\mathcal{H}_1)$ be the linear mappings
$$F_1(A)(u \oplus v) := (Au) \oplus 0, \quad F_2(A)(u \oplus v) := 0 \oplus (Au).$$
Show that $F_{1/2}$ maps $\mathcal{F}_0$ to $\mathcal{F}_1$. Define the measure $\rho_1$ by
$$\rho_1 = \frac{1}{2}((F_1)_* \rho_0 + (F_2)_* \rho_0).$$

(c) Iterate the construction in (ii) and apply (i) to obtain a series of universal measures on $\mathcal{F}$ of fixed total volume and with fixed value of the trace constraint, for which the action and the values of the boundedness constraint tend to zero. Do these universal measures converge? If yes, what is the limit?

Exercise 5.4. (On the spectrum of the closed chain) This exercise is devoted to analyzing general properties of the spectrum of the closed chain.

(a) Let $x, y$ be symmetric operators of finite rank on a Hilbert space $\mathcal{H}$. Show that there is a finite-dimensional subspace $\mathcal{K} \subset \mathcal{H}$ on which both $x, y$ are invariant. By choosing an orthonormal basis of $\mathcal{K}$ and restricting the operators on $\mathcal{K}$, we may represent $x$ and $y$ by Hermitian matrices. Therefore, the remainder of this exercise is formulated for simplicity in terms of Hermitian matrices.

(b) Show that for any matrix $Z$, the characteristic polynomials of $Z$ and of its adjoint $Z^\dagger$ (being the transposed complex conjugated matrix) are related by complex conjugation, i.e.
$$\det(Z^\dagger - \lambda I) = \overline{\det(Z - \lambda I)}.$$

(c) Let $X$ and $Y$ be symmetric matrices. Show that the characteristic polynomials of the matrices $XY$ and $YX$ coincide.

(d) Combine (ii) and (iii) to conclude that the characteristic polynomial of $XY$ has real coefficients, i.e. $\det(XY - \lambda I) = \overline{\det(XY - \lambda I)}$. Infer that the spectrum of the matrix product $XY$ is symmetric about the real axis, i.e.
$$\det(XY - \lambda I) = 0 \implies \det(XY - \bar{\lambda} I) = 0.$$

Exercise 5.5. (Regular points) Let $x \in \mathcal{F}$ have $p(x) \leq n$ negative and $q(x) \leq n$ positive eigenvalues. The pair $(p(x), q(x))$ is referred to as the signature of $x$. The operator $x$ is said to be regular if $\text{sign}(x) = (n, n)$. The goal of this exercise is to show that the set $\mathcal{F}^{\text{reg}}$ of regular points is open in $\mathcal{F}$. Let us define the positive and negative components of $x$ as the operators
$$x_\pm := \frac{x \pm |x|}{2}, \quad |x| := \sqrt{x^2}.$$

From the functional calculus it follows that $x^2 |x| = |x|x$. Prove the following statements.

(a) Bonus: Let $\{e_i, i = 1, \ldots, m\}$ be an orthogonal set. Any vector set $\{h_i, i = 1, \ldots, m\}$ which fulfills the following condition is linearly independent,
$$\|e_i - h_i\| < \frac{\inf\{\|e_i\|, i = 1, \ldots, m\}}{m} \quad \text{for all } i = 1, \ldots, m.$$

(b) For every $x \in \mathcal{F}$,
$$x(\text{im } x_\pm) \subset \text{im } x_\pm \quad \text{and} \quad x_+ x_- = 0.$$
Moreover, $x|_{\text{im } x_-}$ and $x|_{\text{im } x_+}$ are negative and positive definite, respectively.
(c) Let \( x_0 \in \mathcal{F} \). Then there is an orthonormal set \( \{ e_i \mid i = 1 \ldots \dim S_{x_0} \} \) of eigenvectors of \( x_0 \) such that
\[
\langle e_i | x_0 e_i \rangle < 0 \quad \text{for all } i \leq p(x_0), \quad \langle e_i | x_0 e_i \rangle > 0 \quad \text{for all } p(x_0) < i \leq q(x_0).
\]

(d) The following functions are continuous,
\[
f_i : B_r(x_0) \ni x \mapsto f_i(x) := \begin{cases} 
- e_i & i \leq p(x_0) \\
+ e_i & p(x_0) < i \leq p(x_0) + q(x_0). 
\end{cases}
\]

Hint: You can use the general inequality \( \| A - B \| \leq \| A^2 - B^2 \| \)

(e) There is a \( r > 0 \) such that \( p(x) \geq p(x_0) \) and \( q(x) \geq q(x_0) \) for every \( x \in B_r(x_0) \).

Hint: Use the statements above.

(f) Conclude that \( \mathcal{F}^{reg} \) is an open subset of \( \mathcal{F} \).

Exercise 5.6. (On the spectrum of the closed chain - part 2) Let \( (\mathcal{H}, \mathcal{F}, \rho) \) be a causal fermion system and \( x, y \in \mathcal{F} \). For the closed chain
\[
A_{xy} := P(x, y)P(y, x) : (S_x, \langle \cdot, \cdot \rangle_x) \to (S_y, \langle \cdot, \cdot \rangle_y),
\]
the mathematical setting analyzed in Exercise 3.3 is somewhat different, because \( A_{xy} \) is a symmetric operator on an indefinite inner product space. On the other hand, we know that \( A_{xy} \) is isospectral to \( xy \). Indeed, the symmetry result in Exercise 3.3-(iv) can be used to prove a corresponding statement for the closed chain:
\[
\det(A_{xy} - \lambda I) = 0 \iff \det(A_{xy} - \lambda I) = 0.
\]
This result is well-known in the theory of indefinite inner product spaces. In order to derive it from Exercise 3.3-(iv), one can proceed as follows: First, represent the indefinite inner product in the form \( \langle \cdot, \cdot \rangle = \langle \cdot, S \cdot \rangle \), where \( \langle \cdot, \cdot \rangle \) is a scalar product and \( S \) is an invertible operator which is symmetric (with respect to this scalar product). Next, show that the operator \( B := A_{xy}S \) is symmetric (again with respect to this scalar product). Finally, write the closed chain as \( A_{xy} = BS^{-1} \) and apply Exercise 3.3 (iv).

Exercise 5.7. (A causal fermion system on \( \ell_2 \) - part 2) We return to the example of Exercise 5.3. This time we equip it with a Krein structure.

(a) For any \( k \in \mathbb{N} \), construct the spin space \( S_{x_k} \) and its spin scalar product.

(b) Given a vector \( u \in \mathcal{H} \), what is the corresponding wave function \( \psi^u \)? What is the Krein inner product \( \langle \cdot, \cdot \rangle \)?

(c) What is the topology on the Krein space \( \mathcal{K} \)? Does the wave evaluation operator \( \Psi : u \mapsto \psi^u \) give rise to a well-defined and continuous mapping \( \Psi : \mathcal{H} \to \mathcal{K} \)? If yes, is it an embedding? Is it surjective?

(d) Repeat part (c) of this exercise for the causal fermion system obtained if the operators \( x_k \) are multiplied by \( k \), i.e.
\[
x_k u := (0, \ldots, 0, ku_{k+1}, ku_k, 0, \ldots).
\]

Exercise 5.8. (Stability of the causal structure) A binary relation \( P \) on \( \mathcal{F} \) is said to be stable under perturbations if
\[
(x_0, y_0) \in P \implies \exists r > 0 : B_r(x_0) \times B_r(y_0) \subset P.
\]
Two points \( x, y \in \mathcal{F} \) are said to be properly timelike separated if the closed chain \( A_{xy} \) has a strictly positive spectrum and if all eigenspaces are definite subspaces of \( (S_x, \langle \cdot, \cdot \rangle) \).

(a) Show that proper timelike separation implies timelike separation.
(b) Show by a counterexample with $3 \times 3$ matrices that the notion of timelike separation is not stable under perturbations.

(c) Show that the notion of properly timelike separation is stable under perturbations.

Exercise 5.9. (Time direction) The ability to distinguish between past and future can be described in mathematical terms by the existence of an antisymmetric functional $\mathcal{T} : M \times M \to \mathbb{R}$. One then says that

\[
\begin{cases} 
  y \text{ lies in the future of } x & \text{if } \mathcal{T}(x, y) > 0 \\
  y \text{ lies in the past of } x & \text{if } \mathcal{T}(x, y) < 0.
\end{cases}
\]

Can you think of a simple non-trivial example of such a functional which involves only products and linear combinations of the spacetime operators and the orthogonal projections on the corresponding spin spaces?
CHAPTER 6

Causal Variational Principles

The causal action principle as introduced in Section 5.5 has quite a rich structure and is rather involved. Therefore, it is difficult to analyze it in full generality in one step. It is preferable to begin with special cases and simplified situations, and to proceed from there step by step. In fact, doing so leads to a whole class of variational principles, referred to as causal variational principles. These different variational principles capture different features and aspects of the causal action principle. Proceeding in this way also gives a better understanding of the physical meaning of the different structures of a causal fermion system and of the interaction as described by the causal action principle. We now give an overview of the different settings considered so far. This has the advantage that in the later chapters of this book, we can always work in the setting which is most suitable for the particular question in mind. Moreover, for pedagogical reasons, in this book we shall sometimes idealize the setting for example by assuming for technical simplicity that the Lagrangian is smooth.

6.1. The Causal Variational Principle on the Sphere

Clearly, the trace constraint (5.5.4) and the boundedness constraint (5.5.5) complicate the analysis. Therefore, it might be a good idea to consider a simplified setting where these constraints are not needed. This can be accomplished most easily by prescribing the eigenvalues of the operators in \( F \). This method was first proposed in [43], Section 2, in a slightly different formulation. We now explain the method in a way which fits best to our setting. Given \( n \in \mathbb{N} \), we choose real numbers \( \nu_1, \ldots, \nu_{2n} \) with

\[
\nu_1 \leq \cdots \leq \nu_n \leq 0 \leq \nu_{n+1} \leq \cdots \leq \nu_{2n}.
\]

(6.1.1)

We let \( F \) be the set of all symmetric operators on \( \mathcal{H} \) of rank \( 2n \) whose eigenvalues (counted with multiplicities) coincide with \( \nu_1, \ldots, \nu_{2n} \). If \( \mathcal{H} \) is finite-dimensional, the set \( F \) is compact. This is the reason why it is sensible to minimize the causal action (5.5.2) keeping only the volume constraint (5.5.3), which for simplicity we implement by restricting attention to normalized measures,

\[
\rho(F) = 1.
\]

(6.1.2)

The simplest interesting case is obtained by choosing the spin dimension \( n = 1 \) and the Hilbert space \( \mathcal{H} = \mathbb{C}^2 \). We denote the eigenvalues \( \nu_1 \) and \( \nu_2 \) in (6.1.1) by

\[
\nu_{1/2} = 1 \pm \tau \quad \text{with} \quad \tau > 1.
\]

Then \( F \) consists of all Hermitian \( 2 \times 2 \)-matrices \( F \) which have eigenvalues \( \nu_1 \) and \( \nu_2 \). Such a matrix can be represented with the help of Pauli matrices as

\[
F = \tau \bar{x} \sigma + 1 \quad \text{with} \quad \bar{x} \in S^2 \subset \mathbb{R}^3.
\]

Thus the set \( F \) can be identified with the unit sphere \( S^2 \).

The volume constraint (5.5.3) can be implemented most easily by restricting attention to normalized regular Borel measures on \( F \) (i.e. measures with \( \rho(F) = 1 \)). Such a measure
can be both continuous, discrete or a mixture. Examples of continuous measures are obtained by multiplying the Lebesgue measure on the sphere by a non-negative smooth function. By a discrete measure, on the other hand, we here mean a weighted counting measure, i.e. a measure obtained by inserting weight factors into (5.2.2),

\[
\rho = \sum_{i=1}^{L} c_i \delta_{x_i} \quad \text{with} \quad x_i \in \mathcal{F}, \quad c_i \geq 0 \quad \text{and} \quad \sum_{i=1}^{L} c_i = 1 . \quad (6.1.3)
\]

In this setting, a straightforward computation yields for the Lagrangian (5.5.1)

\[
L(x, y) = \max \left(0, D(x, y)\right) \quad \text{with} \quad D(x, y) = 2\tau^2 \left(1 + \langle x, y \rangle\right) \left(2 - \tau^2 \left(1 - \langle x, y \rangle\right)\right) , \quad (6.1.4)
\]

and \(\langle x, y \rangle\) denotes the Euclidean scalar product of two unit vectors \(x, y \in S^2 \subset \mathbb{R}^3\) (thus \(\langle x, y \rangle = \cos \vartheta\), where \(\vartheta\) is the angle between \(x\) and \(y\)).

The resulting so-called causal variational principle on the sphere was introduced in \[43\], Chapter 1] and analyzed in \[79\], Sections 2 and 5] and more recently in \[10\]. We now explain a few results from these papers.

First of all, the causal variational principle on the sphere is well-posed, meaning that the minimum is attained in the class of all normalized regular Borel measures. The minimizing measure is not unique; indeed, there are typically many minimizers. The study in \[79\], Section 2] gives the following numerical result: If \(\tau > \sqrt{2}\), every minimizing measure is a weighted counting measure (6.1.3).

Thus, although we minimize over all regular Borel measures (i.e. measures which can have discrete and continuous components), a minimizing measure always describes a discrete spacetime consisting of a finite number of spacetime points. This result can be interpreted physically as an indication that the causal action principle should give rise to discrete spacetime structures. More details on the numerical findings and the physical interpretation can be found in the review \[57\], Section 4].

To some extent, the above numerical result could be underpinned by analytic results. First, it was proven in \[79\], Section 5.1] that the support has an empty interior:

**Theorem 6.1.1.** If \(\tau > \sqrt{2}\), the support of any minimizing measure does not contain an open subset of \(S^2\).

Intuitively speaking, this result shows that the spacetime points are a subset of the sphere of dimension strictly smaller than two. More recently, it was shown in \[10\] that the dimension of the support is even strictly smaller than one:

**Theorem 6.1.2.** In the case \(\tau > \sqrt{6}\), the support of any minimizing measure is totally disconnected and has Hausdorff dimension at most 6/7.

### 6.2. Causal Variational Principles in the Compact Setting

Generalizing the causal variational principle on the sphere, one can replace \(\mathcal{F}\) by a smooth compact manifold of dimension \(m \geq 1\). The Lagrangian \(L\) can be defined in analogy to (6.1.4) as the positive part of a smooth function \(D\), i.e.

\[
L(x, y) = \max \left(0, D(x, y)\right) \quad \text{with} \quad D \in C^\infty(\mathcal{F} \times \mathcal{F}, \mathbb{R}) . \quad (6.2.1)
\]

We usually assume that
(i) \( \mathcal{L} \) is symmetric: \( \mathcal{L}(x,y) = \mathcal{L}(y,x) \) for all \( x,y \in \mathcal{F} \).

(ii) \( \mathcal{L} \) is strictly positive on the diagonal: \( \mathcal{L}(x,x) > 0 \) for all \( x \in \mathcal{F} \).

The causal variational principle is to minimize the causal action

\[
\mathcal{S} = \int_{\mathcal{F}} d\rho(x) \int_{\mathcal{F}} d\rho(y) \mathcal{L}(x,y)
\]

under variations of the measure \( \rho \) in the class of all regular Borel measures on \( \mathcal{F} \) which are normalized, i.e.

\[
\rho(\mathcal{F}) = 1.
\]

This so-called compact setting was introduced in [79 Section 1.2]. It is the preferable choice for studying phenomena for which the detailed form of the Lagrangian as well as the constraints of the causal action principle are irrelevant.

Giving a minimizing measure \( \rho \), the Lagrangian induces on spacetime \( M := \text{supp} \rho \) a causal structure. Namely, two spacetime points \( x, y \in M \) are said to be timelike and spacelike separated if \( \mathcal{L}(x,y) > 0 \) and \( \mathcal{L}(x,y) = 0 \), respectively. But of course, compared to the causal action principle, spin spaces and physical wave functions (as defined in Section 5.6) are missing in this setting.

We point out that in (6.2.1) we merely assumed that the function \( D \) is smooth. The Lagrangian, however, is Lipschitz continuous. It is in general non-differentiable on the boundary of the light cone as defined by the level set \( \mathcal{D}(x,y) = 0 \). In order to avoid differentiability issues, it is sometimes useful to make the assumptions even stronger by assuming that the Lagrangian itself is smooth,

\[
\mathcal{L} \in C^\infty(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+).
\]

This is the so-called smooth setting. Clearly, the Lagrangian of the causal action (5.5.1) is not smooth, and this fact is indeed responsible for interesting effects like the results on the singular support in [79, 10]. In view of these results, the smoothness assumption (6.2.4) is a mathematical idealization.

The assumptions that \( \mathcal{F} \) is a smooth manifold and that the function \( \mathcal{D} \) in (6.2.1) is smooth are convenient and avoid certain technicalities. But these assumptions much more than what is needed for the analysis. More generally, one can choose \( \mathcal{L} \) as a non-negative continuous function,

\[
\mathcal{L}(x,y) \in C^0(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+).
\]

Going one step further, one may relax the continuity of the Lagrangian by the condition that \( \mathcal{L} \) be lower semi-continuous, i.e. that for all sequences \( x_n \to x \) and \( y_{n'} \to y \),

\[
\mathcal{L}(x,y) \leq \lim \inf_{n,n' \to \infty} \mathcal{L}(x_n,y_{n'}).
\]

Since the Lagrangian of the causal action (5.5.1) is continuous, lower semi-continuity is an unphysical generalization. Nevertheless, this setting is useful for two reasons: First, because from the point of view of calculus of variations, it is a natural generalization to which most methods still apply. And second, because lower semi-continuous Lagrangians are convenient for formulating explicit examples (like the lattice model in [68 Section 5]).

We finally note that also the assumption of \( \mathcal{F} \) being a smooth manifold can be relaxed. From the point of view of calculus of variations, the right setting is to assume that \( \mathcal{F} \) is a compact topological Hausdorff space.

In this book, for pedagogical reasons we do not aim for the highest generality and minimal smoothness and regularity assumptions. An introduction to a more general and more abstract setting can be found in [72 Section 3].
6.3. Causal Variational Principles in the Non-Compact Setting

In the compact setting, spacetime $M$ is a compact subset of $\mathcal{F}$. This is not suitable for describing situations when spacetime has an asymptotic future or past or when spacetime has singularities (like at the big bang or inside a black hole). For studying such situations, it is preferable to work in the so-called *non-compact setting* introduced in [68, Section 2.1].

Thus we now let $\mathcal{F}$ be a non-compact smooth manifold of dimension $m \geq 1$. We choose the Lagrangian again in the form (6.2.1) or (6.2.5), again with the properties (i) and (ii) on page 87. In the non-compact setting, it is not sensible to work with normalized measures. Instead, the total volume $\rho(\mathcal{F})$ is infinite. In order to introduce a well-defined variational principle, we let $\tilde{\rho}$ be another regular Borel measure on $\mathcal{F}$ which satisfies the conditions

$$|\tilde{\rho} - \rho|_\mathcal{F} < \infty \quad \text{and} \quad (\tilde{\rho} - \rho)(\mathcal{F}) = 0$$

(6.3.1)

(where $|.|$ denotes the total variation of a signed measure; see Section 2.3). We then consider the difference of the actions defined by

$$S(\tilde{\rho}) - S(\rho) = \int_\mathcal{F} d(\tilde{\rho} - \rho)(x) \int_\mathcal{F} d\rho(y) \mathcal{L}(x,y)$$

$$+ \int_\mathcal{F} d\rho(x) \int_\mathcal{F} d(\tilde{\rho} - \rho)(y) \mathcal{L}(x,y) + \int_\mathcal{F} d(\tilde{\rho} - \rho)(x) \int_\mathcal{F} d(\tilde{\rho} - \rho)(y) \mathcal{L}(x,y).$$

(6.3.2)

The measure $\rho$ is said to be a *minimizer* of the causal action if this difference is non-negative for all $\tilde{\rho}$ satisfying (6.3.1),

$$S(\tilde{\rho}) - S(\rho) \geq 0.$$

(6.3.3)

Exactly as mentioned at the end of the previous section, the assumption that $\mathcal{F}$ is a smooth manifold can be weakened. From the point of view of calculus of variations, the right setting is to assume that $\mathcal{F}$ is a $\sigma$-locally compact topological Hausdorff space (for details see again [72, Section 3]).

6.4. The Local Trace is Constant

Causal variational principles as introduced in the previous sections are of interest in their own right as a novel class of nonlinear variational principles. But, being primarily interested in causal fermion systems, it is an important task for us to get a concise mathematical connection to the causal action principle. In preparation, we now analyze the trace constraint and derive a first result on minimizing measures of the causal action principle. We do this at such an early stage of this book is because the result can be used to simplify the setup of causal fermion systems, getting the desired connection to causal variational principles (see Section 6.5 below). The following result was first obtained in [13] (albeit with a different method); see also [48, Proposition 1.4.1]. For technical simplicity, we here only consider the finite-dimensional setting.

**Proposition 6.4.1.** Consider the causal variational principle in the finite-dimensional setting $\dim \mathcal{F} < \infty$. Let $\rho$ be a minimizer of finite total volume, $\rho(\mathcal{F}) < \infty$. Then there is a real constant $c$ such that

$$\text{tr}(x) = c \quad \text{for all} \quad x \in M.$$

(6.4.1)

We often refer to $\text{tr}(x)$ as the *local trace* at the point $x$. 
Proof of Proposition 6.4.1. The method is to construct a suitable variation of the measure
\[(\rho_\tau)_{\tau \in (-\delta, \delta)} \quad \text{with} \quad \rho_0 = \rho\]
(for some \(\delta > 0\)). The variation must satisfy the constraints. Then we can make use of the fact that \(\rho\) is a minimizer.

For the construction we combine two different methods. One method is to multiply \(\rho\) by a positive function \(f_\tau : M \to \mathbb{R}^+\),

\[\rho_\tau = f_\tau \rho. \quad (6.4.2)\]

Clearly, this variation does not change the support of the measure. In order to change the support, one can consider a function \(F_\tau : M \to \mathcal{F}\) and take the push-forward measure,

\[\rho_\tau = (F_\tau)_* \rho. \quad (6.4.3)\]

Combining these methods, we are led to considering the family of measures

\[\rho_\tau = (F_\tau)_* (f_\tau \rho). \quad (6.4.4)\]

The condition \(\rho_0 = \rho\) gives rise to the conditions

\[f_0 \equiv 1 \quad \text{and} \quad F_0 \equiv 1. \quad (6.4.5)\]

Finally, we assume that the functions \(f_\tau\) and \(F_\tau\) are defined and smooth in \(\tau\).

The ansatz (6.4.4) is particularly convenient for computations. Namely, by definition of the push-forward measure,

\[\int_\mathcal{F} \mathcal{L}(x, y) \, d\rho(y) = \int_M \mathcal{L}(x, F_\tau(y)) \, f_\tau(y) \, d\rho(y),\]

and similarly for all other integrals. Next we choose \(F_\tau\) in such a way that this integral does not change. To this end we choose

\[F_\tau(x) = \frac{x}{\sqrt{f_\tau(x)}}. \quad (6.4.6)\]

Using that \(\mathcal{L}(x, y)\) is homogeneous in \(y\) of degree two, it follows that

\[\int_M \mathcal{L}(x, F_\tau(y)) \, f_\tau(y) \, d\rho(y) = \int_M \mathcal{L}\left(x, \frac{y}{\sqrt{f_\tau(y)}}\right) \, f_\tau(y) \, d\rho(y)\]

\[= \int_M \mathcal{L}(x, y) \frac{1}{f_\tau(y)} \, f_\tau(y) \, d\rho(y) = \int_M \mathcal{L}(x, y) \, d\rho(y).\]

Arguing similarly in the variable \(x\), one sees that our variation does not change the action. Since the expression \(|xy|^2\) is also homogeneous of degree two in \(x\) and \(y\), the above argument also applies to the functional \(\mathcal{T}\), showing that the boundedness constraint is also respected by our variation.

Let us analyze the constraints. In order to satisfy the volume constraint, we make the ansatz

\[f_\tau = 1 + \tau g, \quad (6.4.7)\]

where \(g\) is a bounded function with zero mean,

\[\int_M g(x) \, d\rho(x) = 0 \quad \text{for all} \ \tau \in (-\delta, \delta). \quad (6.4.8)\]
We finally consider the variation of the trace constraint:

\[
\int_\mathcal{F} \text{tr}(x) \, d(\rho_\tau - \rho)(x) = \int_M \text{tr}(F_\tau(x)) \, f_\tau(x) \, d\rho(x) - \int_M \text{tr}(x) \, d\rho(x) \\
= \int_M \text{tr} \left( \frac{x}{\sqrt{f_\tau(x)}} \right) \, f_\tau(x) \, d\rho(x) - \int_M \text{tr}(x) \, d\rho(x) \\
= \int_M \text{tr} \left( \frac{x}{\sqrt{f_\tau(x)}} \right) \, f_\tau(x) \, d\rho(x).
\]

Employing again the ansatz (6.4.7) and differentiating with respect to \(\tau\), we obtain for the first variation

\[
\frac{d}{d\tau} \int_\mathcal{F} \text{tr}(x) \, d(\rho_\tau - \rho)(x) \bigg|_{\tau=0} = \frac{1}{2} \int_M \text{tr}(x) \, g(x) \, d\rho(x).
\]

If the local trace is constant, then the last integral vanishes in view of (6.4.8). This shows that for the considered first variation, all the constraints are satisfied and the action is stationary. Conversely, if the local trace is not constant, then by choosing \(g\) appropriately, one can arrange that the first variation of the local trace is strictly positive,

\[
\frac{d}{d\tau} \int_\mathcal{F} \text{tr}(x) \, d(\rho_\tau - \rho)(x) \bigg|_{\tau=0} =: d > 0,
\]

whereas the other constraints are satisfied and the action is again stationary. Transforming the measures according to

\[
\rho_\tau \rightarrow (G_\tau)_*(\rho_\tau)
\]

with

\[
G_\tau(x) = x \left( \int_K \text{tr}(x) \, d\rho \right) / \left( \int_K \text{tr}(x) \, d\rho + \tau \, d \right), \tag{6.4.9}
\]

the trace constraint is satisfied to first order. Moreover, as the scaling factor in (6.4.9) is smaller than one, the first variation of the action and of the boundedness constraint are strictly negative, in contradiction to the fact that \(\rho\) is a minimizer. This concludes the proof. \(\square\)

### 6.5. How the Causal Action Principle Fits into the Non-Compact Setting

Under mild technical assumptions on the minimizing measure, the causal action principle for causal fermion systems is a special case of the causal variational principle in the non-compact setting, as we now explain.

Since for minimizers of the causal action principle, all operators in \(M\) have the same trace (see Proposition 6.4.1), we can simplify the setting by restricting attention to linear operators in \(\mathcal{F}\) which have constant trace. Then the trace constraint can be disregarded. We now implement this simplification by modifying our setting. At the same time, we implement the boundedness constraint by a Lagrange multiplier term. Here we apply this method naively by modifying the Lagrangian to

\[
\mathcal{L}_\kappa(x, y) := \frac{1}{4^n} \sum_{i,j=1}^{2n} \left( |\lambda^x_i|^2 - |\lambda^y_j|^2 \right)^2 + \kappa \left( \sum_{i=1}^{2n} |\lambda^x_i|^2 \right)^2, \tag{6.5.1}
\]

where \(\kappa > 0\) is the Lagrange multiplier. The justification for this procedure as given in [13] is a bit subtle, and for brevity we shall not enter these constructions here. It is important to note that, in contrast to the usual Lagrange multiplier, where a minimizer
under constraints merely gives rise to a critical point of the Lagrangian including the Lagrange multipliers, here we obtain again a minimizer of the effective action (for details see [13, Theorem 3.13]).

Finally, we make a mild technical simplification. A spacetime point \( x \in M \) is said to be regular if \( x \) has the maximal possible rank \( 2n \). Otherwise, the spacetime point is singular. In physical applications, all spacetime points are regular, except maybe at singularities like the center of black holes. With this in mind, we here assume that the causal fermion system is regular in the sense that all spacetime points are regular. This assumption has the advantage that the set of all regular points of \( \mathcal{F} \) is a smooth manifold (see Proposition 3.1.3 in Section 3.1). We remark that, in the case that \( \mathcal{H} \) is infinite-dimensional, the set of regular points of \( \mathcal{F} \) can be endowed with the structure of a Banach manifold (for details see [73, Section 3]).

These considerations lead us to the following setting: Let \( (\mathcal{H}, \langle .|.| \rangle_{\mathcal{H}}) \) be a complex Hilbert space. Moreover, we are given parameters \( n \in \mathbb{N} \) (the spin dimension), \( c > 0 \) (the constraint for the local trace) and \( \kappa > 0 \) (the Lagrange multiplier of the boundedness constraint). We let \( \mathcal{F}_{\text{reg}} \subset \mathcal{L}(\mathcal{H}) \) be the set of all symmetric operators \( F \) on \( \mathcal{H} \) with the following properties:

(a) \( F \) has finite rank and (counting multiplicities) has \( n \) positive and \( n \) negative eigenvalues.

(b) The local trace is constant, i.e.

\[
\text{tr}(F) = c.
\]

On \( \mathcal{F}_{\text{reg}} \) we again consider the topology induced by the sup-norm on \( \mathcal{L}(\mathcal{H}) \). We assume that \( \rho \) is a minimizer of the causal action principle, meaning that the action

\[
\mathcal{S}_\kappa(\rho) = \int_{\mathcal{F} \times \mathcal{F}} \mathcal{L}_\kappa(x, y) \, d\rho(x) \, d\rho(y)
\]

is minimal under variations which preserve the total volume. In this way, the causal action principle fits into the framework of causal variational principles in the non-compact setting as introduced in Section 6.4.3 In agreement with (6.2.5), the causal Lagrangian is continuous (in fact, it is even locally Hölder continuous; for details see [73, Section 5.1]). Moreover, it has the desired properties (i) and (ii) on page 87 (it is strictly positive because the Lagrangian can be estimated from below in terms of the local trace; see Exercise 6.2).

In order to avoid misunderstandings, we point out that the above description of causal fermion systems by measures on \( \mathcal{F}_{\text{reg}} \) is not a suitable setting for the existence theory. The reason is that \( \mathcal{F}_{\text{reg}} \) is not closed in \( \mathcal{F} \), because the boundary points in \( \mathcal{F} \) are missing. As a consequence, considering a minimizing sequence \( (\rho_n)_{n \in \mathbb{N}} \) of measures in \( \mathcal{F}_{\text{reg}} \), the limiting measure might well be supported also on \( \mathcal{F} \setminus \mathcal{F}_{\text{reg}} \). For this reason, there is no existence theory for measures on \( \mathcal{F}_{\text{reg}} \). But if a minimizing measure is given, it seems sensible to assume that the resulting causal fermion system is regular. Under this assumption, the analysis of the causal fermion system can be carried out exclusively in \( \mathcal{F}_{\text{reg}} \), whereas \( \mathcal{F} \) is no longer needed. For a convenient and compact notation, in such situations we shall even omit the superscript “reg”, so that \( \mathcal{F} \) denotes the set of all symmetric operators on \( \mathcal{H} \) with the above properties (a) and (b). Moreover, we shall omit the subscript \( \kappa \). Thus, with a slight abuse of notation, we shall denote the Lagrangian including the Lagrange multiplier term (6.5.1) again by \( \mathcal{L} \).
Exercises

Exercise 6.1. (support of a measure)

(a) We return to the example of Exercise ??.
Show that the support of \( \rho \) consists precisely of all the operators \( x_k \).

(b) In order to illustrate how to encode geometric information in the support of a measure, let \( \mathcal{M} \subset \mathbb{R}^3 \) be a smooth surface described in a parametrization \( \Phi \). Thus given an open subset \( \Omega \subset \mathbb{R}^2 \), we consider a smooth injective map

\[
\Phi : \Omega \rightarrow \mathbb{R}^3
\]

with the property that \( D\Phi|_p : \mathbb{R}^2 \rightarrow \mathbb{R}^3 \) has rank two for all \( p \in \Omega \). Then the surface \( \mathcal{M} \) is defined as the image \( \Phi(\Omega) \subset \mathbb{R}^3 \).

We now introduce the measure \( \rho \) as the push-forward measure of the Lebesgue measure on \( \mathbb{R}^2 \): Let \( \mu \) be the Lebesgue measure on \( \mathbb{R}^2 \). We define a set \( U \subset \mathbb{R}^3 \) to be \( \rho \)-measurable if and only if its preimage \( \Phi^{-1}(U) \subset \mathbb{R}^2 \) is \( \mu \)-measurable. On the \( \rho \)-measurable sets we define the measure \( \rho \) by

\[
\rho(U) = \mu(\Phi^{-1}(U)).
\]

Verify that the \( \rho \)-measurable sets form a \( \sigma \)-algebra, and that \( \rho \) is a measure. What are the sets of \( \rho \)-measure zero? What is the support of the measure \( \rho \)?

Suppose that \( \Phi \) is no longer assumed to be injective. Is \( \rho \) still a well-defined measure? Is \( \rho \) well-defined if \( \Phi \) is only assumed to be continuous? What are the minimal regularity assumptions on \( \Phi \) needed for the push-forward measure to be well-defined? What is the support of \( \rho \) in this general setting?

Exercise 6.2. This exercise shows that the trace constraint ensures that the action is non-zero.

Let \( (\mathcal{H}, \mathcal{F}, \rho) \) be a causal fermion system of spin dimension \( n \).

(a) Assume that \( \text{tr}(x) \neq 0 \). Show that \( L(x,x) > 0 \). (For a quantitative statement of this fact in the setting of discrete spacetimes see [41, Proposition 4.3].)

(b) Assume that \( \int \text{tr}(x) \, d\rho \neq 0 \). Show that \( S(\rho) > 0 \).

Exercise 6.3. (The two-dimensional Dirac sphere)

We let \( M = S^2 \subset \mathbb{R}^3 \) with \( d\mu \) the surface area measure, normalized such that \( \mu(S^2) = 1 \). Furthermore, we choose \( n = 1 \) and \( f = 2 \). For a given parameter \( \tau > 1 \) we introduce the mapping \( F : M \rightarrow \mathcal{F} \) by

\[
F(\vec{x}) = \tau \vec{x} \vec{\sigma} + \mathbf{1}.
\]

Then \( \sigma(F(\vec{x})) = \{1+\tau, 1-\tau\} \), and thus \( F(\vec{x}) \) has one positive and one negative eigenvalue. Furthermore, a symmetry argument shows that the identity constraint (C2) is satisfied.

Using identities between Pauli matrices

\[
\sigma^i \sigma^j = \delta^{ij} + i \epsilon^{ijk} \sigma^k,
\]

one obtains

\[
F(\vec{x}) F(\vec{y}) = (1 + \tau^2 \vec{x} \vec{y}) \mathbf{1} + \tau (\vec{x} + \vec{y}) \vec{\sigma} + i \tau^2 (\vec{x} \wedge \vec{y}) \vec{\sigma}.
\]

A straightforward calculation yields for the eigenvalues of this matrix

\[
\lambda_{\pm} = 1 + \tau^2 \cos \vartheta \pm \tau \sqrt{1 + \sin \vartheta} \sqrt{2 - \tau^2 (1 - \cos \vartheta)},
\]

where \( \vartheta \) denotes the angle \( \vartheta \) between \( \vec{x} \) and \( \vec{y} \). If \( \vartheta \) is sufficiently small, the term \( (1 - \cos \vartheta) \) is close to zero, and thus the arguments of the square roots are all positive. However, if \( \vartheta \) becomes so large that

\[
\vartheta \geq \vartheta_{\text{max}} := \arccos \left(1 - \frac{2}{\tau^2}\right),
\]

then the action becomes non-zero.
Using this formula in (5.5.2), we can carry out the integrals to obtain

\[ s = \frac{41}{4}, \]

which simplifies to

\[ S(\lambda_1, \lambda_2) \]

shows that if the \( \lambda_{i/2} \) are both real, then they have the same sign. Hence the Lagrangian simplifies to

\[ \mathcal{L}[A_{xy}] = \mathcal{L} = \frac{\lambda_1 - \lambda_2}{2} \Theta(\vartheta_{\max} - \vartheta) \\
= 2\tau^2 (1 + \cos \vartheta) (2 - \tau^2 (1 - \cos \vartheta)) \Theta(\vartheta_{\max} - \vartheta). \]

Using this formula in (5.5.2), we can carry out the integrals to obtain

\[ S[F] = \frac{1}{2} \int_0^{\vartheta_{\max}} \mathcal{L} \sin \vartheta \, d\vartheta = 4 - \frac{4}{3\tau^2}. \]

Similarly, the functional \( T \) can be computed to be

\[ T[F] = 4\tau^2 (\tau^2 - 2) + 12 - \frac{8}{3\tau^2}. \]

Hence the action (6.5.5) is bounded uniformly in \( \tau \), although the function \( F \), (6.5.2), as well as the functional \( T \), (6.5.6), diverge as \( \tau \to \infty \).

**Exercise 6.4.** This exercise explains why the causal action principle is ill-posed in the case \( \dim \mathcal{H} = \infty \) and \( \rho(\mathcal{F}) < \infty \). The underlying estimates were first given in the setting of discrete spacetimes in [41, Lemma 5.1].

(a) Let \( \mathcal{H}_0 \) be a finite-dimensional Hilbert space of dimension \( n \) and \((\mathcal{H}_0, \rho_0, \mathcal{F}_0)\) be a causal fermion system of finite total volume \( \rho_0(\mathcal{F}_0) \). Let \( \iota : \mathcal{H}_0 \to \mathcal{H} \) be an isometric embedding. Construct a causal fermion system \((\mathcal{H}, \rho, \mathcal{F})\) which has the same action, the same total volume and the same values for the trace and boundedness constraints as the causal fermion system \((\mathcal{H}_0, \rho_0, \mathcal{F}_0)\).

(b) Let \( \mathcal{H}_1 = \mathcal{H}_0 \oplus \mathcal{H}_0 \). Construct a causal fermion system \((\mathcal{H}_1, \rho_1, \mathcal{F}_1)\) which has the same total volume and the same value of the trace constraint as \((\mathcal{H}_0, \rho_0, \mathcal{F}_0)\) but a smaller action and a smaller value of the boundedness constraint. **Hint:** Let \( F_{1/2} : L(\mathcal{H}_0) \to L(\mathcal{H}_1) \) be the linear mappings

\[ (F_1(A))(u \oplus v) = (Au) \oplus 0, \quad (F_2(A))(u \oplus v) = 0 \oplus (Av). \]

Show that \( F_{1/2} \) map \( \mathcal{F}_0 \) to \( \mathcal{F}_1 \). Define \( \rho_1 \) by

\[ \rho_1 = \frac{1}{2} ((F_1)_* \rho + (F_2)_* \rho). \]

(c) Iterate the construction in (b) and apply (a) to obtain a series of measures on \( \mathcal{F} \) of fixed total volume and with fixed value of the trace constraint, for which the action and the values of the boundedness constraint tend to zero. Do these measures converge? If yes, what is the limit?

**Exercise 6.5.** (Embedding of \( S_\ast \mathcal{M} \) into \( S_F(x) \)) The goal of this exercise is to explain how the fibers of the spinor bundle \( S_\ast \mathcal{M} \) are related to the spin spaces \( S_x \) of the corresponding causal fermion system. In order to keep the setting as simple as possible, we let \((\mathcal{M}, g)\) be Minkowski space and \( \mathcal{H} \) a finite-dimensional subspace of the Dirac solution space \( \mathcal{H}_m \), consisting of smooth wave functions of spatially compact support, i.e.,

\[ \mathcal{H} \subset C^\infty(\mathcal{M}, S_\ast \mathcal{M}) \cap \mathcal{H}_m \quad \text{finite-dimensional}. \]
We again let $F(x)$ be the local correlation operators, i.e.

$$\langle \psi | F(x) \psi \rangle = -\langle \psi(x) | \phi(x) \rangle \quad \text{for all } \psi, \phi \in \mathcal{H}$$

(since $\mathcal{H}$ consists of smooth functions, we may leave out the regularization operators).

Defining the measure again by $d\rho = F^\star (d^4x)$, we again obtain a causal fermion system of spin dimension $n = 2$. We next introduce the evaluation map $e_x$ by

$$e_x : \mathcal{H} \to S_x \mathcal{M}, \quad e_x(\psi) = \psi(x).$$

Restricting the evaluation mapping to the spin space $S_{F(x)}$ at the spacetime point $F(x)$ (defined as in the lecture as the image of the operator $F(x)$), we obtain a mapping

$$e_x|_{S_{F(x)}} : S_{F(x)} \to S_x \mathcal{M}.$$

(a) Show that $e_x|_{S_{F(x)}}$ is an isometric embedding.

(b) Show that for all $u \in \mathcal{H}$ and $x \in \mathcal{M}$,

$$e_x|_{S_{F(x)}}(\psi^u(F(x))) = u(x).$$

**Exercise 6.6. (Identification of $S \mathcal{M}$ with $S \mathcal{M}$)** In the setting of the previous exercise, we now make two additional assumptions:

(i) The mapping $F : \mathcal{M} \to \mathcal{F}$ is injective and its image is closed in $\mathcal{F}$.

(ii) The resulting causal fermion system is regular in the sense that for all $x \in \mathcal{M}$, the operator $F(x)$ has rank $2n$.

Using the results of the previous exercise, explain how the following objects can be identified:

(a) $x$ with $F(x)$

(b) $\mathcal{M}$ with $M$

(c) The spinor space $S_x \mathcal{M}$ with the corresponding spin space $S_{F(x)}$

(d) $u \in \mathcal{H}$ with its corresponding physical wave function $\psi^u$

**Exercise 6.7. (The space $C^0(M, SM)$)** In the lecture we introduced a wave function $\psi$ as a mapping from $M$ to $H$ such that $\psi(x) \in S_x$ for all $x \in M$. It is most convenient to define continuity of a wave function by the requirement that for all $x \in M$ and for every $\varepsilon > 0$ there is $\delta > 0$ such that

$$\left\| \sqrt{|y|} \psi(y) - \sqrt{|x|} \psi(x) \right\|_{\mathcal{H}} < \varepsilon \quad \text{for all } y \in M \text{ with } \|y - x\| \leq \delta.$$

Show that, using this definition, every physical wave function is continuous. Thus, denoting the space of continuous wave functions by $C^0(M, SM)$, we obtain an embedding

$$\mathcal{H} \hookrightarrow C^0(M, SM).$$

**Hint:** You may use the inequality

$$\left\| \sqrt{|y|} - \sqrt{|x|} \right\| \leq \|y - x\|^\frac{1}{4} \|y + x\|^\frac{3}{4}.$$

**Exercise 6.8.** This exercise is devoted to analyzing general properties of the spectrum of the closed chain.

(a) As in Definition 5.3.1, we let $x$ and $y$ be symmetric operators of finite rank on a Hilbert space $(\mathcal{H}, \langle . \rangle_{\mathcal{H}})$. Show that there is a finite-dimensional subspace $I \subset \mathcal{H}$ on which both $x$ and $y$ are invariant. By choosing an orthonormal basis of $I$ and restricting the operators to $I$, we may represent both $x$ and $y$ by Hermitian matrices. Therefore, the remainder of this exercise is formulated for simplicity in terms of Hermitian matrices.
(b) Show that for any matrix $Z$, the characteristic polynomials of $Z$ and of its adjoint $Z^*$ (being the transposed complex conjugate matrix) are related by complex conjugation, i.e. $\det(Z^* - \lambda 1) = \overline{\det(Z - \lambda 1)}$.

(c) Let $X$ and $Y$ be symmetric matrices. Show that the characteristic polynomials of the matrices $XY$ and $YX$ coincide.

(d) Combine (b) and (c) to conclude that the characteristic polynomial of $XY$ has real coefficients, i.e. $\det(XY - \lambda 1) = \overline{\det(YX - \lambda 1)}$. Infer that the spectrum of the matrix product $XY$ is symmetric about the real axis, i.e.

$$\det(XY - \lambda 1) = 0 \implies \det(XY - \overline{\lambda} 1) = 0.$$  \hspace{1cm} (6.5.7)

(e) For the closed chain (6.5.3), the mathematical setting is somewhat different, because $A_{xy}$ is a symmetric operator on the indefinite inner product space $(S_x, \langle . | . \rangle_x)$. On the other hand, we saw after (5.6.3) that $A_{xy}$ is isospectral to $xy$. Indeed, the symmetry result (6.5.7) can be used to prove a corresponding statement for $A_{xy}$,

$$\det(A_{xy} - \lambda 1) = 0 \implies \det(A_{xy} - \overline{\lambda} 1) = 0.$$  \hspace{1cm} (6.5.8)

This result is well-known in the theory of indefinite inner product spaces (see for example the textbooks [16], [89] or [41] Section 3]). In order to derive it from (6.5.7), one can proceed as follows: First, represent the indefinite inner product in the form $\langle . | . \rangle = \langle . | S x \rangle$, where $\langle . | . \rangle$ is a scalar product and $S$ is an invertible operator which is symmetric (with respect to this scalar product). Next, show that the operator $B := A_{xy} S$ is symmetric (again with respect to this scalar product). Finally, write the closed chain as $A_{xy} = BS^{-1}$ and apply (6.5.7).
CHAPTER 7

The Euler-Lagrange Equations

In classical field theory, the dynamics of the physical system is revealed by analyzing the Euler-Lagrange equations corresponding to the classical action principle. These Euler-Lagrange equations are the physical equations (like the Maxwell or Einstein equations). They have the mathematical structure of partial differential equations. Likewise, for the causal action principle and causal variational principles, the Euler-Lagrange equations are the main analytic structure describing the dynamics. However, they are no longer differential equations but have a quite different form. In this chapter we shall derive the Euler-Lagrange equations and discuss their general structure.

7.1. The Euler-Lagrange Equations

Let \( \rho \) be a minimizer of the causal variational principle in the non-compact setting (see Section 6.3). We now derive the Euler-Lagrange (EL) equations, following the method in the compact setting [79, Lemma 3.4]. We again define spacetime as the support of \( \rho \),

\[ M := \text{supp} \rho \subset \mathcal{F}. \]

In order to make mathematical sense of the variations, we make the following assumptions:

(i) The measure \( \rho \) is locally finite (meaning that any \( x \in \mathcal{F} \) has an open neighborhood \( U \) with \( \rho(U) < \infty \)).

(ii) The function \( L(x,.) \) is \( \rho \)-integrable for all \( x \in \mathcal{F} \), giving a bounded continuous function on \( \mathcal{F} \).

We introduce the function

\[ \ell(x) = \int_{\mathcal{F}} L(x,y) \, d\rho(y) - s : \mathcal{F} \to \mathbb{R} \quad \text{bounded and lower semi-continuous}, \]

where the parameter \( s \in \mathbb{R} \) will be specified below.

**Theorem 7.1.1.** *(The Euler-Lagrange equations)* Let \( \rho \) be a minimizer of the causal action. Then

\[ \ell|_{\text{supp} \rho} \equiv \inf_{\mathcal{F}} \ell. \]

**Proof.** Given \( x_0 \in \text{supp} \rho \), we choose an open neighborhood \( U \) with \( 0 < \rho(U) < \infty \).

For any \( y \in \mathcal{F} \) we consider the family of measures \( (\hat{\rho}_\tau)_{\tau \in [0,1]} \) given by

\[ \hat{\rho}_\tau = \chi_{M \setminus U} \rho + (1 - \tau) \chi_U \rho + \tau \rho(U) \delta_y \]  

(7.1.3)

where \( \delta_y \) is the Dirac measure supported at \( y \). Then

\[ \hat{\rho}_\tau - \rho = -\tau \chi_U \rho + \tau \rho(U) \delta_y = \tau (\rho(U) \delta_y - \chi_U \rho), \]

implying that \( \hat{\rho}_\tau \) satisfies the volume constraint. Hence

\[ 0 \leq (\mathcal{S}(\hat{\rho}) - \mathcal{S}(\rho)) = 2\tau \left( \rho(U) \left( \ell(y) + s \right) - \int_U \left( \ell(x) + s \right) d\rho(x) \right) + \mathcal{O}(\tau^2). \]

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As a consequence, the linear term must be non-negative,
\[
\ell(y) \geq \frac{1}{\rho(U)} \int_U \ell(x) \, d\rho(x).
\] (7.1.5)

Assume that (7.1.2) is false. Then there is \(x_0 \in \text{supp} \rho\) and \(y \in \mathcal{F}\) such that \(\ell(x_0) > \ell(y)\). Lower semi-continuity of \(\ell\) implies that there is an open neighborhood \(U\) of \(x_0\) such that \(\ell(x) > \ell(y)\) for all \(x \in U\), in contradiction to (7.1.5). This gives the result. □

We always choose the parameter \(s\) such that the infimum of \(\ell\) in (7.1.2) is zero. Then the EL equations read
\[
\ell|_{\text{supp} \rho} \equiv \inf_{\mathcal{F}} \ell = 0.
\] (7.1.6)

### 7.2. The Restricted Euler-Lagrange Equations in the Smooth Setting

The EL equations (7.1.6) make a statement on the function \(\ell\) even at points \(x \in \mathcal{F}\) which are far away from spacetime \(M\) (see the left of Figure 7.1). In this way, the EL equations contain much more information than conventional physical equations formulated in spacetime. At present, it is unclear how this additional information can be used or interpreted. One way of understanding this situation is to take the point of view that all information on the physical system must be obtained by performing observations or measurements in spacetime, which means that the information contained in \(\ell\) away from \(M\) is inaccessible for principal reasons. Here we shall not take sides or discuss whether or to which extent this point of view is correct. Instead, we simply note that it seems preferable and physically sensible to restrict attention to the function \(\ell\) in an arbitrarily small neighborhood \(U\) of \(M\) in \(\mathcal{F}\) (see the right of Figure 7.1). In practice, this means that we shall evaluate \(\ell\) as well as its derivatives only on \(M\). In this way, the causal action principle gives rise to an interaction described by equations in spacetime.

This concept leads us to the so-called restricted EL equations which we now introduce. For technical simplicity, we assume that the Lagrangian is smooth (see (6.2.4) and the discussion thereafter); for a more general derivation see [68, Section 4]. In this case, the minimality of \(\ell\) implies that the derivative of \(\ell\) vanishes on \(M\), i.e.
\[
\ell|_M \equiv 0 \quad \text{and} \quad D\ell|_M \equiv 0 \quad (7.2.1)
\]
(\(D\ell(p) : T_p \mathcal{F} \to \mathbb{R}\) is the derivative). In order to combine these two equations in a compact form, it is convenient to consider a pair \(u := (a, u)\) consisting of a real-valued function \(a\) on \(M\) and a vector field \(u\) on \(T \mathcal{F}\) along \(M\), and to denote the combination of multiplication and directional derivative by
\[
\nabla_u \ell(x) := a(x) \ell(x) + (D_u \ell)(x).
\] (7.2.2)

Then the equations (7.2.1) imply that \(\nabla_u \ell(x)\) vanishes for all \(x \in M\),
\[
\nabla_u \ell|_M = 0 \quad \text{for all} \ u \in \mathfrak{g}_0^\infty.
\] (7.2.3)
These are the so-called restricted EL equations. For brevity, a solution of the restricted EL equations is also referred to as a critical measure. The pair \( u = (a, u) \) is referred to as a jet. This jet is a vector in a corresponding jet space \( J^\infty \) defined by

\[
u = (a, u) \in J^\infty := C^\infty(M, \mathbb{R}) \oplus C^\infty(M, T^* F).
\]

We remark that, in the literature, the restricted EL equations are sometimes also referred to as the weak EL equations. Here we prefer the notion “restricted” in order to avoid potential confusion with weak solutions of these equations (as constructed in [22]; see also Chapter 14 below).

### 7.3. Symmetries and Symmetric Criticality

In many applications, variational principles have an underlying symmetry (for example spherical symmetry or time independence). Typically, it simplifies the variational problem to vary within the class of functions which respect this symmetry. Having found a minimizer within this restricted class, the question arises whether it is also a minimizer of the full variational problem. The general answer to this question is no, simply because the absolute minimizer does not necessarily respect the symmetry of the variational principle. For causal variational principles, the situation is similar, if we only replace “function” by “measure.” As a simple example, we saw in Section 6.1 for the causal variational principle on the sphere that, although the variational principle is spherically symmetric, minimizing measures are typically weighted counting measures, thus breaking spherical symmetry.

Nevertheless, one can hope that minimizers within the class of symmetric functions are critical points of the full variational problem. This statement, referred to as the principle of symmetric criticality, has been formulated and proven under general assumptions in [116]. In this section we explain how the principle of symmetric criticality can be stated and proved in the setting of causal variational principles. As we shall see, the proof is quite simple and rather different from that in the classical calculus of variations. In order to explain the basic idea in the simplest possible setting, we only consider the compact setting and assume that also the symmetry group is compact. Dropping these compactness assumptions make the situation more involved, as is worked out in detail in [82].

As in Section 6.2 we let \( \mathcal{F} \) be a compact manifold. Moreover, we again assume that the Lagrangian \( L \) is continuous (6.2.5), symmetric and strictly positive on the diagonal (see the assumptions (i) and (ii) on page 87). In order to describe the symmetry, we let \( \mathcal{G} \) be a compact Lie group, which should act as a group of diffeomorphisms on \( \mathcal{F} \). To this end, we let \( \Phi : \mathcal{G} \times \mathcal{F} \to \mathcal{F} \) be a continuous mapping with the properties that \( \Phi_g := \Phi(g, \cdot) \) is a diffeomorphism of \( \mathcal{F} \) and that

\[
\Phi_g \circ \Phi_h = \Phi_{gh} \quad \text{for all } g, h \in \mathcal{G}.
\]

Moreover, the Lagrangian is assumed to be invariant under the group action, i.e.

\[
L(\Phi_g x, \Phi_g y) = L(x, y) \quad \text{for all } x, y \in \mathcal{F} \text{ and } g \in \mathcal{G}.
\]

We denote the set of normalized regular Borel measures on \( \mathcal{F} \) by \( \mathcal{M} \). Taking the push-forward of \( \Phi \), we obtain a group action on \( \mathcal{M} \) (for the definition of the push-forward measure see Section 2.3). We denote the measures which are invariant under this group action by \( \mathcal{M}^\mathcal{G} \), i.e.

\[
\mathcal{M}^\mathcal{G} := \{ \rho \in \mathcal{M} \mid (\Phi_g)_* \rho = \rho \quad \text{for all } g \in \mathcal{G} \}.
\]
We also refer to the measures in $\mathcal{M}$ as being equivariant (for more details on equivariant causal variational principles see [13] Section 4).

**Theorem 7.3.1.** Let $\rho$ be a minimizer of the causal action under variations within the class $\mathcal{M}^G$ of equivariant measures. Then $\rho$ is a critical point of the full variational principle in the sense that the EL equations (7.1.6) hold.

**Proof.** We denote the orbits of the group action by $\langle x \rangle := \Phi_G x$ with $x \in \mathcal{F}$. Since $\mathcal{G}$ is compact, so are the orbits. Let $\mu$ be the normalized Haar measure on $\mathcal{G}$. A particular class of equivariant measures are obtained by taking the push-forward of $\mu$ by the mapping $\Phi(\cdot, x)$. More precisely, for given $x \in \mathcal{F}$ we define the measure $\delta_{\langle x \rangle}$ by

$$\delta_{\langle x \rangle}(\Omega) := \mu(\{g \in \mathcal{G} \mid \Phi(g, x) \in \Omega\}).$$

The subscript $\langle x \rangle$ clarifies that, being equivariant, this measure depends only on the orbit.

Given $y \in \mathcal{F}$, we now consider the variation $(\tilde{\rho}_\tau)_{\tau \in [0, 1)}$ within the class of equivariant measures

$$\tilde{\rho}_\tau = (1 - \tau) \rho + \tau \delta_{\langle y \rangle}.$$

Using that $\rho$ is a minimizer within this class, we can follow the procedure in the proof of Theorem 7.1.1 to obtain

$$\int_{\mathcal{F}} \ell(x) \, d\delta_{\langle y \rangle} \geq \int_{\mathcal{F}} \ell(x) \, d\rho. \quad (7.3.2)$$

Moreover, it follows by symmetry that the function $\ell$ is constant on the orbits, because

$$\ell(\Phi_g y) = \int_{\mathcal{F}} \mathcal{L}(\Phi_g y, x) \, d\rho(x) - s = \int_{\mathcal{F}} \mathcal{L}(y, \Phi_g^{-1} x) \, d\rho(x) - s$$

$$= \int_{\mathcal{F}} \mathcal{L}(y, x) \, d\rho(x) - s = \ell(y),$$

where in the last line we used that $\rho$ is equivariant. Hence, integrating over the orbit, we obtain

$$\ell(y) = \int_{\mathcal{F}} \ell(x) \, d\delta_{\langle y \rangle}.$$

Combining this identity with (7.3.2), we conclude that

$$\ell(y) \geq \int_{\mathcal{F}} \ell(x) \, d\rho \quad \text{for all } y \in \mathcal{F}.$$

Now we can argue exactly as in the proof of Theorem 7.1.1 to obtain the result. \qed

**Exercises**
CHAPTER 8

The Linearized Field Equations

The EL equations as derived in the previous chapter (see Theorem 7.1.1 of the restricted EL equations in (7.2.3)) are nonlinear equations. This can be seen immediately from the fact that the measure \( \rho \) enters in a twofold way: It determines the function \( \ell \) via the integration (7.1.1), and it also determines via its support \( M \) where the function \( \ell \) is to be evaluated. As usual, such nonlinear equations are difficult to analyze. Therefore, it is a good idea to simplify these equations by linearization. This leads us to the so-called linearized field equations which describe linear perturbations of the system which respect the EL equations. This chapter is devoted to the derivation of the linearized field equations and to the construction of explicit examples. We remark that the linearized field equations are also a suitable starting point for the analysis of the full EL equations, for example by using perturbative methods (see [53] or Chapter [13]).

8.1. Derivation of the Linearized Field Equations in the Smooth Setting

Linearized fields often appear in physics, for example when describing linearized oscillations or linear waves. In order to derive the linearized equations, one typically considers a family \( G_\tau \) of solutions. The parameter \( \tau \) can be thought of as the “amplitude” of the perturbation, and \( G_\tau |_{\tau = 0} \) describes the unperturbed system (for example an anharmonic oscillator at rest or the vacuum). Then the derivative

\[
\frac{d}{d\tau}G_\tau \bigg|_{\tau = 0}
\]

is the linearized field; it satisfies a linear equation obtained by differentiating the nonlinear equation (like Hamilton’s equation for the anharmonic oscillator or the nonlinear wave equation) with respect to \( \tau \).

The concept of linearization is also fruitful in the context of causal variational principles. Since the system is described by the measure \( \rho \), the above family of solutions now corresponds to a family of measures \( (\tilde{\rho}_\tau)_{\tau \in [0,\delta)} \) which are all critical points of the causal action principle. The basic question is how to vary the measure. Indeed, there are many ways of varying. We begin with a simple method, which we will generalize and discuss afterward. In order to keep the presentation as simple as possible, we again restrict attention to the smooth setting. We choose a family \( f_\tau \) of positive weight functions and a family \( F_\tau \) of mapping from \( M \) to \( \mathcal{F} \). These functions should all be smooth, also in the parameter \( \tau \), i.e.

\[ f \in C^\infty([0,\delta) \times M, \mathbb{R}^+) \quad \text{and} \quad F \in C^\infty([0,\delta) \times M, \mathcal{F}) \, . \]

We multiply \( \rho \) by \( f_\tau \) and then take the push-forward under \( F_\tau \),

\[
\tilde{\rho}_\tau := (F_\tau)_* (f_\tau \rho) \, .
\]  

We assume that for \( \tau = 0 \) the variation is trivial,

\[
f_0 \equiv 1 \quad \text{and} \quad F_0 \equiv 1 \, .
\]
Since multiplying by a positive function leaves the support unchanged, the support of the measure is transformed only by $F_\tau$; more precisely,

$$\text{supp } \tilde{\rho}_\tau = F_\tau(\text{supp } \rho) \quad (8.1.3)$$

(for details see Exercise 8.1).

The assumption that all the measures $\tilde{\rho}_\tau$ are critical means that they all satisfy the restricted EL equations (7.2.3). Taking into account that the support of the measures changes according to (8.1.3), we know that for all $u \in J^0_\infty$ and all $x \in M$,

$$0 = \nabla_u \left( \int_{\mathcal{J}} L(F_\tau(x), y) \, d\tilde{\rho}_\tau(y) - s \right)$$

$$= \nabla_u \left( \int_{\mathcal{J}} L(F_\tau(x), F_\tau(y)) \, f_\tau(y) \, d\rho(y) - s \right),$$

where in the last line we used the definition of the push-forward measure. It is convenient to multiply this equation by $f_\tau(x)$. We can write this factor inside the brackets,

$$0 = \nabla_u \left( \int_{\mathcal{J}} f_\tau(x) \, L(F_\tau(x), F_\tau(y)) \, f_\tau(y) \, d\rho(y) - f_\tau(x) \, s \right),$$

because the terms obtained when the derivative $\nabla_u$ acts on $f_\tau(x)$ vanish in view of the restricted EL equations (7.2.3). Since this equation holds for all $\tau \in [0, \delta)$, we can differentiate at $\tau = 0$ to obtain the equation

$$\nabla_u \left( \int_{\mathcal{J}} (\nabla_1, v + \nabla_2, v) L(x, y) \, d\rho(y) - \nabla_v \, s \right) = 0 \quad (8.1.4)$$

where $v$ is the jet generated by the functions $f_\tau$ and $F_\tau$,

$$v := \left. \frac{d}{d\tau} (f_\tau, F_\tau) \right|_{\tau = 0} \in J^\infty.$$

Here $\nabla_1, v$ and $\nabla_2, v$ denote partial derivatives acting on the first and second argument of the Lagrangian, respectively. Here and throughout this book, we use the following conventions for partial derivatives and jet derivatives:

(i) Partial and jet derivatives with an index $i \in \{1, 2\}$ only act on the respective variable of the function $L$. This implies, for example, that the derivatives commute,

$$\nabla_1, v \nabla_1, u L(x, y) = \nabla_1, u \nabla_1, v L(x, y).$$

(ii) The partial or jet derivatives which do not carry an index act as partial derivatives on the corresponding argument of the Lagrangian. This implies, for example, that

$$\nabla_u \int_{\mathcal{J}} \nabla_1, v L(x, y) \, d\rho(y) = \int_{\mathcal{J}} \nabla_1, u \nabla_1, v L(x, y) \, d\rho(y).$$

We point out that jets are never differentiated. This is a very convenient convention. In (8.1.4) this convention can be used because the terms obtained if the derivative $\nabla_u$ acted on the jet $v$ again vanish as a consequence of the restricted EL equations. We remark that, from a differential geometric perspective, defining higher derivatives on $\mathcal{F}$ would make it necessary to introduce a connection on $\mathcal{F}$. While this could be done, we here use the simpler method that higher derivatives on $\mathcal{F}$ are defined as partial derivatives carried out in distinguished charts. More precisely, around each point $x \in \mathcal{F}$ we distinguish a chart and carry out derivatives as partial derivatives acting on each tensor component in this chart. We remark that, in the setting of causal fermion systems, an atlas of distinguished
8.1. DERIVATION OF THE LINEARIZED FIELD EQUATIONS IN THE SMOOTH SETTING

**Definition 8.1.1.** A jet \( \mathfrak{v} \in \mathfrak{J}_\infty \) is referred to as a solution of the linearized field equations if

\[
(u, \Delta \mathfrak{v})(x) := \nabla u \left( \int_\mathcal{F} (\nabla_{1,\mathfrak{v}} + \nabla_{2,\mathfrak{v}}) \mathcal{L}(x, y) \, d\rho(y) - \nabla_x \mathfrak{v} \right) = 0 \tag{8.1.6}
\]

for all \( u \in \mathfrak{J}_0^\infty \) and all \( x \in M \). The vector space of all linearized solutions is denoted by \( \mathfrak{J}_\text{lin} \subset \mathfrak{J}_\infty \).

We conclude this section with a brief discussion of our ansatz (8.1.1). Intuitively speaking, this ansatz means that the support of the measure is changed smoothly as a whole. In particular, if \( M \) is a smooth four-dimensional submanifold of \( \mathcal{F} \), then the varied measure \( M_\tau \) will again have this property. In physical terms, measures where \( M \) has such a manifold structure describe classical spacetimes. Consequently, the ansatz (8.1.1) and the corresponding linearization (8.1.5) correspond to classical fields in a classical spacetime. In contrast, if the support \( M := \text{supp} \rho \) of the measure does not have the structure of a four-dimensional manifold, then we refer to \( M \) as a quantum spacetime.

The notion "quantum spacetime" appears in the literature in different contexts with rather different meanings. Here we take the above notion as the mathematical definition of what we mean by a quantum spacetime.

In order to give an idea for how such a quantum spacetime may look like, let us consider the example where the unperturbed measure \( \rho \) describes a classical spacetime \( M \) (for example Minkowski space \( M \cong \mathbb{R}^4 \)). As just explained, the ansatz (8.1.1) changes the support of the measure smoothly as a whole (see Figure 8.1 (a)). More generally, one can consider the situation where the measure \( \rho \) "disintegrates" into several "components" which are perturbed differently (see Figure 8.1 (b)). To this end, we choose a parameter \( L \in \mathbb{N} \) (the “number of subsystems”) and consider mappings

\[
f_a \in C^\infty((0, \delta) \times M, \mathbb{R}^+) \ , \quad F_a \in C^\infty((0, \delta) \times M, \mathcal{F}) \quad \text{with} \quad a = 1, \ldots, L .
\]

For the so-called measure with fragmentation, in generalization of (8.1.1) we make the ansatz

\[
\tilde{\rho}_\tau = \frac{1}{L} \sum_{a=1}^L (F_a)_*(f_a \rho) . \tag{8.1.7}
\]

The larger \( L \) is chosen, the more freedom we have in perturbing the measure. We point out that we may choose \( L \) arbitrarily large. In the limit \( L \to \infty \), one can even describe situations where the support of the measure \( \rho \) is “enlarged” by the perturbation, as is shown in Figure 8.1 (c).
Assuming that the family \((\tilde{\rho}_\tau)_{\tau \in [0,\delta)}\) satisfies the restricted EL equations for all \(\tau\), we can again linearize in \(\tau\) to obtain the corresponding linearized field equations. They again have the form (8.1.6), but now with \(v\) being the “averaged jet”

\[
v = \frac{1}{L} \sum_{a=1}^{L} v_a \quad \text{with} \quad v_a = \frac{d}{d\tau}(f_{a,\tau}, F_{a,\tau}) \big|_{\tau=0}.
\]

Therefore, for linearized fields the fragmentation does not give anything essentially new. But on the nonlinear level, fragmentation yields additional effects. We refer the interested reader for more details to [51, Section 5] and [53, Section 5] as well as to the applications to quantum field theory in [60] (see also Chapter 22).

In view of this consideration, the only restriction in describing linear perturbations of a measure \(\rho\) by a jet \(v\) of the form (8.1.5) is that the support of the measure \(\rho\) is changed continuously in \(\tau\), in the sense the support \(\text{supp} \tilde{\rho}_\tau\) lies in a small neighborhood of \(M\) (for details see Exercise 8.2). In particular, we do not cover variations of the form (7.1.3) where part of the measure is “transported” to a point \(y \in F\) which may be far away from \(M\). The reason for disregarding such variations is that, similar as explained before introducing the restricted EL equations in Section 7.2 (see Figure 7.1), analyzing the EL equations outside a small neighborhood of \(M\) does not seem to be of physical relevance.

### 8.2. Commutator Jets in Causal Fermion Systems

In order to illustrate the linearized field equations, we conclude this chapter by deriving explicit classes of solutions. These solutions correspond to inherent symmetries of the system. In this section, we consider commutator jets, which describe the unitary invariance of a causal action principle. In the next section (Section 8.3), we shall derive inner solutions by using the invariance of the measure under diffeomorphisms of \(M\) combined with a suitable multiplication of \(\rho\) by a smooth weight function.

Let \((\mathcal{H}, F, \rho)\) be a causal fermion system. The causal action principle is unitarily invariant in the following sense. Let \(U \in \mathcal{U}(\mathcal{H})\) be a unitary transformation. Given a measure \(\rho\) on \(F\), we can unitarily transform the measure by setting

\[
(U\rho)(\Omega) = \rho(U^{-1}\Omega U) \quad \text{for} \quad \Omega \subset F.
\]

Since the eigenvalues of an operator are invariant under unitary transformations, a measure \(\rho\) is a minimizer or critical point of the causal action principle if and only if \(U\rho\) is.

Infinitesimally, this unitary invariance gives rise to a special class of solutions of the linearized field equations: Let \(\rho\) be a critical measure. We let \(A\) be a symmetric operator on \(\mathcal{H}\), for technical simplicity of finite rank. By exponentiating, we obtain a family of unitary operators \((U_\tau)_{\tau \in \mathbb{R}}\) with

\[
U_\tau := \exp(i\tau A).
\]

According to (8.2.1), the support of the measures \(\tilde{\rho}_\tau := U_\tau \rho\) is given by

\[
\tilde{M}_\tau := \text{supp} \tilde{\rho}_\tau = U_\tau M U_\tau^{-1}.
\]

Due to the unitary invariance of the Lagrangian, the measures \(\tilde{\rho}_\tau\) all satisfy the EL equations. Infinitesimally, the unitary transformations are described by the jet \(v\) given by

\[
v := (0, v) \in \mathfrak{J}^{\text{lin}} \quad \text{with} \quad v(x) = \frac{d}{d\tau}(U_\tau x U_\tau^{-1}) \big|_{\tau=0} = i[A, x].
\]
Due to the commutator in the last equation, we refer to jets of this form as **commutator jets**. The fact that commutator jets generate families of critical measures implies that they satisfy the linearized field equations:

**Lemma 8.2.1.** The commutator jet \( v \) in \((8.2.4)\) satisfies the linearized field equations \((8.1.6)\).

**Proof.** Due to the unitary invariance of the Lagrangian,
\[
L(U_\tau x U_\tau^{-1}, U_\tau y U_\tau^{-1}) = L(x, y).
\]
Differentiating with respect to \( \tau \) and applying the chain rule gives
\[
(D_{1,v} + D_{2,v}) L(x, y) \, d\rho(y) = 0.
\]
Hence the integrand in \((8.1.6)\) vanishes for all \( x, y \in \mathcal{F} \). As a consequence, the integral in \((8.1.6)\) vanishes for all \( x \in \mathcal{F} \). Consequently, also its derivative in the direction of \( u \) vanishes. Using our convention that the jet derivatives act only on the Lagrangian (see the end of Section 8.1), this directional derivative differs from the jet derivative in \((8.1.6)\) by the term \( D_{D,v} \ell(x) \). This term vanishes in view of the restricted EL equations \((7.2.3)\).

### 8.3. Inner Solutions in Smooth Spacetimes

We now return to the setting of causal variational principles. We introduce an additional smoothness assumption and explain why it is useful in some applications.

**Definition 8.3.1.** Spacetime \( M := \text{supp} \rho \) has a **smooth manifold structure** if the following conditions hold:

(i) \( M \) is diffeomorphic to a smooth oriented manifold \( M^k \) of dimension \( k \).

(ii) In a chart \((x, U)\) of \( M^k \), the measure \( \rho \) is absolutely continuous with respect to the Lebesgue measure with a smooth, strictly positive weight function, i.e.
\[
d\rho = h(x) \, d^k x \quad \text{with} \quad h \in C^\infty(M^k, \mathbb{R}^+) \, .
\]

Even though there is no reason why physical spacetime should have a smooth manifold structure on the Planck scale, this assumption is clearly justified on the macroscopic scale of atomic and gravitational physics. With this in mind, the assumption of a smooth manifold structure seems admissible in all applications in which the microscopic structure of spacetime should be irrelevant. Before going on, we point out that one should carefully distinguish the assumption of a smooth manifold structure from the smooth setting introduced in Section 6.2. One should keep in mind that the smoothness of \( L \) does not imply that \( M \) has a smooth manifold structure, nor vice versa.

The fact that \( \rho \) is defined independent of charts implies that the function \( h \) in \((8.3.1)\) transform like a tensor density. Thus on the overlap of two charts \((x, U)\) and \((\tilde{x}, \tilde{U})\), we know that
\[
h(x) \, d^k x = \tilde{h}(\tilde{x}) \, d^k \tilde{x}
\]
and thus
\[
h(x) = \det \left( \frac{\partial x^i}{\partial \tilde{x}^j} \right) \tilde{h}(\tilde{x}) \, .
\]
This transformation law makes it possible to define the covariant divergence of a vector field \( v \) on \( M \cong M^k \) in a local chart by
\[
\text{div} \, v = \frac{1}{h} \partial_j (h v^j)
\]
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(where following the Einstein summation convention we sum over $j = 0, \ldots, k$). Alternatively, the divergence of a vector field $v \in \Gamma(M, TM)$ can be defined independent of charts by the relation

$$\int_M \text{div } v \eta(x) \, d\rho = - \int_M D_v \eta(x) \, d\rho(x),$$

to be satisfied for all test functions $\eta \in C^\infty_0(M, \mathbb{R})$.

**Definition 8.3.2.** An inner solution is a jet $v \in \mathcal{J}^1$ of the form

$$v = (\text{div } v, v) \quad \text{with} \quad v \in \Gamma(M, TM).$$

The vector space of all inner solutions is denoted by $\mathcal{J}^{in} \subset \mathcal{J}^1$.

The name “inner solution” is justified by the following lemma:

**Lemma 8.3.3.** Every inner solution $v \in \mathcal{J}^{in}$ of compact support is a solution of the linearized field equations, i.e.

$$\langle u, \Delta v \rangle_M = 0 \quad \text{for all } u \in \mathcal{J}^\infty_0.$$

**Proof.** Applying the Gauss divergence theorem, one finds that for every $f \in C^1_0(M, \mathbb{R})$,

$$\int_M \nabla_v f \, d\rho = \int_M (\text{div } v f + D_v f) \, d\rho = \int_M \text{div } (fv) \, d\rho = 0.$$

Likewise, in the linearized field equations we may integrate by parts in $y$,

$$\langle u, v \rangle_M = \nabla_u \left( \int_M (\nabla_1 v + \nabla_2 v) \mathcal{L}(x, y) - \nabla_v \mathcal{S} \right)$$

$$= \nabla_u \left( \int_M \nabla_1 v \mathcal{L}(x, y) - \nabla_v \mathcal{S} \right)$$

$$= \nabla_u \nabla_v \ell(x) = \nabla_v (\nabla_u \ell(x)) - \nabla D_v \ell(x) = 0,$$

where the last term vanishes in view of the EL equations. Moreover, we used that the function $\nabla_u \ell$ vanishes identically on $M$ in view of the restricted EL equations. Therefore, this function is differentiable in the direction of every vector field on $M$, and this directional derivative is zero. □

This result also holds for inner solutions $v$ of non-compact support, provided that the vector field $v$ has suitable decay properties at infinity. For details we refer to [59, Section 3].

We now turn to the question which scalar components can be realized by inner solutions. This question can be answered in a more generality by applying Moser’s theorem (see for example [107, Section XVIII, §2] or straightforward generalizations to non-compact manifolds). For simplicity, we here make additional assumptions which make it possible to use hyperbolic methods.

**Theorem 8.3.4.** Assume that $M \simeq \mathcal{M}^k$ has a smooth manifold structure. Moreover, assume that $\mathcal{M}$ is topologically of the form $\mathcal{M} = \mathbb{R} \times \mathcal{N}$ with a manifold $\mathcal{N}$ which admits a complete Riemannian metric $g_N$. Let $a \in C^\infty_sc(\mathcal{M}, \mathbb{R})$ be a smooth function with spatially compact support (meaning that, for all $t \in \mathbb{R}$, the function $a(t, \cdot)$ is compactly supported in $\mathcal{N}$). Then there is a vector field $v \in C^\infty_sc(\mathcal{M}, T\mathcal{M})$, again with spatially compact support, such that the jet $v := (a, v)$ is an inner solution.
Proof. Our task is to solve the equation \( \text{div } v = a \), which can be written equivalently as

\[
\frac{\partial}{\partial j} (hv^j) = ha.
\]

(8.3.3)

We first consider the case that \( a \) has compact support. In order to solve the partial differential equation (8.3.3), it is useful to choose a Lorentzian metric with

\[
ds^2 = dt^2 - g_N.
\]

(8.3.4)

Here the choice of the Riemannian metric \( g_N \) is irrelevant, and the arbitrariness in choosing this metric corresponds to the fact that (8.3.3) is an under-determined equation which admits many different solutions. Let \( \Box \) be the corresponding wave operator. Using for example retarded Green’s operators, there is a solution \( \phi \in C^\infty_c(M, \mathbb{R}) \) with \( \Box \phi = ha \) (for details see Section 13 below). Then the vector field

\[
v^j := \frac{1}{h} g^{jk} \partial_k \phi
\]

(8.3.5)

satisfies (8.3.3) (note that, in view of (8.3.1), we may divide by \( h \) to again obtain a smooth vector field with spatially compact support).

In the case that \( a \) merely has spatially compact support, we decompose \( a \) as

\[
a = a_+ + a_-,
\]

where \( a_+ \) is supported in the set \( \{ t > 0 \} \) and \( a_- \) is supported in the set \( \{ t < 1 \} \). Denoting the advanced and retarded Green’s operators of the scalar wave equation corresponding to the Lorentzian metric (8.3.4) by \( S^\lor \) and \( S^\land \), respectively, the function

\[
\phi := S^\land (ha_+) + S^\lor (ha_-)
\]

is a well-defined solution of the equation \( \Box \phi = ha \) which is smooth and has spatially compact support. Therefore, we can again define the vector field \( v \) by (8.3.5). This gives the result. \( \square \)

Inner solutions can be regarded as infinitesimal generators of transformations of \( M \) which leave the measure \( \rho \) unchanged. Therefore, inner solutions do not change the causal fermion system, but merely describe symmetry transformations of the measure. The result of the previous proposition can be used to change the scalar component of a linearized solution arbitrarily. More specifically, considering a system of finite spatial volume for a scattering process taking place in finite time, all the jets describing the interaction will have spatially compact support. Therefore, we can compensate the scalar components by corresponding inner solutions. With this in mind, if spacetime has a smooth manifold structure (see Definition 8.3.1), then one may always restrict attention to jets with vanishing scalar components. This simplifies many constructions considerably.

Exercises

Exercise 8.1. Let \( F : \mathcal{F} \to \mathcal{F} \) be continuous and \( \rho \) a measure on \( \mathcal{F} \). Show that

\[
supp F_* \rho = \overline{F(supp \rho)}.
\]

Hint: Recall the definition of the support of a measure (2.3.3) and use that the preimage of an open set under a continuous mapping is open.
Exercise 8.2. (a) Assume that $\mathcal{F}$ is locally compact. Moreover, assume that $F \in C^0([0, \delta) \times M, \mathcal{F})$ is continuous and that its preimage of any compact set is compact. Then for any $y \notin M$ there is $\tau_0 \in (0, \delta)$ such that

$$y \notin \text{supp} \tilde{\rho}_\tau \quad \text{for all } \tau \in [0, \tau_0]$$

(where $\tilde{\rho}_\tau$ are again the measures (8.1.1)) Hint: Use the result of Exercise 8.1.

(b) Show that this result remains valid for the variation (8.1.7) with a finite number of subsystems.

(c) What happens for an infinite number of subsystems? Also, is the assumption necessary that the preimage of a compact set under $F$ is compact?

Exercise 8.3. (Linearization of nonlinear partial differential equations) In this exercise you are given two non-linear partial differential equations with corresponding (soliton) solutions. Check that the functions $\phi$ do indeed solve the equations. Then try to figure out what it means to linearize the equations around the given solutions and do it.

(a) The sine-Gordon equation of velocity $v \in (-1, 1)$:

$$\partial_t \phi - \partial_{xx} \phi + \sin \phi = 0, \quad \phi(t, x) = 4 \arctan \left( \frac{x - vt}{\sqrt{1 - v^2}} \right).$$

(b) The Korteweg-de-Vries equation of unit speed:

$$\partial_t \phi + 6 \phi \partial_x \phi + \partial_{xxx} \phi = 0, \quad \phi(t, x) = \frac{1}{2} \sech^2 \left( \frac{x - vt}{2} \right).$$

Hint: You may use the following identities,

$$\sin(4 \arctan(x)) = -4 \frac{x^3 - x}{(1 + x^2)^2}, \quad \tanh(x) - \tanh^3(x) = \sech^2(x) \tanh(x).$$
CHAPTER 9

Surface Layer Integrals and Conservation Laws

In this chapter we introduce surface layer integrals as an adaptation of surface integrals to causal fermion systems and causal variational principles. The mathematical structure of a surface layer integral fits nicely to the analytic structures (namely, the EL equations and the linearized field equations as introduced in Chapter 8). This will become apparent in conservation laws which generalize Noether’s theorem and the symplectic form to the setting of causal variational principles. Next, we introduce a so-called nonlinear surface layer integral which makes it possible to compare two measures $\rho$ and $\tilde{\rho}$ at a given time. Finally, we shall explain how two-dimensional surface integrals can be described by surface layer integrals.

9.1. The Concept of a Surface Layer Integral

In daily life we experience space and objects therein. These objects are usually described by densities, and integrating these densities over space gives particle numbers, charges, the total energy, etc. In mathematical terms, the densities are typically described as the normal components of vector fields on a Cauchy surface, and conservation laws express that the values of these integrals do not depend on the choice of the Cauchy surface, i.e.

$$\int_{\mathcal{N}_1} J^k \nu_k \, d\mu_{\mathcal{N}_1}(x) = \int_{\mathcal{N}_2} J^k \nu_k \, d\mu_{\mathcal{N}_2}(x), \quad (9.1.1)$$

where $\mathcal{N}_1$ and $\mathcal{N}_2$ are two Cauchy surfaces, $\nu$ is the future-directed normal, and $d\mu_{\mathcal{N}_1/2}$ is the induced volume measure.

In the setting of causal variational principles, surface integrals like $(9.1.1)$ are undefined. Instead, one considers so-called surface layer integrals, as we now explain. In general terms, a surface layer integral is a double integral of the form

$$\int_{\Omega} \left( \int_{M \setminus \Omega} \cdots L(x,y) \, d\rho(y) \right) \, d\rho(x), \quad (9.1.2)$$

where one variable is integrated over a subset $\Omega \subset M$, and the other variable is integrated over the complement of $\Omega$. Here $(\cdots)$ stands for a differential operator acting on the Lagrangian. In order to explain the basic idea, we make the assumption that the Lagrangian is of short range in the following sense. We let $d \in C^0(M \times M, \mathbb{R}_0^+)$ be a suitably chosen distance function on $M$. Then the assumption of short range can be quantified by demanding that $L$ should vanish on distances larger than $\delta$, i.e.

$$d(x,y) > \delta \implies L(x,y) = 0. \quad (9.1.3)$$

Under this assumption, the surface layer integral $(9.1.2)$ only involves pairs $(x,y)$ of distance at most $\delta$, where $x$ lies in $\Omega$, whereas $y$ lies in the complement $M \setminus \Omega$. As a consequence, the integral only involves points in a layer around the boundary of $\Omega$ of
width $\delta$, i.e.

$$x, y \in B_\delta(\partial \Omega).$$

Therefore, a double integral of the form \((9.1.2)\) can be regarded as an approximation of a surface integral on the length scale $\delta$, as shown in Figure 9.1. In the setting of causal variational principles, such surface layer integrals take the role of surface integrals. We remark that in applications in Minkowski space or on a Lorentzian manifold, the Lagrangian typically decays on the Compton scale $1/m$ (where $m$ denotes the rest mass of the Dirac particles).

### 9.2. Noether-Like Theorems

In modern physics, the connection between symmetries and conservation laws is of central importance. For continuous symmetries, this connection is made mathematically precise by Noether's theorem (see [113] or the textbooks [90, Section 13.7], [7, Chapter III]). As shown in [67], the connection between symmetries and conservation laws can be extended to the setting of causal variational principles. As we shall see, both the statement and the proof are quite different from the classical Noether theorem; this is why we refer to our result as a Noether-like theorem.

The first step is to introduce a symmetry condition for the Lagrangian $L(x,y)$ of a causal variational principle. Similar to the procedure in Section 7.3, one could describe the symmetry by a group of diffeomorphisms. For the purposes here, the correct setting would be to consider a one-parameter group of diffeomorphisms $\Phi_\tau$, i.e.

$$\Phi : \mathbb{R} \times \mathcal{F} \to \mathcal{F} \quad \text{with} \quad \Phi_\tau \Phi_{\tau'} = \Phi_{\tau + \tau'}, \quad (9.2.1)$$

(we usually write the first argument as a subscript, i.e. $\Phi_\tau(x) \equiv \Phi(\tau, x)$). The symmetry condition could be imposed by demanding that the Lagrangian be invariant under this one-parameter group in the sense that

$$L(x, y) = L(\Phi_\tau(x), \Phi_\tau(y)) \quad \text{for all } \tau \in \mathbb{R} \text{ and } x, y \in \mathcal{F}. \quad (9.2.2)$$

It turns out that this condition is unnecessarily strong for two reasons. First, it suffices to consider families which are defined locally for $\tau \in (-\tau_{\text{max}}, \tau_{\text{max}})$. Second, the mapping $\Phi$ does not need to be defined on all of $\mathcal{F}$. Instead, it is more appropriate to impose the symmetry condition only on spacetime $M \subset \mathcal{F}$. This leads us to consider instead of \((9.2.1)\) a mapping

$$\Phi : (-\tau_{\text{max}}, \tau_{\text{max}}) \times M \to \mathcal{F} \quad \text{with} \quad \Phi_0 = \text{id}_M. \quad (9.2.3)$$

We refer to $\Phi_\tau$ as a variation of $M$ in $\mathcal{F}$. Next, we need to specify what we mean by “smoothness” of this variation. This is a subtle point because, as explained in the example of the causal variational principle on the sphere in Section 6.1, the support of a minimizing measure will in general be singular. Moreover, the function $\ell$ defined...
symmetry of the Lagrangian. Then for any compact subset $\Omega$ if the composition $\ell \circ \Phi : (-\tau_{\text{max}}, \tau_{\text{max}}) \times M \to \mathbb{R}$ is continuous and if its partial derivative $\partial_\tau (\ell \circ \Phi)$ exists and is continuous.

The next question is how to adapt the symmetry condition (9.2.2) to the mapping $\Phi : \mathbb{R} \times \mathcal{F} \to M$. It turns out that the correct procedure is to work with the expression in (9.2.4).

**Definition 9.2.2.** A variation $\Phi_\tau$ of the form (9.2.3) is continuously differentiable if the composition

$$
\ell \circ \Phi : (-\tau_{\text{max}}, \tau_{\text{max}}) \times M \to \mathbb{R}
$$

is continuous and if its partial derivative $\partial_\tau (\ell \circ \Phi)$ exists and is continuous.

Theorem 9.2.3. Let $\rho$ be a critical measure and $\Phi_\tau$ a continuously differentiable symmetry of the Lagrangian. Then for any compact subset $\Omega \subset M$,

$$
\frac{d}{d\tau} \int_\Omega d\rho(x) \int_{M\setminus\Omega} d\rho(y) \left( \mathcal{L}(\Phi_\tau(x), y) - \mathcal{L}(\Phi_{-\tau}(x), y) \right) \bigg|_{\tau=0} = 0. \tag{9.2.6}
$$

**Proof.** Integrating (9.2.5) over $\Omega \times \Omega$ gives

$$
0 = \int_{\Omega \times \Omega} \left( \mathcal{L}(x, \Phi_\tau(y)) - \mathcal{L}(\Phi_{-\tau}(x), y) \right) d\rho(x) d\rho(y)
$$

$$
= \int_{\Omega} d\rho(x) \int_M d\rho(y) \chi(\rho(y)) \left( \mathcal{L}(\Phi_\tau(x), y) - \mathcal{L}(\Phi_{-\tau}(x), y) \right),
$$

where in the last step we used the symmetry of the Lagrangian (9.2.5) as well as the symmetry of the integrand in $x$ and $y$. We rewrite the characteristic function $\chi(y)$ as $1 - (1 - \chi(y))$, multiply out and use the definition of $\ell$, (7.1.1). We thus obtain

$$
0 = \int_{\Omega} \left( \ell(\Phi_\tau(x)) - \ell(\Phi_{-\tau}(x)) \right) d\rho(x)
$$

$$
- \int_{\Omega} d\rho(x) \int_M d\rho(y) \chi_M \chi(y) \left( \mathcal{L}(\Phi_\tau(x), y) - \mathcal{L}(\Phi_{-\tau}(x), y) \right).
$$

We thus obtain the identity

$$
\int_{\Omega} d\rho(x) \int_{M\setminus\Omega} d\rho(y) \left( \mathcal{L}(\Phi_\tau(x), y) - \mathcal{L}(\Phi_{-\tau}(x), y) \right)
$$

$$
= \int_{\Omega} \left( \ell(\Phi_\tau(x)) - \ell(\Phi_{-\tau}(x)) \right) d\rho(x). \tag{9.2.7}
$$

Using that $\ell(\Phi_\tau(x))$ is continuously differentiable (see Definition 9.2.2) and that $\Omega$ is compact, we conclude that the right side of this equation is differentiable at $\tau = 0.$
Moreover, we are allowed to interchange the \( \tau \)-differentiation with integration. The EL equations (7.1.6) imply that
\[
\frac{d}{d\tau} \ell'(\Phi_\tau(x)) \bigg|_{\tau=0} = 0 = \frac{d}{d\tau} \ell'(\Phi_{-\tau}(x)) \bigg|_{\tau=0}.
\]
Hence the right side of (9.2.7) is differentiable at \( \tau = 0 \), and the derivative vanishes. This gives the result.

This theorem requires a detailed explanation. We first clarify the connection to surface layer integrals. To this end, let us assume that \( \Phi_\tau \) and the Lagrangian are differentiable in the sense that the derivatives
\[
\frac{d}{d\tau} \Phi_\tau(x) \bigg|_{\tau=0} =: u(x) \quad \text{and} \quad \frac{d}{d\tau} \mathcal{L}(\Phi_\tau(x), y) \bigg|_{\tau=0}
\]
exist for all \( x, y \in M \) and are continuous on \( M \) respectively \( M \times M \). Then one may exchange differentiation and integration in (9.2.6) and apply the chain rule to obtain
\[
\int \Omega \, d\rho(x) \int_{M \setminus \Omega} d\rho(y) \, D_{1,u} \mathcal{L}(x, y) = 0,
\]
where \( D_{1,u} \) is the partial derivative at \( x \) in the direction of the vector field \( u(x) \). This expression is a surface layer integral as in (9.1.2). In general, the derivatives in (9.2.9) need not exist, because we merely imposed the weaker differentiability assumption of Definition 9.2.1. In this case, the statement of the theorem implies that the derivative of the integral in (9.2.6) exists and vanishes.

We next explain the connection to conservation laws. Let us assume that \( M \) admits a sensible notion of “spatial infinity” and that the vector field \( \partial_\tau \Phi \in \Gamma(M, TF) \) has suitable decay properties at spatial infinity. Then one can choose a sequence \( \Omega_n \subset M \) of compact sets which form an exhaustion of a set \( \Omega \) which extends up to spatial infinity (see Figure 9.2 (a) and (b)). Considering the surface layer integrals for \( \Omega_n \) and passing to limit, one concludes that also the surface layer integral corresponding to \( \Omega \) vanishes. Let us assume that the boundary \( \partial \Omega \) has two components \( N_1 \) and \( N_2 \) (as in Figure 9.2 (b)). Then the above theorem implies that the surface layer integrals over \( N_1 \) and \( N_2 \) coincide (where the surface layer integral over \( N \) is defined as the surface layer integral corresponding to a set \( \Omega_N \) with \( \partial \Omega_N = N \) as shown in Figure 9.2 (c)). In other words, the quantity
\[
\frac{d}{d\tau} \int_{\Omega_N} \, d\rho(x) \int_{M \setminus \Omega_N} d\rho(y) \left( \mathcal{L}(\Phi_\tau(x), y) - \mathcal{L}(\Phi_{-\tau}(x), y) \right) \bigg|_{\tau=0}
\]
is well-defined and independent of the choice of \( N \). In this setting, the surfaces \( N \) can be interpreted as Cauchy surfaces, and the conservation law of Theorem 9.2.3 means that the surface layer integral is preserved under the time evolution.
As a concrete example, the unitary invariance of the causal action principle gives rise to a conservation law, which corresponds to current conservation. This example will be considered in detail in Section 9.4. We finally remark that the conservation laws for energy-momentum can also be obtained from Theorem 9.2.3 assuming that the causal fermion system has symmetries as described by generalized Killing symmetries. We refer the interested reader to [67, Section 6].

9.3. A CLASS OF CONSERVATION LAWS IN THE SMOOTH SETTING

In the previous section we saw that surface layer integrals can be used to formulate a Noether-like theorem which relates symmetries to conservation laws. In this section we shall derive conservation laws even in the absence of symmetries. Instead, these conservation laws are closely tied to the structure of the linearized field equations as derived in Section 8.1. In order to focus on the essence of the construction, we again restrict attention to the smooth setting (6.2.4).

The basic idea of the construction is explained in the following proposition:

**Proposition 9.3.1.** Let \( \Omega \subset M \) be a compact. Then for any solution \( v \in J^{\text{lin}} \) of the linearized field equations (8.1.6),

\[
\gamma_\rho^\Omega(v) := \int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) \left( \nabla_{1,v} - \nabla_{2,v} \right) \mathcal{L}(x,y) = \int_\Omega \nabla_v s \, d\rho.
\]  

**Proof.** In view of the anti-symmetry of the integrand,

\[
\int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) \left( \nabla_{1,v} - \nabla_{2,v} \right) \mathcal{L}(x,y) = 0.
\]

Adding this equation to the left side of (9.3.1), we obtain

\[
\gamma_\rho^\Omega = \int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) \left( \nabla_{1,v} - \nabla_{2,v} \right) \mathcal{L}(x,y) = \int_\Omega \left( 2 \nabla_v \left( \ell(x) + s \right) - \left( \Delta v \right)(x) - \nabla_v s \right),
\]

where in the last line we used the definitions of \( \ell \) and \( \Delta \) (see (7.1.1) and (8.1.6)). Applying the restricted EL equations (7.2.3) and the linearized field equations (8.1.6) gives the result.

We remark that the identity (9.3.1) has a similar structure as the conservation law in the Noether-like theorem (9.2.6). In order to make the connection precise, one describes the symmetry \( \Phi_\tau \) infinitesimally by a jet \( v \) with vanishing scalar component,

\[
v(x) := \frac{d}{d\tau}(0, \Phi_\tau(x)) \bigg|_{\tau=0}.
\]

Using the symmetry property (9.2.5), one verifies similar as in the proof of Lemma 8.2.1 that this jet satisfies the linearized field equations (8.1.6). Therefore, Proposition 9.3.1 applies, and the right side vanishes because \( v \) because \( v \) has no scalar component. We thus recover the identity obtained by carrying out the \( \tau \)-derivative in (9.2.6).

We conclude that Proposition 9.3.1 is a generalization of Theorem 9.2.3. Instead of imposing symmetries, the identity (9.3.1) is a consequence of the linearized field equations. Again choosing \( \Omega \) as the region between two Cauchy surfaces (see Figure 9.2), one obtains a relation between the surface layer integrals around \( N_1 \) and \( N_2 \). If the scalar component
of $v$ vanishes, we obtain a conservation law. Otherwise, the right side of (9.3.1) tells us how the surface layer integral changes in time.

We now generalize Proposition 9.3.1. The basic idea is to integrate anti-symmetric expressions in $x$ and $y$ which involve higher derivatives of the Lagrangian. We again restrict attention to the smooth setting (for the general proof see [69]). Let $\tilde{\rho}_{s,t}$ with $s,t \in (-\delta, \delta)$ be a two-parameter family of measures which are solutions of the restricted EL equations. We assume that these measures are of the form

$$\tilde{\rho}_{s,t} = (F_{s,t})^*(f_{s,t} \rho),$$

(9.3.2)

where $f_{s,t}$ and $F_{s,t}$ are smooth,

$$f \in C^\infty((-\delta, \delta)^2 \times \mathcal{F}, \mathbb{R}^+) \quad \text{and} \quad F \in C^\infty((-\delta, \delta)^2 \times \mathcal{F}, \mathcal{F}),$$

(9.3.3)

and are trivial in the case $s = t = 0$ (6.4.5). Moreover, we make the following regularity assumptions:

(r1) For all $x \in M$, $p,q \geq 0$ and $r \in \{0,1\}$, the following partial derivatives exist and can be interchanged with integration:

$$\int_M \partial^r \partial^p \partial^q L(F_{s+s',t}(x), F_{s,t}(y)) \bigg|_{s'=s=t=0} d\rho(y) = \partial^r \partial^p \partial^q \int_M L(F_{s+s',t}(x), F_{s,t}(y)) \bigg|_{s'=s=t=0} d\rho(y).$$

THEOREM 9.3.2. Let $f$ and $F$ be as in (9.3.3) and (6.4.5) which satisfy the above regularity assumption (r1). Moreover, assume that the measures $\tilde{\rho}_{s,t}$ given by (9.3.2) satisfy the restricted EL equations for all $s$ and $t$. Then for every compact $\Omega \subset M$ and every $k \in \mathbb{N}$,

$$I_{k+1}^\Omega := \int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) \times (\partial_{1,s} - \partial_{2,s})(\partial_{1,t} + \partial_{2,t})^k f_{s,t}(x) L(F_{s,t}(x), F_{s,t}(y)) f_{s,t}(y) \bigg|_{s=t=0}$$

$$= \delta \int_\Omega \partial_{s} \partial_{t}^k f_{s,t}(x) \bigg|_{s=t=0} d\rho(x).$$

(9.3.4)

PROOF. Introducing the short notation

$$L(x_{s,t}, y_{s,t}) = f_{s,t}(x) L(F_{s,t}(x), F_{s,t}(y)) f_{s,t}(y),$$

(9.3.5)

the restricted EL equations (8.1.6) read

$$\nabla_u \left( \int_M L(x_{s,t}, y_{s,t}) \quad d\rho(y) - \delta f_{s,t}(x) \right) = 0 \quad \text{for all } u \in \mathfrak{S}_0^\infty.$$ 

In particular for any $k \geq 0$ and any vector $v = v^s \partial_s + v^t \partial_t$, we obtain

$$\int_M \partial_{1,s}(\partial_{1,v} + \partial_{2,v})^k L(x_{s,t}, y_{s,t}) \quad d\rho(y) \bigg|_{s=t=0} = \delta \partial_{v}^k f_{s,t}(x) \bigg|_{s=t=0}$$

(9.3.6)

$$\int_M (\partial_{1,v} + \partial_{2,v})^{k+1} L(x_{s,t}, y_{s,t}) \quad d\rho(y) \bigg|_{s=t=0} = \delta \partial_{v}^{k+1} f_{s,t}(x) \bigg|_{s=t=0}$$

(9.3.7)

(the derivatives exist and can be exchanged with the integration according to the above regularity assumption (r1)). Differentiating the last equation with respect to $v^s$ and
dividing by \( k + 1 \), we obtain
\[
\int_M (\partial_{1,s} + \partial_{2,s})(\partial_{1,v} + \partial_{2,v})^k L(x_{s,t}, y_{s,t}) \, d\rho(y) = s \partial_s \partial_v^k f_{s,t}(x).
\]
Subtracting twice the identity (9.3.6), we obtain for any \( k \geq 0 \) the equation
\[
\int_M (\partial_{1,s} - \partial_{2,s})(\partial_{1,v} + \partial_{2,v})^k L(x_{s,t}, y_{s,t}) \, d\rho(y) = s \partial_s \partial_v^k f_{s,t}(x).
\]
Integrating the last equation over \( \Omega \) gives
\[
\hat{\Omega} d\rho(x) \hat{M} d\rho(y) (\partial_{1,s} - \partial_{2,s})(\partial_{1,v} + \partial_{2,v})^k L(x_{s,t}, y_{s,t}) = s \hat{\partial}_s \hat{\partial}_v^k f_{s,t}(x) d\rho(x).
\]
On the other hand, since the integrand is anti-symmetric in its arguments \( x \) and \( y \), we also know that
\[
\int_M d\rho(x) \int_M d\rho(y) (\partial_{1,s} - \partial_{2,s})(\partial_{1,v} + \partial_{2,v})^k L(x_{s,t}, y_{s,t}) = 0.
\]
Subtracting this equation from (9.3.8) and evaluating at \( s = t = 0 \) gives the result.

We conclude this section by discussing the conservation law of Proposition 9.3.1 for inner solutions as considered in Section 8.3 (commutator jets will be considered afterward in Section 9.4). To this end, we need to assume again that spacetime has a smooth manifold structure. We first define an integration measure on the boundary of \( \Omega \).

**Definition 9.3.3.** Let \( v = (\text{div } v, v) \in J_{\rho}^\text{in} \) be an inner solution and \( \Omega \subset M \) closed with smooth boundary \( \partial \Omega \). On the boundary, we define the measure \( d\mu(v, x) \) as the contraction of the volume form on \( M \) with \( v \), i.e. in local charts
\[
d\mu(v, x) = h \epsilon_{ijkl} v^i dx^j dx^k dx^l,
\]
where \( \epsilon_{ijkl} \) is the totally anti-symmetric Levi-Civita symbol (normalized by \( \epsilon_{0123} = 1 \)).

Similar as in Lemma 8.3.3, we can integrate by parts with the help of the Gauß divergence theorem. But now boundary terms remain,
\[
\gamma^\Omega_{\mu}(v) = \int_{\partial \Omega} d\mu(v, x) \int_M d\rho(y) \mathcal{L}(x, y) + \mathcal{L}(x, y) = s \int_{\partial \Omega} d\mu(v, x) = s \mu(v, \partial \Omega),
\]
where in the last line we used the symmetry of \( \mathcal{L} \) and employed the EL equations. In this way, the surface layer integral reduces to a usual surface integral over the hypersurface \( \partial \Omega \).

**9.4. The Commutator Inner Product for Causal Fermion Systems**

As a concrete example of a conservation law, we now consider current conservation. To this end, we consider the setting of causal fermion systems. As in Section 8.2 we again let \( \mathcal{A} \) be a symmetric operator of finite rank on \( \mathcal{H} \) and \( \mathcal{U}_\tau \) be the corresponding one-parameter family of unitary transformations (8.2.2). Infinitesimally, this one-parameter family is described by the commutator jet \( v \) (8.2.3). The unitary invariance
of the causal action implies that the commutator jets satisfy the linearized field equations (see Lemma 8.2.1). Moreover, using that the scalar component of commutator jets vanishes, Proposition 9.3.1 gives for any compact $\Omega \subset M$ the conservation law

$$\gamma^\Omega_\rho (v) := \int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) \left( \nabla_{1,y} - \nabla_{2,y} \right) L(x,y) = 0 .$$

In order to understand the significance of this conservation law, it is useful to choose $A$ more specifically as an operator of rank one. More precisely, given a non-zero vector $\psi \in H$, we form the symmetric linear operator

$$A \in \mathcal{L}(H)$$

of rank at one by

$$A u := \langle u | \psi \rangle H \psi$$

(9.4.1)

(thus in bra/ket notation, $A = |\psi\rangle\langle\psi|$). We now form the corresponding commutator jet (8.2.4). Varying the vector $\psi$, we obtain a mapping

$$j : H \to J_{\text{lin}}, \psi \mapsto v .$$

(9.4.2)

Moreover, we choose $\Omega$ again as the past of a Cauchy surface (as shown in Figure 9.2 (c)). We write the resulting conserved quantity as

$$C^\Omega_\rho (u) := \int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) \left( D_{1,j(u)} - D_{2,j(u)} \right) L(x,y) \quad \text{with} \quad u \in H ,$$

(9.4.3)

where for technical simplicity we assume smoothness in order to interchange differentiation with integration. Clearly, the mapping $j$ in (9.4.2), and consequently also the mapping $C^\Omega_\rho$, are homogeneous of degree two, i.e.

$$C^\Omega_\rho (\lambda u) = |\lambda|^2 C^\Omega_\rho (u) \quad \text{for all} \quad u \in H \quad \text{and} \quad \lambda \in \mathbb{C} .$$

Therefore, we can use the polarization formula to define a sesquilinear form on the Hilbert space $H$,

$$\langle u | v \rangle^\rho := \frac{1}{4} \left( C^\Omega_\rho (u + v) - C^\Omega_\rho (u - v) \right) - \frac{i}{4} \left( C^\Omega_\rho (u + iv) - C^\Omega_\rho (u - iv) \right)$$

This sesquilinear form is referred to as the commutator inner product (for details see [61, Section 3]). In [67, Section 5.2] it is shown that for Dirac systems in Minkowski space, this conserved quantity gives back the conservation of the Dirac current (up to an irrelevant prefactor). We thus recover current conservation as special case of a more general conservation law for causal fermion systems. Since in most known examples, the conserved surface layer integral $C^\Omega_\rho (u,v)$ gives back the Hilbert space scalar product, we give this property a name:

**Definition 9.4.1.** Given a critical measure $\rho$ and a subset $\Omega \subset M$, the surface layer integral $C^\Omega_\rho$ is said to represent the scalar product if there is a non-zero real constant $c$ such that for every symmetric operator $A$ of finite rank, the corresponding commutator jet as defined by (8.2.4) has the property that

$$\langle u | u \rangle^\rho_\rho = c \| u \|^2_{\mathcal{H}} .$$

(9.4.4)

In view of the conservation law of Proposition 9.3.1, this property remains valid if $\Omega$ is changed by a compact subset of $M$.

Although being true in most examples, at present there is no general argument why the surface layer integral $C^\Omega_\rho$ should represent the scalar product. Therefore, in this book we shall not assume that this property holds. Instead, we make the following weaker
assumption. We assume that the sesquilinear form $C^\Omega$ is equivalent to the scalar product in the sense that
\[ \langle u | v \rangle^\Omega = \langle u | C^\rho v \rangle_{\mathcal{H}}, \]
where $C^\rho$ is a bounded linear operator on $\mathcal{H}$ with bounded inverse. Under this assumption, the Hilbert space scalar product can be expressed by
\[ \langle u | v \rangle_{\mathcal{H}} = \langle u | C^{-1}_\rho v \rangle^\Omega. \]
In this way, the Hilbert space scalar product can be represented by a surface layer integral involving the physical wave functions in spacetime.

We finally remark that in [61, Section 4] the commutator product was extended to wave functions which are not necessarily the physical wave functions. Current conservation continues to hold for the extension, provided that the wave functions satisfy the so-called dynamical wave equation
\[ \int_M Q^{\text{dyn}}(x, y) \psi(y) \, d\rho(y) = 0. \] (9.4.5)
Here the integral kernel $Q^{\text{dyn}}$ is constructed from first variations of the causal Lagrangian. In this formulation, the commutator inner product takes the form
\[ \langle \psi | \phi \rangle^\rho = -2i \left( \int_{\Omega} d\rho(x) \int_{M \setminus \Omega} d\rho(y) - \int_{M \setminus \Omega} d\rho(x) \int_{\Omega} d\rho(y) \right) \langle \psi(x) | Q^{\text{reg}}(x, y) \phi(y) \rangle_x. \] (9.4.6)

9.5. The Symplectic Form and the Surface Layer Inner Product

For the applications, the most important surface layer integrals are $I^\Omega_1$ (also denoted by $\gamma^\Omega$; see Proposition [9.3.1] and Theorem [9.3.2] in the case $k = 0$) and $I^\Omega_2$ (see Theorem [9.3.2] in the case $k = 1$). We now have a closer look at the surface layer integral $I^\Omega_2$. It is defined by
\[ I^\Omega_2 = \int_{\Omega} d\rho(x) \int_{M \setminus \Omega} d\rho(y) \times (\partial_1 s - \partial_2 s)(\partial_{1, t} + \partial_{2, t}) f_{s, t}(x) L(F_{s, t}(x), F_{s, t}(y)) f_{s, t}(y) \bigg|_{s=t=0} \]
and satisfies for any compact subset $\Omega \subset M$ the identity
\[ I^\Omega_2 = s \int_{\Omega} \partial_1 f_{s, t}(x) \bigg|_{s=t=0} d\rho(x). \] (9.5.2)

These formulas simplify considerably if we anti-symmetrize in the parameters $s$ and $t$. Namely, the formula for $I^\Omega_2$ reduces to the surface layer integral
\[ \int_{\Omega} d\rho(x) \int_{M \setminus \Omega} d\rho(y) (\partial_{1, s} \partial_{2, t} - \partial_{1, t} \partial_{2, s}) f_{s, t}(x) L(F_{s, t}(x), F_{s, t}(y)) f_{s, t}(y) \bigg|_{s=t=0}. \]
Since this expression involves only first partial derivatives, it can be rewritten with jet derivatives as
\[ \sigma^\Omega_{\rho}(u, v) := \int_{\Omega} d\rho(x) \int_{M \setminus \Omega} d\rho(y) (\nabla_u \nabla_v - \nabla_v \nabla_u) L(x, y), \] (9.5.3)
where the jets $u$ and $v$ are the linearized solutions
\[ u = \partial_t f_{s, t} \bigg|_{s=t=0} \quad \text{and} \quad v = \partial_t f_{s, t} \bigg|_{s=t=0}. \] (9.5.4)
Moreover, the right side of (9.5.2) vanishes when anti-symmetrizing in $s$ and $t$. We conclude that

$$\sigma^\Omega_\rho(u,v) = 0 \quad \text{for every compact } \Omega \subset M.$$  

Choosing $\Omega$ again as explained in Figure 9.2, we obtain a conservation law for a surface layer integral over a neighborhood of a hypersurface $N$ which extends to spatial infinity. We refer to $\sigma^\Omega_\rho$ as the symplectic form (the connection to symplectic geometry will be explained after (9.5.6) below).

Symmetrizing $I^\Omega_2$ in the parameters $s$ and $t$ gives the surface layer integral

$$\int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) (\partial_1, s \partial_1, t \partial_2, t) f_{s,t}(x) \mathcal{L}(F_{s,t}(x), F_{s,t}(y)) f_{s,t}(y) \bigg|_{s=t=0}.$$  

This expression has a more difficult structure because it involves second partial derivatives. Such second partial derivatives cannot be expressed directly in terms of second jet derivatives, because the derivatives of the jets also need to be taken into account. In a differential geometric language, defining second derivatives would make it necessary to introduce a connection on $\mathcal{F}$. As explained after (ii) on page 102, we here use the simpler method of taking second partial derivatives in distinguished charts. Then it is useful to introduce the surface layer inner product $(\cdot, \cdot)^\Omega_\rho$ as the contribution to (9.5.5) involving second derivatives of the Lagrangian, i.e.

$$(u,v)^\Omega_\rho := \int_\Omega d\rho(x) \int_{M \setminus \Omega} d\rho(y) (\nabla_1, u \nabla_1, v \nabla_2, \nabla_2) \mathcal{L}(x,y),$$

where the jets $u$ and $v$ are again the linearized solutions (9.5.4). We point out that, in contrast to the symplectic form, the surface layer inner product does not correspond to a conservation law. This has two reasons: First because the right side of (9.5.2) gives rise to a volume term, and second because the derivatives of the jets $u$ and $v$ gives additional correction terms. For the details and the interpretation of these correction terms we refer to [69]. Here we only remark that the significance of the surface layer inner product is that it is an approximate conservation law. In particular, it can be used for estimating solutions of the linearized field equations and for proving existence. We will come back to these applications in Section 14.

We finally comment on the name symplectic form. Clearly, this name is taken from symplectic geometry, where it refers to a closed and non-degenerate two-form $\sigma$ on a manifold which we denote by $\mathcal{B}$. The connection to the surface layer integral (19.1.2) is obtained if we assume that the set of all critical measures of the form (8.1.1) forms a smooth manifold $\mathcal{B}$ (which may be a infinite-dimensional Banach manifold). In this case, a jet $v$ describing first variations of a measure (8.1.5) is a tangent vector in $T_\rho \mathcal{B}$. Consequently, the jet space $\mathcal{J}$ can be identified with the tangent space $T_\rho \mathcal{B}$. The surface layer integral (19.1.2) can be regarded as a mapping

$$\sigma^\Omega_\rho : T_\rho \mathcal{B} \times T_\rho \mathcal{B} \to \mathbb{R}.$$  

Being antisymmetric, it can be regarded as a two-form. Similarly, the conserved surface layer integral $\gamma^\Omega_\rho$ in (9.3.1) can be regarded as a one-form. Moreover, the $t$-derivative in (9.5.1) can be regarded as a directional derivative acting on $I^\Omega_1 = \gamma^\Omega_\rho$. Anti-symmetrizing in $s$ and $t$ corresponds to taking the outer derivative. We thus obtain

$$\sigma^\Omega_\rho = d\gamma^\Omega_\rho,$$
9.6. The Nonlinear Surface Layer Integral

We now introduce a different type of surface layer integral, which can be regarded as a generalization of the surface layer integrals considered so far. In order to explain the basic concept, we return to the general structure of a surface layer integral (9.1.2). The differential operator \((\cdots)\) in the integrand can be regarded as describing first or second variations of the measure \(\rho\). As we saw above, the resulting surface layer integrals give rise to conserved currents, the symplectic form and scalar products. Instead of considering first or second variations of a measure \(\rho\), we now consider an additional measure \(\tilde{\rho}\) which can be thought of as a finite perturbation of the measure \(\rho\). Consequently, we also have two spacetimes \(M := \text{supp} \rho\) and \(\tilde{M} := \text{supp} \tilde{\rho}\).

Choosing two compact subsets \(\Omega \subset M\) and \(\tilde{\Omega} \subset \tilde{M}\) of the corresponding spacetimes, we form the nonlinear surface layer integral by

\[
\gamma^{\tilde{\Omega},\Omega}(\tilde{\rho}, \rho) := \int_{\tilde{\Omega}} d\tilde{\rho}(x) \int_{M \setminus \Omega} d\rho(y) \mathcal{L}(x, y) - \int_{\Omega} d\rho(x) \int_{\tilde{M} \setminus \tilde{\Omega}} d\tilde{\rho}(y) \mathcal{L}(x, y). \tag{9.6.1}
\]

Note that one argument of the Lagrangian is in \(M\), whereas the other is in \(\tilde{M}\). Moreover, one argument lies inside the set \(\Omega\) respectively \(\tilde{\Omega}\), whereas the other argument lies outside. In this way, the nonlinear surface layer integral “compares” the two spacetimes near the boundaries of \(\Omega\) and \(\tilde{\Omega}\), as is illustrated in Figure 9.3. If \(\tilde{\rho}\) is a first or second variation of \(\rho\), one recovers surface layer integrals of the form (9.1.2). In this way, the nonlinear surface layer integral can be regarded as a generating functional for the previous surface layer integrals. Moreover, it has the advantage that it does not rely on continuous variations or a perturbative treatment. Instead, it can be used for comparing two arbitrary measures \(\rho\) and \(\tilde{\rho}\). This nonlinear surface layer integral was introduced in [59]. It plays a central role for getting the connection to quantum field theory (as will be outlined in Chapter 22).

The nonlinear surface layer integral comes with a corresponding conservation law, as we now explain. For technical simplicity, we assume that the measure \(\tilde{\rho}\) can be obtained from \(\rho\) by multiplication with a weight function and a push-forward, i.e.

\[
\tilde{\rho} = F_{\ast}(f \rho) \tag{9.6.2}
\]
with smooth functions \( f \in C^\infty(M, \mathbb{R}^+) \) and \( F \in C^\infty(M, \mathcal{F}) \). We use the mapping \( F \) in order to identify \( M \) with \( \tilde{M} \). In particular, we choose \( \tilde{\Omega} = F(\Omega) \).

Then, using the definition of the push-forward measure, the nonlinear surface layer integral can be written alternatively as

\[
\gamma_{\tilde{\Omega}, \Omega}(\tilde{\rho}, \rho) = \int_{\tilde{\Omega}} d\tilde{\rho}(x) \int_{M \setminus \tilde{\Omega}} d\rho(y) \left( f(x) \mathcal{L}(F(x), y) - \mathcal{L}(x, F(y)) f(y) \right).
\] (9.6.3)

Similar as explained in Section 9.2 in the connection of Noether-like theorems, by a “conservation law” we mean that the nonlinear surface layer integral should vanish for all compact \( \Omega \). In preparation for analyzing how to satisfy this condition, we rewrite the nonlinear surface layer integral as a volume integral by using the antisymmetry of the integrand in (9.6.3),

\[
\gamma_{\tilde{\Omega}, \Omega}(\tilde{\rho}, \rho) = \int_{\tilde{\Omega}} d\tilde{\rho}(x) \int_{M} d\rho(y) \left( f(x) \mathcal{L}(F(x), y) - \mathcal{L}(x, F(y)) f(y) \right).
\] (9.6.4)

In order to write this equation in a simpler form, we introduce a measure \( \nu \) on \( M \) and a measure \( \tilde{\nu} \) on \( \tilde{M} \) by

\[
d\nu(x) := \left( \int_{M} \mathcal{L}(x, y) d\tilde{\rho}(y) \right) d\rho(x) \quad \text{and} \quad d\tilde{\nu}(x) := \left( \int_{M} \mathcal{L}(x, y) d\rho(y) \right) d\tilde{\rho}(x).
\]

Intuitively speaking, these measures describe how the measures \( \rho \) and \( \tilde{\rho} \) are connected to each other by the Lagrangian. We refer to them as the correlation measures. Then we can rewrite (9.6.4) as

\[
\gamma_{\tilde{\Omega}, \Omega}(\tilde{\rho}, \rho) = \tilde{\nu}(F(\Omega)) - \nu(\Omega).
\]

In order to obtain a conservation law, this expression should vanish for all compact \( \Omega \). In other words, the measure \( \nu \) should be the push-forward of the measure \( \tilde{\nu} \) under the mapping \( F \),

\[
\nu = F_* \tilde{\nu}.
\]

In this way, the task of finding a conservation law is reduced to the following abstract problem: Given two measures \( \nu \) on \( M \) and \( \tilde{\nu} \) on \( \tilde{M} \), under which assumptions can one measure be realized as the push-forward of the other? If \( M \) and \( \tilde{M} \) are volume forms on compact manifolds, such a push-forward mapping is obtained from a classical theorem of Jürgen Moser (see for example [107] Section XVIII, §2). In the non-compact case, the existence of \( F \) has been proven under general assumptions in [93]. In this way, the conservation law for the nonlinear surface layer integral can be arranged by adjusting the identification of the spacetimes \( M \) and \( \tilde{M} \).

We finally remark how the nonlinear surface layer integral can be used to “compare” two causal fermion systems \( (\mathcal{H}, \mathcal{F}, \rho) \) and \( (\mathcal{H}, \tilde{\mathcal{F}}, \tilde{\rho}) \). In this setting, one must keep in mind that the causal fermion systems are defined on different Hilbert spaces. Therefore, before forming the nonlinear surface layer integral, we must identify the Hilbert space \( \mathcal{H} \) and \( \tilde{\mathcal{H}} \) by a unitary transformation \( V : \mathcal{H} \to \tilde{H} \). Since this identification is not unique, we are left with the freedom to transform \( V \) according to

\[
V \to V U \quad \text{with} \quad U \in \text{L}(\mathcal{H}) \text{ unitary}.
\]
A possible strategy for getting information independent of this freedom is to integrate over the unitary group. For example, this leads to the so-called partition function

$$Z^{\tilde{\Omega}, \Omega}(\tilde{\rho}, \rho) := \int_G e^{\beta \mathcal{F}^{\tilde{\Omega}, \Omega}(\tilde{\rho}, U\rho)} \, d\mu_G(U),$$

where $\beta$ is a real parameter, and $G$ is a compact subgroup of the unitary group on $\mathcal{H}$ with Haar measure $d\mu_G$. For more details we refer to [60] or Chapter 22.

### 9.7. Two-Dimensional Surface Layer Integrals

The surface layer integrals considered so far were intended to generalize integrals over hypersurfaces. We now explain how lower-dimensional integrals can be described by surface layer integrals. We restrict attention to two-dimensional integrals, noting that the methods can be applied similarly to one-dimensional integrals (i.e. integrals along a curve). It is most convenient to describe a two-dimensional surface $S \subset M$ as

$$S = \partial \Omega \cap \partial V,$$

where $\Omega$ can be thought of as being the past of a Cauchy surface, and $V$ describing a spacetime cylinder. This description has the advantage that the resulting surface layer integrals will be well-defined even in cases when spacetime is singular or discrete, in which case the boundaries $\partial \Omega$ and $\partial V$ are no longer a sensible concept. The most obvious way of introducing a surface layer integral localized in a neighborhood of $S$ is a double integral of the form

$$\int_{\partial \Omega \cap \partial V} \left( \int_{M \setminus (\Omega \cup V)} (\cdots) \mathcal{L}(x, y) \, d\rho(y) \right) \, d\rho(x) \quad (9.7.1)$$

(where $(\cdots)$ again stands for a differential operator acting on the Lagrangian). If the Lagrangian has short range, we only get contributions to this surface layer integral if both $x$ and $y$ are close to the two-dimensional surface $S$ (see the left of Figure 9.4).

The disadvantage of this method is that the surface layer integral (9.7.1) does not seem to fit together with the EL equations and the linearized field equations. Therefore, at present there is no corresponding conservation law. If one considers flows of two-surfaces, it seems preferable to use the following method introduced in [21]. We need to assume that $M$ has a smooth manifold structure and is four-dimensional (see Definition 8.3.1) and that $v$ is a vector field which is transverse to the hypersurface $\partial \Omega$ and tangential to $\partial V$ (see the right of Figure 9.4). Following Definition 9.3.3, the inner solution corresponding to $v$ gives rise to a volume measure $\mu$ on $\partial \Omega$. Thus we can introduce a two-dimensional surface layer integral by

$$A := \int_{\partial \Omega \cap \partial V} d\mu(v, x) \int_{M \setminus V} \, d\rho(y) (\cdots) \mathcal{L}(x, y).$$
Applying the Gauß divergence theorem, this surface layer integral can also be written in the usual way as a double spacetime integral involving jet derivatives of the inner solution,

\[
\hat{\Omega} \setminus V \int_{\Omega \setminus V} d\rho(x) \nabla_{x} \cdot \left( \int_{M \setminus V} d\rho(y) \left( \nabla_{1,s} \pm \nabla_{2,s} \right) \cdots \right) L(x, y) \quad \text{(9.7.2)}
\]

where the notation \( \pm \) means that the formula holds for either choice of the sign (this is because the corresponding term vanishes, as one sees after integrating by parts as in the proof of Lemma 8.3.3 and using that \( v \) is tangential to \( \partial V \)). The obtained surface layer integral (9.7.3) harmonizes with the structures of the EL equations and the linearized field equations, as is exemplified in [21] by a simple connection between area change and matter flux.

**Exercises**

**Exercise 9.1. (Noether-like theorems)** The goal of this exercise is to illustrate the Noether-like theorems mentioned in the lecture. In order to simplify the problem as far as possible, we consider the compact setting and assume furthermore that the Lagrangian is smooth, i.e. \( \mathcal{L} \in C^\infty(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+) \). Let \( \rho \) be a minimizer of the action under variations of \( \rho \) in the class of (positive) normalized regular Borel measures. Let \( u \in T\mathcal{F} \) be a vector field on \( \mathcal{F} \). Assume that \( u \) is a symmetry of the Lagrangian in the sense that

\[
\left( u(x)^i \frac{\partial}{\partial x^j} + u(y)^i \frac{\partial}{\partial y^j} \right) L(x, y) = 0 \quad \text{for all } x, y \in \mathcal{F}. \quad \text{(9.7.4)}
\]

Prove that for any measurable set \( \Omega \subset \mathcal{F} \),

\[
\int_{\Omega} d\rho(x) \int_{\mathcal{F} \setminus \Omega} d\rho(y) u(x)^j \frac{\partial}{\partial x^j} L(x, y) = 0.
\]

*Hint:* Integrate (9.7.4) over \( \Omega \times \Omega \). Transform the integral using the symmetry \( L(x, y) = L(y, x) \). Finally make use of the Euler-Lagrange equations.

**Exercise 9.2. (Commutator jets and conserved surface layer integrals)** Let \( (\mathcal{H}, \mathcal{F}, \rho) \) be a causal fermion system on a finite-dimensional Hilbert space. For any self-adjoint \( S \in L(\mathcal{H}) \), we define the corresponding commutator jet by

\[
\mathcal{C}_S := (0, \mathcal{C}_S), \quad \text{with} \quad \mathcal{C}_S(x) := i[S, x] \quad \text{for all } x \in \mathcal{F}.
\]

Prove the following identity between the conserved one-form and the conserved symplectic form:

\[
\gamma_\rho^\Omega((0, [\mathcal{C}_A, \mathcal{C}_B])) = -\frac{1}{2} \sigma_\rho^\Omega(\mathcal{C}_A, \mathcal{C}_B),
\]

where \([\mathcal{C}_A, \mathcal{C}_B]\) denotes the commutator of vector fields on \( \mathcal{F} \).

**Exercise 9.3. (On the surface-layer inner product)** The goal of this exercise to show that, under a suitable restriction of the jet space, the surface-layer inner product is indeed non-negative. On \( \mathcal{F} = \mathbb{R}^2 \) we define the Lagrangian

\[
L(x, y) = \frac{1}{2} \eta(x_1 - y_1) (x_2 - y_2)^2, \quad \text{where } \eta \in C^\infty_0(\mathbb{R}, \mathbb{R}^+).
\]
Let \( M = \mathbb{R} \subset \mathcal{F} \) equipped with the canonical measure one dimensional Lebesgue measure and consider the set of jets

\[
\mathcal{J} := \left\{ (0, u) \mid u = \sum_{i=1}^{2} u_i \partial_i \in T\mathcal{F} \text{ with } u_1(t, 0) = 0 \text{ and } \partial_1 u_2(t, 0) \leq 0 \text{ for all } t \in \mathbb{R} \right\}.
\]

Let \( \Omega_t := (-\infty, t) \subset M \). Show that the surface-layer inner product \((\cdot, \cdot)_{\Omega_t}|_{\mathcal{J} \times \mathcal{J}}\) is non-negative.

*Hint:* Remember that jets are never differentiated in expressions like \( \nabla_{i,u} \nabla_{j,u} \).
Positive Functionals Arising from Second Variations

10.1. Overview of the Positivity Results

Given a minimizer of a variational principle, second variations are always non-negative. This basic observation goes back to Legendre and Jacobi, who used it to analyze the question whether classical trajectories minimize the action and geodesics minimize arc length \([91]\). In classical field theory (like electrodynamics or general relativity), second variations are less useful because in these theories the action is unbounded from below, so that instead of minimizing one merely seeks for critical points of the action. However, for causal variational principles, the action is indeed bounded from below, and physical spacetime should be described by a minimizer. Therefore, second variations should give rise to positive functionals in spacetime. In this chapter, we show that this is indeed the case, and we also compute the resulting positive functionals.

For technical simplicity, we again restrict attention to causal variational principles in the noncompact smooth setting (see \((6.2.4)\) and Section \(6.3\)). We closely follow the constructions in \([51]\). For the variation we again consider families of measures \((\tilde{\rho}_\tau)_{\tau \in [0, \delta]}\) of the form \((8.1.1)\), but now the measures in this family do not need to be critical. But we assume that \(\rho = \tilde{\rho}_0\) should be a minimizer of the causal action principle. Moreover, the variation should satisfy the volume constraint \((6.3.1)\). For technical simplicity, we restrict attention to variations which are compactly supported (i.e. trivial outside a compact set; for details see Section \(10.3\)). Then the variation of the measure \((8.1.1)\) is described infinitesimally by a pair of a real-valued function and a vector field on \(M\) (for details see Section \(10.2\)), both with compact support,

\[
(a, u) := \left. \frac{d}{d\tau} (f_\tau, F_\tau) \right|_{\tau = 0} \in \mathfrak{J}_0^\infty := C_0^\infty(M, \mathbb{R}) \oplus C_0^\infty(M, T\mathfrak{F}) .
\]

We now state the main result of this chapter:

**Theorem 10.1.1.** Let \(\rho\) be a minimizer of a causal variational principle in the noncompact smooth setting (see \((6.2.4)\) and Section \(6.3\)). Then the following two quadratic functionals on \(\mathfrak{J}_0^\infty\) are positive:

\[
\int_M \nabla^2 \ell(x, u) \, d\rho(x) \geq 0 \quad (10.1.1)
\]

\[
\int_M d\rho(x) \int_M d\rho(y) \nabla_{1,u} \nabla_{2,u} \mathcal{L}(x, y) + \int_M \nabla^2 \ell(x, u) \, d\rho(x) \geq 0 . \quad (10.1.2)
\]

We remark that the following converse of this theorem also holds: Assume that \(\rho\) is a regular Borel measure which satisfies \((7.2.3)\) as well as the inequalities \((10.1.1)\) and \((10.1.2)\). If the inequality \((10.1.2)\) is strict for every non-zero \(u \in \mathfrak{J}_0^\infty\), then \(\rho\) is an isolated local minimum within the class of compactly supported variations with local fragmentation \((8.1.7)\). For the proof we refer to \([51\], Section 5].
The above theorem has various applications. First, the positive bilinear forms can be used to endow jet spaces with a Hilbert space structure where in the scalar product the jets are integrated over spacetime (see Section 10.4). Second, it gives rise to a positive surface layer integral, which in turn induces a spatial scalar product on the linearized solutions (see Section 10.5). Finally, as will be explained in Chapter 14 below, the above functionals are useful for getting energy estimates.

10.2. Positivity of the Hessian of $\ell$

Let $\rho$ be a minimizer of the causal action. According to the EL equations (7.1.2), the function $\ell$ is minimal on $M$. This clearly implies that its Hessian (as computed in any chart) is positive semi-definite, i.e.

$$D^2\ell(x) \geq 0 \quad \text{for all } x \in M := \text{supp } \rho.$$  \hspace{1cm} (10.2.1)

This is the first non-negative quantity obtained from the fact that $\rho$ is a minimizer.

In view of the restricted EL equations (7.2.3), the zero and first order derivatives of $\ell$ vanish for all $x \in M$. Adding such lower derivative terms, we can write (10.2.1) with jet derivatives as

$$\nabla^2\ell|_x(u,u) \geq 0 \quad \text{for all } x \in M,$$

where, following our conventions (i) and (ii) on page 102,

$$\nabla^2\ell|_x(u,u) := a(x)^2 \ell(x) + 2a(x) Du\ell(x) + D^2\ell|_x(u,u).$$

Integrating over $M$ gives the following result:

**Proposition 10.2.1.** Let $\rho$ be a minimizer of the causal action. Then

$$\int_M \nabla^2\ell|_x(u,u) \, d\rho(x) \geq 0 \quad \text{for all } u \in J^\infty_0.$$

10.3. Positivity of Second Variations Generated by Jets

We now analyze second variations for a special class of variations of the measure $\rho$ to obtain another positive functional on jets. Similar as in [68, Section 3] we consider measures of the form

$$\tilde{\rho}_\tau = (F_\tau)_* (f_\tau \rho) \quad \text{for } \tau \in (-\tau_{\text{max}}, \tau_{\text{max}})$$  \hspace{1cm} (10.3.1)

with smooth mappings

$$f \in C^\infty((-\tau_{\text{max}}, \tau_{\text{max}}) \times M, \mathbb{R}^+) \quad \text{and} \quad F \in C^\infty((-\tau_{\text{max}}, \tau_{\text{max}}) \times M, \mathcal{F}),$$

where the star denotes the push-forward measure defined by $((F_\tau)_* \mu)(\Omega) = \mu(F_\tau^{-1}(\Omega))$ (where $\Omega \subset \mathcal{F}$; for basics see for example [15, Section 3.6]). We assume that for $\tau = 0$ the variation is trivial, (8.1.2). Moreover, for technical simplicity we assume that $F_\tau$ and $f_\tau$ are compactly supported, meaning that they are trivial outside a compact set $K \subset M$, i.e.

$$f_\tau|_{M \setminus K} \equiv 1 \quad \text{and} \quad F_\tau|_{M \setminus K} \equiv 1.$$  \hspace{1cm} (10.3.2)

Finally, in order to satisfy the volume constraint on the right side of (6.3.1), we assume that

$$\int_K f_\tau(x) \, d\rho(x) = \rho(K) \quad \text{for all } \tau \in (-\tau_{\text{max}}, \tau_{\text{max}}).$$  \hspace{1cm} (10.3.3)

Then the transformation (10.3.1) is described infinitesimally by a smooth and compactly supported jet,

$$u = (a, u) := (f_0, F_0) \in J^\infty_0.$$
Moreover, differentiating the volume constraint (10.3.3) gives
\[ \int_K a(x) \, d\rho(x) = 0. \quad (10.3.4) \]

We now compute the first and second variation of the action. Combining (6.3.2) with the definition of the push-forward measure, we obtain
\[
S(\tilde{\rho}_\tau) - S(\rho) = 2 \int_K d\rho(x) \int_M d\rho(y) \left( f_{\tau}(x) \, L(F_{\tau}(x), y) \right) \\
+ \int_K d\rho(x) \int_M d\rho(y) \left( f_{\tau}(x) \, f_{\tau}(y) \, L(F_{\tau}(x), F_{\tau}(y)) - L(x, y) \right).
\]

Then the first variation vanishes,
\[ \frac{d}{d\tau} S(\tilde{\rho}_\tau) \Big|_{\tau=0} = 2 \int_K d\rho(x) \int_M d\rho(y) \nabla_{1,u} L(x, y) = 2 \int_K \nabla_u (\ell(x) + s) \, d\rho(x) = 0, \]
where in the last step we used (7.2.3) and (10.3.4) (and similarly for \( \dot{\tilde{\rho}}_0 \)).

Next, the second variation is computed by
\[ \frac{d^2}{d\tau^2} S(\tilde{\rho}_\tau) \Big|_{\tau=0} = 2 \int_K d\rho(x) \int_M d\rho(y) \nabla_{1,u} \nabla_{2,u} L(x, y) \\
+ 2 \int_K d\rho(x) \int_M d\rho(y) \left( (a(x) \, D_{1,u} L(x, y) + D_{1,u} D_{1,u} L(x, y) + (\ddot{f}_0(x) + D_{1,\tilde{\rho}_0}) L(x, y) \right). \]

In the last line we can carry out the \( y \)-integration using (7.1.1). Again combining the EL equations (7.2.3) with (10.3.4) (and similarly for \( \dot{\tilde{f}} \)), we obtain the simple formula
\[ \frac{1}{2} \frac{d^2}{d\tau^2} S(\tilde{\rho}_\tau) \Big|_{\tau=0} = \int_M d\rho(x) \int_M d\rho(y) \nabla_{1,u} \nabla_{2,u} L(x, y) + \int_M \nabla^2 \ell|_x(u, u) \, d\rho(x). \quad (10.3.5) \]

Since \( \rho \) is a minimizer and the first variation vanishes, the second variation is necessarily non-negative, giving rise to the inequality
\[ \int_M d\rho(x) \int_M d\rho(y) \nabla_{1,u} \nabla_{2,u} L(x, y) + \int_M \nabla^2 \ell|_x(u, u) \, d\rho(x) \geq 0, \quad (10.3.6) \]
subject to the condition that the jet \( u \) must satisfy the volume constraint (10.3.4). In the next proposition we remove this condition with a limiting procedure:

**PROPOSITION 10.3.1.** Let \( \rho \) be a minimizer of the causal action. Then
\[ \int_M d\rho(x) \int_M d\rho(y) \nabla_{1,u} \nabla_{2,u} L(x, y) + \int_M \nabla^2 \ell|_x(u, u) \, d\rho(x) \geq 0 \quad \text{for all } u \in J_0^\infty. \]

**Proof.** Let \( u = (a, u) \in J_0^\infty \) be a jet which violates the volume constraint (10.3.4). Then, choosing a compact set \( \Omega \subset M \) with \( \rho(\Omega) > 0 \), the jet \( \hat{u} := (\hat{a}, u) \) with
\[ \hat{a}(x) = a(x) - e(\Omega) \, \chi_\Omega(x) \quad \text{and} \quad e(\Omega) := \frac{1}{\rho(\Omega)} \int_\Omega a(x) \, d\rho(x) \quad (10.3.7) \]
(where \( \chi_\Omega \) is the characteristic function) does satisfy (10.3.4). Choosing the scalar variation \( f_{\tau} = (1 - \tau) + \tau \hat{a} \) and a family of diffeomorphisms \( F_{\tau} \) with \( F_0 = u \), we obtain a variation which satisfies the volume constraint (10.3.3) (note that \( \dot{f} = 0 \)). Clearly, due to the characteristic function, the jet \( \hat{u} \) is no longer smooth, but it has again compact support, and an approximation argument using Lebesgue’s dominated convergence theorem.
shows that the inequality \((10.3.6)\) also holds for \(\hat{u}\). Expanding in powers of \(c\), we thus obtain the inequality
\[
0 \leq \int_K d\rho(x) \int_K d\rho(y) \nabla_{1,u} \nabla_{2,u} \mathcal{L}(x,y) + \int_K \nabla^2 \ell|_x(u,u) \, d\rho(x) - 2c \int_M d\rho(x) \int_K d\rho(y) \chi_\Omega(x) \nabla_{2,u} \mathcal{L}(x,y) + c^2 \int_M d\rho(x) \int_K d\rho(y) \chi_\Omega(x) \chi_\Omega(y) \mathcal{L}(x,y) + \int_M \left( -2c \chi_\Omega(x) \nabla_u \ell(x) + c^2 \chi_\Omega(x)^2 \ell(x) \right) \, d\rho(x)
\]
(the integrand in the last line arises from the contributions to \(\nabla^2 \ell|_x(u,u)\) involving the scalar components of the jets). The last line vanishes due to the restricted EL equations \((7.2.3)\). Hence
\[
\int_K d\rho(x) \int_K d\rho(y) \nabla_{1,u} \nabla_{2,u} \mathcal{L}(x,y) + \int_K \nabla^2 \ell|_x(u,u) \, d\rho(x) - 2c \int_M d\rho(x) \int_K d\rho(y) \nabla_{2,u} \mathcal{L}(x,y) + c^2 \int_M d\rho(x) \int_K d\rho(y) \chi_\Omega(x) \chi_\Omega(y) \mathcal{L}(x,y) =: A(\Omega).
\]

We now let \((\Omega_n)_{n \in \mathbb{N}}\) be an exhaustion of \(M\) by compact sets. We distinguish the two cases when \(\rho(M)\) is finite and infinite and treat these cases separately. If the total volume \(\rho(M)\) is finite, one can take the limit \(n \to \infty\) with Lebesgue’s dominated convergence theorem to obtain
\[
\lim_{n \to \infty} \int_K d\rho(x) \int_{\Omega_n} d\rho(y) \nabla_{1,u} \mathcal{L}(x,y) = \int_K d\rho(x) \int_M d\rho(y) \nabla_{1,u} \mathcal{L}(x,y) = \int_K \nabla_u \left( \ell(x) + s \right) \, d\rho(x) = s \int_K a(x) \, d\rho(x)
\]
\[
\lim_{n \to \infty} A(\Omega_n) = 2c(M)s \int_K a(x) \, d\rho(x) - c(M)^2 \rho(M)s = \frac{\nu}{2\rho(M)} \left( \int_K a(x) \, d\rho(x) \right)^2 \geq 0,
\]
where in the last line we substituted the value of \(c(M)\) in \((10.3.7)\).

In the remaining case that the volume \(\rho(M)\) is infinite, we estimate the terms as follows,
\[
c(\Omega_n)^2 \int_K d\rho(x) \int_K d\rho(y) \mathcal{L}(x,y) \leq c(\Omega_n)^2 \int_K d\rho(x) \int_M d\rho(y) \mathcal{L}(x,y) = c(\Omega_n)^2 s \rho(K) \to 0
\]
\[
\int_K d\rho(x) \int_{\Omega_n} d\rho(y) \nabla_{1,u} \mathcal{L}(x,y) \to \int_K d\rho(x) \int_M d\rho(y) \nabla_{1,u} \mathcal{L}(x,y) = \int_K \nabla_u \left( \ell(x) + s \right) \, d\rho(x) = s \int_K a(x) \, d\rho(x).
\]
As a consequence, \(A(\Omega_n)\) converges to zero as \(n \to \infty\). This concludes the proof. \(\square\)
Proof of Theorem 10.1.1. The inequalities (10.1.1) and (10.1.2) were derived in Propositions 10.2.1 and 10.3.1. □

10.4. Application: Hilbert Spaces of Jets

As an application, we now explain how our positive functionals can be used to endow spaces of jets in spacetime with Hilbert space structures. These Hilbert space structures should be very useful because they make functional analytic tools applicable to the analysis of the jet spaces and the causal action principle.

We introduce the following bilinear forms on $J^\infty_0$,

\[ \langle u, v \rangle := \int_M d\rho(x) \int_M d\rho(y) \nabla_{1,u} \nabla_{2,v} \mathcal{L}(x, y) + \int_M \nabla^2 \ell\big|_x (u, v) \, d\rho(x) \]  
\[ \langle \langle u, v \rangle \rangle := \langle u, v \rangle + \int_M \nabla^2 \ell\big|_x (u, v) \, d\rho(x) \].

(10.4.1)

(10.4.2)

By Theorem 10.1.1 both bilinear forms are positive semi-definite. Thus dividing out the null space and forming the completion gives real Hilbert spaces of jets denoted by $\mathcal{H}\langle \rangle$ and $\mathcal{H}\langle \langle \rangle \rangle$, respectively. Obviously,

\[ \langle u, u \rangle \leq \langle \langle u, u \rangle \rangle , \]

giving rise to a norm-decreasing mapping $\mathcal{H}\langle \langle \rangle \rangle \to \mathcal{H}\langle \rangle$.

For the scalar components of the jets, the two scalar products (10.4.1) and (10.4.2) obviously agree. But they are quite different for the vector components. In order to understand this difference, it is instructive to consider a jet $u = (0, u)$ which describes a symmetry of the Lagrangian, i.e. (for details see [67, Section 3.1])

\[ (D_{1,u} + D_{2,u}) \mathcal{L}(x, y) = 0 \quad \text{for all } x, y \in M . \]

For this jet, a direct computation shows that

\[ \langle u, u \rangle = 0 . \]

Hence symmetry transformations lie in the kernel of the bilinear form $\langle \langle \rangle \rangle$ and thus correspond to the zero vector in $\mathcal{H}\langle \rangle$. Generally speaking, the scalar product $\langle \langle \rangle \rangle$ makes it possible to disregard symmetry transformations of the causal Lagrangian. However, jets describing symmetry transformations do in general correspond to non-zero vectors of the Hilbert space $\mathcal{H}\langle \langle \rangle \rangle$.

10.5. Application: A Positive Surface Layer Integral

We now derive a surface layer integral which is not necessarily conserved, but which has a definite sign. Similar as explained at the beginning of Section 10.4, this can be used to endow the jet space with a Hilbert structure. But in contrast to the scalar products in Section 10.4, where the jets were integrated over spacetime, here the scalar product is given as a surface layer integral. This should be very useful for analyzing the dynamics of jets in spacetime.

Proposition 10.5.1. Assume that $v$ is a solution of the linearized field equations (8.1.6). Then for any compact $\Omega \subset M$, the following surface layer integral is positive,

\[ -\int_{\Omega} d\rho(x) \int_{M \setminus \Omega} d\rho(y) \nabla_{1,v} \nabla_{2,v} \mathcal{L}(x, y) \geq 0 . \]
Proof. Denoting the components of $v$ by $v = (b, \psi)$, we evaluate (8.1.6) for $u = v$ and integrate over $\Omega$. The resulting integrals can be rewritten as follows,

\[
0 = \int_{\Omega} d\rho(x) \int_{M} d\rho(y) \nabla_{1,v} \left( \nabla_{1,v} + \nabla_{2,v} \right) \mathcal{L}(x,y) - s \int_{\Omega} b(x)^2 \, d\rho(x)
\]

\[
= \int_{\Omega} \nabla^2 \ell|_{x}(v, v) \, d\rho(x) + \int_{M} d\rho(x) \int_{\Omega} d\rho(y) \nabla_{1,v} \nabla_{2,v} \mathcal{L}(x,y)
\]

\[
= \int_{\Omega} \nabla^2 \ell|_{x}(v, v) \, d\rho(x) + \int_{M} d\rho(x) \int_{\Omega} d\rho(y) \nabla_{1,v} \nabla_{2,v} \mathcal{L}(x,y)
\]

\[
+ \int_{\Omega} d\rho(x) \int_{M \setminus \Omega} d\rho(y) \nabla_{1,v} \nabla_{2,v} \mathcal{L}(x,y). \tag{10.5.1}
\]

Using characteristic functions, the expression (10.5.1) can be written as

\[
\int_{M} \nabla^2 \ell|_{x}(\chi_{\Omega} v, \chi_{\Omega} v) \, d\rho(x) + \int_{M} d\rho(x) \int_{M} d\rho(y) \nabla_{1,\chi_{\Omega} v} \nabla_{2,\chi_{\Omega} v} \mathcal{L}(x,y). \tag{10.5.2}
\]

Approximating the jet $\chi_{\Omega} v$ by smooth jets with compact support, one finds that the integrals in (10.5.1) are non-negative by Proposition 10.3.1. This gives the result. \qed

We finally remark that in [52, Section 6] the surface layer integral in the last proposition is computed in Minkowski space.

Exercises
CHAPTER 11

Topological and Geometric Structures

This chapter is devoted to the topological and geometric structures of a causal fermion system. We closely follow the presentation in [58] and [55].

11.1. A Topological Vector Bundle

For the topological structures, it is not essential that the operators in $\mathcal{F}$ have at most $n$ positive and at most $n$ negative eigenvalues (see Definition 5.3.1). Instead, it is preferable for the sake of greater generality and broader applicability relax this condition in the following way.

**Definition 11.1.1.** Given a complex Hilbert space $(\mathcal{H}, \langle \cdot, \cdot \rangle_\mathcal{H})$ and parameters $p, q \in \mathbb{N}_0$ with $p \leq q$, we let $\mathcal{F} \subset L(\mathcal{H})$ be the set of all self-adjoint operators on $\mathcal{H}$ of finite rank, which (counting with multiplicities) have at most $p$ positive and at most $q$ negative eigenvalues. On $\mathcal{F}$ we are given a positive measure $\rho$ (defined on a $\sigma$-algebra of subsets of $\mathcal{F}$). We refer to $(\mathcal{H}, \mathcal{F}, \rho)$ as a topological fermion system of spin signature $(p, q)$.

If $p = 0$, we call $(\mathcal{H}, \mathcal{F}, \rho)$ a Riemannian fermion system of spin dimension $n := q$.

Clearly, the case $p = q$ gives back a causal fermion system (see Definition 5.3.1). It should be noted that the assumption $p \leq q$ merely is a convention, because otherwise one may always arrange to replace $\mathcal{F}$ by $-\mathcal{F}$.

In Section 2.5 the notion of a topological vector bundle was introduced (see Definition 2.5.2). Again setting $M = \text{supp} \rho$, we want to construct a topological vector bundle having the spin space $S_x$ as the fiber at the point $x \in M$. To this end, all the spin spaces must have the same dimension and signature, making it necessary to impose the following condition:

**Definition 11.1.2.** The topological fermion system is called **regular** if for all $x \in M$, the operator $x$ has the maximal possible rank $p + q$.

We note that most of our constructions can be extended to non-regular topological fermion systems by decomposing $M$ into subsets on which $x$ has fixed rank and a fixed number of positive and negative eigenvalues (for details see [58 Section 7]).

We define $\mathcal{B}$ as the set of pairs

$$\mathcal{B} = \{(x, \psi) \mid x \in M, \psi \in S_x\}$$

and let $\pi$ be the projection onto the first component. Moreover, we let $(Y, \langle \cdot, \cdot \rangle_Y)$ be an indefinite inner product space of signature $(q, p)$, and choose $G = U(q, p)$ as the group of unitary transformations on $Y$. In order to construct the bundle charts, for any given $x \in M$ we choose a unitary mapping $\sigma : S_x \to Y$. By restricting the projection $\pi_x$ in (5.6.1) to $S_y$, we obtain the mapping

$$\pi_x|_{S_y} : S_y \to S_x.$$
In order to compute its adjoint with respect to the spin scalar product (2.5.2), for \( \psi \in S_x \) and \( \phi \in S_y \) we make the computation
\[
\langle \psi | \pi_x | S_y \phi \rangle_x = -(\langle \psi | x \phi \rangle_{\mathcal{H}}) = -\langle x \psi | \phi \rangle_{\mathcal{H}} = -\langle \pi_y x \psi | \phi \rangle_{\mathcal{H}} = -\langle y (y | S_y)^{-1} \pi_y x \psi | \phi \rangle_{\mathcal{H}} = -\langle (y | S_y)^{-1} \pi_y x \psi | y \rangle_{\mathcal{H}}.
\]
Hence
\[
(\pi_x | S_y)^* = (y | S_y)^{-1} \pi_y x | S_x.
\]
We now introduce the operator
\[
T_{xy} = (\pi_x | S_y)(\pi_x | S_y)^* = \pi_x (y | S_y)^{-1} \pi_y x | S_x : S_x \to S_x.
\]
By construction, this operator is symmetric and \( T_{xx} = 1 \). We now form the polar decomposition of \( T_{xy} \) to obtain a unitary operator \( U_{xy} : S_x \to S_x \). By continuity, there is a neighborhood \( U \) of \( x \) such that for all \( y \in U \), the operator \( T_{xy} \) is invertible and has a unique square root \( \rho_{xy} \) (defined for example by the power series \( \sqrt{T_{xy}} = \sqrt{1 + (T_{xy} - 1)} = 1 + \frac{1}{2}(T_{xy} - 1) + \cdots \)). Introducing the mapping
\[
U_{xy} = \rho_{xy}^{-1} \pi_x | S_y : S_y \to S_x,
\]
the calculation
\[
U_{xy} U_{xy}^* = \rho_{xy}^{-1} \pi_x | S_y (\pi_x | S_y)^* \rho_{xy}^{-1} = \rho_{xy}^{-1} T_{xy} \rho_{xy}^{-1} = 1 | S_x
\]
shows that the mapping \( U_{xy} \) is unitary. Moreover, it clearly depends continuously on \( y \in U \).

We define the bundle chart \( \phi_U \) by
\[
\phi_U(y, v) = (y, (\sigma \circ U_{xy})(v)).
\]
The commutativity of the diagram (2.5.1) is clear by construction. Moreover, the transition functions \( g_{UV} \) in (2.5.2) are in \( G \) because we are working with unitary mappings of the fibers throughout. We choose the topology on \( \mathcal{B} \) such that all the bundle charts are homeomorphisms.

**Definition 11.1.3.** The topological vector bundle \( \mathcal{B} \to \mathcal{M} \) is referred to as the vector bundle associated to the regular topological fermion system \((\mathcal{H}, \mathcal{F}, \rho)\), or simply the associated vector bundle.

The next result shows that every vector bundle over a manifold can be realized as the associated vector bundle of a corresponding topological fermion system.

**Theorem 11.1.4.** Let \( X \to \mathcal{M} \) be a vector bundle over a \( k \)-dimensional topological manifold \( \mathcal{M} \), whose fibers are isomorphic to an indefinite inner product space of signature \((q, p)\). Then there is a regular topological fermion system \((\mathcal{H}, \mathcal{F}, \rho)\) of signature \((p, q)\) such that the associated vector bundle (see Definition 11.1.3) is isomorphic to \( X \). If \( \mathcal{M} \) is compact, the Hilbert space \( \mathcal{H} \) can be chosen to be finite-dimensional.

The proof can be found in [58, Section 3.3].

### 11.2. Geometric Structures of a Causal Fermion System

We now outline constructions from [58] which give general notions of a connection and curvature (see Theorem 11.2.9, Definition 11.2.10 and Definition 11.2.11). So far, these constructions have been carried out only in the case of spin dimension \( n = 2 \). This is the most important case because it allows for the description of Dirac spinors in a four-dimensional spacetime.
11.2.1. Construction of the Spin Connection. Let \((\mathcal{H}, \mathcal{F}, \rho)\) be a causal fermion system of spin dimension \(n = 2\). Moreover, we assume that it is regular (see Definition 11.1.2).

An important structure from spin geometry missing so far is Clifford multiplication. To this end, we need a Clifford algebra represented by symmetric operators on \(S_x\). For convenience, we first consider Clifford algebras with the maximal number of five generators; later we reduce to four spacetime dimensions (see Definition 11.2.13 below). We denote the set of symmetric linear endomorphisms of \(S_x\) by \(\text{Symm}(S_x)\); it is a 16-dimensional real vector space.

**Definition 11.2.1.** A five-dimensional subspace \(K \subset \text{Symm}(S_x)\) is called a Clifford subspace if the following conditions hold:

(i) For any \(u, v \in K\), the anti-commutator \(\{u, v\} \equiv uv + vu\) is a multiple of the identity on \(S_x\).

(ii) The bilinear form \(\langle ., . \rangle\) on \(K\) defined by

\[
\frac{1}{2} \{u, v\} = \langle u, v \rangle \mathbf{1} \quad \text{for all } u, v \in K
\]

is non-degenerate and has signature \((1, 4)\).

In view of the situation in spin geometry, we would like to distinguish a specific Clifford subspace. In order to partially fix the freedom in choosing Clifford subspaces, it is useful to impose that \(K\) should contain a given so-called sign operator.

**Definition 11.2.2.** An operator \(v \in \text{Symm}(S_x)\) is called a sign operator if \(v^2 = \mathbf{1}\) and if the inner product \(\langle . | v . \rangle : S_x \times S_x \to \mathbb{C}\) is positive definite.

**Definition 11.2.3.** Given a sign operator \(v\), the set of Clifford extensions \(T^v\) is defined as the set of all Clifford subspaces containing \(v\),

\[
T^v = \{ K \text{ Clifford subspace with } v \in K \}.
\]

Considering \(x\) as an operator on \(S_x\), this operator has by definition of the spin dimension two positive and two negative eigenvalues. Moreover, the calculation

\[
\langle u | (-x) u \rangle_x = \langle u | x^2 u \rangle_\mathcal{F} > 0 \quad \text{for all } u \in S_x \setminus \{0\}
\]

shows that the operator \((-x)\) is positive definite on \(S_x\). Thus we can introduce a unique sign operator \(s_x\) by demanding that the eigenspaces of \(s_x\) corresponding to the eigenvalues \(\pm 1\) are precisely the positive and negative spectral subspaces of the operator \((-x)\). This sign operator is referred to as the Euclidean sign operator.

A straightforward calculation shows that for two Clifford extensions \(K, \tilde{K} \in T^v\), there is a unitary transformation \(U \in e^{i\mathbb{R}v}\) such that \(\tilde{K} = UKU^{-1}\) (for details see [55, Section 3]). By dividing out this group action, we obtain a five-dimensional vector space, endowed with the inner product \(\langle ., . \rangle\). Taking for \(v\) the Euclidean signature operator, we regard this vector space as a generalization of the usual tangent space.

**Definition 11.2.4.** The tangent space \(T_x\) is defined by

\[
T_x = T^v_x / \exp(i\mathbb{R}s_x).
\]

It is endowed with an inner product \(\langle ., . \rangle\) of signature \((1, 4)\).

We next consider two spacetime points. We need the following assumption.
Definition 11.2.5. Two points \( x, y \in M \) are said to be **properly time-like separated** if the closed chain \( A_{xy} \) has a strictly positive spectrum and if the corresponding eigenspaces are definite subspaces of \( S_x \).

This definition clearly implies that \( x \) and \( y \) are time-like separated (see Definition 5.5.1). Moreover, the eigenspaces of \( A_{xy} \) are definite if and only if those of \( A_{yx} \) are, showing that Definition 11.2.5 is again symmetric in \( x \) and \( y \). As a consequence, the spin space can be decomposed uniquely into an orthogonal direct sum \( S_x = I^+ \oplus I^- \) of a positive definite subspace \( I^+ \) and a negative definite subspace \( I^- \) of \( A_{xy} \). This allows us to introduce a unique sign operator \( v_{xy} \) by demanding that its eigenspaces corresponding to the eigenvalues \( \pm 1 \) are the subspaces \( I^\pm \). This sign operator is referred to as the **directional sign operator** of \( A_{xy} \). Having two sign operators \( s_x \) and \( v_{xy} \) at our disposal, we can distinguish unique corresponding Clifford extensions, provided that the two sign operators satisfy the following generic condition.

Definition 11.2.6. Two sign operators \( v, \tilde{v} \) are said to be **generically separated** if their commutator \([v, \tilde{v}]\) has rank four.

Lemma 11.2.7. Assume that the sign operators \( s_x \) and \( v_{xy} \) are generically separated. Then there are unique Clifford extensions \( K^{(y)}_x \in T^{s_x} \) and \( K_{xy} \in T^{v_{xy}} \) and a unique operator \( \rho \in K^{(y)}_x \cap K_{xy} \) with the following properties:

(i) The relations \( \{s_x, \rho\} = 0 = \{v_{xy}, \rho\} \) hold.

(ii) The operator \( U_{xy} := e^{i\rho} \) transforms one Clifford extension to the other,

\[
K_{xy} = U_{xy} K^{(y)}_x U^{-1}_{xy}.
\]

(iii) If \( \{s_x, v_{xy}\} \) is a multiple of the identity, then \( \rho = 0 \).

The operator \( \rho \) depends continuously on \( s_x \) and \( v_{xy} \).

We refer to \( U_{xy} \) as the **synchronization map**. Exchanging the roles of \( x \) and \( y \), we also have two sign operators \( s_y \) and \( v_{yx} \) at the point \( y \). Assuming that these sign operators are again generically separated, we also obtain a unique Clifford extension \( K_{yx} \in T^{v_{yx}} \).

After these preparations, we can now explain the construction of the spin connection \( D \) (for details see [55, Section 3]). For two spacetime points \( x, y \in M \) with the above properties, we want to introduce an operator

\[
D_{x,y} : S_y \to S_x
\]

(generally speaking, by the subscript \( xy \) we always denote an object at the point \( x \), whereas the additional comma \( x,y \) denotes an operator which maps an object at \( y \) to an object at \( x \)). It is natural to demand that \( D_{x,y} \) is unitary, that \( D_{y,x} \) is its inverse, and that these operators map the directional sign operators at \( x \) and \( y \) to each other,

\[
D_{x,y} = (D_{y,x})^* = (D_{y,x})^{-1} \tag{11.2.1}
\]

\[v_{xy} = D_{x,y} v_{yx} D_{y,x} \tag{11.2.2}\]

The obvious idea for constructing an operator with these properties is to take a polar decomposition of \( P(x,y) \); this amounts to setting

\[
D_{x,y} = A_{xy}^{-1/2} P(x,y) \tag{11.2.3}\]

This definition has the shortcoming that it is not compatible with the chosen Clifford extensions. In particular, it does not give rise to a connection on the corresponding
11.2. GEOMETRIC STRUCTURES OF A CAUSAL FERMION SYSTEM

In order to resolve this problem, we modify (11.2.3) by the ansatz

$$D_{x,y} = e^{i\varphi_{xy} v_{xy}} a_{xy}^{-\frac{1}{2}} P(x,y)$$

(11.2.4)

with a free real parameter $\varphi_{xy}$. In order to comply with (11.2.1), we need to demand that

$$\varphi_{xy} = -\varphi_{yx} \mod 2\pi$$

(11.2.5)

then (11.2.2) is again satisfied. We can now use the freedom in choosing $\varphi_{xy}$ to arrange that the distinguished Clifford subspaces $K_{xy}$ and $K_{yx}$ are mapped onto each other,

$$K_{xy} = D_{x,y} K_{yx} D_{y,x}.$$ 

(11.2.6)

It turns out that this condition determines $\varphi_{xy}$ up to multiples of $\frac{\pi}{2}$. In order to fix $\varphi_{xy}$ uniquely in agreement with (11.2.5), we need to assume that $\varphi_{xy}$ is not a multiple of $\frac{\pi}{4}$. This leads us to the following definition.

**Definition 11.2.8.** Two points $x, y \in M$ are called **spin connectable** if the following conditions hold:

(a) The points $x$ and $y$ are properly timelike separated (note that this already implies that $x$ and $y$ are regular as defined in Section 11.2.1).

(b) The Euclidean sign operators $s_x$ and $s_y$ are generically separated from the directional sign operators $v_{xy}$ and $v_{yx}$, respectively.

(c) Employing the ansatz (11.2.4), the phases $\varphi_{xy}$ which satisfy condition (11.2.6) are not multiples of $\frac{\pi}{4}$.

We denote the set of points which are spin connectable to $x$ by $\mathcal{I}(x)$. It is straightforward to verify that $\mathcal{I}(x)$ is an open subset of $M$.

Under these assumptions, we can fix $\varphi_{xy}$ uniquely by imposing that

$$\varphi_{xy} \in \left(-\frac{\pi}{2}, -\frac{\pi}{4}\right) \cup \left(\frac{\pi}{4}, \frac{\pi}{2}\right),$$

(11.2.7)

giving the following result (for the proofs see [55 Section 3.3]).

**Theorem 11.2.9.** Assume that two points $x, y \in M$ are spin connectable. Then there is a unique spin connection $D_{x,y} : S_y \to S_x$ of the form (11.2.4) having the properties (11.2.1), (11.2.2), (11.2.6) and (11.2.7).

11.2.2. The Metric Connection and Curvature. We now outline a few further constructions from [55 Section 3]. The spin connection induces a connection on the corresponding tangent spaces, as we now explain. Suppose that $u_y \in T_y$. Then, according to Definition 11.2.4 and Lemma 11.2.7, we can consider $u_y$ as a vector of the representative $K^{(y)}(x) \in T^s_y$. By applying the synchronization map, we obtain a vector in $K_{yx}$,

$$u_{yx} := U_{yx} u_y U_{yx}^{-1} \in K_{yx}.$$ 

According to (11.2.6), we can now “parallel transport” the vector to the Clifford subspace $K_{xy}$,

$$u_{xy} := D_{x,y} u_{yx} D_{y,x} \in K_{xy}.$$ 

Finally, we apply the inverse of the synchronization map to obtain the vector

$$u_x := U_{xy}^{-1} u_{xy} U_{xy} \in K^{(y)}_x.$$ 

As $K^{(y)}_x$ is a representative of the tangent space $T_x$ and all transformations were unitary, we obtain an isometry from $T_y$ to $T_x$. 
11. TOPOLOGICAL AND GEOMETRIC STRUCTURES

**Definition 11.2.10.** The isometry between the tangent spaces defined by
\[ \nabla_{x,y} : T_y \to T_x : u_y \mapsto u_x \]
is referred to as the **metric connection** corresponding to the spin connection \( D \).

We next introduce a notion of curvature.

**Definition 11.2.11.** Suppose that three points \( x, y, z \in M \) are pairwise spin connectable. Then the associated **metric curvature** \( R \) is defined by
\[ R(x, y, z) = \nabla_{x,y} \nabla_{y,z} \nabla_{z,x} : T_x \to T_x. \] (11.2.8)

The metric curvature \( R(x, y, z) \) can be thought of as a discrete analog of the holonomy of the Levi-Civita connection on a manifold, where a tangent vector is parallel transported along a loop starting and ending at \( x \). On a manifold, the curvature at \( x \) is immediately obtained from the holonomy by considering the loops in a small neighborhood of \( x \). With this in mind, Definition 11.2.11 indeed generalizes the usual notion of curvature to causal fermion systems.

The following construction relates directional sign operators to vectors of the tangent space. Suppose that \( y \) is spin connectable to \( x \). By synchronizing the directional sign operator \( v_{xy} \), we obtain the vector
\[ \hat{y}_x := U_{xy}^{-1} v_{xy} U_{xy} \in K_x(y). \] (11.2.9)
As \( K_x(y) \in T^{s_x} \) is a representative of the tangent space, we can regard \( \hat{y}_x \) as a tangent vector. We thus obtain a mapping
\[ I(x) \to T_x : y \mapsto \hat{y}_x. \]
We refer to \( \hat{y}_x \) as the **directional tangent vector** of \( y \) in \( T_x \). As \( v_{xy} \) is a sign operator and the transformations in (11.2.9) are unitary, the directional tangent vector is a timelike unit vector with the additional property that the inner product \( \langle \cdot | \cdot \rangle_x \) is positive definite.

We finally explain how to reduce the dimension of the tangent space to four, with the desired Lorentzian signature \((1, 3)\).

**Definition 11.2.12.** The fermion system is called **chirally symmetric** if to every \( x \in M \) we can associate a spacelike vector \( u(x) \in T_x \) which is orthogonal to all directional tangent vectors,
\[ \langle u(x), \hat{y}_x \rangle = 0 \quad \text{for all } y \in I(x), \]
and is parallel with respect to the metric connection, i.e.
\[ u(x) = \nabla_{x,y} u(y) \nabla_{y,x} \quad \text{for all } y \in I(x). \]

**Definition 11.2.13.** For a chirally symmetric fermion system, we introduce the **reduced tangent space** \( T_x^{\text{red}} \) by
\[ T_x^{\text{red}} = \langle u_x \rangle^\perp \subset T_x. \]
Clearly, the reduced tangent space has dimension four and signature \((1, 3)\). Moreover, the operator \( \nabla_{x,y} \) maps the reduced tangent spaces isometrically to each other. The local operator \( \Gamma := -iu/\sqrt{-u^2} \) takes the role of the **pseudo-scalar matrix**.
11.3. Correspondence to Lorentzian Spin Geometry

We also explain how these notions correspond to the usual objects of differential geometry in Minkowski space (Theorem 11.3.1) and on a globally hyperbolic Lorentzian manifold (Theorem 11.3.2). We closely follow the presentation in [57, Section 3.3]; see also the review [50].

We let \((M, g)\) be a time-oriented Lorentzian spin manifold with spinor bundle \(SM\) (for basic definitions see Section 4.5). In order to obtain a corresponding causal fermion system, we adapt the construction in Minkowski space given in Section 5.4: First, we choose a closed subspace \(H\) of the Hilbert space of Dirac solutions \(\mathcal{H}_{m,\langle .|\rangle}\) (as introduced in Section 4.5). Endowed with the induced scalar product \(\langle .|\rangle_{\mathcal{H}} := \langle .|\rangle|_{\mathcal{H}\times\mathcal{H}}\), we obtain a Hilbert space \((H, \langle .|\rangle_{\mathcal{H}})\). Next, one introduces a regularization operator (5.4.1), for example by mollifying the initial data on a Cauchy surface (as is explained in [76, Section 4]). Introducing the local correlation operator \(F_{\varepsilon}(x)\) for every \(x \in M\) again by (5.4.4), we define the measure \(\rho\) on \(\mathcal{F}\) as the push-forward of the volume measure \(\mu\) on \(M\), i.e.

\[
\rho = (F_{\varepsilon})_{\ast} \mu,
\]

where, in local coordinates, the measure \(\mu\) has the form

\[
d\mu = \sqrt{|\det g|} \, d^4x.
\]

We thus obtain a causal fermion system \((\mathcal{F}, \mathcal{H}, \rho)\) describing the curved spacetime \((M, g)\).

Let \(\gamma(t)\) be a smooth, future-directed and timelike curve, for simplicity parametrized by the arc length, defined on the interval \([0, T]\) with \(\gamma(0) = y\) and \(\gamma(T) = x\). Then the parallel transport of tangent vectors along \(\gamma\) with respect to the Levi-Civita connection \(\nabla_{LC}\) gives rise to the isometry

\[
\nabla_{LC}^{x,y} : T_y \to T_x.
\]

In order to compare with the metric connection \(\nabla\) of Definition 11.2.10, we subdivide \(\gamma\) (for simplicity with equal spacing, although a non-uniform spacing would work just as well). Thus for any given \(N\), we define the points \(x_0, \ldots, x_N\) by

\[
x_n = \gamma(t_n) \quad \text{with} \quad t_n = \frac{nT}{N}.
\]

We define the parallel transport \(\nabla_{x,y}^N\) by successively composing the parallel transport between neighboring points,

\[
\nabla_{x,y}^N := \nabla_{x_N, x_{N-1}} \nabla_{x_{N-1}, x_{N-2}} \cdots \nabla_{x_1, x_0} : T_y \to T_x.
\]

We first state a result in the Minkowski vacuum. We choose \(\mathcal{H}\) as the subspace of all negative-energy solutions of the Dirac equation (describing the Dirac sea; see the preliminaries in Section 1.5). For technical simplicity, we choose the \(i\varepsilon\)-regularization, where the regularization operator (5.4.1) is the multiplication operator by \(e^{i\varepsilon\omega}\) in momentum space in (5.4.2).

**Theorem 11.3.1.** For every \(\varepsilon > 0\) we consider the causal fermion systems \((\mathcal{F}, \mathcal{H}, \rho)\) describing the vacuum with \(i\varepsilon\)-regularization. Then for a generic curve \(\gamma\) and for every \(N \in \mathbb{N}\), there is \(\varepsilon_0\) such that for all \(\varepsilon \in (0, \varepsilon_0]\) and all \(n = 1, \ldots, N\), the points \(x_n\) and \(x_{n-1}\) are spin connectable. Moreover,

\[
\nabla_{x,y}^{LC} = \lim_{N \to \infty} \lim_{\varepsilon \to 0^+} \nabla_{x,y}^N.
\]
By a generic curve we mean that the admissible curves are dense in the $C^\infty$-topology (i.e., for any smooth $\gamma$ and every $K \in \mathbb{N}$, there is a sequence $\gamma_\ell$ of admissible curves such that $D^k \gamma_\ell \to D^k \gamma$ uniformly for all $k = 0, \ldots, K$). The restriction to generic curves is needed in order to ensure that the Euclidean and directional sign operators are generically separated (see Definition 11.2.8 (b)). The proof of the above theorem is given in [55, Section 4].

Clearly, in this theorem the connection $\nabla^\text{LC}_{x,y}$ is trivial. In order to show that our connection also coincides with the Levi-Civita connection in the case with curvature, in [55, Section 5] a globally hyperbolic Lorentzian manifold is considered. For technical simplicity, we assume that the manifold is flat Minkowski space in the past of a given Cauchy hypersurface.

**Theorem 11.3.2.** Let $(\mathcal{M}, g)$ be a globally hyperbolic manifold which is isometric to Minkowski space in the past of a given Cauchy-hypersurface $\mathcal{N}$. For given $\gamma$, for any $\varepsilon > 0$ we consider the causal fermion system $(\mathcal{H}, \mathcal{F}, \rho)$ which in the past of $\mathcal{N}$ coincides with the causal fermion system in Minkowski space considered in Theorem 11.3.1. Then for a generic curve $\gamma$ and for every sufficiently large $N$, there is $\varepsilon_0$ such that for all $\varepsilon \in (0, \varepsilon_0]$ and all $n = 1, \ldots, N$, the points $x_n$ and $x_{n-1}$ are spin connectable. Moreover,

$$\lim_{N \to \infty} \lim_{\varepsilon \to 0} \nabla^N_{x,y} - \nabla^\text{LC}_{x,y} = O \left( \frac{L(\gamma)}{m^2} \right) \left( 1 + O\left( \frac{\text{scal}}{m^2} \right) \right),$$

where $R$ denotes the Riemann curvature tensor, $\text{scal}$ is scalar curvature, and $L(\gamma)$ is the length of the curve $\gamma$.

Thus the metric connection of Definition 11.2.10 indeed coincides with the Levi-Civita connection, up to higher order curvature corrections. For detailed explanations and the proof we refer to [55, Section 5].

We conclude this section with a few remarks on further constructions [55]. First, there is the subtle point that the unitary transformation $U \in \exp(iR s_x)$ which is used to identify two representatives $K, \tilde{K} \in T_x$ via the relation $\tilde{K} = UKU^{-1}$ (see Definition 11.2.4) is not unique. More precisely, the operator $U$ can be transformed according to $U \to -U$ and $U \to s_x U$.

As a consequence, the metric connection (see Definition 11.2.10) is defined only up to the transformation

$$\nabla_{x,y} u \to s_x (\nabla_{x,y} u) s_x.$$

Note that this transformation maps representatives of the same tangent vector into each other, so that $\nabla_{x,y} u \in T_x$ is still a well-defined tangent vector. But we get an ambiguity when composing the metric connection several times (as for example in the expression for the metric curvature in Definition 11.2.11). This ambiguity can be removed by considering parity-preserving systems as introduced in [55, Section 3.4].

At first sight, one might conjecture that Theorem 11.3.2 should also apply to the spin connection in the sense that

$$D^\text{LC}_{x,y} = \lim_{N \to \infty} \lim_{\varepsilon \to 0} D^N_{x,y},$$

where $D^\text{LC}$ is the spin connection on $S\mathcal{M}$ induced by the Levi-Civita connection and

$$D^N_{x,y} := D_{x_N,x_{N-1}} D_{x_{N-1},x_{N-2}} \cdots D_{x_1,x_0} : S_y \to S_x$$

(11.3.1)
(and $D$ is the spin connection of Theorem 11.2.9). It turns out that this conjecture is false. But the conjecture becomes true if we replace (11.3.1) by the operator product

$$D^N_{(x,y)} := D_{x_N,x_{N-1}} U^{(x_N|x_{N-2})}_{x_N,x_{N-1}} D_{x_{N-1},x_{N-2}} U^{(x_{N-1}|x_{N-3})}_{x_{N-1},x_{N-2}} \cdots U^{(x_2|x_0)}_{x_2,x_0} D_{x_1,x_0}.$$ 

Here the intermediate factors $U^{(1)}$ are the so-called splice maps given by

$$U^{(z|y)} = U_{xz} V U_{xy},$$

where $U_{xz}$ and $U_{xy}$ are synchronization maps, and $V \in \exp(iR_{s_x})$ is an operator which identifies the representatives $K_{xy}, K_{xz} \in T_x$ (for details see [55], Section 3.7 and Section 5). The splice maps also enter the spin curvature $\mathcal{R}$, which is defined in analogy to the metric curvature (11.2.8) by

$$\mathcal{R}(x,y,z) = U^{(z|y)}_x D_{x,y} U^{(x|z)}_y D_{y,z} U^{(y|x)}_z D_{z,x} : S_x \to S_x.$$ 

Exercises
Part 3

Mathematical Methods and Analytic Constructions
CHAPTER 12

Measure-Theoretic Methods

The main goal of this chapter is to prove the existence of minimizers for the causal action principle in the case that $\mathcal{H}$ is finite dimensional and $\rho$ is normalized, i.e.

$$\dim \mathcal{H} =: f < \infty \quad \text{and} \quad \rho(\mathcal{F}) = 1.$$  \hspace{1cm} (12.0.1)

After introducing the necessary methods (Sections 12.1 and 12.2) we first apply them in order to prove existence of minimizers for causal variational principles in the compact setting (Section 12.3). In preparation for the proof for the causal action principle, we illustrate the constraints by a few examples (Section 12.4). The difficulties revealed by these examples can be resolved by working with the so-called moment measures. After introducing the needed mathematical methods (Section 12.5), the moment measures are introduced (Section 12.6). Then the existence proof is completed (Section 12.7). In order to give a first idea for how to deal with an infinite total volume, we finally prove existence of minimizers for causal variational principles in the non-compact setting (Section 12.8).

Our general strategy is to apply the direct method of the calculus of variations, which can be summarized as follows:

(a) Choose a minimizing sequence, i.e. a sequence of measures $(\rho_k)$ which satisfy the constraint such that

$$S(\rho_k) \rightarrow \inf_{\rho} S(\rho).$$

Such a minimizing sequence always exists by definition of the infimum (note that the action and therefore also its infimum are non-negative).

(b) Show that a subsequence of the measures converges in a suitable sense,

$$\rho_{k_l} \xrightarrow{\text{in some sense}} \rho.$$  

Here the quotation marks indicate that we still need to specify in which sense the sequence should converge (convergence in which space, strong or weak convergence, etc.).

(c) Finally, one must show that the action is lower semi-continuous, i.e.

$$S(\rho) \leq \liminf_{l \to \infty} S(\rho_{k_l}).$$

Also, one must prove that the limit measure $\rho$ satisfies the constraints.

Once these three steps have been carried out, the measure $\rho$ is a desired minimizer. We point out that this procedure does not give a unique minimizer, simply because there may be different minimizing sequences, and because the choice of the subsequences may involve an arbitrariness. Indeed, for the causal action principle we do not expect uniqueness. There should be many different minimizers, which describe different physical systems (like the vacuum, a system involving particles and fields, etc.). This intuitive picture is confirmed by the numerical study in \cite{79} which shows that, even if the dimension of $\mathcal{H}$ is small (equal to two), there are indeed many different minimizers.
12.1. The Banach-Alaoglu Theorem

For our purposes, it suffices to consider the case that the Banach space is separable, in which case the theorem was first proved by Banach (Alaoglu proved the generalization to non-separable Banach spaces; this makes use of Tychonoff’s theorem and goes beyond what we need here). Indeed, the idea of proof of the theorem can be traced back to Eduard Helly’s doctoral thesis in 1912, where the closely related “Helly’s selection theorem” is proved (of course without reference to Banach spaces, which were introduced later). We closely follow the presentation in [110], Section 10.3.

Let \((E, \|\cdot\|_E)\) be a separable (real or complex) Banach space and \((E^*, \|\cdot\|_{E^*})\) its dual space with the usual sup-norm, i.e.

\[
\|\phi\|_{E^*} = \sup_{u \in E, \|u\| = 1} |\phi(u)|.
\]

(12.1.1)

A sequence \((\phi_n)_{n \in \mathbb{N}}\) in \(E^*\) is said to be weak*-convergent to \(\phi \in E^*\) if

\[
\lim_{n \to \infty} \phi_n(u) = \phi(u) \quad \text{for all } u \in E.
\]

Theorem 12.1.1. (Banach-Alaoglu in the separable case) Let \(E\) be a separable Banach space. Then every bounded sequence in \(E^*\) has a weak*-convergent subsequence.

Proof. Let \(\phi_n\) be a bounded sequence in \(E^*\), meaning that there is a constant \(c > 0\) with

\[
\|\phi_n\|_{E^*} \leq c \quad \text{for all } n \in \mathbb{N}.
\]

(12.1.2)

We let \((u_\ell)_{\ell \in \mathbb{N}}\) be a sequence in \(E\) which is dense in \(E\). Combining (12.1.2) with (12.1.1), the estimate

\[
|\phi_n(u_1)| \leq \|\phi_n\|_{E^*} \|u_1\|_E \leq c \|u_1\|_E
\]

(12.1.3)

shows that \((\phi_n(u_1))_{n \in \mathbb{N}}\) is a bounded sequence. Thus we can choose a convergent subsequence. By inductively choosing subsequences and taking the diagonal sequence, we obtain a subsequence \((\phi_{n_j})\) such that the limit \(\lim_{j \to \infty} \phi_{n_j}(u_\ell)\) exists for all \(\ell \in \mathbb{N}\). Hence setting

\[
\phi(u_\ell) := \lim_{j \to \infty} \phi_{n_j}(u_\ell),
\]

we obtain a densely defined functional. Taking the limit in (12.1.3) (and the similar inequalities for \(u_2, u_3, \ldots\)), one sees that this functional is again continuous. Therefore, it has a unique continuous extension to \(E\). By continuity, the resulting functional \(\phi \in E^*\) satisfies the relations

\[
\phi(u) = \lim_{j \to \infty} \phi_{n_j}(u) \quad \text{for all } u \in \mathcal{H}.
\]

This concludes the proof. □

12.2. The Riesz Representation Theorem

In this section and Section 12.5, we shall introduce the methods from measure theory needed for the existence proofs. Apart from the books already mentioned in the preliminaries (Section 2.3), we recommend [32] is also a good reference. However, this book is only concerned with measures in \(\mathbb{R}^n\), but otherwise goes far beyond what we need here.

For our purposes, it suffices to restrict attention to the case that the base space \(\mathcal{K}\) is a compact topological space. We always consider bounded regular Borel measures on \(\mathcal{K}\) (for the preliminaries see Section 2.3). In order to avoid confusion, we note that by a measure we always mean a positive measure (signed measures will not at all be considered in this
12.2. THE RIESZ REPRESENTATION THEOREM

A bounded measure, also referred to as a measure of finite total volume, often, we normalize the measure such that \( \mu(K) = 1 \).

In words, the Riesz representation theorem makes it possible to represent a linear functionals on the Banach space of continuous functions of a topological space by a regular Borel measures on this topological space. We remark that we already came across the Riesz representation theorem in the functional analysis course, where it was used for the construction of spectral measures. However, back then we only proved it in the special case that the topological space was an interval of the real line. We now state the general theorem and outline its proof, mainly following the presentation in [97, \$56].

As a simple example, one can choose \( K \) as the closed unit ball in \( \mathbb{R}^n \). Restricting the Lebesgue measure to the Borel subsets of \( K \) gives a Radon measure. The Lebesgue measure itself is a completion of this Radon measure obtained by extending the \( \sigma \)-algebra of measurable sets by all subsets of Borel sets of measure zero. Since this completion is a rather trivial extension, in what follows we prefer to work with Radon measures or, equivalently, with normalized regular Borel measures.

**Theorem 12.2.1. (Riesz representation theorem)** Let \( K \) be a compact topological space, and \( E = C^0(K, \mathbb{R}) \) the Banach space of continuous functions on \( K \) with the usual sup-norm,

\[
\|f\| = \sup_{x \in K} |f(x)|.
\]

Let \( \Lambda \in E^* \) be a continuous linear functions which is positive in the sense that

\[
\Lambda(f) \geq 0 \quad \text{for all nonnegative function } f \in C^0(K, \mathbb{R}).
\]

Then there is a unique regular Borel measure \( \mu \) such that

\[
\Lambda(f) = \int_K f \, d\mu \quad \text{for all } f \in C^0(K, \mathbb{R}).
\]

Outline of the Proof. We follow the strategy in [97, \$56]. Given a Borel set \( A \subset K \), we set

\[
\lambda(A) = \inf \{ \Lambda(f) \mid f \in C^0(K, \mathbb{R}) \text{ and } f \geq \chi_A \} \in \mathbb{R}_0^+.
\]

Intuitively speaking, \( \lambda \) gives us the desired “volume” of the set \( A \). But there is the technical problem that \( \lambda \) is in general not a regular Borel measure. Instead, it merely is a content, meaning that it has the following properties:

(i) non-negative and finite: \( 0 \leq \lambda(A) < \infty \)
(ii) monotone: \( C, D \text{ compact and } C \subset D \implies \lambda(C) \leq \lambda(D) \)
(iii) additive: \( C, D \text{ compact and disjoint } \implies \lambda(C \cup D) = \lambda(C) + \lambda(D) \)
(iv) subadditive: \( C, D \text{ compact } \implies \lambda(C \cup D) \leq \lambda(C) + \lambda(D) \)

At this stage, we are in a similar situation as in the elementary measure theory course after saying that a cube of length \( \ell \) in \( \mathbb{R}^3 \) should have volume \( \ell^3 \). In order to get from this “volume measure” to a measure in the mathematical sense, one has to proceed in several steps invoking the subtle and clever constructions of measure theory (due to Lebesgue, Hahn, Carathéodory and others) in order to get a mapping from a \( \sigma \)-algebra to the non-negative real numbers which is \( \sigma \)-additive. In simple terms, repeating these constructions starting from the above content gives the desired Borel measure \( \mu \). For brevity, we here merely outline the constructions and refer for details to text books on measure theory (like for example [97, Chapter X]).
The first step is to approximate (or exhaust) from inside by compact sets. Thus one introduces the inner content $\lambda_*$ by
\[
\lambda_*(U) = \sup \{ \lambda(C) \mid C \subset U \text{ compact} \}.
\]
This inner content is countably subadditive and countably additive. The second step is to exhaust from outside by open sets. This gives the outer measure $\mu^*$,
\[
\mu^*(U) = \inf \{ \lambda_*(\Omega) \mid \Omega \supset U \text{ open} \}.
\]
The outer measure is defined for any subset of $\mathcal{K}$. Therefore, it remains to distinguish the measurable sets. This is accomplished by Carathéodory’s criterion, which defines a set $A \subset \mathcal{K}$ to be measurable if
\[
\mu^*(A) = \mu^*(A \cap B) + \mu^*(A \setminus B)
\]
for every subset $B \subset \mathcal{K}$. Then Carathéodory’s lemma (for a concise proof see for example [18, Lemma 2.8]) implies that the measurable sets form a $\sigma$-algebra, and that the restriction of $\mu^*$ to the measurable sets is indeed a measure, denoted by $\mu$.

In order to complete the proof, one still needs to verify that every Borel set is $\mu$-measurable. Moreover, it remains to show that the resulting Borel measure is regular. To this end, one first needs to show that the content $\lambda$ is regular in the following sense:

(v) regular: For every compact $C$,
\[
\lambda(C) = \inf \{ \lambda(D) \mid D \text{ compact and } C \subset \overset{\circ}{D} \}.
\]
As the proofs of these remaining points are rather straightforward and not very instructive, we refer for the details to [97, §54–§56].

12.3. Existence of Minimizers for Causal Variational Principles in the Compact Setting

We now apply the above methods in order to prove existence for causal variational principles in the compact setting. Our strategy is to apply the Banach-Alaoglu theorem to a specific Banach space, namely the continuous functions on a compact metric space. We first verify that this Banach space is indeed separable.

**Proposition 12.3.1.** Let $\mathcal{K}$ be a compact metric space. Then $C^0(\mathcal{K}, \mathbb{R})$ is a separable Banach space.

**Proof.** The proposition is a consequence of the Stone-Weierstrass theorem, whose proof can be found for example in [26, 7.3.1] We closely follow the proof given in [26, 7.4.4].

Covering $\mathcal{K}$ by a finite number of open balls of radii 1, 1/2, 1/3, ..., one gets an enumerable basis of the open sets $(U_n)_{n \in \mathbb{N}}$. For any $n \in \mathbb{N}$, we let $g_n$ be the continuous function
\[
g_n(x) = d(x, \mathcal{K} \setminus U_n).
\]
Clearly, the algebra generated by these functions (by taking finite products and finite linear combinations) is again separable. Therefore, it suffices to show that this algebra is dense in $C^0(\mathcal{K}, \mathbb{R})$. To this end, we need to verify the assumptions of the Stone-Weierstrass theorem. The only assumption which is not obvious is that the algebra separates the points. This can be seen as follows: Let $x$ and $y$ be two distinct points in $\mathcal{K}$. Since the $(U_n)$ are a basis of the topology, there is $U_n$ with $x \in U_n$ and $y \notin U_n$. As a consequence, $g_n(x) > 0$ but $g_n(y) = 0$. □
12.3. EXISTENCE OF MINIMIZERS IN THE COMPACT SETTING

We proceed by proving a compactness result for Radon measures.

**Theorem 12.3.2.** Let \( \rho_n \) be a series of regular Borel measures on \( C^0(\mathcal{K}, \mathbb{R}) \) which are bounded in the sense that there is a constant \( c > 0 \) with

\[
\rho_n(\mathcal{K}) \leq c \quad \text{for all } n.
\]

Then there is a subsequence \((\rho_{n_k})\) which converges as a measure, i.e.

\[
\lim_{k \to \infty} \int_{\mathcal{K}} f \, d\rho_{n_k} = \int_{\mathcal{K}} f \, d\rho \quad \text{for all } f \in C^0(\mathcal{K}, \mathbb{R}).
\]

(12.3.1)

Moreover, the total volume converges, i.e.

\[
\rho(\mathcal{K}) = \lim_{k \to \infty} \rho_{n_k}(\mathcal{K}).
\]

(12.3.2)

**Proof.** Via

\[
\phi_n(f) := \int_{\mathcal{K}} f \, d\rho_n,
\]

every measure can be identified with a positive linear functional on \( E := C^0(\mathcal{K}, \mathbb{R}) \). Since \( E \) is separable (Proposition 12.3.1), we can apply the Banach-Alaoglu theorem in the separable case (Theorem 12.1.1) to conclude that there is a weak*-convergent subsequence, i.e.

\[
\lim_{k \to \infty} \phi_{n_k}(f) = \phi(f) \quad \text{for all } f \in C^0(\mathcal{K}, \mathbb{R}).
\]

(12.3.3)

Clearly, since all \( \phi_{n_k} \) are positive, the same is true for the limit \( \phi \). Therefore, the Riesz representation theorem (Theorem 12.2.1) makes it possible to represent \( \phi \) by a regular Borel measure \( \rho \), i.e.

\[
\phi(f) = \int_{\mathcal{K}} f \, d\rho \quad \text{for all } f \in C^0(\mathcal{K}, \mathbb{R}).
\]

(12.3.4)

Choosing \( f \) as the constant function, one obtains (12.3.2). This concludes the proof. \( \Box \)

**Theorem 12.3.3.** Assume that \( \mathcal{F} \) is a compact topological space and the Lagrangian is continuous,

\[
\mathcal{L} \in C^0(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+_0).
\]

Then the causal variational principle where the causal action (6.2.2) is minimized in the class of regular Borel measures under the volume constraint (6.2.3) is well-posed in the sense that every minimizing sequence \((\rho_n)_{n \in \mathbb{N}}\) has a subsequence which converges as a measure to a minimizer \( \rho \).

**Proof.** The existence of a convergent subsequence \((\rho_{n_k})_{k \in \mathbb{N}}\) is proven in Theorem 12.3.2. It remains to show that the action is continuous, i.e.

\[
\lim_{k \to \infty} \mathcal{S}(\rho_{n_k}) = \mathcal{S}(\rho).
\]

This is verified in detail as follows. Using that the Lagrangian is continuous in its second argument, we know that

\[
\lim_{k \to \infty} \int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho_{n_k}(y) = \int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho(y) \quad \text{for all } x \in \mathcal{F}.
\]

(12.3.5)

Next, since \( \mathcal{F} \) is compact, the Lagrangian is even uniformly continuous on \( \mathcal{F} \times \mathcal{F} \). Therefore, given \( \varepsilon > 0 \), every point \( x \in \mathcal{F} \) has an open neighborhood \( U(x) \subset \mathcal{F} \) such that

\[
|\mathcal{L}(\hat{x}, y) - \mathcal{L}(x, y)| < \varepsilon \quad \text{for all } \hat{x} \in U(x) \text{ and } y \in \mathcal{F}.
\]
Integrating over $y$ with respect to any normed regular Borel measure $\tilde{\rho}$, it follows that

$$\left| \int_{\mathcal{F}} \mathcal{L}(\hat{x}, y) \, d\hat{\rho}(y) - \int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho(y) \right| \leq \varepsilon \quad \text{for all } \hat{x} \in U(x).$$

(12.3.6)

Covering $\mathcal{F}$ by a finite number of such neighborhoods $U(x_1), \ldots, U(x_N)$, one can combine the pointwise convergence (12.3.5) for $x = x_1, \ldots, x_N$ with the estimate (12.3.6) to conclude that for any $\varepsilon > 0$ there is $k_0 \in \mathbb{N}$ such that

$$\left| \int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho_{n_k}(y) - \int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho(y) \right| \leq 3\varepsilon \quad \text{for all } x \in \mathcal{F} \text{ and } k \geq k_0.$$

Integrating over $y$ with respect to $\rho_{n_k}$ and $\rho$ gives for all $k \geq k_0$ the respective inequalities

$$\left| S(\rho_{n_k}) - \int_{\mathcal{F}} d\rho_{n_k}(x) \int_{\mathcal{F}} d\rho(y) \mathcal{L}(x, y) \right| \leq 3\varepsilon,$n

$$\int_{\mathcal{F}} d\rho(x) \int_{\mathcal{F}} d\rho_{n_k}(y) \mathcal{L}(x, y) - S(\rho) \right| \leq 3\varepsilon.$$

Combining these inequalities and using that the Lagrangian is symmetric in its two arguments, we conclude that

$$\left| S(\rho_{n_k}) - S(\rho) \right| \leq 6\varepsilon.$$

This gives the result.

We finally remark that the statement of this theorem also holds if the Lagrangian merely is lower semi-continuous, as is worked out in [72, Section 3.2].

12.4. Examples Illustrating the Constraints

Compared to causal variational principles in the compact setting, the existence proof for the causal action principle is considerably harder because we need to handle the constraints (5.5.3)–(5.5.5) and must face the difficulty that the set $\mathcal{F}$ is unbounded and therefore non-compact. We now explain the constraints in a few examples. The necessity of the volume constraint is quite obvious: If we dropped the constraint of fixed total volume (5.5.3), the measure $\rho = 0$ would be a trivial minimizer. The role of the trace constraint is already less obvious. It is explained in the next two examples.

**Example 12.4.1. (necessity of the trace constraint)** Let $x$ be the operator with the matrix representation

$$x = \text{diag}(1, \ldots, 1, -1, \ldots, -1, 0, 0, \ldots)$$

(12.4.1)

Moreover, we choose $\rho$ as a multiple of the Dirac measure supported at $x$. Then $T > 0$ but $S = 0$.

**Example 12.4.2. (non-triviality of the action with trace constraint)** Let $\rho$ be a normalized measure which satisfies the trace constraint in a non-trivial way, i.e.

$$\int_{\mathcal{F}} \text{tr}(x) \, d\rho(x) = \text{const} \neq 0.$$

Let us prove that the action is non-zero. This will show that the trace constraint really avoids trivial minimizers of the causal action principle.
(a) Since the integral over the trace is non-zero, there is a point $x$ in the support of $\rho$ with $\text{tr}(x) \neq 0$. We denote the non-trivial eigenvalues of $x$ by $\nu_1, \ldots, \nu_{2n}$, and order them according to

$$\nu_1 \leq \cdots \leq \nu_n \leq 0 \leq \nu_{n+1} \leq \cdots \leq \nu_{2n}.$$  

(12.4.2)

The fact that the trace of $x$ is non-zero clearly implies that the $\nu_i$ do not all have the same absolute value. As a consequence, the nontrivial eigenvalues of the operator product $x^2$ given by $\lambda^{xx}_j = \nu_j^2$ are all non-negative and not all equal. Using the form of the Lagrangian in [5.5.1], we conclude that $\mathcal{L}(x, x) > 0$.

(b) Since the Lagrangian is continuous in both arguments, there is an open neighborhood $U \subset \mathcal{F}$ of $x$ such that $\mathcal{L}(y, z) > 0$ for all $y, z \in U$. Since $x$ lies in the support of $\rho$, we know that $\rho(U) > 0$. As a consequence,

$$S \geq \int_U d\rho(x) \int_U d\rho(y) \mathcal{L}(x, y) > 0$$

(because if the integrals vanished, then the integrand would have to be zero almost everywhere, a contradiction).

We now come to the boundedness constraint. In order to explain how it comes about, we give an explicit example with $(4 \times 4)$-matrices (for a similar example with $(2 \times 2)$-matrices see Exercise 6.3).

**Example 12.4.3. (necessity of the boundedness constraint)** The following example explains why the boundedness constraint (5.5.5) is needed to ensure the existence of minimizers. It was first given in [43, Example 2.9]. Let $\mathcal{H} = \mathbb{C}^4$. We introduce the four $4 \times 4$-matrices acting on $\mathcal{H}$ by

$$\gamma^\alpha = \begin{pmatrix} \sigma^\alpha & 0 \\ 0 & -\sigma^\alpha \end{pmatrix}, \quad \alpha = 1, 2, 3 \quad \text{and} \quad \gamma^4 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}$$

(where the $\sigma^\alpha$ are again the Pauli matrices [1.3.1]). Given a parameter $\tau > 1$, we consider the following mapping from the sphere $S^3 \subset \mathbb{R}^4$ to the linear operators on $\mathcal{H}$,

$$F : S^3 \to L(\mathcal{H}) \; \text{,} \quad F(x) = \sum_{i=1}^{4} \tau x^i \gamma^i + \mathbb{1}.$$ 

(a) The matrices $F(x)$ have two positive and two negative eigenvalues:

Since the computation of the eigenvalues of $4 \times 4$-matrices is tedious, it is preferable to proceed as follows. The matrices $\gamma^j$ are the Dirac matrices of Euclidean $\mathbb{R}^4$, satisfying the anti-commutation relations

$$\{\gamma^i, \gamma^j\} = 2\delta^{ij} \mathbb{1} \quad (i, j = 1, \ldots, 4).$$
As a consequence,
\[ F(x) - 1 = \sum_{i=1}^{4} \tau x^i \gamma^i \]
\[ (F(x) - 1)^2 = \sum_{i,j=1}^{4} \tau^2 x^i x^j \gamma^i \gamma^j = \frac{\tau^2}{2} \sum_{i,j=1}^{4} x^i x^j \{ \gamma^i, \gamma^j \} \]
\[ = \frac{\tau^2}{2} \sum_{i,j=1}^{4} x^i x^j 2 \delta_{ij} 1 = \tau^2 1. \]

Hence the matrix \( F(x) \) satisfies the polynomial equation
\[ (F(x) - 1)^2 = \tau^2 1. \]

We conclude that \( F(x) \) has the eigenvalues
\[ \nu_{\pm} = 1 \pm \tau. \]
Since \( F(x) - 1 \) is trace-free, each eigenvalue must appear with multiplicity two. Using that \( \tau > 1 \), we conclude that \( F(x) \) really has two positive and two negative eigenvalues.

(b) Construction of a causal fermion system:
Let \( \mu \) be the normalized Lebesgue measure on \( S^3 \subset \mathbb{R}^4 \). Setting \( \rho = F_* \mu \) defines a causal fermion system of spin dimension two and total volume one. Since the matrices \( F(x) \) all have trace four, we also know that
\[ \int \text{tr}(x) \ d\rho(x) = \int \text{tr}(F(x)) \ d\mu(x) = 4. \]

Therefore, the volume constraint (5.5.3) and the trace constraint (5.5.4) are satisfied, both with constants independent of \( \tau \).

(c) Computation of the eigenvalues of \( F(x) F(y) \):
Again, this can be computed most conveniently using the Clifford relations.
\[ F(x) F(y) = \left( \sum_{i=1}^{4} \tau x^i \gamma^i + 1 \right) \left( \sum_{j=1}^{4} \tau y^j \gamma^j + 1 \right) \]
\[ = (1 + \tau^2 \langle x, y \rangle) + \tau \sum_{i=1}^{4} (x^i + y^i) \gamma^i + \frac{\tau^2}{2} \sum_{i,j=1}^{4} x^i y^j \{ \gamma^i, \gamma^j \}. \quad (12.4.3) \]
Using that
\[ \gamma^i [\gamma^i, \gamma^j] = -[\gamma^i, \gamma^j] \gamma^i, \]
we conclude that
\[ (F(x) F(y) - (1 + \tau^2 \langle x, y \rangle) 1)^2 = \tau^2 \sum_{i=1}^{4} (x^i + y^i)^2 + \left( \frac{\tau^2}{2} \sum_{i,j=1}^{4} x^i y^j [\gamma^i, \gamma^j] \right)^2. \]
This can be simplified with the help of the relations

\[ \sum_{i=1}^{5} (x^i + y^i)^2 = 2 + 2 \langle x, y \rangle \]  
\[ (\sum_{i,j=1}^{4} x^i y^j [\gamma^i, \gamma^j])^2 = -4 \sin^2 \vartheta = -4 \left(1 - \langle x, y \rangle^2\right), \]  
where \( \vartheta \) is the angle between the vectors \( x, y \in \mathbb{R}^4 \). The relation (12.4.5) can be verified in detail as follows. The rotational symmetry of the Euclidean Dirac operator on \( \mathbb{R}^5 \) means that for every rotation \( O \in SO(4) \) there is a unitary operator \( U \in SU(4) \) such that

\[ O^j \gamma^j = U \gamma^i U^{-1}. \]

Making use of this rotational symmetry, we can arrange that the vector \( x \) is the basis vector \( e_1 \) and that \( y = \cos \vartheta e_1 + \sin \vartheta e_2 \). As a consequence,

\[ \sum_{i,j=1}^{5} x^i y^j [\gamma^i, \gamma^j] = \sin \vartheta [\gamma^1, \gamma^2] = 2 \sin \vartheta \gamma^1 \gamma^2 \]

and applying the anti-commutation relations gives (12.4.5). Combining the above equations, we conclude that the product \( F(x) F(y) \) satisfies the polynomial equation

\[ \left( F(x) F(y) - (1 + \tau^2 \langle x, y \rangle) 1 \right)^2 = 2 \tau^2 (1 + \langle x, y \rangle) - \tau^4 \left(1 - \langle x, y \rangle^2\right) \]

Taking the square root, the zeros of this polynomial are computed by

\[ \lambda_{1/2} = 1 + \tau^2 \langle x, y \rangle \pm \tau \sqrt{1 + \langle x, y \rangle} \sqrt{2 - \tau^2 (1 - \langle x, y \rangle)}. \]  

Moreover, taking the trace of (12.4.3), one finds

\[ \text{tr} \left( F(x) F(y) \right) = 4 (1 + \tau^2 \langle x, y \rangle). \]

This implies that each eigenvalue in (12.4.6) has algebraic multiplicity two.

(d) Computation of the Lagrangian:

Again denoting the angle between the vectors \( x, y \in \mathbb{R}^4 \) by \( \vartheta \). Then if \( \vartheta \) is sufficiently small, then the term \( (1 - \langle x, y \rangle) \) is close to zero, and thus the arguments of the square roots are all positive. However, if \( \vartheta \) becomes so large that

\[ \vartheta \geq \vartheta_{\text{max}} := \arccos \left(1 - \frac{2}{\tau^2}\right), \]

then the argument of the last square root in (12.4.6) becomes negative, so that the \( \lambda_{1/2} \) form a complex conjugate pair. Moreover, a short calculation shows that

\[ \lambda_1 \lambda_2 = (1 + \tau)^2 (1 - \tau)^2 > 0, \]
implying that if the $\lambda_{1/2}$ are both real, then they have the same sign. Using this information, the Lagrangian simplifies to

$$\mathcal{L}(F(x), F(y)) = \frac{1}{8} \sum_{i,j=1}^{4} \left( |\lambda_{i}^x| - |\lambda_{j}^y| \right)^2 = \frac{1}{2} \sum_{i,j=1}^{2} \left( |\lambda_i| - |\lambda_j| \right)^2$$

$$= \frac{1}{2} \Theta(\vartheta_{\text{max}} - \vartheta) \sum_{i,j=1}^{2} (\lambda_i - \lambda_j)^2 = \Theta(\vartheta_{\text{max}} - \vartheta) (\lambda_1 - \lambda_2)^2$$

$$= 4\tau^2 (1 + \cos \vartheta) (2 - \tau^2 (1 - \cos \vartheta)) \Theta(\vartheta_{\text{max}} - \vartheta).$$

(e) **Computation the action:**

Inserting this Lagrangian in (5.5.2) and using the definition of the push-forward measure, we obtain

$$S = \int_{S^3} d\mu(x) \int_{S^3} d\mu(y) \mathcal{L}(F(x), F(y))$$

$$= \int_{S^3} d\mu(y) \mathcal{L}(F(x), F(y)) = \frac{2}{\pi} \int_{0}^{\vartheta_{\text{max}}} \mathcal{L}(\cos \vartheta) \sin^2 \vartheta d\vartheta$$

$$= \frac{512}{15\pi} \frac{1}{\tau} + O(\tau^{-2}).$$

Thus setting $F_k = F|_{\tau=k}$, we have constructed a divergent minimizing sequence. However, the integral in the boundedness constraint (5.5.5) also diverges as $k \to \infty$. This example shows that leaving out the boundedness constraint, there is no minimizer. ◊

We finally remark that this example is not as artificial or academic as it might appear at first sight. Indeed, as observed by Niki Kilbertus in his master thesis [104], when discretizing a Dirac system in $\mathbb{R} \times S^3$ (where the sphere can be thought of as a spatial compactification of Minkowski space), then in the simplest case of four occupied Dirac states (referred to as “one shell,” i.e $\text{dim} \mathcal{H} = 4$), this system reduces precisely to the last example. In simple terms, this observation can be summarized by saying that Clifford structures tend to make the causal action small.

### 12.5. The Radon-Nikodym Theorem

As already mentioned at the beginning of the previous section, one difficulty in the existence proof for the causal action principle is the fact that the set $\mathcal{F}$ is unbounded and thus non-compact. In order to deal with this difficulty, we need one more mathematical tool: the Radon-Nikodym theorem. We now give the proof of the Radon-Nikodym theorem by von Neumann following the presentation in [126], Chapter 6. An alternative method of proof is given in [97, 32]. As in Section 12.2 it again suffices to consider the case that the base space $\mathcal{X}$ is a compact topological space.

**Definition 12.5.1.** A Radon measure $\lambda$ is absolutely continuous with respect to another Radon measure $\nu$, denoted by

$$\lambda \ll \nu,$$

if the implication

$$\nu(E) = 0 \implies \lambda(E) = 0$$
12.5. THE RADON-NIKODYM THEOREM

holds for any Borel set \( E \). The measure \( \lambda \) is **concentrated** on the Borel set \( A \) if \( \lambda(E) = \lambda(E \cap A) \) for all Borel sets \( E \). The measures \( \lambda \) and \( \mu \) are **mutually singular**, denoted by

\[ \lambda \perp \nu, \]

if there are disjoint Borel sets \( A \) and \( B \) such that \( \lambda \) is concentrated in \( A \) and \( \nu \) is concentrated in \( B \).

In order to avoid confusion, we point out that the supports of two mutually singular measures are not necessarily disjoint, as one sees in the simple example of the Lebesgue measure on \((0, 1)\) and the Dirac measure supported at the origin,

\[ \lambda := dx|_{(0, 1)} \quad \text{and} \quad \mu = \delta_0. \]

Since the support is by definition a closed set (see (2.3.3)), the support of \( dx|_{(0, 1)} \) contains the origin, which is precisely the support of the Dirac measure. But clearly, the two measures are concentrated on the sets \((0, 1)\) and \(\{0\}\), respectively, and are thus mutually singular.

**Theorem 12.5.2. (Radon-Nikodym)** Let \( \mu \) and \( \lambda \) be Radon measures on the compact topological space \( K \).

(a) There is a unique pair of Borel measures \( \lambda_a \) and \( \lambda_s \) such that

\[ \lambda = \lambda_a + \lambda_s \quad \text{and} \quad \lambda_a \ll \mu, \lambda_s \perp \mu. \]

(b) There is a unique function \( h \in L^1(K, d\mu) \) such that

\[ \lambda_a(E) = \int_E h \ d\mu \quad \text{for every Borel set } E. \quad (12.5.1) \]

The pair \((\lambda_a, \lambda_s)\) is also referred to as the **Lebesgue decomposition** of \( \lambda \) with respect to \( \mu \).

**Proof of Theorem 12.5.2.** The uniqueness of the decomposition is easily seen as follows: Suppose that \((\lambda'_a, \lambda'_s)\) is another Lebesgue decomposition. Then

\[ \lambda'_a - \lambda_a = \lambda_s - \lambda'_s. \quad (12.5.2) \]

Since \( \lambda_s \perp \mu \) and \( \lambda'_s \perp \mu \), the measures \( \lambda_s \) and \( \lambda'_s \) are concentrated in a Borel set \( A \) with \( \mu(A) = 0 \). Evaluating \((12.5.2)\) on Borel subsets of \( A \), the left side vanishes, because \( \lambda_a \) and \( \lambda'_a \) are both absolutely continuous with respect to \( \mu \). Hence \( \lambda'_a - \lambda_a = 0 \). Using this relation in \((12.5.2)\), we also conclude that \( \lambda'_s - \lambda_s = 0 \). This proves uniqueness.

For the existence proof, we let \( \rho \) be the measure \( \rho = \lambda + \mu \). Then

\[ \int_K f \ d\rho = \int_K f \ d\lambda + \int_K f \ d\mu \quad (12.5.3) \]

for any non-negative Borel function \( f \). If \( f \in L^2(K, d\rho) \), the Schwarz inequality gives

\[ \left| \int_K f \ d\lambda \right| \leq \int_K |f| \ d\rho \leq \sqrt{\mu(K)} \ ||f||_{L^2(K, d\rho)}. \]

Therefore, the mapping \( f \mapsto \int_K f \ d\lambda \) is a bounded linear functional on \( L^2(K, d\rho) \). By the Fréchet-Riesz theorem, we can represent this linear functional by a function \( g \in L^2(K, d\rho) \), i.e.

\[ \int_K f \ d\lambda = \int_K g f \ d\rho \quad \text{for all } f \in L^2(K, d\rho). \quad (12.5.4) \]
We next want to show that, by modifying $g$ on a set of $\rho$-measure zero, we can arrange that $g$ takes values in the interval $[0, 1]$. To this end, we let $E$ be any Borel set with $\rho(E) > 0$. Evaluating (12.5.4) for $f = \chi_E$, we obtain
\[
0 \leq \frac{1}{\rho(E)} \int_E g \, d\rho = \frac{\lambda(E)}{\rho(E)} \leq 1.
\]
Now the claim follows from elementary measure theory (see for example [126, Theorem 1.40]).

Using (12.5.3), we can rewrite (12.5.4) as
\[
\int_{\mathcal{X}} (1 - g) \, f \, d\lambda = \int_{\mathcal{X}} g \, f \, d\mu \quad \text{for all non-negative } f \in L^2(\mathcal{X}, d\rho). \tag{12.5.5}
\]
We introduce the Borel sets
\[
A = \{ x \in \mathcal{X} \mid 0 \leq g(x) < 1 \} \quad \text{and} \quad B = \{ x \in \mathcal{X} \mid g(x) = 1 \}.
\]
and define the measures $\lambda_a$ and $\lambda_s$ by
\[
d\lambda_a = \chi_A \, d\lambda \quad \text{and} \quad d\lambda_s = \chi_B \, d\lambda.
\]
Choosing $f = \chi_B$ in (12.5.5), one sees that $\mu(B) = 0$, implying that $\lambda_s \perp \mu$.

Moreover, since $g$ is bounded, we can evaluate (12.5.5) for
\[
f = (1 + g + \cdots + g^n) \, \chi_E
\]
for any $n \in \mathbb{N}$ and any Borel set $E$. Using the same transformation with “telescopic sums” as in the evaluation of the geometric or Neumann series, we obtain
\[
\int_E (1 - g^{n+1}) \, d\lambda = \int_E g \, (1 + g + \cdots + g^n) \, d\mu. \tag{12.5.6}
\]
At every point of $B$, the factor $(1 - g^{n+1})$ in the integrand on the left vanishes. At every point of $A$, on the other hand, the factor $(1 - g^{n+1})$ is monotone increasing in $n$ and converges to one. Hence Lebesgue’s monotone convergence theorem implies that the left side of (12.5.6) converges to
\[
\lim_{n \to \infty} \int_E (1 - g^{n+1}) \, d\lambda = \lambda(E \cap A).
\]
The integrand on the right side of (12.5.6), on the other hand, is monotone increasing in $n$, so that the limit
\[
h(x) := \lim_{n \to \infty} g(x) \, (1 + g(x) + \cdots + g^n(x)) \quad \text{exists in } \mathbb{R}_0^+ \cup \{\infty\}.
\]
Moreover, the monotone convergence theorem implies that
\[
\lim_{n \to \infty} \int_E g \, (1 + g + \cdots + g^n) \, d\mu = \int_E h \, d\mu \in \mathbb{R}_0^+ \cup \{\infty\}.
\]
We conclude that, in the limit $n \to \infty$, the relation (12.5.6) yields
\[
\lambda_a(E) = \lambda(E \cap A) = \int_E h \, d\mu \quad \text{for any Borel set } E.
\]
Choosing $E = \mathcal{X}$, one sees that $h \in L^1(\mathcal{X}, d\mu)$. This concludes the proof of (12.5.1).

Finally, the representation (12.5.1) implies that $\lambda_a \ll \mu$. $\Box$
12.6. Moment Measures

We now introduce an important concept needed for the existence proof: the moment measures. We again assume that the Hilbert space is finite-dimensional and that the measure $\rho$ is normalized (12.0.1). We consider $\mathcal{F}$ with the metric induced by the sup-norm on $L(\mathcal{F})$, i.e.

$$d(p, q) = \|p - q\|$$

(and $\|\cdot\|$ as in (5.5.6)). The basic difficulty in applying the abstract theorems is that $\mathcal{F}$ is not compact (indeed, it is a star-shaped in the sense that $p \in \mathcal{F}$ implies $\lambda p \in \mathcal{F}$ for all $\lambda \in \mathbb{R}$). If the metric space is non-compact, our abstract results no longer apply, as becomes clear in the following simple example.

**Example 12.6.1.** Consider the Banach space $C_0^0(\mathbb{R}, \mathbb{R})$ of compactly supported continuous functions. Let $\rho_n = \delta_n$ be the sequence of Dirac measures supported at $n \in \mathbb{N}$. Then for any $f \in C_0^0(\mathbb{R}, \mathbb{R})$,

$$\lim_{n \to \infty} \int_{-\infty}^{\infty} f \, d\rho_n = \lim_{n \to \infty} f(n) = 0 .$$

Hence the sequence $(\rho_n)_{n \in \mathbb{R}}$ converges as a measure to zero. Thus the limiting measure is no longer normalized. This shows that Theorem [12.3.2 fails to hold if the base space is non-compact. ♦

The way out is to make use of the fact that the causal action as well as the constraints are formed of functionals which are homogeneous under the scaling $p \to \lambda p$ of degree zero, one or two. This makes it possible to restrict attention to a compact subset of $\mathcal{F}$, and to consider three measures on this compact set.

**Definition 12.6.2.** Let $\mathcal{K}$ be the compact metric space

$$\mathcal{K} = \{ p \in \mathcal{F} \text{ with } \|p\| = 1 \} \cup \{0\} .$$

For a given measure $\rho$ on $\mathcal{F}$, we define the measurable sets of $\mathcal{K}$ by the requirement that the sets $\mathbb{R}^+ \Omega = \{ \lambda p \mid \lambda \in \mathbb{R}^+, p \in \Omega \}$ and $\mathbb{R}^- \Omega$ should be $\rho$-measurable in $\mathcal{F}$. We introduce the measures $m^{(0)}$, $m^{(1)}_\pm$, and $m^{(2)}$ by

$$m^{(0)}(\Omega) = \frac{1}{2} \rho(\mathbb{R}^+ \Omega \setminus \{0\}) + \frac{1}{2} \rho(\mathbb{R}^- \Omega \setminus \{0\}) + \rho(\Omega \cap \{0\}) \quad (12.6.1)$$

$$m^{(1)}_+(\Omega) = \frac{1}{2} \int_{\mathbb{R}^+ \Omega} \|p\| \, d\rho(p) \quad (12.6.2)$$

$$m^{(1)}_- (\Omega) = \frac{1}{2} \int_{\mathbb{R}^- \Omega} \|p\| \, d\rho(p) \quad (12.6.3)$$

$$m^{(2)}(\Omega) = \frac{1}{2} \int_{\mathbb{R}^+ \Omega} \|p\|^2 \, d\rho(p) + \frac{1}{2} \int_{\mathbb{R}^- \Omega} \|p\|^2 \, d\rho(p) . \quad (12.6.4)$$

The measures $m^{(l)}$ and $m^{(l)}_\pm$ are referred to as the $l$th moment measure.

We remark that, as a short notation, it is convenient to abbreviate the difference of the first moment measures by

$$m^{(1)}(\Omega) := m^{(1)}_+(\Omega) - m^{(1)}_- (\Omega) . \quad (12.6.5)$$

In mathematical terms, $m^{(1)}$ can be regarded as a signed measure (see for example [97], §28) or [126] Chapter 6]). For simplicity, we here avoid the concept of signed measures.
by working instead with the measures \( m_{\pm} \). Nevertheless, we introduce an \( m^{(1)} \)-integral as a short notation for the difference of the integrals with respect to \( m_{+}^{(1)} \) and \( m_{-}^{(1)} \), i.e.

\[
\int_{\mathcal{X}} h \, dm^{(1)} := \int_{\mathcal{X}} h \, dm_{+}^{(1)} - \int_{\mathcal{X}} h \, dm_{-}^{(1)},
\]

where for simplicity we always assume that \( h \) is continuous.

The \( \rho \)-integrals of homogeneous functions can be rewritten as integrals over \( K \) using the moment measures, as we now make precise.

**Definition 12.6.3.** A function \( h \in C^0(\mathcal{F}) \) is called **homogeneous of degree** \( \ell \) with \( \ell \in \{0, 1, 2\} \) if

\[
h(\nu x) = \nu^{\ell} h(x) \quad \text{for all } \nu \in \mathbb{R} \text{ and } x \in \mathcal{F}.
\]

(12.6.6)

**Lemma 12.6.4.** Let \( h \in C^0(\mathcal{F}) \) be a function which is homogeneous of degree \( \ell \in \{0, 1, 2\} \). Then

\[
\int_{\mathcal{F}} h \, d\rho = \int_{\mathcal{X}} h \, dm^{(\ell)}.
\]

**Proof.** We first note that, using the homogeneity (12.6.6), the function \( h \) is uniquely determined by its restriction to \( K \). Moreover, using an approximation argument with Lebesgue’s dominated convergence theorem, it suffices to consider a function \( h \) which is homogeneous of degree \( \ell \) and simple in the sense that its restriction to \( K \) takes a finite number of values, i.e.

\[
h|_K = \sum_{i=1}^{N} c_i \chi_{\Omega_i}
\]

with Borel sets \( \Omega_1, \ldots, \Omega_N \subset \mathcal{K} \) and real coefficients \( c_1, \ldots, c_N \). For such functions,

\[
\int_{\mathcal{F}} h \, d\rho = \sum_{i=1}^{N} c_i \int_{\mathbb{R}^+ \Omega_i} \|p\|^\ell \, dp(p) = \sum_{i=1}^{N} c_i \, m^\ell(\Omega_i) = \int_{\mathcal{X}} h \, dm^{(\ell)}.
\]

This concludes the proof. \( \square \)

Applying this lemma, the normalization \( \rho(\mathcal{F}) = 1 \) can be expressed in terms of the moment measures as

\[
m^{(0)}(\mathcal{K}) = 1,
\]

(12.6.7)

whereas the action (5.5.2) as well as the functionals in the constraints (5.5.5) and (5.5.4) can be written as

\[
S(\rho) = \int_{\mathcal{K} \times \mathcal{K}} \mathcal{L}(p, q) \, dm^{(2)}(p) \, dm^{(2)}(q)
\]

(12.6.8)

\[
T(\rho) = \int_{\mathcal{K} \times \mathcal{K}} |pq|^2 \, dm^{(2)}(p) \, dm^{(2)}(q)
\]

(12.6.9)

\[
\int_{\mathcal{F}} \text{tr}(x) \, d\rho(x) = \int_{\mathcal{X}} \text{tr}(p) \, dm^{(1)}_+(p) - \int_{\mathcal{X}} \text{tr}(p) \, dm^{(1)}_-(p).
\]

(12.6.10)

Working with the moment measures has the advantage that they are measures on the compact space \( \mathcal{K} \). We also learn that two measures \( \rho \) and \( \tilde{\rho} \) whose moment measures coincide yield the same values for the functionals \( S \) and \( T \) as well as for the integral (12.6.10) entering the trace constraint. It is most convenient to work exclusively with the moment measures. At the end, we shall construct a suitable representative \( \rho \) of the limiting
12.7. Existence of Minimizers for the Causal Action Principle

A key step for making this method work is the following a-priori estimate.

**Lemma 12.6.5.** There is a constant \( \varepsilon = \varepsilon (f, n) > 0 \) such that for every measure \( \rho \) on \( \mathcal{F} \) the corresponding moment measures (see Definition 12.6.2) satisfy for all measurable \( \Omega \subset \mathcal{K} \) the following inequalities:

\[
\left( m_+^{(1)}(\Omega) + m_-^{(1)}(\Omega) \right)^2 \leq m_+^{(0)}(\Omega) m^{(2)}(\Omega) \tag{12.6.11}
\]

\[
m^{(2)}(\mathcal{K}) \leq \frac{\sqrt{T(\rho)}}{\varepsilon}. \tag{12.6.12}
\]

**Proof.** The inequality (12.6.11) follows immediately from Hölder’s inequality,

\[
|2(m_+^{(1)}(\Omega) + m_-^{(1)}(\Omega))|^2 \leq \left( \int_{\mathbb{R}^\Omega} \|p\| d\rho(p) \right)^2 \leq \rho(\mathbb{R}^\Omega) \int_{\mathbb{R}^\Omega} \|p\|^2 d\rho(p) \leq 4m_+^{(0)}(\Omega) m^{(2)}(\Omega).
\]

In order to prove (12.6.12), we introduce the mapping

\[
\phi : \mathcal{K} \times \mathcal{K} \to \mathbb{R} : (p, q) \mapsto |pq|.
\]

Clearly, \( \phi \) is continuous and

\[
\phi(p, p) = |p|^2 = \text{Tr}(p^2) = \|p\|^2 = 1.
\]

Thus every point \( r \in \mathcal{K} \) has a neighborhood \( U(r) \subset \mathcal{K} \) with

\[
\phi(p, q) \geq \frac{1}{2} \quad \text{for all } p, q \in U(r). \tag{12.6.13}
\]

Since \( \mathcal{K} \) is compact, there is a finite number of points \( r_1, \ldots, r_N \) such that the corresponding sets \( U_i := U(r_i) \) cover \( \mathcal{K} \). Due to the additivity property of measures, there is an index \( i \in \{1, \ldots, N\} \) such that

\[
m^{(2)}(U_i) \geq \frac{m^{(2)}(\mathcal{K})}{N}. \tag{12.6.14}
\]

We write \( T \) in the form (12.6.9) and apply (12.6.13) as well as (12.6.14) to obtain

\[
T(\rho) \geq \iint_{U_i \times U_i} |pq|^2 dm^{(2)}(p) dm^{(2)}(q) \geq \frac{1}{2} m^{(2)}(U_i)^2 \geq \frac{m^{(2)}(\mathcal{K})^2}{2N^2}.
\]

Setting \( \varepsilon = 1/(\sqrt{2N}) \), the result follows. \( \square \)

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After the above preparations, we can follow the strategy of the direct method of the calculus of variations described at the beginning of Chapter 12 to obtain the following result.

**Theorem 12.7.1.** Let \( \mathcal{K} \) be a finite-dimensional Hilbert space and \( n \in \mathbb{N} \). Let \( \rho_n \) be a minimizing sequence of regular Borel measures on \( \mathcal{F} \) satisfying our constraints (5.5.3), (5.5.4) and (5.5.5) (for fixed and finite constants). Then there is a regular Borel measure \( \rho \) which also satisfies the constraints (again with the same constants) and

\[
S(\rho) = \liminf_{n \to \infty} S(\rho_n).
\]
In short, the method for constructing \( \rho \) is to take a limit of the moment measures of the \( \rho_k \) and to realize this limit by the measure \( \rho \). In more detail, we proceed as follows. In view of Lemma 12.6.5, we know that the moment measures are uniformly bounded measures on a compact metric space \( K \). Applying the compactness result of Theorem 12.3.2 (based on the Banach-Alaoglu theorem and the Riesz representation theorem) we conclude that for a suitable subsequence (which we again denote by \( (\rho_k) \)), the moment measures converge in the \( C_0^\infty(K)^* \)-topology to regular Borel measures,

\[
m_k^{(0)} \to m^{(0)} , \quad m_k^{(1)} \to m_\pm^{(1)} \quad \text{and} \quad m_k^{(2)} \to m^{(2)} ,
\]

which again have the properties (12.6.7), (12.6.11) and (12.6.12).

We next form the Radon-Nikodym decompositions of \( m_\pm^{(1)} \) and \( m^{(2)} \) with respect to \( m^{(0)} \). The inequality (12.6.11) shows that every set of \( m^{(0)} \)-measure zero is also a set of measure zero with respect to \( m_\pm^{(1)} \) and \( m^{(1)} \). In other words, the measures \( m_\pm^{(1)} \) are absolutely continuous with respect to \( m^{(0)} \). Hence, applying Theorem 12.5.2, we obtain the Radon-Nikodym decompositions

\[
\begin{align*}
dm_\pm^{(1)} &= f_\pm \, dm^{(0)} \quad \text{with} \quad f_\pm \in L^1(K, dm^{(0)}) .
\end{align*}
\]

As a consequence, the signed measure \( m^{(1)} \) in (12.6.5) has the decomposition

\[
\begin{align*}
dm^{(1)} &= f^{(1)} \, dm^{(0)} \quad \text{with} \quad f^{(1)} := f_+ - f_- \in L^1(K, dm^{(0)}) .
\end{align*}
\]

(12.7.1)

As we do not know whether also \( m^{(2)} \) is absolutely continuous with respect to \( m^{(0)} \), Theorem 12.5.2 gives the decomposition

\[
\begin{align*}
dm^{(2)} &= f^{(2)} \, dm^{(0)} + dm^{(2)}_{\text{sing}} \quad \text{with} \quad f^{(2)} \in L^1(K, dm^{(0)}) ,
\end{align*}
\]

(12.7.2)

where the measure \( m^{(2)}_{\text{sing}} \) is singular with respect to \( m^{(0)} \).

**Lemma 12.7.2.** The functions \( f^{(1)} \) and \( f^{(2)} \) in the Radon-Nikodym decompositions (12.7.1) and (12.7.2) can be chosen such that

\[
|f^{(1)}|^2 \leq f^{(2)} .
\]

**Proof.** Since \( m^{(2)}_{\text{sing}} \perp m^{(0)} \), there is a Borel set \( V \) such that

\[
\chi_V \, dm^{(0)} = dm^{(0)} \quad \text{and} \quad \chi_V \, dm^{(2)}_{\text{sing}} = 0 .
\]

Then, using the Radon-Nikodym decompositions (12.7.1) and (12.7.2) in (12.6.11), we obtain for any Borel set \( U \subset V \) the inequality

\[
\left| \int_U f^{(1)} \, dm^{(0)} \right|^2 \leq m^{(0)}(U) \int_U f^{(2)} \, dm^{(0)} .
\]

If the function \( f^{(1)} \) does not change signs on \( U \), we conclude that

\[
\inf_U |f^{(1)}|^2 \leq \sup_U f^{(2)} .
\]

By decomposing \( U \) into the two sets where \( f^{(1)} \) is positive and negative, respectively, one readily sees that this inequality even holds for any Borel set \( U \subset V \). As a consequence, the inequality \(|f^{(1)}|^2 \leq f^{(2)}\) holds almost everywhere (with respect to the measure \( m^{(0)} \)), concluding the proof. \( \square \)
In particular, we conclude that \( f^{(1)} \) even lies in \( L^2(\mathcal{K}, dm^{(0)}) \). Setting \( f = f^{(1)} \) and \( dn = (f^{(2)} - |f|^2) \, dm^{(0)} + dm^{(2)}_{\text{sing}} \), we obtain the decomposition

\[
dm^{(1)} = f \, dm^{(0)}, \quad dm^{(2)} = |f|^2 \, dm^{(0)} + dn,
\]

where \( f \in L^2(\mathcal{K}, dm^{(0)}) \), and \( n \) is a positive measure which need not be absolutely continuous with respect to \( m^{(0)} \). From the definition \( \text{(12.6.5)} \) it is clear that \( f \) is odd, i.e.

\[
f(-p) = -f(p) \quad \text{for all } p \in \mathcal{K}.
\]

The remaining task is to represent the limiting moment measures \( m^{(i)} \) in \( \text{(12.7.3)} \) by a measure \( \rho \). Unfortunately, there is the basic problem that such a measure can exist only if \( m^{(2)} \) is absolutely continuous with respect to \( m^{(0)} \), as the following consideration shows: Assume conversely that \( m^{(2)} \) is not absolutely continuous with respect to \( m^{(0)} \). Then there is a measurable set \( \Omega \subset \mathcal{K} \) with \( m^{(0)}(\Omega) = 0 \) and \( m^{(2)}(\Omega) \neq 0 \). Assume furthermore that there is a measure \( \rho \) on \( \mathcal{F} \) which represents the limiting moment measures in the sense that \( \text{(12.6.1)} - \text{(12.6.4)} \) hold. From \( \text{(12.6.1)} \) we conclude that the set \( \mathbb{R}\Omega \subset \mathcal{F} \) has \( \rho \)-measure zero. But then the integral \( \text{(12.6.4)} \) also vanishes, a contradiction.

This problem can also be understood in terms of the limiting sequence \( \rho_k \). We cannot exclude that there is a star-shaped region \( \mathbb{R}\Omega \subset \mathcal{F} \) such that the measures \( \rho_k(\mathbb{R}\Omega) \) tend to zero, but the corresponding moment integrals \( \text{(12.6.4)} \) have a non-zero limit. Using a notion from the calculus of variations for curvature functionals, we refer to this phenomenon as the possibility of bubbling. This bubbling effect is illustrated by the following example.

**Example 12.7.3. (Bubbling)** We choose \( f = 2 \) and \( n = 1 \). Furthermore, we let \( \mathcal{M} = [0, 1] \) with \( \mu \) the Lebesgue measure. For any parameters \( \kappa \geq 0 \) and \( \varepsilon \in (0, \frac{1}{2}) \), we introduce the mapping \( F_{\varepsilon} : \mathcal{M} \to \mathcal{F} \) by

\[
F_{\varepsilon}(x) = \frac{1}{1 - 2\varepsilon} \begin{cases}
-\kappa \varepsilon^{-\frac{1}{2}} \sigma^3 & \text{if } x \leq \varepsilon \\
1 + \sigma^1 \cos(\nu x) + \sigma^2 \sin(\nu x) & \text{if } \varepsilon < x \leq 1 - \varepsilon \\
\kappa \varepsilon^{-\frac{1}{2}} \sigma^3 & \text{if } x > 1 - \varepsilon,
\end{cases}
\]

where we set \( \nu = 2\pi/(1 - 2\varepsilon) \) (and \( \sigma^j \) are the Pauli matrices). The corresponding measure \( \rho_{\varepsilon} \) on \( \mathcal{F} \) has the following properties. On the set

\[
S := \{ 1 + \sigma^1 \sigma^1 + \sigma^2 \sigma^2 \text{ with } (\sigma^1)^2 + (\sigma^2)^2 = 1 \},
\]

which can be identified with a circle \( S^1 \), \( \rho_{\varepsilon} \) is a multiple of the Lebesgue measure. Moreover, \( \rho_{\varepsilon} \) is supported at the two points

\[
p_{\pm} := \pm \frac{\kappa \varepsilon^{-\frac{1}{2}}}{1 - 2\varepsilon} \sigma^3 \quad \text{with} \quad \rho_{\varepsilon}(\{p_+\}) = \rho_{\varepsilon}(\{p_-\}) = \varepsilon. \text{ (12.7.5)}
\]

A short calculation shows that the trace constraint is satisfied. Furthermore, the separations of the points \( p_+ \) and \( p_- \) from each other and from \( S \) are either spacelike or just in the boundary case between spacelike and timelike. Thus for computing the action, we only need to take into account the pairs \( (p_+, p_+), (p_-, p_-) \) as well as pairs \( (x, y) \) with \( x, y \in S \). A straightforward computation yields

\[
S(\rho_{\varepsilon}) = \frac{3}{(1 - 2\varepsilon)^2}, \quad \mathcal{T}(\rho_{\varepsilon}) = \frac{6}{(1 - 2\varepsilon)^2} + \frac{16\kappa^2}{(1 - 2\varepsilon)^3} + \frac{16\kappa^4}{(1 - 2\varepsilon)^4}. \text{ (12.7.6)}
\]
Let us consider the limit $\varepsilon \searrow 0$. From (12.7.6) we see that the functionals $S$ and $T$ converge,

$$
\lim_{\varepsilon \searrow 0} S = 3, \quad \lim_{\varepsilon \searrow 0} T = 6 + 16(\kappa^2 + \kappa^4).
$$

Moreover, there are clearly no convergence problems on the set $S$. Thus it remains to consider the situation at the two points $p_{\pm}$, (12.7.5), which move to infinity as $\varepsilon$ tends to zero. These two points enter the moment measures only at the corresponding normalized points $\hat{p}_{\pm} = \frac{p_{\pm}}{\|p_{\pm}\|} \in K$. A short calculation shows that the limiting moment measures $m(l) = \lim_{\varepsilon \searrow 0} m(l)$ satisfy the relations

$$
m(0)(\{\hat{p}_{\pm}\}) = m(1)(\{\hat{p}_{\pm}\}) = 0 \quad \text{but} \quad m(2)(\{\hat{p}_{\pm}\}) = \kappa^2 > 0.
$$

Hence $m(2)$ is indeed not absolutely continuous with respect to $m(0)$.

In order to avoid misunderstandings, we point out that this example does not show that bubbling really occurs for minimizing sequences, because we do not know whether the family $(\rho_\varepsilon)_{0<\varepsilon<1/2}$ is minimizing. But at least, our example shows that bubbling makes it possible to arrange arbitrary large values of $T$ without increasing the action $S$ (see (12.7.7) for large $\kappa$).

In order to handle possible bubbling phenomena, it is important to observe that the second moment measure does not enter the trace constraint. Therefore, by taking out the term $dn$ in (12.7.3) we decrease the functionals $S$ and $T$ (see (12.6.8) and (12.6.9)), without affecting the trace constraint. It remains to show that the resulting moment measure can indeed be realized by a measure $\rho$. This is shown in the next lemma.

**Lemma 12.7.4.** For any normalized regular Borel measure $m(0)$ on $K$ and any function $f \in L^2(K, \mathbb{R})$, there is a normalized regular Borel measure $\hat{\rho}$ whose moment measures $\hat{m}(l)$ are given by

$$
\hat{m}(0) = m(0), \quad \hat{m}(1) = f dm(0), \quad \hat{m}(2) = |f|^2 dm(0).
$$

**Proof.** We introduce the mapping

$$
F : K \to \mathcal{F}, \quad F(x) = f(x) x.
$$

Choosing $\hat{\rho} := F_* m(0)$, a direct computation shows that the corresponding moment measures indeed satisfy (12.7.8).

This concludes the proof of Theorem (12.7.1). We finally remark that the fact that we dropped the measure $dn$ in (12.7.3) implies that the value of $T$ might decrease in the limit. This is the only reason why the boundedness constraint (5.5.5) is formulated as an inequality, and not as an equality. It is not clear if the causal action principle also admits minimizers if the inequality in (5.5.5) is replaced by an equality. We conjecture is that the answer is yes. But at present, there is no proof. For the physical applications, it makes no difference if (5.5.5) is an equality or an inequality, because then one works with the corresponding Euler-Lagrange equations, into which the constraints enter via Lagrange multiplier terms (for details see [13]).

12.8. Existence of Minimizers for Causal Variational Principles in the Non-Compact Setting

In Theorem (12.7.1) the existence of minimizers was established in the case that the Hilbert space $K$ is finite-dimensional and the total volume $\rho(F)$ of spacetime is finite. From the physical point of view, this existence result is quite satisfying, because one
12.8. EXISTENCE OF MINIMIZERS IN THE NON-COMPACT SETTING

can take the point of view that it should be possible to describe our universe at least approximately by a causal fermion system with \( \dim \mathcal{H} < \infty \) and \( \rho(\mathcal{F}) < \infty \). From the mathematical point of view, however, it is interesting and important to also study the cases of an infinite-dimensional Hilbert space and/or infinite total volume. The case \( \dim \mathcal{H} < \infty \) and \( \rho(\mathcal{F}) = \infty \) is not sensible (see Exercise 12.6). In the case \( \dim \mathcal{H} = \infty \) and \( \rho(\mathcal{F}) < \infty \), on the other hand, there are minimizing sequences which converge to zero. Therefore, it remains to study the \textit{infinite-dimensional setting} \( \dim \mathcal{H} < \infty \) and \( \rho(\mathcal{F}) < \infty \) already mentioned in Section 5.5. In this setting, the existence theory is difficult and has not yet been developed. Therefore, our strategy is to approach the problem in two steps. The first step is to deal with infinite total volume; this has been carried out in [72]. The second step, which involves the difficulty of dealing with non-locally compact spaces, is currently under investigation (for first results see [108]).

We now outline the basic strategy in the simplest possible case (more details and a more general treatment can be found in [72]. We consider \textit{causal variational principles in the non-compact setting} as introduced in Section 6.3. We assume that the Lagrangian is smooth,

\[ \mathcal{L} \in C^\infty(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+), \]

and has the properties (i) and (ii) on page 87. Moreover, we make the following simplifying assumption.

**Definition 12.8.1.** The Lagrangian has **compact range** if for every compact set \( K \subset \mathcal{F} \) there is a compact set \( K' \subset \mathcal{F} \) such that

\[ \mathcal{L}(x, y) = 0 \quad \text{for all } x \in K \text{ and } y \notin K'. \]

The goal of this section is to prove the following theorem.

**Theorem 12.8.2.** Under the above assumptions, there is a regular Borel measure \( \rho \) on \( \mathcal{F} \) (not necessarily bounded) which satisfies the EL equations

\[ \ell\big|_{\text{supp } \rho} \equiv \inf_M \ell = 0 \quad \text{with} \quad \ell(x) := \int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho(y) - 1. \tag{12.8.1} \]

For the proof, we first exhaust \( \mathcal{F} \) by compact sets \( (K_j)_{j \in \mathbb{N}}, \text{i.e.} \)

\[ K_1 \subset K_2 \subset \cdots \subset \mathcal{F} \quad \text{and} \quad \bigcup_{j=1}^{\infty} K_j = \mathcal{F}. \]

On each \( K_j \) we consider the restricted variational principle where we minimize the action

\[ \mathcal{S}_{K_j}(\rho) = \int_{K_j} d\rho(x) \int_{K_j} d\rho(y) \mathcal{L}(x, y) \]

under variations of \( \rho \) within the class of normalized regular Borel measure on \( K_j \). Using the existence theory in the compact setting (see Theorem 12.3.3), each of these restricted variational principles has a minimizer, which we denote by \( \rho_j \). Each of these measures satisfies the EL equations stated in Theorem 7.1.1. Thus, introducing the functions

\[ \ell_j \in C^0(K_j, \mathbb{R}), \quad \ell_j(x) := \int_{K_j} \mathcal{L}(x, y) \, d\rho_j(y) - s_j, \]

one can choose the parameters \( s_j > 0 \) such that

\[ \ell_j\big|_{\text{supp } \rho_j} \equiv \inf_{K_j} \ell_j = 0. \]
Typically, the support of the measures $\rho_j$ will be “spread out” over larger and larger subsets of $\mathcal{F}$. This also means that, working with normalized measures, the measures $\rho_j$ typically converge to the trivial measure $\rho = 0$. In order to ensure a non-trivial measure, we must perform a suitable rescaling. To this end, we introduce the measures

$$\tilde{\rho}_j = \frac{\rho_j}{\tilde{s}_j}.$$ 

These new measures are no longer normalized, but they satisfy the EL equations with $\tilde{s}_j = 1$, i.e.

$$\tilde{\ell}_j|_{\text{supp} \tilde{\rho}_j} \equiv \inf_{K_j} \tilde{\ell}_j = 0 \quad \text{with} \quad \tilde{\ell}_j(x) := \int_{K_j} \mathcal{L}(x, y) \, d\tilde{\rho}_j(y) - 1. \quad (12.8.2)$$

Our next task is to construct a limit measure $\rho$ of the measures $\tilde{\rho}_j$. We first extend the measures $\tilde{\rho}_j$ by zero to all of $\mathcal{F}$ and denote them by $\rho^{[j]}$, 

$$\rho^{[j]}(U) := \rho_j(U \cap K_j) \quad \text{for any Borel subset } U \subset \mathcal{F}.$$ 

In the next lemma we show that these measures are bounded on every compact set.

**Lemma 12.8.3.** For every compact subset $K \subset \mathcal{F}$ there is a constant $C_K > 0$ such that

$$\rho^{[j]}(K) \leq C_K \quad \text{for all } j \in \mathbb{N}. \quad (12.8.3)$$

**Proof.** Since $\mathcal{L}(x, \cdot)$ is continuous and strictly positive at $x$, there is an open neighborhood $U(x)$ of $x$ with

$$\mathcal{L}(y, z) \geq \frac{\mathcal{L}(x, x)}{2} > 0 \quad \text{for all } y, z \in U(x).$$

Covering $K$ by a finite number of such neighborhoods $U(x_1), \ldots, U(x_L)$, it suffices to show the inequality (12.8.3) for the sets $K \cap U(x_\ell)$ for any $\ell \in \{1, \ldots, L\}$. Moreover, we choose $N$ so large that $K_N \supseteq K$ and fix $k \geq N$. If $K \cap \text{supp} \rho^{[k]} = \emptyset$, there is nothing to prove. Otherwise, there is a point $z \in K \cap \text{supp} \rho^{[k]}$. Using the EL equations (12.8.2) at $z$, it follows that

$$1 = \int_\mathcal{F} \mathcal{L}(z, y) \, d\rho^{[k]}(y) \geq \int_{U(x_\ell)} \mathcal{L}(z, y) \, d\rho^{[k]}(y) \geq \frac{\mathcal{L}(x_\ell, x_\ell)}{2} \rho^{[k]}(U(x_\ell)).$$

Hence

$$\rho^{[k]}(U(x_\ell)) \leq \frac{2}{\mathcal{L}(x_\ell, x_\ell)}. \quad (12.8.4)$$

This inequality holds for any $k \geq N$. We introduce the constants $c(x_\ell)$ as the maximum of $2/\mathcal{L}(x_\ell, x_\ell)$ and $\rho^{[1]}(U(x_\ell)), \ldots, \rho^{[N-1]}(U(x_\ell))$. Since the open sets $U(x_1), \ldots, U(x_L)$ cover $K$, we finally introduce the constant $C_K$ as the sum of the constants $c(x_1), \ldots, c(x_L)$.

Given a compact set $K$, combining the result of the previous lemma with the compactness of measures on compact topological spaces (see Theorem (12.3.2)), we conclude that there is a subsequence $(\rho^{[j_n]})$ whose restrictions to $K$ converge as a measure (i.e. in the sense (12.3.1)). Proceeding inductively for the compact sets $K_1, K_2, \ldots$ and choosing a diagonal sequence, one gets a subsequence of measures on $\mathcal{F}$, denoted by $\rho^{(k)}$, whose restriction to any compact set $K_j$ converges, i.e.

$$\rho^{(k)}|_{K_j} \text{ converges as } k \to \infty \text{ to } \rho|_{K_j} \text{ for all } j \in \mathbb{N}, \quad (12.8.5)$$
where $\rho$ is a regular Borel measure on $\mathcal{F}$ (typically of infinite total volume). The convergence of measures in (12.8.5) is referred to as *vague convergence* (for more details see [8, Definition 30.1] or [72, Section 4.1]).

It remains to show that the obtained measure $\rho$ is a non-trivial minimizer. In order to show that it is non-trivial, we make use of the EL equations (12.8.2). Let $x \in \mathcal{F}$. Then (12.8.2) implies that

$$\int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho^{(k)}(y) \geq 1.$$  

Since $\mathcal{L}$ has compact range, we may pass to the limit to conclude that

$$\int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho(y) \geq 1. \quad (12.8.6)$$

This shows (in a quantitative way) that the measure $\rho$ is non-zero.

Our final step for proving the EL equations (12.8.1) is to show that the EL equations (12.8.2) are preserved in the limit. In view of the lower bound (12.8.6), it remains to show that $\ell$ vanishes on the support of $\rho$. Thus let $x \in \text{supp} \rho$. We choose a compact subset $K \subset \mathcal{F}$ such that $x$ lies in its interior. Again using that the Lagrangian has compact range, there is another compact subset $K' \subset \mathcal{F}$ such that (12.8.1) holds. The fact that $x$ lies in the support and that the measures $\rho^{(k)}$ converge vaguely to $\rho$ implies that there is a sequence $x_k \in \text{supp} \rho^{(k)}$ which converges to $x$. The EL equations for each $\rho^{(k)}$ imply that, for sufficiently large $k$,

$$\int_{K'} \mathcal{L}(x_k, y) \, d\rho^{(k)}(y) = 1.$$  

Taking the limit is a bit subtle because both the argument $x_k$ of the Lagrangian and the integration measure depend on $k$. Therefore, we begin with the estimate

$$\left| \int_{K'} \mathcal{L}(x, y) \, d\rho(y) - \int_{K'} \mathcal{L}(x_k, y) \, d\rho^{(k)}(y) \right| \leq \left| \ell(x) - \ell^{(k)}(x) \right| + \sup_j \left| \ell^{(j)}(x) - \ell^{(j)}(x_k) \right|, \quad (12.8.7)$$

where we set

$$\ell^{(k)}(z) := \int_{K'} \mathcal{L}(z, y) \, d\rho^{(k)}(y) - 1.$$  

The first summand on the right side of (12.8.7) tends to zero because the measures $\rho^{(k)}$ converge vaguely to $\rho$. The second summand, on the other hand, tends to zero because the functions $\ell^{(j)}$ are equicontinuous (for more details on this argument see [72, Section 4.2]). This concludes the proof of Theorem 12.8.2.

We finally note that, starting from the EL equations (12.8.1), one can also show that $\rho$ is a *minimizer under variations of finite volume*, meaning that for every regular Borel measure $\tilde{\rho}$ satisfying (6.3.1), the difference of actions (6.3.2) is nonnegative (6.3.3). The proof can be found in [72, Section 4.3].

**Exercises**

**Exercise 12.1.** Let $\Lambda$ be the functional

$$\Lambda : C^0([0, 1], \mathbb{R}) \to \mathbb{R}, \quad \Lambda(f) = \sup_{x \in [0, 1]} f(x).$$
Can this functional be represented by a measure? Analyze how your findings are compatible with the Riesz representation theorem.

**Exercise 12.2.** Let \( \rho \) be the Borel measure on \([0, \pi]\) given by
\[
\rho(\Omega) = \int_{\Omega} \sin x \, dx + \sum_{n=1}^{\infty} \frac{1}{n^2} \chi_{\Omega} \left( \frac{1}{n} \right).
\]
Compute the Lebesgue decomposition of \( \rho \) with respect to the Lebesgue measure.

**Exercise 12.3.** Let \( F = \mathbb{R}^2 \) and \( K = S^1 \cup \{0\} \) be a compact subset of \( F \). Given a Borel measure \( \rho \) on \( F \), the moment measures \( m^{(0)}, m^{(1)} \) and \( m^{(2)} \) can be defined just as in the lecture. Compute these moment measures for the following choices of \( \rho \):
(a) \( \rho = F_\ast \mu_{S^1} \), where \( F : S^1 \hookrightarrow \mathbb{R}^2 \) is the natural injection and \( \mu_{S^1} \) the normalized Lebesgue measure on \( S^1 \).
(b) \( \rho = \delta_{(0,0)} + \delta_{(1,1)} + \delta_{(5,0)} \) (where \( \delta_{(x,y)} \) denotes the Dirac measure supported at \((x, y) \in \mathbb{R}^2\)).
(c) \( \rho = F_\ast \mu_{\mathbb{R}} \), where \( F : \mathbb{R} \to \mathbb{R}^2 \quad F(x) = (x, 2) \)
and \( \mu_{\mathbb{R}} \) is the Lebesgue measure.

**Exercise 12.4.** (a) Let \((\rho_n)_{n \in \mathbb{N}}\) be a sequence of normalized regular Borel measures on \( \mathbb{R} \) with the property that there is a constant \( c > 0 \) such that
\[
\int_{-\infty}^{\infty} x^2 \, d\rho_n(x) \leq c \quad \text{for all } n.
\]
Show that a subsequence converges again to a normalized Borel measure on \( \mathbb{R} \).
*Hint:* Apply the compactness result of the lecture to the measures restricted to the interval \([-L, L]\) and analyze the behavior as \( L \to \infty \).
(b) More generally, assume that for a given non-negative function \( f(x) \),
\[
\int_{-\infty}^{\infty} f(x) \, d\rho_n(x) \leq c \quad \text{for all } n.
\]
Which condition on \( f \) ensures that the a subsequence of the measures converges to a normalized Borel measure? Justify your result by a counter example.

**Exercise 12.5.** Let \( \mathcal{M} \subset \mathbb{R} \) be a closed embedded submanifold of \( \mathbb{R}^3 \). We choose a compact set \( K \subset \mathbb{R}^3 \) which contains \( \mathcal{M} \). On \( C^0(K, \mathbb{R}) \) we introduce the functional
\[
\Lambda : C^0(K, \mathbb{R}) \to \mathbb{R}, \quad \Lambda(f) = \int_{\mathcal{M}} f(x) \, d\mu_{\mathcal{M}}(x)
\]
(where \( d\mu_{\mathcal{M}} \) is the volume measure corresponding to the induced Riemannian metric on \( \mathcal{M} \)). Show that this functional is linear, bounded and positive. Apply the Riesz representation theorem to represent this functional by a Borel measure on \( K \). What is the support of this measure?

**Exercise 12.6.** This exercise explains why the causal action principle is ill-posed in the case \( \dim \mathcal{H} = \infty \) and \( \rho(\mathcal{H}) < \infty \). The underlying estimates were first given in the setting of discrete spacetimes in [11], Lemma 5.1.
(a) Let $\mathcal{H}_0$ be a finite-dimensional Hilbert space of dimension $n$ and $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$ be a causal fermion system of finite total volume $\rho_0(\mathcal{F}_0)$. Let $\iota : \mathcal{H}_0 \rightarrow \mathcal{H}$ be an isometric embedding. Construct a causal fermion system $(\mathcal{H}, \mathcal{F}, \rho)$ which has the same action, the same total volume and the same values for the trace and boundedness constraints as the causal fermion system $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$.

(b) Let $\mathcal{H}_1 = \mathcal{H}_0 \oplus \mathcal{H}_0$. Construct a causal fermion system $(\mathcal{H}_1, \mathcal{F}_1, \rho_1)$ which has the same total volume and the same value of the trace constraint as $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$ but a smaller action and a smaller value of the boundedness constraint. Hint: Let $F_1/2 : \mathcal{L}(\mathcal{H}_0) \rightarrow \mathcal{L}(\mathcal{H}_1)$ be the linear mappings

$$(F_1(A))(u \oplus v) = (Au) \oplus 0, \quad (F_2(A))(u \oplus v) = 0 \oplus (Av).$$

Show that $F_1/2$ map $\mathcal{F}_0$ to $\mathcal{F}_1$. Define $\rho_1$ by

$$\rho_1 = \frac{1}{2}((F_1)_* \rho + (F_2)_* \rho).$$

(c) Iterate the construction in (b) and apply (a) to obtain a series of universal measures on $\mathcal{F}$ of fixed total volume and with fixed value of the trace constraint, for which the action and the values of the boundedness constraint tend to zero. Do these universal measures converge? If yes, what is the limit?

EXERCISE 12.7. (Riesz representation theorem - part 1) Let $\Lambda$ be the functional

$$\Lambda : C^0([0,1], \mathbb{R}) \rightarrow \mathbb{R}, \quad \Lambda(f) := \sup_{x \in [0,1]} f(x).$$

Can this functional be represented by a measure? Analyze how your findings are compatible with the Riesz representation theorem.

EXERCISE 12.8. (Riesz representation theorem - part 2) Let $\mathbb{M}$ be a closed embedded submanifold of $\mathbb{R}^3$. We choose a compact set $K \subset \mathbb{R}^3$ which contain $\mathbb{M}$. On $C^0(K, \mathbb{R})$ we introduce the functional

$$\Lambda : C^0(K, \mathbb{R}) \rightarrow \mathbb{R}, \quad \Lambda(f) = \int_{\mathbb{M}} f(x) d\mu_{\mathbb{M}}(x),$$

where $d\mu_{\mathbb{M}}$ is the volume measure corresponding to the induced Riemannian metric on $\mathbb{M}$. Show that this functional is linear, bounded and positive. Apply the Riesz representation theorem to represent this functional by a Borel measure on $K$. What is the support of this measure?

EXERCISE 12.9. (Radon-Nikodym decomposition) Let $\rho$ be the Borel measure on $[0, \pi]$ given by

$$\rho(\Omega) := \int_{\Omega} \sin x \, dx + \sum_{n=1}^{\infty} \frac{1}{n^2} \chi_{\Omega}(\frac{1}{n}).$$

Compute the Radon-Nikodym decomposition of $\rho$ with respect to the Lebesgue measure.

EXERCISE 12.10. (Normalized regular Borel measures: compactness results)

(a) Let $(\rho_n)_{n \in \mathbb{N}}$ be a sequence of normalized regular Borel measures on $\mathbb{R}$ with the property that there is a constant $c > 0$ such that

$$\int_{\mathbb{R}} x^2 \, d\rho_n(x) \leq c \quad \text{for all } n \in \mathbb{N}.$$
Show that a subsequence converges again to a normalized Borel measure on \( \mathbb{R} \).

**Hint:** Apply the compactness result of the lecture to the measures restricted to the interval \([-L, L]\) and analyze the behavior as \( L \to \infty \).

(b) More generally, assume that for a given non-negative function \( f \in C^0(\mathbb{R}, \mathbb{R}) \),

\[
\int_{\mathbb{R}} f(x) \, d\rho_n(x) \leq c \quad \text{for all } n \in \mathbb{N}.
\]

Which condition on \( f \) ensures that a subsequence of the measures converges to a normalized Borel measure? Justify your results by a counterexample.

**Exercise 12.11.** *(Derivative of measures)* Let \( \mu \) be the counting measure on the \( \sigma \)-algebra \( \mathcal{P}(\mathbb{N}) \). Consider the measure

\[
\lambda(\emptyset) := 0, \quad \lambda(E) := \sum_{n \in E} (1 + n)^2, \quad E \in \mathcal{P}(\mathbb{N}).
\]

Show that \( \mu \) and \( \lambda \) are equivalent (one is absolutely continuous with respect to the other) and determine the Radon-Nikodym derivative \( \frac{d\mu}{d\lambda} \).

**Exercise 12.12.** *(Minimizers)* Let \( M \) denote the 2-sphere \( S^2 \subset \mathbb{R}^3 \) and let \( d\mu_M \) be the normalized canonical surface measure. Consider a Lagrangian on \( M \times M \) defined by

\[
\mathcal{L}(x, y) := \frac{1}{1 + \|x - y\|_{\mathbb{R}^3}} \quad \text{for all } x, y \in M.
\]

Show that the action \( \mathcal{S}(\mu_M) \) is minimal under variations of the form

\[
d\rho_{x_0, t} := (1 - t)d\mu_M + td\delta_{x_0}, \quad \text{with } t \in [0, 1),
\]

where \( \delta_{x_0} \) is the Dirac measure centered at \( x_0 \in M \).

**Exercise 12.13.** *(Moment measures)* Let \( \mathcal{F} = \mathbb{R}^2 \) and \( K = S^1 \cup \{0\} \) be a compact subset of \( \mathcal{F} \). Given a Borel measure \( \rho \) on \( \mathcal{F} \), the moment measures \( m^{(0)}, m^{(1)} \) and \( m^{(2)} \) can be defined just as in the lecture. Compute these moment measures for the following choices of \( \rho \):

(a) \( \rho = F_* (\mu_{S^1}) \), where \( F : S^1 \to \mathbb{R}^2 \) is the natural injection and \( \mu_{S^1} \) is the normalized Lebesgue measure on \( S^1 \).

(b) \( \rho = \delta_{(0,0)} + \delta_{(1,1)} + \delta_{(5,0)} \) (where \( \delta_{(x,y)} \) denote the Dirac measure supported at \( (x, y) \in \mathbb{R}^2 \)).

(c) \( \rho = F_* (\mu_\mathbb{R}) \), where \( \mu_\mathbb{R} \) is the Lebesgue measure on \( \mathbb{R} \) and

\[
F : \mathbb{R} \to \mathbb{R}^2, \quad F(x) = (x, 2).
\]
Methods of Hyperbolic Partial Differential Equations

13.1. The Cauchy Problem, Linear Symmetric Hyperbolic Systems

In this section, we shall prove that the Cauchy problem for the Dirac equation in the presence of an external potential has a unique global solution. Moreover, we want to show that the finite speed of propagation as postulated by special relativity is indeed respected by the solutions of the Dirac equation. For later purposes, it is preferable to include an inhomogeneity. Thus we consider the Cauchy problem in Minkowski space

$$(i\partial + B - m)\psi = \phi \in C^\infty(M, S.M), \quad \psi|_{t_0} = \psi_0 \in C^\infty(\mathbb{R}^3, S.M) \quad (13.1.1)$$

for a given inhomogeneity $\phi$ and initial data $\psi_0$. In order to make the standard methods available, we multiply the equation by $-i\gamma^0$,

$$1_{\mathbb{C}^4}\partial_t \psi + \gamma^0 \nabla \psi - i\gamma^0 (B - m)\psi = -i\gamma^0 \phi. \quad (13.1.2)$$

Now the matrices in front of the derivatives are all Hermitian (with respect to the standard scalar product on $\mathbb{C}^4$). Moreover, the matrix in front of the time derivative is positive definite. Kurt Otto Friedrichs [86] observed that these properties are precisely what is needed in order to get a well-posed Cauchy problem. He combined these properties in the notion of a symmetric hyperbolic system. We now give its general definition. More specifically, we consider a system of $N$ complex-valued equations with spatial coordinates $\vec{x} \in \mathbb{R}^m$ and time $t$ in an interval $[0, T]$ with $T > 0$. The initial data will always be prescribed at time $t = 0$. For notational clarity, we denote partial derivatives in spatial direction by $\nabla$.

**Definition 13.1.1.** A linear system of differential equations of the form

$$A^0(t, \vec{x}) \partial_t u(t, \vec{x}) + \sum_{\alpha=1}^m A^\alpha(t, \vec{x}) \nabla_\alpha u(t, \vec{x}) + B(t, \vec{x}) u(t, \vec{x}) = w(t, \vec{x}) \quad (13.1.3)$$

with

$$A^0, A^\alpha, B \in C^\infty([0, T] \times \mathbb{R}^m, L(\mathbb{C}^N)), \quad w \in C^\infty([0, T] \times \mathbb{R}^m, \mathbb{C}^N).$$

is called symmetric hyperbolic if

(i) The matrices $A^0$ and $A^\alpha$ are Hermitian,

$$(A^0)^\dagger = A^0 \quad \text{and} \quad (A^\alpha)^\dagger = A^\alpha$$

(where $\dagger$ denotes the adjoint with respect to the canonical scalar product on $\mathbb{C}^N$).

(ii) The matrix $A^0$ is uniformly positive definite, i.e. there is a positive constant $C$ such that

$$A^0(t, \vec{x}) > C 1_{\mathbb{C}^N} \quad \text{for all} \ (t, \vec{x}) \in ([0, T] \times \mathbb{R}^m).$$

In the case $w \equiv 0$, the linear system is called homogeneous.
A good reference for linear symmetric hyperbolic systems is the book by Fritz John [102, Section 5.3] (who was Friedrichs’ colleague at the Courant Institute). Our presentation was also influenced by [123, Chapter 8]. We remark that the concept of symmetric hyperbolic systems can be extended to nonlinear equations of the form

\[ A^0(t, \vec{x}, u) \partial_t u(t, \vec{x}) + \sum_{\alpha=1}^m A^\alpha(t, \vec{x}, u) \nabla_\alpha u(t, \vec{x}) + B(t, \vec{x}, u) = 0, \]

where the matrices \( A^0 \) and \( A^\alpha \) should again satisfy the above conditions (i) and (ii). For details we refer to [133, Section 16] or [125, Section 7]. Having the Dirac equation in mind, we always restrict attention to linear systems. We also note that an alternative method for proving existence of fundamental solutions is to work with the so-called Riesz distributions (for a good textbook see [6]). Yet another method is to work with estimates in the interaction picture [24]. For completeness we finally remark that the concept of symmetric hyperbolic systems was extended by Friedrichs to so-called symmetric positive systems [87].

It is a remarkable fact that all partial differential equations in relativistic physics as well as most wave-type equations can be rewritten as a symmetric hyperbolic system. As an illustration, we now explain this reformulation in the example of a scalar hyperbolic equation.

**Example 13.1.2.** Consider a scalar hyperbolic equation of the form

\[ \partial_{tt} \phi(t, \vec{x}) = \sum_{\alpha, \beta=1}^m a_{\alpha \beta} \partial_t \phi(t, \vec{x}) \nabla_{\alpha \beta} \phi + \sum_{\alpha=1}^m b_\alpha(t, \vec{x}) \nabla_\alpha \phi + c(t, \vec{x}) \partial_t \phi + d(t, \vec{x}) \phi \quad (13.1.4) \]

with \((a_{\alpha \beta})\) a symmetric, uniformly positive matrix (in the case \( a_{\alpha \beta} = \delta_{\alpha \beta} \) and \( b,c,d = 0 \), one gets the scalar wave equation). Now the initial data prescribes \( \phi \) and its first time derivatives,

\[ \phi|_{t=0} = \phi_0 \in C^\infty(M), \quad \partial_t \phi|_{t=0} = \phi_1 \in C^\infty(M). \quad (13.1.5) \]

In order to rewrite the equation as a symmetric hyperbolic system, we introduce the vector \( u \) with \( k := m + 2 \) components by

\[ u_1 = \nabla_1 \phi, \ldots, u_m = \nabla_m \phi, \quad u_{m+1} = \partial_t \phi, \quad u_{m+2} = \phi. \quad (13.1.6) \]

Then the system

\[
\begin{align*}
\sum_{\beta=1}^m a_{\alpha \beta} \partial_t u_{\beta} - \sum_{\beta=1}^m a_{\alpha \beta} \nabla_\beta u_{m+1} = 0 \\
- \sum_{\alpha, \beta=1}^m a_{\alpha \beta} \nabla_\beta u_\alpha - \sum_{\alpha=1}^m b_\alpha u_\alpha + \partial_t u_{m+1} - c u_{m+1} - d u_{m+2} = 0 \\
0 - u_{m+1} + \partial_t u_{m+2} = 0
\end{align*}
\quad (13.1.7)
\]

is symmetric hyperbolic (as one verifies by direct inspection). Also, a short calculation shows that if \( \phi \) is a smooth solution of the scalar equation \((13.1.4)\), then the corresponding vector \( u \) is a solution of the system \((13.1.7)\). Conversely, assume that \( u \) is a smooth solution of \((13.1.7)\) which satisfies the initial condition \( u|_{t=0} = u_0 \), where \( u_0 \) is determined by \( \phi_0 \) and \( \phi_1 \) via \((13.1.6)\). Setting \( \phi = u_{m+2} \), the last line in \((13.1.7)\) shows that \( u_{m+1} = \partial_t \phi \). Moreover, the first line in \((13.1.7)\) implies that \( \partial_t u_\beta = \nabla_\beta u_{m+1} = \partial_t \nabla_\beta \phi \). Integrating over \( t \) and using that the relation \( u_\beta = \nabla_\beta \phi \) holds initially, we conclude that
this relation holds for all times. Finally, the second line in (13.1.7) yields that \( \phi \) satisfies the scalar hyperbolic equation (13.1.4). In this sense, the Cauchy problem for the system (13.1.7) is equivalent to that for the scalar equation (13.1.5).

This procedure works similarly for other physical equations like the Klein-Gordon or Maxwell equations. Exercise 13.1 is concerned with the example of the homogeneous Maxwell equations.

### 13.2. Finite Propagation Speed and Uniqueness of Solutions

For what follows, it is convenient to combine the time and spatial coordinates to a spacetime vector \( x = (t, \vec{x}) \in [0, T] \times \mathbb{R}^m \). We denote the spacetime dimension by \( n = m + 1 \). Moreover, setting \( \partial_0 \equiv \partial_t \), we use latin spacetime indices \( i \in \{0, \ldots, m\} \) and use the Einstein summation convention. Then our linear system (13.1.3) can be written in the compact form

\[
A^j(x) \partial_j u(x) + B(x) u(x) = w(x) .
\]

Next, a direction in spacetime can be described by a vector \( \xi = (\xi_i)_{i=0, \ldots, m} \in \mathbb{R}^{m+1} \). Contracting with the matrices \( A^j(x) \), we obtain the Hermitian \( N \times N \)-matrix

\[
A(x, \xi) := A^j(x) \xi_j ,
\]

referred to as the characteristic matrix. Note that in the example of the Dirac equation (13.1.2), the index \( i \) is a vector index in Minkowski space, and \( \xi \) should be regarded as a co-vector (i.e. a vector in the cotangent bundle). One should keep in mind that, despite the suggestive notation, the equation (13.2.1) should not be considered as being manifestly covariant, because it corresponds to the Hamiltonian formulation (13.1.2), where a time direction is distinguished.

The determinant of the characteristic matrix is referred to as the characteristic polynomial, being a polynomial in the components \( \xi_i \). For our purposes, it is most helpful to consider whether the characteristic matrix is positive or negative definite. If the vector \( \xi = (\tau, \vec{0}) \) points in the time direction, then \( A(x, \xi) = \tau A^0 \), which in view of Definition 13.2.1 is a definite matrix. By continuity, \( A(x, \xi) \) is definite if the spatial components of \( \xi \) are sufficiently small. In the example of the Dirac equation (13.1.2), the fact that

\[
A(x, \xi) = 1\xi_0 + \tau^0 \vec{\gamma} \vec{\xi}
\]

has eigenvalues \( \xi^0 \pm |\vec{\xi}| \)

shows that \( A(x, \xi) \) is definite if and only if \( \xi \) is a timelike vector. Moreover, it is positive definite if and only if \( \xi \) is future-directed and timelike. This suggests that the causal properties of the equation are encoded in the positivity of the characteristic matrix. We simply use this connection to define the causal structure for a general symmetric hyperbolic system.

**Definition 13.2.1.** The vector \( \xi \in \mathbb{R}^{m+1} \) is called timelike at the spacetime point \( x \) if the characteristic matrix \( A(x, \xi) \) is definite. A timelike vector is called future-directed if \( A(x, \xi) \) is positive definite. If the characteristic polynomial vanishes, then the vector \( \xi \) is called lightlike. A hypersurface \( \mathcal{H} \subset [0, T] \times \mathbb{R}^m \) with normal \( \nu \) is called spacelike if the matrix \( A(x, \nu) \) is positive definite for all \( x \in \mathcal{H} \).

The notion “normal” used here requires an explanation. The simplest method is to represent the hypersurface locally as the zero set of a function \( \phi(x) \). Then the normal can be defined as the gradient of \( \phi \). In this way, the gradient is a co-vector, so that the contraction \( A^j \nu_j = A^j \partial_j \phi \) is well-defined without referring to a scalar product. In particular, the last definition is independent of the choice of a scalar product on spacetime...
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Figure 13.1. Lens-shaped regions.

vectors in \( \mathbb{R}^n \). We always choose the normal to be future-directed, and we normalize it with respect to the Euclidean scalar product on \( \mathbb{R}^{m+1} \), but these are merely conventions.

We shall now explain why and in which sense the solutions of symmetric hyperbolic systems comply with this notion of causality.

**Definition 13.2.2.** Let \( u \) be a smooth solution of the linear symmetric hyperbolic system (13.2.1). A subset \( K \) of the initial value surface \( \{ t = 0 \} \) determines the solution at a spacetime point \( x \in [0, T] \times \mathbb{R}^m \) if every smooth solution of the system which coincides on \( K \) with \( u \), also coincides with \( u \) at \( x \). The **domain of determination** of \( K \) is the set of all spacetime points at which the solution is determined by the initial data on \( K \).

**Definition 13.2.3.** An open subset \( L \subset (0, T) \times \mathbb{R}^m \) is called a **lens-shaped region** if \( L \) is relatively compact in \( \mathbb{R}^n \) and if its boundary \( \partial L \) is contained in the union of two smooth hypersurfaces \( S_0 \) and \( S_1 \) whose intersection with \( L \) is spacelike. We set \( (\partial L)_+ = \partial L \cap S_1 \) and \( (\partial L)_- = \partial L \cap S_0 \), where we adopt the convention that \( (\partial L)_+ \) lies to the future of \( (\partial L)_- \).

Figure 13.1 shows typical examples of lens-shaped regions. Often, one chooses \( S_0 = \{ t = 0 \} \) as the initial data surface. Moreover, it is often convenient to write the hypersurface \( S_1 \) as a graph \( S_1 = \{(t, \vec{x}) \mid t = f(\vec{x})\} \). In this case, \( S_1 \) is the zero set of the function \( \phi(t, \vec{x}) = t - f(\vec{x}) \), and the normal \( \nu \) is the gradient of this function, i.e.

\[
(\nu_j)_{j=0,\ldots,m} = (1, \nabla_1 f, \ldots, \nabla_m f).
\]

We first consider the homogeneous equation

\[
(A^j \partial_j + B) u = 0.
\]  

The idea for analyzing the domain of determination is to multiply this equation by a suitable test function and to integrate over a lens-shaped region. More precisely, we consider the equation

\[
0 = \int_L e^{-Kt} 2 \text{Re} \langle u, (A^j \partial_j + B) u \rangle \, d^n x,
\]  

where \( \langle \cdot, \cdot \rangle \) denotes the canonical scalar product on \( \mathbb{C}^N \), and \( K > 0 \) a positive parameter to be determined later. Since the \( A^j \) are Hermitian, we have

\[
\partial_j \langle u, A^j u \rangle = 2 \text{Re} \langle u, A^j \partial_j u \rangle + \langle u, (\partial_j A^j) u \rangle,
\]

and using this identity in (13.2.4) gives

\[
0 = \int_L e^{-Kt} \left( \partial_j \langle u, A^j u \rangle + \langle u, (B + B^* - (\partial_j A^j)) u \rangle \right) \, d^n x.
\]  

(13.2.6)
In the first term we integrate by parts with the Gauß divergence theorem,
\[ \int_{L} e^{-Kt} \partial_{j} \langle u, A^{j} u \rangle \, d^{n}x = K \int_{L} e^{-Kt} \langle u, A_{0}^{0} u \rangle \, d^{n}x \]
\[ + \int_{(\partial L)_{+}} e^{-Kt} \langle u, \nu_{j} A^{j} u \rangle \, d\mu_{\partial L_{+}} - \int_{(\partial L)_{-}} e^{-Kt} \langle u, \nu_{j} A^{j} u \rangle \, d\mu_{\partial L_{-}} \].
\[ (13.2.7) \]

We now use (13.2.7) in (13.2.6) and solve for the surface integral over \( (\partial L)_{+} \),
\[ \int_{(\partial L)_{+}} e^{-Kt} \langle u, \nu_{j} A^{j} u \rangle \, d\mu_{\partial L_{+}} = \int_{(\partial L)_{-}} e^{-Kt} \langle u, \nu_{j} A^{j} u \rangle \, d\mu_{\partial L_{-}} \]
\[ + \int_{L} e^{-Kt} \langle u, \nu_{j} (A^{j} - K - B - B^{*} + (\partial_{j} A^{j})) u \rangle \, d^{n}x. \]
\[ (13.2.8) \]

This identity is the basis for the following uniqueness results.

**Theorem 13.2.4.** Let \( u_{1} \) and \( u_{2} \) be two smooth solutions of the linear symmetric hyperbolic system \((13.1.3)\) which coincide on the past boundary of a lens-shaped region \( L \),
\[ u_{1} |_{(\partial L)_{-}} = u_{2} |_{(\partial L)_{-}}. \]

Then \( u_{1} \) and \( u_{2} \) coincide in the whole set \( L \).

**Proof.** The function \( u := u_{1} - u_{2} \) is a solution of the homogeneous system \((13.2.3)\) with \( u |_{(\partial L)_{-}} = 0 \). Hence \((13.2.8)\) simplifies to
\[ \int_{(\partial L)_{+}} e^{-Kt} \langle u, \nu_{j} A^{j} u \rangle \, d\mu_{\partial L_{+}} = \int_{L} e^{-Kt} \langle u, \nu_{j} (A^{j} - K - B - B^{*} + (\partial_{j} A^{j})) u \rangle. \]
Assume that \( u \) does not vanish identically in \( L \). By choosing \( K \) sufficiently large, we can then arrange that the right side becomes negative. However, since \( \partial L_{+} \) is a spacelike hypersurface, the left side is non-negative. This is a contradiction. \( \square \)

As an immediate consequence, we obtain the following uniqueness result for solutions of the Cauchy problem.

**Corollary 13.2.5.** Let \( u_{1} \) and \( u_{2} \) be two smooth solutions of the linear symmetric hyperbolic system \((13.1.3)\) with the same initial at time \( t = 0 \). Then \( u_{1} \equiv u_{2} \) in a neighborhood of the initial data surface.

If the matrices \( A^{j} \) are uniformly bounded and \( A_{0} \) is uniformly positive, then \( u_{1} \equiv u_{2} \) in \([0, T] \times \mathbb{R}^{m} \).

**Proof.** The local uniqueness result follows immediately by covering the initial data surface by lens-shaped regions (see the left of Figure 13.2). For the global uniqueness, for any \( x_{0} = (t_{0}, \vec{x}_{0}) \in [0, T] \times \mathbb{R}^{m} \) our task is to choose a lens-shaped region which contains \( x_{0} \) and whose past boundary \( S_{0} \) is contained in the surface \( \{ t = 0 \} \). We need to rule out the situation shown on the right of Figure 13.2 that the hypersurface \( S_{1} \) does not intersect \( S_{0} \), in which case we would not get a relatively compact lens-shaped region. To this end, we must use that the matrices \( A^{j} \) are uniformly bounded by assumption.
As a consequence, there is $\varepsilon > 0$ such that the inequality $\|\nabla f\| \leq \varepsilon$ implies that the hypersurface $S_1 = \{(t = f(\vec{x}), \vec{x})\}$ is spacelike. Possibly after decreasing $\varepsilon$, we may choose
\[
    f(\vec{x}) = t_0 + \varepsilon \left(1 - \sqrt{1 + \|\vec{x} - \vec{x}_0\|^2}\right).
\]
This concludes the proof.

By a suitable choice of lens-shaped region one can get an upper bound for the maximal propagation speed. For the Dirac equation, where the causal structure of Definition 13.2.1 coincides in view of (13.2.2) with that of Minkowski space, one can choose for $S_1$ a family of spacelike hypersurfaces which converge to the boundary of a light cone (see Figure 13.3). This shows that the maximal propagation speed for Dirac waves is indeed the speed of light (which, according to our conventions, is equal to one).

### 13.3. Global Existence of Smooth Solutions

In this section we will show that, by refining the above uniqueness argument, we even obtain an existence proof. The close connection between existence and uniqueness for linear equations is a familiar theme in mathematics. The simplest setting where it occurs is in the study of the linear equation $Au = v$ with a given vector $v \in \mathbb{R}^n$ and a quadratic matrix $A$. In this case, the uniqueness of the solution implies that the matrix $A$ is invertible, which in turn ensures existence. A more interesting example is Fredholm’s alternative for compact operators (see for example [121], Section VI.5). The procedure for globally hyperbolic systems follows somewhat similar ideas. Here the general strategy is to construct a bounded linear functional on a Hilbert space in such a way that the Fréchet-Riesz theorem (see Theorem 2.2.4) gives the desired solution.

Before beginning, we point out that, in view of uniqueness and finite propagation speed, it suffices to consider the problem in a bounded spatial region. Indeed, once we have constructed “local solutions” in small lens-shaped regions as shown on the left of Figure 13.2, uniqueness implies that these solutions agree in the overlap of the lens-shaped regions, making it possible to “glue them together” to obtain the desired solution which is global in space. We will come back to this construction in more detail in the context of the Dirac equation in Sections 13.4 and 13.6 (see also Figure 13.4). Having this construction in mind, we may start from a local problem and to extend the coefficients of the symmetric hyperbolic system in an arbitrary way outside. Therefore, it is no loss of generality to consider a problem in the whole space $\mathbb{R}^m$. Choosing a bounded time interval $t \in [0, T]$ (where $t = 0$ is the initial time), we are led to considering the time strip
\[
    R_T := [0, T] \times \mathbb{R}^m.
\]
We now write the linear system (13.2.1) as
\[
    Lu = w \quad \text{with} \quad L := A^j \partial_j + B, \quad (13.3.1)
\]
where we again sum over $j = 0, \ldots, m$. Again using that the system can be extended arbitrarily outside a bounded spatial region, we may assume that that the functions $A^j, B$ and $w$ are uniformly bounded in $R_T$ and that $w$ has spatially compact support (meaning that $w(t, \cdot) \in C^\infty_c(\mathbb{R}^m)$ for all $t \in [0, T]$). Moreover, for convenience we again assume smoothness of $A^j, B$ and $w$. In the \emph{Cauchy problem} one seeks for a solution of the equation \( (13.3.1) \) with prescribed initial data $u_0 \in C^\infty(\mathbb{R}^m)$ at time $t = 0$, 

\[
Lu = w, \quad u|_{t=0} = u_0 \in C^0_0(\mathbb{R}^m)
\]  

in $C^s(R_T)$. First of all, we may restrict attention to the case $u_0 \equiv 0$, 

\[
Lu = w, \quad u|_{t=0} = 0.
\]  

In order to see this, let $u$ be a solution of the above Cauchy problem. Choosing a function $v \in C^\infty(R_T)$ which vanish both at times $t = 0$ and $t = \lambda$, respectively. As the space of test functions we choose the space of test functions we choose a test function $v$ with spatially compact support by $C^\infty(\mathbb{R}^m)$ and vanishes at $t = 0$ if conversely $\bar{u}$ is a solution of the corresponding Cauchy problem with zero initial data, then $u := \bar{u} + v$ is a solution of the original problem \( (13.3.2) \).

In preparation of the existence proof, we need to introduce the notion of a weak solution. In order to get into the weak formulation, we multiply the equation \( (13.3.1) \) by a test function $v(t, \vec{x})$ and integrate over $R_T$, giving rise to the equation

\[
\langle v, Lu \rangle_{L^2(R_T)} = \langle v, w \rangle_{L^2(R_T)}
\]

with the $L^2$-scalar product defined by

\[
\langle v, v' \rangle_{L^2(R_T)} := \int_0^T d \int_{\mathbb{R}^m} \langle v(t, \vec{x}), v'(t, \vec{x}) \rangle \, d^m x.
\]

The next step is to integrate by parts, so that the derivatives act on the test function $v$. Before doing so, we need to specify the regularity of the test functions. To this end, for $\lambda \in [0, T]$ we consider the time strip

\[
R_\lambda := [0, \lambda] \times \mathbb{R}^m.
\]

We denote the $s$-times continuously differential functions on $R_\lambda$ with spatially compact support by $C^s(R_\lambda)$. The function spaces

\[
\underline{C^s(R_\lambda)} \quad \text{and} \quad \overline{C^s(R_\lambda)}
\]

are defined as the functions which in addition vanish at $t = 0$ and $t = \lambda$, respectively. As the space of test functions we choose $C^1(R_T)$; this guarantees that integrating by parts does not yield boundary terms at $t = T$. For a classical solution $u \in C^1(R_T)$ (i.e. a solution with zero Cauchy data \( (13.3.3) \)), also the boundary term at $t = 0$ vanishes. We thus obtain

\[
\langle v, w \rangle_{L^2(R_T)} = \langle \bar{L}v, u \rangle_{L^2(R_T)} \quad \text{for all } v \in C^1(R_T),
\]

where $\bar{L}$ is the formal adjoint of $L$ with respect to the scalar product \( (14.2.9) \), i.e.

\[
\bar{L} := \bar{A}^j \partial_j + \bar{B} \quad \text{with} \quad \bar{A}^j = -A^j \quad \text{and} \quad \bar{B} = B^+ - (\partial_j A^j).
\]

Now suppose that a function $u \in C^1(R_T)$ satisfies \( (13.3.5) \). Testing with functions $v \in \overline{C^1(R_T)} \cap C^1(R_T)$ which vanish both at times $t = 0$ and $t = T$, we can integrate by parts without boundary terms. Using a standard denseness argument, one finds that $u$ solves
the symmetric hyperbolic system (13.3.1). Next, testing with a function \( v \in C^1(R_T) \) which does not vanish at \( t = 0 \), only the boundary remains, giving the equation
\[
\int_{\mathbb{R}^m} \langle v(t, \vec{x}), u(t, \vec{x}) \rangle \, d^m x = 0 \quad \text{for all } v \in C^1(R_T),
\]
which in turn implies that \( u \) vanishes initially. Thus \( u \) is a solution of the Cauchy problem (13.3.3). To summarize, for functions \( u \in C^1(R_T) \), the weak formulation (13.3.5) is equivalent to our Cauchy problem (13.3.1) and (13.3.3). Therefore, it is sensible to take (13.3.5) as the definition of a weak solution of the Cauchy problem. The main advantage of the weak formulation (13.3.5) is that it is well-defined even for functions which are not differentiable.

Our next step is to derive so-called energy estimates for a given solution \( u \in C^1(R_T) \). To this end, we return to the formula for the divergence (13.2.5) and using the equation (13.3.1), we obtain
\[
\partial_j \langle u, A_j u \rangle + \langle u, Cu \rangle = 2 \text{Re} \langle u, w \rangle - \langle u, Cu \rangle,
\]
where \( C := B + B^* - (\partial_j A^j) \).

Next, we integrate (13.3.7) over \( R_\lambda \), integrate by parts and use that the initial values at \( t = 0 \) vanish. We thus obtain
\[
E(\lambda) := \int_{t=\lambda}^\lambda \langle u, A^0 u \rangle \, d^m x = \int_0^\lambda d\tau \int_{R_\tau} \left( 2 \text{Re} \langle u, w \rangle - \langle u, Cu \rangle \right) \, d^m x.
\]
Since the matrix \( C \) is uniformly bounded and \( A_0 \) is uniformly positive, there is a constant \( K > 1 \) such that
\[
|\langle u, Cu \rangle| \leq K |\langle u, A^0 u \rangle|.
\]
Moreover, the linear term in \( u \) can be estimated with the Schwarz inequality by
\[
2 \text{Re} \langle u, w \rangle \leq \mu |\langle u, w \rangle| + \frac{1}{\mu} \langle w, w \rangle \leq \langle u, A^0 w \rangle + \frac{1}{\mu} \langle w, A^0 w \rangle
\]
with a suitable constant \( \mu > 0 \). Applying these estimates in (13.3.9) gives
\[
E(\lambda) \leq (K + 1) \int_0^\lambda E(t) \, dt + \frac{1}{\mu^2} \int_{R_\lambda} \langle w, A^0 w \rangle \, d^m x.
\]
Writing this inequality as
\[
\frac{d}{d\lambda} e^{-(K+1)\lambda} \int_0^\lambda E(t) \, dt \leq e^{-(K+1)\lambda} \frac{1}{\mu^2} \int_{R_T} \langle w, A^0 w \rangle \, d^m x,
\]
we can integrate over \( \lambda \) to obtain
\[
\int_0^T E(\lambda) \, d\lambda \leq \frac{e^{(K+1)T} - 1}{K + 1} \frac{1}{\mu^2} \int_{R_T} \langle w, A^0 w \rangle \, d^m x.
\]
Finally, we apply the mean value theorem and use that the exponential function is monotone to conclude that
\[
\int_0^T E(\lambda) \, d\lambda \leq \frac{T}{\mu^2} e^{(K+1)T} \int_{R_T} \langle w, A^0 w \rangle \, d^m x.
\]
(13.3.10)
This is the desired energy estimate. Before going on, we point out that the notion of “energy” used for the quantity \( E(\lambda) \) does in general not coincide with the physical energy.
13.3. GLOBAL EXISTENCE OF SMOOTH SOLUTIONS

In fact, for the Dirac equation (13.1.2), \( E(\lambda) \) has the interpretation as the electric charge. Following Example 13.1.2, for the scalar wave equation \( \Box \phi = 0 \), we find

\[
E(\lambda) = \int_{\mathbb{R}^m} \left( |\partial_t \phi|^2 + |\nabla \phi|^2 + |\phi|^2 \right) d^m x.
\] (13.3.11)

This differs from the physical energy by the last summand \( |\phi|^2 \) (and an overall factor of two). The name “energy” for \( E(\lambda) \) was motivated by the fact, considering only the highest derivative terms, the expression (13.3.11) is indeed the physical energy. We point out that, in contrast to the physical energy, the quantity \( E(\lambda) \) does in general depend on time. The point is that (13.3.10) gives an a-priori control of the energy in terms of the inhomogeneity. The exponential factor in (13.3.10) can be understood in analogy to a Grönwall estimate.

For the following construction, it is convenient to introduce on \( C^1(R_T) \) the scalar product

\[
(u, v) = \int_{R_T} \langle u, A^0 v \rangle \, d^m x.
\] (13.3.12)

We denote the corresponding norm by \( \| \cdot \| \). Setting furthermore

\[
\Gamma^2 = \frac{T}{\mu^2} e^{(K+1)T},
\]

the energy estimate can be written in the compact form

\[
(u, u) \leq \Gamma^2 (w, w).
\] (13.3.13)

This inequality holds for every solution \( u \) of the differential equation \( Lu = w \) which vanishes at \( t = 0 \). Noting that every function \( u \in C^1(R_T) \) is a solution of this differential equation with inhomogeneity \( w := L u \), we obtain

\[
\| u \| \leq \Gamma \| Lu \| \quad \text{for all } u \in C^1(R_T).
\] (13.3.13)

This is the form of the energy estimates suitable for an abstract existence proof.

Note that the operator \( \tilde{L} \) in (13.3.6) is also symmetric hyperbolic and has the same boundedness and positivity properties as \( L \). Hence, repeating the above arguments, we obtain similar to (13.3.13) the “dual estimate”

\[
\| v \| \leq \tilde{\Gamma} \| \tilde{L} v \| \quad \text{for all } v \in C^1(R_T).
\] (13.3.14)

We now want to show the existence of weak solutions with the help of the Fréchet-Riesz theorem (for basics on functional analysis see [121] or [110]). To this end, we first introduce on \( C^1(R_T) \) yet another scalar product denoted by

\[
\langle v, v' \rangle = (\tilde{L} v, \tilde{L} v).
\] (13.3.15)

This scalar product is indeed positive definite, because for any \( v \neq 0 \),

\[
\langle v, v \rangle = (\tilde{L} v, \tilde{L} v) \geq \tilde{\Gamma}^{-2} (v, v) \neq 0,
\]

where in the last step we applied (13.3.14). Forming the completion, we obtain the Hilbert space \( \mathcal{H}, (\cdot, \cdot) \). We denote the corresponding norm by \( \| \cdot \| \). In view of (13.3.14) and (13.3.15), we know that every vector \( v \in \mathcal{H} \) is a function in \( L^2(R_T, d^m x) \). Moreover, we know from (13.3.15) that \( \tilde{L} v \) is also in \( L^2(R_T, d^m x) \). We remark that, in the language of functional analysis, the space \( \mathcal{H} \) can be identified with the Sobolev space \( W^{1,2}(R_T) \), but we do not need this here.
We now consider for \( w \in C^0(\mathbb{R}^T) \) and \( v \in C^1(\mathbb{R}^T) \) the linear functional \( \langle v, w \rangle_{L^2(\mathbb{R}^T)} \). In view of the estimate
\[
|\langle v, w \rangle_{L^2(\mathbb{R}^T)}| \leq \|v\|_{L^2(\mathbb{R}^T)} \|w\|_{L^2(\mathbb{R}^T)} \leq \frac{\Gamma}{C} \|v\|_{L^2(\mathbb{R}^T)} \|w\|_{L^2(\mathbb{R}^T)},
\]
this functional is continuous in \( v \in H \). The Fréchet-Riesz theorem shows that there is \( U \in H \) with
\[
\langle v, w \rangle_{L^2(\mathbb{R}^T)} = \langle v, U \rangle = (\tilde{L}v, \tilde{U}) \quad \text{for all } v \in \mathcal{H}.
\]
Rewriting the last scalar product as \( (\tilde{L}v, \tilde{U}) = \langle Lv, A^0 \tilde{U} \rangle_{L^2(\mathbb{R}^T)} \), one sees that the function \( u := A^0 \tilde{U} \in L^2(\mathbb{R}^T, d^n x) \) satisfies the equation \( (13.3.5) \) and is thus the desired weak solution. Note that all our methods apply for arbitrarily large \( T \).

We have thus proved the global existence of weak solutions. We next want to show that the solutions are smooth. Thus our task is to show that our constructed weak solution \( u \) is of the class \( C^s(\mathbb{R}_\lambda) \), where \( s \geq 1 \) can be chosen arbitrarily large. We first show that a linear symmetric hyperbolic system can be “enlarged” to include the partial derivatives of \( \phi \).

**Lemma 13.3.1.** Suppose that the system \( A^j \partial_j u + Bu = w \) is symmetric hyperbolic. Then there is a symmetric hyperbolic system of the form
\[
\tilde{A}^j \partial_j \Psi + \tilde{B} \Psi = \tilde{w}
\]
for the vector \( \Psi := (\partial_t u, \nabla_1 u, \ldots, \nabla_m u, u) \in \mathbb{C}^{(n+1)N} \).

**Proof.** Let \( i \) be a fixed spacetime index. We differentiate the equation \( Lu = w \),
\[
\partial_i w = \partial_i Lu = L \partial_i u + (\partial_i A^j) \partial_j u + (\partial_i B) u.
\]
This equation can be written as
\[
A^j \partial_j \Psi_i + \sum_{j=1}^n \tilde{B}^j_i \Psi_j + (\partial_i B) u = \tilde{w}_i,
\]
where we set
\[
\tilde{B}^j_i = B \delta^j_i + (\partial_i A^j) \quad \text{and} \quad \tilde{w}_i = \partial_i w.
\]
Combining these equations with the equation \( Lu = w \), we obtain a system of the form \( (13.3.16) \), where the matrices \( \tilde{A}^j \) are block diagonal in the sense that
\[
\tilde{A}^j = ((\tilde{A}^j)_{\alpha,\beta})_{\alpha,\beta=0,\ldots,m+1} \quad \text{with} \quad (\tilde{A}^j)_{\alpha,\beta} = A^j \delta^\alpha_\beta.
\]
Obviously, this system is again symmetric hyperbolic. \( \square \)

Iterating this lemma, we obtain (at least in principle) a symmetric hyperbolic system for \( u \) and all its partial derivatives up to any given order \( s \). Since the corresponding weak solution is in \( L^2(\mathbb{R}_T) \), we conclude that \( u \) and all its weak partial derivatives are square integrable. The next lemma, which is a special case of the general Sobolev embedding theorems (see for example \[31\] Section II.5. or \[132\] Section 4), gives smoothness of the solution.
Lemma 13.3.2. Let \( s > \frac{m}{2} \) be an integer. If a function \( g \) on \( \mathbb{R}^m \) is \( s \) times weakly differentiable and
\[
\int_{\mathbb{R}^m} |\nabla^\alpha g|^2 \, d^m x < C \tag{13.3.17}
\]
for all multi-indices \( \alpha \) with \(|\alpha| \leq s\), then \( g \) is bounded, \( g \in L^\infty(\mathbb{R}^m) \). Likewise, if \( g \) is \( s + l + 1 \) times weakly differentiable with \( l \geq 1 \) and \( |\alpha| \leq s + l + 1 \), then \( g \in \mathcal{C}^l(\mathbb{R}^m) \).

**Proof.** We apply the Schwarz inequality to the Fourier transform,
\[
|g(x)|^2 = \left| \int_{\mathbb{R}^m} \frac{d^m k}{(2\pi)^m} \hat{g}(k) e^{-ikx} \right|^2 \\
= \left| \int_{\mathbb{R}^m} \frac{d^m k}{(2\pi)^m} (1 + |k|^2)^{-\frac{s}{2}} (1 + |k|^2)^{\frac{s}{2}} \hat{g}(k) e^{-ikx} \right|^2 \\
\leq c_m \int_{\mathbb{R}^m} \frac{d^m k}{(2\pi)^m} (1 + |k|^2)^s |\hat{g}(k)|^2,
\]
where the constant \( c_m \) is finite due to our choice of \( s \),
\[
c_m = \int_{\mathbb{R}^m} \frac{d^m k}{(2\pi)^m} (1 + |k|^2)^{-s} < \infty.
\]
Using the Plancherel formula together with the fact that a factor \( k^2 \) corresponds to a Laplacian in position space, we obtain
\[
\int_{\mathbb{R}^m} \frac{d^m k}{(2\pi)^m} (1 + |k|^2)^s |\hat{g}(k)|^2 = \sum_{\ell=0}^{s} \binom{n}{\ell} \|\nabla^\ell g\|^2_{L^2(\mathbb{R}^m)} < c.
\]
Hence \( \sqrt{c_m}c \) is an \( L^\infty \)-bound for \( g \).

Next, if \( g \) is \( s + l + 1 \) times weakly differentiable, then \( \|Dg\|_{L^\infty(\mathbb{R}^m)} < c \). As a consequence, the mean value theorem yields \(|g(x) - g(y)| \leq c|x - y|\), so that \( g \) is Lipschitz continuous. Finally, if \( g \) is \( s + l + 1 \) times weakly differentiable, then all partial derivatives \( \nabla^\alpha g \) of order \(|\alpha| \leq l\) are Lipschitz continuous, so that \( g \in \mathcal{C}^l(\mathbb{R}^m) \). \( \square \)

More precisely, in order to apply this lemma, we fix a time \( t \) and consider the solution \( u(\lambda,.) \). The identity \( \ref{1339} \) implies that \( E(\lambda) \) is controlled in terms of \( \|w\| \) and \( \|u\| \). After iteratively applying Lemma \( \ref{1331} \), we conclude that the weak derivatives of \( u(\lambda,.) \) exist to any order and are in \( L^2(\mathbb{R}^m) \). It follows that \( u(\lambda,.) \) is smooth. Finally, one uses the equation to conclude that \( u \) is also smooth in the time variable.

The results of this section can be summarized as follows.

**Theorem 13.3.3.** Consider the Cauchy problem
\[
(A^0 \partial_t + \sum_{\alpha=1}^{m} A^\alpha \nabla_\alpha + B) u = w \in \mathcal{C}^0_{\infty}([0, T] \times \mathbb{R}^m), \quad u|_{t=0} = u_0 \in \mathcal{C}^0_{\infty}(\mathbb{R}^m).
\]
Assume that the matrices \( A^0, A^\lambda \) and \( B \) as well as the functions \( w \) and \( u_0 \) are smooth. Moreover, assume that all these functions as well as all their partial derivatives are uniformly bounded (where the bound may depend on the order of the derivatives). Then the Cauchy problem has a smooth solution on \([0, T] \times \mathbb{R}^m \).

This theorem also applies in the case \( T = \infty \), giving global existence of a smooth solution.

We finally show that the solutions depend smoothly an parameters.
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**Corollary 13.3.4.** Suppose that the matrices $A^j, B$ and the functions $w, u_0$ depend smoothly on a parameter $\lambda$. Then the family of solutions $u(\lambda)$ is also smooth in $\lambda$.

**Proof.** First, similar as explained after (13.3.3), we may restrict attention to the case $u_0 = 0$. Differentiating the equation $Lu = w$ with respect to $\lambda$, we obtain

$$Lu_\lambda = (\partial_\lambda L)u + \partial_\lambda w =: \tilde{w},$$

where $u_\lambda$ stands for the formal derivative $\partial_\lambda u$. This is a symmetric hyperbolic system for $u_\lambda$. According to Theorem 13.3.3, we know that $u$ and therefore $\tilde{w}$ are smooth. Considering the limit of the difference quotients, one verifies that $u_\lambda$ really coincides with $\partial_\lambda u(\lambda)$ for our given family of solutions $u(\lambda)$. The higher $\lambda$-derivatives can be treated inductively. $\square$

13.4. The Causal Dirac Green’s Operators in Minkowski Space

We now want to apply the previous general existence and uniqueness results to the Cauchy problem (13.1.1) for the Dirac equation in Minkowski space in the presence of an external potential $B$.

**Theorem 13.4.1.** Consider the Cauchy problem for the Dirac equation (13.1.1) for smooth initial data $\psi_0$, a smooth inhomogeneity $\phi$ and a smooth matrix-valued potential $B \in C^\infty(\mathcal{M}, C^{4 \times 4})$. Then there is a unique global smooth solution $\psi \in C^\infty(\mathcal{M}, S\mathcal{M})$.

**Proof.** Writing the Dirac equation in the Hamiltonian form (13.1.2), we obtain a symmetric hyperbolic system. In view of the uniqueness result for smooth solutions of Corollary 13.2.5, it suffices to construct a smooth solution at any given time $T \in \mathbb{R}$. It suffices to consider the case $T > t_0$, because otherwise we reverse the time direction. Moreover, we can arrange by a time shift that $t_0 = 0$.

We cannot apply Theorem 13.3.3 directly because the coefficient functions in (13.1.2) do not need to be bounded, nor are our initial values compactly supported. For this reason, we need to construct local solutions and “glue them together” using linearity: We first extend the initial data $\psi_0$ smoothly to the time strip $\mathbb{R}T$ and consider the Cauchy problem for $\tilde{\psi} := \psi - \psi_0$,

$$(i\partial + B - m)\tilde{\psi} = \tilde{\phi} \in C^\infty(\mathcal{M}, S\mathcal{M}), \quad \tilde{\psi}|_{t_0} = 0.$$ (13.4.1)

We let $(\eta_k)_{k \in \mathbb{N}}$ be a smooth partition of unity of $\mathbb{R}^m$ with $\eta_k \in C^\infty_0(\mathbb{R}^m)$ (for details see for example [126, Theorem 2.13]). We extend these functions to static functions on $R_T$ (i.e. $\eta_k(t, \vec{x}) := \eta_k(\vec{x})$). Given $k \in \mathbb{N}$, we first solve the Cauchy problem for the inhomogeneity $\eta_k \tilde{\phi}$. We choose a compact set $K \subset \mathbb{R}^m$ such that $[0, T] \times K$ contains the causal future of the support of $(\eta_k \tilde{\phi})$ (see Figure 13.4 more specifically, we could choose $K = B_{2T}(\text{supp} \eta_k)$). Next, we choose a smooth, compactly supported function $\theta \in C^\infty_0(\mathbb{R}^m)$ with $\theta|_K \equiv 1$. We again extend $\theta$ to a static function on $R_T$. 

![Figure 13.4. Construction of local solutions.](image)
We now consider the modified Cauchy problem

\[
\left(1_{C^1} \partial_t \psi + \gamma^0 \gamma^\mu + \theta (- i \gamma^0 (B - m)) \right) \tilde{\psi}_k = -i \gamma^0 \eta_k \tilde{\phi}, \quad \tilde{\psi}_k|_{t_0} = 0.
\]

Now the coefficients in the PDE are uniformly bounded, and the inhomogeneity has compact support. Therefore, we can apply Theorem 13.3.3 to obtain a global smooth solution. Due to finite propagation speed (see Theorem 13.2.4, where we choose lens-shaped regions \(L\) as shown in Figure 13.4), this solution vanishes outside \(K\). Therefore, it is also a solution of the unmodified Dirac equation, with initial data \(\eta_k \tilde{\phi}\).

Finally, summing over \(k\) gives the desired solution of the original Cauchy problem,

\[
\psi := \sum_{k=1}^{\infty} \tilde{\psi}_k.
\]

Here the series converges because, again due to finite propagation speed, it is locally finite.

We next explain how the previous existence and uniqueness results give rise to the existence of causal Green’s functions, being defined as distributions. Our main tool is the Schwartz kernel theorem. We do not give a proof of this more advanced result of distribution theory but refer instead to [100, Section 5.2] or [132, Section 4.6]. For better consistency with the notation in the perturbative treatment in Section 18, from now on we denote the objects in the presence of an external potential with an additional tilde. We begin with a representation formula for the solution of the Cauchy problem in terms of a distribution.

**Theorem 13.4.2.** Assume that the external potential \(B\) is smooth and that \(B\) and all its partial derivatives are uniformly bounded in Minkowski space. Then for any \(t, t_0\) there is a unique distribution \(\tilde{k}_m(t, .; t_0, .) \in D'(\mathbb{R}^3 \times \mathbb{R}^3)\) such that the solution of the Cauchy problem (17.0.1) has the representation

\[
\psi(t, \vec{x}) = 2\pi \int_N \tilde{k}_m(t, \vec{x}; t_0, \vec{y}) \gamma^0 \psi_0(\vec{y}) \, d^3y.
\]

The integral kernel \(k_m\) is also a distribution in spacetime, \(k_m \in D'(M \times M)\) It is a distributional solution of the Dirac equation,

\[
(i \partial_\mu + B - m) \tilde{k}_m(x, y) = 0.
\]

**Proof.** The energy estimates combined with the Sobolev embedding of Lemma [13.3.2] showed that there is \(k \in \mathbb{N}\) and a constant \(C = C(t, t_0, \vec{x}, B)\) such that the solution \(\psi(t, .)\) of the Cauchy problem is bounded in terms of the initial data by

\[
|\psi(t, \vec{x})| \leq C |\psi_0|_{C^k},
\]

where \(|\psi|^2 := \langle \psi | \gamma^0 \psi \rangle\), and the \(C^k\)-norm is defined by

\[
|\psi_0|_{C^k} = \max_{|\beta| \leq k} \sup_{\vec{x} \in \mathbb{R}^3} |\nabla^\beta \psi_0(\vec{x})|.
\]

Moreover, this estimate is locally uniform in \(\vec{x}\), meaning that for any compact set \(K \subset \mathbb{R}^3\), there is a constant \(C\) such that (13.4.4) holds for all \(\vec{x} \in K\). This makes it possible to apply the the Schwartz kernel theorem [100, Theorem 5.2.1], showing that \(\tilde{k}_m(t, .; t_0, .) \in D'(\mathbb{R}^3 \times \mathbb{R}^3)\).

Next, we note that the constant \(C\) in (13.4.4) can also be chosen locally uniformly in \(t\) and \(t_0\). Thus, after evaluating weakly in \(t\) and \(t_0\), we may again apply the Schwartz kernel
Theorem 13.4.3. Assume that the external potential \( B \) is smooth and that \( B \) and all its partial derivatives are uniformly bounded in Minkowski space. Then there are unique distributions
\[
\tilde{s}_m^\wedge, \tilde{s}_m^\vee \in \mathcal{D}'(M \times M)
\]
which satisfy the distributional equations
\[
(i\partial_x + B - m) \tilde{s}_m(x, y) = \delta^4(x - y) \quad (13.4.6)
\]
and are supported in the upper respectively lower light cone,
\[
\text{supp} \tilde{s}_m^\wedge(x, .) \subset J^\wedge_x, \quad \text{supp} \tilde{s}_m^\vee(x, .) \subset J^\vee_x. \quad (13.4.7)
\]

Proof. It is clear by construction and the fact that the constant \( C \) in (13.4.4) can be chosen locally uniformly in \( x \) and \( y \) that the causal Green’s functions are well-defined distributions in \( \mathcal{D}'(M \times M) \). The support property (13.4.7) follows immediately from finite propagation speed as explained at the end of Section 13.2. The uniqueness of the Green’s functions is clear from the uniqueness of solutions of the Cauchy problem. In order to derive the distributional equations (13.4.6), we only consider the retarded Green’s function (the argument for the advanced Green’s function is analogous). Then, according to (13.4.2) and (13.4.5),
\[
\Theta(t - t_0) \psi(t, \bar{x}) = i \int_N \tilde{s}_m^\wedge(t, \bar{x}; t_0, \bar{y}) \gamma^0 \psi_0(\bar{y}) \, d^3y,
\]
where \( \psi \) is the solution of the corresponding Cauchy problem. Applying the Dirac operator in the distributional sense yields
\[
i\gamma^0 \delta(t - t_0) \psi_0(t, \bar{x}) = i(D_x - m) \int_N \tilde{s}_m^\wedge(t, \bar{x}; t_0, \bar{y}) \gamma^0 \psi_0(\bar{y}) \, d^3y.
\]
We now choose the initial values as the restriction of a test function in spacetime, \( \psi_0 = \phi|_{t=t_0} \) with \( \phi \in C_0^\infty(M, SM) \). Then we can integrate over \( t_0 \) to obtain
\[
i\gamma^0 \phi(x) = (D_x - m) \int_M \tilde{s}_m^\wedge(x, y) i\gamma^0 \phi(y) \, d^4y.
\]
This gives the result. \( \square \)
We remark that, turning the above argument around, we can also use the causal Green’s functions in order to define the causal fundamental solution as a bi-distribution in spacetime,
\[ \tilde{k}_m := \frac{1}{2\pi i} (\tilde{s}_m^\gamma - \tilde{s}_m^\delta) \in \mathcal{D}'(\mathcal{M} \times \bar{\mathcal{M}}). \tag{13.4.8} \]

We conclude this section by collecting a few useful properties of the causal fundamental solutions, which will be needed later in this book. First, the causal fundamental solution has the remarkable property that it relates the scalar product with the spacetime inner product. The following proposition goes back to John Dimock (see [27, Proposition 2.2]).

**Proposition 13.4.4.** For any \( \psi_m \in \mathcal{H}_m \) and \( \phi \in C_0^\infty(\mathcal{M}, S\mathcal{M}) \),
\[ (\psi_m | \tilde{k}_m \phi)_m = \langle \psi_m | \phi \rangle. \tag{13.4.9} \]

**Proof.** We first give the proof under the additional assumption that \( \psi_m \in C_0^\infty(\mathcal{M}, S\mathcal{M}) \). We choose Cauchy surfaces \( \mathcal{N}_+ \) and \( \mathcal{N}_- \) lying in the future and past of \( \text{supp} \phi \), respectively. Let \( \Omega \) be the spacetime region between these two Cauchy surfaces, i.e. \( \partial \Omega = \mathcal{N}_+ \cup \mathcal{N}_- \). Then, according to (18.2.10),
\[ (\psi_m | \tilde{k}_m \phi)_m = (\psi_m | \tilde{k}_m \phi)_{N_+} = \frac{i}{2\pi} (\psi_m | \tilde{s}_m^\gamma \phi)_{N_+} \]
\[ = \frac{i}{2\pi} \left[ (\psi_m | \tilde{s}_m^\gamma \phi)_{N_+} - (\psi_m | \tilde{s}_m^\gamma \phi)_{N_-} \right] \]
\[ = i \int \nabla_\gamma \langle \psi_m | \gamma^j \tilde{s}_m^j \phi \rangle_x d\mu(x), \]
where in the last line we applied the Gauß divergence theorem and used (15.1.2). Using that \( \psi_m \) satisfies the Dirac equation, a calculation similar to (13.4.10) yields
\[ (\psi_m | \tilde{k}_m \phi)_m = \int_\Omega \langle \psi_m | (\mathcal{D} - m) \tilde{s}_m^\gamma \phi \rangle_x d\mu(x) \overset{16.2.1}{=} \int_\Omega \langle \psi_m | \phi \rangle_x d\mu(x). \]

As \( \phi \) is supported in \( \Omega \), we can extend the last integration to all of \( \mathcal{M} \), giving the result.

In order to extend the result to general \( \psi_m \in \mathcal{H}_m \), we use the following approximation argument. Let \( \psi_m^n \in \mathcal{H}_m \cap C_0^\infty(\mathcal{M}, S\mathcal{M}) \) be a sequence which converges in \( \mathcal{H}_m \) to \( \psi_m \). Then obviously \( (\psi_m^n | \tilde{k}_m \phi)_m \to (\psi_m | \tilde{k}_m \phi)_m \). In order to show that the right side of (13.4.9) also converges, it suffices to prove that \( \psi_m^n \) converges in \( L^2_{\text{loc}}(\mathcal{M}, S\mathcal{M}) \) to \( \psi_m \). Thus let \( K \subset \mathcal{M} \) be a compact set contained in the domain of a chart \( (x, U) \). Using Fubini’s theorem, we obtain for any \( \psi \in \mathcal{H}_m \cap C_0^\infty(\mathcal{M}, S\mathcal{M}) \) the estimate
\[ \int_K \langle \psi | \psi \rangle d\mu_{\mathcal{M}} = \int dx_0 \int \langle \psi | \psi \rangle \sqrt{|g|} d^3x \leq C(K) \langle \psi | \psi \rangle_m. \]

Applying this estimate to the functions \( \psi = \psi_m^n - \psi_m^{n'} \), we see that \( \psi_m^n \) converges in \( L^2(K, S\mathcal{M}) \) to a function \( \tilde{\psi} \). This implies that \( \psi_m^n \) converges to \( \tilde{\psi} \) pointwise almost everywhere (with respect to the measure \( d\mu_{\mathcal{M}} \)). Moreover, the convergence of \( \psi_m^{n} \) in \( \mathcal{H}_m \) to \( \psi_m \) implies that the restriction of \( \psi_m^n \) to any Cauchy surface \( \mathcal{N} \) converges to \( \psi_m |_{\mathcal{N}} \) pointwise almost everywhere (with respect to the measure \( d\mu_{\mathcal{N}} \)). It follows that \( \tilde{\psi} = \psi_m |_K \), concluding the proof. \( \square \)
Corollary 13.4.5. The operator \( \tilde{k}_m \), (18.2.10), is symmetric with respect to the inner product (15.1.3).

Proof. Using Proposition 13.4.4 we obtain for all \( \phi, \psi \in C_0^\infty(\mathcal{M}, S\mathcal{M}) \),
\[
<\tilde{k}_m \phi | \psi> = (\tilde{k}_m \phi | \tilde{k}_m \psi)_m = <\phi | \tilde{k}_m \psi> ,
\]
concluding the proof. \( \square \)

13.5. A Polynomial Estimate in Time

We now derive an estimate which shows that the solutions of the Dirac equation increase at most polynomially in time. This result will be needed in Section 17.2.1. For the proof we adapt standard methods of the theory of partial differential equations to the Dirac equation. In generalization of (16.3.4), we denote the spatial Sobolev norms by
\[
\| \phi \|_{W^{a,2}} = \sum_{\alpha \text{ with } |\alpha| \leq a} \int_{\mathbb{R}^3} |\nabla^\alpha \phi(\vec{x})|^2 \, d^3x .
\]

Lemma 13.5.1. We are given two non-negative integers \( a \) and \( b \) as well as a smooth time-dependent potential \( B \). In the case \( a > 0 \) and \( b \geq 0 \), we assume furthermore that the spatial derivatives of \( B \) decay faster than linearly for large times in the sense that
\[
|\nabla B(t)|_{C^{a-1}} \leq \frac{c}{1+|t|^{1+\varepsilon}} \quad (13.5.1)
\]
for suitable constants \( c, \varepsilon > 0 \). Then there is a constant \( C = C(c, \varepsilon, a, b) \) such that every family of solutions \( \psi \in H^\infty \) of the Dirac equation (13.14) for varying mass parameter can be estimated for all times in terms of the boundary values at \( t = 0 \) by
\[
\| \partial_t^b \psi_m |_t \|_{W^{a,2}} \leq C \left( 1 + |t|^b \right) \sum_{p=0}^b \| \partial_t^p \psi_m |_{t=0} \|_{W^{a,2}} .
\]

Proof. We choose a multi-index \( \alpha \) of length \( a := |\alpha| \) and a non-negative integer \( b \). Differentiating the Dirac equation (1.3.14) with respect to the mass parameter and to the spatial variables gives
\[
(i\partial_\gamma + B - m) \nabla^\alpha \partial_m \psi_m = b \nabla^\alpha \partial_m^{-1} \psi_m - \nabla^\alpha (B \partial_m \psi_m) + B \nabla^\alpha \partial_m \psi_m .
\]
Introducing the abbreviations
\[
\Xi := \nabla^\alpha \partial_m \psi_m \quad \text{and} \quad \phi := b \nabla^\alpha \partial_m^{-1} \psi_m - \nabla^\alpha (B \partial_m \psi_m) + B \nabla^\alpha \partial_m \psi_m ,
\]
we rewrite this equation as the inhomogeneous Dirac equation
\[
(D - m) \Xi = \phi .
\]
A calculation similar to current conservation yields
\[
-i \partial_\gamma <\Xi | \gamma \Xi> = <(D - m) \Xi | \Xi> - <\Xi | (D - m) \Xi> = <\phi | \Xi> - <\Xi | \phi> .
\]
Integrating over the equal time hypersurfaces and using the Schwarz inequality, we obtain
\[
|\partial_t (\Xi | \Xi)_t| \leq 2 \| \Xi |_t \| \| \phi |_t \|
\]
and thus
\[
|\partial_t \| \Xi |_t \| \leq \| \phi |_t \|. 
\]
Substituting the specific forms of Ξ and ϕ and using the Schwarz and triangle inequalities, we obtain the estimate

\[ \left| \partial_t \| \nabla^\alpha \partial_m^b \psi_m \|_t \right| \leq b \left\| \nabla^\alpha \partial_m^{b-1} \psi_m \|_t \right\| + c a \left| \nabla \mathcal{B}(t) \right|_{C^{a-1}} \left\| \partial_m^b \psi_m \|_t \right\|_{W^{a-1,2}}, \]  

(13.5.2)

where we used the notation (17.2.2).

We now proceed inductively in the maximal total order \( a + b \) of the derivatives. In the case \( a = b = 0 \), the claim follows immediately from the unitarity of the time evolution. In order to prove the induction step, we note that in (13.5.2), the order of differentiation of the wave function on the right hand side is smaller than that on the left hand side at least by one. In the case \( a = 0 \) and \( b \geq 0 \), the induction hypothesis yields the inequality

\[ \left| \partial_t \| \partial_m^b \psi_m \|_t \right| \leq b \left\| \partial_m^{b-1} \psi_m \|_t \right\| \leq b C \left( 1 + |t|^b \right) \sum_{p=0}^{b-1} \left\| \partial_m^p \psi_m \|_{t=0} \right\|, \]

and integrating this inequality from 0 to \( t \) gives the result. In the case \( a > 0 \) and \( b \geq 0 \), we apply (13.5.1) together with the induction hypothesis to obtain

\[ \left| \partial_t \| \partial_m^b \psi_m \|_t \right|_{W^{a,2}} \leq b C \left( 1 + |t|^b \right) \sum_{p=0}^{b-1} \left\| \partial_m^p \psi_m \|_{t=0} \right\|_{W^{a,2}} \]

\[ + c C \frac{1 + |t|^b}{1 + |t|^{b+\varepsilon}} \sum_{p=0}^{b} \left\| \partial_m^p \psi_m \|_{t=0} \right\|_{W^{a-1,2}}. \]

Again integrating over \( t \) gives the result. \( \square \)

13.6. The Cauchy Problem in Globally Hyperbolic Spacetimes

We conclude this chapter by extending the global existence and uniqueness result for the Dirac equation to curved spacetime. These results were already stated in Section 4.5. We are now in the position for giving the proof. The reader not interested in or not familiar with curved spacetime may skip this section. We note that more details on the geometric properties of globally hyperbolic spacetimes can be found in [6, Section 3.2].

Proof of Theorem 4.5.1. Exactly as explained in the proof of Theorem 13.4.1 by considering the Cauchy problem for \( \psi - \psi_0 \) one may reduce the problem to that of zero initial data zero. Moreover, choosing a partition of unity \( (\eta_k) \) subordinate to the charts of a given atlas, it suffices to consider the compactly supported inhomogeneity \( \eta_k \phi \) (the sum over \( k \) is again locally finite, similar as explained in the proof of Theorem 13.4.1). In view of these constructions, it remains to consider the Cauchy problem

\[ (\mathcal{D} - m)\psi = \phi \in C_0^\infty(M, S M), \quad \psi|_{\mathcal{N}_{t_0}} = 0. \]  

(13.6.1)

We denote the support of \( \phi \) by \( K \).

Clearly, in local charts the Dirac equation can be written as a symmetric hyperbolic system. Therefore, the results in Sections 13.2 and 13.3 yield existence and uniqueness of solutions of the Cauchy problem in local charts. This also yields global uniqueness: Let \( \psi \) and \( \tilde{\psi} \) be two smooth solutions of the Cauchy problem (13.6.1). Then their difference \( \Xi := \tilde{\psi} - \psi \) is a homogeneous solution which vanishes at time \( t_0 \). In view of a possible time
reversal, it suffices to consider the solution in the future of \(t_0\). Thus let \(x \in \mathcal{M}\) be in the future of \(t_0\). Then the past light cone \(J^\wedge(x)\) intersects the future of \(t_0\) in a compact set,

\[
J^\wedge(x) \cap \left( \bigcup_{t \geq t_0} \mathcal{N}_t \right)
\]

is compact.

Therefore, we can choose \(\delta > 0\) such that for every \(\hat{t}\), there is a finite number of lens-shaped regions which cover the time strip

\[
J^\wedge(x) \cap \left( \bigcup_{t = \hat{t}}^{\hat{t} + \delta} \mathcal{N}_t \right).
\]

On each lens-shaped regions, the solution for the Cauchy problem with zero initial data vanishes identically. Therefore, we can proceed inductively to conclude that \(\Xi(x) = 0\).

Since \(x\) is arbitrary, the solution \(\Xi\) vanishes identically on \(\mathcal{M}\).

In order to prove global existence, we proceed indirectly. In view of a possible time reversal, it suffices to consider the Cauchy problem to the future. Thus suppose that the solution exists only up to finite time \(t_{\text{max}}\) (see Figure 13.5). Due to finite propagation speed, the solution is supported in the domain of causal dependence of \(K\),

\[
\text{supp } \psi \subset J^\vee(K) \cup J^\wedge(K).
\]

By properties of globally hyperbolic spacetimes, the intersection \(D\) of this set with the Cauchy surface \(\mathcal{N}_{t_{\text{max}}}\) is compact. Covering \(D\) by a finite number of charts, we choose \(\delta\) such that the sets \(J^\vee(D) \cup J^\wedge(D) \cap \mathcal{N}_t\) lie in the domain of these charts for all \(t \in [t_{\text{max}} - \delta, t_{\text{max}} + \delta]\). Next we choose a finite number of lens-shaped regions \(L_\ell\) which also cover all these sets (see again Figure 13.5). In each of these lens-shaped regions we can solve the Cauchy problem with initial data at time \(t_{\text{max}} - \delta\). In this way, we get a solution up to time \(t_{\text{max}} + \delta\). This is a contradiction, thereby proving that the solution must exist for all times.

**Proof of Theorem 13.5.2** By extending the initial data \(\psi_0\) to a smooth and compactly supported function in spacetime and considering the Cauchy problem for \(\psi - \psi_0\), it again suffices to consider the case of zero initial data (13.6.1). The solution constructed subsequently the proof of Theorem thmcauchy was supported in \(\mathfrak{F}^\vee(K) \subset \mathfrak{F}^\wedge(K)\). By general properties of globally hyperbolic manifolds, the intersection of this set with every Cauchy surface is compact. This concludes the proof.

Similar as explained in Section 13.4 in Minkowski space, also in curved spacetime the solution of the Cauchy problem can be expressed in terms of the causal fundamental

![Figure 13.5. Global solutions in globally hyperbolic spacetimes.](image-url)
solution \( k_m \), as we now explain. Similar as explained in Minkowski space in Section 13.4, the retarded and advanced Green’s operators \( s^\wedge_m \) and \( s^\vee_m \) are linear mappings (for details see for example [27, 6])

\[
s^\wedge_m, s^\vee_m : C_0^\infty(\mathcal{M}, S\mathcal{M}) \rightarrow C_c^\infty(\mathcal{M}, S\mathcal{M}).
\]

They satisfy the defining equation of the Green’s operator

\[
(\mathcal{D} - m) \left( s^\wedge_m \phi \right) = \phi.
\]

(13.6.2)

Moreover, they are uniquely determined by the condition that the support of \( s^\wedge_m \phi \) (or \( s^\vee_m \phi \)) lies in the future (respectively the past) of \( \text{supp} \phi \). The causal fundamental solution \( k_m \) is introduced by

\[
k_m := \frac{1}{2\pi i} \left( s^\vee_m - s^\wedge_m \right) : C_0^\infty(\mathcal{M}, S\mathcal{M}) \rightarrow C_c^\infty(\mathcal{M}, S\mathcal{M}) \cap \mathcal{H}_m.
\]

(13.6.3)

Note that it maps to solutions of the Dirac equation.

**Proposition 13.6.1.** The solution of the Cauchy problem (13.5.1) has the representation

\[
\psi(x) = 2\pi \int_N k_m(x, y) \psi(y) d\mu_N(y),
\]

(13.6.4)

where \( k_m(x, y) \) is the causal fundamental solution (18.2.10).

**Proof.** Let us consider a point \( x \) in the future of \( \mathcal{N} \) (the case for the past is analogous). In this case, due to (18.2.10), the lemma simplifies to

\[
\psi(x) = i \int_N s^\wedge_m(x, y) \psi(y) d\mu_N(y).
\]

(13.6.5)

In preparation, we want to prove that for any \( \phi \in C^\infty(\mathcal{M}, S\mathcal{M}) \) which has compact support to the past of \( \mathcal{N} \) and with the property that \( (\mathcal{D} - m)\phi \) has compact support the equation

\[
\phi = s^\wedge_m ((\mathcal{D} - m)\phi)
\]

(13.6.6)

holds. To this end, we consider the function

\[
\Xi(x) := \phi - s^\wedge_m ((\mathcal{D} - m)\phi).
\]

Applying the operator \( (\mathcal{D} - m) \) and using the defining equation of the Green’s operators, one sees that \( \Xi \) is a solution of the Dirac equation. Moreover, \( \Xi \) obviously vanishes in the past of the support of \( \phi \). The uniqueness of the solution of the Cauchy problem implies that \( \Xi \) vanishes identically, proving (13.6.6).

In order to derive equation (13.6.5), we let \( \eta \in C^\infty(\mathcal{M}) \) be a function which is identically equal to one at \( x \) and on \( \mathcal{N} \), but such that the function \( \eta \psi \) has compact support. (For example, in a foliation \((\mathcal{N}_t)_{t \in \mathbb{R}}\) with \( \mathcal{N} = \mathcal{N}_{t_0} \) one can take \( \eta = \eta(t) \) as a smooth function with \( \eta|_{[t_0, \infty)} \equiv 1 \) which vanishes if \( t < t_0 - 1 \).) Then we can apply (13.6.6) to the wave function \( \phi = \eta \psi \). We thus obtain for any \( x \) in the future of \( \mathcal{N} \) the relations

\[
\psi(x) = (\eta \psi)(x) = \left( s^\wedge_m ((\mathcal{D} - m)(\eta \psi)) \right)(x) = \left( s^\wedge_m (i\gamma^j(\partial_j \eta) \psi) \right)(x),
\]

(13.6.7)

where we have used that \( \psi \) is a solution of the Dirac equation.

To conclude the proof, for \( \eta \) in (13.6.7) we choose a sequence \( \eta_j \) which converges in the distributional sense to the function which in the future \( \mathcal{N} \) is equal to one and in the past of \( \mathcal{N} \) is equal to zero. This yields \( \partial_j \eta_j \rightarrow \nu \), and thus the right-hand-side of (13.6.7) is equal to the right-hand-side of (13.6.5). \( \square \)
Exercises

Exercise 13.1. The homogeneous Maxwell equations for the electric field $E : \mathbb{R}^3 \to \mathbb{R}^3$ and the magnetic field $B : \mathbb{R}^3 \to \mathbb{R}^3$ read
\[
\nabla \times B = \partial_t E, \quad \nabla \times E = -\partial_t B.
\]
Rewrite these equations as a symmetric hyperbolic system. Remark: We here ignore the equations $\text{div} \ E = \text{div} \ B = 0$. The reason is that these equations hold automatically if they are satisfied initially.

Exercise 13.2. Consider the scalar wave equation \((\partial_{tt} - \Delta_{\mathbb{R}^n}) \phi(t,x) = 0\).
(a) Rewrite the equation as a symmetric hyperbolic system
\[
A^0 \partial_t u + \sum_{\alpha=1}^m A^\alpha \nabla_{\alpha} u + Bu = 0.
\]
(b) Determine the timelike and future-directed directions. Which directions $\xi$ are characteristic (meaning that the characteristic polynomial $\det A(x,\xi)$ vanishes)?
(c) Express the “energy”
\[
E(t) = \int_{\mathbb{R}^n} \langle u, A^0 u \rangle \, d^m x
\]
in terms of $\phi(t,x)$. Compare the resulting expression with the conserved physical energy
\[
\int_{\mathbb{R}^n} (|\partial_t \phi|^2 + |\nabla \phi|^2) \, d^m x.
\]
(d) Compute $\frac{dE(t)}{dt}$. Prove the inequality
\[
\frac{dE(t)}{dt} \leq E(t)
\]
and integrate it (Grönwall’s lemma).

Exercise 13.3. Consider the solution of the homogeneous wave equation
\[
(\partial_{tt} - \Delta_{\mathbb{R}^n}) \phi(t,x) = 0
\]
for smooth initial data $\phi(0,x) = f(x)$ and $\partial_t \phi(0,x) = g(x)$.
Show by a suitable choice of lens-shaped regions that $\phi(t_0,x_0)$ depends only on the initial data in the closed ball $\{ x \in \mathbb{R}^n : |x-x_0| \leq t_0 \}$.

Exercise 13.4. We consider the system
\[
\begin{align*}
\partial_t u_1(t, x) + \partial_x u_1(t, x) + 4 \partial_x u_2(t, x) &= 0 \\
\partial_t u_2(t, x) + 4 \partial_x u_1(t, x) + \partial_x u_2(t, x) &= 0
\end{align*}
\]
(a) Write the system in symmetric hyperbolic form.
(b) Compute the solution of the Cauchy problem for initial data $u_1(0, x) = \sin x$ and $u_2(0, x) = \cos x$.

Exercise 13.5. (The Euler equations) The evolution equation for an isentropic compressible fluid reads
\[
\begin{align*}
\partial_t v + \nabla_v v + \frac{1}{\rho} \text{grad}(p) &= 0 \\
\partial_t \rho + \nabla_v \rho + \rho \text{div}(v) &= 0.
\end{align*}
\] (13.6.8)
Here $v : \mathbb{R}^+ \times \mathbb{R}^3 \to \mathbb{R}^3$ is the velocity vector field, $\rho : \mathbb{R}^+ \times \mathbb{R}^3 \to (0, \infty)$ the density and $p = A \rho^{\gamma}$ the pressure (where $A > 0$ and $\gamma > 1$).
(a) Show that (13.6.8) is equivalent to a quasilinear symmetric hyperbolic system, provided that \( \rho \) is bounded away from zero.

(b) Show that for smooth solutions, the system (13.6.8) is equivalent to

\[
\begin{align*}
\partial_t v + \nabla_v v + \text{grad}(h(\rho)) &= 0 \\
\partial_t \rho + \text{div}(\rho v) &= 0,
\end{align*}
\]

(13.6.9)

where \( h \in C^\infty(\mathbb{R}) \) satisfies the equation \( h'(\rho) = \rho^{-1}p'(\rho) \).

(c) Let \((v, \rho)\) be a solution of (13.6.9) with \( v(t,x) = \nabla_x \varphi(t,x) \) for a real-valued potential \( \varphi \). Prove Bernoulli’s law: If \( \varphi \) and \( \rho \) decay at infinity sufficiently fast and if \( h(0) = 0 \), then

\[
\partial_t \varphi + \frac{1}{2} |\nabla_x \varphi|^2 + h(\rho) = 0.
\]

(13.6.10)

(d) Show that (13.6.8) can also be rewritten as a system for \((p, v)\),

\[
\begin{align*}
\partial_t v + \nabla_v v + \rho(p)^{-1} \text{grad}(p) &= 0 \\
\partial_t p + \nabla_v p + (\gamma p) \text{div}(v) &= 0. 
\end{align*}
\]

(13.6.11)

Rewrite this system in symmetric hyperbolic form.

**Exercise 13.6.** Let \( \lambda > 0 \). A symmetric hyperbolic system of the form

\[
\partial_t u + A^\alpha(u) \partial_\alpha u + \lambda u = 0,
\]

where the matrices \( A^\alpha \) are smooth, uniformly bounded and uniformly positive, is an example of a so-called dissipative system.

(a) Prove that for spatially compact solutions, the following energy estimate holds:

\[
\frac{d}{dt} \|u(t)\|_{H^p}^2 \leq \left( -2\lambda + c\|u(t)\|_{C^1} \right) \|u(t)\|_{H^p}^2.
\]

(b) Prove: If the initial data \( u_0 \) is sufficiently small in the \( C^1 \)-norm, then there exists a global solution.

**Hint:** Choose \( p \) sufficiently large and use the Sobolev embedding theorem.

**Exercise 13.7.** (Causality in the setting of symmetric hyperbolic systems) The Dirac equation \((i\partial - m)\psi = 0\) can be rewritten as a symmetric hyperbolic system, i.e. in the form \( c > 0 \)

\[
(A^0(x) \partial_0 + A^\alpha(x) \partial_\alpha + B(x))\psi = 0, \quad \text{with} \quad (A^i)^\dagger = A^i \quad \text{and} \quad A^0(x) \geq cI.
\]

For such systems a notion of causality can be introduced: a vector \( \xi \in \mathbb{R}^4 \) is said to be time-like or light-like at \( x \in \mathbb{R}^4 \), if the matrix \( A(x, \xi) := A^i(x) \xi_i \) is definite (either positive or negative) or singular, respectively.

Find the matrices \( A^i \) and \( B \) for the Dirac equation and show that the above notions of time-like and light-like vectors coincide with the corresponding notions in Minkowski space as explained in the lecture.
CHAPTER 14

Energy Methods for the Linearized Field Equations

In the previous chapter, we used energy methods in order to study the Cauchy problem for linear symmetric hyperbolic systems. We now briefly explain how these methods can be adapted to the linearized field equations for causal variational principles as introduced in Chapter 8. These constructions are carried out in detail in [22]; for later developments see [61, 70]. Here we do not aim for the largest generality, but instead explain the basic ideas in the simplest possible setting.

14.1. Local Foliations by Surface Layers

We consider causal variational principles in the compact setting (see Section 6.3). Moreover, for technical simplicity we again assume that the Lagrangian is smooth (6.2.4). Following our procedure for symmetric hyperbolic systems, we want to analyze the initial problem “locally” in an open subset \( U \) of space-time \( M \). In analogy to the time function in a lens-shaped region \( L \) (see Section 13.2) we here choose a foliation of a compact subset \( L \subset U \) by surface layers. This motivates the following definition.

Definition 14.1.1. Let \( U \subset M \) be an open subset of space-time and \( I := [t_{\min}, t_{\max}] \) a compact interval. Moreover, we let \( \eta \in C^\infty(I \times U, \mathbb{R}) \) be a function with \( 0 \leq \eta \leq 1 \) which for all \( t \in I \) has the following properties:

(i) The function \( \theta(t,.) := \partial_t \eta(t,.) \) is non-negative and compactly supported in \( U \).

(ii) For all \( x \in \text{supp} \theta(t,.) \) and all \( y \in M \setminus U \), the function \( \mathcal{L}(x,y) \) as well as its first and second derivatives.

We also write \( \eta(t,x) \) as \( \eta_t(x) \) and \( \theta(t,x) \) as \( \theta_t(x) \). We refer to \( \eta_t \in I \) as a local foliation inside \( U \).

The situation in mind is shown in Figure 14.1. The parameter \( t \) can be thought of as the time of a local observer and will be referred to simply as time. The support of the function \( \theta_t \) is a surface layer. The function \( \eta_t \) should be thought of as being equal to one in the past and equal to zero in the future of this surface layer. The condition (i) implies
that the set \( L \) defined by
\[
L := \bigcup_{t \in I} \text{supp} \theta_t
\] (14.1.1)
is compact. It is the region of space-time described by the local foliation. The condition (ii) has the purpose to ensure that the dynamics in the region \( L \) does not depend on the jets outside \( U \), making it possible to restrict attention to the space-time region \( U \). Sometimes, we refer to this property that \( L \) is \( L \)-localized in \( U \).

Following the procedure for hyperbolic partial differential equations, our first goal is to analyze the initial value problem. For the following constructions, it will be useful to combine the functions \( \eta_t \) and \( \theta_t \) with the measure \( \rho \) such as to form new measures: The measure
\[
d\rho_t(x) := \theta_t(x) \, d\rho(x)
\] (14.1.2)
with \( t \in I \) is supported in the surface layer at time \( t \). Likewise, the measures \( \eta_t \) \( d\rho \) and \( (1 - \eta_t) \) \( d\rho \) are supported in the past respectively future of the surface layer at time \( t \). For the measures supported in a space-time strip, we use the notation
\[
\eta_{[t_0,t_1]} \, d\rho \quad \text{with} \quad \eta_{[t_0,t_1]} := \eta_{t_1} - \eta_{t_0} \in C_0^\infty(U),
\] (14.1.3)
where we always choose \( t_0, t_1 \in I \) with \( t_0 \leq t_1 \). Note that the function \( \eta_{[t_0,t_1]} \) is supported in \( L \).

### 14.2. Energy Estimates and Hyperbolicity Conditions

For the analysis of the linearized field equations it is helpful to study the surface layer integrals as introduced in Section 9.5 for our local foliation \( (\eta_t)_{t \in I} \). It is useful to “soften” these surface layer integrals by surface layer integrals by rewriting the integration domains with characteristic functions and replacing the characteristic functions by smooth cutoff functions formed of \( \eta_t \), i.e. symbolically
\[
\int \Omega d\rho(x) \int_M \Omega d\rho(y) \cdots = \int_M d\rho(x) \int_M d\rho(y) \chi_\Omega(x) (1 - \chi_\Omega(y)) \cdots \\
\rightarrow \int_M d\rho(x) \int_M d\rho(y) \eta_t(x) (1 - \eta_t(y)) \cdots .
\]
We thus define the softened symplectic form and the softened surface layer inner product by
\[
(u,v)^t = \int_U d\rho(x) \, \eta_t(x) \int_U d\rho(y) \, (1 - \eta_t(y)) \left( \nabla_{1,u} \nabla_{1,v} - \nabla_{2,u} \nabla_{2,v} \right) \mathcal{L}(x,y) \] (14.2.1)
\[
\sigma^t(u,v) = \int_U d\rho(x) \, \eta_t(x) \int_U d\rho(y) \, (1 - \eta_t(y)) \left( \nabla_{1,u} \nabla_{2,v} - \nabla_{1,v} \nabla_{2,u} \right) \mathcal{L}(x,y) .
\] (14.2.2)

The quantity \( (u,u)^t \) is of central importance for the following constructions, because it will play the role of the energy used in our energy estimates. In preparation of these estimates, we derive an energy identity:
Using the property in Definition 14.1.1 (ii), in the solutions (see Section 13.3). With this in mind, our strategy is to express the hyperbolicity condition. This notion can be understood as follows. As explained in Chapter 13, in the theory of hyperbolic partial differential equations the hyperbolicity of the equations gives rise to a positive energy (see $E(\lambda)$ in (13.3.9)). The positivity of this energy was a consequence of the structure of the equations (more precisely, for symmetric linear hyperbolic systems, it is a consequence of the positivity statement in Definition 13.1.1 (ii)). The positivity of the energy is crucial for the analysis of hyperbolic equations, because it gives both uniqueness (see Section 13.2) and existence of weak solutions (see Section 13.3). With this in mind, our strategy is to express the hyperbolic

\[ \frac{d}{dt} \langle u, u \rangle^t = 2 \int_U \langle u, \Delta u \rangle(x) \, dp_t(x) \]

\[ - 2 \int_U \Delta_2[u, u] \, dp_t(x) + s \int_U a(x)^2 \, dp_t(x), \]  

(14.2.3)

where the operator $\Delta_2 : \mathcal{J} \times \mathcal{J} \to \mathcal{J}^*$ defined by

\[ \langle u, \Delta_2[u_1, u_2] \rangle(x) \]

\[ = \frac{1}{2} \nabla_u \left( \int_M (\nabla_{1,u_1} + \nabla_{2,u_1})(\nabla_{1,u_2} + \nabla_{2,u_2}) \mathcal{L}(x, y) \, dp(y) - \nabla_{u_1} \nabla_{u_2} s \right). \]  

(14.2.4)

Proof. Differentiating (14.2.1) with respect to $t$ gives

\[ \frac{d}{dt} \langle u, u \rangle^t = \int_U dp(x) \theta_t(x) \int_U dp(y) (1 - \eta_t(y)) (\nabla_{1,u}^2 - \nabla_{2,u}^2) \mathcal{L}(x, y) \]

\[ - \int_U dp(x) \theta_t(x) \int_U dp(y) \eta_t(y) (\nabla_{1,u}^2 - \nabla_{2,u}^2) \mathcal{L}(x, y) \]

\[ = \int_U dp(x) \theta_t(x) \int_U dp(y) (\nabla_{1,u}^2 - \nabla_{2,u}^2) \mathcal{L}(x, y). \]  

(14.2.5)

Next, for all $x \in M$ we may use Definition 14.1.1 (ii) to change the integration range in (8.1.6) from $M$ to $U$,

\[ \langle u, \Delta u \rangle(x) = \int_U \nabla_{1,u}(\nabla_{1,u} + \nabla_{2,u}) \mathcal{L}(x, y) \, dp(y) - s \, a(x)^2. \]

Multiplying by $\theta_t$ and integrating, we obtain

\[ 0 = \int_U \theta_t(x) \langle u, \Delta u \rangle(x) \, dp(x) + s \int_U \theta_t(x) a(x)^2 \, dp(x) \]

\[ - \int_U dp(x) \theta_t(x) \int_U dp(y) (\nabla_{1,u}^2 + \nabla_{1,u} \nabla_{2,u}) \mathcal{L}(x, y). \]

We multiply this equation by two and add (14.2.5). This gives

\[ \frac{d}{dt} \langle u, u \rangle^t = - \int_U dp(x) \theta_t(x) \int_U dp(y) (\nabla_{1,u}^2 + \nabla_{2,u})^2 \mathcal{L}(x, y) \]

\[ + 2 \int_U \theta_t(x) \langle u, \Delta u \rangle(x) \, dp(x) + 2s \int_U \theta_t(x) a(x)^2 \, dp(x). \]

Using the property in Definition 14.1.1 (ii), in the $y$-integral we may replace the integration range $U$ by $M$, making it possible to apply (14.2.4). Rewriting the obtained integrals using the notation (14.1.2) gives (14.2.3). \(\square\)

In order to make use of this energy identity, we need to impose a condition which we call hyperbolicity condition. This notion can be understood as follows. As explained in Chapter 13, in the theory of hyperbolic partial differential equations the hyperbolicity of the equations gives rise to a positive energy (see $E(\lambda)$ in (13.3.9)). The positivity of this energy was a consequence of the structure of the equations (more precisely, for symmetric linear hyperbolic systems, it is a consequence of the positivity statement in Definition 13.1.1 (ii)). The positivity of the energy is crucial for the analysis of hyperbolic equations, because it gives both uniqueness (see Section 13.2) and existence of weak solutions (see Section 13.3). With this in mind, our strategy is to express the hyperbolic
nature of the linearized field equations by imposing a positivity condition for our “energy” \((u,u)^t\). As we shall see, this so-called hyperbolicity condition is precisely what is needed in order to obtain existence and uniqueness of solutions. For Dirac systems in Minkowski space, the hyperbolicity conditions can be verified by direct computation (for details see [52]). With this in mind, our hyperbolicity conditions are physically sensible.

For all \(x \in M\) we choose the subspace of the tangent space spanned by the test jets,

\[
\Gamma_x := \{ u(x) \mid u \in \Gamma^{\text{test}} \} \subset T_x \mathcal{F}.
\]

We introduce a Riemannian metric \(g_x\) on \(\Gamma_x\). The choice of the Riemannian metric is arbitrary; the resulting freedom can be used in order to satisfy the hyperbolicity conditions below (note, however, that for causal fermion systems a canonical Riemannian metric is obtained form the Hilbert-Schmidt scalar product; see [66, 73]). This Riemannian metric also induces a pointwise scalar product on the jets. Namely, setting

\[
\mathfrak{J}_x := \mathbb{R} \oplus \Gamma_x,
\]

we obtain the scalar product on \(\mathfrak{J}_x\)

\[
\langle \cdot, \cdot \rangle_x : \mathfrak{J}_x \times \mathfrak{J}_x \to \mathbb{R}, \quad \langle u, \tilde{u} \rangle_x := a(x) \tilde{\alpha}(x) + g_x(u(x), \tilde{u}(x))
\]

(14.2.6)

(where we again denote the scalar and vector components of the jet by \(u = (a,u)\)). We denote the corresponding norm by \(\|u\|_x\). We begin with a simple estimate of the energy identity in Lemma [14.2.1]

\[
\text{Lemma 14.2.3. Assume that the hyperbolicity condition of Definition [14.2.2] holds. Then for every } t \in I \text{ and all } u \in \mathfrak{J},
\]

\[
\frac{d}{dt} \|u\|^t \leq C \|\Delta u\|_{L^2(U,d\rho_t)} + c \|u\|^t
\]

(14.2.8)

with

\[
c := c^2 + \frac{C^2}{2}.
\]

**Proof.** Applying (14.2.7) in (14.2.3), we obtain

\[
\frac{d}{dt} \langle u, u \rangle^t \leq 2 \int_U \langle u, \Delta u \rangle_x \, d\rho_t(x) - 2 \int_U \Delta_2[u,u] \, d\rho_t(x) + \frac{s}{2} \int_U b(x)^2 \, d\rho_t(x)
\]

\[
\leq 2 \int_U \langle u, \Delta u \rangle_x \, d\rho_t(x) + \left(2c^2 + c^2 s\right) \langle u, u \rangle^t
\]

\[
\leq 2 \|u\|_{L^2(U,d\rho_t)} \|\Delta u\|_{L^2(U,d\rho_t)} + 2c \langle u, u \rangle^t
\]

\[
\leq 2C \|u\|^t \|\Delta u\|_{L^2(U,d\rho_t)} + 2c \langle u, u \rangle^t,
\]

where in the last line we applied (14.2.7). Using the relation \(\partial_t \|u\|^t = \partial_t(u,u)^t/(2\|u\|^t)\) gives the result. \(\square\)
Applying Grönwall-type estimates, the inequality (14.2.8) shows that $\|u\|^t$ grows at most exponentially in time, provided that $\Delta u$ decays in time sufficiently fast. We here make this statement precise by estimates in Hilbert spaces of jets with zero initial values. In the lens-shaped region $L$ we work with the $L^2$-scalar product

$$\langle u, v \rangle_{L^2(L)} := \int_L \langle u(x), v(x) \rangle_x \eta_I(x) \, d\rho(x) ,$$

(14.2.9)

which, according to (14.1.2) and (14.1.3), can also be written in terms of a time integral,

$$\langle u, v \rangle_{L^2(L)} = \int_{t_0}^{t_{\text{max}}} \langle u, v \rangle_{L^2(U, d\rho_t)} \, dt .$$

(14.2.10)

The corresponding norm is denoted by $\|\cdot\|_{L^2(L)}$.

**Proposition 14.2.4. (energy estimate)** Assume that the hyperbolicity condition of Definition 14.2.2 holds. Then, choosing

$$\Gamma = 2 C e^{2c (t_{\text{max}} - t_0)} \left( t_{\text{max}} - t_0 \right) ,$$

(14.2.11)

the following estimate holds,

$$\|u\|_{L^2(L)} \leq \Gamma \|\Delta u\|_{L^2(L)} \quad \text{for all } u \in \mathcal{J} \text{ with } \|u\|_{t_0} = 0 .$$

**Proof.** We write the energy estimate of Lemma 14.2.3 as

$$\frac{d}{dt} \left( e^{-2ct} (u, u)^t \right) \leq 2 e^{-2ct} C \|u\|^t \|\Delta u\|_{L^2(U, d\rho_t)} .$$

Integrating over $t$ from $t_0$ to some $t \in I$ and using the hyperbolicity condition (14.2.7), we obtain

$$e^{-2ct} (u, u)^t = \int_{t_0}^{t} \frac{d}{dt'} \left( e^{-2ct'} (u, u)^{t'} \right) \, dt'$$

$$\leq 2 C \int_{t_0}^{t} e^{-2ct'} \|u\|^{t'} \|\Delta u\|_{L^2(U, d\rho_{t'})} \, dt' .$$

Multiplying by $e^{2ct}$ gives the inequality

$$(u, u)^t \leq 2 C \int_{t_0}^{t} e^{2c (t-t')} \|u\|^{t'} \|\Delta u\|_{L^2(U, d\rho_{t'})} \, dt'$$

$$\leq 2 C e^{2c (t_{\text{max}} - t_0)} \int_{t_0}^{t_{\text{max}}} \|u\|^{t'} \|\Delta u\|_{L^2(U, d\rho_{t'})} \, dt'$$

$$\leq 2 C e^{2c (t_{\text{max}} - t_0)} \|\Delta u\|_{L^2(L)} \left( \int_{t_0}^{t_{\text{max}}} (u, u)^{t'} \, dt' \right)^{\frac{1}{2}} ,$$

where in the last step we used the Schwarz inequality and (14.2.10). Integrating once again over $t$ from $t_0$ to $t_{\text{max}}$ gives

$$\left( \int_{t_0}^{t_{\text{max}}} (u, u)^{t} \, dt \right)^{\frac{1}{2}} \leq 2 C e^{2c (t_{\text{max}} - t_0)} (t_{\text{max}} - t_0) \|\Delta u\|_{L^2(L)} .$$

(14.2.12)

Finally, we apply the hyperbolicity condition (14.2.7) in (14.2.10),

$$\|v\|_{L^2(L)} = \left( \int_{t_0}^{t_{\text{max}}} \|u\|_{L^2(U, d\rho_t)}^2 \, dt \right)^{\frac{1}{2}} \leq C \left( \int_{t_0}^{t_{\text{max}}} (u, u)^{t} \, dt \right)^{\frac{1}{2}} .$$

Combining this inequality with (14.2.12) gives the result. \qed
14.3. Uniqueness of Strong Solutions

Based on the above energy estimates, we can now prove uniqueness of strong solutions of the Cauchy problem. The methods is quite similar to that employed in Section 13.2 for symmetric hyperbolic systems. In preparation of formulating the Cauchy problem, we need to introduce jets which vanish at initial time $t_{\text{min}}$. To this end, we demand that the jet vanishes in the surface layer and that the corresponding softened surface layer integrals (14.2.1) and (14.2.2) vanish,

$$\mathfrak{J}_{t_{\text{min}}} := \{ u \in \mathfrak{J} \mid \eta_{t_{\text{min}}} u \equiv 0 \quad \text{and} \quad (u, v)_{t_{\text{min}}} = 0 = \sigma^{t_{\text{min}}}(u, v) \quad \text{for all} \quad v \in \mathfrak{J} \}.$$ 

Similarly, we define the space of jets which vanish at final time $t_{\text{max}}$ by

$$\mathfrak{J}_{t_{\text{max}}^U} := \{ u \in \mathfrak{J} \mid (1 - \eta_{t_{\text{max}}}) u \equiv 0 \quad \text{and} \quad (u, v)_{t_{\text{max}}^U} = 0 = \sigma^{t_{\text{max}}}(u, v) \quad \text{for all} \quad v \in \mathfrak{J} \}.$$ 

A strong solution of the Cauchy problem is a jet $u \in \mathfrak{J}_{t_{\text{max}}^U}$ which satisfies the equations

$$\Delta u = w \quad \text{in} \quad L \quad \text{and} \quad u - u_0 \in \mathfrak{J}_{t_{\text{min}}},$$

where $u_0 \in \mathfrak{J}$ is the initial data and $w$ is the inhomogeneity.

**Proposition 14.3.1. (uniqueness of strong solutions)** Let $(\eta_t)_{t \in I}$ be a local foliation inside $U$ which satisfies the hyperbolicity conditions (see Definitions 14.1.1 and 14.2.2). Then the Cauchy problem (14.3.3) with $u_0, w \in \mathfrak{J}_{t_{\text{max}}^U}$ has at most one solution $u$ in $L$.

**Proof.** Let $u$ be the difference of two solutions. Then $u$ is a solution of the homogeneous equation with zero initial data. Applying Lemma 14.2.3, we obtain

$$\frac{d}{dt} \|u\|^t \leq c \|u\|^t \quad \text{and thus} \quad \frac{d}{dt} (e^{-ct} \|u\|^t) \leq 0 .$$

It follows that $\|u\|^t$ vanishes for all $t$ in the respective interval. Using (14.2.7), we conclude that $u$ vanishes identically in $L$. This gives the result. $\square$

Similar as explained in Section 13.2 for symmetric hyperbolic systems, this uniqueness statement also gives information on the speed of propagation and the resulting causal structure. For details we refer to [22, 23].

14.4. Existence of Weak Solutions

Our existence proof is inspired by the method invented by K.O. Friedrichs for symmetric hyperbolic systems in [86] as outlined in Section 13.3. Our first step is to formulate the linearized field equations weakly. To this end, we need to “integrate by parts” with the help of the following lemma.

**Lemma 14.4.1. (Green’s formula)** For all $u, v \in \mathfrak{J}$,

$$\sigma^{t_{\text{max}}}(u, v) - \sigma^{t_{\text{min}}}(u, v) = \langle u, \Delta v \rangle_{L^2(L)} - \langle \Delta u, v \rangle_{L^2(L)} .$$

(14.4.1)
Proof. Using the definition of the $L^2$-scalar product in (14.2.9) and the definition of the linearized field operator $[8.1.0]$, we obtain

$$\langle u, Δv \rangle_{L^2(L)} - \langle Δu, v \rangle_{L^2(L)} = \int_U (\langle u, Δv \rangle - \langle Δu, v \rangle) \eta_I d\rho$$

$$= \int_U d\rho(x) \eta_I(x) \nabla_u \left( \int_M (\nabla_{1,v} + \nabla_{2,v})L(x,y) \; d\rho(y) - \nabla_v s \right)$$

$$- \int_U d\rho(x) \eta_I(x) \nabla_v \left( \int_M (\nabla_{1,u} + \nabla_{2,u})L(x,y) \; d\rho(y) - \nabla_u s \right).$$

Here the space-time point $x$ is in $L$. Using Definition [14.1.1] (ii), we get a contribution to the integrals only if $y ∈ U$. Therefore, we may replace the integration range $M$ by $U$. We thus obtain

$$\langle u, Δv \rangle_{L^2(L)} - \langle Δu, v \rangle_{L^2(L)}$$

$$= \int_U d\rho(x) \eta_I(x) \int_U d\rho(y)(\nabla_{1,u} \nabla_{2,v} - \nabla_{2,u} \nabla_{1,v})L(x,y),$$

where we used that, following our convention $[8.1]$, the second derivatives of the Lagrangian are symmetric. Using the definition $[14.1.3]$ as well as the anti-symmetry of the integrand, the term $[15.4.4]$ can be rewritten as

$$\int_U d\rho(x) \eta_I(x) \int_U d\rho(y)(\nabla_{1,u} \nabla_{2,v} - \nabla_{2,u} \nabla_{1,v})L(x,y)$$

$$= \int_U d\rho(x) \int_U d\rho(y) \eta_I(x) (\nabla_{1,u} \nabla_{2,v} - \nabla_{2,u} \nabla_{1,v})L(x,y)\bigg|^{t_{max}}_{t_0}$$

$$= \int_U d\rho(x) \int_U d\rho(y) (\eta_I(x) - \eta_I(x) \eta_I(y)) (\nabla_{1,u} \nabla_{2,v} - \nabla_{2,u} \nabla_{1,v})L(x,y)\bigg|^{t_{max}}_{t_0}$$

$$= \int_U d\rho(x) \int_U d\rho(y) \eta_I(x) (1 - \eta_I(y)) (\nabla_{1,u} \nabla_{2,v} - \nabla_{2,u} \nabla_{1,v})L(x,y)\bigg|^{t_{max}}_{t_0}$$

$$= \sigma_{t_{max}}(u,v) - \sigma_{t_{min}}(u,v).$$

This gives the result. \[\square\]

Assume that $u$ is a strong solution of the Cauchy problem $[14.3.3]$. As usual, replacing $u$ by $u - u_0$ and $w$ by $w - Δu_0 ∈ J$, it suffices to consider the Cauchy problem for zero initial data, i.e.,

$$Δu = w \quad \text{in } U \quad \text{and} \quad u ∈ J_{U_{t_{min}}}. \quad (14.4.3)$$

Then, applying the above Green’s formula, we obtain for any $v ∈ J$,

$$\langle v, w \rangle_{L^2(L)} = \langle u, Δu \rangle_{L^2(L)} - \langle Δv, u \rangle_{L^2(L)} - \sigma_{t_{max}}(v, u) + \sigma_{t_{min}}(v, u).$$

Having implemented the vanishing initial data by the condition $u ∈ J_{U_{t_{min}}}$, the symplectic form vanishes at time $t_{min}$. In order to also get rid of the boundary values at time $t_{max}$, we restrict attention to test jets which vanish at $t_{max}$. This leads us to the following definition:

**Definition 14.4.2.** A jet $u ∈ L^2(L)$ is a **weak solution** of the Cauchy problem $Δu = w$ with zero initial data if

$$\langle Δv, u \rangle_{L^2(L)} = \langle v, w \rangle_{L^2(L)} \quad \text{for all } v ∈ J_{t_{max}}. \quad (14.4.4)$$
Clearly, the energy estimate of Proposition 14.2.4 also holds if we exchange the roles of $t_{\text{max}}$ and $t_{\text{min}}$, i.e.,

$$\|u\|_{L^2(L)} \leq \Gamma \|\Delta u\|_{L^2(L)} \quad \text{for all } u \in \overline{J}^{t_{\text{max}}}$$

(14.4.5)

(where the constant $\Gamma$ is again given by (14.2.11)).

We introduce the positive semi-definite bilinear form

$$<[.,.]: \overline{J}^{t_{\text{max}}} \times \overline{J}^{t_{\text{max}}} \rightarrow \mathbb{R}, \quad <u,v> = \langle \Delta u, \Delta v \rangle_{L^2(L)}.$$

Dividing out the null space and forming the completion, we obtain a Hilbert space $(\mathcal{H}, <.,.>).$ The corresponding norm is denoted by $\|\cdot\|.$

We now consider the linear functional $\langle w, . \rangle_{L^2(L)}$ on $\overline{J}^{t_{\text{max}}}$. Applying the Schwarz inequality and (14.4.5), we obtain

$$|\langle w, u \rangle_{L^2(L)}| \leq \|w\|_{L^2(L)} \|u\|_{L^2(L)} \leq \Gamma \|w\|_{L^2(L)} \|u\|,$$

proving that the linear functional $\langle w, . \rangle_{L^2(L)}$ on $\overline{J}^{t_{\text{max}}}$ is bounded on $\mathcal{H}$. Therefore, it can be extended uniquely to a bounded linear functional on all of $\mathcal{H}$. Moreover, by the Fréchet-Riesz theorem there is a unique vector $U \in \mathcal{H}$ with

$$\langle w, v \rangle_{L^2(L)} = <U, v> = \langle \Delta U, \Delta v \rangle_{L^2(L)} \quad \text{for all } v \in \overline{J}^{t_{\text{max}}}.$$

Hence $u := \Delta U \in L^2(L)$ is the desired weak solution. We point out that in the above estimates, the inhomogeneity $w$ enters only via its $L^2$-norm, making it possible to generalize our methods to $w \in L^2(L)$. We have obtained the following result:

**Theorem 14.4.3.** Assume that $(\eta_t)_{t \in I}$ is a local foliation satisfying the hyperbolicity conditions (see Definitions 14.1.1 and 14.2.2). Then for every $w \in L^2(L)$ there is a weak solution $u \in L^2(L)$ of the Cauchy problem (14.4.4).

**Exercises**
CHAPTER 15

Functional Analytic Methods in Space-Time

When constructing a causal fermion system in Minkowski space in Section 5.4, we chose $\mathcal{H}$ as a subspace of the solution space $\mathcal{H}_m$ of the Dirac equation. In principle, one can choose $\mathcal{H}$ as one likes, and different choices give rise to different causal fermion systems. However, if one wants to describe a given physical system, one must specify the subspace $\mathcal{H} \subset \mathcal{H}_m$, and it important to do it right. It is not obvious what “right” and “wrong” should be. Generally speaking, $\mathcal{H}$ can be thought of as the “occupied states” of the physical system in mind. If we want to describe the vacuum in Minkowski space (i.e. no particles and no interaction is present), then the natural and only physically reasonable choice is to let $\mathcal{H}$ be the subspace of all negative-frequency solutions of the Dirac equation. This choice corresponds to the physical concept of the Dirac sea as introduced by Dirac in 1930, which led to the prediction of anti-particles (discovered shortly afterward in 1932, earning Dirac the Nobel prize in 1933). Following these physical concepts, it is also clear that if particles and anti-particles (but no interaction of the matter) is present, then $\mathcal{H}$ is obtained from the subspace of all negative-frequency solutions by occupying additional particle states and by creating “holes” in the sea corresponding to the anti-particle states. Once an interaction (for example an electromagnetic field) is present, it is no longer clear how $\mathcal{H}$ is to be chosen. The reason is that as soon as the fields are time-dependent, the notion of positive and negative frequency solutions breaks down, so that there is no obvious decomposition of the solution space into two subspaces. But for the description of the physical system, a decomposition of the solution space is needed, and taking the “wrong” decomposition leads to artificial mathematical and physical difficulties.

We now explain a functional analytic method which gives rise to a canonical decomposition of the solution space into two subspaces, even in the time-dependent situation. In the static situation, this decomposition reduces to the canonical frequency splitting. This splitting is “right” in the sense that it gives rise to a physically sensible ground state of the system (a so-called Hadamard state, as we will learn in Chapter 19 later in this lecture). Moreover, when performing our construction perturbatively, one can compute the singularities of $P(x,y)$ explicitly working exclusively with bounded line integrals. These explicit computations are the backbone of the analysis of the continuum limit in 48. Before outlining the perturbative constructions (see Chapter 18), we now explain the functional analytic construction.

15.1. General Setting and Basic Ideas

We now recall our setting and explain the idea of the construction. We first summarize the structures of Section 1 using a more general notation, which has the advantage that it can be used just as well if Minkowski space is replaced by a globally hyperbolic spacetime. Thus the reader who is familiar with general relativity and Lorentzian geometry, in what follows can consider $(\mathcal{M},g)$ as a globally hyperbolic Lorentzian manifold with spinor
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bundle \((S\mathcal{M},\prec,\succ)\). The Dirac equation is written as
\[
(D - m)\psi_m = 0 \quad (15.1.1)
\]
(here the subscript \(m\) indicates the mass of the solution; this is of advantage because later on, we shall consider families of solutions with a varying mass parameter). In Minkowski space, one chooses \(D = i\partial^\gamma + B\) such as to get back to \((1.3.14)\). In a globally hyperbolic spacetime, the Dirac operator is a first order differential operator, but the coefficients depend on the metric (for details see Chapter 4). Next, we let \(\mathcal{N}\) be any Cauchy surface. Then the scalar product \((1.3.12)\) on the solutions can be written more generally as
\[
(\psi_m|\phi_m)_m = 2\pi \hat{\mathcal{N}} \prec \psi_m|\phi_m\succ_x d\mu_N(x), \quad (15.1.2)
\]
where \(\nu\) its future-directed normal and \(d\mu_N\) is the volume measure given by the induced Riemannian metric on \(\mathcal{N}\) (in Minkowski space and \(\mathcal{N} = \{t = \text{const}\}\), the normal has the components \(\nu^i = (1,0,0,0)\) and \(d\mu_N = d^3x\), giving back \((1.3.12)\)). Similar to the computation \((1.3.10)\), the vector field \(\prec \psi_m|\phi_m\succ_x\) is again divergence-free, implying that this scalar product is independent of the choice of the Cauchy surface (for details see \([76\text{, Section 2}]\)). Forming the completion gives the Hilbert space \((\mathcal{H}_m,\langle\cdot,\cdot\rangle_m)\).

For the following constructions, we make use of another structure which was not used so far. Namely, given two wave functions \(\psi\) and \(\phi\) (not necessarily solutions of the Dirac equation), one can integrate their pointwise inner product \(\prec \psi|\phi\succ_x\) over spacetime. In order for this integral to be well-defined, one can proceed for example as follows. We denote the smooth sections of the spinor bundle by \(C^\infty(M, SM)\). Similarly, \(C^\infty_0(M, SM)\) denotes the smooth sections with compact support. On the wave functions, one has the Lorentz invariant inner product
\[
\langle\cdot,\cdot\rangle : C^\infty(M, SM) \times C^\infty_0(M, SM) \to \mathbb{C},
\]
\[
\langle\psi|\phi\rangle = \int_M \langle\psi|\phi\rangle_x d\mu_M.
\quad (15.1.3)
\]
In order to explain the basic idea of the construction as first given in \([76]\), let us assume for simplicity that the integral in \((15.1.3)\) exists for all solutions \(\psi_m, \phi_m \in \mathcal{H}_m\). This condition is not satisfied in Minkowski space because the time integral in \((15.1.3)\) in general diverges. But it is indeed satisfied in spacetimes of finite lifetime (for details see \([76\text{, Section 3.2}]\)). Then the spacetime inner product can be extended by continuity to a bilinear form
\[
\langle\cdot,\cdot\rangle : \mathcal{H}_m \times \mathcal{H}_m \to \mathbb{C},
\]
which is bounded, i.e.
\[
|\langle\phi_m|\psi_m\rangle| \leq c \|\phi_m\|_m \|\psi_m\|_m \quad (15.1.4)
\]
(where \(\|\cdot\|_m = (\cdot|\cdot)_m^{1/2}\) is the norm on \(\mathcal{H}_m\)). Then, applying the Fréchet-Riesz theorem, we can uniquely represent this inner product with a signature operator \(S\),
\[
S : \mathcal{H}_m \to \mathcal{H}_m \quad \text{with} \quad \langle\phi_m|\psi_m\rangle = (\phi_m|S \psi_m)_m. \quad (15.1.5)
\]
We refer to \(S\) as the fermionic signature operator. It is obviously a symmetric operator. Moreover, it is bounded according to \((15.1.4)\). We conclude that it is self-adjoint. The spectral theorem gives the spectral decomposition
\[
S = \int_{\sigma(S)} \lambda dE_\lambda,
\]
where $E_\lambda$ is the spectral measure (see for example [121]). The spectral measure gives rise to the spectral calculus

$$f(S) = \int_{\sigma(S)} f(\lambda) \, dE_\lambda : \mathcal{H}_m \to \mathcal{H}_m,$$

where $f$ is a bounded Borel function on $\sigma(S) \subset \mathbb{R}$. Choosing $f$ as a characteristic function, one obtains the operators $\chi_{(0,\infty)}(S)$ and $\chi_{(-\infty,0)}(S)$. Their images are referred to as the positive and negative spectral subspace of $\mathcal{H}_m$, respectively. In this way, one obtains the desired decomposition of the solution space into two subspaces. We remark that the fermionic signature operator also gives a setting for doing spectral geometry and index theory with Lorentzian signature. We will not enter this subject here but refer the interested reader to the paper [55, 49].

The basic shortcoming of the above construction is that in many physically interesting spacetimes (like Minkowski space) the inequality (15.1.4) fails to be true. The idea to bypass this problem is to make use of the fact that a typical solution $\psi \in C^\infty(\mathcal{M}, S\mathcal{M}) \cap \mathcal{H}_m$ of the Dirac equation oscillates for large times. If we consider families of solutions with varying mass parameter, then the wave functions for different values of $m$ typically have different phases. Therefore, integrating over the mass parameter leads to destructive interference, giving rise to decay in time. In order to make this idea mathematically precise, one considers families of solutions $\{\psi_m\}_{m \in I}$ of the family of Dirac equations (15.1.1) with the mass parameter $m$ varying in an open interval $I$. We need to assume that $I$ does not contain the origin, because our methods for dealing with infinite lifetime do not apply in the massless case $m = 0$ (this seems no physical restriction because all known fermions in nature have a non-zero rest mass). By symmetry, it suffices to consider positive masses. Thus we choose

$$I := (m_L, m_R) \subset \mathbb{R} \quad \text{with parameters} \quad m_L, m_R > 0. \quad (15.1.6)$$

We always choose the family of solutions $\{\psi_m\}_{m \in I}$ in the class $C^\infty_{sc,0}(\mathcal{M} \times I, S\mathcal{M})$ of smooth solutions with spatially compact support in Minkowski space $\mathcal{M}$ which depend smoothly on $m$ and vanish identically for $m$ outside a compact subset of $I$. Then the “decay due to destructive interference” can be made precise by demanding that there is a constant $c > 0$ such that

$$| \langle \int_I \phi_m \, dm | \int_I \psi_m \, dm' \rangle | \leq c \int_I \|\phi_m\|_m \|\psi_m\|_m \, dm \quad (15.1.7)$$

for all families of solutions $\{\psi_m\}_{m \in I}, \{\phi_m\}_{m \in I} \in C^\infty_{sc,0}(\mathcal{M} \times I, S\mathcal{M})$. The point is that we integrate over the mass parameter before taking the spacetime inner product. Intuitively speaking, integrating over the mass parameter generates a decay of the wave function, making sure that the time integral converges. The inequality (15.1.7) is one variant of the so-called mass oscillation property. If (15.1.7) holds, we shall prove that there is a representation

$$\langle \int_I \phi_m \, dm | \int_I \psi_m \, dm' \rangle = \int_I (\phi_m | \tilde{\delta}_m \psi_m) \, dm,$$

which for every $m \in I$ uniquely defines the fermionic signature operator $\tilde{\delta}_m$. This operator is bounded and symmetric with respect to the scalar product (15.1.2). Moreover, it does not depend on the choice of the interval $I$. Now the positive and negative spectral subspaces of the operator $\tilde{\delta}_m$ again yield the desired splitting of the solution space into two subspaces.
15.2. The Mass Oscillation Properties

In a spacetime of infinite life time, the spacetime inner product $\langle \psi_m | \phi_m \rangle$ of two solutions $\psi_m, \phi_m \in \mathcal{H}_m$ is in general ill-defined, because the time integral in (15.1.3) may diverge. In order to avoid this difficulty, we shall consider families of solutions with a variable mass parameter. The so-called mass oscillation property will make sense of the spacetime integral in (15.1.3) after integrating over the mass parameter.

We consider the mass parameter in a bounded open interval $I$ (15.1.6). For a given Cauchy surface $\mathcal{N}$, we consider a function $\psi_N(x, m) \in S_x \mathcal{M}$ with $x \in \mathcal{N}$ and $m \in I$. We assume that this wave function is smooth and has compact support in both variables, $\psi_N \in C^\infty_0(\mathcal{N} \times I, \mathcal{M})$. For every $m \in I$, we let $\psi(., m)$ be the solution of the Cauchy problem for initial data $\psi_N(., m)$,

$$(D - m) \psi(x, m) = 0, \quad \psi(x, m) = \psi_N(x, m) \quad \forall x \in \mathcal{N}. \quad (15.2.1)$$

Since the solution of the Cauchy problem is smooth and depends smoothly on parameters, we know that $\psi \in C^\infty(\mathcal{M} \times I, \mathcal{M})$. Moreover, due to finite propagation speed, $\psi(., m)$ has spatially compact support. Finally, the solution is clearly compactly supported in the mass parameter $m$. We summarize these properties by writing

$$\psi \in C^\infty_{\text{sc}, 0}(\mathcal{M} \times I, \mathcal{M}), \quad (15.2.2)$$

where $C^\infty_{\text{sc}, 0}(\mathcal{M} \times I, \mathcal{M})$ denotes the smooth wave functions with spatially compact support which are also compactly supported in $I$. We often denote the dependence on $m$ by a subscript, $\psi_m(x) := \psi(x, m)$. Then for any fixed $m$, we can take the scalar product (15.1.2). On families of solutions $\psi, \phi \in C^\infty_{\text{sc}, 0}(\mathcal{M} \times I, \mathcal{M})$ of (15.2.1), we introduce a scalar product by integrating over the mass parameter,

$$(\psi | \phi) := \int_I \langle \psi_m | \phi_m \rangle_m \ dm \quad (15.2.3)$$

(where $dm$ is the Lebesgue measure). Forming the completion gives the Hilbert space $(\mathcal{H}, (\cdot | \cdot))$. It consists of measurable functions $\psi(x, m)$ such that for almost all $m \in I$, the function $\psi(., m)$ is a weak solution of the Dirac equation which is square integrable over any Cauchy surface. Moreover, this spatial integral is integrable over $m \in I$, so that the scalar product (15.2.3) is well-defined. We denote the norm on $\mathcal{H}$ by $\| \cdot \|$.

For the applications, it is useful to introduce a subspace of the solutions of the form (15.2.2):

**Definition 15.2.1.** We let $\mathcal{H}^\infty \subset C^\infty_{\text{sc}, 0}(\mathcal{M} \times I, \mathcal{M}) \cap \mathcal{H}$ be a subspace of the smooth solutions with the following properties:

(i) $\mathcal{H}^\infty$ is invariant under multiplication by smooth functions in the mass parameter,

$$\eta(m) \psi(x, m) \in \mathcal{H}^\infty \quad \forall \psi \in \mathcal{H}^\infty, \eta \in C^\infty(I).$$

(ii) For every $m \in I$, the set $\mathcal{H}_m := \{ \psi(., m) \ | \ \psi \in \mathcal{H}^\infty \}$ is a dense subspace of $\mathcal{H}_m$,

$$\overline{\mathcal{H}_m^{\infty}(\cdot)} = \mathcal{H}_m \quad \forall m \in I.$$  

We refer to $\mathcal{H}^\infty$ as the domain for the mass oscillation property.

The simplest choice is to set $\mathcal{H}^\infty = C^\infty_{\text{sc}, 0}(\mathcal{M} \times I, \mathcal{M}) \cap \mathcal{H}$, but in some applications it is preferable to choose $\mathcal{H}^\infty$ as a proper subspace of $C^\infty_{\text{sc}, 0}(\mathcal{M} \times I, \mathcal{M}) \cap \mathcal{H}$.

Our motivation for considering a variable mass parameter is that integrating over the mass parameter should improve the decay properties of the wave function for large times (similar as explained in the introduction in the vacuum Minkowski space). This decay
for large times should also make it possible to integrate the Dirac operator in the inner product \([15.1.3]\) by parts without boundary terms,
\[
\langle Dp\psi | p\phi \rangle = \langle p\psi | Dp\phi \rangle ,
\]
implying that the solutions for different mass parameters should be orthogonal with respect to this inner product. Instead of acting with the Dirac operator, it is technically easier to work with the operator of multiplication by \(m\), which we denote by
\[
T : \mathcal{H} \to \mathcal{H} , \quad (T\psi)_m = m\psi_m .
\]
In view of property (i) in Definition \([15.2.1]\) this operator leaves \(\mathcal{H}^\infty\) invariant,
\[
T|_{\mathcal{H}^\infty} : \mathcal{H}^\infty \to \mathcal{H}^\infty .
\]
Moreover, \(T\) is a symmetric operator, and it is bounded because the interval \(I\) is,
\[
T^* = T \in \mathcal{L}(\mathcal{H}) . \tag{15.2.4}
\]
Finally, integrating over \(m\) gives the operation
\[
p : \mathcal{H}^\infty \to C^\infty_{sc}(M,S\mathcal{M}) , \quad p\psi = \int_I \psi_m \, dm . \tag{15.2.5}
\]
We point out for clarity that \(p\psi\) no longer satisfies a Dirac equation. The following notions were introduced in \([77]\), and we refer the reader to this paper for more details.

**Definition 15.2.2.** The Dirac operator \(D = i\partial + \mathcal{B}\) on Minkowski space \(\mathcal{M}\) has the **weak mass oscillation property** in the interval \(I = (m_L,m_R)\) with domain \(\mathcal{H}^\infty\) if the following conditions hold:

(a) For every \(\psi,\phi \in \mathcal{H}^\infty\), the function \(\langle p\phi | p\psi \rangle\) is integrable on \(\mathcal{M}\). Moreover, there is a constant \(c = c(\psi)\) such that
\[
|\langle p\phi | p\psi \rangle| \leq c \|\phi\| \quad \text{for all } \phi \in \mathcal{H}^\infty . \tag{15.2.6}
\]

(b) For all \(\psi,\phi \in \mathcal{H}^\infty\),
\[
\langle pT\psi | p\phi \rangle = \langle p\psi | pT\phi \rangle . \tag{15.2.7}
\]

**Definition 15.2.3.** The Dirac operator \(D = i\partial + \mathcal{B}\) on Minkowski space \(\mathcal{M}\) has the **strong mass oscillation property** in the interval \(I = (m_L,m_R)\) with domain \(\mathcal{H}^\infty\) if there is a constant \(c > 0\) such that
\[
|\langle p\psi | p\phi \rangle| \leq c \int_I \|\phi_m\|_m \|\psi_m\|_m \, dm \quad \text{for all } \psi,\phi \in \mathcal{H}^\infty . \tag{15.2.8}
\]

**15.3. The Fermionic Signature Operator**

In this section we give abstract constructions based on the mass oscillation property. We first assume that the **weak mass oscillation property** of Definition \([15.2.2]\) holds. Then, in view of the inequality \([15.2.6]\), every \(\psi \in \mathcal{H}^\infty\) gives rise to a bounded linear functional on \(\mathcal{H}\). By continuity, this linear functional can be uniquely extended to \(\mathcal{H}\). The Fréchet-Riesz theorem allows us to represent this linear functional by a vector \(u \in \mathcal{H}\), i.e.
\[
(u|\phi) = \langle p\psi | p\phi \rangle \quad \forall \phi \in \mathcal{H} .
\]
Varying \(\psi\), we obtain the linear mapping
\[
S : \mathcal{H}^\infty \to \mathcal{H} , \quad (S\psi|\phi) = \langle p\psi | p\phi \rangle \quad \forall \phi \in \mathcal{H} .
\]
This operator is symmetric because
\[
(S\psi|\phi) = \langle p\psi | p\phi \rangle = \langle \psi | S\phi \rangle \quad \forall \phi,\psi \in \mathcal{H}^\infty .
\]
Moreover, (15.2.7) implies that the operators $S$ and $T$ commute,

$$ST = TS : \mathcal{H}^\infty \rightarrow \mathcal{H}.$$  (15.3.1)

Thus the weak mass oscillation property makes it possible to introduce $S$ as a densely defined symmetric operator on $\mathcal{H}^\infty$. It is indeed possible to construct a self-adjoint extension of the operator $S^2$ (using the Friedrich’s extension), giving rise to a functional calculus with corresponding spectral measure (for details see [77, Section 3]). In this setting the operator $S$ and the spectral measure are operators on the Hilbert space $\mathcal{H}$ which involves an integration over the mass parameter. In simple terms, this implies that all objects are defined only for almost all values of $m$ (with respect to the Lebesgue measure on $I \subset \mathbb{R}$), and they can be modified arbitrarily on subsets of $I$ of measure zero. But it does not seem possible to “evaluate pointwise in the mass” by constructing operators $S_m$ which act on the Hilbert space $\mathcal{H}_m$ for fixed mass.

In view of this shortcoming, we shall not enter the spectral calculus based on the weak mass oscillation operator. Instead, we move on to the strong mass oscillation property, which makes life much easier because it implies that $S$ is a bounded operator.

**Theorem 15.3.1.** The following statements are equivalent:

(i) The strong mass oscillation property holds.

(ii) There is a constant $c > 0$ such that for all $\psi, \phi \in \mathcal{H}^\infty$, the following two relations hold:

$$|<p\psi|p\phi>| \leq c \|\psi\| \|\phi\| \quad (15.3.2)$$

$$<pT\psi|p\phi> = <p\psi|pT\phi>.$$  (15.3.3)

(iii) There is a family of linear operators $S_m \in L(\mathcal{H}_m)$ which are uniformly bounded,

$$\sup_{m \in I} \|S_m\| < \infty,$$

such that

$$<p\psi|p\phi> = \int_I (\psi_m | S_m \phi_m)_m \, dm \quad \forall \psi, \phi \in \mathcal{H}^\infty.$$  (15.3.4)

**Proof.** The implication (iii)$\Rightarrow$(i) follows immediately from the estimate

$$|<p\psi|p\phi>| \leq \int_I |(\psi_m | S_m \phi_m)_m| \, dm \leq \sup_{m \in I} \|S_m\| \int_I \|\psi_m\| \|\phi_m\| \, dm .$$

In order to prove the implication (i)$\Rightarrow$(ii), we first apply the Schwarz inequality to (15.2.8) to obtain

$$|<p\psi|p\phi>| \leq c \int_I \|\phi_m\|_m \|\psi_m\|_m \, dm$$

$$\leq c \left( \int_I \|\phi_m\|_m^2 \, dm \right)^{1/2} \left( \int_I \|\psi_m\|_m^2 \, dm \right)^{1/2} = c \|\phi\| \|\psi\| ,$$

proving (15.3.2). Next, for given $N \in \mathbb{N}$ we subdivide the interval $I = (m_L, m_R)$ by choosing the intermediate points

$$m_\ell = \frac{\ell}{N} (m_R - m_L) + m_L, \quad \ell = 0, \ldots, N.$$
Moreover, we choose non-negative test functions \( \eta_1, \ldots, \eta_N \in C_0^\infty(\mathbb{R}) \) which form a partition of unity and are supported in small sub-intervals, meaning that

\[
\sum_{\ell=1}^N \eta_\ell|_I = 1|_I \quad \text{and} \quad \text{supp } \eta_\ell \subset (m_{\ell-1}, m_{\ell+1}) ,
\]

where we set \( m_{-1} = m_L - 1 \) and \( m_{N+1} = m_R + 1 \). For any smooth function \( \eta \in C_0^\infty(\mathbb{R}) \) we define the bounded linear operator \( \eta(T) : \mathcal{H}^\infty \rightarrow \mathcal{H}^\infty \) by

\[
(\eta(T)\psi)_m = \eta(m) \psi_m .
\]

Then by linearity,

\[
\langle pT\psi|p\phi \rangle - \langle p\psi|pT\phi \rangle = \sum_{\ell,\ell'=1}^N \left( \langle pT \eta_\ell(T) \psi \mid p \eta_{\ell'}(T) \phi \rangle - \langle p \eta_\ell(T) \psi \mid p T \eta_{\ell'}(T) \phi \rangle \right)
\]

\[
= \sum_{\ell,\ell'=1}^N \left( \langle p (T - m_\ell) \eta_\ell(T) \psi \mid p \eta_{\ell'}(T) \phi \rangle - \langle p \eta_\ell(T) \psi \mid p (T - m_\ell) \eta_{\ell'}(T) \phi \rangle \right).
\]

Taking the absolute value and applying (15.2.8), we obtain

\[
|\langle pT\psi|p\phi \rangle - \langle p\psi|pT\phi \rangle| \leq c \sum_{\ell,\ell'=1}^N \int_I |m - m_\ell| \eta_\ell(m) \eta_{\ell'}(m) ||\phi_m||_m ||\psi_m||_m dm .
\]

In view of the second property in (15.3.5), we only get a contribution if \( |\ell - \ell'| \leq 1 \). Moreover, we know that \( |m - m_\ell| \leq 2 |I|/N \) on the support of \( \eta_\ell \). Thus

\[
|\langle pT\psi|p\phi \rangle - \langle p\psi|pT\phi \rangle| \leq \frac{6c |I|}{N} \sum_{\ell=1}^N \int_I \eta_\ell(m) ||\phi_m||_m ||\psi_m||_m dm
\]

\[
= \frac{6c |I|}{N} \int_I ||\phi_m||_m ||\psi_m||_m dm .
\]

Since \( N \) is arbitrary, we obtain (15.3.3).

It remains to prove the implication (ii) \( \Rightarrow \) (iii). Combining (15.3.2) with the Fréchet-Riesz theorem, there is a bounded operator \( S \in L(\mathcal{H}) \) with

\[
\langle p\psi|p\phi \rangle = (\psi|S\phi) \quad \forall \psi, \phi \in \mathcal{H}^\infty .
\]

(15.3.6)

The relation (15.3.3) implies that the operators \( S \) and \( T \) commute. Moreover, these two operators are obviously symmetric and thus self-adjoint. Hence the spectral theorem for commuting self-adjoint operators implies that there is a spectral measure \( F \) on \( \sigma(S) \times I \) such that

\[
S^p T^q = \int_{\sigma(S) \times I} \nu^p m^q dF_{\nu,m} \quad \forall p, q \in \mathbb{N} .
\]

(15.3.7)

For given \( \psi, \phi \in \mathcal{H}^\infty \), we introduce the Borel measure \( \mu_{\psi,\phi} \) on \( I \) by

\[
\mu_{\psi,\phi}(\Omega) = \int_{\sigma(S) \times \Omega} \nu d(\psi|F_{\nu,m}\phi) .
\]

(15.3.8)

Then \( \mu_{\psi,\phi}(I) = (\psi|S\phi) \) and

\[
\mu_{\psi,\phi}(\Omega) = \int_{\sigma(S) \times I} \nu d(\chi_{\Omega}(T) \psi \mid F_{\nu,m} \chi_{\Omega}(T) \phi) = (\chi_{\Omega}(T) \psi \mid S \chi_{\Omega}(T) \phi) .
\]
Since the operator $S$ is bounded, we conclude that
\[ |\mu_{\psi,\phi}(\Omega)| \leq c \|\chi_\Omega(T)\psi\| \|\chi_\Omega(T)\phi\| \leq c \left( \int_\Omega \|\psi\|_m^2 \, dm \int_\Omega \|\phi\|_{m'}^2 \, dm' \right)^{\frac{1}{2}}, \]  
(15.3.9)

This shows that the measure $\mu$ is absolutely continuous with respect to the Lebesgue measure. The Radon-Nikodym theorem (see Theorem 12.5.2) implies that there is a unique function $f_{\psi,\phi} \in L^1(I, dm)$ such that
\[ \mu_{\psi,\phi}(\Omega) = \int_I f_{\psi,\phi}(m) \, dm. \]  
(15.3.10)

Using this representation in (15.3.9), we conclude that for any $\varphi \in \mathbb{R}$,
\[ \text{Re} \left( e^{i\varphi} \int_\Omega f_{\psi,\phi}(m) \, dm \right) \leq |\mu_{\psi,\phi}(\Omega)| \leq c |\Omega| \left( \sup_{m \in \Omega} \|\psi_m\|_m \right) \left( \sup_{m' \in \Omega} \|\phi_{m'}\|_{m'} \right). \]

As a consequence, for almost all $m \in I$ (with respect to the Lebesgue measure $dm$),
\[ \text{Re} \left( e^{i\varphi} f_{\psi,\phi}(m) \right) \leq c \|\psi_m\|_m \|\phi_{m'}\|_{m'} \]

Since the phase factor is arbitrary, we obtain the pointwise bound
\[ |f_{\psi,\phi}(m)| \leq c \|\psi_m\|_m \|\phi_{m'}\|_{m'} \quad \text{for almost all } m \in I. \]

Using this inequality, we can apply the Fréchet-Riesz theorem to obtain a unique operator $\tilde{S}_m \in L(H_m)$ such that
\[ f_{\psi,\phi}(m) = (\psi_m|\tilde{S}_m\phi_m)_m \quad \text{and} \quad \|\tilde{S}_m\| \leq c. \]  
(15.3.11)

Combining the above results, for any $\psi, \phi \in H^\infty$, we obtain
\[ <p\psi|p\phi> = (\psi|S\phi) = \int_{\sigma(S) \times I} \nu \, d\psi(\phi) = \int_I \int_I d\mu_{\psi,\phi} = \int_I f_{\psi,\phi}(m) \, dm = \int_I (\psi_m|\tilde{S}_m\phi_m)_m \, dm. \]

This concludes the proof.

Comparing the statement of Theorem 15.3.1 (ii) with Definition 15.2.2, we immediately obtain the following result.

**Corollary 15.3.2.** The strong mass oscillation property implies the weak mass oscillation property.

We next show uniqueness as well as the independence of the choice of the interval $I$.

**Proposition 15.3.3.** (uniqueness of $S_m$) The family $(S_m)_{m \in I}$ in the statement of Theorem 15.3.1 can be chosen such that for all $\psi, \phi \in H^\infty$, the expectation value $f_{\psi,\phi}(m) := (\psi_m|S_m\phi_m)_m$ is continuous in $m$,
\[ f_{\psi,\phi} \in C^0(I). \]  
(15.3.12)

The family $(S_m)_{m \in I}$ with the properties (15.3.4) and (15.3.12) is unique. Moreover, choosing two intervals $\hat{I}$ and $I$ with $m \in \hat{I} \subset I$ and $0 \notin \hat{T}$, and denoting all the objects constructed in $\hat{I}$ with an additional check, we have
\[ \tilde{S}_m = S_m. \]
(15.3.13)
**Proof.** Let us show that the function \( f_{\psi,\phi} \) is continuous. To this end, we choose a function \( \eta \in C_0^\infty(I) \). Then for any \( \varepsilon > 0 \) which is so small that \( B_\varepsilon(\supp \eta) \subset I \), we obtain

\[
\int_I \left( f_{\psi,\phi}(m + \varepsilon) - f_{\psi,\phi}(m) \right) \eta(m) \, dm = \int_I f_{\psi,\phi}(m) \left( \eta(m - \varepsilon) - \eta(m) \right) \, dm \\
=s< \int_I \left( \eta(m - \varepsilon) - \eta(m) \right) \psi_m \, dm \mid p\phi> = < \int_I \eta(m) \left( \psi_{m+\varepsilon} - \psi_m \right) \, dm \mid p\phi>,
\]
where in (s) we used \( (15.3.7) \) and \( (15.3.8) \). Applying \( (15.3.2) \), we obtain

\[
\left| \int_I \left( f_{\psi,\phi}(m + \varepsilon) - f_{\psi,\phi}(m) \right) \eta(m) \, dm \right| \leq c \left\| \psi_{+\varepsilon} - \psi \right\| \sup_I |\eta|,
\]
where the vector \( \psi_{+\varepsilon} \in \mathcal{H}^\infty \) is defined by \( (\psi_{+\varepsilon})_m := \psi_{m+\varepsilon} \). Since \( \lim_{\varepsilon \searrow 0} \left\| \psi_{+\varepsilon} - \psi \right\| = 0 \) and \( \eta \) is arbitrary, we conclude that \( f_{\psi,\phi} \) is continuous \( (15.3.12) \). This continuity is important because it implies that the function \( f_{\psi,\phi} \) is uniquely defined pointwise (whereas in \( (15.3.10) \) this function could be modified arbitrarily on sets of measure zero).

In order to prove \( (15.3.13) \), we note that the representation \( (15.3.6) \) implies that

\[
(\psi | S\phi) = (\psi | \hat{S}\phi) \quad \text{for all } \psi, \phi \in \mathcal{H}^\infty.
\]

Using \( (15.3.8) \) and \( (15.3.10) \), it follows that

\[
\int_\Omega \hat{f}_{\psi,\phi}(m) \, dm = \int_\Omega f_{\psi,\phi}(m) \, dm \quad \text{for all } \Omega \subset \hat{I}.
\]
Choosing \( \hat{f}_{\psi,\phi}(m) \) and \( f_{\psi,\phi}(m) \) as continuous functions, we conclude that they coincide for every \( m \in \hat{I} \). It follows from \( (15.4.1) \) that the operators \( \hat{S}_m \) and \( S_m \) coincide. This concludes the proof. \( \square \)

### 15.4. The Unregularized Kernel of the Fermionic Projector

From Definition \( 15.3.4 \), the operator \( S_m \) is obviously symmetric. Thus the spectral theorem gives rise to the spectral decomposition

\[
S_m = \int_{\sigma(S_m)} \nu \, dE_\nu,
\]
where \( E_\nu \) is the spectral measure (see for example \( [121] \)). The spectral measure gives rise to the spectral calculus

\[
f(S_m) = \int_{\sigma(S_m)} f(\nu) \, dE_\nu,
\]
where \( f \) is a bounded Borel function.

**Definition 15.4.1.** Assume that the Dirac operator \( D \) on \( (M, g) \) satisfies the strong mass oscillation property (see Definition \( 15.2.3 \)). We define the operators \( P_\pm : C_0^\infty(M, SM) \to \mathcal{H}_m \) by

\[
P_+ = \chi_{[0,\infty)}(S_m) k_m \quad \text{and} \quad P_- = -\chi_{(-\infty,0)}(S_m) k_m \quad (15.4.1)
\]
(where \( \chi \) denotes the characteristic function). The **fermionic projector** \( P \) is defined by \( P = P_- \).
The proof for two-point distribution on \( \mathcal{M} \)

Since the norm of the operator \( \chi \) with an integral kernel

Exactly as explained in \[ \square \] and a local trivialization.

Theorem 5.2.1, keeping in mind that this theorem applies just as well to bundle-valued

\( D \) right side is continuous on \( (\mathcal{M} \times \mathcal{M}) \), the operators \( P_{\pm} \) are symmetric,

Moreover, the image of \( P_{\pm} \) is the positive respectively negative spectral subspace of \( S_m \), i.e.

\[
\begin{align*}
P_+(C_0^\infty(\mathcal{M}, S\mathcal{M})) &= E_{(0,\infty)}(\mathcal{H}_m), & P_-(C_0^\infty(\mathcal{M}, S\mathcal{M})) &= E_{(-\infty,0)}(\mathcal{H}_m).
\end{align*}
\]

Proof. According to Proposition \[ 13.4.4 \] and Definition \[ 15.4.1 \]

We define the space of distributions \( C_0^\infty(\mathcal{M}, S\mathcal{M}) \) by

\[
\begin{align*}
\left<\phi | P \psi \right> &= (k_m \phi | P \psi)_m = -\chi_{(-\infty,0)}(S_m) k_m \phi | k_m \psi)_m \\
&= -\left< k_m \phi | \chi_{(-\infty,0)}(S_m) k_m \psi \right>_m = \left< \phi | P_\psi \right>.
\end{align*}
\]

The proof for \( P_+ \) is similar. The relations \[ 15.4.2 \] follow immediately from the fact that

k_m(C_0^\infty(\mathcal{M}, S\mathcal{M})) \) is dense in \( \mathcal{H}_m \).

Proposition 15.4.2. For all \( \phi, \psi \in C_0^\infty(\mathcal{M}, S\mathcal{M}) \), the operators \( P_{\pm} \) are symmetric,

\[
\left< P_{\pm} \phi | \psi \right> = \left< \phi | P_{\pm} \psi \right>.
\]

Since the norm of the operator \( \chi \) with an integral kernel

Exactly as explained in \[ \square \] and a local trivialization.

Theorem 5.2.1, keeping in mind that this theorem applies just as well to bundle-valued

\( D \) right side is continuous on \( (\mathcal{M} \times \mathcal{M}) \), the operators \( P_{\pm} \) are symmetric,

Moreover, the image of \( P_{\pm} \) is the positive respectively negative spectral subspace of \( S_m \), i.e.

\[
\begin{align*}
P_+(C_0^\infty(\mathcal{M}, S\mathcal{M})) &= E_{(0,\infty)}(\mathcal{H}_m), & P_-(C_0^\infty(\mathcal{M}, S\mathcal{M})) &= E_{(-\infty,0)}(\mathcal{H}_m).
\end{align*}
\]

Proof. According to Proposition \[ 13.4.4 \] and Definition \[ 15.4.1 \]

\[
\begin{align*}
\left<\phi | P \psi \right> &= (k_m \phi | P \psi)_m = -\chi_{(-\infty,0)}(S_m) k_m \phi | k_m \psi)_m \\
&= -\left< k_m \phi | \chi_{(-\infty,0)}(S_m) k_m \psi \right>_m = \left< \phi | P_\psi \right>.
\end{align*}
\]

The proof for \( P_+ \) is similar. The relations \[ 15.4.2 \] follow immediately from the fact that

k_m(C_0^\infty(\mathcal{M}, S\mathcal{M})) \) is dense in \( \mathcal{H}_m \).

Similar as in \[ 76 \] Theorem 3.12, the fermionic projector can be represented by a two-point distribution on \( \mathcal{M} \). As usual, we denote the space of test functions (with the Fréchet topology) by \( D \) and define the space of distributions \( D' \) as its dual space.

Theorem 15.4.3. Assume that the strong mass oscillation property holds. Then there is a unique distribution \( P \in D'(\mathcal{M} \times \mathcal{M}) \) such that for all \( \phi, \psi \in C_0^\infty(\mathcal{M}, S\mathcal{M}) \),

\[
\left< \phi | P \psi \right> = P(\phi \otimes \psi).
\]

Proof. According to Proposition \[ 13.4.4 \] and Definition \[ 15.4.1 \]

\[
\begin{align*}
\left<\phi | P \psi \right> &= (k_m \phi | P \psi)_m = -\chi_{(-\infty,0)}(S_m) k_m \phi | k_m \psi)_m \\
&= -\left< k_m \phi | \chi_{(-\infty,0)}(S_m) k_m \psi \right>_m = \left< \phi | P_\psi \right>.
\end{align*}
\]

Since the norm of the operator \( \chi_{(-\infty,0)}(S_m) \) is bounded by one, we conclude that

\[
|\left< \phi | P \psi \right>| \leq \|k_m \phi\| \|k_m \psi\| = \left< \phi | k_m \phi \right> \left< \psi | k_m \psi \right> \frac{1}{2},
\]

where in the last step we again applied Proposition \[ 13.4.4 \] As \( k_m \in D'((\mathcal{M} \times \mathcal{M}) \), the right side is continuous on \( D((\mathcal{M} \times \mathcal{M}) \). We conclude that also \( \left< \phi | P \psi \right> \) is continuous on \( D((\mathcal{M} \times \mathcal{M}) \). The result now follows from the Schwartz kernel theorem (see \[ 100 \] Theorem 5.2.1), keeping in mind that this theorem applies just as well to bundle-valued distributions on a manifold simply by working with the components in local coordinates and a local trivialization).

Exactly as explained in \[ 76 \] Section 3.5 \], it is convenient to use the standard notation with an integral kernel \( P(x, y) \),

\[
\begin{align*}
\left<\phi | P \psi \right> &= \int_{\mathcal{M} \times \mathcal{M}} \left<\phi(x) | P(x, y) \psi(y)\right>_x d\mu_\mathcal{M}(x) d\mu_\mathcal{M}(y) \\
(P \psi)(x) &= \int_{\mathcal{M}} P(x, y) \psi(y) d\mu_\mathcal{M}(y)
\end{align*}
\]

(where \( P(.,.) \) coincides with the distribution \( P \) above). In view of Proposition \[ 15.4.2 \] we know that the last integral is not only a distribution, but a function which is square integrable over every Cauchy surface. Moreover, the symmetry of \( P \) shown in Proposition \[ 15.4.4 \] implies that

\[
P(x, y)^* = P(y, x),
\]

where the star denotes the adjoint with respect to the spin inner product. Finally, exactly as shown in \[ 76 \] Proposition 3.13 \], the spatial normalization property of Proposition \[ 15.4.4 \] makes it possible to obtain a representation of the fermionic projector in
15.4. THE UNREGULARIZED KERNEL OF THE FERMIONIC PROJECTOR

terms of one-particle states. To this end, one chooses an orthonormal basis \((\psi_j)_j \in \mathbb{N}\) of the subspace \(\chi_{(-\infty, 0)}(S_m) \subset \mathcal{H}_m\). Then

\[
P(x, y) = -\sum_{j=1}^{\infty} |\psi_j(x)\rangle \langle \psi_j(y)|
\]

with convergence in \(D'(\mathcal{M} \times \mathcal{M})\).

We now specify the normalization of the fermionic projector. We introduce an operator \(\Pi\) by

\[
\Pi : \mathcal{H}_m \to \mathcal{H}_m, \quad (\Pi \psi_m)(x) = \frac{-2\pi}{\hat{N}} P(x, y) \psi(y) |N(y)\, d\mu_N(y),
\]

where \(N\) is any Cauchy surface.

**Proposition 15.4.4. (spatial normalization)** The operator \(\Pi\) is a projection operator on \(\mathcal{H}_m\).

**Proof.** According to Theorem 13.4.2 the spatial integral in (15.4.3) can be combined with the factor \(k_m\) in (15.4.1) to give the solution of the corresponding Cauchy problem. Thus

\[
\Pi : \mathcal{H}_m \to \mathcal{H}_m, \quad (\Pi \psi_m)(x) = \chi_{(-\infty, 0)}(S_m) \psi_m,
\]

showing that \(\Pi\) is a projection operator. \(\square\)

Instead of the spatial normalization, one could also consider the mass normalization (for details on the different normalization methods see [81], Section 2). To this end, one needs to consider families of fermionic projectors \(P_m\) indexed by the mass parameter. Then for all \(\phi, \psi \in C^\infty_0(\mathcal{M}, S\mathcal{M})\), we can use (15.3.4) and Proposition 13.4.4 to obtain

\[
<\mathbf{p}(P_m\phi) | \mathbf{p}(P_{m'}\psi)> = \int_I (P_m\phi | S_m P_m\psi)_m \, dm = \int_I (k_m\phi | S_m \chi_{(-\infty, 0)}(S_m) k_m\psi)_m \, dm = \int_I <\phi | S_m \chi_{(-\infty, 0)}(S_m) k_m\psi> \, dm = -<\phi | \mathbf{p}(S_m P_m\psi)>,
\]

which can be written in a compact formal notation as

\[
P_m P_{m'} = \delta(m - m') (-S_m) P_m.
\]

Due to the factor \((-S_m)\) on the right, in general the fermionic projector does not satisfy the mass normalization condition. The mass normalization condition could be arranged by modifying the definition (15.4.1) to

\[
S_m^{-1} \chi_{(-\infty, 0)}(S_m) k_m.
\]

Here we prefer to work with the spatial normalization. For a detailed discussion of the different normalization methods we refer to [81] Section 2.

We finally remark that corresponding causal fermion systems can be constructed exactly as in [76] Section 4] by introducing regularization operators \(\mathcal{R}_\epsilon\) \(\epsilon > 0\), computing the local correlation operators \(F^\epsilon(x)\) and defining the universal measure by \(d\rho = F^\epsilon d\mu_{\mathcal{M}}\).
Exercises

Exercise 15.1. This exercise illustrates the mass oscillation property. Let \( 0 < m_L < m_R \) and \( \eta \in C_0^\infty((m_L, m_R)) \). Show that the function \( f \) given by
\[
f(t) = \int_{m_L}^{m_R} \eta(m) e^{-i \sqrt{1 + m^2} t} \, dm
\]
has rapid decay. Does this result remain valid if \( m_L \) and \( m_R \) are chosen to have opposite signs? Justify your finding by a proof or a counter example.

Exercise 15.2. Let \( \mathcal{M} \) be the “spacetime strip”
\[
\mathcal{M} = \{(t, \vec{x}) \in \mathbb{R}^{1,3} \text{ with } 0 < t < T\}
\]
Show that for any solution \( \psi \in C_\text{sc}^\infty(\mathcal{M}, S\mathcal{M}) \cap \mathcal{H}_m \) of the Dirac equation, the following inequality holds,
\[
|\langle \psi | \phi \rangle| \leq T \| \psi \|_m \| \phi \|_m.
\]
This estimate illustrates how in spacetimes of finite lifetime, the spacetime inner product is a bounded sesquilinear form on \( \mathcal{H}_m \).

Exercise 15.3. Let \( \mathcal{M} \) again be the “spacetime strip” of the previous exercise. As in the lecture, let \( \psi, \phi \in \mathcal{H} \cap C_\text{sc}^\infty(\mathcal{M} \times I, S\mathcal{M}) \) be families of smooth Dirac solutions of spatially compact support, with compact support in the mass parameter. Moreover, we again define the operators \( p, T : \mathcal{H} \to \mathcal{H} \) as in the lecture. Does the equation
\[
\langle pT \psi | p \phi \rangle = \langle p \psi | pT \phi \rangle
\]
(which appears in the weak mass oscillation property) in general hold? Justify your answer by a proof or a counter example.

Exercise 15.4. Let \( \mathcal{M} \) again be the “spacetime strip” of the previous exercises. Moreover, as in Exercise 6.5 we again let \( \mathcal{H} \subset \mathcal{H}_m \) be a finite-dimensional subspace of the Dirac solution space \( \mathcal{H}_m \), consisting of smooth wave functions of spatially compact support, i.e.
\[
\mathcal{H} \subset C_\text{sc}^\infty(\mathcal{M}, S\mathcal{M}) \cap H_m \quad \text{finite-dimensional}.
\]
Show that the fermionic signature operator \( S \in L(\mathcal{H}) \) defined by
\[
\langle \psi | \phi \rangle = (\psi | S\phi)_m \quad \text{for all } \psi, \phi \in \mathcal{H}
\]
can be expressed within the causal fermion system by
\[
S = - \int_M x \, d\rho(x)
\]
(where \( \rho \) is again the push-forward of \( d\mu_{\mathcal{M}} \)).

Exercise 15.5. Let \( E \) be the Banach space \( E = C^0([0, 1], \mathbb{C}) \) and \( \Lambda : E \times E \to \mathbb{C} \) be sesquilinear, bounded and positive semi-definite.

(a) Assume that \( \Lambda \) satisfies for a suitable constant \( c > 0 \) and all \( f, g \in E \) the inequality
\[
|\Lambda(f, g)| \leq c \sup_{x \in [0,1]} |f(x) g(x)|.
\]
Show that there is a regular bounded Borel measure \( \mu \) such that
\[
\Lambda(f, g) = \int_0^1 \overline{f(x)} g(x) \, d\mu(x).
\]
(b) Now make the stronger assumption that $\Lambda$ satisfies for a suitable constant $\tilde{c} > 0$ and all $f, g \in E$ the inequality

$$|\Lambda(f, g)| \leq \tilde{c} \int_0^1 |f(x) g(x)| \, dx.$$  \hfill (15.4.5)

Show that $\mu$ is absolutely continuous w.r.t. to the Lebesgue measure. Show that there is a non-negative function $h \in L^1([0, 1], dx)$ such that

$$\Lambda(f, g) = \int_0^1 f(x) g(x) h(x) \, dx.$$  \hfill (15.4.6)

Show that $h$ is pointwise bounded by $c$.

(c) In order to clarify the different assumptions in this exercise, give an example for a sesquilinear, bounded and positive semi-definite functional $\Lambda$ which violates (15.4.4). Give an example which satisfies (15.4.4) but violates (15.4.5).

Exercise 15.6. (Towards the mass oscillation property - part 1) This exercise illustrates the mass oscillation property. Let $0 < m_L < m_R$ and $\eta \in C^\infty_0((m_L, m_R))$. Show that the function $f$ given by

$$f(t) = \int_{m_L}^{m_R} \eta(m) e^{-\sqrt{1+m^2} t} \, dm$$

has rapid decay. Does this result remain valid if $m_L$ and $m_R$ are chosen to have opposite signs? Justify your finding by a proof or a counter example.

Exercise 15.7. (Towards the mass oscillation property - part 2) Let $R_T$ be the “spacetime strip”

$$R_T = \{(t, \vec{x}) \in \mathbb{R}^{1,3} : 0 < t < T\}.$$  \hfill (15.4.7)

Show that for any solutions $\psi, \phi \in C^\infty_0(\mathbb{R}^4, \mathbb{C}^4) \cap \mathcal{H}_m$ of the Dirac equation, the following inequality holds,

$$|\langle \psi | \phi \rangle| \leq T \|\psi\|_m \|\phi\|_m,$$  \hfill (15.4.8)

where $\langle \psi | \phi \rangle := \int_{R_T} \langle \psi(x) | \phi(x) \rangle \, d^4x$.

This estimate illustrates how in spacetimes of finite lifetime, the spacetime inner product is a bounded sesquilinear form on $\mathcal{H}_m$.

Exercise 15.8. (Towards the mass oscillation property - part 3) Let $R_T$ again be the “spacetime strip” of the previous exercises. Moreover, we again let $\mathcal{H} \subset \mathcal{H}_m$ be a finite-dimensional subspace of the Dirac solution space $\mathcal{H}_m$, consisting of smooth wave functions of spatially compact support, i.e.

$$\mathcal{H} \subset C^\infty_0(\mathbb{R}^4, \mathbb{C}^4) \cap \mathcal{H}_m \quad \text{finite-dimensional}.$$  \hfill (15.4.9)

Show that the fermionic signature operator $\mathcal{S} \in L(\mathcal{H})$ defined by

$$\langle \psi | \phi \rangle = (\psi | S \phi)_m \quad \text{for all } \psi, \phi \in \mathcal{H}$$

can be expressed within the causal fermion system by

$$\mathcal{S} = -\int_{R_T} x \, d\rho(x)$$

(where $\rho$ is again the push-forward of $d^4x$).


Exercise 15.9. (Bonus: External field problem) In physics, the notion of “particle” and “anti-particle” is often introduced as follows: Solutions of the Dirac equation with positive frequency are called “particles” and solutions with negative frequency “anti-particles”. In this exercise, we will check in how far this makes sense. To this end, take a look at the Dirac equation in an external field:

\[(i\partial_t + \mathcal{B} - m)\psi = 0.\]  

(15.4.6)

Assume that \(\mathcal{B}\) is time-dependent and has the following form:

\[B(t,x) = V \Theta(t - t_0)\Theta(t_1 - t),\]

where \(V \in \mathbb{R}\), \(\Theta\) denotes the Heaviside step function and \(t_0 = 0, t_1 = 1\). In order to construct a solution thereof, for a given momentum \(\vec{k}\), we use plane wave solutions of the Dirac equation,

\[\psi(t, \vec{x}) = e^{-i\omega t + i\vec{k}\vec{x}}\chi_{\vec{k}},\]

where \(\chi_{\vec{k}}\) is a spinor \(\in \mathbb{C}^4\), and patch them together suitably. (The quantity \(\omega\) is called the “frequency” or “energy”, and \(\vec{k}\) the “momentum”.) To simplify the calculation, we set \(\vec{k} = (k_1, 0, 0)^T\). Proceed as follows:

(a) First, take a look at the region \(t < t_0\). Reformulate (15.4.6) such that there is only the time derivative on the left hand side. (Hint: Multiply by \(\gamma_0\).)

(b) Insert the plane wave ansatz with \(\vec{k} = (k_1, 0, 0)^T\) into the equation. Your equation now has the form \(\omega \psi = H(k_1)\psi\). Show that the eigenvalues of \(H(k_1)\) are \(\pm \omega_0\) with \(\omega_0 : = \sqrt{(k_1)^2 + m^2}\).

(c) Show that one eigenvector belonging to \(+ \omega_0\) is \(\chi_0^+ := \frac{m + \omega_0}{k_1}0, 0, 1)^T\) and that one eigenvector belonging to \(- \omega_0\) is \(\chi_0^- := \frac{m - \omega_0}{k_1}0, 0, 1)^T\). (Both eigenvalues have multiplicity 2, but we don’t need the other two eigenvectors here.)

(d) With this, you have constructed plane wave solutions \(e^{-i(\pm \omega_0)t + i\vec{k}\vec{x}}\chi_{\vec{k}}^\pm\) for \(t < t_0\) and also for \(t > t_1\). By transforming \(m \rightarrow (m - V)\), you immediately obtain plane wave solutions also for \(t_0 < t < t_1\). Denote the respective quantities by \(\omega_1\) and \(\chi_{\vec{k}}^\pm\).

(e) Assume that for \(t < t_0\) there is one “particle” present, i.e. set

\[\psi(t, \vec{x}) = e^{-i\omega_0 t + i\vec{k}\vec{x}}\chi_{\vec{k}}^+\]  

for \(t < t_0\).

Assume that the solution for \(t_0 < t < t_1\) takes the form

\[Ae^{-i\omega_1 t + i\vec{k}\vec{x}}\chi_{\vec{k}}^+ + Be^{-i(-\omega_1)t + i\vec{k}\vec{x}}\chi_{\vec{k}}^-\]  

with \(A, B \in \mathbb{R}\).

Calculate \(A\) and \(B\) for the case \(k_1 = 1\) and \(V = m\) by demanding continuity of the solution at \(t = t_0\).

(f) Assume that for \(t > t_1\) the solution takes the form

\[Ce^{-i\omega_1 t + i\vec{k}\vec{x}}\chi_{\vec{k}}^+ + De^{-i(-\omega_1)t + i\vec{k}\vec{x}}\chi_{\vec{k}}^-\]  

with \(C, D \in \mathbb{C}\).

Calculate \(C\) and \(D\) for \(m = 2\) by demanding continuity of the solution at \(t = t_1\) (here you may want to use computer algebra).

(g) Interpret what you have found. Why could this be called the “external field problem”? 

16.1. The Causal Green’s Operators

We already encountered Green’s operators in Chapter 13 when solving the Cauchy problem with methods of hyperbolic partial differential equations (see Theorem 13.4.3). In Minkowski space, these Green’s operators can be computed in more detail with Fourier methods. Our starting point is the definition of the Green’s function \( s_m(x,y) \) of the vacuum Dirac equation by the distributional equation

\[
(i\partial_x - m) s_m(x,y) = \delta^4(x-y),
\]

(16.1.1)

where \( \delta^4(x,y) \) denotes the 4-dimensional Dirac distribution. Taking the Fourier transform of (16.1.1),

\[
s_m(x,y) = \frac{d^4k}{(2\pi)^4} \hat{s}_m(k) e^{-ik(x-y)}
\]

(16.1.2)

(where \( x, y \in \mathcal{M} \) are spacetime points, \( k \) is the four-momentum, and \( k(x-y) \) denotes the Minkowski inner product) we obtain the algebraic equation

\[
(\slashed{k} - m) s_m(k) = 1.
\]

(16.1.3)

Multiplying by \( \slashed{k} + m \) and using the identity \( (\slashed{k} - m)(\slashed{k} + m) = k^2 - m^2 \), one sees that if \( k^2 \neq m^2 \), the matrix \( \slashed{k} - m \) is invertible. If conversely \( k^2 = m^2 \), we have \( (\slashed{k} - m)^2 = -2m(\slashed{k} - m) \), which shows that the matrix \( \slashed{k} - m \) is diagonalizable with eigenvalues \( -2m \) and zero. Since the Dirac matrices \( (1.3.3) \) are trace-free, we have \( \text{Tr}(\slashed{k} - m) = -4m \). It follows that the matrix \( \slashed{k} - m \) has a two-dimensional kernel if \( k \) is on the mass shell. This shows that the Green’s function of the Dirac equation is not unique. If we add to it any vector in the kernel of \( \slashed{k} - m \), i.e. if we add to it a solution of the homogeneous Dirac equation, it still satisfies the defining equation (16.1.1) (for details see [14].)

A convenient method for solving the equation (16.1.3) for \( s_m(k) \) is to use a \( \pm i\varepsilon \)-regularization on the mass shell. Common choices are the advanced and the retarded Green’s functions, which are defined by

\[
s_m^\vee(k) = \lim_{\varepsilon \searrow 0} \frac{\slashed{k} + m}{\varepsilon^2 - (k^2 - m^2 - i\varepsilon k^0)} \quad \text{and} \quad s_m^\wedge(k) = \lim_{\varepsilon \searrow 0} \frac{\slashed{k} + m}{\varepsilon^2 - (k^2 - m^2 + i\varepsilon k^0)},
\]

(16.1.4)

respectively (with the limit \( \varepsilon \searrow 0 \) taken in the distributional sense). Computing their Fourier transform (27.3.1), one sees that they are causal in the sense that their supports lie in the upper and lower light cone, respectively,

\[
\text{supp} s_m^\vee(x,.) \subset J_x^\vee, \quad \text{supp} s_m^\wedge(x,.) \subset J_x^\wedge.
\]

(16.1.5)

Mathematically, the formulas in (16.1.4) define the Green’s functions in momentum space as tempered distributions. Taking their Fourier transform (27.3.1), the advanced and retarded Green’s functions are tempered distributions in the variable \( \xi := y - x \). We
also regard these distributions as integral kernels of corresponding operators on the wave functions, i.e.

$$(s_m(\psi))(x) := \int_{\mathcal{M}} s_m(x, y) \psi(y) \, d^4y.$$ 

We thus obtain operators

$$s^\wedge_m, s^\vee_m : C_0^\infty(\mathcal{M}, S\mathcal{M}) \to C_\infty(\mathcal{M}, S\mathcal{M}) .$$

(16.1.6)

Here $C_0^\infty(\mathcal{M}, S\mathcal{M})$ denote the smooth functions with compact support in $\mathcal{M}$, taking values in the spinors, and $C_\infty$ denotes the smooth functions with spatially compact support.

### 16.2. The Causal Fundamental Solution and Time Evolution

We now state a few properties of the Green’s operators and explain why they are useful. The considerations in this section are valid more generally in the presence of an external potential. Then the defining equation of the Green’s function (16.1.1) is modified similar to (1.3.14) to

$$(i\partial_x + \mathcal{B} - m) s_m(x, y) = \delta^4(x - y),$$

(16.2.1)

Then the existence of Green’s functions can no longer be proven by Fourier transformation. Instead, one can use methods of hyperbolic PDEs (symmetric hyperbolic systems) which we learned in the PDE2 course. Here we shall not enter these methods again. Instead, we simply assume that we are given advanced and retarded Green’s functions.

The causal fundamental solution $k_m$ is defined as the difference of the advanced and the retarded Green’s function,

$$k_m(x, y) := \frac{1}{2\pi i} \left( s^\vee_m(x, y) - s^\wedge_m(x, y) \right).$$

(16.2.2)

It is a distribution which is causal in the sense that it vanishes if $x$ and $y$ have spacelike separation. Moreover, it is a distributional solution of the homogeneous Dirac equation,

$$(i\partial_x + \mathcal{B} - m) k_m(x, y) = 0.$$ 

The unique solvability of the Cauchy problem allows us to introduce the time evolution operator of the Dirac equation as follows. Solving the Cauchy problem with initial data at time $t$ and evaluating the solution with the help of Proposition 13.6.1 at some other time $t'$ gives rise to a mapping

$$U^{t', t} : \mathcal{H}_t \to \mathcal{H}_{t'},$$

which we call time-evolution operator. Since the scalar product (15.1.2) is time independent, the operator $U^{t', t}$ is unitary, and since the Cauchy problem can be solved forwards and backwards in time, the unitary time evolution operators form a representation of the group $(\mathbb{R}, +)$. More precisely,

$$U^{t, t} = 1 \quad \text{and} \quad U^{t', t'} U^{t', t} = U^{t', t}.$$ 

Proposition 13.6.1 immediately gives the following representation of $U^{t', t}$:

$$(U^{t', t} \psi)(\vec{y}) = \int_{\mathbb{R}^3} U^{t', t}(\vec{y}, \vec{x}) \psi(t, \vec{x}) \, d^3x ,$$

(16.2.3)

where the kernel $U^{t', t}(\vec{y}, \vec{x})$ is defined as

$$U^{t', t}(\vec{y}, \vec{x}) = 2\pi k_m((t', \vec{y}), (t, \vec{x})) \gamma^0 .$$

(16.2.4)
16.3. Proof of the Weak Mass Oscillation Property in the Minkowski Vacuum

In the remainder of this chapter, we return to the Dirac equation in Minkowski space \([16.1.1]\). An external potential will be considered in the next chapter (Chapter \([17]\).

The mass oscillation property in the Minkowski vacuum can be proved using Fourier methods. Here we shall give two different approaches in detail. The method of the first space \((16.1.1)\). An external potential will be considered in the next chapter (Chapter \([17]\).

interval \([I]\) of constant time Cauchy hypersurfaces in a fixed reference frame \((t, \vec{x})\) and a variable mass parameter \(m\) in the interval \(I = (m_L, m_R)\) with \(m_L, m_R > 0\). The families of solutions \(\psi = (\psi_m)_{m \in I}\) of the Dirac equations \((i \not \partial - m)\psi_m = 0\) are contained in the Hilbert space \((\mathcal{H}, \langle ., . \rangle)\) with scalar product \([15.2.3]\). The subspace \(\mathcal{H}^\infty \subset \mathcal{H}\) in Definition \([15.2.1]\) is chosen as

\[
\mathcal{H}^\infty = C^{\infty}_{sc, 0}(\mathcal{M} \times I, S\mathcal{M}) \cap \mathcal{H}.
\]

For what follows, it is convenient to work with the Fourier transform in space, i.e.

\[
\hat{\psi}(t, \vec{k}) = \int_{\mathbb{R}^3} \psi(t, \vec{x}) e^{-i\vec{k}\cdot\vec{x}} \, d^3x, \quad \psi(t, \vec{x}) = \int_{\mathbb{R}^3} \frac{d^3k}{(2\pi)^3} \hat{\psi}(t, \vec{k}) e^{i\vec{k}\cdot\vec{x}}.
\]

Then a family of solutions \(\hat{\psi} \in \mathcal{H}^\infty\) has the representation

\[
\hat{\psi}_m(t, \vec{k}) = c_+(\vec{k}, m) e^{-i\omega(\vec{k}, m)t} + c_-(\vec{k}, m) e^{i\omega(\vec{k}, m)t} \quad \text{for all } m \in I
\]

with suitable spinor-valued coefficients \(c_\pm(\vec{k}, m)\) and \(\omega(\vec{k}, m) := \sqrt{\vec{k}^2 + m^2}\). Integrating over the mass parameter, we obtain a superposition of waves oscillating at different frequencies. Intuitively speaking, this leads to destructive interference for large \(t\), giving rise to decay in time. This picture can be made precise using integration by parts in \(m\), as we now explain. Integrating \([16.3.2]\) over the mass by applying the operator \(p\), \([15.2.5]\), we obtain

\[
p\hat{\psi}(t, \vec{k}) = \int_I (c_+ e^{-i\omega t} + c_- e^{i\omega t}) \, dm
\]

\[
= \int_I \frac{i}{t} \frac{e^{-i\omega t}}{\partial_m \omega} \left( c_+ \partial_m e^{-i\omega t} - c_- \partial_m e^{i\omega t} \right) \, dm
\]

\[
= - \frac{i}{t} \int_I \left[ \partial_m \left( \frac{c_+}{\partial_m \omega} \right) e^{-i\omega t} - \partial_m \left( \frac{c_-}{\partial_m \omega} \right) e^{i\omega t} \right] \, dm
\]

(we do not get boundary terms because \(\psi \in \mathcal{H}^\infty\) has compact support in \(m\)). With \(\partial_m \omega = m/\omega\), we conclude that

\[
p\hat{\psi}(t, \vec{k}) = - \frac{i}{t} \int_I \left[ \partial_m \left( \frac{\omega c_+}{m} \right) e^{-i\omega t} - \partial_m \left( \frac{\omega c_-}{m} \right) e^{i\omega t} \right] \, dm.
\]

Since the coefficients \(c_\pm\) depend smoothly on \(m\), the resulting integrand is bounded uniformly in time, giving a decay at least like \(1/t\), i.e. \(|p\hat{\psi}(t, \vec{k})| \lesssim 1/t\). Iterating this procedure, one even can prove decay rates \(\lesssim 1/t^2, 1/t^3, \ldots\). The price one pays is that higher and higher powers in \(\omega\) come up in the integrand, which means that in order for the spatial Fourier integral to exist, one needs a faster decay of \(c_\pm\) in \(|\vec{k}|\). Expressed in terms of the initial data, this means that every factor \(1/t\) gives rise to an additional
Braucht man denn hier die $W^{2,2}$-norm?
Genügt es nicht, dass man für jedes $\psi \in \mathcal{H}^\infty$ den quadratischen Abfall hat? Falls ja, vereinfache dies?

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spatial derivative acting on the initial data. This motivates the following basic estimate.

**Lemma 16.3.1.** For any $\psi \in \mathcal{H}^\infty$, there is a constant $C = C(m_L)$ such that

$$\| (p\psi) | I \|_t \leq \frac{C|I|}{1 + t^2} \sup_{m \in I} \sum_{b=0}^2 \| (\partial_{\mu}^b \psi_m) | t = 0 \|_{W^{2,2}},$$

(16.3.3)

where $\| . \|_t$ is the norm corresponding to the scalar product

$$(.,.)_t := 2\pi \int_{\mathbb{R}^3} <\cdot, \gamma^0 \cdot x> d^3 x : L^2(\mathcal{N}, \mathcal{M}) \times L^2(\mathcal{N}, \mathcal{M}) \to \mathbb{C}$$

(which is similar to (15.1.2), but now applied to wave functions which do not need to be solutions), and $\| . \|_{W^{2,2}}$ is the spatial Sobolev norm

$$\| \phi \|_{W^{2,2}}^2 := \sum_{\alpha \text{ with } |\alpha| \leq 2} \int_{\mathbb{R}^3} |\nabla^\alpha \phi(x)|^2 d^3 x,$$

(16.3.4)

where $\alpha$ is a multi-index.

The absolute value in (16.3.4) is the norm $\| . \| := \sqrt{<\cdot, \gamma^0 \cdot x>}$ on the spinors. If we again identify all spinor spaces in the Dirac representation with $\mathbb{C}^4$, this simply is the standard Euclidean norm on $\mathbb{C}^4$.

The proof of this lemma will be given later in this section. Before, we infer the weak mass oscillation property.

**Corollary 16.3.2.** The vacuum Dirac operator $i\partial$ in Minkowski space has the weak mass oscillation property with domain [16.3.1].

**Proof.** For every $\psi, \phi \in \mathcal{H}^\infty$, the Schwarz inequality gives

$$|<p\psi|p\phi>| = \frac{1}{2\pi} \left| \int_{-\infty}^{\infty} \left( (p\psi)|_t \gamma^0 (p\phi)|_t \right)_t dt \right| \leq \int_{-\infty}^{\infty} \| (p\psi)|_t \| \| (p\phi)|_t \| dt.$$

(16.3.5)

Applying Lemma [16.3.1] together with the estimate

$$\| (p\phi)|_t \|_t^2 = \int_{I \times I} \phi_m |_t \phi_{m'} |_t dm dm'$$

$$\leq \frac{1}{2} \int_{I \times I} \left( \| \phi_m \|_m^2 + \| \phi_{m'} \|_{m'}^2 \right) dm dm' = |I| \| \phi \|_t^2,$$

we obtain inequality (15.2.6) with

$$c = C|I|^{\frac{3}{2}} \sup_{m \in I} \sum_{b=0}^2 \| \partial_{\mu}^b \psi_m | t = 0 \|_{W^{2,2}} \frac{1}{1 + t^2} dt < \infty.$$  

(16.3.6)

The identity (15.2.7) follows by integrating the Dirac operator by parts,

$$<pT\psi|p\phi> = <pD\psi|p\phi> = <Dp\psi|p\phi> = \int_{\mathcal{M}} <Dp\psi|p\phi>_x d^4 x$$

(16.3.7)

In (*), we used that the Dirac operator is formally self-adjoint with respect to the inner product $<.,.>$. Moreover, we do not get boundary terms because of the time decay in Lemma [16.3.1].
The remainder of this section is devoted to the proof of Lemma 16.3.1. Using the result of Proposition 13.6.1, we can express the solution \( \psi_m \) of the Cauchy problem in terms of the causal fundamental solution \( k_m \). In order to bring \( k_m \) into a more explicit form, we use (18.2.10) together with formulas for the advanced and retarded Green’s functions. Indeed, these Green’s functions are the multiplication operators in momentum space

\[
s^\vee_m(p) = \lim_{\varepsilon \to 0} \frac{p + m}{k^2 - m^2 - i\varepsilon p^0} \quad \text{and} \quad s^\wedge_m(p) = \lim_{\varepsilon \to 0} \frac{p + m}{p^2 - m^2 + i\varepsilon k^0}
\]

(with the limit \( \varepsilon \to 0 \) taken in the distributional sense, and where the vector \( p \) is the four-momentum). We thus obtain in momentum space

\[
k_m(p) = \frac{1}{2\pi i} (p + m) \lim_{\varepsilon \to 0} \left[ \frac{1}{p^2 - m^2 - i\varepsilon p^0} - \frac{1}{p^2 - m^2 + i\varepsilon p^0} \right]
\]

(16.3.8)

It is convenient to transform spatial coordinates of the time evolution operator to momentum space. First, in the Minkowski vacuum, the time evolution operator can be represented as in (16.2.3) with an integral kernel \( U_{t,t'}(\vec{y}, \vec{x}) \) which depends only on the difference vector \( \vec{y} - \vec{x} \). We set

\[
U_{t,t'}(\vec{k}) := \hat{R}_3 U_{t,t'}(\vec{y}, 0) e^{-i\vec{k} \cdot \vec{y}} d^3 y.
\]

Combining (16.2.3) with (16.3.8) yields

\[
U_{t,t'}(\vec{k}) = \sum_{\pm} \Pi_{\pm}(\vec{k}) e^{\mp i\omega(t-t')} \int_{-\infty}^{\infty} e^{i\omega(k^0)} \epsilon(k^0) e^{-i\omega(t-t')} d\omega.
\]

Carrying out the \( \omega \)-integral, we get

\[
U_{t,t'}(\vec{k}) = \sum_{\pm} \Pi_{\pm}(\vec{k}) e^{\mp i\omega(t-t')},
\]

(16.3.9)

where we set

\[
\Pi_{\pm}(\vec{k}) := \pm \frac{1}{2\omega(\vec{k})} (k_\pm + m) \gamma^0
\]

(16.3.10)

with \( \omega(\vec{k}) = \sqrt{|\vec{k}|^2 + m^2} \) and \( k_\pm = (\pm \omega(\vec{k}), \vec{k}) \).

Moreover, applying Plancherel’s theorem, the scalar product (15.1.2) can be written in momentum space as

\[
(\psi_m | \phi_m)_m = (2\pi)^{-2} \int_{\mathbb{R}^3} \langle \psi_m(t, \vec{k}) | \gamma^0 \phi_m(t, \vec{k}) \rangle \, d^3 k.
\]
The unitarity of the time evolution operator in position space implies that the matrix \( U_{t,t'}^{m}(\vec{k}) \) is unitary (with respect to the scalar product \( \langle \cdot, \cdot \rangle_{C^2} \equiv \langle \cdot | \gamma^0 \cdot \rangle \)), meaning that its eigenvalues are on the unit circle and the corresponding eigenspaces are orthogonal. It follows that the operators \( \Pi_{\pm}(\vec{k}) \) in (27.104) are the orthogonal projection operators to the eigenspaces corresponding to the eigenvalues \( e^{\mp i\omega(t-t')} \), i.e.

\[
\gamma^0 \Pi^*_{s} \gamma^0 = \Pi_{s} \quad \text{and} \quad \Pi_{s}(\vec{k}) \Pi_{s'}(\vec{k}) = \delta_{s,s'} \Pi_{s}(\vec{k}) \quad \text{for} \quad s,s' \in \{+,-\}.
\]

(16.3.11)

(These relations can also be verified by straightforward computations using (16.3.10); see Exercise 16.9.)

The next two lemmas involve derivatives with respect to the mass parameter \( m \). For clarity, we again denote the \( m \)-dependence of the operators by the subscript \( m \).

**Lemma 16.3.3.** The time evolution operator in the vacuum satisfies the relation

\[
(t-t') U_{t,t'}^{m}(\vec{k}) = \frac{\partial}{\partial m} V_{t,t'}^{m}(\vec{k}) + W_{t,t'}^{m}(\vec{k}),
\]

(16.3.12)

where

\[
V_{t,t'}^{m}(\vec{k}) = \sum_{\pm} \frac{i}{2m} \left( k_{\pm} + m \right) \gamma^0 e^{\mp i\omega(t-t')} \quad \text{(16.3.13)}
\]

\[
W_{t,t'}^{m}(\vec{k}) = \sum_{\pm} \frac{i}{2} \left( k_{\pm} \gamma^0 + \frac{1}{m} \right) e^{\mp i\omega(t-t')} \quad \text{(16.3.14)}
\]

The operators \( V_{t,t'}^{m} \) and \( W_{t,t'}^{m} \) are estimated uniformly by

\[
\| V_{t,t'}^{m}(\vec{k}) \| + \| W_{t,t'}^{m}(\vec{k}) \| \leq C \left( 1 + \frac{|\vec{k}|}{m} \right),
\]

(16.3.15)

where the constant \( C \) is independent of \( m, \vec{k}, t \) and \( t' \) (and \( \| \cdot \| \) is any norm on the \( 2 \times 2 \)-matrices).

**Proof.** First, we generate the factor \( t-t' \) by differentiating the exponential in (27.104) with respect to \( \omega \),

\[
(t-t') U_{t,t'}^{m}(\vec{k}) = \sum_{\pm} \Pi_{\pm}(\vec{k}) \left( \pm i \frac{\partial}{\partial \omega} e^{\mp i\omega(t-t')} \right).
\]

Next, we want to rewrite the \( \omega \)-derivative as a derivative with respect to \( m \). Taking the total differential of the dispersion relation \( \omega^2 - |\vec{k}|^2 = m^2 \) for fixed \( \vec{k} \), one finds that

\[
\frac{\partial}{\partial \omega} = \frac{\omega}{m} \frac{\partial}{\partial m}.
\]

Hence

\[
(t-t') U_{t,t'}^{m} = \sum_{\pm} \Pi_{\pm} \left( \pm i \frac{\omega}{m} \frac{\partial}{\partial m} e^{\mp i\omega(t-t')} \right)
\]

\[
= \frac{\partial}{\partial m} \sum_{\pm} \left( \pm i \frac{\omega}{m} \Pi_{\pm} e^{\mp i\omega(t-t')} \right) - \sum_{\pm} \left( \frac{\partial}{\partial m} \left[ \pm i \frac{\omega}{m} \Pi_{\pm} \right] \right) e^{\mp i\omega(t-t')}.
\]

Computing the operators in the round brackets using (16.3.10) gives the identities (16.3.13) and (16.3.14). Estimating these formulas, one obtains bounds which are at most linear in \( |\vec{k}| \), proving (16.3.15). \( \square \)
This method can be iterated to generate more factors of $t - t'$. In the next lemma, we prove at least quadratic decay in time. For later use, it is preferable to formulate the result in position space.

**Lemma 16.3.4.** The time evolution operator in the vacuum has the representation

$$U_{m}^{t,t'} = \frac{1}{(t-t')^{2}} \left( \frac{\partial^{2}}{\partial m^{2}} A_{m}^{t,t'} + \frac{\partial}{\partial m} B_{m}^{t,t'} + C_{m}^{t,t'} \right)$$

(16.3.17)

with operators

$$A_{m}^{t,t'}, B_{m}^{t,t'}, C_{m}^{t,t'} : W^{2,2}(\mathcal{N}_{t}, \mathcal{S}\mathcal{M}) \to L^{2}(\mathcal{N}_{t}, \mathcal{S}\mathcal{M})$$

which are bounded uniformly in time by

$$\| A_{m}^{t,t'}(\phi) \|_{t} + \| B_{m}^{t,t'}(\phi) \|_{t} + \| C_{m}^{t,t'}(\phi) \|_{t} \leq c \| \phi \|_{W^{2,2}}$$

(16.3.18)

where $c$ is a constant which depends only on $m$.

**Proof.** A straightforward computation using exactly the same methods as in Lemma 16.3.3 yields the representation

$$(t-t')^{2} U_{m}^{t,t'}(\vec{k}) = \frac{\partial^{2}}{\partial m^{2}} A_{m}^{t,t'}(\vec{k}) + \frac{\partial}{\partial m} B_{m}^{t,t'}(\vec{k}) + C_{m}^{t,t'}(\vec{k}),$$

(16.3.19)

where the operators $A_{m}^{t,t'}, B_{m}^{t,t'}, C_{m}^{t,t'}$ are bounded by

$$\| A_{m}^{t,t'}(\vec{k}) \| + \| B_{m}^{t,t'}(\vec{k}) \| + \| C_{m}^{t,t'}(\vec{k}) \| \leq \frac{C}{m} \left( 1 + \frac{\| \vec{k} \|}{m} + \frac{\| \vec{k} \|^{2}}{m^{2}} \right),$$

(16.3.20)

with a numerical constant $C > 0$. We remark that, compared to (16.3.12), the right of (16.3.20) involves an additional $1/m$. This prefactor is necessary for dimensional reasons, because the additional factor $t - t'$ in (16.3.19) (compared to (16.3.12)) brings in an additional dimension of length (and in natural units, the factor $1/m$ also has the dimension of length). The additional summand $\| \vec{k} \|^{2}/m^{2}$ in (16.3.20) can be understood from the fact that applying (16.3.16) generates a factor of $\omega/m$ which for large $\| \vec{k} \|$ scales like $\| \vec{k} \|/m$.

Translating this result to position space and keeping in mind that the vector $\vec{k}$ corresponds to the derivative $-i\vec{\nabla}$, we obtain the result.

**Proof of Lemma 16.3.1.** First of all, the Schwarz inequality gives

$$\| (p\psi)_{l} \|_{l} \leq \int_{I} \| \psi_{m} \|_{m} dm \leq \sqrt{|I|} \| \psi \|.$$

Thus it remains to show the decay for large $t$, i.e.

$$\| (p\psi)_{l} \|_{l} \leq \frac{C |I|}{t^{2}} \sum_{m=1}^{2} \sup_{b=0} \| \partial_{m}^{b} (\psi_{m})_{l=0} \|_{W^{2,2}}.$$

(16.3.21)

We apply Lemma 16.3.3 and integrate by parts in $m$ to obtain

$$(p\psi)_{l} = \int_{I} U_{m}^{t,0} \psi_{m} dm = \frac{1}{t^{2}} \int_{I} (\partial_{m}^{2} A_{m}^{t,0} + \partial_{m} B_{m}^{t,0} + C_{m}^{t,0}) \psi_{m} dm$$

$$ = \frac{1}{t^{2}} \int_{I} \left( A_{m}^{t,0} (\partial_{m} \psi_{m})_{l=0} - B_{m}^{t,0} (\partial_{m} \psi_{m})_{l=0} + C_{m}^{t,0} \psi_{m} \right) dm.$$

Taking the norm and using (16.3.18) gives (16.3.21).
We finally note that the previous estimates are not optimal for two reasons. First, the pointwise quadratic decay in (16.3.3) is more than what is needed for the convergence of the integral in (16.3.6). Second and more importantly, the Schwarz inequality (16.3.5) does not catch the optimal scaling behavior in \( \vec{k} \). This is the reason why the constant in (15.2.3) involves derivatives of \( \psi_m \) (cf. (16.3.6)), making it impossible to prove the inequality (15.2.8) which arises in the strong mass oscillation property. In order to improve the estimates, one needs to use Fourier methods both in space and time, as will be explained in the next section.

16.4. Proof of the Strong Mass Oscillation Property in the Minkowski Vacuum

**Theorem 16.4.1.** The vacuum Dirac operator in Minkowski space has the strong mass oscillation property with domain (16.3.1).

Our proof relies on a Plancherel argument in spacetime. It also provides an alternative method for establishing the weak mass oscillation property.

**Proof of Theorem 16.4.1.** Let \( \psi = (\psi_m)_{m \in I} \in \mathcal{H}^\infty \) be a family of solutions of the Dirac equation for a varying mass parameter in the Minkowski vacuum. Using Proposition 13.6.1, one can express \( \psi_m \) in terms of its values at time \( t = 0 \) by

\[
\psi_m(x) = 2\pi \int_{\mathbb{R}^3} \kappa_m(x, (0, \vec{y})) \gamma^0 \psi_m|_{t=0}(\vec{y}) \, d^3 y.
\]

We now take the Fourier transform, denoting the four-momentum by \( k \). Using (16.3.8), we obtain

\[
\hat{\psi}_m(k) = 2\pi k_m(k) \gamma^0 \hat{\psi}_m^0(\vec{k}) = 2\pi \delta(k^2 - m^2) \epsilon(k^0) (\vec{k} + m) \gamma_0 \hat{\psi}_m^0(\vec{k}),
\]

where \( \hat{\psi}_m^0(\vec{k}) \) denotes the spatial Fourier transform of \( \psi_m|_{t=0} \) (in order to avoid an ambiguity of notation, the hat of the Fourier transform in spacetime was omitted). Obviously, this is a distribution supported on the mass shell. In particular, it is not square integrable over \( \mathbb{R}^4 \).

Integrating over \( m \), we obtain the following function

\[
(p \psi)(k) = 2\pi \chi_I(m) \frac{1}{2m} \epsilon(k^0) (\vec{k} + m) \gamma^0 \hat{\psi}_m^0(\vec{k}) \bigg|_{m=\sqrt{k^2}},
\]

where \( m \) now is a function of the momentum variables. Since the function \( \hat{\psi}_m|_{t=0} \) is compactly supported and smooth in the spatial variables, its Fourier transform \( \hat{\psi}_m^0(\vec{k}) \) has rapid decay. This shows that the function (16.4.1) is indeed square integrable. Using Plancherel, we see that condition (a) in Definition 15.2.2 is satisfied. Moreover, the operator \( T \) is simply the operator of multiplication by \( \sqrt{\vec{k}^2} \), so that condition (b) obviously holds. This again shows the weak mass oscillation property.

In order to prove the strong mass oscillation property, we need to compute the inner product \( \langle p \psi | p \phi \rangle \). To this end, we first write this inner product in momentum space as

\[
\langle p \psi | p \phi \rangle = \int \frac{d^4k}{(2\pi)^4} \chi_I(m) \frac{1}{4m^2} \langle (\vec{k} + m) \gamma^0 \hat{\psi}_m^0(\vec{k}) | (\vec{k} + m) \gamma^0 \hat{\psi}_m^0(\vec{k}) \rangle \bigg|_{m=\sqrt{k^2}}
\]

\[
= \int \frac{d^4k}{4\pi^2} \chi_I(m) \frac{1}{2m} \langle \gamma^0 \hat{\psi}_m^0(\vec{k}) | (\vec{k} + m) \gamma^0 \hat{\psi}_m^0(\vec{k}) \rangle \bigg|_{m=\sqrt{k^2}}.
\]
Reparametrizing the $k^0$-integral as an integral over $m$, we obtain
\[
\langle \mathbf{p} \psi | \phi \rangle = \frac{1}{4\pi^2} \int_I \int_{\mathbb{R}^3} \frac{d^3k}{2|k^0|} \left< \gamma^0 \psi_m^0(\vec{k}) \left( (\vec{k} + m) \gamma^0 \phi^0_m(\vec{k}) \right) \right|_{k^0 = \pm \sqrt{|\vec{k}|^2 + m^2}}. \quad (16.4.2)
\]
Estimating the inner product with the Schwarz inequality and applying Plancherel’s theorem, one finds
\[
|\langle \mathbf{p} \psi | \phi \rangle| \leq \frac{1}{4\pi^2} \int_I \int_{\mathbb{R}^3} \|\psi_m^0(\vec{k})\| \|\phi^0_m(\vec{k})\| d^3k \leq 2\pi \int_I \|\psi_m\| \|\phi_m\| dm.
\]
Thus the inequality (15.2.8) holds. □

• Explain how to read off the eigenvalues and eigenspaces of the fermionic signature operator.

**Exercises**

**Exercise 16.1.** This exercise is devoted to a clean proof of the distributional relation (16.4.5) in one dimension. More precisely, we want to prove the slightly more general statement that for any function $\eta \in C^1(\mathbb{R}) \cap L^1(\mathbb{R})$,
\[
\lim_{\varepsilon \searrow 0} \int_{\mathbb{R}} \eta(x) \left( \frac{1}{x - i\varepsilon} - \frac{1}{x + i\varepsilon} \right) dx = 2\pi i \eta(0). \quad (16.4.3)
\]
(a) Let $\eta \in C^1(\mathbb{R}) \cap L^1(\mathbb{R})$ with $\eta(0) = 0$. Show with the help of Lebesgue’s dominated convergence theorem that (16.4.3) holds.
(b) Show with residues that (16.4.3) holds for the function $\eta(x) = 1/(x^2 + 1)$.
(c) Combine the results of (a) and (b) to prove (16.4.3) for general $\eta \in C^1(\mathbb{R}) \cap L^1(\mathbb{R})$.

**Exercise 16.2.** This exercise recalls basics on the principal value in one dimension
\[
\frac{1}{2} \lim_{\varepsilon \searrow 0} \int_{\mathbb{R}} \eta(x) \left( \frac{1}{x - i\varepsilon} + \frac{1}{x + i\varepsilon} \right) dx =: \text{PP} \frac{\eta(x)}{x}. \quad (16.4.4)
\]
(a) Repeat the method in Exercise 16.1 to show that the limit of the left side of (16.4.4) exist for any $\eta \in C^1(\mathbb{R}) \cap L^1(\mathbb{R})$. Derive a corresponding estimate which shows that PP is a well-defined tempered distribution.
(b) Show that for any $\eta \in C^1(\mathbb{R}) \cap L^1(\mathbb{R})$,
\[
\text{PP}(\eta) = \lim_{\varepsilon \searrow 0} \left( \int_{-\infty}^{-\varepsilon} + \int_{\varepsilon}^{\infty} \right) \frac{\eta(x)}{x} dx.
\]

**Exercise 16.3.** The goal of this exercise is to justify that the one-dimensional relations
\[
\lim_{\varepsilon \searrow 0} \left( \frac{1}{x - i\varepsilon} - \frac{1}{x + i\varepsilon} \right) = 2\pi i \delta(x) \quad (16.4.5)
\]
\[
\frac{1}{2} \lim_{\varepsilon \searrow 0} \left( \frac{1}{x - i\varepsilon} + \frac{1}{x + i\varepsilon} \right) =: \frac{\text{PP}}{x}. \quad (16.4.6)
\]
apply four-dimensional setting to obtain the relation
\[
\lim_{\varepsilon \searrow 0} \frac{1}{\tau^2 + (\varepsilon + it)^2} = \lim_{\varepsilon \searrow 0} \frac{1}{\tau^2 - t^2 + i\varepsilon t} = -\frac{\text{PP}}{\xi^2} - i\pi \delta(\xi^2) \epsilon(\xi^0), \quad (16.4.7)
\]
(a) Let $T$ be a distribution on $\mathbb{R}$, $\Omega \subset \mathcal{M}$ be an open subset of Minkowski space and $f : \Omega \to \mathbb{R}$ a smooth function with nowhere vanishing gradient. Show that the relation

$$(f^* T)(\eta) := T(\phi_f(\eta)) , \quad \eta \in C_0^\infty(\Omega)$$

with

$$\phi_f(\eta)(t) := \frac{\partial}{\partial t} \int_{\Omega} \Theta(t - f(x)) \eta(x) \, d^4x$$

(where $\Theta$ is the Heaviside function) defines $f^* T$ as a distribution on $\Omega$ (this is the so-called pullback of $T$ under $f$; for details see [84, Section 7.2]).

(b) Choosing $\Omega$ as the half space in the future, $\Omega = \{ x \in \mathcal{M}, x^0 > 0 \}$, one can rewrite the expression on the left of (16.4.7) as

$$\lim_{\varepsilon \to 0} \frac{1}{r^2 - t^2 + i\varepsilon}.$$ 

Use (a) to conclude that this expression is a well-defined distribution for any $\varepsilon > 0$. Show that the limit $\varepsilon \to 0$ exist in the distributional sense.

(c) Repeating the procedure of (b) for the half space in the past, one obtains a distribution on $\mathcal{M} \setminus \{ t = 0 \}$. Show that this distribution coincides with the limit in (16.4.7).

Hint: Similar as in Exercise 16.1, one can estimate the behavior at the origin with Lebesgue’s dominated convergence theorem.

EXERCISE 16.4. This exercise is devoted to the advanced Green’s function $s^\vee_m$.

(a) Assume that $m > 0$. Show that the limit $\nu \to 0$ in (16.1.4) exist in the distributional sense.

(b) Show that the limit $\nu \to 0$ in (16.1.4) also exists in the massless case $m = 0$ and that

$$\lim_{m \to 0} s^\vee_m(k) = s^\vee_0(k)$$

as a distribution.

Hint: Proceed similar as in Exercise 16.3.

(c) Consider the Fourier integral in the $q^0$-variable

$$\int_{-\infty}^{\infty} \frac{1}{q^2 - m^2 - i\nu q^0} e^{iq^0 t} \, dq^0.$$

Show with residues that this integral vanishes for sufficiently small $\nu$ if $t < 0$.

(d) Argue with Lorentz invariance to prove the left side of (16.1.5).

EXERCISE 16.5. Modifying the location of the poles in (16.1.4) gives rise to the distribution

$$s^F_m(k) := \lim_{\nu \to 0} \frac{k + m}{k^2 - m^2 + i\nu}.$$ 

This is the well-known Feynman propagator, which is often described intuitively by saying that “positive frequencies move to the future and negative frequencies move to the past.” Make this sentence precise by a computation similar to that in Exercise 16.4 (c).

EXERCISE 16.6. Given $\omega \in \mathbb{R}$, we consider the ordinary differential operator $D = i\partial_t + \omega$.

(a) Construct the advanced and retarded Green’s functions, which satisfy in analogy to the kernels in the lecture the equation

$$D_t s(t, t') = \delta(t - t').$$
(b) Using the formula of the lecture, compute the resulting causal fundamental solution. How is it related to the time evolution operator $U^{t,t'}$? On which Hilbert space does the time evolution operator act as a unitary operator?

**EXERCISE 16.7.** Consider the massless Dirac equation $D\psi = 0$ in the two-dimensional spacetime cylinder $\mathbb{R} \times S^1$, i.e.

$$D = i \begin{pmatrix} 0 & 1 \\
1 & 0 
\end{pmatrix} \partial_t + i \begin{pmatrix} 0 & 1 \\
1 & 0 
\end{pmatrix} \partial_\varphi$$

with $t \in \mathbb{R}$ and $\varphi \in (0, 2\pi)$.

(a) Choose the spin inner product such that the Dirac matrices become symmetric. What is the resulting spacetime inner product $\langle ., . \rangle$? What is the scalar product $(., .)$?

(b) Employ for $k \in \mathbb{Z}$ the separation ansatz

$$\psi(t, \varphi) = e^{-ik\varphi} \chi(t) \quad \text{with} \quad \chi(t) \in \mathbb{C}^2.$$  

Derive the resulting ODE for $\chi$. Compute the time evolution operator for this ODE. *Hint:* Use the result of Exercise [16.6](#).

(c) Use a Fourier series decomposition in order to deduce a series representation of the time evolution operator of the Dirac operator on $\mathbb{R} \times S^1$. Try to carry out the infinite series to obtain a closed expression for $U^{t,t'}$. How can one see finite propagation speed?

**EXERCISE 16.8.** As in Exercise [16.7](#) we consider the two-dimensional massless Dirac equation.

(a) Adapt the formulas of the lecture for the advanced and retarded Green’s operators in momentum space to the two-dimensional massless case.

(b) Compute the Fourier transform to obtain $s^\vee(x,y)$ and $s^\wedge(x,y)$.

(c) Use the result of (b) to compute the causal fundamental solution and the time evolution operator.

(d) How can one see finite propagation speed? How is the obtained formula related to the formula in Exercise [16.7](#)(c)?

**EXERCISE 16.9.** Verify the relations [16.3.11](#) by direct computation starting from the definition [16.3.10](#).

**EXERCISE 16.10.** Verify by formal computation that in the Minkowski vacuum, the fundamental solution $k_m$ and the Green’s operator $s_m$ defined by

$$s_m := \frac{1}{2} \left( s_m^\vee + s_m^\wedge \right)$$

satisfy the distributional relations in the mass parameters $m$ and $m'$

$$k_m k_{m'} = \delta(m - m') p_m, \quad k_m s_{m'} k_m = \frac{\text{PP}}{m - m'} k_m,$$

where PP denotes the principal part, and $p_m$ is the distribution

$$p_m(k) = (\hat{k} + m) \delta(k^2 - m^2).$$

*Hint:* By a “formal computation” we mean that you do not need to evaluate weakly in the mass with test functions.
Exercise 16.11. Proceed similar as in Exercise 16.10 to derive a relation for the operator product $s_m s_{m'}$. Derive the relation

$$s_m s_{m'} = \frac{PP}{m - m'} (s_m - s_{m'}) + \pi^2 \delta(m - m') p_m.$$
CHAPTER 17

Methods of Scattering Theory

We return to the Cauchy problem in the presence of an external potential,

\[ (\mathcal{D} - m) \psi_m = 0, \quad \psi_m|_{t_0} = \psi_0 \in C^\infty(\mathcal{N}_{t_0} \simeq \mathbb{R}^3, S\mathcal{M}), \]  

(17.0.1)

with \( \mathcal{D} \) as in (1.3.14). For notational clarity, we shall often denote the objects in the presence of the external potential by a tilde (the “interacting objects”), whereas the objects without tilde refer to the Minkowski vacuum.

17.1. The Lippmann-Schwinger Equation

The Dirac dynamics can be rewritten in terms of a symmetric operator \( \tilde{H} \). To this end, we multiply the Dirac equation (1.3.14) by \( \gamma^0 \) and bring the \( t \)-derivative separately on one side of the equation,

\[ i\partial_t \psi_m = \tilde{H} \psi_m, \quad \text{where} \quad \tilde{H} := -\gamma^0 (i\gamma^\nabla + \mathcal{B} - m) \]  

(17.1.1)

(note that \( \gamma^j \partial_j = \gamma^0 \partial_t + \gamma^\nabla \)). We refer to (17.1.1) as the Dirac equation in Hamiltonian form. The fact that the scalar product (15.1.2) is time independent implies that for any two solutions \( \phi_m, \psi_m \in C^\infty_{sc}(\mathcal{M}, S\mathcal{M}) \cap \mathcal{H}_m \),

\[ 0 = \partial_t (\phi_m | \psi_m) = i((\tilde{H} \phi_m | \psi_m)_m - (\phi_m | \tilde{H} \psi_m)_m), \]

showing that the Hamiltonian is a symmetric operator on \( \mathcal{H}_m \). The Lippmann-Schwinger equation can be used to compare the dynamics in the Minkowski vacuum with the dynamics in the presence of an external potential. We denote the time evolution operator in the Minkowski vacuum by \( U_{t,t_0}^m \).

**Proposition 17.1.1.** The Cauchy problem (17.0.1) has a solution \( \psi_m \) which satisfies the equation

\[ \psi_m|_t = U_{t,t_0}^m \psi_0 + i \int_{t_0}^t U_{t,m}^{t,\tau} (\gamma^0 \mathcal{B} \psi_m)|_{\tau} \, d\tau, \]  

(17.1.2)

referred to as the Lippmann-Schwinger equation.

**Proof.** Obviously, the wave function \( \psi_m|_t \) given by (17.1.2) has the correct initial values at \( t = t_0 \). Thus it remains to show that \( \psi_m|_t \) satisfies the Dirac equation. To this end, we rewrite the Dirac equation in the Hamiltonian form (17.1.1), and separate the vacuum Hamiltonian \( H \) from the term involving the external potential,

\[ (i\partial_t - H) \psi_m = -\gamma^0 \mathcal{B} \psi_m \quad \text{with} \quad H = -i\gamma^0 \gamma^\nabla + \gamma^0 m. \]  

(17.1.3)

Applying the operator \( i\partial_t - H \) to (17.1.2) and observing that the time evolution operator maps to solutions of the vacuum Dirac equation, only the derivative of the upper limit of integration contributes,

\[ (i\partial_t - H) \psi_m|_t = -U_{t,m}^{t,\tau} (\gamma^0 \mathcal{B} \psi_m)|_{\tau=t} = -\gamma^0 \mathcal{B} \psi_m|_t, \]

so that (17.1.3) is indeed satisfied. \( \square \)
17.2. The Mass Oscillation Property in the Presence of an External Potential

The goal of this section is to prove the following result:

**Theorem 17.2.1.** Assume that the external potential $B$ is smooth and for large times decays faster than quadratically in the sense that

$$|B(t)|_{C^2} \leq \frac{c}{1+|t|^{2+\varepsilon}}$$

for suitable constants $\varepsilon, c > 0$. Then the strong mass oscillation property holds.

The $C^2$-norm in (17.2.1) is defined as follows. We denote spatial derivatives by $\nabla$ and use the notation with multi-indices, i.e. for a multi-index $\alpha = (\alpha_1, \ldots, \alpha_p)$ we set $\nabla^\alpha = \partial_{\alpha_1} \cdots \partial_{\alpha_p}$ and denote the length of the multi-index by $|\alpha| = p$. Then the spatial $C^k$-norms of the potential are defined by

$$|B(t)|_{C^k} := \max_{|\alpha| \leq k} \sup_{x \in \mathbb{R}^3} |\nabla^\alpha B(t, x)|,$$

where $|.|$ is the sup-norm corresponding to the norm $|\phi|^2 := \langle \phi | \gamma^0 \phi \rangle$ on the spinors.

**17.2.1. Proof of the Weak Mass Oscillation Property.** In this section, we prove the following theorem.

**Theorem 17.2.2.** Assume that the time-dependent external potential $B$ is smooth and decays faster than quadratically for large times in the sense that (17.2.1) holds for suitable constants $c, \varepsilon > 0$. Then the Dirac operator $D = i\partial / \partial t + B$ has the weak mass oscillation property.

We expect that this theorem could be improved by weakening the decay assumptions on the potential. However, this would require refinements of our methods which would go beyond the scope of this paper. Also, using that Dirac solutions dissipate, the pointwise decay in time could probably be replaced or partially compensated by suitable spatial decay assumptions. Moreover, one could probably refine the result of the above theorem by working with other norms (like weighted $C^k$- or Sobolev norms).

The main step is the following basic estimate, which is the analog of Lemma 16.3.1 in the presence of an external potential.

**Proposition 17.2.3.** Under the decay assumptions (17.2.1) on the external potential $B$, there are constants $c, \varepsilon > 0$ such that for every family $\psi \in \mathcal{H}^\infty$ of solutions of the Dirac equation (1.3.14) with varying mass,

$$\left\| (p \psi)|_t \right\|_t \leq \frac{c}{1+|t|^{1+\varepsilon}} \sup_{m \in I} \sum_{b=0}^{2} \left\| (\partial^b_m \phi_m)|_t \right\|_{W^{2,2}}.$$  

(17.2.3)

We first show that this proposition implies the weak mass oscillation property.

**Proof of Theorem 17.2.2 under the assumption that Proposition 17.2.3 holds.** In order to derive the inequality (15.2.6), we begin with the estimate

$$|\langle p \psi | p \phi \rangle| \leq \frac{1}{2\pi} \int_{-\infty}^{\infty} \left| (p \psi)|_t \right| \left| p \phi|_t \right| dt \leq \sup_{t \in \mathbb{R}} \left\| p \phi|_t \right\|_t \int_{-\infty}^{\infty} \left\| p \psi|_t \right\|_t dt.$$
The last integral is finite by Proposition \[17.2.3\]. The supremum can be bounded by the Hilbert space norm using the Hölder inequality,

\[
\|p\phi(t)\|_l = \left\| \int_I \phi_m |t \ dm \right\|_l \leq \int_I \|\phi_m(t)\|_l \ dm \leq \sqrt{|I|} \left( \int_I \|\phi_m(t)\|_{l}^2 \ dm \right)^{1/2} = \sqrt{|I|} \|\phi\|,
\]

giving \[(15.2.6)\).

Using \[(13.3.13)\], the Dirac operator $i\partial + B$ is formally self-adjoint with respect to the inner product $<.,.>$. Therefore, the identity \[(15.2.7)\] can be obtained just as in \[(16.3.7)\] by integrating the Dirac operator in spacetime by parts, noting that we do not get boundary terms in view of the time decay in Proposition \[17.2.3\].

The remainder of this section is devoted to the proof of Proposition \[17.2.3\]. We make use of the Lippmann-Schwinger equation \[(17.1.2)\],

\[
\psi_m |t = U_{m}^{t,0} \psi_m |_{t=0} + i \int_0^t U_{m}^{t,\tau} (\gamma^0 B \psi_m)|_{\tau} d\tau.
\]

Since the first summand of this equation is controlled by Lemma \[16.3.1\] it remains to estimate the second summand. Again using \[(16.3.17)\] and integrating by parts with respect to the mass, we obtain

\[
\int_I U_{m}^{t,\tau} (\gamma^0 B \psi_m) |_{\tau} \ dm = \frac{1}{(t-\tau)^2} \int_I (A_{m}^{t,\tau} \partial^2_m - B_{m}^{t,\tau} \partial_m + C_{m}^{t,\tau}) (\gamma^0 B \psi_m)|_{\tau} \ dm
\]

and thus

\[
\left\| \int_I U_{m}^{t,\tau} (\gamma^0 B \psi_m) |_{\tau} \ dm \right\|_l \leq \frac{c|I|}{(t-\tau)^2} \sup_{m \in I} \sum_{b=0}^{2} \|B(\tau) (\partial^b_m \psi_m)|_{\tau}\|_{W^{2,2}}
\]

\[
\leq \frac{c|I|}{(t-\tau)^2} |B(\tau)\|_{C^2} \sup_{m \in I} \sum_{b=0}^{2} \|\partial^b_m \psi_m|_{\tau}\|_{W^{2,2}}.
\]

We now bound $B(\tau)$ with the help of \[(17.2.1)\] and estimate the Sobolev norm $\|\partial^b_m \psi_m|_{\tau}\|_{W^{2,2}}$ at time $\tau$ by means of Lemma \[13.5.1\]. This gives rise to the inequality

\[
\left\| \int_I U_{m}^{t,\tau} (\gamma^0 B \psi_m) |_{\tau} \ dm \right\|_l \leq \frac{c^2 C |I|}{(t-\tau)^2} \frac{1 + |\tau|^2}{1 + |\tau|^{2+\varepsilon}} \sup_{m \in I} \sum_{b=0}^{2} \|\partial^b_m \psi_m|_{\tau=0}\|_{W^{2,2}},
\]

which yields the desired decay provided that $\tau$ and $t$ are not close to each other. More precisely, we shall apply this inequality in the case $|\tau| \leq |t|/2$. Then the estimate simplifies to

\[
\left\| \int_I U_{m}^{t,\tau} (\gamma^0 B \psi_m) |_{\tau} \ dm \right\|_t \leq \frac{\tilde{C}}{t^2 (1 + |\tau|^\varepsilon) \sup_{m \in I} \sum_{b=0}^{2} \|\partial^b_m \psi_m|_{\tau=0}\|_{W^{2,2}}},
\]

if $|\tau| \leq |t|/2$ \[(17.2.5)\]

with a new constant $\tilde{C} > 0$. In the remaining case $|\tau| > |t|/2$, we use the unitarity of $U_{m}^{t,\tau}$ to obtain

\[
\left\| \int_I U_{m}^{t,\tau} (\gamma^0 B \psi_m) |_{\tau} \ dm \right\|_t \leq |I| |B(\tau)|_{C^0} \sup_{m \in I} \|\psi_m\|.
\]
Applying (17.2.1) together with the inequality |τ| > |t|/2, this gives
\[ \left\| \int_I U^{t, \tau}_m (\gamma^0 B \psi_m) | \tau \right\| dm \leq \frac{\tilde{C}}{t^{2+\varepsilon}} \sup_{m \in I} \| \psi_m \| \quad \text{if } |\tau| > |t|/2 . \] (17.2.6)
This again decays for large \( t \) because \( \tau \) is close to \( t \) and \( |B(\tau)| \sim 0 \) decays for large \( \tau \).

Comparing (17.2.5) and (17.2.6), we find that the inequality in (17.2.5) even holds for all \( \tau \). Thus integrating this inequality over \( \tau \in [0, t] \), we obtain the following estimate for the second summand in (17.2.4),
\[ \left\| \int_I \int_0^t U^{t, \tau}_m (\gamma^0 B \psi_m) | \tau \right\| d\tau \right\| \leq \frac{C'}{t^{1+\varepsilon}} \sup_{m \in I} \sum_{b=0}^2 \left\| \partial^b_m \psi | t=0 \right\| W^{2,2} . \]

(17.2.7)
(where \( C' > 0 \) is a new constant). Combining this inequality with the estimate (16.3.3) of the first summand in (17.2.4), we obtain the desired inequality (17.2.3). This concludes the proof of Proposition 17.2.3.

17.2.2. Proof of the Strong Mass Oscillation Property. In this section, we prove the following result.

**Theorem 17.2.4.** Assume that the weak mass oscillation property holds and that the external potential \( B \) satisfies the condition
\[ \int_{-\infty}^{\infty} |B(\tau)| \sim 0 \quad \text{d} \tau < \infty . \] (17.2.7)
Then the Dirac operator \( \mathcal{D} = i \partial + B \) has the strong mass oscillation property.

Combining this theorem with Theorem 17.2.2, one immediately obtains Theorem 17.2.1.

For the proof we shall derive an explicit formula for the fermionic signature operator (Proposition 17.2.5). This formula is obtained by comparing the dynamics in the presence of the external potential with that in the Minkowski vacuum using the Lippmann-Schwinger equation, and by employing distributional relations for products of fundamental solutions and Green’s functions (Lemma 17.2.8).

We first return to the formula (16.4.2) in the Minkowski vacuum. Applying Plancherel’s theorem and using (15.1.2), we conclude that
\[ \langle \mathcal{S}_m(\vec{k}) \phi | \phi \rangle = \int_I (\psi_m^0 | \mathcal{S}_m(\vec{k}) \phi^0_m) dm , \] (17.2.8)
where
\[ \mathcal{S}_m(\vec{k}) := \sum_{k^0 = \pm \omega(\vec{k})} \frac{\vec{k} + m}{2 \omega(\vec{k})} \gamma^0 = \frac{\vec{k} \gamma + m}{\omega(\vec{k})} \gamma^0 . \] (17.2.9)
Comparing (17.2.8) with (15.3.4), one sees that the matrix \( \mathcal{S}_m(\vec{k}) \) is indeed the fermionic signature operator, considered as a multiplication operator in momentum space. By direct computation, one verifies that the matrix \( \mathcal{S}_m(\vec{k}) \) has eigenvalues ±1.

In order to compare the dynamics in the presence of the external potential with that in the Minkowski vacuum, we work with the Hamiltonian formulation. We decompose the Dirac Hamiltonian (17.1.1) into the Hamiltonian in the Minkowski vacuum (17.1.3) plus a potential,
\[ \tilde{H} = H + \mathcal{V} \quad \text{with} \quad \mathcal{V} := -\gamma^0 B . \]
Proposition 17.2.5. Assume that the potential $\mathcal{B}$ satisfies the condition (17.2.7). Then for every $\psi, \phi \in \mathcal{H}^\infty$,

$$\langle p \psi | p \phi \rangle = \int_t (\psi_m | \tilde{S}_m \phi_m)_m \, dm,$$

where $\tilde{S}_m : \mathcal{H}_m \to \mathcal{H}_m$ are bounded linear operators which act on the wave functions at time $t_0$ by

$$\tilde{S}_m = S_m - \frac{i}{2} \int_{-\infty}^{\infty} \epsilon(t - t_0) \left[ S_m U_{m,t}^0 \, \nabla(t) \tilde{U}_{m,t}^0 - \tilde{U}_{m,t}^0 \, \nabla(t) S_m U_{m,t}^0 \right] \, dt \quad (17.2.11)$$

$$+ \frac{1}{2} \left( \int_{t_0}^{\infty} \int_{t_0}^{\infty} + \int_{-\infty}^{t_0} \int_{-\infty}^{t_0} \right) \tilde{U}_{m,t}^0 \, \nabla(t) S_m U_{m,t'} \, \nabla(t') \tilde{U}_{m,t}^0 \, dt \, dt' \quad (17.2.12)$$

(and $S_m$ is again the fermionic signature operator of the vacuum (17.2.3)).

Before entering the proof of this proposition, it is instructive to verify that the above formula for $\tilde{S}_m$ does not depend on the choice of $t_0$.

Remark 17.2.6. (Independence of $\tilde{S}_m$ on $t_0$) Our strategy is to differentiate the above formula for $\tilde{S}_m$ with respect to $t_0$ and to verify that we obtain zero. We first observe that taking a solution $\phi_m \in \mathcal{H}_m$ of the Dirac equation in the presence of $\mathcal{B}$, evaluating at time $t_0$ and applying the time evolution operator $\tilde{U}_{m,t}^0$ gives $\phi_m$ at time $t$, i.e. $\tilde{U}_{m,t}^0 \phi_m|_{t_0} = \phi_m|_t$. Differentiating with respect to $t_0$ yields

$$\partial_{t_0} \tilde{U}_{m,t}^0 \phi_m|_{t_0} = 0.$$

The situation is different when one considers the time evolution operator of the vacuum. Namely, in the expression $U_{m,t}^0 \phi_m|_{t_0}$, the wave function $\phi_m$ satisfies the Dirac equation $(i\partial_t - H)\phi_m = \nabla\phi_m$, whereas the time evolution operator solves the Dirac equation with $\nabla \equiv 0$. As a consequence,

$$\partial_{t_0} U_{m,t}^0 \phi_m|_{t_0} = -i U_{m,t}^0 (\nabla\phi_m)|_{t_0}.$$ 

Using these formulas together with $U_{t_0, t_0} = 1 = \tilde{U}_{t_0, t_0}$, a straightforward computation gives

$$\partial_{t_0} (\psi_m | (17.2.11) \phi_m)|_{t_0} = -i (\psi_m | [S_m, \nabla] \phi_m)|_{t_0}$$

$$- \frac{i}{2} (-2) (\psi_m | (S_m \nabla(t_0) - \nabla(t_0) S_m) \phi_m)|_{t_0}$$

$$- \frac{i}{2} \int_{-\infty}^{\infty} \epsilon(t - t_0) \left( - (\nabla(t_0)) \psi_m | S_m U_{m,t}^{0,t} \nabla(t) \tilde{U}_{m,t}^{0,t} \phi_m \right)|_{t_0} \, dt$$

$$+ \frac{i}{2} \int_{-\infty}^{\infty} \epsilon(t - t_0) \left( \psi_m | \tilde{U}_{m,t}^{0,t} \nabla(t) S_m U_{m,t}^{0,t} (-i \nabla(t_0)) \phi_m \right)|_{t_0} \, dt$$

$$\partial_{t_0} (\psi_m | (17.2.12) \phi_m)|_{t_0} = - \frac{1}{2} \int_{-\infty}^{\infty} \epsilon(t' - t_0) \left( \psi_m | \nabla(t) S_m U_{m,t}^{t',t} \nabla(t') \tilde{U}_{m,t}^{t',t} \phi_m \right)|_{t_0} \, dt'$$

$$- \frac{1}{2} \int_{-\infty}^{\infty} \epsilon(t - t_0) \left( \psi_m | \tilde{U}_{m,t}^{t,t} \nabla(t) S_m U_{m,t}^{t,t} \nabla(t_0) \phi_m \right)|_{t_0} \, dt,$$

where for notational simplicity we here omitted the restrictions $|t_0$ for the solutions $\psi_m$ and $\phi_m$. Adding the terms gives zero. \(\Diamond\)
The remainder of this section is devoted to the proof of Proposition 17.2.5. Our strategy is to combine the Lippmann-Schwinger equation with estimates in momentum space. We begin with two technical lemmas.

Lemma 17.2.7. Assume that the external potential \( B \) satisfies condition (17.2.7). For any \( t_0 \in \mathbb{R} \), we denote the characteristic functions in the future respectively past of this hypersurface \( t = t_0 \) by \( \chi^+_{t_0}(x) \) (i.e. \( \chi^+_{t_0}(x) = \Theta(x^0 - t_0) \)), where \( \Theta \) is the Heaviside function. Then for any \( \psi_m \in \mathcal{C}_c^\infty(\mathcal{M}, \mathcal{S}^\mathcal{M}) \cap \mathcal{H}_m \), the wave function \( k_m(\chi^+_{t_0} B \psi_m) \) is a well-defined vector in \( \mathcal{H}_{t_0} \) and

\[
\|k_m(\chi^+_{t_0} B \psi_m)\|_{t_0} \leq \frac{1}{2\pi} \|\psi_m\|_m \int_{-\infty}^{\infty} \chi^+_{t_0}(\tau) |B(\tau)|_{C^0} d\tau .
\]

Proof. Using the integral kernel representation (16.2.3) and (16.2.4) together with the fact that the time evolution in the vacuum is unitary, we obtain

\[
2\pi \left\| \int_{\mathbb{R}^3} k_m((t_0, \cdot), (\tau, \vec{y})) (\chi^+_{t_0} B \psi_m)(\tau, \vec{y}) d^3y \right\|_{t_0} = \left\| \int_{t_0}^{t_{\tau}} U_{t_{\tau}}^{t} \gamma^0 (\chi^+_{t_0} B \psi_m) \right\|_{t_0} = \left\| \gamma^0 (\chi^+_{t_0} B \psi_m) \right\|_{t_0} \leq |B(\tau)|_{C^0} \|\psi_m\|_m .
\]

Integrating over \( \tau \) and using (17.2.7) gives the result. \( \square \)

The following lemma is proved in [54, Eqs. (2.13)–(2.17)] (see Exercises 16.10 and 16.11).

**Lemma 17.2.8.** In the Minkowski vacuum, the fundamental solution \( k_m \) and the Green’s function \( s_m \) defined by

\[
s_m := \frac{1}{2} (s^\vee_m + s^\wedge_m) \tag{17.2.13}
\]

satisfy the distributional relations in the mass parameters \( m \) and \( m' \)

\[
k_m k_{m'} = \delta(m - m') p_m \quad \quad k_m s_{m'} = s_{m'} k_m = \frac{PP}{m - m'} k_m \quad \quad s_m s_{m'} = \frac{PP}{m - m'} (s_m - s_{m'}) + \pi^2 \delta(m - m') p_m ,
\]

where \( PP \) denotes the principal part, and \( p_m \) is the distribution

\[
p_m(k) = (\mathring{k} + m) \delta(k^2 - m^2) . \tag{17.2.14}
\]

Proof of Proposition 17.2.5. Let \( \psi \in \mathcal{H}^\infty \) be a family of solutions of the Dirac equation for varying mass. We denote the boundary values at time \( t_0 \) by \( \psi_m \). Then we can write the Lippmann-Schwinger equation (17.1.2) as

\[
\psi_m|_t = U_{t_0}^{t, \tau} \psi_m^0 + i \int_{t_0}^{t} U_{t_0}^{t, \tau} (\gamma^0 B \psi_m) d\tau .
\]

We now bring this equation into a more useful form. Expressing the time evolution operator with the help of (16.2.4) in terms of the fundamental solution, we obtain

\[
\psi_m(x) = 2\pi \int_{\mathbb{R}^3} k_m(x, (t_0, \vec{y})) \gamma^0 \psi_m^0(t_0, \vec{y}) d^3y + 2\pi i \int_{t_0}^{t} dy^0 \int_{\mathbb{R}^3} d^3y k_m(x, y)(B \psi_m)(y) .
\]
Applying (18.2.10) and using that the advanced and retarded Green’s functions are supported in the future and past light cones, respectively, we can rewrite the last integral in terms of the advanced and retarded Green’s functions,
\[
\psi_m = 2\pi k_m (\gamma^0 \delta(t_0)\psi^0_m - s^\wedge_m (\chi^+_{t_0}B\psi_m) - s^\vee_m (\chi^-_{t_0}B\psi_m)),
\]
where \(\delta(t_0(x) := \delta(t_0 - x^0)\) is the Dirac distribution supported on the hypersurface \(x^0 = t_0\).

Next, we express the advanced and retarded Green’s functions in terms of the Green’s function \(\chi\) : According to (18.2.10), we have the relations
\[
\psi_m = 2\pi \gamma^0 \delta(t_0)\psi^0_m + i\pi \epsilon_{t_0} B\psi_m,
\]
and thus
\[
\psi_m = k_m g_m - s_m B\psi_m \quad \text{with} \quad g_m := 2\pi \gamma^0 \delta(t_0)\psi^0_m + i\pi \epsilon_{t_0} B\psi_m,
\]
where \(\epsilon_{t_0}\) is the step function
\[
\epsilon_{t_0}(x) := \epsilon(x^0 - t_0)
\]
(and we omitted the brackets in expressions like \(k_m g_m \equiv k_m (g_m)\)). Note that the expression \(k_m g_m\) is well-defined according to Lemma 17.2.7. We also remark that by applying the operator \((i\partial_m - m)\) to the distribution \(g_m\) in (17.2.15), one immediately verifies that \(\psi_m\) indeed satisfies the Dirac equation \((i\partial_m - m)\psi_m = -B\psi_m\).

Now we can compute the inner product \(\langle p\psi | p\psi \rangle\) with the help of Lemma 17.2.8 Namely, using (17.2.15),
\[
\langle p\psi | p\psi \rangle = \int_{I \times I} \left< k_m g_m - s_m B\psi_m \mid k_{m'} g_{m'} - s_{m'} B\psi_{m'} \right> dm \, dm'
\]
\[
= \int_I \left< g_m | p_m g_m \right> + \pi^2 \left< B\psi_m | p_m B\psi_m \right> \, dm
\]
\[
+ \int_{I \times I} \frac{PP}{m - m'} \left< B\psi_m | k_{m'} g_{m'} \right> - \left< k_m g_m | B\psi_{m'} \right>
\]
\[
+ \left< B\psi_m \mid (s_m - s_{m'}) B\psi_{m'} \right> \, dm \, dm'.
\]
Note that this computation is mathematically well-defined in the distributional sense because \(\psi_m\) and \(g_m\) are smooth and compactly supported in the mass parameter \(m\).

Employing the explicit formula for \(g_m\) in (17.2.15), we obtain
\[
\langle p\psi | p\psi \rangle = \int_I \left< g_m | p_m g_m \right> + \pi^2 \left< B\psi_m | p_m B\psi_m \right> \, dm.
\]
Comparing (16.3.8) with (18.2.9) and taking into account that the operator \(S_m\) defined by (17.2.9) gives a minus sign for the states of negative frequency, we get
\[
p_m = S_m k_m.
\]
Using this identity together with Proposition 13.4.4 in the vacuum yields the relations
\[
\left< g_m | p_m g_m \right> = (k_m g_m \mid S_m k_m g_m)|_{t_0}
\]
\[
\left< B\psi_m | p_m B\psi_m \right> = (k_m B\psi_m \mid S_m k_m B\psi_m)|_{t_0}.
\]
We finally apply Proposition 13.6.1 to obtain the representation
\[
\langle p\psi | p\psi \rangle = \int_I \left( (h_m \mid S_m h_m)|_{t_0} + \pi^2 (k_m B\psi_m \mid S_m k_m B\psi_m)|_{t_0} \right) \, dm,
\]
where
\[
h_m := \psi_m + i\pi k_m (\epsilon_{t_0} B\psi_m).\]
Comparing (17.2.10) with (17.2.16), we get
\[
\langle \psi_m | \tilde{S}_m \psi_m \rangle = (h_m | S_m h_m)_{t_0} + \pi^2 (k_m \mathcal{B} \psi_m | S_m k_m \mathcal{B} \psi_m)_{t_0}.
\]
Expressing the operators \(k_m\) according to (16.2.4) by the time evolution operator and writing \(\psi_m\) in terms of the initial data as
\[
\psi_m|_t = \tilde{U}^{t,t_0} \psi|_{t_0},
\]
we obtain
\[
\langle \psi_m | \tilde{S}_m \psi_m \rangle = \langle \psi | S_m \psi \rangle_{t_0} - \frac{i}{2} \int_{-\infty}^{\infty} \epsilon(t-t_0) \langle \psi | S_m U^{t_0,t} \mathcal{V}(t) \tilde{U}^{t,t_0} \psi \rangle_{t_0} dt
\]
\[
+ \frac{i}{2} \int_{-\infty}^{\infty} \epsilon(t-t_0) \langle U^{t_0,t} \mathcal{V}(t) \tilde{U}^{t,t_0} \psi | S_m \psi \rangle_{t_0} dt dt
\]
\[
+ \frac{1}{4} \int_{\mathbb{R} \times \mathbb{R}} \epsilon(t-t_0) \epsilon(t'-t_0) \langle U^{t_0,t} \mathcal{V}(t) \tilde{U}^{t,t_0} \psi | S_m U^{t_0,t'} \mathcal{V}(t') \tilde{U}^{t',t_0} \psi \rangle_{t_0} dt dt' dt dt'
\]
Rearranging the terms and polarizing gives the result.

**Proof of Theorem 17.2.4.** Since the time evolution operators are unitary and the operators \(\tilde{S}_m\) have norm one (see (17.2.31)), the representation (17.2.11) and (17.2.12) gives rise to the following estimate for the sup-norm of \(\tilde{S}_m\),
\[
\| \tilde{S}_m \| \leq 1 + \int_{\mathbb{R}} |\mathcal{V}(t)_{C^0}| dt + \int_{\mathbb{R} \times \mathbb{R}} |\mathcal{V}(t)|_{C^0} |\mathcal{V}(t')|_{C^0} dt dt'.
\]
The decay assumption (17.2.7) implies that the sup-norm of \(\tilde{S}_m\) is bounded uniformly in \(m\). Using this fact in (17.2.10) gives the inequality (15.2.8), thereby establishing the strong mass oscillation property.

We finally remark that the uniqueness statement in Proposition 15.3.3 implies that (17.2.11) and (17.2.12) yields an explicit representation of the fermionic signature operator in the presence of a time-dependent external potential.

**Exercises**

**Exercise 17.1.** For a smooth one-parameter family of matrices \(F(\alpha), \alpha \in \mathbb{R}\), the **ordered exponential** \(\text{Pexp}(\int_a^b F(\alpha) \, d\alpha)\) is given by the Dyson series
\[
\text{Pexp} \left( \int_a^b F(\alpha) \, d\alpha \right) = 1 + \int_a^b F(t_0) \, dt_0 + \int_a^b \int_t^b F(t_0) \, dt_0 \, F(t_1) \, dt_1 + \int_a^b \int_t^b \int_t_0^b F(t_0) \, dt_0 \, F(t_1) \, dt_1 \, F(t_2) \, dt_2 + \cdots.
\]
In this exercise we collect a few elementary properties of the ordered exponential.

(a) Assume that the matrix-valued function \(F\) is commutative in the sense that
\[
[F(\alpha), F(\beta)] = 0 \quad \text{for all } \alpha, \beta \in [a, b].
\]
Show that the ordered exponential reduces to the ordinary exponential,
\[
\text{Pexp} \left( \int_a^b F(\alpha) \, d\alpha \right) = \exp \left( \int_a^b F(\alpha) \, d\alpha \right).
\]
Hint: Show inductively that
\[
\int_a^b dt_0 F(t_0) \int_{t_0}^b dt_1 F(t_1) \cdots \int_{t_{n-1}}^b dt_n F(t_n) = \frac{1}{(n+1)!} \left( \int_a^b F(t) dt \right)^{n+1}.
\]
(b) Assume that \( F \) is continuous on \([a,b]\). Show that the Dyson series converges absolutely and that

\[
\left\| \text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right) \right\| \leq \exp \left( \int_a^b \| F(\alpha) \| d\alpha \right).
\]

Hint: Estimate the integrals and apply (a).

(c) Show by direct computation that the ordered exponential satisfies the equations

\[
\frac{d}{da} \text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right) = -F(a) \text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right)
\]

(17.2.17)

\[
\text{Pexp} \left( \int_a^a F(\alpha) d\alpha \right) = 1.
\]

(17.2.18)

Use the uniqueness theorem for solutions of ordinary differential equations to give an alternative definition in terms of the solution of an initial-value problem. Use this reformulation to show the group property

\[
\text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right) \text{Pexp} \left( \int_b^c F(\alpha) d\alpha \right) = \text{Pexp} \left( \int_a^c F(\alpha) d\alpha \right).
\]

(17.2.19)

(d) Show that

\[
\frac{d}{db} \text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right) = \text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right) F(b).
\]

(17.2.20)

Hint: Differentiate the identity (17.2.19) in the case \( c = a \) and use the group properties (17.2.18) and (17.2.19).

(e) Show that

\[
\text{Pexp} \left( \int_a^b F(\alpha) d\alpha \right)^* = \text{Pexp} \left( \int_b^a (-F(\alpha)^*) d\alpha \right).
\]

Deduce that if \( F(\alpha) \) is an anti-Hermitian matrix, then the ordered exponential is a unitary matrix. Hint: There are two alternative methods. One method is to argue using the differential equations (17.2.17) and (17.2.20) or with the group property. A more computational approach is to take the adjoint of the Dyson series and reparametrize the integrals.

**Exercise 17.2.** Given \( \omega \in \mathbb{R} \) and a smooth function \( V(t) \), we consider the ordinary differential equation

\[
(i\partial_t + \omega) \phi(t) = V(t) \phi(t).
\]

(a) Write down the Lippmann-Schwinger equation, taking the right side of the equation as the perturbation. Hint: The free time evolution operator \( U^{t,t'} \) was computed in Exercise 16.6.

(b) Express the Lippmann-Schwinger equation in the case \( \omega = 0 \) explicitly as an integral equation. How is it related to the integral equation used in the Picard iteration (in the proof of the Picard-Lindelöf theorem)?
CHAPTER 18

Methods of Perturbation Theory

In Chapter 15, the unregularized kernel of the fermionic projector was constructed abstractly with functional analytic methods. In order to fill these constructions with life, one can analyze the kernel with methods of perturbation theory. The resulting explicit formulas give a detailed understanding of the structure of this kernel. We now outline the perturbative methods; more details can be found in [48, Chapter 2] or in the original papers [35, 54, 81].

In preparation, we rewrite the definition of the fermionic signature operator in a way suitable for the perturbative treatment. Our starting point is the representation (15.3.4) of the spacetime inner product in terms of the scalar product,

\[ \langle p \psi | p \psi' \rangle = \hat{M} \left( \psi_m | S_m \psi'_m \right)_m \ dm. \]  

(18.0.1)

Here \( \psi = (\psi_m)_{m \in I} \) and similarly \( \psi' \) are families of solutions of the Dirac equation for a varying mass parameter. Now we consider specifically families obtained by acting with the causal fundamental solution on given test wave functions, i.e.

\[ \psi_m = \tilde{k}_m \phi \quad \text{and} \quad \psi'_m = \tilde{k}_m \phi' \quad \text{with} \quad \phi, \phi' \in C^\infty_{sc}(\mathcal{M}, S\mathcal{M}). \]

Using this ansatz in (18.0.1) and pulling the mass integrals outside, we obtain the formula

\[ \int_I \hat{I} dm \int_I \hat{I} dm' \langle \tilde{k}_m \phi | \tilde{k}_m \phi' \rangle = \int_I (\tilde{k}_m \phi | S_m \tilde{k}_m \phi')_m dm. \]  

(18.0.2)

Next, we rewrite the integrand on the left side by using that the fundamental solution is symmetric with respect to the spacetime inner product (see Corollary 13.4.5),

\[ \langle \tilde{k}_m \phi | \tilde{k}_m \phi' \rangle = \langle \phi | \tilde{k}_m \tilde{k}_m \phi' \rangle. \]

Moreover, the integrand on the right can be rewritten with the help of Proposition 13.4.4 as

\[ (\tilde{k}_m \phi | S_m \tilde{k}_m \phi')_m = \langle \phi | S_m \tilde{k}_m \phi' \rangle. \]

Thus (18.0.2) becomes

\[ \int_I dm \int_I dm' \langle \phi | \tilde{k}_m \tilde{k}_m \phi' \rangle = \int_I \langle \phi | S_m \tilde{k}_m \phi' \rangle dm. \]

Here one should keep in mind the product \( \tilde{k}_m \tilde{k}_m \) is an operator product in spacetime,

\[ (\tilde{k}_m \tilde{k}_m')(x, y) = \int_\mathcal{M} \tilde{k}_m(x, z) \tilde{k}_m'(z, y) \ d^4 z, \]  

(18.0.3)

whereas in the product \( S_m \tilde{k}_m \) we multiply by an operator on the Hilbert space \( \mathcal{H}_m \) defined for example at time \( t \). In order to clarify the notation, we write this product as

\[ S_m \mid_t \tilde{k}_m. \]
Then the relation (18.0.3) can be written in the short form
\[ \tilde{k}_m \tilde{k}_{m'} = \delta(m - m') S_m | \tilde{k}_m. \] (18.0.4)
In this way, one is led to considering products of operators in spacetime which involve the mass as a parameter. Carrying out the products gives rise to \( \delta \)-distributions in the respective mass parameters.

This computational procedure was introduced in [35]. In the Minkowski vacuum, it can be carried out most conveniently in momentum space. We begin with the formula for the causal fundamental solution in momentum space (16.3.8),
\[ k_m(p) = \left( \frac{1}{p^2} + m \right) \delta(p^2 - m^2) \epsilon(p^0). \] (18.0.5)
Then, using Plancherel together as well as the anti-commutation relations of the Dirac matrices, we obtain
\[ (k_m k_{m'}) (p) = (p + m) \delta(p^2 - m^2) \epsilon(p^0) \delta(p^2 - (m')^2) \epsilon(p^0) = (p^2 + (m + m') \phi + mm') \delta(m^2 - (m')^2) \delta(k^2 - m^2) = \delta(m - m') \phi \delta(p^2 - m^2) = \delta(m - m') \epsilon(p^0) k_m(p). \]
Comparing with (18.0.4), we can read off that that the fermionic signature operator simply is the operator of multiplication operator by the sign of the frequency,
\[ S_m(p) = \epsilon(p^0). \]
This computation is an efficient way of seeing that, in the Minkowski vacuum, the fermionic signature operator gives back the frequency splitting.

We proceed by explaining how this computation can be extended to the situation when an external potential \( B \) is present. We want to proceed order by order in a perturbation expansion in \( B \). Before entering the details, we point out that by a “perturbation expansion” we mean a formal expansion in powers of \( B \). The resulting formulas will be well-defined and finite to every order. But it is unknown whether the power series converges. This procedure is convincing because we already know from our functional analytic constructions in Chapter 15 that the fermionic signature operator and the unregularized fermionic projector are well-defined mathematical objects. With this in mind, the only purpose of the constructions in this chapter is to compute these objects more explicitly. For this purpose, a perturbative treatment order by order in perturbation theory is most suitable.

### 18.1. Perturbation Expansion of the Causal Green’s Operators

We already encountered the causal Green’s operators for the Dirac equation several times in this book. In Section 13.4, they were constructed with methods of hyperbolic partial differential equations (see Theorem 13.4.3). In Chapter 16.1 on the other hand, we used Fourier methods to derive explicit formulas for the causal Green’s operator in the Minkowski vacuum (see (16.1.4)). Taking these explicit formulas as the starting point, one can also write down closed formulas for the causal Green’s functions in the presence of an external potential. In order to state these formulas, we consider the Dirac equation (1.3.14) in the presence of an external potential \( B \). We always denote the objects in the presence of the external potential with a tilde, whereas the objects without tilde
18.2. THE CAUSAL PERTURBATION EXPANSION OF THE FERMIONIC PROJECTOR

refer to the vacuum. Then the advanced and retarded Dirac Green’s functions have the perturbation expansions

\[ s^\vee_m = \sum_{n=0}^{\infty} (-s_m B)^n s^\vee_m, \quad s^\wedge_m = \sum_{n=0}^{\infty} (-s_m B)^n s^\wedge_m. \]  

These formulas can be understood as follows. First, one sees by direct computation using the defining equation of the Green’s operator (16.1.1) that they are formal solutions of (16.2.1). For example, for the advanced Green’s operator,

\[ (i\partial / x + B - m) \left( \sum_{n=0}^{\infty} (-s_m B)^n s_m \right) \]

\[ = (i\partial / x - m) s^\vee_m \left( \sum_{n=0}^{\infty} (-B s^\vee_m)^n \right) + B \left( \sum_{n=0}^{\infty} (-s_m B)^n s^\wedge_m \right) \]

\[ = \sum_{n=0}^{\infty} (-B s^\vee_m)^n + B \left( \sum_{n=0}^{\infty} (-s_m B)^n s^\wedge_m \right) = 1. \]

Second, the fact that the Green’s operators in (27.3.4) are either all advanced or all retarded implements the causal properties of the respective Green’s operators. Let us consider for example the integral kernel of the first order contribution to the advanced Green’s operator

\[ (i\partial / x + B - m)(s_m(x, y)) = -\int_M s_m(x, z) B(z) s_m(z, y) d^4z. \]  

The integrand vanishes unless \( z \) lies in the causal future of \( y \) and \( x \) lies in the causal future of \( z \). Using transitivity of the causal relations, one concludes that the integral is zero unless \( x \) lies in the causal future of \( y \). In this sense, the expression (18.1.2) is again causal and retarded. The higher orders can be treated similarly by induction.

We finally explain in which sense the perturbation series (27.3.4) are mathematically well-defined. To every order in perturbation theory, the operator products are well-defined and finite, provided that the potential \( B \) is smooth and decays so fast at infinity that the functions \( B(x), x^i B(x), \) and \( x^i x^j B(x) \) are integrable (for an inductive proof see [48, Lemma 2.1.2]). Knowing that the Green’s operators are well-defined non-perturbatively (see Chapter 13), we disregard the issue of convergence of the perturbation series.

18.2. The Causal Perturbation Expansion of the Fermionic Projector

Using (13.4.8), we also have a unique perturbation expansion for the causal fundamental solution,

\[ \tilde{k}_m = \frac{1}{2\pi i} (s^\vee_m - s^\wedge_m). \]  

Using the identities

\[ s^\vee_m = s_m + i\pi k_m, \quad s^\wedge_m = s_m - i\pi k_m, \]  

where we introduced the symmetric Green’s operator

\[ s_m := \frac{1}{2} (s^\vee_m + s^\wedge_m), \]  

\[ s^\vee_m = \sum_{n=0}^{\infty} (-s_m B)^n s^\vee_m, \quad s^\wedge_m = \sum_{n=0}^{\infty} (-s_m B)^n s^\wedge_m, \]  

\[ (i\partial / x + B - m)(s_m(x, y)) = -\int_M s_m(x, z) B(z) s_m(z, y) d^4z. \]  

The integrand vanishes unless \( z \) lies in the causal future of \( y \) and \( x \) lies in the causal future of \( z \). Using transitivity of the causal relations, one concludes that the integral is zero unless \( x \) lies in the causal future of \( y \). In this sense, the expression (18.1.2) is again causal and retarded. The higher orders can be treated similarly by induction.

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one can write the above perturbation series as operator product expansions. More precisely,
\[ \tilde{k}_m = \sum_{\beta=0}^{\infty} (i\pi)^{2\beta} b^<_m k_m (b_m k_m)^{2\beta} b^>_m, \]  
(18.2.4)
where the factors \( b^*_m \) are defined by
\[ b^<_m = \sum_{n=0}^{\infty} (-s_m B)^n, \quad b_m = \sum_{n=0}^{\infty} (-B s_m)^n B, \quad b^>_m = \sum_{n=0}^{\infty} (-B s_m)^n. \]  
(18.2.5)

In the following constructions, we need to multiply the operator products in (18.2.4). These products have a mathematical meaning as distributions in the involved mass parameters,
\[ p_m p_m' = k_m k_{m'} = \delta(m - m') p_m \]  
(18.2.6)
\[ p_m k_{m'} = k_m p_{m'} = \delta(m - m') k_m \]  
(18.2.7)
\[ k_m b^>_m b^<_m = \delta(m - m') \left( p_m + \pi^2 k_m b_m p_m b_m k_m \right), \]  
(18.2.8)
where
\[ p_m(q) = (q + m) \delta(q^2 - m^2) \]  
(18.2.9)
\[ k_m(q) = (q + m) \delta(q^2 - m^2) \epsilon(q^0). \]  
(18.2.10)

Since all these formulas involve a common prefactor \( \delta(m - m') \), we can introduce a convenient notation by leaving out this factor and omitting the mass indices. For clarity, we denote this short notation with a dot, i.e. symbolically
\[ A \cdot B = C \quad \text{stands for} \quad A_m B_{m'} = \delta(m - m') C_m. \]  
(18.2.11)

With this short notation, the multiplication rules can be written in the compact form
\[ p \cdot p = k \cdot k = p, \quad p \cdot k = k \cdot p = k, \quad k b^> \cdot b^<_k = p + \pi^2 k b p k. \]  
(18.2.12)

In all the subsequent calculations, the operator products are well-defined provided that the potential \( B \) is sufficiently smooth and has suitable decay properties at infinity (for details see again [48, Lemma 2.1.2]). But again, all infinite series are to be understood merely as formal power series in the potential \( B \).

Using this notation, we can write (18.2.4) as
\[ \tilde{k} = k + \Delta k \quad \text{with} \quad \Delta \tilde{k} = \sum_{\beta=0}^{\infty} (i\pi)^{2\beta} b^<_k k (b k)^{2\beta} b^>_k - k, \]  
(18.2.13)
(note that \( \Delta \tilde{k} \) is at least linear in \( B \)). Powers of the operator \( \tilde{k} \) with the product (18.2.11) are well-defined using the multiplication rules (18.2.12). This makes it possible to develop a spectral calculus for \( \tilde{k} \), which is formulated most conveniently with contour integrals.

To this end, we introduce the resolvent by
\[ \tilde{R}_\lambda = (\tilde{k} - \lambda)^{-1}. \]  
(18.2.14)
We choose a contour \( \Gamma_+ \) which encloses the point 1 in counter-clockwise direction and does not enclose the points \(-1\) and 0. Likewise, \( \Gamma_- \) is chosen as a contour which encloses
the point $-1$ in counter-clockwise direction and does not enclose the points 1 and 0. Given a holomorphic function $f$ we define $f(\tilde{k})$ by

$$f(\tilde{k}) := \frac{1}{2\pi i} \oint_{\Gamma_+ \cup \Gamma_-} f(\lambda) \tilde{R}_\lambda \, d\lambda. \quad (18.2.15)$$

Before going on, we need to explain how the resolvent and these contour integrals are to be understood mathematically. First, the resolvent can be expressed in terms of the vacuum resolvent with a perturbation series being a formal Neumann series,

$$\tilde{R}_\lambda = (k - \lambda + \Delta k)^{-1} = (1 + R_\lambda \cdot \Delta k)^{-1} \cdot R_\lambda = \sum_{n=0}^{\infty} (-R_\lambda \cdot \Delta k)^n \cdot R_\lambda. \quad (18.2.16)$$

In order to define $R_\lambda$, we note that, according to (18.2.12), the operator $k$ has the eigenvalues $\pm 1$ and 0 with corresponding spectral projectors $(p \pm k)/2$ and $1 - p$. Hence we can write the free resolvent as

$$R_\lambda = \frac{p + k}{2} \left( \frac{1}{1 - \lambda} \right) + \frac{p - k}{2} \left( \frac{1}{1 - \lambda} \right) - \frac{1 - p}{\lambda}. \quad (18.2.17)$$

Substituting this formula in (18.2.16), to every order in perturbation theory we obtain a meromorphic function in $\lambda$ having poles only at $\lambda = 0$ and $\lambda = \pm 1$. Therefore, the contour integral in (18.2.15) can be computed with residues, and the result is independent of the choice of the contours $\Gamma_-$ and $\Gamma_+$. In this way, the operator $f(\tilde{k})$ is uniquely defined as a formal perturbation series which is well-defined and finite to every order.

We now establish the functional calculus.

**Theorem 18.2.1. (functional calculus)** For any functions $f, g$ which are holomorphic in discs around $\pm 1$ which contain the contours $\Gamma_\pm$,

$$(i\partial + B - m) f(\tilde{k}) = 0 \quad (18.2.18)$$

$$f(\tilde{k})^* = \overline{f(\tilde{k})} \quad (18.2.19)$$

$$f(\tilde{k}) \cdot g(\tilde{k}) = (fg)(\tilde{k}) \quad (18.2.20)$$

**Proof.** Since the image of the operator $\tilde{k}$ lies in the kernel of the Dirac operator, we know that

$$(i\partial + B - m) \tilde{R}_\lambda = (i\partial + B - m) (\lambda^{-1}).$$

Taking the contour integral (18.2.15) gives (18.2.18).

The operators $p_n, k_m$ and $s_m$ are obviously symmetric (see (18.2.9), (18.2.10) and (18.2.3)). According to (18.2.4), the operator $k_m$ is also symmetric. Hence the resolvent $\tilde{R}_\lambda$ defined by (18.2.14) has the property

$$\tilde{R}_\lambda^* = \tilde{R}_\lambda^\dagger.$$

The relation (18.2.19) follows by taking the adjoint of (18.2.15) and reparametrizing the integral.

The starting point for proving (18.2.20) is the resolvent identity

$$\tilde{R}_\lambda \cdot \tilde{R}_{\lambda'} = \frac{1}{\lambda - \lambda'} \left( \tilde{R}_\lambda - \tilde{R}_{\lambda'} \right). \quad (18.2.21)$$

We set $\Gamma = \Gamma_+ \cup \Gamma_-$ and denote the corresponding contour for $\lambda'$ by $\Gamma'$. Since the integral (18.2.15) is independent of the precise choice of the contour, we may choose

$$\Gamma = \partial B_\delta(1) \cup \partial B_\delta(-1) \quad \text{and} \quad \Gamma' = \partial B_\delta(1) \cup \partial B_\delta(-1).$$
for sufficiently small $\delta < 1/2$. Then $\Gamma$ does not enclose any point of $\Gamma'$, implying that
\[ \oint_{\Gamma} \frac{f(\lambda)}{\lambda - \lambda'} d\lambda = 0 \quad \text{for all } \lambda' \in \Gamma'. \tag{18.2.22} \]
On the other hand, $\Gamma'$ encloses every point of $\Gamma$, so that
\[ \oint_{\Gamma'} f(\lambda) g(\lambda') \frac{\tilde{R}_\lambda}{\lambda - \lambda'} d\lambda' = -2\pi i f(\lambda) g(\lambda) \tilde{R}_\lambda \quad \text{for all } \lambda \in \Gamma. \tag{18.2.23} \]
Combining (18.2.21) with (18.2.22) and (18.2.23), we obtain
\[ f(\tilde{k}) \cdot g(\tilde{k}) = -\frac{1}{4\pi^2} \oint_{\Gamma} f(\lambda) d\lambda \oint_{\Gamma'} g(\lambda') d\lambda' \frac{1}{\lambda - \lambda'} \left( \tilde{R}_\lambda - \tilde{R}_{\lambda'} \right) = -\frac{1}{2\pi i} \oint_{\Gamma} f(\lambda) g(\lambda) \tilde{R}_\lambda d\lambda = (fg)(\tilde{k}). \]
This concludes the proof. □

This functional calculus makes it possible to compute the unregularized kernel of the fermionic projector, as we now explain. Our starting point is the defining equation for the fermionic signature operator (18.0.4), which we can now write in the short form
\[ \tilde{k} \cdot \tilde{k} = \tilde{S}_m |_{t} \tilde{k}. \]
Iterating this relation, we obtain for any $p \in \mathbb{N}$
\[ (\tilde{k} \cdot \tilde{k})^p \tilde{k} = (\tilde{S}_m |_{t})^p \tilde{k} = (\tilde{S}_m)^p |_{t} \tilde{k} \quad \text{for all } p \in \mathbb{N}. \]
Consequently, this formula also holds for the functional calculus, i.e.
\[ f(\tilde{k}) \cdot \tilde{k} = f(\tilde{S}_m) |_{t} \tilde{k}. \]
This formula makes it possible to express the unregularized kernel $\tilde{P}_-$ in (15.4.1) by
\[ \tilde{P}_- = -\chi_{(-\infty,0)} (\tilde{S}_m) \tilde{k}_m = -\chi_{(-\infty,0)} (\tilde{S}_m) |_{t} \tilde{k} = \chi_{(-\infty,0)} (\tilde{k}) \cdot \tilde{k} = -\chi_{(0,\infty)} (\tilde{S}_m) \tilde{k} \cdot \tilde{k} = -\frac{1}{2\pi i} \oint_{\Gamma_-} \tilde{R}_\lambda d\lambda \cdot \tilde{k} = -\frac{1}{2\pi i} \oint_{\Gamma_-} (-\lambda) \tilde{R}_\lambda d\lambda \]
Substituting the perturbation expansion for $\tilde{R}_\lambda$ in (18.2.16) and writing the vacuum resolvent in the form (18.2.17), one can carry out the contour integral with residues. This gives the desired perturbation expansion for $P_-$. More details on this method and explicit formulas can be found in [81] Section 3.3 and Appendix A].

Exercises
CHAPTER 19

Methods of Microlocal Analysis

19.1. The Hadamard Expansion in Minkowski Space

In Chapter 15, the unregularized kernel of the fermionic projector $P(x, y) = P_-(x, y)$ was constructed abstractly. In Chapter 18, we saw how this kernel can be expanded in a perturbation series in powers of the external potential. In order to gain more explicit information on the form of the unregularized kernel, it is very useful to analyze its singularity structure on the light cone. It turns out that $P(x, y)$ has singularities on the light cone, which can be described by the so-called Hadamard expansion of the form

$$ P(x, y) = \lim_{\varepsilon \to 0} i \frac{\partial}{\partial x} \left( \frac{U(x, y)}{\Gamma_{\varepsilon}(x, y)} + V(x, y) \log \Gamma_{\varepsilon}(x, y) + W(x, y) \right), \quad (19.1.1) $$

where

$$ \Gamma_{\varepsilon}(x, y) := (y - x)^i (y - x)_j - i \varepsilon (y - x)_0, \quad (19.1.2) $$

and $U$, $V$ and $W$ are smooth functions on $\mathcal{M} \times \mathcal{M}$ taking values in the $4 \times 4$-matrices acting on the spinors (we always denote spacetime indices by latin letters running from $0, \ldots, 3$). This local expansion is based on the method of integration along characteristics which will be explained below (see after (19.1.9) or also [95, 83] or [6]). In Minkowski space, the light-cone expansion [37, 38] (see also [48, Section 2.2]) gives an efficient procedure for computing an infinite number of Hadamard coefficients in one step. The Hadamard form (19.1.1) carries over to curved spacetime. Moreover, there is an interesting connection to the so-called wave front set in microlocal analysis. These generalizations will be briefly outlined in Section 19.3 below. In all the other sections of this chapter, we restrict attention to Minkowski space.

It turns out that for an external potential in Minkowski space, the kernel of the fermionic projector is indeed of Hadamard form.

**Theorem 19.1.1.** Assume that the external potential $\mathcal{B}$ is smooth, and that its time derivatives decay at infinity in the sense that (17.2.1) holds and in addition that

$$ \int_{-\infty}^{\infty} |\partial^p_t \mathcal{B}(t)|_{C^0} \, dt < \infty \quad \text{for all } p \in \mathbb{N} $$

(with the $C^0$-norm as defined in (17.2.2)). Moreover, assume that the potential satisfies the bound

$$ \int_{-\infty}^{\infty} |\mathcal{B}(t)|_{C^0} \, dt < \sqrt{2} - 1. \quad (19.1.3) $$

Then the fermionic projector $P(x, y)$ is of Hadamard form.

The proof of this theorem will be given in Section 19.4 below.

We conclude this section by explaining how the expansion (19.1.1) comes about and how the involved functions $U$, $V$ and $W$, at least in principle, can be computed iteratively.
using the method of integration along characteristics. We begin by computing the unregularized kernel in the Minkowski vacuum. To this end, one rewrites the factor \((\mathbf{k} + m)\) in terms of a differential operator in position space,

\[
P(x, y) = (i \partial_x + m) T_{m^2}(x, y),
\]

where \(T_{m^2}\) is the scalar bi-distribution

\[
T_{m^2}(x, y) := \int \frac{d^4 k}{(2 \pi)^4} \delta(k^2 - m^2) \Theta(-k^0) e^{-ik(x-y)}. \tag{19.1.5}
\]

Away from the light cone (i.e. for \(\xi^2 \neq 0\)), \(T_{m^2}(x, y)\) is a smooth function given by

\[
T_{m^2}(x, y) = \begin{cases} 
\frac{m}{16 \pi^2} \frac{Y_1(m \sqrt{\xi^2})}{\sqrt{\xi^2}} + \frac{im}{16 \pi^2} \frac{J_1(m \sqrt{\xi^2})}{\sqrt{\xi^2}} \epsilon(\xi^0) & \text{if } \xi \text{ is timelike} \\
\frac{m}{8 \pi^3} \frac{K_1(m \sqrt{-\xi^2})}{\sqrt{-\xi^2}} & \text{if } \xi \text{ is spacelike}
\end{cases}
\tag{19.1.6}
\]

where \(J_1, Y_1\) and \(K_1\) are Bessel functions. Expanding the Bessel functions in (19.1.6) in a power series, one obtains (see [114, (10.2.2), (10.8.1) and (10.25.2), (10.31.1)])

\[
T_{m^2}(x, y) = -\frac{1}{8 \pi^3} \frac{\mathrm{PP}}{(y-x)^2} + i \pi \delta((y-x)^2) \epsilon((y-x)^0) \\
+ \frac{m^2}{32 \pi^3} \sum_{j=0}^{\infty} \frac{(-1)^j}{j! (j+1)!} \left( \frac{m^2(y-x)^2)^j}{4^j} \right) \\
\times \left( \log |m^2(y-x)^2| + c_j + i \pi \Theta((y-x)^2) \epsilon((y-x)^0) \right)
\]

with real coefficients \(c_j\) (here \(\Theta\) and \(\epsilon\) are again the Heaviside and the sign function, respectively). In particular, one sees that \(T_{m^2}\) is a distribution which is singular on the light cone. These singularities can be written in a shorter form using residues as

\[
T_{m^2}(x, y) = \lim_{\epsilon \to 0} \left( -\frac{1}{8 \pi^3} \frac{1}{\Gamma_\epsilon(x, y)} \\
+ \frac{m^2}{32 \pi^3} \sum_{j=0}^{\infty} \frac{(-1)^j}{j! (j+1)!} \left( \frac{m^2 \Gamma_\epsilon(x, y)^j}{4^j} \right) \right) \log \left( m^2 \Gamma_\epsilon(x, y) + c_j \right). \tag{19.1.7}
\]

Noting that the series converge, one obtains a function of the desired form as in the brackets in (19.1.1). This shows that the term \(i \partial_x T_{m^2}(x, y)\) in (19.1.4) is of Hadamard form. For the term \(m T_{m^2}(x, y)\) this can be shown by pulling one derivative and working with matrix-valued kernels. Indeed,

\[
m T_{m^2}(x, y) = -\frac{1}{m} \Box x \left( T_{m^2}(x, y) - T_0(x, y) \right) = \partial_x \left\{ -\frac{1}{m} \partial_x \left( T_{m^2}(x, y) - T_0(x, y) \right) \right\},
\]

and computing the curly brackets by differentiating (19.1.7) one obtains again an expression of the Hadamard form (19.1.4).
The summands in (19.1.7) can be understood by verifying that $T_{m^2}$ satisfies the Klein-Gordon equation. Indeed, using the abbreviation $\xi := y - x$, we obtain

$$
\frac{\partial}{\partial x^j} \left( \frac{1}{\Gamma_\epsilon (x, y)} \right) = \frac{\partial_j \Gamma_\epsilon (x, y)}{\Gamma_\epsilon (x, y) \Gamma_\epsilon (x, y)^2} = \frac{1}{\Gamma_\epsilon(x,y)^2} \left( 2\xi_j - i\varepsilon \delta_{j,0} \right) 
$$

and this tends to zero as $\varepsilon \searrow 0$. Thus the leading term in (19.1.7) satisfies the scalar wave equation. In the Klein-Gordon equation, however, the term involving the mass remains,

$$
- m^2 \frac{1}{\Gamma_\epsilon(x,y)}. \tag{19.1.8} 
$$

This term is compensated by the next term in the expansion (19.1.7), because

$$
\frac{\partial}{\partial x^j} \log \left( \Gamma_\epsilon (x, y) \right) = \frac{\partial_j \Gamma_\epsilon (x, y)}{\Gamma_\epsilon (x, y)^2} = \frac{1}{\Gamma_\epsilon(x,y)^2} \left( 2\xi_j - i\varepsilon \delta_{j,0} \right) 
$$

$$
\square x \log \left( \Gamma_\epsilon (x, y) \right) = -\frac{1}{\Gamma_\epsilon(x,y)^2} \left( 2\xi_j - i\varepsilon \delta_{j,0} \right) \left( 2\xi_j - i\varepsilon \delta_{j,0}^\epsilon \right) + \frac{8}{\Gamma_\epsilon(x,y)} 
$$

$$
= -\frac{4}{\Gamma_\epsilon(x,y)^2} \left( 4\xi_j^2 - 4i\varepsilon \xi_0 - \varepsilon^2 \right) + \frac{8}{\Gamma_\epsilon(x,y)} 
$$

Now the first summand in the last line cancels the term (19.1.8) in the Klein-Gordon equation. Proceeding order by order in powers of $\Gamma_\epsilon(x,y)$, one can verify all the coefficients in (19.1.7).

This method of applying the wave operator term by term is also useful for computing the functions $U$, $V$ and $W$ in (19.1.1) in the case that an external potential is present. In fact, these functions can be expressed in terms of line integrals along the light cone. This method of integration along characteristics goes back to Hadamard [95] and is described in the classic textbook [83] in curved spacetime. In order to explain the method in the simplest possible context, let us assume that we consider the wave equation with an external scalar potential $a(x)$, i.e.

$$
\left( - \square_x - a(x) \right) \tilde{T}(x, y) = 0
$$

(the Dirac equation will be treated more systematically in Section 19.2). In modification of the series in (19.1.7) we make the ansatz

$$
\tilde{T}(x, y) = \lim_{\varepsilon \searrow 0} \left( \frac{1}{\Gamma_\epsilon(x,y)} + \sum_{n=1}^{\infty} f_n (x, y) \Gamma_\epsilon(x,y)^n \log \left( \Gamma_\epsilon(x,y) \right) \right) \tag{19.1.9} 
$$

Compared to (19.1.8), now the the error term of the first summand involves the potential $a(x)$,

$$
- \frac{a(x)}{\Gamma_\epsilon(x,y)}. \tag{19.1.10} 
$$
The hope is to compensate this term by a suitable choice of \( f_1(x, y) \). Indeed,
\[
\partial_{x^j} \left( f_1(x, y) \log \Gamma_\varepsilon(x, y) \right) = f_1(x, y) \frac{\partial_j \Gamma_\varepsilon(x, y)}{\Gamma_\varepsilon(x, y)} + \partial_j f_1(x, y) \log \Gamma_\varepsilon(x, y)
\]
\[
\Box_x f_1(x, y) \log \Gamma_\varepsilon(x, y) = f_1(x, y) \frac{4}{\Gamma_\varepsilon(x, y)} - 2 \partial_j f_1(x, y) \frac{2 \xi_j}{\Gamma_\varepsilon(x, y)} + \cdots ,
\]
where \( \cdots \) stands for all terms which either have a lower order singularity on the light cone or tend to zero as \( \varepsilon \to 0 \). In order for this contribution to compensate (19.1.10), the function \( f_1 \) must satisfy the equation
\[
4 f_1(x, y) - 4 \xi^j \partial_j f_1(x, y) = a(x).
\]
Such a differential equation of first order can be solved with the method of characteristics (see for example [31], Section I.3.2). More specifically, the solution is an integral along the straight line \( \xi \mathbb{R} \). In order to describe the singular behavior on the light cone, it suffices to consider the case that \( \xi \) is tangential to the light cone. Similarly, also to higher order in the expansion parameter \( n \), we obtain transport equations along the light cone, which can be solved iteratively order by order.

### 19.2. The Light-Cone Expansion

We first give the basic definition of the light-cone expansion and explain it afterward.

**Definition 19.2.1.** A distribution \( A(x, y) \) on \( \mathcal{M} \times \mathcal{M} \) is of the order \( O((y - x)^{2p}) \) for \( p \in \mathbb{Z} \) if the product
\[
(y - x)^{-2p} A(x, y)
\]
is a regular distribution (i.e., a locally integrable function). An expansion of the form
\[
A(x, y) = \sum_{j=g}^{\infty} A^{[j]}(x, y)
\]
with \( g \in \mathbb{Z} \) is called light-cone expansion if the \( A^{[j]}(x, y) \) are distributions of the order \( O((y - x)^{2j}) \) and if \( A \) is approximated by the partial sums in the sense that for all \( p \geq g \),
\[
A(x, y) - \sum_{j=g}^{p} A^{[j]}(x, y)
\]
is of the order \( O((y - x)^{2p+2}) \).

The parameter \( g \) gives the leading order of the singularity of \( A(x, y) \) on the light cone. We point out that we do not demand that the infinite series in (19.2.1) converges. Thus, similar to a formal Taylor series, the series in (19.2.1) is defined only via the approximation by the partial sums (19.2.2). The notion of the light-cone expansion is illustrated in Exercise 19.1.

As a concrete example, due to the factors \( \Gamma_\varepsilon(x, y) \), the series (19.1.7) is a light-cone expansion. The term with the leading singularity becomes integrable after multiplying by \( (y - x)^2 \), showing that \( g = -1 \).

Our task is to perform the light-cone expansion of the unregularized kernel of the fermionic projector. Schematically, this construction consists of several steps:

1. Perform the light-cone expansion of the causal Green’s operators \( \tilde{s}_m^\vee \) and \( \tilde{s}_m^\wedge \). Here one proceeds inductively for each summand of the perturbation series (27.3.4).
2. Using the relation (18.2.1), one obtains a corresponding light-cone expansion for the causal fundamental solution \( \tilde{k}_m \).
The so-called residual argument relates the sought-after light-cone expansion of \( \tilde{P}(x, y) \) to that of \( \tilde{k}_m \).

This procedure is described in detail in [48, Chapter 2]. In order to avoid an unnecessary overlap, we here focus on the light-cone expansion of the causal Green’s operators and only introduce the concepts needed for the basics on the continuum limit in Chapter 21. Before doing so, we illustrate the light-cone expansion by a simple example.

Example 19.2.2. Consider the massless Dirac equation in the presence of an external electromagnetic potential \( A \),

\[
(i\slashed{\partial} + \mathcal{A}) \hat{P}(x, y) .
\]

For simplicity assume that \( A \) is smooth and compactly supported in spacetime. Then, to first order in perturbation theory, the light-cone expansion of the unregularized kernel \( \hat{P}(x, y) \) takes the form

\[
\hat{P}(x, y) = \frac{i}{2} \exp \left( -i \int_0^1 A_j \big|_{\alpha y + (1-\alpha)x} \xi^j \, d\alpha \right) P(x, y) \quad (19.2.3)
\]

\[
- \frac{1}{2} \xi \xi \int_0^1 (\alpha - \alpha^2) j^i \big|_{\alpha y + (1-\alpha)x} \, d\alpha \, T^{(0)} \quad (19.2.4)
\]

\[
+ \frac{1}{4} \xi \int_0^1 F^{ij} \big|_{\alpha y + (1-\alpha)x} \gamma_i \gamma_j \, d\alpha \, T^{(0)} \quad (19.2.5)
\]

\[
- \xi \int_0^1 (1 - \alpha) F^{ij} \big|_{\alpha y + (1-\alpha)x} \gamma_j \, d\alpha \, T^{(0)} \quad (19.2.6)
\]

\[
- \xi \int_0^1 (1 - \alpha) (\alpha - \alpha^2) \slashed{\partial} j^i \big|_{\alpha y + (1-\alpha)x} \, d\alpha \, T^{(1)} \quad (19.2.7)
\]

\[
+ \xi (\deg < 1) + (\deg < 0) + O(A^2) ,
\]

where \( F^{jk} = \partial^j A^k - \partial^k A^j \) is the field tensor and \( j^k = \partial^k A^j - \Box A^k \) is the corresponding Maxwell current. Moreover, the factors \( T^{(0)} \) and \( T^{(1)} \) are the leading summands in (19.1.7); more precisely,

\[
T^{(0)}(x, y) = -\frac{1}{8\pi^3} \lim_{\varepsilon \searrow 0} \frac{1}{\Gamma_\varepsilon(x, y)} \quad (19.2.9)
\]

\[
T^{(1)}(x, y) = \frac{1}{32\pi^3} \lim_{\varepsilon \searrow 0} \log \Gamma_\varepsilon(x, y) .
\]

Each summand has the general structure of being the product of a smooth function and a distribution which is singular on the light cone. The smooth factor is an integral along the straight line segment joining the points \( x \) and \( y \). The integrand involves the electromagnetic potential and its partial derivatives. We remark for clarity that the term \( (19.2.3) \) involves a gauge phase as needed for gauge invariance (as already mentioned in (5.7.5) in Section 5.7). All the other integrands are gauge invariant, as is obvious from the fact that they are expressed in terms of the electromagnetic field tensor and the Maxwell current.

To higher order on the light cone or to higher order in the mass or the external potentials, the formulas of the light-cone expansions have a similar structure. More
detailed formulas can be found in the original papers [37, 38], in [39, Appendix B] and [48, Appendix B].

We now explain how to perform the light-cone expansion of the causal Green’s operators. In order to get a first idea for how to proceed, we begin by considering the free advanced Green’s function \( s^\lor_m \) of the Dirac equation of mass \( m \) in position space: Similar to (19.1.4), it is again convenient to pull the Dirac matrices out of explicit formula (19.2.13). Apart from giving a motivation for the desired form (19.2.15)

\[
s^\lor_m(x, y) = (i\dot{\vec{\gamma}}_x + m) S^\lor_{m^2}(x, y),
\]

where \( S^\lor_{m^2} \) is the advanced Green’s function of the Klein-Gordon operator,

\[
S^\lor_{m^2}(x, y) = \lim_{\nu \to 0} \int \frac{d^4p}{(2\pi)^4} \frac{1}{p^2 - m^2 - i\nu} e^{-ip(x-y)}.
\]

Computing this Fourier integral and expanding the resulting Bessel function in a power series gives (for details see Exercise 19.2)

\[
S^\lor_{m^2}(x, y) = -\frac{1}{2\pi} \delta((y-x)^2) \Theta(y^0 - x^0)
+ \frac{m^2}{4\pi} J_1 \left( \frac{m^2(y-x)^2}{\sqrt{m^2(y-x)^2}} \right) \Theta((y-x)^2) \Theta(y^0 - x^0)
= -\frac{1}{2\pi} \delta((y-x)^2) \Theta(y^0 - x^0)
+ \frac{m^2}{8\pi} \sum_{j=0}^{\infty} \frac{(-1)^j}{j!(j+1)!} \left( \frac{m^2(y-x)^2}{4} \right)^j \Theta((y-x)^2) \Theta(y^0 - x^0).
\]

This computation shows that \( S^\lor_{m^2}(x, y) \) has a \( \delta((y-x)^2) \)-like singularity on the light cone. Furthermore, one sees that \( S^\lor_{m^2} \) is a power series in \( m^2 \). The important point for what follows is that the higher order contributions in \( m^2 \) contain more factors \( (y-x)^2 \) and are thus of higher order on the light cone. More precisely,

\[
\left( \frac{d}{dm^2} \right)^n S^\lor_{m^2}(x, y) \bigg|_{m=0} \quad \text{is of the order } \mathcal{O}((y-x)^{2n-2}).
\]

According to (19.2.10), the Dirac Green’s function is obtained by computing the first partial derivatives of (19.2.13). Therefore, \( s^\lor_m(x, y) \) has a singularity on the light cone which is even \( \sim \delta'((y-x)^2) \). The higher order contributions in \( m \) are again of increasing order on the light cone. This means that we can view the Taylor expansion of (19.2.10) in \( m \),

\[
s^\lor_m(x, y) = \sum_{n=0}^{\infty} (i\dot{\vec{\gamma}} + m) \frac{m^{2n}}{n!} \left( \frac{d}{dm^2} \right)^n S^\lor_{m^2}(x, y) \bigg|_{m=0},
\]

as a light-cone expansion of the free Green’s function. Our idea is to generalize this formula to the case with interaction. More precisely, we want to express the perturbed Green’s function in the form

\[
\hat{s}^\lor(x, y) = \sum_{n=0}^{\infty} F_n(x, y) \left( \frac{d}{dm^2} \right)^n S^\lor_{m^2}(x, y) \bigg|_{m=0}
\]

with factors \( F_n \) which depend on the external potential. We will see that this method is very convenient; especially, we can in this way avoid working with the rather complicated explicit formula (19.2.13). Apart from giving a motivation for the desired form (19.2.15)
of the formulas of the light-cone expansion, the mass expansion (19.2.13) leads to the conjecture that even the higher order contributions in the mass to the perturbed Green’s functions might be of higher order on the light cone. If this conjecture was true, it would be a good idea to expand the perturbation expansion of \( \tilde{s} \) with respect to the parameter \( m \). Therefore, our strategy is to first expand (27.3.4) with respect to the mass and to try to express the contributions to the resulting expansion in a form similar to (19.2.15).

The expansion of (27.3.4) with respect to \( m \) gives a double sum over the orders in the mass parameter and in the external potential. It is convenient to combine these two expansions in a single perturbation series. To this end, we rewrite the Dirac operator as

\[
i\partial + B - m = i\partial + B \quad \text{with} \quad B := B - m.
\]

(19.2.16)

For the light-cone expansion of the Green’s functions, we will always view \( B \) as the perturbation of the Dirac operator. This has the advantage that the unperturbed objects are massless. Expanding in powers of \( B \) gives the mass expansion and the perturbation expansion in one step. In order to simplify the notation, for the massless objects we usually omit the index \( m \). Thus we write the Green’s function of the massless Dirac equation in the Minkowski vacuum as

\[
\begin{align*}
s^{\vee}(x, y) &= i\partial_x S^{\vee}_{m^2}(x, y) \big|_{m=0}, \\
s^{\wedge}(x, y) &= i\partial_x S^{\wedge}_{m^2}(x, y) \big|_{m=0}.
\end{align*}
\]

(19.2.17)

Then the interacting Green’s functions are given by the perturbation series

\[
\begin{align*}
\tilde{s}^{\vee} &= \sum_{k=0}^{\infty} (-s^{\vee}B)^k s^{\vee}, \\
\tilde{s}^{\wedge} &= \sum_{k=0}^{\infty} (-s^{\wedge}B)^k s^{\wedge}.
\end{align*}
\]

(19.2.18)

The constructions of the following subsections are exactly the same for the advanced and retarded Green’s functions. In order to treat both cases at once, in the remainder of this section we will omit all superscripts ‘\( \vee \)’, ‘\( \wedge \)’. The formulas for the advanced and retarded Green’s functions are obtained by either adding ‘\( \vee \)’ or ‘\( \wedge \)’ to all factors \( s, S \).

We now explain how the individual contributions to the perturbation expansion (19.2.18) can be written similar to the right side of (19.2.15) as a sum of terms of increasing order on the light cone. For the mass expansion of \( S^{m^2} \), we set \( a = m^2 \) and use the notation

\[
S^{(l)} = \left( \frac{d}{da} \right)^l S_a \big|_{a=0}.
\]

(19.2.19)

In preparation, we derive some computation rules for the \( S^{(l)} \): \( S_a \) satisfies the defining equation of a Klein-Gordon Green’s function

\[
(-\Box - a) S_a(x, y) = \delta^4(x - y).
\]

Differentiating with respect to \( a \) and setting \( a = 0 \) gives

\[
-\Box S^{(l)}(x, y) = \delta_{l,0} \delta^4(x - y) + l S^{(l-1)}(x, y), \quad l \geq 0.
\]

(19.2.20)

(For \( l = 0 \), this formula does not seem to make sense because \( S^{(-1)} \) is undefined. The expression is meaningful, however, if one keeps in mind that in this case the factor \( l \) is zero, and thus the whole second summand vanishes. We will also use this convention in the following calculations.) Next, we differentiate the formulas for \( S_a \) in momentum space,

\[
S_a^{\vee}(p) = \frac{1}{p^2 - a - i\nu p^0}, \quad S_a^{\wedge}(p) = \frac{1}{p^2 - a + i\nu p^0}
\]

(19.2.21)
with respect to both $p$ and $a$. Comparing the results gives the relation

$$\frac{\partial}{\partial p^k} S_a(p) = -2p_k \frac{d}{da} S_a(p),$$

or, after expanding in the parameter $a$,

$$\frac{\partial}{\partial p^k} S(l)^{(i)}(p) = -2p_k S(l+1)^{(i)}(p), \quad l \geq 0. \quad (19.2.22)$$

This formula also determines the derivatives of $S(l)^{(i)}$ in position space; namely

$$\frac{\partial}{\partial x^k} S(l)^{(i)}(x, y) = \int \frac{d^4 p}{(2\pi)^4} S(l)^{(i)}(p) (-ip_k) e^{-ip(x-y)}$$

$$= \frac{i}{2} \int \frac{d^4 p}{(2\pi)^4} \frac{\partial}{\partial p^k} S(l-1)^{(i)}(p) e^{-ip(x-y)}$$

$$= -i \frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} S(l-1)^{(i)}(p) \frac{\partial}{\partial p^k} e^{-ip(x-y)}$$

$$= \frac{1}{2} (y - x)_k S(l-1)^{(i)}(x, y), \quad l \geq 1. \quad (19.2.23)$$

We iterate this relation to calculate the Laplacian,

$$-\Box x S(l)^{(i)}(x, y) = -\frac{1}{2} \frac{\partial}{\partial x^k} \left( (y - x)^k S(l-1)^{(i)}(x, y) \right)$$

$$= 2 S(l-1)^{(i)}(x, y) + \frac{1}{4} (y - x)^2 S(l-2)^{(i)}(x, y), \quad l \geq 2.$$

After comparing with (19.2.20), we conclude that

$$(y - x)^2 S(l)^{(i)}(x, y) = -4l S(l+1)^{(i)}(x, y), \quad l \geq 0. \quad (19.2.24)$$

Finally, $S(l)^{(i)}(x, y)$ is only a function of $(y - x)$, which implies that

$$\frac{\partial}{\partial x^k} S(l)^{(i)}(x, y) = -\frac{\partial}{\partial y^k} S(l)^{(i)}(x, y), \quad l \geq 0. \quad (19.2.25)$$

The following lemma gives the light-cone expansion of an operator product which is linear in the external potential. We will later use it for the iterative light-cone expansion of more complicated operator products; in this case, the potential will be a composite expression in $B$ and its partial derivatives. In order to avoid confusion then, we denote the external potential by $V$.

**Lemma 19.2.3.** (light-cone expansion to first order) For any $l, r \geq 0$, the operator product $S(l)^{(i)} V S(r)^{(j)}$ has the light-cone expansion

$$(S(l)^{(i)} V S(r)^{(j)})(x, y) = \sum_{n=0}^{\infty} \frac{1}{n!} \int_0^1 \alpha^l (1 - \alpha)^r (\alpha \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S(n+l+r+1)(x, y). \quad (19.2.26)$$

**Proof.** The method of proof is to first compute the Laplacian of both sides of (19.2.26). The resulting formulas will have a similar structure, making it possible to proceed inductively.

On the left side of (19.2.26), we calculate the Laplacian with the help of (19.2.20) to obtain

$$-\Box x (S(l)^{(i)} V S(r)^{(j)})(x, y) = \delta_{l,0} V(x) S(r)^{(j)}(x, y) + l (S(l-1)^{(i)} V S(r)^{(j)})(x, y). \quad (19.2.27)$$
The Laplacian of the integral on the right side of (19.2.26) can be computed with (19.2.23) and (19.2.20),

\[ -\Box_x \int_0^1 \alpha^l (1 - \alpha)^r (\alpha - \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r+1)}(x, y) \]

(19.2.28)

\[ = - \int_0^1 \alpha^l (1 - \alpha)^{r+2} (\alpha - \alpha^2)^n (\Box^{n+1} V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r+1)}(x, y) \]

\[ - \int_0^1 \alpha^l (1 - \alpha)^{r+1} (\alpha - \alpha^2)^n (\partial_k \Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, (y - x)^k \, S^{(n+l+r)}(x, y) \]

\[ + (n + l + r + 1) \int_0^1 \alpha^l (1 - \alpha)^r (\alpha - \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r)}(x, y). \]

In the second summand, we rewrite the partial derivative as a derivative with respect to \( \alpha \),

\[ (y - x)^k (\partial_k \Box^n V)_{\alpha y + (1 - \alpha)x} = \frac{d}{d\alpha} (\Box^n V)_{\alpha y + (1 - \alpha)x} \]

(as is verified immediately by computing the right side with the chain rule). This makes it possible to integrate in \( \alpha \) by parts. We thus obtain

\[ \int_0^1 \alpha^l (1 - \alpha)^{r+1} (\alpha - \alpha^2)^n (\partial_k \Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, (y - x)^k \]

\[ = \int_0^1 \alpha^l (1 - \alpha)^{r+1} (\alpha - \alpha^2)^n \frac{d}{d\alpha} (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \]

\[ = -\delta_{n,0} \delta_{l,0} V(x) - (n + l) \int_0^1 \alpha^l (1 - \alpha)^{r+2} (\alpha - \alpha^2)^{n-1} (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \]

\[ + (n + r + 1) \int_0^1 \alpha^l (1 - \alpha)^r (\alpha - \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \]

\[ = -\delta_{n,0} \delta_{l,0} V(x) \]

\[ - n \int_0^1 \alpha^l (1 - \alpha)^{r+2} (\alpha - \alpha^2)^{n-1} (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \]

\[ + (n + l + r + 1) \int_0^1 \alpha^l (1 - \alpha)^r (\alpha - \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \]

\[ - l \int_0^1 \alpha^{l-1} (1 - \alpha)^r (\alpha - \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha. \]

We substitute back into the original equation to obtain

\[ (19.2.28) = \delta_{n,0} \delta_{l,0} V(x) \, S^{(r)}(x, y) \]

\[ + l \int_0^1 \alpha^{l-1} (1 - \alpha)^r (\alpha - \alpha^2)^n (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r)}(x, y) \]

\[ - \int_0^1 \alpha^l (1 - \alpha)^{r+2} (\alpha - \alpha^2)^n (\Box^{n+1} V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r+1)}(x, y) \]

\[ + n \int_0^1 \alpha^l (1 - \alpha)^{r+2} (\alpha - \alpha^2)^{n-1} (\Box^n V)_{\alpha y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r)}(x, y). \]
After dividing by \(n!\) and summation over \(n\), the last two summands are telescopic and cancel each other. Thus one gets

\[
- \Box \sum_{n=0}^{\infty} \frac{1}{n!} \int_{0}^{1} \alpha^l \left(1 - \alpha\right)^r \left(\alpha - \alpha^2\right)^n \left(\Box^n V\right)_{\partial y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r+1)}(x, y)
\]

\[
= \delta_{l,0} V(x) \, S^{(r)}(x, y)
\]

\[
+ \sum_{n=0}^{\infty} \frac{1}{n!} \int_{0}^{1} \alpha^{l-1} \left(1 - \alpha\right)^r \left(\alpha - \alpha^2\right)^n \left(\Box^n V\right)_{\partial y + (1 - \alpha)x} \, d\alpha \, S^{(n+l+r)}(x, y) .
\]

(19.2.29)

We now compare the formulas (19.2.27) and (19.2.29) for the Laplacian of both sides of (19.2.26). In the special case \(l = 0\), these formulas coincide, and we can use a uniqueness argument for the solutions of the wave equation to prove (19.2.26): We assume that we consider the advanced Green’s function (for the retarded Green’s function, the argument is analogous). For given \(y\), we denote the difference of both sides of (19.2.26) by \(F(x)\).

Since the support of \(F(x)\) is in the past light cone \(x \in L_0^\wedge y\), \(F\) vanishes in a neighborhood of the hypersurface \(\mathcal{H} = \{z \in \mathbb{R}^4 | z^0 = y^0 + 1\}\). Moreover, the Laplacian of \(F\) is identically zero according to (19.2.27) and (19.2.29). We conclude that

\[
\Box F = 0 \quad \text{and} \quad F|_{\mathcal{H}} = \partial_{k} F|_{\mathcal{H}} = 0 .
\]

Since the wave equation has a unique solution for given initial data on the Cauchy surface \(\mathcal{H}\), \(F\) vanishes identically.

The general case follows by induction in \(l\): Suppose that (19.2.26) holds for given \(\hat{l}\) (and arbitrary \(r\)). Then, according to (19.2.27), (19.2.29), and the induction hypothesis, the Laplacian of both sides of (19.2.26) coincides for \(l = \hat{l} + 1\). The above uniqueness argument for the solutions of the wave equation again gives (19.2.26). \(\square\)

We recall for clarity that, according to (19.2.14), the higher \(\alpha\)-derivatives of \(S_{\alpha}(x, y)\) are of higher order on the light cone. Thus the summands in (19.2.26) are of increasing order on the light cone, and the infinite sum is mathematically well-defined in the sense of Definition (19.2.1) via the approximation by the partial sums (19.2.2).

19.3. The Hadamard Form in Curved Spacetime and the Wave Front Set

The Hadamard expansion (19.1.1) can also be formulated in curved spacetime. To this end, one simply replaces the function (19.1.2) by

\[
\Gamma_\varepsilon(x, y) := \Gamma(x, y) - i\varepsilon \left(t(y) - t(x)\right) ,
\]

where \(t\) is a time function and \(\Gamma(x, y)\) is the geodesic distance squared, with the sign convention that \(\Gamma\) is positive in timelike and negative in spacelike directions. If a bi-distribution is of Hadamard form in one chart, it is also of Hadamard form in another chart. More details on the Hadamard expansion for Dirac fields can be found in [127, 94] or [55, Appendix A].

Moreover, the Hadamard form can be formulated alternatively in terms of the wave front set, as we will now briefly mention. We work in an open subset \(U \subset \mathbb{R}^n\). We denote the distributions in \(U\) by \(\mathcal{D}'(U)\) (being the dual space of \(C^\infty(U, \mathbb{C})\) with the topology induced by the \(C^k\)-norms). An open conic neighborhood of a point \(\xi \in \mathbb{R}^n\) is defined to
be an open neighborhood which is invariant under the action of \( \mathbb{R}^+ \) by multiplication. Thus an open conic neighborhood is of the form
\[
\{ \lambda x \mid x \in S, \lambda \in \mathbb{R}^+ \},
\]
where \( S \) is an open subset of \( S^{n-1} \subset \mathbb{R}^n \).

**Definition 19.3.1.** Let \( \phi \in \mathcal{D}'(U) \). The wave front set \( \text{WF}(\phi) \) is the complement in \( U \times \mathbb{R}^n \setminus \{0\} \) of all points \((x, \xi) \in U \times \mathbb{R}^n \setminus \{0\}\) with the following property: There exists a function \( f \in C^\infty(U, \mathbb{R}) \) with \( f(x) = 1 \) and an open conic neighborhood \( V \) of \( \xi \) such that
\[
\sup_{\zeta \in V} (1 + |\zeta|)^N \left| (\hat{f}\phi)(\zeta) \right| < \infty \quad \text{for all } N \in \mathbb{N}.
\] (19.3.1)
This definition can be understood as follows. First, in view of taking the complement, the condition (19.3.1) ensures that the point \((x, \xi)\) does not lie in the wave front set. With the help of the cutoff function \( f \) one can disregard the behavior of \( \phi \) away from \( x \). In other words, the condition 19.3.1 only depends on the behavior of \( \phi \) in a arbitrarily small neighborhood of \( x \). This condition states that the Fourier transform has rapid decay in a cone around \( \xi \). Since decay properties of the Fourier transform correspond to smoothness properties in position space, we obtain a smoothness statement for \( \phi \) at \( x \), but only along the “wave front” described by \( \xi \).

- Eventuell Beispiele oder Übungsaufgaben? Beispielsweise \( \phi(x) = \delta^2(x) \) for \( U = \mathbb{R}^2 \)? Oder auch \( P(x, y) \) im Minkowski-Vakuum?

Definition 19.3.1 readily extends to a distribution \( \phi \) on a manifold \( M \), in which case the wave front set is a subset of the cotangent bundle,
\[
\text{WF}(\phi) \subset T^*M \setminus 0
\]
(where \( 0 \) is the zero section). The wave front set can also be defined for bundle-valued distributions by choosing a local trivialization and taking the wave front sets of the component functions. The unregularized kernel of the fermionic projector \( P \) is a bi-distribution on \( M \times M \). Therefore, its wave front set takes values in the product of the cotangent bundles,
\[
\text{WF} P \subset (T^*M \setminus 0) \times (T^*M \setminus 0).
\]

**Definition 19.3.2.** The unregularized kernel \( P(x, y) \) is said to be of **Hadamard form** if its wave front set has the property
\[
\text{WF} P \subset \left\{ (x_1, \xi_1, x_2, -\xi_2) \mid \text{there is a null geodesic } \gamma : I \to M \text{ with } a, b \in I, \gamma(a) = x_1, \gamma(b) = x_2 \text{ and } \xi_1 = \dot{\gamma}(a), \xi_2 = \dot{\gamma}(b) \text{ past-directed} \right\}.
\]
Intuitively, this definition means that there are singularities only on the light cone, and that these singularities are formed only of negative frequencies. The equivalence of this definition with the local Hadamard expansion (19.1.1) has been established in [119]. Good references on microlocal analysis and the wave front set are [100] and [4, Chapter 4].

### 19.4. Proof of the Hadamard Property in an External Potential

In this section, we give the proof of Theorem [19.1.1]. We closely follow the presentation in [75]. In preparation, we derive so-called frequency splitting estimates which give control of the “mixing” of the positive and negative frequencies in the solutions of the Dirac equation as caused by the time-dependent external potential (Theorem [19.4.1]).
Based on these estimates, we will complete the proof of Theorem 19.1.1 at the end of Section 19.4.2.

19.4.1. Frequency Mixing Estimates. For the following constructions, we again choose the hypersurface \( \mathcal{N} := \mathcal{N}_{t_0} \) at some given time \( t_0 \). Moreover, we always fix the mass parameter \( m > 0 \). Since we are no longer considering families of solutions, for ease in notation we omit the index \( m \) at the Dirac wave functions, the scalar products and the corresponding norms. We also identify the solution space \( \mathcal{H}_m \) with the Hilbert space \( \mathcal{H}_{t_0} \) of square integrable wave functions on \( \mathcal{N} \). On \( \mathcal{H}_{t_0} \), we can act with the Hamiltonian \( H \) of the vacuum, and using the above identification, the operator \( H \) becomes an operator on \( \mathcal{H}_m \) (which clearly depends on the choice of \( t_0 \)).

We work with a so-called frequency splitting with respect to the vacuum dynamics. To this end, we decompose the Hilbert space \( \mathcal{H}_m \) as
\[
\mathcal{H}_m = \mathcal{H}_m^+ \oplus \mathcal{H}_m^-
\]
with
\[
\mathcal{H}_\pm = \chi_\pm(H) \mathcal{H}_m,
\]
where \( \chi_\pm \) are the characteristic functions
\[
\chi_+ := \chi_{[0, \infty)} \quad \text{and} \quad \chi_- := \chi_{(-\infty, 0)}.
\]
(19.4.1)

For convenience, we write this decomposition in components and use a block matrix notation for operators, i.e.
\[
\psi = \begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} \quad \text{and} \quad A = \begin{pmatrix} A_+^+ & A_+^- \\ A_-^+ & A_-^- \end{pmatrix},
\]
where \( A_s^s' = \chi^s(H)A \chi^{s'}(H) \) and \( s, s' \in \{\pm\} \).

The representation in Proposition 17.2.5 makes it possible to let the fermionic signature operator \( \tilde{S}_m \) act on the Hilbert space \( \mathcal{H}_m \) (for fixed \( m \)). We decompose this operator with respect to the above frequency splitting,
\[
\tilde{S}_m = S^D + \Delta \tilde{S}, \quad \text{where} \quad S^D := \tilde{S}_+^+ + \tilde{S}_-^- \quad \text{and} \quad \Delta \tilde{S} := \tilde{S}_+^- + \tilde{S}_-^+.
\]
Thus the operator \( S^D \) maps positive to positive and negative to negative frequencies. The operator \( \Delta \tilde{S} \), on the other hand, mixes positive and negative frequencies. In the next theorem, it is shown under a suitable smallness assumption on \( \mathcal{B} \) that the operators \( \chi_\pm(\tilde{S}_m) \) coincide with the projections \( \chi_\pm(H) \), up to smooth contributions. The main task in the proof is to control the “frequency mixing” as described by the operator \( \Delta \tilde{S} \).

**Theorem 19.4.1.** Under the assumptions of Theorem 19.1.1, the operators \( \chi_\pm(\tilde{S}_m) \) have the representations
\[
\chi_\pm(\tilde{S}_m) = \chi_\pm(H) + \frac{1}{2\pi i} \oint_{\partial B_{\frac{1}{2}} (\pm 1)} (\tilde{S}_m - \lambda)^{-1} \Delta \tilde{S} (S^D - \lambda)^{-1} d\lambda,
\]
(19.4.2)
where the contour integral is an integral operator with a smooth integral kernel.

Here \( B_{\frac{1}{2}} \) denotes the open ball of radius \( 1/2 \). The operator \( (\tilde{S}_m - \lambda)^{-1} \) is also referred to as the resolvent of \( \tilde{S}_m \).

This theorem will be proved in several steps. We begin with a preparatory lemma.

**Lemma 19.4.2.** Under the assumptions (17.2.1) and (19.1.3), the spectrum of \( S^D \) is located in the set
\[
\sigma(S^D) \subset \left[ -\frac{3}{2}, -\frac{1}{2} \right] \cup \left[ \frac{1}{2}, \frac{3}{2} \right],
\]
(19.4.3)
Moreover, 
\[ \chi^\pm(S^D) = \chi^\pm(H), \] 
and the operators \( \chi^\pm(\hat{S}_m) \) have the representations \( \text{(19.4.2)} \).

**Proof.** Since the subspaces \( \mathcal{H}^\pm \) are invariant under the action of \( S^D \), our task is to show that the spectrum of \( S^D|_{\mathcal{H}^\pm} \) is positive and negative, respectively. This statement would certainly be true if we replaced \( S^D \) by \( S_m \), because the operator \( S_m \) has the eigenvalues \( \pm 1 \) with \( \mathcal{H}^\pm \) as the corresponding eigenspaces. Estimating the representation in Proposition \( 17.2.5 \) with the Schwarz inequality, we obtain
\[ \left| \langle \psi | S^D \phi \rangle - \langle \psi | S_m \phi \rangle \right| \leq \left( c + \frac{c^2}{2} \right) \| \psi \| \| \phi \| \quad \text{with} \quad c := \int_{-\infty}^{\infty} |B(\tau)|_{C^0} d\tau. \]

Using the assumption \( \text{(19.1.3)} \), we conclude that
\[ \left| \langle \psi | S^D \phi \rangle - \langle \psi | S_m \phi \rangle \right| < \frac{1}{2} \| \psi \| \| \phi \| \quad \text{for all} \quad \psi, \phi \in \mathcal{H}_m. \]

Standard estimates on the continuity of the spectrum (see for example \([103], \text{§IV.3}\)) yield that the spectrum of \( S^D \) differs by that of the operator \( S_m \) at most by \( 1/2 \). This gives \( \text{(19.4.3)} \) and \( \text{(19.4.4)} \).

In order to prove the representation \( \text{(19.4.2)} \), we take the resolvent identity
\[ (\hat{S}_m - \lambda)^{-1} = (S^D - \lambda)^{-1} - (\hat{S}_m - \lambda)^{-1} \Delta \hat{S} (S^D - \lambda)^{-1}, \]
form the contour integral and apply \( \text{(19.4.3)} \). This gives the result. \( \square \)

The next lemma relates the smoothness of an integral kernel to the boundedness of the product of the operator with powers of the vacuum Hamiltonian.

**Lemma 19.4.3.** Let \( A \in L(\mathcal{H}_m) \) be an operator which maps smooth functions to smooth functions and has the property that for all \( p, q \in \mathbb{N} \), the operator product
\[ H^q A H^p : C^\infty(N, SM) \to C^\infty(N, SM) \] 
extends to a bounded linear operator on \( \mathcal{H}_m \). Then, considering \( A \) as an operator on \( \mathcal{H}_m \), this operator can be represented as an integral operator with a smooth integral kernel, i.e.
\[ (A\psi)(x) = \int_N \mathcal{A}(x, (t_0, \vec{y})) \gamma^0 \psi(t_0, \vec{y}) d^3y \quad \text{with} \quad \mathcal{A} \in C^\infty(M \times M). \]

**Proof.** Since in momentum space, the square of the Hamiltonian takes the form
\[ H(\vec{k})^2 = \left( \gamma^0 (\vec{\gamma} \vec{k} + m) \right)^2 = \left( -\vec{\gamma} \vec{k} + m \right) \left( \vec{\gamma} \vec{k} + m \right) = |\vec{k}|^2 + m^2, \]
the wave function \( \hat{\psi} \) defined by
\[ \hat{\psi}(\vec{k}) := \frac{1}{|\vec{k}|^2 + m^2} e^{i\vec{k} \vec{x}_0} \Xi \]
for a constant spinor \( \Xi \) and \( \vec{x}_0 \in \mathbb{R}^3 \), satisfies the equation
\[ H^2 \psi(\vec{x}) = \delta^3(\vec{x} - \vec{x}_0) \Xi. \]
Moreover, one verifies immediately that \( \psi \in \mathcal{H}_{t_0} \) is square-integrable. Using the last equation together with \( \text{(19.4.5)} \), we conclude that
\[ H^q A(\delta^3(\vec{x} - \vec{x}_0) \Xi) = H^q AH^2 \psi \in \mathcal{H}_{t_0}. \]
Since \( q \) is arbitrary, it follows that \( A \) has an integral representation in the spatial variables,

\[
(A\phi)(\mathbf{x}) = \int_{\mathcal{N}} A(\mathbf{x}, \mathbf{y})\gamma^0 \phi(\mathbf{y})\,d^3y \quad \text{with} \quad A \in C^\infty(\mathcal{N} \times \mathcal{N}) .
\]

We now extend this integral kernel to \( \mathcal{M} \times \mathcal{M} \) by solving the Cauchy problem in the variables \( x \) and \( y \). This preserves smoothness by the global existence and regularity results for linear hyperbolic equations, giving the result.

**Lemma 19.4.4.** Under the assumptions of Theorem [19.7.1] for all \( p \in \mathbb{N} \) the iterated commutator

\[
S^{(p)} := \underbrace{[H, [H, \ldots, [H, S_m] \ldots]]}_{p \text{ factors}}
\]

is a bounded operator on \( \mathcal{H}_m \).

**Proof.** In the vacuum, the Hamiltonian clearly commutes with the time evolution operator,

\[
[H, U_m^{t,t'}] = 0 . \tag{19.4.6}
\]

In order to derive a corresponding commutator relation in the presence of the external potential, one must take into account that \( \tilde{H} \) is time-dependent. For ease in notation, we do not write out this dependence, but instead understand that the Hamiltonian is to be evaluated at the correct time, i.e.

\[
\tilde{U}_m^{t,t'} \tilde{H} \equiv \tilde{U}_m^{t,t'} \tilde{H}(t') \quad \text{and} \quad \tilde{H} \tilde{U}_m^{t,t'} \equiv \tilde{H}(t) \tilde{U}_m^{t,t'} .
\]

Then

\[
(i\partial_t - \tilde{H})(\tilde{H} \tilde{U}_m^{t,t'} - \tilde{U}_m^{t,t'} \tilde{H}) = i\dot{\tilde{H}} \tilde{U}_m^{t,t'} \quad \text{and} \quad \tilde{H} \tilde{U}_m^{t,t'} - \tilde{U}_m^{t,t'} \tilde{H} \big|_{t=t'} = 0
\]

(here and in what follows the dot denotes the partial derivative with respect to \( t \)). Solving the corresponding Cauchy problem gives

\[
[H, \tilde{U}_m^{t,t'}] = \int_\nu^t \tilde{U}_m^{t,\tau} \dot{\tilde{U}}_m^{\tau,t'} d\tau . \tag{19.4.7}
\]

In order to compute the commutator of \( H \) with the operator products in [17.2.11] and [17.2.12], we first differentiate the expression \( U_m^{t',t} \nabla \tilde{U}_m^{t,t'} \) with respect to \( t \),

\[
i\partial_t (U_m^{t',t} \nabla \tilde{U}_m^{t,t'}) = iU_m^{t',t} \dot{\nabla} \tilde{U}_m^{t,t'} + U_m^{t',t} \nabla \dot{\tilde{U}}_m^{t,t'} - U_m^{t',t} \nabla \dot{U}_m^{t,t'} . \tag{19.4.8}
\]

Moreover, using the commutation relations (19.4.6) and (19.4.7), we obtain

\[
H \left( U_m^{t',t} \nabla \tilde{U}_m^{t,t'} \right) = U_m^{t',t} H \nabla \tilde{U}_m^{t,t'} - U_m^{t',t} \nabla \dot{\tilde{U}}_m^{t,t'} + U_m^{t',t} \nabla \dot{U}_m^{t,t'} + \int_\nu^t U_m^{t',t} \nabla \tilde{U}_m^{t,\tau} \dot{\tilde{U}}_m^{\tau,t'} d\tau ,
\]

where in the last step we applied (19.4.8). It follows that

\[
[H, U_m^{t',t} \nabla \tilde{U}_m^{t,t'}] = H \left( U_m^{t',t} \nabla \tilde{U}_m^{t,t'} \right) = H (U_m^{t',t} \nabla \tilde{U}_m^{t,t'}) + \int_\nu^t U_m^{t',t} \nabla \tilde{U}_m^{t,\tau} \dot{\tilde{U}}_m^{\tau,t'} d\tau .
\]
Proceeding in this way, one can calculate the commutator of $H$ with all the terms in (17.2.11) and (17.2.12). We write the result symbolically as

$$[H, \tilde{S}_m] = S^{(1)},$$

where $S^{(1)}$ is a bounded operator. Higher commutators can be computed inductively, giving the result. □

We point out that this lemma only makes a statement on the iterative commutators. Expressions like $[H^p, \tilde{S}_m]$ or $H^q \tilde{S}_m H^p$ will not be bounded operators in general. However, the next lemma shows that the operator $\Delta \tilde{S}$ has the remarkable property that multiplying by powers of $H$ from the left and/or right again gives a bounded operator.

**Lemma 19.4.5.** Under the assumptions of Theorem 19.1.1, for all $p, q \in \mathbb{N} \cup \{0\}$ the product $H^q \Delta \tilde{S} H^p$ is a bounded operator on $\mathcal{H}_m$.

**Proof.** We only consider the products $H^q \tilde{S}_m$ because the operator $\tilde{S}_m^+$ can be treated similarly. Multiplying (19.4.7) from the left and right by the resolvent of $H$, we obtain

$$\left[ (H - \mu)^{-1}, \tilde{S}_m \right] = -(H - \mu)^{-1} S^{(1)} (H - \mu)^{-1}.$$ 

Writing the result of Lemma 19.4.4 as

$$[H, S^{(p)}] = S^{(p+1)} \quad \text{with} \quad S^{(p+1)} \in \mathcal{L}(\mathcal{H})$$

yields more generally the commutation relations

$$\left[ (H - \mu)^{-1}, S^{(p)} \right] = -(H - \mu)^{-1} S^{(p+1)} (H - \mu)^{-1} \quad \text{for} \ p \in \mathbb{N}. \quad (19.4.9)$$

Choosing a contour $\gamma$ which encloses the interval $(-\infty, -m]$ as shown in Figure 19.1, one finds

$$H \tilde{S}_m^- = -\frac{1}{2\pi i} \int_\gamma \mu (H - \mu)^{-1} \tilde{S}_m \chi^+(H) \, d\mu$$

$$= \mathcal{S} H \chi^-(H) \chi^+(H) + \frac{1}{2\pi i} \int_\gamma \mu (H - \mu)^{-1} S^{(1)} (H - \mu)^{-1} \chi^+(H) \, d\mu$$

$$= \frac{1}{2\pi i} \int_\gamma \mu (H - \mu)^{-1} S^{(1)} (H - \mu)^{-1} \chi^+(H) \, d\mu,$$

where in the last step we used that $\chi^-(H) \chi^+(H) = 0$. In order to show that this operator product is bounded, it is useful to employ the spectral theorem for $H$, which we write as

$$f(H) = \int_{\mathbb{R} \setminus [-m,m]} f(\lambda) \, dE_\lambda,$$ 

(19.4.10)
where \( dE_\lambda \) is the spectral measure of \( H \). This gives

\[
H S_+^- = \iint_{\mathbb{R} \times \mathbb{R}} \left( \frac{1}{2\pi i} \int \frac{\mu}{\lambda - \mu} \frac{1}{\lambda' - \mu} \chi^+(\lambda') \, dE_\lambda \right) S^{(1)} dE_{\lambda'} \, d\mu
= - \iint_{\mathbb{R} \times \mathbb{R}} \frac{\lambda}{\lambda - \lambda'} \chi^- (\lambda) \chi^+ (\lambda') \, dE_\lambda \, S^{(1)} \, dE_{\lambda'}.
\]

(19.4.11)

Note that the term \( \lambda - \lambda' \) is bounded away from zero. Thus the factor \( \lambda / (\lambda - \lambda') \) is bounded, showing that the operator \( HS_+^- \) is in \( L(\mathcal{H}_m) \).

This method can be iterated. To this end, we first rewrite the product with commutators,

\[
H^q S_+^- = \chi^- (H) \left( H^- \chi^- (H) \right)^p \tilde{S}_m \chi^+ (H)
= \chi^- (H) \left[ H^- , [H^- , \ldots , [H^- , \tilde{S}_m] \ldots ] \right] \chi^+ (H),
\]

where we used the abbreviation \( H^- := H \chi^- (H) \). Multiplying from the right by \( H^p \), we can commute factors \( H^+ := H \chi^+ (H) \) to the left to obtain

\[
H^q S_+^+ H^p = (-1)^p \chi^- (H) \left[ H^+ , \ldots , [H^+ , [H^- , \ldots , [H^- , \tilde{S}_m] \ldots ] \ldots ] \right] \chi^+ (H).
\]

Representing each factor \( H^\pm \) by a contour integral, one can compute the commutators inductively with the help (19.4.9). Applying the spectral theorem (19.4.10) to the left and right of the resulting factor \( S^{(p+q)} \) yields a constant times the expression

\[
\iint_{\mathbb{R} \times \mathbb{R}} \chi^- (\lambda) \chi^+ (\lambda') \, dE_\lambda \, S^{(p+q)} \, dE_{\lambda'}
\times \int_{\gamma_1} \frac{\mu_1 \, d\mu_1}{(\lambda - \mu_1)(\lambda' - \mu_1)} \cdots \int_{\gamma_{p+q}} \frac{\mu_{p+q} \, d\mu_{p+q}}{(\lambda - \mu_{p+q})(\lambda' - \mu_{p+q})}.
\]

Carrying out the contour integrals with residues, we obtain similar to (19.4.11) an expression of the form

\[
H^q S_+^+ H^p = \iint_{\mathbb{R} \times \mathbb{R}} f(\lambda, \lambda') \chi^- (\lambda) \chi^+ (\lambda') \, dE_\lambda \, S^{(p+q)} \, dE_{\lambda'}
\]

with a bounded function \( f \). This concludes the proof. \( \square \)

**Proof of Theorem 19.4.1.** It remains to be shown that the contour integral in (19.4.2) has a smooth integral kernel. To this end, we multiply the integrand from the left by \( H^q \) and from the right by \( H^p \) and commute the factors \( H \) iteratively to the inside. More precisely, we use the formula

\[
H^q (\tilde{S}_m - \lambda)^{-1} = \sum_{a=0}^q \left[ H, \ldots, [H, (\tilde{S}_m - \lambda)^{-1}] \ldots \right] H^{q-a}
\]
(note that the sum is telescopic; here we use the convention that the summand for \(a = 0\) is simply \((\tilde{s}_m - \lambda)^{-1}Hq\)). Hence

\[
H^q(\tilde{s}_m - \lambda)^{-1} \Delta \tilde{s}(S^D - \lambda)^{-1} H^p
\]

\[
= \sum_{a=0}^{q} \sum_{b=0}^{p} \underbrace{[H, \ldots, [H, (\tilde{s}_m - \lambda)^{-1}] \ldots]}_{a \text{ factors}} H^{q-a} \Delta \tilde{s} H^{p-b} \underbrace{[(S^D - \lambda)^{-1}, H], \ldots, H]}_{b \text{ factors}}.
\]

According to Lemma [19.4.3], the intermediate product \(H^{q-a} \Delta \tilde{s} H^{p-b}\) is a bounded operator. Moreover, the commutators can be computed inductively with the help of Lemma [19.4.4] and the formula

\[
[H, (\tilde{s}_m - \lambda)^{-1}] = - (\tilde{s}_m - \lambda^{-1}) [H, \tilde{s}_m] (\tilde{s}_m - \lambda^{-1})
\]

(and similarly for \(S^D\)). This gives operators which are all bounded for \(\lambda \in \partial B_{\frac{1}{2}}(\pm 1)\).

Since the integration contour is compact, the result follows. □

19.4.2. Proof of the Hadamard Form. Relying on the frequency mixing estimates of the previous section, we can now give the proof of Theorem [19.1.1]. Recall that the fermionic projector is given by \((see (15.4.1))

\[
P = -\chi^-(\tilde{s}_m) \tilde{k}_m,
\]

(19.4.12)

where we again used the short notation [19.4.4]. Here again the operator \(\chi^-(\tilde{s}_m)\) acts on the solution space \(H_m\) of the Dirac equation, which can be identified with the space \(H_{t_0}\) of square integrable wave functions at time \(t_0\) (see the beginning of Section [19.4.1]). For the following arguments, it is important to note that this identification can be made at any time \(t_0\).

In order to prove that the bi-distribution corresponding to \(P\) is of Hadamard form, we compare the fermionic projectors for three different Dirac operators and use the theorem on the propagation of singularities in \([127]\). More precisely, we consider the following three fermionic projectors:

1. The fermionic projector \(P^{\text{vac}}\) in the Minkowski vacuum.
2. The fermionic projector \(\tilde{P}\) in the presence of the external potential \(\tilde{\mathcal{B}}(x) := \eta(x^0) \mathcal{B}(x)\),

where \(\eta \geq 0\) is a smooth function with \(\eta \equiv 0\) and \(\eta \equiv 1\).
3. The fermionic projector \(P\) in the presence of the external potential \(\mathcal{B}(x)\).

The potential \(\tilde{\mathcal{B}}\) vanishes for negative times, whereas for times \(x^0 > 1\) it coincides with \(\mathcal{B}\). Thus it smoothly interpolates between the dynamics with and without external potential. The specific form of the potential \(\tilde{\mathcal{B}}\) in the transition region \(0 \leq x^0 \leq 1\) is of no relevance for our arguments.

In the Minkowski vacuum, the relation (19.4.12) gives the usual two-point function composed of all negative-frequency solutions of the Dirac equation. It is therefore obvious that the bi-distribution \(P^{\text{vac}}(x, y)\) is of Hadamard form.

We now compare \(P^{\text{vac}}\) with \(\tilde{P}\). To this end, we choose an arbitrary time \(t_0 < 0\). Then, applying the result of Theorem [19.4.1] to (19.4.12), we get

\[
P^{\text{vac}} = -\chi^-(H) k_m \quad \text{and} \quad \tilde{P} = -\chi^-(H) \tilde{k}_m + \text{(smooth)},
\]
where \( \tilde{k}_m \) is the causal fundamental solution in the presence of the potential \( \tilde{B} \). Since \( \tilde{B} \) vanishes in a neighborhood of the Cauchy surface at time \( t_0 \), we conclude that \( P^{\text{vac}} \) and \( \tilde{P} \) coincide in this neighborhood up to a smooth contribution. It follows that also \( \tilde{P}(x, y) \) is of Hadamard form in this neighborhood. Using the theorem on the propagation of singularities [127, Theorem 5.5], we conclude that \( \tilde{P}(x, y) \) is of Hadamard form for all \( x, y \in M \).

Next, we compare \( \tilde{P} \) with \( P \). Thus we choose an arbitrary time \( t_0 > 1 \). Using again the result of Theorem 19.4.1 in (19.4.12), we obtain

\[
\tilde{P} = -\chi^-(H) \tilde{k}_m + (\text{smooth}) \quad \text{and} \quad P = -\chi^-(H) \tilde{k}_m + (\text{smooth})
\]

(where the smooth contributions may of course be different). Since \( \tilde{B} \) and \( B \) coincide in a neighborhood of the Cauchy surface at time \( t_0 \), we infer that \( \tilde{P} \) and \( P \) coincide in this neighborhood up to a smooth contribution. As a consequence, \( P(x, y) \) is of Hadamard form in this neighborhood. Again applying [127, Theorem 5.5], it follows that \( P(x, y) \) is of Hadamard form for all \( x, y \in M \). This concludes the proof of Theorem 19.1.1.

Exercises

Exercise 19.1. This exercise explains the notion of the light-cone expansion in simple examples.

(a) What is the light-cone expansion of a smooth function on \( M \times M \)? In which sense is it trivial? In which sense is it non-unique?

(b) Show that \( A(x, y) = \log |y - x|^2 \) is a well-defined distribution on \( M \times M \). What is the order on the light cone? Write down a light-cone expansion.

(c) Now consider the distributional derivatives

\[
\left( \frac{\partial}{\partial x^m} \right)^p A(x, y) \quad \text{with} \quad p \in \mathbb{N}
\]

and \( A(x, y) \) as in part (b). What is the order on the light cone? Write down a light-cone expansion.

(d) Consider the function

\[
E(x, y) = \sin ((y - x)^2) \log |y - x|^2.
\]

Determine the order on the light cone and give a light-cone expansion.

(e) Consider the function

\[
E(x, y) = \begin{cases} 
  e^{-\frac{1}{(y-x)^2}} & \text{if } (y-x)^2 \geq 0 \\
  0 & \text{otherwise .}
\end{cases}
\]

Determine the order on the light cone and give a light-cone expansion.

(f) Show that the expression

\[
\lim_{\varepsilon \searrow 0} \frac{\log |y - x|^2}{(y - x)^2 + i\varepsilon}
\]

is a well-defined distribution on \( M \times M \). Derive its light-cone expansion.

Exercise 19.2. This exercise is devoted to computing the Fourier transform of the advanced Green’s function \([19.2.11]\) and deriving the series expansion \([19.2.13]\).
(a) As in Lemma 19.2, we set \( \xi = y - x \) and \( \xi = (t, \vec{\xi}) \) with \( t > 0 \). Moreover, we choose polar coordinates \( r = (|\vec{\xi}|, \vartheta, \varphi) \). Carry out the \( \omega \)-integration with residues and compute the angular integrals to obtain

\[
S_{m^2}^\vee(x, y) = \frac{i}{8\pi} \int_0^\infty \frac{p}{\omega(p)} \left( e^{-ipr} - e^{ipr} \right) \left( e^{i\omega(p)t} - e^{-i\omega(p)t} \right) dp,
\]

where \( p = |\vec{p}| \) and \( \omega(p) := \sqrt{|\vec{p}|^2 + m^2} \). Justify this integral as the Fourier transform of a distribution and show that

\[
S_{m^2}^\vee(x, y) = \lim_{\epsilon \to 0} \int_0^\infty e^{-\epsilon p} \frac{p}{\omega(p)} \left( e^{-ipr} - e^{ipr} \right) \left( e^{i\omega(p)t} - e^{-i\omega(p)t} \right) dp
\]

with convergence as a distribution.

(b) Verify (19.2.12) in the case \( m = 0 \) by setting \( \omega(p) = p \) and using (16.4.5).

(c) In order to analyze the behavior away from the light cone, it is most convenient to take the limit \( r \to 0 \) and use Lorentz invariance. Show that in this limit,

\[
S_{m^2}^\vee(x, y) = \lim_{\epsilon \to 0} \int_0^\infty e^{-\epsilon p} \frac{p^2}{\omega(p)} \left( e^{i\omega(p)t} - e^{-i\omega(p)t} \right) dp = \lim_{\epsilon \to 0} \int_0^\infty e^{-\epsilon p} \sqrt{\omega^2 - m^2} \left( e^{i\omega t} - e^{-i\omega t} \right) d\omega.
\]

Compute this integral using [92] formula (3.961.1) (similar as in the proof of Lemma 19.2). Use the relations between Bessel functions [114] (10.27.6), (10.27.11) to obtain (19.2.12) away from the light cone.

As an alternative method for computing the Fourier integral, one can begin from the integral representation for \( J_0 \) in [114] (10.9.12)], differentiate with respect to \( x \) and use [114] (10.6.3)].

(d) Combine the results of (b) and (c) to prove (19.2.12). Why is there no additional contribution at \( \xi = 0 \)?

(e) Use the series expansion [114] (10.2.2)] to derive (19.2.13).

(f) The series expansion (19.2.13) can also be derived without using Bessel functions. To this end, one expands (19.4.13) in powers of \( m^2 \) and computes the Fourier transform term by term. Verify explicitly that this procedure really gives (19.2.13).

**Exercise 19.3. (Understanding the light cone expansion)** This exercise aims to familiarize you with some of the particularities of the light cone expansion.

(a) Let \( A(x, y) := (x - y)^{2k_0} \) with \( k_0 \in \mathbb{Z} \). Which order(s) on the light cone is this? (Prove your answer.) Construct a light-cone expansion of \( A(x, y) \) and prove that it is one.

(b) Let \( B(x, y) := (x - y)^{2k_0} + (x - y)^{2k_1} \), where \( k_0, k_1 \in \mathbb{Z} \) and \( k_0 < k_1 \). Which order(s) on the light cone is this? (Prove your answer.) Construct a light-cone expansion of \( B(x, y) \) and prove that it is one.

(c) Let \( C(x, y) := (x - y)^{2k_0} f(x, y) + (x - y)^{2k_1} g(x, y) \), where \( f \) and \( g \) are smooth functions in \( x \) and \( y \) and \( k_0, k_1 \) as above. Construct a light-cone expansion of \( C(x, y) \) and prove that it is one.

(d) Let \( D(x, y) := \sin ((x - y)^2) (x - y)^2 \). Use your results from ii.) and iii.) to construct two different light-cone expansions of \( D(x, y) \). Why might this non-uniqueness not be a problem for the scope of the lecture?
(e) Finally, consider the function

\[ E(x, y) = \sin ((y - x)^2) + \begin{cases} 
  e^{-\frac{1}{(y-x)^2}} & \text{if } (y-x)^2 \geq 0 \\
  0 & \text{else}
\end{cases} , \]

Determine its order(s) on the light cone and a light cone expansion. (Prove your answer.)

*Hint:* For (d) and (e): Expand the sine function.
Part 4

Applications and Outlook
CHAPTER 20

A Few Explicit Examples of Causal Variational Principles

In this chapter we introduce a few examples of causal variational principles and analyze them in detail. These examples are too simple for being of of direct physical interest. Instead, they are chosen in order to illustrate the different mathematical structures introduced previously. It is a specific feature of these examples that a minimizing measure can be given in closed form, making it possible to analyze the system explicitly. Similar examples were first given in [74].

20.1. A One-Dimensional Gaussian

We let \( \mathcal{F} = \mathbb{R} \) and choose the Lagrangian as the Gaussian

\[
\mathcal{L}(x, y) = \frac{1}{\sqrt{\pi}} e^{-(x-y)^2}. \tag{20.1.1}
\]

Lemma 20.1.1. The Lebesgue measure

\[ d\rho = dx \]

is a minimizer of the causal action principle for the Lagrangian \( \mathcal{L}(x, y) \) in the class of variations of finite volume (see (6.3.2) and (6.3.1)). It is the unique minimizer within this class of variations.

Proof. Writing the difference of the actions as in (6.3.2), we can carry out the integrals over \( \rho \) using that the Gaussian is normalized,

\[
\int_{\mathcal{F}} \mathcal{L}(x, y) \, d\rho(y) = 1.
\]

We thus obtain

\[
S(\rho) - S(\hat{\rho}) = 2 \int_N d(\rho - \hat{\rho})(x) + \int_N d(\rho - \hat{\rho})(x) \int_N d(\rho - \hat{\rho})(y) \, \mathcal{L}(x, y)
\]

\[
= \int_N d(\rho - \hat{\rho})(x) \int_N d(\rho - \hat{\rho})(y) \, \mathcal{L}(x, y),
\]

where in the last line we used the volume constraint (6.3.1). In order to show that the last double integral is positive, we take the Fourier transform and use that the Fourier transform of a Gaussian is again a Gaussian. More precisely,

\[
\int_N e^{ipx} \, \mathcal{L}(x, y) \, dx = e^{-\frac{p^2}{4}} =: f(p).
\]

Moreover, the estimate

\[
\left| \int_N e^{ipx} \, d(\rho - \hat{\rho})(x) \right| \leq |\hat{\rho} - \rho|_{(\mathcal{F})} < \infty
\]

shows that the Fourier transform of the signed measure \( \hat{\rho} - \rho \) is a bounded function \( g \in L^\infty(\mathbb{R}) \). Approximating this function in \( L^2(\mathbb{R}) \), we can apply Plancherel’s theorem and
use the fact that convolution in position space corresponds to multiplication in momentum space. We thus obtain
\[
\hat{N}_d (\rho - \tilde{\rho}) (x) (\rho - \tilde{\rho}) (y) L(x, y) = \int_{-\infty}^{\infty} g(p) e^{-\frac{p^2}{4}} g(p) \geq 0 ,
\]
and the inequality is strict unless \(\tilde{\rho} = \rho\). This concludes the proof. \(\Box\)

The EL equations read
\[
\int_{\mathcal{F}} L(x, y) \, d\rho(y) = 1 \quad \text{for all} \quad x \in \mathbb{R} .
\]
We now specify the jet spaces. Since the Lagrangian is smooth, it is obvious that
\[
\mathcal{J}^{\text{diff}} = \mathcal{J} = C^\infty(\mathbb{R}) \oplus C^\infty(\mathbb{R})
\]
(where we identify a vector field \(a(x) \partial_x\) on \(\mathbb{R}\) with the function \(a(x)\)). The choice of \(\mathcal{J}^{\text{test}}\) is less obvious. For simplicity, we restrict attention to functions which are bounded together with all their derivatives, denoted by
\[
C^\infty_b := \{ f \in C^\infty(\mathbb{R}) \mid f^{(n)} \in L^\infty \text{ for all } n \in \mathbb{N}_0 \}.
\]
Now different choices are possible. Our first choice is to consider jets whose scalar components are compactly supported,
\[
\mathcal{J}^{\text{test}} = C^\infty_0(\mathbb{R}) \oplus C^\infty_b(\mathbb{R}) .
\]
The linearized field equations (8.1.6) reduce to the scalar equation
\[
\int_N (\nabla_1, v + \nabla_2, v) L(x, y) \, d\rho(y) - \nabla_2 1 = 0 \quad \text{for all} \quad x \in \mathbb{R} ,
\]
because if this equation holds, then the \(x\)-derivative of the left side is also zero. Using the EL equations (20.1.3), the linearized field simplify to
\[
\int_N \nabla_2, v L(x, y) \, d\rho(y) = 0 \quad \text{for all} \quad x \in \mathbb{R} .
\]
A specific class of solutions can be given explicitly. Indeed, choosing
\[
u = (a, A) \quad \text{with} \quad a \in C^\infty_0(\mathbb{R}) \text{ and } A(x) := \int_0^x a(t) \, dt \in C^\infty_b(\mathbb{R}) ,
\]
integration by parts yields
\[
\int_N \nabla_2, v L(x, y) \, d\rho(y) = \int_N (A'(y) + A(y) \partial_y) L(x, y) \, dy = 0 .
\]
These linearized solutions are referred to as inner solutions, as introduced in a more general context in [59]. Inner solutions can be regarded as infinitesimal generators of transformations of \(M\) which leave the measure \(\rho\) unchanged. Therefore, inner solutions do not change the causal fermion system, but merely describe symmetry transformations of the measure. With this in mind, inner solutions are not of interest by themselves. But they can be used in order to simplify the form of the jet spaces. For example, by adding suitable inner solutions one can arrange that the test jets have vanishing scalar components, as can be arranged by the transformation
\[
v = (b, v) \mapsto \tilde{v} := v + u \quad \text{with} \quad u = (-b, -B) ,
\]
20.2. A MINIMIZING MEASURE SUPPORTED ON A HYPERPLANE 263

where \( B \) is an indefinite integral of \( b \).

In our example, we can also consider the inner solutions alternatively in order to eliminate the vector component of the test jets. To this end, it is preferable to choose the space of test jets as

\[
\mathcal{J}^{\text{test}} = C_0^\infty(\mathbb{R}) \oplus C_0^\infty(\mathbb{R}) .
\]  

(20.1.6)

Now the vector component disappears under the transformation

\[
v = (b, v) \mapsto \tilde{v} := v + u \quad \text{with} \quad u = (-v', -v) \in \mathcal{J}^{\text{test}} .
\]

Therefore, it remains to consider the scalar components of jets. For technical simplicity, we restrict attention to compactly supported functions. Thus we choose the jet space \( \mathcal{J}^{\text{vary}} \) as

\[
\mathcal{J}^{\text{vary}} = C_0^\infty(\mathbb{R}) \oplus \{0\} .
\]

Then the Laplacian reduces to the integral operator with kernel \( L(x, y) \),

\[
(\Delta(b, 0))(x) = \int_{\mathcal{J}} L(x, y) b(y) \, dy .
\]

**20.2. A Minimizing Measure Supported on a Hyperplane**

In the previous example, the support of the minimizing measure was the whole space \( \mathcal{F} \). In most examples motivated from the physical applications, however, the minimizing measure will be supported on a low-dimensional subset of \( \mathcal{F} \) (see for instance the minimizers with singular support in \([79, 10]\)). We now give a simple example where the minimizing measure is supported on a hyperplane of \( \mathcal{F} \). We let \( \mathcal{F} = \mathbb{R}^2 \) and choose the Lagrangian as

\[
\mathcal{L}(x, y; x', y') = \frac{1}{\sqrt{\pi}} e^{-(x-x')^2} (1 + y^2) (1 + y'^2) ,
\]

(20.2.1)

where \((x, y), (x', y') \in \mathcal{F}\).

**Lemma 20.2.1.** The measure

\[
d\rho = dx \times \delta_y
\]

(20.2.2)

(where \( \delta_y \) is the Dirac measure) is the unique minimizer of the causal action principle for the Lagrangian (20.2.1) under variations of finite volume (see (6.3.2) and (6.3.1)).

Note that this measure is supported on the \( x \)-axis,

\[
M := \text{supp} \rho = \mathbb{R} \times \{0\} .
\]

**Proof of Lemma 20.2.1.** Let \( \tilde{\rho} \) be a regular Borel measure on \( \mathcal{F} \) satisfying (6.3.1). Then the difference of actions (6.3.2) is computed by

\[
\mathcal{S}(\tilde{\rho}) - \mathcal{S}(\rho) = \frac{2}{\sqrt{\pi}} \int_{\mathcal{F}} d(\tilde{\rho} - \rho)(x, y) \int_N \, dx' \, e^{-(x-x')^2} (1 + y^2)
\]

\[+ \frac{1}{\sqrt{\pi}} \int_{\mathcal{F}} d(\tilde{\rho} - \rho)(x, y) \int_{\mathcal{F}} d(\tilde{\rho} - \rho)(x', y') \, e^{-(x-x')^2} (1 + y^2) (1 + y'^2) .
\]

(20.2.3)

(20.2.4)

Using that the negative part of the measure \( \tilde{\rho} - \rho \) is supported on the \( x \)-axis, the first term (20.2.3) can be estimated by

\[
\geq \frac{2}{\sqrt{\pi}} \int_{\mathcal{F}} d(\tilde{\rho} - \rho)(x, y) \int_N \, dx' \, e^{-(x-x')^2} (1 + y^2)
\]

\[\geq \frac{2}{\sqrt{\pi}} \int_{\mathcal{F}} d(\tilde{\rho} - \rho)(x, y) \int_N \, dx' \, e^{-(x-x')^2} = \int_{\mathcal{F}} d(\tilde{\rho} - \rho)(x, y) = 0 ,
\]
where in the last step we used the volume constraint. The second term \((20.2.4)\), on the other hand, can be rewritten as

\[
\frac{1}{\sqrt{\pi}} \int_{\mathcal{F}} d\mu(x, y) \int_{\mathcal{F}} d\mu(x', y') e^{-(x-x')^2}
\]

with the signed measure \(\rho\) defined by

\[
d\mu(x, y) := (1 + y^2) \, d(\hat{\rho} - \rho)(x, y).
\]

Now we can proceed as in the proof of Lemma \(20.1.1\) and use that the Fourier transform of the integral kernel is strictly positive. For the uniqueness statement one uses that the inequality in \((*)\) is strict unless \(\hat{\rho}\) is supported on the \(x\)-axis. Then one can argue as in the proof of Lemma \(20.1.1\).

\[\square\]

For the minimizing measure \((20.2.2)\), the function \(\ell\) takes the form

\[
\ell(x, y) = \int_{\mathcal{F}} \mathcal{L}(x, y; x', y') \, d\rho(x', y') - 1 = y^2,
\]

showing that the EL equations \((7.1.2)\) are indeed satisfied. We now specify the jet spaces. Since the Lagrangian is smooth, it is obvious that

\[
\mathcal{J}^{\text{diff}} = \mathcal{J} = C^\infty(\mathbb{R}) \oplus C^\infty(\mathbb{R}, \mathbb{R}^2) \quad \text{(20.2.5)}
\]

where \(C^\infty(\mathbb{R}, \mathbb{R}^2)\) should be regarded as the space of two-dimensional vector fields along the \(x\)-axis. Similar as explained after \((20.1.6)\), we want to use the inner solutions for simplifying the vector components of the jets. To this end, in analogy to \((20.1.6)\) we choose

\[
\mathcal{J}^{\text{test}} = C^\infty_b(\mathbb{R}) \oplus C^\infty_b(\mathbb{R}, \mathbb{R}^2) \quad \text{(20.2.6)}
\]

The linearized field equations \((8.1.6)\) read

\[
\nabla_u \left( \int_{-\infty}^{\infty} (\nabla_{1, v} + \nabla_{2, v}) e^{-(x-x')^2} \left(1 + y^2\right) \left(1 + y'^2\right) \, dx' - \nabla_v \sqrt{\pi} \right) \bigg|_{y=y'=0} = 0. \quad \text{(20.2.7)}
\]

Now the inner solutions are generated by the vector fields tangential to the \(x\)-axis. More precisely, in analogy to \((20.1.3)\), we consider the jet

\[
v = (b, (B, 0)) \quad \text{with} \quad b \in C^\infty_0 \quad \text{and} \quad B(x) := \int_{-\infty}^{x} b(t) \, dt \in C^\infty_b(\mathbb{R}). \quad \text{(20.2.8)}
\]

Substituting this jet into \((20.2.7)\), the linearized field equations simplify to

\[
\nabla_u \left( (1 + y^2) \int_{-\infty}^{\infty} (\nabla_{1, v} + \nabla_{2, v}) e^{-(x-x')^2} \, dx' - \nabla_v \sqrt{\pi} \right) \bigg|_{y=y'=0} = 0.
\]

The second component of the vector field \(u\) yields a \(y\)-derivative, giving rise to a factor \(2y\), which vanishes at \(y = 0\). Therefore, it suffices to test with a vector field \(u\) which is tangential to the \(x\)-axis. Now we are back in the example of the one-dimensional Gaussian. Integrating by parts as in \((20.1.5)\) one sees that the jet \(v\) indeed satisfies the linearized field equations.

By suitably subtracting inner solutions, we can compensate the tangential components of the jets. This leads us to choose

\[
\mathcal{J}^{\text{vary}} = C^\infty_0(\mathbb{R}) \oplus \left( \{0\} \oplus C^\infty_0(\mathbb{R}) \right). \quad \text{(20.2.9)}
\]
Then the Laplacian simplifies as follows,

\[
\langle u, \Delta v \rangle (x) = \frac{1}{\sqrt{\pi}} \nabla u \bigg\rvert_{y=y'=0} \left( \int_{-\infty}^{\infty} \left( \nabla_{1,0} + \nabla_{2,0} \right) e^{-\langle x-x' \rangle^2} \left( 1 + y' \right)^2 \left( 1 + y'^2 \right) dx' - \nabla_y \sqrt{\pi} \right) \bigg\rvert_{y=y'=0}
\]

\[
= \frac{2}{\sqrt{\pi}} u(x) v(x) \int_{-\infty}^{\infty} e^{-(x-x')^2} dx' + \frac{1}{\sqrt{\pi}} a(x) \int_{-\infty}^{\infty} b(x') dx' 
\]

\[
+ a(x) \left( \frac{1}{\sqrt{\pi}} \int_{-\infty}^{\infty} b(x) e^{-(x-x')^2} dx' - b(x) \right) 
\]

\[
= 2 u(x) v(x) + \frac{1}{\sqrt{\pi}} a(x) \int_{-\infty}^{\infty} e^{-(x-x')^2} b(x') dx',
\]

where \( u = (a, (0, u)) \) and \( v = (b, (0, v)) \). Hence the inhomogeneous linearized field equations (20.2.9) with \( w = (e, w) \) give rise to separate equations for the scalar and vector components,

\[
\frac{1}{\sqrt{\pi}} \int_{-\infty}^{\infty} e^{-(x-x')^2} b(x') dx' = e(x), \quad v(x) = \frac{w(x)}{2}.
\]  

(20.2.10)

20.3. A Non-Homogeneous Minimizing Measure

In the previous examples, the minimizing measure \( \rho \) was translation invariant in the direction of the \( x \)-axis. We now give a general procedure for constructing examples of causal variational principles where the minimizing measure has no translational symmetry. In order to work in a concrete example, our starting point is again the one-dimensional Gaussian (20.1.1). But the method can be adapted to other kernels in a straightforward way. In view of these generalizations, we begin with the following abstract result.

**Lemma 20.3.1.** Let \( \mu \) be a measure on the \( m \)-dimensional manifold \( \mathcal{F} \) whose support is the whole manifold,

\[
\text{supp} \mu = \mathcal{F}.
\]

Moreover, let \( \mathcal{L}(x,y) \in L^1_{\text{loc}}(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+) \) be a symmetric, non-negative kernel on \( \mathcal{F} \times \mathcal{F} \). Next, let \( h \in C^0(\mathcal{F}, \mathbb{R}^+) \) be a strictly positive, continuous function on \( \mathcal{F} \). Assume that:

(i) \( \int_{\mathcal{F}} \mathcal{L}(x,y) h(y) \, d\mu(y) = 1 \) for all \( x \in \mathcal{F} \).

(ii) For all compactly supported bounded functions with zero mean,

\[
g \in L^\infty_0(\mathcal{F}, \mathbb{R}^+) \quad \text{and} \quad \int_{\mathcal{F}} g \, d\mu = 0,
\]

the following inequality holds,

\[
\int_{\mathcal{F}} d\mu(x) \int_{\mathcal{F}} d\mu(y) \mathcal{L}(x,y) g(x) g(y) \geq 0.
\]

(20.3.1)

Then the measure \( d\rho := h \, d\mu \) is a minimizer of the causal action principle under variations of finite volume (see (6.3.2) and (6.3.1)). If the inequality (20.3.1) is strict for any non-zero \( g \), then the minimizing measure is unique within the class of such variations.

**Proof.** We consider the variation

\[
\tilde{\rho}_\tau = \rho + \tau g \, d\mu = (h + \tau g) \, d\mu.
\]

(20.3.2)
Since $h$ is continuous and strictly positive and $g$ is continuous and compactly supported, the function $h + \tau g$ is non-negative for sufficiently small $|\tau|$. Furthermore, using that $g$ has mean zero, we conclude that (20.3.2) is an admissible variation of finite volume (6.3.1). Moreover, the difference of the actions (6.3.2) is well-defined and computed by

\[
S(\tilde{\mu}_\tau) - S(\rho) = 2\tau \int d\rho(x) g(x) \int N d\rho(y) h(y) L(x, y) + \tau^2 \int N d\rho(x) \int N d\rho(y) g(x) g(y)
\]

\[
\geq 2\tau \int N g(y) d\rho(y) = 0 ,
\]

where in the second step we used the above assumptions (i) and (ii). The last step follows from the fact that $g$ has mean zero.

We conclude that the measure $\rho$ is a minimizer under variations of the form (20.3.2).

In order to treat a general variation of finite volume (6.3.1), we approximate $\tilde{\rho}$ by a sequence of functions $g_n$ with the property that the measures $g_n \rho$ converge vaguely to $\tilde{\rho}$. For details see [30].

Our goal is to apply this lemma to kernels of the form

\[
L(x, y) = f(x) e^{-(x-y)^2} f(y)
\]

with a strictly positive function $f$, which for convenience we again choose as a Gaussian,

\[
f(x) = e^{\alpha x^2} \quad \text{with } \alpha \in \mathbb{R} .
\]

This kernel has the property (ii) because for all non-trivial $g \in L_0^\infty(\mathcal{F}, \mathbb{R}^+),$

\[
\int_\mathcal{F} d\mu(x) \int_\mathcal{F} d\mu(y) L(x, y) g(x) g(y) = \int_\mathcal{F} d\mu(x) \int_\mathcal{F} d\mu(y) e^{-(x-y)^2} (fg)(x) (fg)(y) > 0 ,
\]

where the last step is proved exactly as in the example of the Gaussian (see (20.1.2)). In order to arrange (i), for $h$ we make an ansatz again with a Gaussian,

\[
h(x) = c e^{\beta x^2}.
\]

Then

\[
\int_\mathcal{F} L(x, y) h(y) d\mu(y) = c \int_{-\infty}^{\infty} e^{\alpha x^2} e^{-(x-y)^2} e^{(\alpha+\beta)y^2} dy
\]

\[
= c \exp \left\{ \alpha x^2 - x^2 - \frac{x^2}{\alpha + \beta - 1} \right\} \int_{-\infty}^{\infty} \exp \left\{ (\alpha + \beta - 1) \left( y - \frac{x}{\alpha + \beta - 1} \right)^2 \right\} dy
\]

\[
= c \sqrt{\frac{\pi}{1 - \alpha - \beta}} \exp \left( \alpha x^2 - x^2 - \frac{x^2}{\alpha + \beta - 1} \right).
\]

In order to arrange that this function is constant one, we choose

\[
c = \sqrt{\frac{1 - \alpha - \beta}{\pi}} \quad \text{and} \quad \beta = -\frac{\alpha(2 - \alpha)}{1 - \alpha} .
\]

In order for the above Gaussian integral to converge, we need to ensure that $1 - \alpha - \beta > 0.$ In view of the formula

\[
1 - \alpha - \beta = \frac{1}{1 - \alpha} ,
\]

this can be arranged simply by choosing $\alpha < 1.$ Our finding is summarized as follows.
Proposition 20.3.2. For any $\alpha < 1$, we let $f$ and $h$ be the Gaussians (20.3.4) and (20.3.5) with $c$ and $\beta$ according to (20.3.6). Then the measure $d\rho = h\,dx$ is the unique minimizer of the causal action corresponding to the Lagrangian (20.3.3) within the class of variations of finite volume.

As a concrete example, we consider the well-known Mehler kernel (see for example [SS Section 1.5])

$$E(x, y) = \frac{1}{\sqrt{1-\mu^2}} \exp\left(-\frac{\mu^2(x^2 + y^2) - 2\mu xy}{(1-\mu^2)}\right)$$

with $\mu > 0$. Rescaling $x$ and $y$ according to

$$x, y \to \sqrt{\frac{1-\mu^2}{\mu}} x, y,$$

the Mehler kernel becomes

$$E(x, y) = \frac{1}{\sqrt{1-\mu^2}} \exp\left(-\frac{\mu(x^2 + y^2) - 2xy}{1-\mu^2}\right).$$

This kernel is of the desired form (20.3.3) if we choose

$$\alpha = 1 - \mu < 1, \quad \beta = \frac{\mu^2 - 1}{\mu}.$$

We finally remark that this non-homogeneous example can be used as the starting point for the construction of higher-dimensional examples with minimizing measures supported on lower-dimensional subsets, exactly as explained for the Gaussian in Section 20.2.

Exercises
CHAPTER 21

Basics on the Continuum Limit

In the continuum limit one analyzes the EL equations of the causal action principle for systems of Dirac seas in the presence of classical bosonic fields. As worked out in detail in [48, Chapters 3-5], this limiting case yields the interactions of the standard model and gravity on the level of second-quantized fermionic fields interacting with classical bosonic fields. In this chapter we explain schematically how the analysis of the continuum limit works and give an overview of the obtained results.

21.1. The Regularized Light-Cone Expansion

In Chapters 18 and 19 it was explained how to construct and analyze the unregularized kernel of the fermionic projector \( \tilde{P}(x,y) \) in Minkowski space in the presence of an external potential \( \mathcal{B} \). Our goal is to answer the question whether the causal fermion system corresponding to this kernel satisfies the EL equations corresponding to the causal action principle. Thus we would like to evaluate the EL equations as stated abstractly in Theorem 7.1.1. The basic procedure is to form the closed chain (see (5.6.3)) and to compute its eigenvalues \( \lambda_1^{xy}, \ldots \lambda_{2n}^{xy} \in \mathbb{C} \). This, in turn, makes it possible to compute the causal action and the constraints (see (5.5.1)-(5.5.5)). Considering first variations of \( P(x,y) \), one then obtains the EL equations.

The main obstacle before one can carry out this procedure is that, in order to obtain mathematically well-defined quantities, one needs to introduce an ultraviolet regularization. As explained in detail in Chapter 5, this regularization is not merely a technical procedure, but it corresponds to implementing a specific microscopic structure of spacetime. In the vacuum, the regularization was introduced with the help of a regularization operator \( \mathcal{R}_\varepsilon \) (see (5.4.1)). Different choices of regularization operators correspond to different microscopic structures of spacetime. Since the structure of our physical spacetime on the Planck scale is largely unknown, the strategy is to allow for a general class of regularization operators, making it possible to analyze later on how the results depend on the regularization (for more details on this so-called method of variable regularization see [48, §1.2.1]).

Having chosen a regularization operator \( \mathcal{R}_\varepsilon \) in the vacuum, we can follow the procedure in Chapter 5 and construct the corresponding causal fermion system, which also gives rise to the kernel of the fermionic projector (5.6.2). After suitable identifications, this kernel can be represented in Minkowski space by the regularized vacuum kernel \( P^\varepsilon(x,y) \), being the integral kernel of the operator

\[
P^\varepsilon := -\mathcal{R}_\varepsilon \pi_{\mathcal{M}} \mathcal{R}_\varepsilon^* k_m : C_0^\infty(\mathcal{M}, S\mathcal{M}) \rightarrow \mathcal{H}_m.
\]

The next question is whether \( \mathcal{R}_\varepsilon \) also determines the regularization of the kernel \( \tilde{P}(x,y) \) in the presence of a classical potential. This question is rather subtle, because it is not clear a-priori whether the regularized objects should still satisfy the Dirac equation.
In [48, Appendix F] and [39, Appendix D] a procedure is given for \textit{regularizing the light-cone expansion} (see [71] for related constructions in curved spacetime). We denote the regularized kernel in the presence of the bosonic potential $B$ by $\tilde{P}_\varepsilon(x, y)$.

### 21.2. The Formalism of the Continuum Limit

Having introduced the regularized kernel of the fermionic projector $\tilde{P}_\varepsilon(x, y)$, we can form the closed chain

$$ A_{xy}^\varepsilon := \tilde{P}_\varepsilon(x, y) \tilde{P}_\varepsilon(y, x), $$

and compute its eigenvalues and proceed by analyzing the EL equations. In the continuum limit, one focuses on the limiting case $\varepsilon \searrow 0$ when the ultraviolet regularization is removed. This limiting case is comparatively easy to analyze. This can be understood from the fact that, in the limit $\varepsilon \searrow 0$, the closed chain $A_{xy}^\varepsilon$ becomes singular on the light cone. Therefore, asymptotically for small $\varepsilon$, it suffices to take into account the contributions to $A_{xy}^\varepsilon$ on the light cone. These contributions, on the other hand, are captured precisely by the light-cone expansion of the unregularized kernel $\tilde{P}(x, y)$ (see Section 19.2 or the explicit formulas in Example 19.2.2). This is the basic reason why, in the continuum limit, the EL equations can be rewritten as field equations involving fermionic wave functions as well as derivatives of the bosonic potentials.

More specifically, the asymptotics $\varepsilon \searrow 0$ is captured by the \textit{formalism of the continuum limit}, which we now outline (for more details see [48, Section 2.4] or the derivation of the formalism in [39, Chapter 4]). In the first step, one regularizes the light-cone expansion symbolically by leaving all smooth contributions unchanged, whereas the singular factors $T^{(n)}$ of the factors $T^{(n)}$, we employ the replacement rule

$$ m^p T^{(n)} \rightarrow m^p T^{(n)}_{[p]}. \quad (21.2.1) $$

Thus for the formulas of Example 19.2.2, the factors $T^{(n)}$ get an additional index $[0]$. If the light-cone expansion involves powers of the rest mass, these powers are taken into account in the lower index. The resulting factors $T^{(n)}_{[p]}$ are smooth functions, making all the subsequent computations well-defined. The detailed form of these functions does not need to be specified, because we can get along with the following computation rules. In computations one may treat the $T^{(n)}_{[p]}$ like complex functions. However, one must be careful when tensor indices of factors $\xi$ are contracted with each other. Naively, this gives a factor $\xi^2$ which vanishes on the light cone and thus changes the singular behavior on the light cone. In order to describe this effect correctly, we first write every summand of the light cone expansion such that it involves at most one factor $\xi$ (this can always be arranged using the anti-commutation relations of the Dirac matrices). We now associate every factor $\xi$ to the corresponding factor $T^{(n)}_{[p]}$. In short calculations, this can be indicated by putting brackets around the two factors, whereas in the general situation we add corresponding indices to the factor $\xi$, giving rise to the replacement rule

$$ m^p \xi T^{(n)} \rightarrow m^p \xi^{(n)} T^{(n)}_{[p]}. \quad (21.2.2) $$

For example, we write the regularized fermionic projector of the vacuum as

$$ P^\varepsilon = i \frac{1}{2} \sum_{n=0}^{\infty} \frac{m^{2n}}{n!} \xi^{(-1+n)} T^{(-1+n)}_{[2n]} + \sum_{n=0}^{\infty} \frac{m^{2n+1}}{n!} T^{(n)}_{[2n+1]}.$$
The kernel $P(y, x)$ is obtained by taking the conjugate (see (21.2.6)). The conjugates of the factors $T_{[p]}^{(n)}$ and $\xi_{[p]}^{(n)}$ are the complex conjugates,

$$T_{[p]}^{(n)} := (T_{[p]}^{(n)})^* \quad \text{and} \quad \xi_{[p]}^{(n)} := (\xi_{[p]}^{(n)})^*. $$

One must carefully distinguish between these factors with and without complex conjugation. In particular, the factors $\xi_{[p]}^{(n)}$ need not be symmetric,

$$(\xi_{[p]}^{(n)})^* \neq \xi_{[p]}^{(n)} \quad \text{in general}. $$

When forming composite expressions, the tensor indices of the factors $\xi$ are contracted to other tensor indices. The factors $\xi$ which are contracted to other factors $\xi$ are called inner factors. The contractions of the inner factors are handled with the so-called contraction rules

$$ (\xi_{[p]}^{(n)})^j (\xi_{[p']}^{(n')})_j = \frac{1}{2} \left( z_{[p]}^{(n)} + z_{[p']}^{(n')} \right), \tag{21.2.3} $$

$$ (\xi_{[p]}^{(n)})^j (\xi_{[p']}^{(n')})_j = \frac{1}{2} \left( z_{[p]}^{(n)} + z_{[p']}^{(n')} \right), \tag{21.2.4} $$

$$ z_{[p]}^{(n)} T_{[p]}^{(n)} = -4 \left( n T_{[p]}^{(n+1)} + T_{[p]}^{(n+2)} \right), \tag{21.2.5} $$

which are to be complemented by the complex conjugates of these equations. Here the factors $z_{[p]}^{(n)}$ can be regarded simply as a book-keeping device to ensure the correct application of the rule (27.5.43). The factors $T_{[p]}^{(n)}$ have the same scaling behavior as the $T_{[p]}^{(n)}$, but their detailed form is somewhat different; we simply treat them as a new class of symbols. In cases where the lower index does not need to be specified we write $T^{(n)}_0$. After applying the contraction rules, all inner factors $\xi$ have disappeared. The remaining so-called outer factors $\xi$ need no special attention and are treated like smooth functions.

Next, to any factor $T^{(n)}_0$ we associate the degree $\deg T^{(n)}_0$ by

$$ \deg T^{(n)}_0 = 1 - n. $$

The degree is additive in products, whereas the degree of a quotient is defined as the difference of the degrees of numerator and denominator. The degree of an expression can be thought of as describing the order of its singularity on the light cone, in the sense that a larger degree corresponds to a stronger singularity (for example, the contraction rule (21.2.5) increments $n$ and thus decrements the degree, in agreement with the naive observation that the function $z = \xi^2$ vanishes on the light cone). Using formal Taylor series, we can expand in the degree. In all our applications, this will give rise to terms of the form

$$ \eta(x, y) \frac{T^{(a_1)}_0 \cdots T^{(a_m)}_0}{T^{(b_1)}_0 \cdots T^{(b_d)}_0} \frac{T^{(c_1)}_0 \cdots T^{(c_m)}_0}{T^{(d_1)}_0 \cdots T^{(d_d)}_0} \quad \text{with} \ \eta(x, y) \ \text{smooth}. \tag{21.2.6} $$

The quotient of the two monomials in this equation is referred to as a simple fraction.

A simple fraction can be given a quantitative meaning by considering one-dimensional integrals along curves which cross the light cone transversely away from the origin $\xi = 0$. This procedure is called weak evaluation on the light cone. For our purpose, it suffices to integrate over the time coordinate $t = \xi^0$ for fixed $\vec{\xi} \neq 0$. Moreover, using the symmetry under reflections $\xi \to -\xi$, it suffices to consider the upper light cone $t \approx |\vec{\xi}|$. The resulting
integrals diverge if the regularization is removed. The leading contribution for small \( \varepsilon \) can be written as

\[
\int_{|\vec{\xi}|-\varepsilon}^{|\vec{\xi}|+\varepsilon} dt \eta(t,\vec{\xi}) \frac{T_0^{(a_1)} \cdots T_0^{(a_\alpha)} T_0^{(b_1)} \cdots T_0^{(b_\beta)}}{T_0^{(c_1)} \cdots T_0^{(c_\gamma)} T_0^{(d_1)} \cdots T_0^{(d_\delta)}} \approx \eta(|\vec{\xi}|,\vec{\xi}) \frac{c_{\text{reg}}}{(i|\vec{\xi}|)_L} \log^r(\varepsilon|\vec{\xi}|) \varepsilon^{L-1},
\]

where \( L \) is the degree of the simple fraction and \( c_{\text{reg}} \), the so-called regularization parameter, is a real-valued function of the spatial direction \( \vec{\xi}/|\vec{\xi}| \) which also depends on the simple fraction and on the regularization details (the error of the approximation will be specified below). The integer \( r \) describes a possible logarithmic divergence. Apart from this logarithmic divergence, the scalings in both \( \xi \) and \( \varepsilon \) are described by the degree.

When analyzing a sum of expressions of the form (21.2.6), one must know if the corresponding regularization parameters are related to each other. In this respect, the integration-by-parts rules are important, which are described symbolically as follows. On the factors \( T_0^{(n)} \) we introduce a derivation \( \nabla \) by

\[
\nabla T_0^{(n)} = T_0^{(n-1)}.
\]

Extending this derivation with the Leibniz and quotient rules to simple fractions, the integration-by-parts rules state that

\[
\nabla \left( \frac{T_0^{(a_1)} \cdots T_0^{(a_\alpha)} T_0^{(b_1)} \cdots T_0^{(b_\beta)}}{T_0^{(c_1)} \cdots T_0^{(c_\gamma)} T_0^{(d_1)} \cdots T_0^{(d_\delta)}} \right) = 0.
\]

These rules give relations between simple fractions. Simple fractions which are not related to each other by the integration-by-parts rules are called basic fractions. As shown in [39] Appendix E], there are no further relations between the basic fractions. Thus the corresponding basic regularization parameters are linearly independent.

The above symbolic computation rules give a convenient procedure to evaluate composite expressions in the fermionic projector, referred to as the analysis in the continuum limit: After applying the contraction rules and expanding in the degree, the EL equations can be rewritten as equations involving a finite number of terms of the form (21.2.6). By applying the integration-by-parts rules, we can arrange that all simple fractions are basic fractions. We evaluate weakly on the light cone (21.2.7) and collect the terms according to their scaling in \( \xi \). Taking for every given scaling in \( \xi \) only the leading pole in \( \varepsilon \), we obtain equations which involve linear combinations of smooth functions and basic regularization parameters. We consider the basic regularization parameters as empirical parameters describing the unknown microscopic structure of spacetime. We thus end up with equations involving smooth functions and a finite number of free parameters.

We finally specify the error of the above expansions. By not regularizing the bosonic potentials and fermionic wave functions, we clearly disregard the higher orders in \( \varepsilon/\ell_{\text{macro}} \).

Furthermore, in (21.2.7) we must stay away from the origin, meaning that we neglect the higher orders in \( \varepsilon/|\vec{\xi}| \).

The higher order corrections in \( \varepsilon/|\vec{\xi}| \) depend on the fine structure of the regularization and thus seem unknown for principal reasons. Neglecting the terms in (21.2.9) and (21.2.10) also justifies the formal Taylor expansion in the degree. Clearly, leaving
out the terms \((21.2.10)\) is justified only if \(|\vec{\xi}| \gg \varepsilon\). Therefore, whenever using the above formalism, we must always ensure that \(|\vec{\xi}|\) is much larger than \(\varepsilon\).

We finally remark that, when working out the Einstein equations, one must go beyond error terms of the form \((21.2.9)\) and \((21.2.10)\). The reason is that the gravitational scales like \(\kappa \sim \delta^2 \approx \varepsilon^2\). In order not to lose the relevant terms in the error terms, one must take certain higher order contributions into account. This is done by using the so-called \(\iota\)-formalism. Here we do not enter the details but merely refer the interested reader to [48 §4.2.7].

21.3. Overview of Results of the Continuum Limit Analysis

We now outline the main results of the continuum limit analysis as obtained in [48 Chapters 3-5]. The main input is to specify the regularized kernel \(P^\varepsilon(x, y)\) of the vacuum. This involves:

- The fermion configuration in the vacuum, including the masses of the leptons and quarks. Moreover, it is built in that the neutrinos break the chiral symmetry.
- The vacuum kernel should satisfy the EL equations. This poses a few constraints on the regularization operator.

The output of the continuum limit are the following results:

- The structure of the interaction on the level of classical gauge theory.
- The gauge groups and their coupling to the fermions.
- The equations of linearized gravity.

In [48] the continuum limit is worked out in three steps for systems of increasing complexity. In Chapter 3, a system formed of a sum of three Dirac seas is considered. This configuration, referred to as a sector, can be thought of a simplified model describing the three generations of charged leptons \((e, \mu, \tau)\). In the continuum limit, we obtain the following results for the interaction as described by the causal action principle:

- The fermions interact via an axial gauge field.
- This axial gauge field is massive, with the mass determined by the masses of the fermions and the regularization.
- We find that the field equations for the axial gauge field arise in the continuum limit only if the number of generations equals three. For one or two generations, the resulting equations are overdetermined, whereas for more than three generations, the equations are under-determined (which means in particular that there is no well-posed Cauchy problem).
- We obtain nonlocal corrections to the classical field equations described by integral kernels which decay on the Compton scale. It seems that these nonlocal corrections capture certain features of the underlying quantum field theory. But the detailed connection has not been worked out.
- There is no gravitational field and no Higgs field.

In Chapter 4, a system formed as a direct sum of two sectors is considered. This system is referred to as a block. The first sectors looks as in Chapter 3. In the second sector, however, the chiral symmetry is broke. This system can be regarded as a model for the leptons, including the three generations of neutrinos. In the continuum limit, we obtain the following results for the interaction as described by the causal action principle:

- The fermions interact via a SU(2) gauge field, which couples only to one chirality (say, the left-handed fermions).
The corresponding gauge field is again massive.
Moreover, the fermions interact linearly via the linearized Einstein equations, where the coupling constant is related to the regularization length.

Finally, in Chapter 5 a realistic system involving leptons and quarks is considered. To this end, one considers a direct sum of eight sectors, one of which with broken chiral symmetry (the neutrino sector). These eight sectors form pairs, referred to as blocks. The block containing the neutrino sector describes the leptons, whereas the other blocks describe the quarks. Moreover, we obtain the following results:

- The fermions interact via the gauge group \( U(1) \times SU(2)_L \times SU(3) \). The corresponding gauge fields couple to the fermions as in the standard model. The \( SU(2) \)-field couples on to the left-handed component and is massive. The other gauge fields are massless.
- Moreover, the fermions interact linearly via the linearized Einstein equations.
- The EL equations corresponding to the causal action principle coincide with those of the standard model after spontaneous symmetry breaking, plus linearized gravity.
- There are scalar degrees of freedom which can be identified with the Higgs potential. However, the corresponding dynamical equations have not yet been worked out.

We conclude this section by discussing a few aspects of the derivation of these results. We begin with the system of one sector as considered in [48], Chapter 3. In this case, the kernel of the fermionic projector is the sum of \( g \in \mathbb{N} \) Dirac seas of masses \( m_1, \ldots, m_g \), i.e.

\[
P(x, y) = \sum_{\beta=1}^{g} P_{m_\alpha}(x, y),
\]

where again

\[
P_{m}(x, y) = \int \frac{d^4k}{(2\pi)^4} \left( \frac{k^2}{k^2} + m \right) \Theta(-k^0) e^{-ik(x-y)}.
\]

In order to perturb the system by gauge potentials, we first introduce the kernel of the auxiliary fermionic projector \( P_{aux}(x, y) \), which is obtained from \( P(x, y) \) if one replaces the sums by direct sums,

\[
P_{aux}(x, y) = \bigoplus_{\beta=1}^{g} P_{m_\alpha}(x, y)
\]

(this means that \( P_{aux}(x, y) \) is represented by a \((4g \times 4g)\)-matrix). The auxiliary kernel satisfies the Dirac equation

\[
\left( i \partial_x - \begin{pmatrix} m_1 & 0 & 0 \\ 0 & \ddots & 0 \\ 0 & 0 & m_g \end{pmatrix} \right) P_{aux}(x, y) = 0.
\]

Therefore, it can be perturbed as usual by inserting a potential \( B \) into the Dirac equation

\[
\left( i \partial_x + B(x) - \begin{pmatrix} m_1 & 0 & 0 \\ 0 & \ddots & 0 \\ 0 & 0 & m_g \end{pmatrix} \right) \tilde{P}_{aux}(x, y) = 0 \quad (21.3.2)
\]

(where \( B(x) \) is a matrix potential acting on \( \mathbb{C}^{4g} \)). The perturbed kernel \( \tilde{P} \) can be computed with the methods explained in Chapters [18] and [19]. Finally, we obtain the perturbed
kernel of the fermionic projector by summing over the generation indices in an operation referred to as the \textit{sectorial projection},

\begin{equation}
\tilde{P}(x, y) := \sum_{\alpha, \beta=1}^{g} \tilde{P}_{\alpha \beta}(x, y).
\end{equation}

After introducing an ultraviolet regularization, this kernel can be analyzed in the EL equations of the causal action principle, exactly as outlined in Section 21.2 above.

In order to gain the largest possible freedom in perturbing the system, the operator $B$ should be chosen as general as possible. For this reason, in [48, Chapter 4] a general class of potential was considered, including nonlocal potentials (i.e. integral operators).

A general conclusion of the analysis is that, in order to satisfy the EL equations, the potential $B$ must be local, i.e. a differential operator or a multiplication operator by a potential which may involve left- and right-handed potentials, but also bilinear, scalar or pseudo-scalar potentials,

\begin{equation}
B(x) = \chi_L A_R(x) + \chi_R A_L(x) + \sigma^{ij} A_{ij}(x) + \Phi(x) + i\gamma^5 \Xi(x) \tag{21.3.3}
\end{equation}

(where each of the potentials is a $g \times g$-matrix acting on the generations). Analyzing the continuum limit for such multiplication operators, one gets the above-mentioned results.

One feature which at first sight might be surprising is that, despite local gauge symmetry, we get \textit{massive gauge fields}. In order to understand how this comes about, we need to consider local gauge symmetries in connection with the chiral gauge potentials in (21.3.3). On the fundamental level of the causal fermion system, \textit{local gauge transformations} arise from the freedom in choosing bases of the spin spaces (see (5.8.1) and (5.8.2) in Section 5.8). In the present setting with four-component Dirac spinors, the local gauge transformations take the form

\begin{equation}
\psi(x) \rightarrow U(x) \psi(x) \quad \text{with} \quad U(x) \in U(2, 2), \tag{21.3.4}
\end{equation}

where $U(2, 2)$ is the group of unitary transformations of the spinors at the spacetime point $x$. The causal action principle is gauge invariant in the sense that the causal action is invariant under such gauge transformations. The group $U(2, 2)$ can be used to describe gravity as a gauge theory (for details see Section 4.2 or [36]). Restricting attention to flat spacetime, the main interest is that $U(2, 2)$ contains the gauge group $U(1)$ of electrodynamics as a subgroup. In other words, the causal action principle is gauge invariant under local phase transformations

\begin{equation}
\psi(x) \rightarrow e^{-i\Lambda(x)} \psi(x)
\end{equation}

with a real-valued function $\Lambda$.

The chiral potentials in (21.3.3) also give rise to generalized phase transformations. This can be seen for example by working out the leading term to the light-cone expansion (similar to (19.2.3) for the electromagnetic potential). One finds that the chiral gauge potentials give rise to phase transformations of the left- and right-handed components of the wave functions, i.e.

\begin{equation}
\psi(x) \rightarrow U(x) \psi(x) \quad \text{with} \quad U(x) := \chi_L e^{-i\Lambda_L(x)} + \chi_R e^{-i\Lambda_R(x)} \tag{21.3.5}
\end{equation}

(again with real-valued functions $\Lambda_L$ and $\Lambda_R$). The point is that this transformation is \textit{not unitary} with respect to the spin inner product, because the chirality flips when taking the adjoint

\begin{equation}
U^* = \chi_R e^{i\Lambda_L(x)} + \chi_L e^{-i\Lambda_R(x)} \quad \text{but} \quad U^{-1} = \chi_L e^{i\Lambda_L(x)} + \chi_R e^{-i\Lambda_R(x)}
\end{equation}
BASICS ON THE CONTINUUM LIMIT

(note that \( \chi^*_L = (1 - \gamma^5)^*/2 = (1 + \gamma^5)/2 = \chi_R \) because \( (\gamma^5)^* = -\gamma^5 \)). Therefore, as soon as \( A_L \neq A_R \), the generalized phase transformation \( U(x) \) in (21.3.4) is not a local transformation of the form (21.3.1). Consequently, the local transformation in (21.3.5) does not correspond to a symmetry of the causal action principle. Therefore, it is not a contradiction if these gauge potentials arise in the effective field equations as mass terms.

More specifically, the relative phases between left- and right-handed potentials do come up in the closed chain \( A_{xy} = P(x, y)P(y, x) \), as one sees immediately from the fact that, if \( P(x, y) \) is vectorial, then the chirality at the corresponding factor, i.e.

\[
\chi_L A_{xy} = \chi_L P(x, y) \chi_R P(y, x) \\
\quad \rightarrow \exp \left( -i(\Lambda_L(x) - \Lambda_R(x)) \right) \exp \left( i(\Lambda_L(y) - \Lambda_R(y)) \right) A_{xy}.
\]

(21.3.6)

Working out the corresponding contribution to the EL equations in the continuum limit, one finds that the axial current and a corresponding axial mass term come up in the effective field equations. The coupling constant and the bosonic mass depend on the detailed form of the regularization. But they can be computed for specific choices of the regularization, as is exemplified in [48, Chapter 3] by a hard cutoff in momentum space and the \( \varepsilon \)-regularization.

We now move on to the system of two sectors as analyzed in [48, Chapter 4]. The vacuum is described by a kernel of the fermionic projector \( P(x, y) \) being a direct sum of two summands, each of which is of the form (21.3.1), where we choose the number of generations as \( g = 3 \). Hence \( P(x, y) \) is a \( 8 \times 8 \)-matrix. Replacing the sums by direct sums, one obtains the corresponding auxiliary kernel \( P_{\text{aux}}(x, y) \) (being represented by a \( 24 \times 24 \)-matrix). In order to account for the observational fact that neutrinos are observed only as left-handed particles, one must break the chiral symmetry of one of the sectors (the neutrino sector). To this end, we assume that the regularization of the neutrino sectors is different from that of the other sector (the charged sector) by contributions which are not left-right invariant. The relevant length scale is denoted by \( \delta \gg \varepsilon \). This procedure is very general and seems the right thing to do, because the regularization effects on the scale \( \delta \) are also needed in order to obtain the correct form of the curvature term in the Einstein equations. In fact, the obtained linearized Einstein equations involve the coupling constant \( \kappa \sim \delta^2 \). As briefly mentioned at the end of Section 10.2, the derivation of the Einstein equations uses the \( \iota \)-formalism, which goes beyond the standard formalism of the continuum limit.

The system analyzed in [48, Chapter 5] is obtained similarly by adding direct summands to \( P(x, y) \) describing the three generations of quarks. We begin with eight sectors. These eight sectors form pairs, giving rise to four blocks. We conclude by outlining how this mechanism of spontaneous block formation comes about. To this end, we return to the gauge phases as already mentioned in (21.3.4) and (21.3.5). We already saw in (21.3.6) that, if the kernel of the fermionic projector is vectorial, then the relative phases (i.e. the difference of left- and right-handed gauge phases) show up in the eigenvalues of the closed chain. Such phase factors drop out of the causal Lagrangian because of the absolute values in (5.5.1). However, the situation becomes more involved if the kernel of the fermionic projector is not vectorial. Indeed, expanding the vacuum kernel in powers of the rest mass, the zero order contribution to \( P(x, y) \) is vectorial, whereas the first order contribution is scalar (more generally, one sees from (19.1.1) that the even orders in the mass are vectorial, whereas the odd orders are scalar). As a consequence, the absolute values of the eigenvalues \( |\lambda_i^{xy}| \) depend on a rather complicated way on the chiral gauge
phases. Moreover, considering a direct sum of Dirac seas, one must keep into account that the gauge phases in the above formulas must be replaced by generalized phases which can be described as ordered exponentials of the gauge potentials. Evaluating the causal Lagrangian (5.3.1), one gets conditions for chiral gauge conditions. In simple terms, these conditions can be stated by demanding that matrices formed of ordered exponentials of the gauge potentials must have degeneracies. Qualitatively speaking these degeneracies mean that the left-handed gauge potential must be the same in each block, and this condition even makes it possible to explain why such blocks form. A more detailed and more precise explanation can be found in [48, Chapter 5].

Exercises
CHAPTER 22

Connection to Quantum Field Theory

In this chapter we give an outlook on how to get a connection between the causal action principle and the dynamics of quantum fields. Since this direction of research is still work in progress, we do not enter any details but instead try to explain a few basic concepts and ideas. Our presentation is based on the recent research papers \[59, 60, 62\]. Moreover, we also mention a few constructions from the alternative approach in \[47\], which is more closely tied to the analysis of the continuum limit (as outlined in Chapter 21).

22.1. General Concepts

Before beginning, we point out that in all examples of causal fermion systems considered so far in this book (maybe with the exception of the causal variational principle on the sphere in Section 6.1), the measure \(\rho\) was the push-forward of the volume measure on Minkowski space or a Lorentzian manifold. Thus we first constructed a local correlation map (see (5.4.5))

\[
F^\varepsilon : \mathcal{M} \to \mathcal{F},
\]

and the introduced the measure \(\rho\) on \(\mathcal{F}\) by (see (5.4.6))

\[
\rho = (F^\varepsilon)_* \mu_{\mathcal{M}},
\]

(22.1.1)

where \(\mu_{\mathcal{M}}\) is the four-dimensional volume measure on \(\mathcal{M}\). In all these examples, the measure \(\rho\) had the special property that it was supported on a smooth four-dimensional subset of \(\mathcal{F}\) given by (for details see Exercise 8.1)

\[
M := \text{supp} \rho = F^\varepsilon(\mathcal{M}).
\]

Also when varying the measure in the derivation of the linearized field equations or in the study of interacting systems in the continuum limit, we always restricted attention to measures having this property (see \[8.1.3\] in Section 8.1 or Chapter 21). While this procedure seems a good starting point for the analysis of the causal action principle and gives good approximate solutions of the EL equations, we cannot expect that true minimizers are of this particular form.

With this in mind, our strategy is to allow for more general measures on \(\mathcal{F}\), and to analyze the causal action principle for these general measures. As we will see, this analysis gives rise to close connections to quantum field theory. We proceed step by step and begin by explaining a construction which explains why, going beyond push-forward measures of the form (22.1.1), makes it possible to further decrease the causal action. In other words, the following argument shows that minimizers of the causal action will not have the form of a push-forward measures (22.1.1), but will have more complicated structure. This argument is given in more detail in \[48, \S 1.5.3\]. Assume that we are given \(L\) measures \(\rho_1, \ldots, \rho_L\) be positive measures on \(\mathcal{F}\). Then the convex combination \(\tilde{\rho}\)
given by
\[ \tilde{\rho} := \frac{1}{L} \sum_{a=1}^{L} \rho_a \tag{22.1.2} \]
is again a positive measure on \( F \). Moreover, if the \( \rho_a \) satisfy the linear constraints (i.e. the volume constraint \( (5.5.3) \) and the trace constraint \( (5.5.4) \)), then these constraints are again respected by \( \tilde{\rho} \).

Next, we let \( \rho \) be a minimizing measure (describing for example the vacuum). Choosing unitary transformations \( U_1, \ldots, U_L \), we introduce the measures \( \rho_a \) in (22.1.2) as
\[ \rho_a(\Omega) := \rho(U^{-1}_a \Omega U_a) . \]
Thus, in words, the measures \( \rho_a \) are obtained from \( \rho \) by taking the unitary transformation by \( U_a \). Since the causal action and the constraints are unitarily invariant, each of the measures \( \rho_a \) is again minimizing. Let us compute the action of the convex combination (22.1.2). First, by (5.5.2),
\[ S(\tilde{\rho}) = \frac{1}{L^2} \sum_{a,b=1}^{L} \int_{F \times F} L(x,y) \, d\rho_a(x) \, d\rho_b(y) . \]
If \( a = b \), we obtain the action of the measure \( \rho_a \) which, due to unitary invariance, is equal to the action of \( \rho \). We thus obtain
\[ S(\tilde{\rho}) = \frac{S(\rho)}{L} + \frac{1}{L^2} \sum_{a \neq b} \int_{F \times F} L(x,y) \, d\rho_a(x) \, d\rho_b(y) . \tag{22.1.3} \]

Let us consider the contributions for \( a \neq b \) in more detail. In order to simplify the explanations, it is convenient to assume that the measures \( \rho_a \) have mutually disjoint supports (this can typically be arranged by a suitable choice of the unitary transformations \( U_a \)). Then the spacetime \( \tilde{M} := \text{supp} \, \tilde{\rho} \) can be decomposed into \( L \) “sub-spacetimes” \( M_a := \text{supp} \, \rho_a \),
\[ \tilde{M} = M_1 \cup \cdots \cup M_L \quad \text{and} \quad M_a \cap M_b = \emptyset \quad \text{if} \ a \neq b . \]
The Lagrangian of the last summand in (22.1.3) is computed from the fermionic projector \( P_{a,b}(x,y) \) where \( x \in M_a \) and \( y \in M_b \) are in different sub-spacetimes. Similar to (5.6.8), it can be expressed in terms of the physical wave functions by (for details see [48, Lemma 1.5.2])
\[ P_{a,b}(x,y) = - \sum_{i,j} |\psi_i^e(x)\rangle \langle (U_a U_b^*)_j^i | \psi_j^e(y) \rangle . \tag{22.1.4} \]
The point is that this fermionic projector involves the operator product \( U_a U_b^* \). By choosing the unitary operators \( U_a \) and \( U_b \) suitably, one can arrange that this operator product involve many phase factors. Moreover, one can arrange that, carrying out the sums in (22.1.4), these phases cancel each other due to destructive interference. In this way, the kernel \( P(x,y) \) can be made small if \( x \) and \( y \) lie in different sub-spacetimes. As a consequence, taking into account the factor \( 1/N \) in the first summand in (22.1.3), also the causal action of \( \tilde{\rho} \) becomes small. Clearly, this argument works only if the number \( L \) of sub-spacetimes is not too large, because otherwise it becomes more and more difficult to arrange destructive interference for all summands of the sum in (22.1.3) (estimating the optimal number \( L \) of subsystem is a difficult problem which we do not enter here). Also, we cannot expect that the simple ansatz (22.1.2) will already give a minimizer.
But at least, the above argument explains why it is too naive to think of a minimizing measure as being the push-forward measure of a volume measure under a smooth local correlation map. Instead, a minimizing measure could be composed of a large number of sub-spacetimes.

In the above construction, the fact that kernel $P_{a,b}(x,y)$ is very small for $a \neq b$ also means that the sub-spacetimes hardly interact with each other due to decoherence effects. Therefore, one can take the point of view that, in order to describe all physical phenomena, it suffices to restrict attention to one sub-spacetime. The appearance of many sub-spacetimes is an intriguing mathematical effect which may have interesting philosophical implications, but it is of no relevance as far as physical predictions are concerned. For this reason, here we shall not discuss these decoherent sub-spacetimes further. Also, we leave the question open whether they really occur for minimizing measures. Instead, we take the point of view that, in case our minimizing measure consists of several decoherent sub-spacetimes, we restrict it to one sub-spacetime and denote the resulting measure by $\rho$.

In order to understand the dynamics of a causal fermion system, it is more interesting to consider convex combinations of measures which are not decoherent. In order to explain the idea in a simple example, suppose we choose electromagnetic potentials $A_1, \ldots, A_L$ in Minkowski space (which do not need to satisfy Maxwell’s equations). Constructing the regularized kernels $P^\varepsilon_a(x,y)$ (as explained in Chapters 18 and 21), one gets corresponding causal fermion systems described by measures $\rho_a$. Abstractly, these measures can be written similar as explained in the context of the linearized field equations (see (8.1.7) in Section 8.1) as

$$\tilde{\rho} = \sum_{a=1}^L (F_a)_* (f_a \rho),$$

(22.1.5)

where $F_a$ is the corresponding local correlation map, and $f_a$ is a weight function. Since these measure are obtained from each other by small perturbations, it seems a good idea to depict the corresponding supports $M_a := \text{supp} \rho_a$ as being close together (see Figure 8.1 (b)). The convex combination of these measures (22.1.2) is referred to as a measure with fragmentation (see [51, Sections 1 and 5] or [53, Section 5]). In the limit when $N$ gets large, the fragmented measure $\tilde{\rho}$ goes over to a measure with enlarged support (see Figure 8.1 (e)). Integrating over this measure also involves an integration over the “internal degrees of freedom” corresponding to the directions which are transverse to $M := \text{supp} \rho$ (see the left of Figure 22.1). Integrating with respect to $\tilde{\rho}$ bears similarity to the path integral formulation of quantum field theory if one identifies the above “internal degrees of freedom” with field configurations.

For the mathematical description of the interacting measure $\tilde{\rho}$, working with fragmented measures does not seem to be the best method. One difficulty is that it is
a-priori not clear how large $N$ is to be chosen. Moreover, mechanisms where $N$ changes dynamically are difficult to implement, at least perturbatively. For these reasons, it seems preferable to describe $\tilde{\rho}$ with a different method, referred to as holographic mixing. At first sight, this method seems very different from fragmentation. However, as we will explain at the very end of this section, fragmentation also allows for the description of fragmented measures, at least if the construction is carried out in sufficiently large generality. The general idea and a few related constructions will be explained in the next section.

### 22.2. The Mechanism of Holographic Mixing

Let $(\mathcal{H}, \mathcal{F}, \rho)$ be a causal fermion system (for example describing the Minkowski vacuum). The wave evaluation operator $\Psi$ is a mapping which to every vector in $\mathcal{H}$ associates the corresponding physical wave function (for more details see for example [48 §1.1.4])

$$\Psi : \mathcal{H} \to C^0(M, SM), \quad u \mapsto \psi^u, \quad (22.2.1)$$

where the physical wave function $\psi^u$ is again given by (5.6.7). Evaluating at a fixed space-time point gives the mapping

$$\Psi(x) : \mathcal{H} \to S_x M, \quad u \mapsto \psi^u(x).$$

Working with the wave evaluation operator makes it possible to write the kernel of the fermionic projector (5.6.8) in the short form (for a detailed proof see [48, Lemma 1.1.3])

$$P(x, y) = -\Psi(x) \Psi(y)^*. \quad (22.2.2)$$

The general procedure of holographic mixing is to replace the wave evaluation operator by a linear combination of wave evaluation operators $\Psi_a$,

$$\tilde{\Psi} := \sum_{a=1}^L \Psi_a, \quad (22.2.3)$$

which in turn are all obtained by perturbing $\Psi$ (more details see below). Now we form the corresponding local correlation map,

$$\tilde{F} : M \to \mathcal{F}, \quad \tilde{F}(x) := -\tilde{\Psi}(x)^* \tilde{\Psi}(x),$$

and take the corresponding push-forward measure,

$$\tilde{\rho} := \tilde{F}_* \rho. \quad (22.2.4)$$

In this way, we have constructed a new measure $\rho$ which incorporates the perturbations described all the wave evaluation operators $\Psi_1, \ldots, \Psi_L$. However, in contrast to the convex combination of measures (22.1.5), the support of the measure (22.2.4) in general does not decompose into several fragments. In fact, if the mapping $\tilde{F}$ is continuous, injective and closed, the support of $\tilde{\rho}$ will again be homeomorphic to $M$. In other words, the topological structure of spacetime remains unchanged by the above procedure.

More concretely, the perturbed wave evaluation operators $\Psi_a$ can be obtained as follows. Suppose that the causal fermion system $(\mathcal{H}, \mathcal{F}, \rho)$ was constructed similar as in Section 5.4 from a system of Dirac wave functions satisfying for example the Dirac equation

$$(\mathcal{D} - m)\psi = 0.$$
Then one can perturb the system by considering the Dirac equation in the presence of classical potentials $\mathcal{B}_1, \ldots, \mathcal{B}_L$,

$$ (\mathcal{D} + \mathcal{B}_a - m) \psi_a = 0 . \quad (22.2.5) $$

The corresponding wave evaluation operators $\Psi_a$ are built up of all these Dirac solutions. In this way, the resulting wave evaluation operator (22.2.3) involves all the classical potentials $\mathcal{B}_a$. Qualitatively speaking, the resulting spacetime $\tilde{\mathcal{M}}$ can be thought of as being in a “superposition” of all these potentials. But this analogy does not carry over to a more technical level.

As already mentioned after (22.2.4), taking the push-forward with respect to a mapping $F$ does not change the topological structure of spacetime. Even more, if $F$ is smooth and varies only on macroscopic scales, then all microscopic structures of spacetime structures remain unchanged. This does not account for the picture of a measure $\tilde{\rho}$ which accounts for additional “internal degrees of freedom” as shown in Figure 8.1 (c) and the left of Figure 22.1. In order to allow the description of such measures, one needs to consider mapping $F$ which are not smooth but instead “fluctuate” on the microscopic scale (as is shown symbolically on the right of Figure 22.1). If we allow for such fluctuations even on the Planck scale, then the procedure (22.2.3) does allow for the description of all measures described previously with fragmentation (22.1.5). This considerations explains why the wave evaluation operators $\Psi_a$ should be constructed not only by introducing classical potentials (22.2.5), but in addition by introducing small-scale fluctuations. This leads us to the concept of **holographic mixing**, where in the first step we choose operators $A_a$ on $\mathcal{H}$ which add up to the identity,

$$ \sum_{a=1}^{N} A_a = \mathbb{1} , $$

and then decompose the local correlation operator by multiplying from the right with $A_a$,

$$ \Psi_a := \Psi A_a . $$

In the second step, the physical wave functions in $\Psi_a$ are perturbed by classical potentials $A_a$ by considering again the Dirac equation (22.2.5). In the last step, we again take the sum of the wave evaluation operators (22.2.3) and form the push-forward measure (22.2.4). This procedure is referred to as **holographic mixing**.

The resulting wave evaluation operator $\tilde{\Psi}$ involves both the operators $A_a$ and the potentials $\mathcal{B}_a$. Similar as explained in (22.1.4) in the context of fragmentation, the operators $A_a$ enter the kernel of the fermionic projector,

$$ P(x, y) = - \sum_{a, b=1}^{N} |\psi^{e_i}(x)\rangle \langle A_a A_b^*|_{j}^{e_i} \langle \psi^{e_i}(y) | . \quad (22.2.6) $$

In this way, one can build in phase factors into this kernel, possibly giving rise to destructive interference. In other words $\tilde{\Psi}$ is a sum of many, possibly decoherent components. The name “holographic mixing” is inspired by the similarity to a hologram in which several pictures are stored, each of which becomes visible only when looking at the hologram in the corresponding coherent light. For more details on the holographic mixing and related constructions we refer to [62].
22.3. A Distinguished Quantum State

The constructions outlined in the previous sections make it possible to construct general measures \( \tilde{\rho} \) which go beyond measures describing a classical spacetime with classical bosonic fields. Moreover, the EL equations for these measures can be understood as equations describing the dynamics in spacetime. The remaining question is how to interpret the resulting measure \( \tilde{\rho} \). Can it be understood in terms of an interaction via quantum fields? Or, in more physical terms, what does the measure \( \tilde{\rho} \) tell us about measurements performed in the corresponding spacetime? In order to address these questions in a systematic way, in \([60]\) a distinguished quantum state is constructed. It describes how the interacting measure \( \tilde{\rho} \) looks like if measurements are performed at a given time using the objects of a causal fermion describing the vacuum. This “measurement” can also be understood more generally as a “comparison” of the measures \( \tilde{\rho} \) and \( \rho \) at time \( t \). In technical terms, the quantum state, denoted by \( \omega^t \) is a positive linear functional on the algebra of fields \( A \) of the non-interacting spacetime,

\[
\omega : A \to \mathbb{C} \quad \text{with} \quad \omega(A^*A) \geq 0 \quad \text{for all} \quad A \in A.
\]

Here we use the language of algebraic quantum field theory (as introduced for example in the textbooks \([4, 19, 122]\)) which seems most suitable for describing quantum fields in the needed generality.

We now outline the construction of the quantum state as given in \([62]\). We are given two causal fermion systems \((\tilde{\mathcal{H}}, \tilde{\mathcal{F}}, \tilde{\rho})\) and \((\mathcal{H}, \mathcal{F}, \rho)\) describing the interacting system and the vacuum, respectively. Our goal is to “compare” these causal fermion systems at a given time. In order to specify the time, we choose sets \( \tilde{\Omega} \subset \tilde{\mathcal{M}} := \text{supp} \, \tilde{\rho} \) and \( \Omega \subset \mathcal{M} := \text{supp} \, \rho \) which can be thought of as the past of this time in the respective spacetimes. We want to relate the two causal fermions systems with the help of the nonlinear surface layer integral \([9.6.1]\) introduced in Section 9.6. However, we need to take into account that the causal fermion systems are defined on two different Hilbert spaces \( \tilde{\mathcal{H}} \) and \( \mathcal{H} \). Therefore, in order to make sense of the nonlinear surface layer integral, we need to identify the Hilbert spaces \( \mathcal{H} \) and \( \tilde{\mathcal{H}} \) by a unitary transformation denoted by \( V \),

\[
V : \mathcal{H} \to \tilde{\mathcal{H}} \quad \text{unitary}.
\]

Then the operators in \( \tilde{\mathcal{F}} \) can be identified with operators in \( \mathcal{F} \) by the unitary transformation,

\[
\mathcal{F} = V^{-1} \tilde{\mathcal{F}} V.
\]

An important point to keep in mind is that this identification is not canonical, but it leaves the freedom to transform the operator \( V \) according to

\[
V \to V \mathcal{U} \quad \text{with} \quad \mathcal{U} \in L(\mathcal{H}) \quad \text{unitary}.
\]

The freedom in choosing \( \mathcal{U} \) must be taken into account in the nonlinear surface layer integral, which now takes the form

\[
\gamma_{\tilde{\Omega}, \Omega}(\tilde{\rho}, \mathcal{U}\rho) = \int_{\tilde{\Omega}} d\tilde{\rho}(x) \int_{\mathcal{M} \setminus \Omega} d\rho(y) \mathcal{L}(x, \mathcal{U}y\mathcal{U}^{-1}) - \int_{\Omega} d\rho(x) \int_{\tilde{\mathcal{M}} \setminus \tilde{\Omega}} d\tilde{\rho}(y) \mathcal{L}(\mathcal{U}y\mathcal{U}^{-1}, y).
\]

The method for dealing with the freedom in choosing \( \mathcal{U} \) is to integrate over the unitary group. Moreover, it is preferable to consider the exponential of the nonlinear surface layer integral,
integral. This leads us to introduce the partition function $Z_{\tilde{\Omega}, \Omega}$ by

$$Z_{\tilde{\Omega}, \Omega}(\beta, \tilde{\rho}) = \int \exp \left( \beta \gamma_{\tilde{\Omega}, \Omega}(\tilde{\rho}, \mathcal{U}\rho) \right) d\mu_{G}(\mathcal{U}),$$

where $\mu_{G}$ is the normalized Haar measure on the unitary group (in order for this Haar measure to be well-defined, one needs to assume that the Hilbert space $\mathcal{H}$ is finite-dimensional, or else one must exhaust $\mathcal{H}$ by finite-dimensional subspaces).

Similar to the path integral formulation of quantum field theory, the quantum state is obtained by introducing insertions into the integrand of the partition function, i.e. symbolically,

$$\omega(\cdots) = \frac{1}{Z_{\Omega, \Omega}(\beta, \tilde{\rho})} \int \exp \left( \beta \gamma_{\tilde{\Omega}, \Omega}(\tilde{\rho}, \mathcal{U}\rho) \right) d\mu_{G}(\mathcal{U}).$$

These insertions have the structure of surface layer integrals involving linearized solutions in the vacuum spacetime. Likewise, the argument of the state on the left side is formed of operators which are parametrized by the same linearized solutions which enter the insertions on the right side. More precisely, they are operators of the field algebra $\mathcal{A}$, being defined as the $\ast$-algebra generated by the linearized solutions, subject to the canonical commutation and anti-commutation relations. The commutation relations involve the causal fundamental solution of the linearized solutions which can be constructed with energy methods as outlined in Section 14 (for details see [22]). Likewise, for the anti-commutation relations, we use the causal fundamental solutions of the dynamical wave equation mentioned at the end of Section 9.4 in (9.4.5) (for more details see [61]). The positivity property of the state is ensured by the specific form of the insertions. We refer the interested reader to [60].

**Exercises**
Part 5

Old Things
CHAPTER 23

Physical Background

23.1. Axiomatic Formulation of Quantum Mechanics

Quantum mechanics is the result of one of the largest revolutions in physics, which took place in the early 20th century, roughly simultaneously with the revolution in our understanding of space and time as explained in Section 1.2. The result of this revolution was captured in elegant and useful mathematical form in 1932 by John von Neumann [135]. In this section, we will introduce the formalism which he introduced using somewhat updated notation. Even though there have been many developments since this time, which we cannot explain in this book, von Neumann’s formalism is still largely in use and is still thought in many, if not most, undergraduate quantum mechanics courses.

According to current consensus among physicists, a physical theory, which describes a system (which could be as small as an atom or as large as the whole universe) has to consist at least of the following:

- A space of states of the system.
- A set of observables: To each measurement which can be performed on the system, there corresponds an observable which specifies which measurement outcomes can occur given a certain state, as well as the probability distributions of all outcomes.
- A law which describes the evolution of the states, i.e. a curve in the corresponding state space.

We now state the axioms of quantum mechanics in the conventional form. An explanation, as well as an example will follow. The mathematical objects are introduced in detail in Chapter 2.

1. The space of states consists of all one-dimensional subspaces of a complex separable Hilbert space \( \mathcal{H}, \langle \cdot | \cdot \rangle \).
2. The observables associated to measurements are the self-adjoint operators on \( \mathcal{H} \).
3. The possible measurement outcomes of a measurement which is described by an observable \( A \) are given by the spectrum of \( A \), denoted \( \sigma(A) \).
4. If the system is in state \( \psi \), the probability to find a measurement outcome \( \lambda_0 \in \sigma(A) \) with corresponding eigenvector \( \lambda \in \mathcal{H} \) is given by
   \[
   |\langle \lambda | \psi \rangle|^2 \in \mathbb{R}_0^+.
   \]
5. The unperturbed time evolution of a closed system is given by the Schrödinger equation, ...
6. After measurement...

Remarks:

\[1\] This is, of course, a huge idealisation. The application of a physical theory to calculate, i.e. predict, measurement results consists of many rules which are not available in axiomatized form, e.g. to determine just which observable one has to associate to a particular apparatus in the laboratory.
• states: only pure states
• superposition: source of many of the interesting properties of quantum theory
• observables if whole universe is described: effective (also relevant for CFS)
• Dirac notation

23.1.1. Interpretation of the axioms.
1. Axioms of QT
2. Little bit of meaning
3. Schrödinger equation for free particle (next section)
4. Special Relativity including Minkowski space
5. Dirac equation in Minkowski space
6. Spinors Bundle (on curved spacetime?)
CHAPTER 24

Mathematical Preliminaries

24.1. “Old” version: Distributions and Fourier Transformation

Consider the following process of “concentration” of a mass density of fixed total mass:

**Figure 24.1.** A sequence of functions concentrating in one point.

Viewed as functions, the values converge to zero everywhere except at the origin, where they diverge to infinity. Hence the limiting object cannot sensibly be described as a function. The mathematical theory of distributions, also called “singular functions”, provides a generalization of the usual concept of a function of \( n \) real variables which is capable to encapture also limiting objects of processes as the one illustrated above. This framework has proven to be extremely useful in the study of partial differential equations as well in the last century up to today.

24.1.1. Test function spaces. Let \( \Omega \subset \mathbb{R}^n \) be an open subset. All functions are assumed to be complex-valued in the following.

We set \( \mathcal{E}(\Omega) := C^\infty(\Omega) \) and \( \mathcal{D}(\Omega) := C^\infty_0(\Omega) := \{ f \in C^\infty(\Omega) \mid \text{supp } f \text{ is compact} \} \).

Recall here that the support of a function \( f \in C^\infty(\Omega) \) is defined as

\[
\text{supp } f := \overline{\{ x \in \Omega \mid f(x) \neq 0 \}},
\]

where the overline denotes the closure in \( \Omega \). We refer to elements of \( \mathcal{E}(\Omega) \) as test functions, and elements of \( \mathcal{D}(\Omega) \) as compactly supported test functions. Polynomials and functions defined by convergent power series are examples of test functions. Concerning compactly supported smooth functions, exercise 27.54 contains a common basic construction of such a function, a so-called smooth cutoff function. Note that both \( \mathcal{E}(\Omega) \) and \( \mathcal{D}(\Omega) \) are complex vector spaces with respect to pointwise addition and scalar multiplication.

The definition of distributions also requires a notion of convergence of sequence of test functions.

**Definition 24.1.1 (Convergence of test functions).**

1. We say a sequence \( (u_n)_{n \in \mathbb{N}} \) of test functions \( \mathcal{E} \)-converges to \( u \in \mathcal{E}(\Omega) \) if all partial derivatives of the sequence converge uniformly on all compact subsets of \( \Omega \) to the corresponding partial derivatives of \( u \).

2. We say a sequence \( (u_n)_{n \in \mathbb{N}} \) of compactly supported test functions \( \mathcal{D} \)-converges to \( u \in \mathcal{D}(\Omega) \) if there exists a compact subset \( K \subset \Omega \) with \( \text{supp } u, \text{supp } u_n \subset K \) for all \( n \in \mathbb{N} \) and if all partial derivatives of the sequence converge uniformly to the corresponding partial derivatives of \( u \).

Written out explicitly, one has \( u_n \to u \) in \( \mathcal{E}(\Omega) \) if

\[
\forall K \subset \Omega \text{ compact: } \forall \alpha \in \mathbb{N}^n : \sup_{x \in K} |\partial^\alpha u_n(x) - \partial^\alpha u(x)| \to 0.
\]

(24.1.2)
For a sequence of compactly supported test functions clearly $\mathcal{D}$-convergence implies $\mathcal{E}$-convergence. The converse is false, since the supports of a merely $\mathcal{E}$-convergent sequence may “expand to infinity” in the limit. A concrete, simple example of this phenomenon is contained in exercise 27.55.

Apart from test functions and compactly supported test functions, there is a third class of test functions which is important in particular in the context of the Fourier transform. It lies somewhat “in between” the previously introduced spaces in the sense of decay properties at infinity.

**Definition 24.1.2 (Schwartz functions).** A smooth function $u : \mathbb{R}^n \rightarrow \mathbb{C}$ is said to be **rapidly decaying** or a **Schwartz function** if

$$\forall \alpha, \beta \in \mathbb{N}^n : \sup_{x \in \mathbb{R}^n} |x^\alpha \partial^\beta u(x)| < \infty. \quad (24.1.3)$$

The vector space of all Schwartz functions is denoted by $\mathcal{S}(\mathbb{R}^n)$.

The notion “rapidly decaying” is perhaps best understood if one rewrites (24.1.3) as follows: For any $\alpha, \beta \in \mathbb{N}^n$ there exists $C > 0$ such that

$$\forall x \in \mathbb{R}^n \setminus \{0\} : |\partial^\beta u(x)| \leq \frac{C}{|x|^\alpha}. \quad (24.1.3')$$

This shows that any derivative of $u$ is required to decay fast than any polynomial at infinity. Since this is a condition at infinity, we restrict to functions defined on all of $\mathbb{R}^n$. In exercise 27.56 you may explore some further reformulations of this condition.

Note that any compactly supported test function is trivially also a Schwartz function. The standard example of a noncompactly supported Schwartz function are Gaussians.

**Example 24.1.3 (Gaussians are rapidly decaying).** The Gaussian $u : \mathbb{R}^n \rightarrow \mathbb{C}$ defined by $u(x) = e^{-|x|^2}$ lies in $\mathcal{S}(\mathbb{R}^n)$. For any $\alpha \in \mathbb{N}^n$ one has $\partial^\alpha u = p_\alpha u$ for some polynomial $p_\alpha$, as can be shown by induction on $|\alpha|$. Using this, (24.1.3) is easily established by l’Hopitals rule for $n = 1$, the case $n > 1$ can be traced back to the one-dimensional case.

More generally, for any polynomial $p$ and any positive definite, symmetric matrix $A \in \text{Mat}_n(\mathbb{R})$, the function $v : \mathbb{R}^n \rightarrow \mathbb{C}$ defined by $v(x) = p(x)e^{-x^tAx}$ lies in $\mathcal{S}(\mathbb{R}^n)$. The details are left to the reader as exercise 27.57.

As before we now specify a notion of convergence for Schwartz functions which will be needed for the definition of the corresponding space of distributions.

**Definition 24.1.4.** We say that a sequence $(u_n)_n$ of Schwartz functions $\mathcal{S}$-**converges** to some $u \in \mathcal{S}(\mathbb{R}^n)$ if for all $\alpha, \beta \in \mathbb{N}^n$ the functions $x^\alpha \partial^\beta u_n$ converge uniformly to $x^\alpha \partial^\beta u$.

Note the following obvious inclusions of the three test function spaces:

$$\mathcal{D}(\mathbb{R}^n) \subset \mathcal{S}(\mathbb{R}^n) \subset \mathcal{E}(\mathbb{R}^n). \quad (24.1.4)$$

Note further that a $\mathcal{D}$-convergent sequence also $\mathcal{S}$-converges to the same limit, and an $\mathcal{S}$-convergent sequence also $\mathcal{E}$-converges to the same limit. Therefore the inclusions (24.1.4) are both sequentially continuous. Clearly for any $\Omega \subset \mathbb{R}^n$ open the inclusion $\mathcal{D}(\Omega) \subset \mathcal{E}(\Omega)$ still holds and is sequentially continuous.

Since Schwartz functions decay rapidly at infinity, they are integrable on all of $\mathbb{R}^n$. They even include continuously into all $L^p$-spaces, as proven by the following Lemma. Its proof contains a simple but very useful argument which is worth remembering.
Lemma 24.1.5. If $u \in \mathcal{S}(\mathbb{R}^n)$, then $u \in L^p(\mathbb{R}^n)$ for all $p \in [1, \infty]$. Moreover, the inclusion $\mathcal{S}(\mathbb{R}^n) \to L^p(\mathbb{R}^n)$ is sequentially continuous.

Proof. We focus on the cases $p \in [1, \infty)$ and leave the case $p = \infty$ as exercise. Let $u \in \mathcal{S}(\mathbb{R}^n)$ be given. For $p \in [1, \infty)$ given, choose $m \in \mathbb{N}$ with $pm > n$. Then the function $x \mapsto (1 + |x|^2)^{-\frac{pm}{2}}$ is integrable on $\mathbb{R}^n$ (use polar coordinates). Therefore

$$
\|u\|_{L^p(\mathbb{R}^n)}^p = \int_{\mathbb{R}^n} |u(x)|^p \, dx = \int_{\mathbb{R}^n} |u(x)|^p(1 + |x|^2)^{\frac{pm}{2}}(1 + |x|^2)^{-\frac{pm}{2}} \, dx = \int_{\mathbb{R}^n} \left|1 + |x|^2\right|^{\frac{pm}{2}}|u(x)|^p (1 + |x|^2)^{-\frac{pm}{2}} \, dx \leq \sup_{x \in \mathbb{R}^n} \left|1 + |x|^2\right|^{\frac{pm}{2}}u(x) \int_{\mathbb{R}^n} (1 + |x|^2)^{-\frac{pm}{2}} \, dx$$

$$< \infty. $$

Concerning sequential continuity, let $(u_n)_n$ be a sequence in $\mathcal{S}(\mathbb{R}^n)$ with $u_n \to u \in \mathcal{S}(\mathbb{R}^n)$. Applying the previous computation to $u_n - u$ shows that

$$
\|u_n - u\|_{L^p(\mathbb{R}^n)}^p \leq C \sup_{x \in \mathbb{R}^n} \left|1 + |x|^2\right|^{\frac{pm}{2}}(u_n - u)$$

for some $C > 0$. By definition 24.1.4 this tends to zero for $n \to \infty$ since $u_n \to u$ in $\mathcal{S}$. □

The special function

$$
\mathbb{R}^n \ni x \mapsto \langle x \rangle := (1 + |x|^2)^{\frac{1}{2}} \in \mathbb{R}
$$

(24.1.5)

used in the proof is called Japanese bracket. For large $|x|$, its value is comparable to $|x|$, whereas for small $|x|$ its value is comparable to 1. The latter has the advantage that one can take negative powers of it without having to worry about creating singularities.

Remark 24.1.6. It is possible to equip each test function space with a topology such that the corresponding notion of convergence coincides with the one introduced above. Since for most practical purposes this does not give any real benefits, we chose to not include it here. The interested reader is referred to [Friedlander].

24.1.2. Definition and examples of distributions. Now we are set up for the definition of a distribution. To each of the three test function spaces corresponds a space of distributions.

Definition 24.1.7 (Distributions). Let $\Omega \subset \mathbb{R}^n$ be an open subset.

1. A distribution on $\Omega$ is a sequentially continuous linear map $T : \mathcal{D}(\Omega) \to \mathbb{C}$. The vector space of all distributions on $\Omega$ is denoted by $\mathcal{D}'(\Omega)$.

2. A compactly supported distribution on $\Omega$ is a sequentially continuous linear map $T : \mathcal{E}(\Omega) \to \mathbb{C}$. The vector space of all distributions on $\Omega$ is denoted by $\mathcal{E}'(\Omega)$.

3. A tempered distribution is a sequentially continuous linear map $T : \mathcal{S}(\mathbb{R}^n) \to \mathbb{C}$. The vector space of all tempered distributions is denoted by $\mathcal{S}'(\mathbb{R}^n)$.

For a distribution $T$ and a test function $u$, we will often write $\langle T, u \rangle$ instead of $T(u)$.
MATHEMATICAL PRELIMINARIES

To be precise, the requirement of sequential continuity of a distribution $T$ means that whenever $u_n \to u$ holds in the corresponding space of test functions, then $(T, u_n) \to (T, u)$ holds in $\mathcal{C}$. Due to the sequentially continuous inclusion (24.1.4) of test function spaces one has the reverse inclusion of distribution spaces:

$$\mathcal{E}'(\mathbb{R}^n) \subset \mathcal{S}'(\mathbb{R}^n) \subset \mathcal{D}'(\mathbb{R}^n). \quad (24.1.6)$$

Again, for a domain $\Omega \subset \mathbb{R}^n$ one still has the inclusion $\mathcal{E}'(\Omega) \subset \mathcal{D}'(\Omega)$.

We now go over some examples of distributions, further examples can be found in the exercises.

Example 24.1.8 (Dirac $\delta$-distribution). For $\Omega \subset \mathbb{R}^n$ open and $x_0 \in \Omega$, the (Dirac) $\delta$-distribution at $x_0$ is the distribution $\delta_{x_0} \in \mathcal{E}'(\Omega)$ defined by

$$\delta_{x_0}(u) := u(x_0) \quad (24.1.7)$$

for any $u \in \mathcal{E}(\Omega)$. Clearly $\delta_{x_0}$ is a linear function. Concerning sequential continuity, suppose that we have $u_n \to u$ in $\mathcal{E}(\Omega)$. This means in particular that $u_n \to u$ uniformly on any compact subset of $\Omega$ which implies pointwise convergence as well. Therefore

$$\delta_{x_0}(u_n) = u_n(x_0) \to u(x_0) = \delta_{x_0}(u).$$

Since test functions are differentiable and also the notions of convergence of test functions takes into account derivatives, one may also use derivatives of test functions for the construction of distributions. For instance, for any $\Omega \subset \mathbb{R}^n$ open, any $x_0 \in \Omega$, and any $\alpha \in \mathbb{N}^n$, the map

$$\mathcal{E}(\Omega) \ni u \mapsto (\partial^\alpha u)(x_0) \in \mathbb{C}$$

defines a distribution. This can be checked similarly as for the $\delta$-distribution.

Example 24.1.9 (Integration as a distribution). Consider integration over $\mathbb{R}$ as a map $I : \mathcal{D}(\mathbb{R}) \to \mathbb{C}$, $I(u) := \int_\mathbb{R} u(x) \, dx$. Since any $u \in \mathcal{D}(\mathbb{R})$ has compact support and is continuous, $I(u)$ is well-defined. Clearly $I$ is linear as well. Concerning sequential continuity, suppose that $u_n \to u$ in $\mathcal{D}(\mathbb{R})$. This means that there exists some $R > 0$ such that $\text{supp } u_n, \text{supp } u \subset [-R, R]$ and (in particular) $u_n \to u$ uniformly on $[-R, R]$. Hence

$$I(u_n) = \int_{-R}^R u_n(x) \, dx \to \int_{-R}^R u(x) \, dx = \int_{-R}^R u(x) \, dx = I(u),$$

correspondence of the integrals following for instance from uniform convergence.

The next example is of particular importance since it describes in which sense a function may be understood as a distribution as well, and hence why distributions indeed generalize the notion of functions.

Example 24.1.10 (Functions as distributions). A function $f : \Omega \to \mathbb{C}$ is said to be **locally integrable** if $\int_K |f(x)| \, dx < \infty$ for any compact subset $K \subset \Omega$. If $f$ is such a function, then for any $u \in \mathcal{D}(\Omega)$ the integral

$$T_f(u) := \int_\Omega f(x)u(x) \, dx \quad (24.1.8)$$

is well-defined. Clearly the map $T_f : \mathcal{D}(\Omega) \to \mathbb{C}$ is linear. It is also sequentially continuous: If $u_n \to u$ in $\mathcal{D}(\Omega)$, there exists $K \subset \Omega$ compact with $\text{supp } u_n, \text{supp } u \subset K$ for all $n \in \mathbb{N}$ and $u_n \to u$ uniformly. By uniform convergence, we have $|u_n| \leq C + |u|$ for some $C > 0$ and all $n \in \mathbb{N}$. Since all these functions are supported in the compact set $K$, we also have $|u_n f| \leq (C + |u|)|f| \in L^1(\Omega)$. It now follows from the dominated convergence theorem that $T_f(u_n) \to T_f(u)$ holds.
Distributions of the form \((24.1.8)\) are called regular distributions. By the so-called fundamental theorem in the calculus of variations (see [] for \(f\) and \(g\) locally integrable, \(T_f = T_g\) implies that \(f = g\) almost everywhere (with respect to Lebesgue measure). For continuous functions you are asked to prove this in exercise \(27.62\). Hence, if as usual in integration theory we identify functions which agree almost everywhere and denote by \(L^1_{\text{loc}}(\Omega)\) the space of the corresponding equivalence classes of all locally integrable functions on \(\Omega\) (all with respect to Lebesgue measure), then the map

\[
L^1_{\text{loc}}(\Omega) \ni f \mapsto T_f \in \mathcal{D}'(\Omega)
\]  

(24.1.9)
is injective. Therefore one can view \(L^1_{\text{loc}}(\Omega)\) as a subspace of \(\mathcal{D}'(\Omega)\). For this reason, we often simply write \(f\) instead of \(T_f\) if it is clear from the context that we want to understand \(f\) as a distribution. Note that in particular all continuous functions and all \(L^p\)-functions for \(p \in [1, \infty]\) are contained in \(L^1_{\text{loc}}\). Be aware that, just as \(L^p\)-functions, locally integrable functions may also be unbounded even on compact sets, i.e. “have poles”, as long as these poles are integrable, see the exercises \(27.64\) and \(27.65\).

The \(\delta\)-distribution is not a regular distribution. To see this in the case of \(\Omega = \mathbb{R}\), fix some \(u \in \mathcal{D}(\mathbb{R})\) with \(0 \leq f \leq 1\), \(f(0) = 1\), and \(\text{supp} u \subset [-1, 1]\). Set \(u_n(x) := u(nx)\), then we have \(\delta_0(u_n) = u_n(0) = u(0) = 1\) for all \(n \in \mathbb{N}\). However, for any locally integrable function \(f : \mathbb{R} \to \mathbb{C}\) we have \(u_n(x)f(x) \to 0\) for all \(x \neq 0\) and \(|u_n f| \leq 1_{[-1,1]} f \in L^1(\mathbb{R})\), so by dominated convergence it follows that

\[
T_f(u_n) = \int_{\mathbb{R}} u_n(x)f(x) \, dx \to 0.
\]

Hence \(T_f \neq \delta_0\) for all \(f \in L^1_{\text{loc}}(\mathbb{R})\).

The following exercise connects distributions with another possible generalization of the notion of a function, namely that of a measure. In a certain sense (which is made precise by the so-called Riesz representation theorem), these are all distributions which only use the values of a test function but not if its derivatives.

**Example 24.1.11** (Locally finite measures as distributions). Let \(\mu\) be a Borel measure on \(\Omega\), i.e. a measure on the Borel \(\sigma\)-algebra of \(\Omega\). Then \(\mu\) is called locally finite if every point \(x \in \Omega\) possesses an open neighborhood \(U_x \subset \Omega\) such that \(\mu(U_x) < \infty\). Using that any open cover of a compact set has a finite subcover, any locally finite Borel \(\mu\) measure defines a distribution \(T_\mu \in \mathcal{D}'(\Omega)\) by integration, i.e. by \(u \mapsto \int_{\Omega} u(x) \, d\mu(x)\). You are asked to fill in the details in exercise \(27.67\).

In contrast to the previous example, in exercise \(27.68\) you find a distribution which uses arbitrarily many derivatives (i.e., of arbitrary high order).

We finish this section with a convenient reformulation of sequential continuity of a linear functional on a test function space in terms of an estimate. If you are familiar with basic functional analysis you may compare this to the result that a linear map on a Banach space is continuous if and only if it is bounded.

**Proposition 24.1.12** (Sequential continuity of linear functionals on \(\mathcal{D}\)).

For a linear map \(A : \mathcal{D}(\Omega) \to \mathbb{C}\) the following are equivalent:

1. \(A\) is sequentially continuous.
2. \(A\) is sequentially continuous at \(0\).
3. For every compact set \(K \subset \Omega\) there exist \(C > 0\) and \(m \in \mathbb{N}\) such that

\[
|Au| \leq C \sum_{|\alpha| \leq m} \|\partial^\alpha u\|_{L^\infty(K)}
\]

(24.1.10)
for all \( u \in \mathcal{D}(\Omega) \) with \( \text{supp} \, u \subset K \).

**Proof.** Clearly (1) implies (2). Conversely, if (2) holds and if \( u_k \to u \), then we have \( v_k := u_k - u \to 0 \), so \( Au_k \to 0 \) by (2). By linearity it follows that \( Au_k = Av_k + Au \to Au \).

Next we show that (3) implies (2). To this end, suppose that \( u_k \to 0 \). Then there exists \( K \subset \Omega \) compact such that \( \text{supp} \, u_k \subset K \) for all \( k \in \mathbb{N} \) and \( \| \partial^\alpha u_k \|_{L^\infty(\Omega)} \to 0 \) for all \( \alpha \in \mathbb{N}^n \). Let \( C > 0 \) and \( m \in \mathbb{N} \) be such that (24.1.10) holds for this \( K \). Then we have

\[
|Au_k| \leq C \sum_{|\alpha| \leq m} \| \partial^\alpha u_k \|_{L^\infty(\Omega)} \to 0.
\]

This shows that (2) holds.

Conversely, suppose that (2) holds. Suppose for the sake of a contradiction that (3) does not hold. Then there exist \( K \subset \Omega \) compact and a sequence \((u_k)_k \in \mathcal{D}(\Omega) \) such that

\[
|Au_k| > k \sum_{|\alpha| \leq k} \| \partial^\alpha u_k \|_{L^\infty(\Omega)},
\]

holds for all \( k \in \mathbb{N} \). This implies that \( \|u_k\|_{L^\infty(\Omega)} > 0 \) for all \( k \in \mathbb{N} \) since otherwise \( u_k = 0 \) and so \( Au_k = 0 \) by linearity, which contradicts (2). Therefore we can rescale and set

\[
v_k := \frac{u_k}{k \sum_{|\alpha| \leq k} \| \partial^\alpha u_k \|_{L^\infty(\Omega)}} \in \mathcal{D}(\Omega).
\]

Clearly \( \text{supp} \, v_k \subset K \). We claim that \( v_k \to 0 \) in \( \mathcal{D} \). To see this, let \( \alpha \in \mathbb{N}^n \). For every \( k \in \mathbb{N} \) with \( k > |\alpha| \) we have

\[
\| \partial^\alpha v_k \|_{L^\infty(\Omega)} = \frac{1}{k \sum_{|\beta| \leq k} \| \partial^\beta u_k \|_{L^\infty(\Omega)}} \| \partial^\alpha u \|_{L^\infty(\Omega)} < \frac{1}{k}.
\]

Hence \( \| \partial^\alpha v_k \|_{L^\infty(\Omega)} \to 0 \) and so \( v_k \to 0 \) in \( \mathcal{D} \). However, by the choice of \( v_k \) and (2) it follows that \( |Av_k| > 1 \) holds for all \( k \in \mathbb{N} \). This contradicts \( Av_k \to 0 \). Hence (3) must hold after all. \( \square \)

Similar results of course also hold for linear functions on \( \mathcal{E} \) and \( \mathcal{S} \), as you are asked to prove in the exercises \[27.69\] and \[27.70\]. Let us illustrate the use of this result in an example.

**Example 24.1.13.** The function \( \frac{1}{x} \) is not locally integrable on \( \mathbb{R} \), so it cannot directly be interpreted as a distribution. However, exploiting the fact that it is an odd function, we can define a distribution \( \langle \text{v.p.} \left( \frac{1}{x} \right), u \rangle := \text{V.P.} \int_{-\infty}^{\infty} \frac{u(x)}{x} \, dx := \lim_{\varepsilon \to 0} \int_{|x| \geq \varepsilon} \frac{u(x)}{x} \, dx \) (24.1.11) for any \( u \in \mathcal{D}(\mathbb{R}) \). To see that this integral is well-defined, note that \( \frac{1}{x} \) being odd we have

\[
\int_{|x| \geq \varepsilon} \frac{u(x)}{x} \, dx = \frac{1}{2} \int_{|x| \geq \varepsilon} \frac{u(x) - u(-x)}{x} \, dx,
\]

and that the function \( \frac{u(x) - u(-x)}{x} \) has a continuous extension to \( x = 0 \) by \( 2u'(0) \) while still being compactly supported. It is clear that (24.1.11) is linear in \( u \), and we are going to use Proposition \[24.1.12\] to show that it is also sequentially continuous. To this end, let \( K \subset \mathbb{R} \) be compact and \( u \in \mathcal{D}(\mathbb{R}) \) with \( \text{supp} \, u \subset K \). Pick \( N > 0 \) with \( K \subset [-N, N] \). By the mean value theorem, for any \( x \in \mathbb{R} \setminus \{0\} \) there exists \( \theta \) between \(-x\) and \( x \) such
that \( \frac{u(x) - u(-x)}{2x} = \delta(0) \). Therefore the function \( \frac{u(x) - u(-x)}{2x} \) is bounded in absolute value by \( \|u'\|_{L^\infty(K)} \). It follows that for any \( \epsilon > 0 \) we have

\[
\left| \int_{|x| \geq \epsilon} \frac{u(x)}{x} \, dx \right| = \left| \int_{N \geq |x| \geq \epsilon} \frac{u(x)}{x} \, dx \right| = \left| \int_{N \geq |x| \geq \epsilon} \frac{u(x) - u(-x)}{2x} \, dx \right| \leq 2N\|u'\|_{L^\infty(K)}.
\]

It follows that also \( \langle \psi, u \rangle \leq 2N\|u'\|_{L^\infty(K)} \) holds which by Proposition 24.1.12 implies sequential continuity of \( \psi \).

24.1.3. Convergence of distributions. We started with the example of a “collapsing” sequence of functions as motivation for the notion of distributions. We now introduce a notion of convergence of distributions with respect to which such a sequence of functions has a distributional limit.

**Definition 24.1.14 (Convergence of distributions).** A sequence \( (T_n)_{n \in \mathbb{N}} \) in \( \mathcal{D}'(\Omega) \) is said to \( \mathcal{D}' \)-converge to some \( T \in \mathcal{D}'(\Omega) \) if

\[
\forall u \in \mathcal{D}(\Omega) : \ T_n(u) \to T(u). \tag{24.1.12}
\]

Convergence of compactly supported and tempered distributions is defined analogously.

In physics one often says that one first has to “smear with a test function” and then take the limit. Mathematically, (24.1.12) is of course nothing but pointwise convergence of the sequence when viewed as functions from \( \mathcal{D}(\Omega) \) to \( \mathbb{C} \). With respect to these notions of convergence, the inclusions \( \mathcal{D}' \subseteq \mathcal{S}' \subseteq \mathcal{E}' \) become sequentially continuous.

To illustrate convergence in the sense of distributions, we describe a common way of approximating the \( \delta \)-distribution by regular (even smooth) distributions. This will also be important later for showing that any distribution can be approximated by smooth functions.

**Definition 24.1.15 (Dirac \( \delta \)-sequences).** A sequence \( (\psi_k)_{k \in \mathbb{N}} \) in \( C^\infty(\mathbb{R}^n) \cap L^1(\mathbb{R}^n) \) is called a (Dirac) \( \delta \)-sequence or approximate identity if it satisfies the following properties:

1. \( L^1 \)-equiboundedness: \( \sup_{k \in \mathbb{N}} \left\| \psi_k \right\|_{L^1(\mathbb{R}^n)} < \infty \).

2. Normalization: For every \( k \in \mathbb{N} \) one has \( \int_{\mathbb{R}^n} \psi_k(x) \, dx = 1 \).

3. Concentration at 0: For every \( R > 0 \) one has \( \lim_{k \to \infty} \left( \int_{\{ |x| > R \}} |\psi_k(x)| \, dx \right) = 0 \).

If additionally \( \psi_k \in C^\infty_0(\mathbb{R}^n) \) for all \( k \in \mathbb{N} \) and \( \text{supp} \psi_k \to \{0\} \) in the sense that for every \( R > 0 \) there exists \( N \in \mathbb{N} \) such that \( \text{supp} \psi_k \subseteq B_R(0) \) for all \( k \geq N \), then one speaks of a strict \( \delta \)-sequence.

We now show that any such sequence converges to the \( \delta \)-distribution. In exercise 27.71 and 27.72 you find concrete examples of \( \delta \)-sequences.

**Theorem 24.1.16 (Smooth approximation of the \( \delta \)-distribution).** Let \( (\psi_k) \) be a \( \delta \)-sequence. Then \( \psi_k \to \delta_0 \) in \( \mathcal{D}' \) and in \( \mathcal{S}' \). If \( (\psi_k) \) is a strict \( \delta \)-sequence, then also \( \psi_k \to \delta_0 \) in \( \mathcal{E}' \).

**Proof.** Let \( u \in \mathcal{D}(\mathbb{R}^n) \) be given. We wish to show that \( \int_{\mathbb{R}^n} \psi_k(x) f(x) \, dx \to f(0) \). By the normalization condition we can write

\[
\int_{\mathbb{R}^n} \psi_k(x) f(x) \, dx - f(0) = \int_{\mathbb{R}^n} \psi_k(x) (f(x) - f(0)) \, dx.
\]
Now we split this integral into large and small \(|x|\) and use that \(|f(x) - f(0)|\) will be small for small \(|x|\), whereas \(|\psi_k(x)|\) will be small for large \(|x|\) (and large \(k\)). More precisely, let \(\epsilon > 0\) be given. Choose \(\delta > 0\) such that \(|f(x) - f(0)| < \epsilon\) for \(|x| < \delta\). Moreover, choose \(N \in \mathbb{N}\) such that \(\int_{\mathbb{R}^n \setminus B_3(0)} |\psi_k(x)| \, dx < \epsilon\) for all \(k \geq N\). Then for all \(k \geq N\) we have

\[
\left| \int_{\mathbb{R}^n} \psi_k(x)(f(x) - f(0)) \, dx \right| \leq \int_{B_3(0)} |\psi_k(x)| \cdot |f(x) - f(0)| \, dx + \int_{\mathbb{R}^n \setminus B_3(0)} |\psi_k(x)| \cdot |f(x) - f(0)| \, dx
\]

\[
\leq \epsilon (\|\psi_k\|_{L^1} + 2 \|f\|_{\infty}).
\]

As \(\epsilon > 0\) was arbitrary, this shows convergence of the left-hand side to 0.

For \(u \in \mathcal{S}(\mathbb{R}^n)\) the exact same argument works as well. For \(u \in \mathcal{E}(\mathbb{R}^n)\) and \((\psi_k)_k\) a strict \(\delta\)-sequence, one replaces \(u\) by \(\chi u\) for some \(\chi \in C_0^\infty(\mathbb{R}^n)\) with \(\chi(0) = 1\) and \(\text{supp} \psi_k \subseteq \text{supp} \chi\) for all sufficiently large \(k\).

Dirac \(\delta\)-sequences illustrate how a sequence of functions which concentrate more and more “mass” into a small space may have a limit in the sense of distributions. The following example shows another way in which a sequence of functions, which clearly do not converge as functions, may have a distributional limit.

**Example 24.1.17.** For \(k \in \mathbb{N}\), let \(\psi_k \in C^\infty(\mathbb{R})\) be given by \(\psi_k(x) = \sin(kx)\). For \(u \in \mathcal{D}'(\mathbb{R})\), by integration by parts we have

\[
\langle \psi_k, u(x) \rangle = \int_{\mathbb{R}} \sin(kx)u(x) \, dx = \int_{\mathbb{R}} -\frac{1}{k} \left( \cos(kx) \right)'u(x) \, dx = \frac{1}{k} \int_{\mathbb{R}} \cos(kx)u'(x) \, dx \to 0.
\]

This shows that \(\psi_k \to 0\) in \(\mathcal{D}'(\mathbb{R}^n)\) for \(k \to \infty\). In intuitive terms, the ever faster oscillations of \(\psi_k\) as \(k\) increases lead to cancellations in the integral which results in the convergence to 0.

Finally, as yet another illustration for computing distributional limits, we prove the following famous formula which relates two distributional “renormalizations” of \(\frac{1}{x}\).

**Theorem 24.1.18 (Sokhotski-Plemelj formula).** The family of functions \((\frac{1}{x \pm \nu})_{\nu > 0}\) has a distributional limit in \(\mathcal{D}'(\mathbb{R})\) for \(\nu \to 0\), denoted by \(\frac{1}{x} \pm \nu\cdot\pi\delta\). This limit satisfies

\[
\frac{1}{x \pm i0} := \lim_{\nu \to 0^+} \frac{1}{x \pm i\nu} = \text{v.p.} \left( \frac{1}{x} \right) \mp i\pi\delta,
\]

where \(\text{v.p.} \left( \frac{1}{x} \right)\) denotes the Cauchy principal value of \(\frac{1}{x}\) defined in Example 24.1.13.

**Proof.** Let \(u \in \mathcal{D}(\mathbb{R})\) be given. As in example 24.1.13 we are going to cleverly exploit that \(\frac{1}{x}\) is odd. Observe that as a consequence we have \(\langle \text{v.p.} \left( \frac{1}{x} \right), \varphi \rangle = 0\) for any even function \(\varphi \in \mathcal{D}(\mathbb{R})\). Since the function \(\varphi(x) := \frac{u(0)}{1+x^2}\) is even we have \(\langle \text{v.p.} \left( \frac{1}{x} \right), u \rangle = \langle \text{v.p.} \left( \frac{1}{x} \right), u - \varphi \rangle\). On the other hand,

\[
\int_{\mathbb{R}} \frac{1}{x \pm i\nu} u(x) \, dx = \int_{\mathbb{R}} \frac{u(x) - \varphi(x)}{x \pm i\nu} \, dx + u(0) \int_{\mathbb{R}} \frac{1}{(x \pm i\nu)(1 + x^2)} \, dx.
\]

Regarding the first integral, note that by l’Hopitals rule

\[
\lim_{x \to 0} \frac{u(x) - \varphi(x)}{x} = \lim_{x \to 0} \frac{u'(x) - \varphi'(x)}{1} = u'(0),
\]

where \(\varphi(x) := \frac{u(0)}{1+x^2}\), \(\varphi'(x) = \frac{-2u(0)x}{(1+x^2)^2}\) and \(\lim_{x \to 0} \frac{-2u(0)x}{(1+x^2)^2} = 0\). Therefore

\[
\lim_{\nu \to 0^+} \int_{\mathbb{R}} \frac{1}{x \pm i\nu} u(x) \, dx = u'(0).
\]
so the function \( \frac{u(x) - \varphi(x)}{x} \) extends continuously to all of \( \mathbb{R} \). Moreover, since \( u \) has compact support and \( \varphi \) decays as \( \frac{1}{x^2} \) at infinity, it is also integrable. Therefore, an application of dominated convergence yields that

\[
\lim_{\epsilon \to 0} \int_{\mathbb{R}} \frac{u(x) - \varphi(x)}{x \pm i\epsilon} \, dx = \int_{\mathbb{R}} \frac{u(x) - \varphi(x)}{x} \, dx
\]

\[
= \lim_{\delta \to 0} \int_{|x| \geq \delta} \frac{u(x) - \varphi(x)}{x} \, dx
\]

\[
= \langle \text{v.p.} \left( \frac{1}{x} \right), u - \varphi \rangle
\]

\[
= \langle \text{v.p.} \left( \frac{1}{x} \right), u \rangle .
\]

For any \( \epsilon > 0 \), the second integral in \((*)\) can easily be computed by the residue theorem, closing the contour in the upper half-plane. In the limit \( \epsilon \to 0 \) this leaves \( \pm i\pi \), just as desired.

\[\square\]

Next we show that the inclusions of various functions into distributions (as regular distributions) are sequentially continuous. Concretely this simply means that convergence as (test) functions implies convergence as distribution.

**Theorem 24.1.19.** All of the following maps, which map a function \( f \) to the regular distribution \( T_f \), are linear, injective and sequentially continuous:

\[
\mathcal{E} (\Omega) \to \mathcal{D}' (\Omega), \quad \mathcal{D} (\Omega) \to \mathcal{E}' (\Omega), \quad L^p (\Omega) \to \mathcal{D}' (\Omega), \quad L^p (\mathbb{R}^n) \to \mathcal{S}' (\Omega).
\]

**Proof.** Linearity is obvious and injectivity follows from the fundamental theorem in the calculus of variations. We show sequential continuity in the first case, the other cases can be handled similarly and are left as exercise. Suppose that \( f_n \to f \) in \( \mathcal{E} (\Omega) \). By linearity we may assume \( f = 0 \), otherwise we consider the difference \( f_n - f \). Let \( u \in \mathcal{D} (\Omega) \) be given. Since \( \text{supp } u \subset \Omega \) is compact, by definition of \( \mathcal{E} \)-convergence the functions \( f_n \) converge uniformly to zero on \( \text{supp } f \). This implies that there exists a constant \( C > 0 \) such that \( |f_n| < C \) on \( \text{supp } u \) for all \( n \in \mathbb{N} \), and thus \( |f_n u| \leq C |u| \in C^0_0 (\Omega) \subset L^1 (\Omega) \).

Since clearly also \( f_n (x) \to 0 \) for every \( x \in \Omega \) it now follows by dominated convergence that \( \int_{\Omega} f_n (x) u (x) \, dx \to 0 \). Hence \( f_n \to 0 \) also holds in \( \mathcal{D}' (\Omega) \). \(\square\)

**24.1.4. Calculus with distributions.** One motivation for studying distributions is to solve (linear) differential equations. To this end, one first needs to specify how a (linear) partial differential operator acts on a distribution, in particular how to differentiate a distribution and how to multiply it by a smooth function. The guideline for this is that on smooth functions, viewed as regular distributions, differentiation and multiplication by smooth functions should work in the same way one is used to. To explain what this means, let \( f \in C^\infty (\mathbb{R}) \) be a smooth function and let \( T_f \in \mathcal{D}' (\mathbb{R}) \) be the corresponding (regular) distribution. Concerning differentiation, compatibility means that we would like to define \( (T_f)' \) in such a way that \( (T_f)' = T_{f'} \) holds. Since for any test function \( u \in \mathcal{D} (\mathbb{R}) \) integration by parts gives

\[
\langle T_{f'}, u \rangle = \int_{\mathbb{R}} f'(x) u(x) \, dx = - \int_{\mathbb{R}} f(x) u'(x) \, dx = - \langle T_f, u' \rangle ,
\]

we should define \( (T_f)' \) by the very right-hand side in this equation. Note that in this right-hand side we can replace \( T_f \) by any, not necessarily regular, distribution \( T \).
Definition 24.1.20. Let $T \in \mathcal{D}'(\Omega)$. For $\alpha \in \mathbb{N}^n$ and for $f \in C^\infty(\Omega)$ we define $\partial^\alpha T \in \mathcal{D}'(\Omega)$ and $fT \in \mathcal{D}'(\Omega)$ as the distributions acting on a test function $u \in \mathcal{D}(\Omega)$ by
\[
\langle \partial^\alpha T, u \rangle := (-1)^{|\alpha|} \langle T, \partial^\alpha u \rangle, \quad (24.1.14)
\]
\[
\langle fT, u \rangle := \langle T, fu \rangle. \quad (24.1.15)
\]
For $\mathcal{E}'(\Omega)$ and $\mathcal{S}'(\mathbb{R}^n)$ one uses the same definitions, except that tempered distributions may only be multiplied by functions which grow at most polynomially at infinity (so-called functions of “moderate growth”, see exercise 27.59).

Of course one has to verify that (24.1.14) and (24.1.15) indeed define distributions, i.e., sequentially continuous linear functionals. This is not difficult and left as exercise.

Note that any distribution can be infinitely often differentiated. Therefore, if viewed as distribution, also a function which in the classical sense is not differentiable, may be differentiated as a distribution. As before, we illustrate these new notions by concrete examples. More examples can be found in the exercises.

Example 24.1.21. The Heaviside function $\Theta : \mathbb{R} \to \mathbb{R}$ is defined by
\[
\Theta(x) = \begin{cases} 
0 & x \leq 0 \\
1 & x > 0.
\end{cases} \quad (24.1.16)
\]
Clearly $\Theta \in L^1_{\text{loc}}(\mathbb{R})$ since it is bounded, so we can view $\Theta$ as a distribution. Computing the derivative $\Theta'$ as a distribution is actually quite easily: For any $u \in \mathcal{D}(\mathbb{R})$ we have
\[
\langle \Theta', u \rangle := -\int_\mathbb{R} \Theta(x)u'(x) \, dx = -\int_0^\infty u'(x) \, dx = u(0) = \langle \delta_0, u \rangle.
\]
Hence $\Theta' = \delta_0$, as intuition might have suggested.

Example 24.1.22. For $x_0 \in \mathbb{R}^n$ and $f \in C^\infty(\mathbb{R}^n)$, we compute $f\delta_{x_0} \in \mathcal{D}'(\mathbb{R}^n)$. To this end, note that for any $u \in \mathcal{D}(\mathbb{R}^n)$ we have
\[
\langle f\delta_{x_0}, u \rangle := \langle \delta_{x_0}, fu \rangle = f(x_0)u(x_0) = \langle f(x_0)\delta_{x_0}, u \rangle.
\]
This shows that $f\delta_{x_0} = f(x_0)\delta_{x_0}$.

The usual rules of differentiation also apply to distributions as far as they make sense. They are easy to prove by simply going back to the definitions, therefore the verifications are left as exercise.

Proposition 24.1.23 (Calculus rules for distributions). Let $T \in \mathcal{D}'(\Omega)$.

1. For any $\alpha, \beta \in \mathbb{N}^n$ one has $\partial^\alpha (\partial^\beta T) = \partial^{\alpha+\beta} T$.
2. For any $f, g \in C^\infty(\Omega)$ one has $f(gT) = (fg)T$.
3. For any $\alpha \in \mathbb{N}^n$ and $f \in C^\infty(\Omega)$ the Leibniz rule holds:
\[
\partial^\alpha (fT) = \sum_{\beta \leq \alpha} \binom{\alpha}{\beta} (\partial^{\alpha-\beta} f)(\partial^\beta T) \quad (24.1.17)
\]

By the fundamental theorem of calculus, any (continuously) differentiable function can be reconstructed from its derivative up to an additive constant. We now prove the same statement for distributions, the proof may also serve as further illustration of how to work with distributions.

Proposition 24.1.24. Let $S, T \in \mathcal{D}'(\mathbb{R})$ and let $f \in C(\mathbb{R})$. 

(1) If \( T' = T_f \), then \( T = T_F \) for some antiderivative \( F \in C^1(\mathbb{R}) \) of \( f \).

(2) If \( S' = T' \), then \( S \) and \( T \) coincide "up to a constant", i.e., there exists \( c \in \mathbb{R} \) such that \( \langle S, u \rangle = (T, u) + \int_{\mathbb{R}} cu(x) \, dx \) for all \( u \in D(\mathbb{R}) \).

**Proof.** *Hint:* It might be helpful to show first that \( v \in D(\mathbb{R}) \) satisfies \( v = v' \) for some other test function \( w \in D(\mathbb{R}) \) if and only if \( \int_{\mathbb{R}} v(x) \, dx = 0 \). Then decompose a general \( u \in D(\mathbb{R}) \) as \( u = u_1 + u_2 \) where \( \int_{\mathbb{R}} u_1(x) \, dx = 0 \) by using a cutoff function \( \chi \) with \( \int_{\mathbb{R}} \chi(x) \, dx = 1 \).

Note that the first part of the previous proposition may be viewed as a (very simple) "regularity result": If a distribution \( T \) satisfies the differential equation \( T' = w \) with a regular (continuous) right-hand side, then \( T \) itself is regular (actually one degree more regular than the right-hand side).

Having defined differentiation and multiplication by smooth functions, we can now define how a linear differential operator (with smooth coefficients) acts on distributions. Concretely, if \( A = \sum_{|\alpha| \leq m} a_\alpha \partial^\alpha \) is a linear differential operator with certain coefficients \( a_\alpha \in C^\infty(\Omega) \), then for \( T \in D'(\Omega) \) we define

\[
AT := \sum_{|\alpha| \leq m} a_\alpha (\partial^\alpha T).
\]  

(24.1.18)

With this definition it now makes sense to say that a distribution \( T \) satisfies the Laplace equation \( \Delta T = 0 \) or the wave equation \( (\partial^2_t - D)T = 0 \), or any other linear partial differential equation, also replacing \( 0 \) on the right-hand side by an inhomogeneity. Be careful however that at this point it is not yet clear how one might impose initial or boundary values. In concrete calculations one should also be careful that when unwinding \( \langle AT, u \rangle \) for some \( u \in D(\Omega) \), one obtains

\[
\langle AT, u \rangle = \sum_{\alpha \leq m} \langle a_\alpha (\partial^\alpha T), u \rangle = \sum_{\alpha \leq m} \langle \partial^\alpha T, a_\alpha u \rangle = \sum_{\alpha \leq m} (-1)^{|\alpha|} \langle T, \partial^\alpha (a_\alpha u) \rangle,
\]

where on the right-hand side also the coefficients of \( A \) are differentiated. This is of course just as it would happen in an integration by parts computation.

Particularly important distributional solutions of linear partial differential equations are obtained by setting the right-hand side equal to the \( \delta \)-distribution.

**Definition 24.1.25** (Green’s function). If \( A \) is linear differential operator on some open set \( \Omega \subset \mathbb{R}^n \), and \( x_0 \in \Omega \), then a distribution \( G_{x_0} \in D'(\Omega) \) satisfying \( AG_{x_0} = \delta_{x_0} \), where \( \delta_{x_0} \) is the \( \delta \)-distribution at \( x_0 \), is called Green’s function in \( x_0 \) for \( A \).

In the exercises 27.84 - 27.89 you find examples of Green’s functions for various equations. Note that Green’s functions are in general only unique up to addition of solutions of the homogeneous equation \( AT = 0 \).

We close this section with the remark that that differentiation and multiplication are sequentially continuous linear maps on test function spaces and distribution spaces. The proofs, which are left to the reader in exercises 27.83 provides an excellent opportunity to consolidate the various previous definitions.

**24.1.5. Fourier transform.** The Fourier transform is an important tool, for instance in the study of differential equations, and it is very useful that it can also be generalized to (tempered) distributions. Before doing so, we recall the definition and some important properties about the Fourier transform of functions. For further reading one may for instance consult [Stein and Shakarchi, Fourier Analysis].
For \( f \in L^1(\mathbb{R}^n) \), the Fourier transform is the function \( \hat{f} : \mathbb{R}^n \to \mathbb{C} \) defined by
\[
\hat{f}(\xi) := \int_{\mathbb{R}^n} f(x)e^{-ix \cdot \xi} \, dx,
\]
where \( x \cdot \xi = x^1\xi^1 + \ldots + x^n\xi^n \). Instead of \( \hat{f} \) we will also write \( F(f) \) for the Fourier transform of \( f \). The basic properties of the Fourier transform are most easily stated for Schwartz functions, but many of the following properties however also hold for more general classes of functions. Most of these properties can be proved by straightforward manipulations of integrals, and are left to the reader in exercise 27.91.

**Theorem 24.1.26 (Basic properties of the Fourier transform on \( S(\mathbb{R}^n) \)).**

Let \( f, g \in S(\mathbb{R}^n) \). Then the following hold:

1. We have \( \hat{\partial_\alpha f}(\xi) = (i\xi)^\alpha \hat{f}(\xi) \) and \( \hat{x^\alpha f}(\xi) = i|\alpha| \partial_\alpha \xi \hat{f}(\xi) \). (24.1.20)

2. It holds that
\[
\int_{\mathbb{R}^n} f(x) \overline{g(x)} \, dx = \int_{\mathbb{R}^n} \hat{f}(\xi) \overline{\hat{g}(\xi)} \, d\xi.
\]

3. As a map \( F : S(\mathbb{R}^n) \to S(\mathbb{R}^n) \), the Fourier transform is a sequentially continuous isomorphism. The inverse is given explicitly by the Fourier inversion formula:
\[
f(x) = \frac{1}{(2\pi)^n} \int_{\mathbb{R}^n} \hat{f}(\xi)e^{ix \cdot \xi} \, d\xi.
\]

4. Parseval’s relation holds:
\[
\int_{\mathbb{R}^n} f(x) \overline{g(x)} \, dx = \int_{\mathbb{R}^n} \hat{f}(\xi) \overline{\hat{g}(\xi)} \, d\xi.
\]

5. Plancharel’s formula holds:
\[
\|f\|_{L^2(\mathbb{R}^n)} = \|\hat{f}\|_{L^2(\mathbb{R}^n)}.
\]

6. The convolution theorem holds:
\[
\hat{f \ast g} = \hat{f} \cdot \hat{g}.
\]

□

Now we move on to the Fourier transform of a distribution. For this purpose, there are actually two possible definitions: On the one hand the definition (24.1.19) of the Fourier transform may be read as the action of the (regular) distribution \( T_f \) on the function \( x \mapsto e^{-ix \cdot \xi} \). On the other hand, the identity (24.1.21) suggests to define the Fourier transform of a distribution by “moving the Fourier transform to the test function” (similar as for differentiation). We will take the second approach since this allows to define the Fourier transform for a larger class of distributions, namely all tempered distributions, whereas the first approach is restricted to compactly supported distributions (since the imaginary exponential function does not decay). In the end we are going to show that for compactly supported distributions both approaches coincide.

**Definition 24.1.27 (Fourier transform of tempered distributions).** For a tempered distribution \( T \in S'(\mathbb{R}^n) \), its Fourier transform \( \hat{T} \in S'(\mathbb{R}^n) \) is defined by
\[
\langle \hat{T}, u \rangle := \langle T, \hat{u} \rangle
\]
for all \( u \in S(\mathbb{R}^n) \). Instead of \( \hat{T} \) we will also write \( F(T) \).
Note that \( \hat{T} \) is indeed linear and sequentially continuous again since it is simply the composition \( \hat{T} = T \circ F \) of two linear and sequentially continuous maps on \( S \). Even more, since \( F : S(\mathbb{R}^n) \to S(\mathbb{R}^n) \) is an isomorphism, the same holds for the Fourier transform on \( S'(\mathbb{R}^n) \). It is also not difficult to verify that the “calculus rules” (24.1.20) continue to hold, we leave the proof to the reader in exercise 27.92.

**Theorem 24.1.28 (Fourier transform on tempered distributions).** The Fourier transform is a sequentially continuous linear isomorphism \( F : S'(\mathbb{R}^n) \to S'(\mathbb{R}^n) \). Moreover, for any \( T \in S'(\mathbb{R}^n) \) and \( \alpha \in \mathbb{N}^n \) we have

\[
\hat{(\partial_x^\alpha T)} = (i \xi)^\alpha \hat{T}(\xi) \quad \text{and} \quad \hat{(x^\alpha T)} = i^{\vert \alpha \vert} \partial_x^\alpha \hat{T}.
\] (24.1.27)

Concerning (24.1.27), note that the multiplication of a tempered distribution by a monomial such as \( x^\alpha \) or \( \xi^\alpha \) is well-defined since monomials are clearly of moderate growth (cf. exercise 27.59).

Let us make some examples. The second one shows that the extension of the Fourier transform to distributions can also help to compute the Fourier transforms of functions.

**Example 24.1.29.** Let \( \delta \in S'(\mathbb{R}^n) \) be the \( \delta \)-distribution (at the origin). For \( u \in S(\mathbb{R}^n) \) we have

\[
\langle \hat{\delta}, u \rangle \overset{\text{Def.}}{=} \langle \delta, \hat{u} \rangle = \hat{u}(0) = \int_{\mathbb{R}^n} e^{i0 \cdot x} u(x) \, dx = \int_{\mathbb{R}^n} 1 \cdot u(x) \, dx = \langle 1, u \rangle.
\]

This shows that \( \hat{\delta} = 1 \). In particular, the Fourier transform of \( \hat{\delta} \) is a regular distribution.

**Example 24.1.30.** Consider the function \( f \in L^1(\mathbb{R}) \) defined by \( f(t) = -e^t \) for \( t \leq 0 \) and \( f(t) = e^{-t} \) for \( t > 0 \). It is easy to verify that \( f'' = 2\delta' + f \) holds in \( S' \). From (24.1.27) and the previous example it follows that \( -\xi^2 \hat{f} = 2i \xi + \hat{f} \) holds, so \( \hat{f} = \frac{-2i \xi}{1+i \xi^2} \).

We close this section by connecting our definition of the distributional Fourier transform to the alternative approach of defining the distributional Fourier transform by viewing the definition of the Fourier transform of a function as application of a regular distribution to the test function \( e^{-i x \cdot \xi} \). While this only works for compactly supported distributions, this connection proves that the Fourier transform of such a distribution is in fact a smooth (even analytic) function.

**Theorem 24.1.31.** Let \( T \in \mathcal{E}'(\mathbb{R}^n) \subset S'(\mathbb{R}^n) \). Then its Fourier transform \( \hat{T} \) is actually a smooth function of moderate growth, and is explicitly given by

\[
\hat{T}(\xi) = \langle T(x), e^{-i x \cdot \xi} \rangle.
\] (24.1.28)

Moreover, \( \hat{T} \) can be extended from \( \mathbb{R}^n \) to \( \mathbb{C}^n \) as an entire function by this formula.

**Proof.** Can probably be done by first approximating \( T \) and then controlling everything in the limit. \( \Box \)

It is a general and important theme that smoothness properties of a function or distribution and decay properties of its Fourier transform are related. This already shines through from how the Fourier transform interchanges differentiation and multiplication by monomials, cf. (24.1.27).
24.1.6. Support and singular support of a distribution. Intuitively the δ-distribution (at the origin) is supported in the origin since it is “zero everywhere but at the origin (where it is infinite)” . Note however that it is of course not sensible to speak about the value of a distribution “at a point”, rather one always has to “smear with a test function”. It is, however, still possible to define the notion of support of a distribution.

Definition 24.1.32 (Localization and support of distributions). Let $S, T \in \mathcal{D}'(\Omega)$.

1. For an open subset $U \subset \Omega$ we say that $S$ and $T$ coincide on $U$ and write $S|_U = T|_U$ if $\langle S, u \rangle = \langle T, u \rangle$ for all $u \in \mathcal{D}(\Omega)$ with $\text{supp} u \subset U$.
2. We say that $T$ vanishes in a point $x \in \Omega$ if there exists an open neighborhood $U \subset \Omega$ of $x$ such that $T|_U = 0$.
3. The set
   \[
   \text{supp} T := \Omega \setminus \{ x \in \Omega \mid T \text{ vanishes at } x \}
   \]
   (24.1.29)
   is called the support of $T$.

As first example we show that for a continuous function, its support as a function and as a distribution coincide. In the example we further illustrate an important technique for proving local properties of distributions by suitable choices of cutoff functions.

Example 24.1.33. For $f \in C(\Omega)$ we show that $\text{supp} T_f = \text{supp} f$. Recall that
\[
\text{supp} f = \overline{\{ x \in \Omega \mid f(x) \neq 0 \}},
\]
where the overline denotes closure in $\Omega$. One the one hand, if $x \notin \text{supp} f$, then since $\text{supp} f \subset \Omega$ is closed there exists an open neighborhood $U \subset \Omega$ of $x$ such that $f|_U = 0$. Now clearly $\langle T_f, u \rangle = 0$ for any $u \in \mathcal{D}(\Omega)$ with $\text{supp} u \subset U$, hence $x \notin \text{supp} T_f$. This shows that $\text{supp} T_f \subset \text{supp} f$ holds. Conversely, let $x \in \text{supp} f$. Choose any open neighborhood $U \subset \Omega$ of $x$ and $x' \in U$ with $f(x') \neq 0$ (note that $f(x)$ may be zero). After possibly rescaling $f$, we may assume that $f(x') = 2$. By continuity, there exists $\epsilon > 0$ sufficiently small such that $B_\epsilon(x') \subset U$ and $f|_{B_\epsilon(x')} > 1$. Now pick a cutoff function $\chi \in \mathcal{D}(\Omega)$ with $\text{supp} \chi \subset B_\epsilon(x) \subset U$, $0 \leq \chi \leq 1$, and $\chi|_{B_{\epsilon/2}(x)} = 1$. Then $\langle T_f, \chi \rangle \geq |B_{\epsilon/2}(x)| > 0$. Therefore $x \in \text{supp} T_f$.

Exercise [27.3] explains the terminology “compactly supported distribution” for the elements of $\mathcal{E}'$. These are in fact precisely those elements of $\mathcal{D}'$ whose support is compact.

As another illustration of how to use the concept of support of a distribution we completely characterize all those distributions whose support consists of a single point.

Theorem 24.1.34 (Distributions supported at a point). Let $T \in \mathcal{D}'(\Omega)$ with $\text{supp} T = \{ x_0 \}$. Then there exist (unique) $m \in \mathbb{N}$ and $c_\alpha \in \mathbb{C}$ for all $\alpha \in \mathbb{N}$ with $|\alpha| \leq m$ such that
\[
T = \sum_{|\alpha| \leq m} c_\alpha \partial^\alpha \delta_{x_0}.
\]
(24.1.30)

Proof. Reference, or exercise. \[ \square \]

Besides the support, another local concept is that of the singular support which describes where a distribution is not smooth. Similar as the support, this is expressed via the action on test functions.

Definition 24.1.35. Let $T \in \mathcal{D}'(\Omega)$. We say that $T$ is smooth at a point $x \in \Omega$ if there exists an open neighborhood $U \subset \Omega$ of $x$ and $f \in C^\infty(U)$ with $T|_U = T_f$. The set
\[
\text{singsupp} u := \Omega \setminus \{ x \in \Omega \mid u \text{ is smooth at } x \}
\]
(24.1.31)
is called the singular support of $u$. 
Be careful that even a regular distribution can have nonempty singular support since regular distributions need only be locally integrable functions whereas the singular support is about (non-)smoothness (cf. exercise 27.9).

**Example 24.1.36.** We show that $\text{singsupp} \delta_0 = \{0\}$. Clearly $\text{singsupp} \delta_0 \subset \{0\}$ since for any $x \neq 0$ and any $f \in \mathcal{D}(\mathbb{R})$ with $\text{supp} f \subset B_{|x|/2}(x)$ we have $\langle \delta, f \rangle = f(0) = 0$, hence $\delta_0$ coincides on $B_{|x|/2}(x)$ with the smooth constant function 0. On the other hand, let $U \subset \mathbb{R}$ be any open neighborhood of 0. Pick $\epsilon > 0$ with $[\epsilon, \epsilon] \subset U$ and a sequence $(u_n)_n$ of test functions with $\text{supp} u_n \subset (-\epsilon, \epsilon)$, $\|u_n\|_{L^1} \leq 1$ and $u_n(0) = n$. Then $\langle \delta_0, u_n \rangle = u_n(0) = n \to \infty$, whereas for any $f \in C^\infty(U)$ we have

$$\left| \langle T_f, u_n \rangle \right| = \left| \int_{-\epsilon}^{\epsilon} f(x)u_n(x) \, dx \right| \leq \|f\|_{C^0([-\epsilon, \epsilon])} \|u_n\|_{L^1(\mathbb{R})} \leq \|f\|_{C^0([-\epsilon, \epsilon])}.$$  

Since this remains bounded for $n \to \infty$, it follows that $\delta_0|_U = T_f$. Hence $0 \in \text{singsupp} \delta_0$.

**24.1.7. Convolution and regularization of distributions.** We have already seen in Theorem 24.1.16 that the $\delta$-distribution can be approximated (in the sense of distributions) by smooth functions. The same is actually true for any distribution. Similarly as the approximation of non-smooth functions by smooth ones, this can be shown by the technique of convolution.

We start by recalling some important results about convolution of functions. First of all, the **convolution** of two functions $f$ and $g$ is the function $f * g$ defined by

$$(f * g)(x) := \int_{\mathbb{R}^n} f(x - y)g(y) \, dy. \quad (24.1.32)$$

In order for this integral to be well-defined, one of course needs to assume that $f$ and $g$ satisfy suitable integrability properties. One obvious sufficient condition is for instance that $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ and $g \in C_0(\mathbb{R}^n)$. We will focus on this case in the following, but also other conditions are sufficient (see ??).

One of the most interesting features of convolution is that $f * g$ is as smooth as the smoother one of the two functions $f$ and $g$. Therefore convolution can be used to “smoothen” non-smooth functions. This and other important properties of the convolution of functions are summarized in the following statement.

**Theorem 24.1.37 (Properties of convolution).**

1. Convolutions are commutative: For any $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ and $\psi \in C_0(\mathbb{R}^n)$ one has

$$(f * \psi)(x) = \int_{\mathbb{R}^n} f(x - y)\psi(y) \, dy = \int_{\mathbb{R}^n} f(y)\psi(x - y) \, dy = (\psi * f)(x). \quad (24.1.33)$$

2. Convolutions increase support: For any $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ and $\psi \in C_0(\mathbb{R}^n)$ one has

$$\text{supp}(f * \psi) \subset \text{supp} f + \text{supp} \psi. \quad (24.1.34)$$

3. Convolutions smooth: If $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ and $\psi \in C^k(\mathbb{R}^n)$ for $k \in \mathbb{N} \cup \{\infty\}$, then also $f * \psi \in C^k(\mathbb{R}^n)$. Moreover, for any $\alpha \in \mathbb{N}^n$ with $|\alpha| \leq k$ we have

$$\partial^\alpha(f * \psi) = f * \partial^\alpha \psi. \quad (24.1.35)$$

4. Smooth approximation by convolution: If $f \in L^1_{\text{loc}}(\mathbb{R}^n)$ and $(\psi_k)_k \subset C_0^{\infty}(\mathbb{R}^n)$ is a strict delta sequence (cf. Def. 24.1.15), then $f * \psi_k \to f$ pointwise almost everywhere for $k \to \infty$. Moreover, if $f$ is continuous then $f * \psi_k \to f$ uniformly on compact sets, if $f \in \mathcal{D}(\mathbb{R}^n)$ (or $\mathcal{E}, \mathcal{S}$), then $f * \psi_k \to f$ in $\mathcal{D}$ (or $\mathcal{E}, \mathcal{S}$).
For the proof see Exercise 27.17.

We now extend the process of convolution from test functions to distributions. To see how this may be reasonably defined, note that we can view (24.1.32) as action of the distribution \( T_g \) on the “shifted test function” \( y \mapsto f(x - y) \).

**Definition 24.1.38 (Convolution of distribution with test function).** For \( T \in \mathcal{D}(\mathbb{R}^n) \) and \( \psi \in C_0^\infty(\mathbb{R}^n) \), the convolution of \( u \) with \( \psi \) is the function \( T \ast \psi : \mathbb{R}^n \to \mathbb{C} \) defined by

\[
(T \ast \psi)(x) := \langle T(y), \psi(x - y) \rangle.
\]

The expression in the middle makes more explicit on which variable the distribution \( T \) acts.

Note that the convolution of a distribution with a smooth function is a function. Note also that by this definition we cannot take the convolution of two distributions (this is possible but requires further thought).

We now want to show that the convolution of a distribution with a smooth function is again a smooth function, and use this to show that any distribution can be approximated by a sequence of smooth functions. To do so, we need to study the dependence of \( \langle T, \psi(x - \cdot) \rangle \) on the “parameter” \( x \). Here the following more general result is useful, which can be viewed as a sort of generalization of “differentiating under the integral”. The proof proceeds by working out difference quotients explicitly without using any special tricks. We leave it as an exercise.

We now apply Theorem 24.1.47 to the convolution of a distribution with a test function.

**Corollary 24.1.39.** For any \( T \in \mathcal{D}(\mathbb{R}^n) \) and \( \psi \in C_0^\infty(\mathbb{R}^n) \), we have \( T \ast \psi \in C^\infty(\mathbb{R}^n) \).

**Proof.** If \( \psi \in C_0^\infty(\mathbb{R}^n) \), then the function \( \varphi \in C^\infty(\mathbb{R}^n \times \mathbb{R}^n) \) given by \( \varphi(x, y) = \psi(x - y) \) satisfies the assumptions of Theorem 24.1.47. Therefore \( T \ast \psi \) is smooth. \( \square \)

Our next aim is to show that if \( T \in \mathcal{D}'(\Omega) \) and \( (\psi_k)_k \) is a (model) delta sequence, then \( T \ast \psi_k \to T \) in \( \mathcal{D}' \). This shows that every distribution can be approximated by smooth functions. By definition of \( \mathcal{D}' \)-convergence, we have to show that

\[
\forall u \in \mathcal{D}(\Omega) : \langle T \ast \psi_k, u \rangle \to \langle T, u \rangle
\]

The idea is simply to “move the convolution to the other side” and use that \( u \ast \psi_k \to u \) in \( \mathcal{D}(\Omega) \) by Theorem 24.1.37. To this end, a sort of generalization of Fubini’s theorem is required. To see this, consider first the case that \( T = T_f \) is a regular distribution. Then using Fubini’s Theorem one has

\[
\langle T_f \ast \psi, u \rangle = \int_{\mathbb{R}^n} (T_f \ast \psi)(x)u(x) \, dx = \int_{\mathbb{R}^n} \left( \int_{\mathbb{R}^n} f(y)\psi(x - y) \, dy \right) u(x) \, dx
\]

Fubini

\[
= \int_{\mathbb{R}^n} f(y) \left( \int_{\mathbb{R}^n} u(x)\psi(x - y) \, dx \right) \, dy
\]

where in the last line we used the notation \( \tilde{\psi}(x) := \psi(-x) \) for the “reflected function” (note the order of variables in the inner integral in the third line). We now show that this still holds true if \( u \) is a distribution.
**Proposition 24.1.40.** Let \( T \in \mathcal{D}'(\mathbb{R}^n) \) and \( \varphi \in \mathcal{D}(\mathbb{R}^n) \). Then for all \( \psi \in \mathcal{D}(\mathbb{R}^n) \) we have

\[
\langle T * \varphi, \psi \rangle = \langle T, \hat{\varphi} * \psi \rangle,
\]

where \( \hat{\varphi} \in \mathcal{D}(\mathbb{R}^n) \) denotes the “reflected function” with \( \hat{\varphi}(x) = \varphi(-x) \).

**Proof.** Choose \( R > 0 \) sufficiently large such that \( \text{supp} \, \psi \subset [-R, R]^n \). Then \( (T * \varphi) \psi \in C_0^\infty([-R, R]^n) \). Making a Riemann sum approximation we have **Have to say more precisely how** \([R, R]^n\) **is cut apart and what the** \( x_j \) **are**

\[
\langle T * \varphi, \psi \rangle = \int_{[-R,R]^n} (T * \varphi)(x) \psi(x) \, dx \\
= \int_{[-R,R]^n} (T(y), \varphi(x-y)) \psi(x) \, dx \\
= \lim_{N \to \infty} \sum_{j=1}^N (T(y), \varphi(x_j-y)) \psi(x_j) \left( \frac{2R}{N} \right)^n \\
= \lim_{N \to \infty} (T(y), \sum_{j=1}^N \hat{\varphi}(y-x_j) \psi(x_j) \left( \frac{2R}{N} \right)^n).
\]

The sum in the second factor converges to \( \hat{\varphi} * \psi \) (pointwise, i.e., for fixed \( y \)). If we can argue that we can take the limit inside the second factor, then we are done. For this we need to show that the sum converges in \( \mathcal{D} \) (cf. Definition 24.1.1 for \( \mathcal{D} \)-convergence). To this end, note first that for any \( y \in \mathbb{R}^n \) the support of the function \( x \mapsto \hat{\varphi}(y-x) \psi(x) \) is contained in the fixed compact set \( \text{supp} \, \varphi + \text{supp} \, \psi \). Moreover, since \( x \mapsto \hat{\varphi}(y-x) \psi(x) \) and also all of its derivatives are continuous and compactly supported, all of them are actually uniformly continuous functions. Using this one can show that the Riemann sum converges uniformly to \( \hat{\varphi} * \psi \), and equally for all derivatives. We leave the details as an exercise. \( \square \)

It now follows immediately that any distribution can be approximated by a sequence of smooth functions, even compactly supported ones by using a simple cutoff trick.

**Corollary 24.1.41 (Distributions can be approximated by smooth functions).** \( \mathcal{D}(\mathbb{R}^n) \subset \mathcal{D}'(\mathbb{R}^n) \) is dense, i.e., for any \( T \in \mathcal{D}'(\mathbb{R}^n) \) there exists a sequence \( (f_k)_k \subset \mathcal{D}(\mathbb{R}^n) \) such that \( T_{f_k} \to T \) in \( \mathcal{D}' \).

**Proof.** Let \( T \in \mathcal{D}'(\mathbb{R}^n) \). Choose a strict delta sequence \( (\psi_k)_k \) and a bump function \( \chi \in C_0^\infty(\mathbb{R}^n) \) with \( \chi|_{B_1(0)} = 1 \). Set

\[
T_k := \left( \chi \left( \frac{x}{k} \right) T \right) * \psi_k \in \mathcal{D}(\mathbb{R}^n),
\]

then for any \( \varphi \in \mathcal{D}(\mathbb{R}^n) \) we have

\[
\langle T_k, \varphi \rangle = \langle \chi \left( \frac{x}{k} \right) T, \psi_k * \varphi \rangle = \langle T, \chi \left( \frac{x}{k} \right) (\psi_k * \varphi) \rangle.
\]

It is easy to verify that \( \chi \left( \frac{x}{k} \right) (\psi_k * \varphi) \to \varphi \) in \( \mathcal{D} \), so by continuity of \( T \) the right-hand side converges to \( \langle T, \varphi \rangle \). This shows that \( T * \psi_k \to T \) in \( \mathcal{D}'(\mathbb{R}^n) \) as desired. \( \square \)

**The following can probably also be left out.** So far we have only defined the convolution of a distribution with a smooth function. Building on this it is not difficult to define also the convolution of two distributions. The definition for this can again be motivated by regarding smooth functions as (regular) distributions. So let \( f \in C^\infty(\mathbb{R}^n) \) and \( g \in C_0^\infty(\mathbb{R}^n) \), and
let $\varphi \in \mathcal{D}(\mathbb{R}^n)$ be a test function. Then the action of the distribution $T_f \ast g$ on $\varphi$ can be rewritten using Fubinis theorem as

$$
\langle T_f \ast g, \varphi \rangle = \int_{\mathbb{R}^n} (f \ast g)(x) \varphi(x) \, dx
$$

Here we again used the “check-notation” for the reflection of a function. With the last line we have almost arrived at an expression where the regular distributions $T_f$ and $T_g$ can be replaced by general distributions, we only have to explain first how to extend the “check-operation” to distributions. We leave this as a simple exercise to the reader.

**Definition 24.1.42 (Convolution of two distributions).** For $S \in \mathcal{D}'(\mathbb{R}^n)$ and $T \in \mathcal{E}'(\mathbb{R}^n)$, the convolution of $S$ with $T$ is the distribution $S \ast T \in \mathcal{D}'(\mathbb{R}^n)$ defined by

$$
\langle S \ast T, \varphi \rangle := \langle S, \check{T} \ast \varphi \rangle \quad \forall \varphi \in \mathcal{D}(\mathbb{R}^n).
$$

(24.1.38)

Here the distribution $\check{T}$ is defined by $\langle \check{T}, \varphi \rangle := \langle T, \check{\varphi} \rangle$ for all $\varphi \in \mathcal{D}(\mathbb{R}^n)$, cf. Exercise 27.20.

The verification that (24.1.38) indeed defines a distribution is left as an exercise (see Exercise 27.21).

Using convolution of two distributions one can show an interesting structural result about distributions, namely that locally, any distribution is the (distributional) derivative of a continuous function.

**Theorem 24.1.43 (Local structure of distributions).** Let $T \in \mathcal{D}'(\mathbb{R}^n)$. Then for any bounded domain $\Omega \subset \mathbb{R}^n$ there exists $f \in C(\mathbb{R}^n)$ and $\alpha \in \mathbb{N}^n$ such that

$$
T|_{\Omega} = \partial^\alpha f|_{\Omega}.
$$

(24.1.39)

Here the right-hand side has to be understood in the distributional sense.

**Proof.** The idea is to write $u = u \ast \delta$ and then write $\delta = \partial^\alpha E$ with some continuous $E$ such that $u \ast E$ is continuous. \qed

**Corollary 24.1.44 (Global structure of compactly supported distributions).** Let $u \in \mathcal{E}'(\mathbb{R}^n)$, and let $\Omega \subset \mathbb{R}^n$ be an open neighborhood of $\text{supp } u$. Then there exists $m \in \mathbb{N}$ and functions $f_\alpha \in C_0(\Omega)$ for all $\alpha \in \mathbb{N}^n$ with $|\alpha| \leq m$ such that

$$
u = \sum_{|\alpha| \leq m} \partial^\alpha f_\alpha.
$$

(24.1.40)

The right-hand side again has to be understood in the distributional sense.

- convolution of two distributions
- some examples, or nice formulae
- maybe solve some evolution equations by convolution as motivation

The following only makes sense if the convolution of two distributions has been defined before. We finish this section by showing that the convolution theorem (24.1.25)
about the Fourier transform of a convolution still holds for the convolution of a tempered distribution and a compactly supported one. Note that without the compact support assumption the multiplication of the two Fourier transforms wouldn’t even be defined.

**Theorem 24.1.45.** Let \( u \in S'(\mathbb{R}^n) \) and \( v \in E'(\mathbb{R}^n) \). Then \( u * v \in S'(\mathbb{R}^n) \) and

\[
\hat{u} * \hat{v} = \hat{u} \cdot \hat{v}
\]

in the sense of tempered distributions.

**Proof.** To be done! \( \square \)

Compact support seems like a somewhat brutal decay assumption, and one can indeed relax it. For instance, the convolution theorem also holds for \( u \in S'(\mathbb{R}^n) \) and \( \varphi \in S(\mathbb{R}^n) \subset S'(\mathbb{R}^n) \). In this case one can show that \( u * \varphi \in \mathcal{O}_M(\mathbb{R}^n) \subset S'(\mathbb{R}^n) \) and \( \hat{u} * \hat{\varphi} = \hat{u} \cdot \hat{\varphi} \) again holds in the sense of tempered distributions.

The convolution theorem is important for instance because one can use it to produce solution formulas for certain (linear) evolution equations. We illustrate this by one example and give some further exercises for you to work out.

**Example 24.1.46.** We use the convolution theorem to solve the Cauchy problem for the free Schrödinger equation, i.e.

\[
\begin{aligned}
\mathcal{I} \partial_t \psi(t, x) &= -\Delta \psi(t, x) \\
\psi(0, x) &= \psi_0(x)
\end{aligned}
\]

Hereby we assume that \( \psi_0 \in S(\mathbb{R}^n) \). Moreover, we use that . . . Taking the Fourier transform in \( x \), we obtain the equation

\[
\mathcal{I} \partial_t \hat{\psi}(t, \xi) = |\xi|^2 \hat{\psi}(t, \xi).
\]

This implies that

\[
\hat{\psi}(t, \xi) = e^{-|\xi|^2 t} \hat{\psi}_0(\xi).
\]

Now notice that \( \hat{\psi}_0 \in S(\mathbb{R}^n) \), whereas the smooth and bounded function \( e^{-|\xi|^2 t} \) may be viewed as tempered distribution. Therefore it follows from the convolution theorem that

\[
\psi(t, x) = \mathcal{F}_x^{-1} \left( e^{-|\xi|^2 t} \right) * \psi_0(x) = \frac{1}{(4\pi \mathcal{I} t)^n} \int_{\mathbb{R}^n} e^{-\frac{|x-y|^2}{4\mathcal{I} t}} \psi_0(y) \, dy.
\]  

(24.1.42)

Here we used that the inverse Fourier transform of the complex Gaussian \( \exp(-\mathcal{I} |\xi|^2 t) \) is the complex Gaussian \( (4\pi \mathcal{I} t)^{-n} \exp(-\mathcal{I} |\cdot|^2 t) \).

**24.1.8. The technical toolbox.**

**Theorem 24.1.47 (”Differentiation under the integral”).** Let \( \Omega_1 \subset \mathbb{R}^n \) and \( \Omega_2 \subset \mathbb{R}^n \) be open and \( u \in C^\infty(\Omega_1 \times \Omega_2) \). Suppose that \( u \) is compactly supported in the first variable with “locally uniformly supports in the second variable”, i.e., suppose that for every \( y_0 \in \Omega_2 \) there exists an open neighborhood \( U \subset \Omega_2 \) of \( y_0 \) and a compact set \( K \subset \Omega_1 \) such that

\[
\forall y \in U : \quad \text{supp} \, u(\cdot, y) \subset K.
\]  

(24.1.43)

Then for every distribution \( T \in \mathcal{D}'(\Omega_1) \), the function

\[
\Omega_2 \ni y \mapsto \langle T, u(\cdot, y) \rangle \in \mathbb{C}
\]

is smooth, and for every \( \alpha \in \mathbb{N}^n \) one has

\[
\partial_\alpha^y \langle T, u(\cdot, y) \rangle = \langle T, \partial_\alpha^y u(\cdot, y) \rangle.
\]  

(24.1.45)

**Proof.** One can explicitly refer to difference quotients and estimate everything appropriately. The mean value theorem is often useful. \( \square \)

\footnote{Note that \( S(\mathbb{R}^n) \not\subset E'(\mathbb{R}^n) \).}
In the previous theorem, the support assumption (24.1.43) is automatically satisfied if the test function \( u \) has compact support. If on the other hand the distribution \( T \) has compact support, one can in fact drop the support assumption on the test function altogether. You may prove this in exercise 27.100.

**Proposition 24.1.48.** Let \( \Omega_1 \subset \mathbb{R}^n \) and \( \Omega_2 \subset \mathbb{R}^m \) be open. Then

\[
\text{span}\{f \otimes g \mid f \in \mathcal{D}(\Omega_1), g \in \mathcal{D}(\Omega_2)\} \subset \mathcal{D}(\Omega_1 \times \Omega_2)
\]

is dense.

**Theorem 24.1.49 (Tensor product of distributions).** Let \( \Omega_1 \subset \mathbb{R}^n \) and \( \Omega_2 \subset \mathbb{R}^m \) be open. For any \( S \in \mathcal{D}'(\Omega_1) \) and \( T \in \mathcal{D}'(\Omega_2) \) there exists a unique \( S \otimes T \in \mathcal{D}'(\Omega_1 \times \Omega_2) \) such that

\[
\langle S \otimes T, u \otimes v \rangle = \langle S, u \rangle \langle T, v \rangle
\]

(24.1.46)

holds for all \( u \in \mathcal{D}(\Omega_1) \) and \( v \in \mathcal{D}(\Omega_2) \). Moreover, the following properties hold.

1. For any \( u \in \mathcal{D}(\Omega_1 \times \Omega_2) \) we have “Fubini’s theorem”:
   \[
   \langle S \otimes T, u \rangle = \langle S(x), \langle T(y), u(x,y) \rangle \rangle = \langle \langle S(x), u(x,y) \rangle, T(y) \rangle
   \]

(24.1.47)

2. \( \text{supp}(S \otimes T) = \text{supp}(S) \times \text{supp}(T) \)
3. \( \partial_x^a \partial_y^b (S \otimes T) = \partial_x^a S \otimes (\partial_y^b T) \)
4. The map \( (S, T) \mapsto S \otimes T \) is bilinear and (jointly) sequentially continuous.

**24.1.9. Further topics about distributions.** There is some material here which I commented out.

### 24.2. The Spectral Theorem for Self-Adjoint Operators

The main objective of this chapter is to explain to you the so-called spectral theorem for selfadjoint operators on Hilbert spaces.

**Theorem 24.2.1 (Spectral theorem).** Let \( \mathcal{H} \) be a complex Hilbert space and let \( A \in L(\mathcal{H}) \) be a (bounded) selfadjoint linear operator on \( \mathcal{H} \). Then there exists a unique spectral measure \( E : \mathcal{B}(\sigma(A)) \to L(\mathcal{H}) \) on the Borel \( \sigma \)-algebra over the spectrum of \( A \) such that

\[
A = \int_{\sigma(A)} \lambda dE(\lambda).
\]

(24.2.1)

Moreover, for any bounded measurable function \( f : \sigma(A) \to \mathbb{C} \) integration against \( E \) defines a bounded operator \( f(A) \in L(\mathcal{H}) \) via \( f(A) := \int_{\sigma(A)} f(\lambda) dE(\lambda) \). These operators satisfy the following properties.

This important result is a generalization of the well-known fact that a Hermitian matrix may be diagonalized to the case of infinite-dimensional (Hilbert) spaces. All the notions appearing in this theorem will be explained in the following sections.

**24.2.1. The finite-dimensional case.** Let \( n \in \mathbb{N} \) and \( A \in \mathbb{C}^{n \times n} \) be given, and assume that \( A \) is hermitian, i.e., \( A^* = A \), where \( A^* = \overline{A}^t \) is the hermitian transpose, i.e., complex conjugate and transpose of \( A \). The following facts should be familiar from linear algebra:

- All eigenvalues of \( A \) are real numbers,
- eigenvectors belonging to different eigenvalues are always orthogonal,
- there exists an orthonormal basis of \( \mathbb{C}^n \) consisting of eigenvectors of \( A \).

Counting with multiplicities, we enumerate the eigenvalues of \( A \) by \( \lambda_1, \ldots, \lambda_n \in \mathbb{R} \) and choose a corresponding orthonormal basis of eigenvectors \( v_1, \ldots, v_n \in \mathbb{C}^n \). If we denote by \( U \in \mathbb{C}^{n \times n} \) the matrix whose columns are precisely \( v_1, \ldots, v_n \), then \( U \) is unitary and the matrix \( U^*AU \) is the diagonal matrix with entries \( \lambda_1, \ldots, \lambda_n \) on the diagonal.
24.2. THE SPECTRAL THEOREM FOR SELF-ADJOINT OPERATORS

We now reformulate this in such a way that it becomes apparent why this is precisely the same statement as the spectral theorem from above. To this end, we write \( A = UDU^* \) and symbolically compute with matrices as follows:

\[
A = UDU^* = \begin{pmatrix} v_1 & \cdots & v_n \end{pmatrix} \begin{pmatrix} \lambda_1 & & \\ & \ddots & \\ & & \lambda_n \end{pmatrix} \begin{pmatrix} v_1^* \\ \vdots \\ v_n^* \end{pmatrix} = \lambda_1 v_1 v_1^* + \cdots + \lambda_n v_n v_n^* .
\]

(24.2.2)

To make the connection to the spectral theorem we observe that this is exactly the same as \( \sigma \) if we define the spectral measure \( E : \mathcal{B}(\sigma(A)) \to \mathbb{C}^{n \times n} \) by

\[
E(\Omega) := \sum_{i=1}^n v_i v_i^* .
\]

(24.2.3)

Here the sum is over all eigenvalues (with multiplicities) lying in \( \Omega \). In other words, \( E(\Omega) \) is the sum of the projection on all those eigenspaces whose eigenvalue lies in the set \( \Omega \). We leave it as an exercise to write out the integral \( \int_\Omega \) for this spectral measure concretely (one should first read the part about spectral measures).

24.2.3. Hilbert spaces.

24.2.3.1. Bounded operators and their spectra.

**Definition 24.2.2.** Let \( \mathcal{H} \) be a Hilbert space. A bounded operator on \( \mathcal{H} \) is a linear continuous map from \( \mathcal{H} \) to \( \mathcal{H} \). The set of all bounded operators on \( \mathcal{H} \) is denoted by \( L(\mathcal{H}) \).

Since (pointwise) sums and multiples of linear, continuous maps are again linear and continuous again, the set \( L(\mathcal{H}) \) is in fact a vector space with respect to (pointwise) addition and scaling. The terminology “bounded” is due to the fact that a linear map \( A : \mathcal{H} \to \mathcal{H} \) is continuous if and only there exists a constant \( C > 0 \) such that

\[
\forall u \in \mathcal{H} : \quad \|Au\| \leq C \cdot \|u\| .
\]

(24.2.4)

You are asked to show this in exercise 27.32. Assigning to each bounded operator the optimal constant \( C > 0 \) for which \( \|A\| \) holds yields a norm on \( L(\mathcal{H}) \) which turns this space into a Banach space.

**Theorem 24.2.3 (Operator norm).** For \( A \in L(\mathcal{H}) \), define

\[
\|A\|_\infty := \sup_{x \in \mathcal{H}, \|x\|=1} \|Ax\| .
\]

(24.2.5)

Then \( \|A\|_\infty < \infty \), and this number is the smallest possible constant \( C \geq 0 \) such that \( \|A\| \infty \) holds. Moreover, \( L(\mathcal{H}) \), \( L(\mathcal{H}), \|\cdot\|_\infty \) is a Banach space.

The proof of this result is straight-forward and you are asked for it in Exercise 27.33.

Now we turn towards the spectrum of a bounded operator which generalizes the set of all eigenvalues of a matrix.

**Definition 24.2.4.** Let \( A \in L(\mathcal{H}) \) be given. Set

\[
\rho(A) := \{ \lambda \in \mathbb{C} : \lambda \mathbb{I}_\mathcal{H} - A \text{ is invertible} \} ,
\]

(24.2.6)

\[
\sigma(A) := \mathbb{C} \setminus \rho(A) .
\]

(24.2.7)

One calls \( \rho(A) \) the resolvent set and \( \sigma(A) \) the spectrum of \( A \).

Let us make some remarks about the spectrum of an operator. For given \( A \in L(\mathcal{H}) \) and \( \lambda \in \mathbb{C} \), there are a priori two different reasons why one might have \( \lambda \in \sigma(A) \):

\[\footnote{In BraKet-notation the last line of this calculation would read \( A = \lambda_1|v_1><v_1| + \cdots + \lambda_n|v_n><v_n| \).} \]
The map $\lambda \text{id}_H - A$ is not injective. Due to linearity, in this case there exists $x \in H$ with $(\lambda \text{id}_H - A)x = 0$, i.e., $Ax = \lambda x$. Hence $\lambda$ is an eigenvalue of $A$.

The map $\lambda \text{id}_H - A$ is injective but not surjective. In this case, $\lambda$ is not an eigenvalue of $A$.

For $\dim H < \infty$ the second case is excluded by the rank-nullity-theorem. For $\dim H = \infty$ the second case is possible, however, so $\sigma(A)$ may consist of more than just eigenvalues.

Let us remark that as a consequence of the so-called open mapping theorem (see ??), the inverse of a bounded operator on a Hilbert space is automatically bounded again. Hence one has $(\lambda \text{id}_H - A)^{-1} \in L(H)$ for any $\lambda \in \rho(A)$. One calls the family of operators $R(\lambda) := (\lambda \text{id}_H - A)^{-1}$ with $\lambda \in \rho(A)$ the resolvent of $A$. The resolvent plays an important role in the proof of some of the following properties of the spectrum of a bounded operator, most notably the fact that by a geometric series argument it can be expressed as a power series in $\lambda$ around any point of $\rho(A)$.

**Theorem 24.2.5** (Basic facts about the spectrum). Let $A \in L(H)$ be given. The following properties are satisfied:

1. The resolvent set $\rho(A) \subseteq \mathbb{C}$ is open.
2. The spectrum $\sigma(A) \subseteq \mathbb{C}$ is compact and $|\lambda| \leq \|A\|_\infty$ for any $\lambda \in \sigma(T)$.
3. If $H$ is a complex Hilbert space, then $\sigma(A) \neq \emptyset$.

**Proof.** Let $\lambda_0 \in \rho(A)$ be given. For $\lambda \in \mathbb{C}$ we have ("common denominator")

$$\text{id}_H - (\lambda_0 - \lambda) \cdot (\lambda_0 \text{id}_H - A)^{-1} = (\text{id}_H - A) \cdot (\lambda_0 \text{id}_H - A)^{-1},$$

and thus

$$\lambda \text{id}_H - A = (\text{id}_H - (\lambda - \lambda_0)(\lambda_0 \text{id}_H - A)^{-1}) \cdot (\lambda_0 \text{id}_H - A).$$

If $|\lambda - \lambda_0| < \|((\lambda_0 \text{id}_H - A)^{-1})\|$, then the Neumann series (cf. Exercise 27.35) provides the inverse to the first factor on the right-hand side. Hence for $|\lambda - \lambda_0| < \|((\lambda_0 \text{id}_H - A)^{-1})\|$ we have $\lambda \in \rho(A)$ as well. This shows the first part, and of course also that $\sigma(A) \subseteq \mathbb{C}$ is open. By a similar argument one can show that for any $\lambda \in \mathbb{C}$, if $|\lambda| > \|A\|_1$, then $\lambda \in \rho(A)$. This concludes the proof of the second part. For a proof of the third part, see ??.

**Theorem 24.2.6.** Let $T \in L(H)$ be given. Then $T$ is invertible if and only if $T(H) \subseteq H$ is dense and there exists $\epsilon > 0$ such that $\|Tu\| \geq \epsilon \cdot \|u\|$ holds for all $u \in H$.

**Proof.** If $T$ is invertible, then clearly its range is dense in $H$. Setting $\epsilon := \|A^{-1}\|_1^{-1}$ it follows from boundedness that for any $u \in H$ one has

$$\|u\| = \|A^{-1}Au\| \leq \|A^{-1}\| \cdot \|Au\| \quad \Rightarrow \quad \epsilon \cdot \|u\| \leq \|Au\|.

Conversely, we first show that $T$ is bijective. Injectivity follows immediately from the given estimate. Concerning surjectivity, it suffices to show that the range of $T$ is closed. To this end, let $w_n \in \text{Ran}(T_n)$ be a sequence with $w_n \to w \in H$. Choose $u_n \in H$ with $w_n = Tu_n$. Note that by the give estimate

$$\|u_n - u_m\| \leq \frac{1}{\epsilon} \|T u_n - T u_m\| = \frac{1}{\epsilon} \|w_n - w_m\|.

Since $(w_n)_n$ converges, it is a Cauchy sequence. The previous estimate shows that then also $(u_n)_n$ is a Cauchy sequence, hence $u_n \to u$ for some $u \in H$. By continuity of $T$ it follows that $Tu = \lim_{n \to \infty} T u_n = \lim_{n \to \infty} w_n = w$, so $w \in \text{Ran}(T)$. Continuity of the inverse map is also an immediate consequence of the given estimate.

An extremely useful concept for linear operators on Hilbert spaces is the notion of the adjoint of an operator.

**Theorem 24.2.7** (The adjoint of an operator). Let $A \in L(H)$ be given. There exists a unique operator $A^* \in L(H)$ such that

$$\forall u, v \in H: \quad \langle Au, v \rangle = \langle u, A^* v \rangle .$$

(24.2.8)
This operator is called the adjoint operator of \( A \) or, in short, the adjoint of \( A \).

Proof. For fixed \( v \in \mathcal{H} \), by Cauchy-Schwarz the map \( \mathcal{H} \ni u \mapsto \langle Au, v \rangle \in \mathbb{C} \) is linear and bounded. Therefore, by the Riesz representation theorem \(^{24.2.3} \) there exists a unique vector \( A^*v \in \mathcal{H} \) such that \( \langle Au, v \rangle = \langle u, A^*v \rangle \) for any \( u \in \mathcal{H} \). As you may check, the map \( v \mapsto A^*v \) defines a bounded operator on \( \mathcal{H} \). \( \square \)

Exercise \(^{24.2.18} \) shows that taking the adjoint of an operator has certain similarities to taking the complex conjugate of a complex number.

Definition 24.2.8. Let \( A \in L(\mathcal{H}) \).

1. \( A \) is called selfadjoint if \( A = A^* \).
2. \( A \) is called unitary if \( A^*A = \text{id}_\mathcal{H} \) and \( AA^* = \text{id}_\mathcal{H} \).
3. \( A \) is called normal if \( AA^* = A^*A \).

Note that any selfadjoint and any unitary operator is a normal operator. Note also that any unitary operator \( A \) is invertible and \( A^{-1} = A^* \).

Normal operators are precisely those for which the spectral theorem holds. Using the following result it will suffice to focus on selfadjoint operators first.

Proposition 24.2.9. An operator \( A \in L(\mathcal{H}) \) is normal if and only if the (selfadjoint) operators \( S, T \in L(\mathcal{H}) \) defined by

\[
S := \frac{1}{2}(A + A^*), \quad T := \frac{1}{2i}(A - A^*)
\]

commute. One always has \( A = S + iT \).

Being selfadjoint, unitary or normal has certain consequences for the spectrum of an operator. The “algebraic” consequences are probably familiar from Linear Algebra.

Theorem 24.2.10 (Spectral properties of selfadjoint, unitary and normal operators).

Let \( A \in L(\mathcal{H}) \) be given.

1. If \( A \) is selfadjoint, then \( \sigma(A) \subseteq \mathbb{R} \).
2. If \( A \) is unitary, then \( \sigma(A) \subseteq S^1 := \{ z \in \mathbb{C} \mid |z| = 1 \} \).
3. If \( A \) is normal, then eigenvectors corresponding to different eigenvalues are orthogonal.
4. If \( A \) is normal, then the norm of \( A \) is equal to the spectral radius, i.e., we have \( \|A\| = \sup\{ |\lambda| : \lambda \in \sigma(A) \} \). Moreover, if \( \mathcal{H} \) is a complex Hilbert space then there exists \( \lambda \in \sigma(A) \) such that \( |\lambda| = \|A\| \).

24.2.4. Spectral measures. Having covered the basics about operators, we now discuss the integration part of the spectral theorem, i.e., what a spectral measure is and how one may integrate with respect to such a measure. In general, all the notions are very similar as in ordinary (Lebesgue) measure and integration theory.

Definition 24.2.11 (Spectral measure). Let \( \Omega \) be a set and \( \mathcal{A} \subseteq \mathcal{P}(\Omega) \) a \( \sigma \)-algebra on \( \Omega \). A spectral measure or also projection-valued measure on \( (\Omega, \mathcal{A}) \) is a map \( E : \mathcal{A} \rightarrow L(\mathcal{H}) \) satisfying the following properties:

1. \( E(A) \) is an orthogonal (selfadjoint) projection for every \( A \in \mathcal{A} \),
2. \( E(\Omega) = \text{id}_\mathcal{H} \),
3. If \( A_1, A_2, \ldots \in \mathcal{A} \) are pairwise disjoint, and \( A := \bigcup_{j \in \mathbb{N}} A_j \), then it holds that

\[
\forall u \in \mathcal{H} : \quad E(A)u = \sum_{j=1}^{\infty} E(A_j)u .
\]

Thus a spectral measure is just like an ordinary (positive) measure except for the fact that its values are orthogonal projections of some Hilbert space. There is also a very simple way in which one can extract “ordinary”, i.e., real- or complex-valued measures from a given spectral
measure \( E \). Namely, one may form “expectation-values” and “transition probabilities”, i.e., for given \( u,v \in \mathcal{H} \) one defines

\[
E_u : \mathcal{A} \to [0, \infty), \quad A \mapsto E_u(A) := \langle E(A)u, u \rangle, \tag{24.2.11}
\]

\[
E_{u,v} : \mathcal{A} \to \mathbb{C}, \quad A \mapsto E_{u,v}(A) := \langle E(A)u, v \rangle. \tag{24.2.12}
\]

In Exercise 24.21 you are asked to show that these functions are (positive, respectively complex) measures. Since \( E(A) \) is an orthogonal projection, for any \( u \in \mathcal{H} \) one has

\[
E_u(A) = \langle E(A)u, u \rangle = \langle E(A)u, E(A)u \rangle = \| E(A)u \|^2.
\]

The fact that the values of a spectral measure are orthogonal projections has certain very useful consequences.

**Theorem 24.2.12 (Properties of spectral measures).** Let \( E : \mathcal{A} \to L(\mathcal{H}) \) be a spectral measure, and let \( A, B \in \mathcal{A} \) as well as \( A_1, A_2, \ldots \in \mathcal{A} \) be given. Then the following hold:

1. \( E(A \cap B) = E(A)E(B) = E(B)E(A) \)
2. \( A \cap B = \emptyset \), then \( E(A)E(B) = 0 \) and therefore \( \text{Ran}(E(A)) \perp \text{Ran}(E(B)) \).
3. If \( A \subseteq B \), then \( E(A)E(B) = E(A) \) and therefore \( \text{Ran}(E(A)) \subseteq \text{Ran}(E(B)) \).
4. If \( A_1 \subseteq A_2 \subseteq \ldots \subseteq A_3 \subseteq \ldots \) and \( A := \bigcup_{n=1}^{\infty} A_j \), then \( E(A_n)u \to E(A)u \) for all \( u \in \mathcal{H} \).
5. If \( A_1 \supseteq A_2 \supseteq \ldots \supseteq A_3 \supseteq \ldots \) and \( A := \bigcap_{n=1}^{\infty} A_j \), then \( E(A_n)u \to E(A)u \) for all \( u \in \mathcal{H} \).

Now we come to integration with respect to a given spectral measure. This is done via the usual “Lebesgue-approach” starting from simple functions. In the following, let \( E : \mathcal{A} \to L(\mathcal{H}) \) denote a given spectral measure on a measurable space \( \Omega \).

**Step 1: Integration of Simple Functions**

Recall that a **simple function** is a measurable function \( f : \Omega \to \mathbb{C} \) which only takes finitely many distinct values. Such a function \( f \) can always be written as

\[
f = \sum_{j=1}^{n} a_j 1_{A_j}, \tag{24.2.13}
\]

for certain \( a_1, \ldots, a_n \in \mathbb{C} \) and \( A_1, \ldots, A_n \in \mathcal{A} \). For instance, one may take \( a_1, \ldots, a_n \) to be the distinct values of \( f \) and set \( A_1 := f^{-1}(\{a_1\}), \ldots, A_n := f^{-1}(\{a_n\}) \). For this choice the sets \( A_1, \ldots, A_n \) are also disjoint. We now define the **integral of \( f \) over \( \Omega \)** as

\[
T_E(f) := \int_{\Omega} f(x) \, dE(x) := \sum_{j=1}^{n} a_j E(A_j) \in L(\mathcal{H}), \tag{24.2.14}
\]

where \( f = \sum_{j=1}^{n} a_j 1_{A_j} \) is any such representation of \( f \) (the outcome in (24.2.14) is independent of this choice). Note that the integral (24.2.14) yields an **operator**.

**Step 2: Integration of Bounded, Measurable Functions**

The extension of the integral to more general functions will be done by means of approximation, using the fact that every bounded, measurable function can be uniformly approximated by a sequence of simple functions. We start with some terminology.

**Definition 24.2.13.** Let \((\Omega, \mathcal{A})\) be a measurable space and \( E : \mathcal{A} \to L(\mathcal{H}) \) a spectral measure. For a bounded, measurable function \( f : \Omega \to \mathbb{C} \) one calls

\[
\text{ess sup} \{ f \} := \inf \left\{ M \geq 0 : E(\{ x \in \Omega : |f(x)| > M \}) = 0 \right\}, \tag{24.2.15}
\]

the **essential supremum** of \( f \) (with respect to \( E \)). We denote the **space of bounded, measurable functions** by \( L^\infty(\Omega) \), and set \( L^\infty(\Omega) := L^\infty(\Omega) / \{ f \in L^\infty(\Omega) | \text{ess sup} \{ f \} = 0 \} \).

As in usual Lebesgue measure theory, the space \( L^\infty(\Omega) \) is a Banach space. Note that it depends on the measure \( E \) through (24.2.13), although this is not depicted in the notation \( L^\infty(\Omega) \). The space \( L^\infty(\Omega) \) does not depend on \( E \), but is not a Banach space since \( \text{ess sup} | \cdot | \) is not a norm on \( L^\infty(\Omega) \).
Lemma 24.2.14. Let \( f : \Omega \to \mathbb{C} \) be a simple function and let \( T_E(f) \in L(H) \) be defined by (24.2.14). Then
\[
\|T_E(f)\|_\infty = \text{ess sup}|f|.
\]

Proof. Write \( f = \sum_{j=1}^{n} a_j \mathbf{1}_{A_j} \) for disjoint \( A_1, \ldots, A_n \). For the sake of this proof we may assume that \( E(A_1), \ldots, E(A_n) \neq 0 \), otherwise we can alter \( f \) in such way that this holds and both \( T_E(f) \) and \( \text{ess sup}|f| \) remain unchanged. Then \( \text{ess sup}|f| = \max\{|a_1|, \ldots, |a_n|\} \), and without loss of generality we may assume that \( \text{ess sup}|f| = |a_1| \).

Since \( E(A_1) \neq 0 \) we may choose \( u \in E(A_1)(H) \) with \( \|u\| = 1 \). Since \( A_1, \ldots, A_n \) are disjoint, \( E(A_1), \ldots, E(A_n) \) project onto pairwise orthogonal subspaces. Therefore
\[
T_E(f)u = \sum_{j=1}^{n} a_j E(A_j)u = a_1 u,
\]
and thus
\[
\|T_E(f)u\|^2 = \langle T_E(f)u, T_E(f)u \rangle = |a_1|^2 \cdot \|u\|^2 = \text{ess sup}|f|^2.
\]
This shows that \( \|T_E(f)\|_\infty \geq \text{ess sup}|f| \).

Concerning the reverse estimate, note that for any \( u \in H \) the vectors \( E(A_1)u, \ldots, E(A_n)u \) are orthonormal. Therefore, setting \( A := \bigcup A_j \) one has
\[
\|T_E(f)u\|^2 = \left\| \sum_{j=1}^{n} a_j E(A_j)u \right\|^2 = \sum_{j=1}^{n} a_j E(A_j)u \sum_{j=1}^{n} a_j E(A_j)u = \sum_{j=1}^{n} |a_j|^2 \|E(A_j)u\|^2 \\
\leq \|f\|_\infty^2 \cdot \sum_{j=1}^{n} |E(A_j)u\|^2 = \|f\|_\infty^2 \cdot \left\| \sum_{j=1}^{n} E(A_j)u \right\|^2 = \|f\|_\infty^2 \cdot \|E(A)u\|^2 \\
\leq \|f\|_\infty^2 \cdot \|u\|^2.
\]
Note for the last step that \( u = E(A)u + E(A^c)u \) is an orthogonal decomposition of \( u \), and hence \( \|u\|^2 = \|E(A)u\|^2 + \|E(A^c)u\|^2 \geq \|E(A)u\|^2 \).

By the obvious inequality \( \text{ess sup}|f| \leq \sup|f| \) which holds for all \( f \in \mathcal{L}_\infty(\Omega) \) we have \( \|T_E(f)\|_\infty \leq \sup|f| \) for all simple functions \( f \in \mathcal{L}_\infty(\Omega) \).

Combining (24.2.16) with the standard fact that any bounded measurable function can be uniformly approximated by simple functions (cf. [Folland, Real Analysis, Theorem 2.10]), it is easy to extend integration from simple functions to bounded measurable functions.

Theorem 24.2.15 (Integration with respect to a projection-valued measures).
Let \( (\Omega, \mathcal{A}) \) be a measurable space and let \( E : \mathcal{A} \to L(H) \) be a projection-valued measure. Then there exists a unique map \( T_E : \mathcal{L}_\infty(\Omega) \to L(H) \), assigning to each bounded measurable function a bounded operator, with the following properties:

1. The assignment \( T_E : \mathcal{L}_\infty(\Omega) \to L(H) \) is an isometric, involutive homomorphism of unital \(*\)-algebras, i.e., for any \( f, g \in \mathcal{L}_\infty(\Omega) \) and \( \alpha, \beta \in \mathbb{C} \) one has
\[
T_E(\alpha f + \beta g) = \alpha T_E(f) + \beta T_E(g),
\]
\[
T_E(fg) = T_E(f)T_E(g),
\]
\[
T_E(\overline{f}) = T_E(f)^*,
\]
\[
T_E(1) = \text{id}_H,
\]
\[
\|T_E(f)\|_\infty = \text{ess sup}|f|.
\]
\[
\begin{aligned}
T_E(\alpha f + \beta g) &= \alpha T_E(f) + \beta T_E(g), \\
T_E(fg) &= T_E(f)T_E(g), \\
T_E(\overline{f}) &= T_E(f)^*, \\
T_E(1) &= \text{id}_H, \\
\|T_E(f)\|_\infty &= \text{ess sup}|f|.
\end{aligned}
\]
(2) For any \( f \in \mathcal{L}^\infty(\Omega) \) and any \( u, v \in \mathcal{H} \) it holds that
\[
\langle T_E(f)u, v \rangle = \int_\Omega f(x) \, dE_{u,v}(x),
\]
\[
\langle T_E(f)u, u \rangle = \int_\Omega f(x) \, dE_u(x),
\]
\[
\|T_E(f)u\|^2 = \int_\Omega |f(x)|^2 \, dE_u(x).
\]

One also writes \( \int_\Omega f(x) \, dE(x) \) for the operator \( T_E(f) \).

**Proof.** We first verify all the claimed properties for simple functions. Then we will define \( T_E(f) \) also for general bounded, measurable functions, and verify the claimed properties also in this case. Uniqueness of the map \( T_E \) is left as exercise.

For simple functions, define \( T_E \) as in \((24.2.14)\). All properties stated in (1) are then easily verified, the last one has already been shown in Lemma \((24.2.14)\). Concerning the properties in (2), we verify the first one exemplary and leave the others as exercises. Given a simple function \( f : \Omega \to \mathbb{C} \) and \( u, v \in \mathcal{H} \), write \( f = \sum_{j=1}^n a_j \mathbb{1}_{A_j} \) and calculate
\[
\langle T_E(f)u, v \rangle = \sum_{j=1}^n a_j \langle E(A_j)u, v \rangle = \sum_{j=1}^n a_j E_{u,v}(A_j)
\]
\[
= \sum_{j=1}^n a_j \int_\Omega \mathbb{1}_{A_j}(x) \, dE_{u,v}(x) = \int_\Omega f(x) \, dE_{u,v}(x),
\]
where the last step uses linearity of integration with respect to a complex measure.

Next we define \( T_E \) also for general bounded, measurable function. To this end, let \( f : \Omega \to \mathbb{C} \) be bounded and measurable. By [Folland, Real Analysis, Theorem 2.10] we can choose a sequence \( (f_n)_n \) of simple functions converging uniformly to \( f \). In particular, we have \( \text{ess sup}|f_n - f| \to 0 \), so by Lemma \((24.2.14)\) it follows that
\[
\|T_E(f_n) - T_E(f_m)\|_\infty = \|T_E(f_n - f_m)\|_\infty = \text{ess sup}|f_n - f_m|.
\]

Since \( \text{ess sup}|f_n - f| \to 0 \), the sequence \( (f_n)_n \) is a Cauchy sequence with respect to \( \text{ess sup}| \cdot | \). By \((*)\), so is the sequence \( (T_E(f_n))_n \). Since \( L(\mathcal{H}) \) is complete, it thus converges to some operator which we denote by \( T_E(f) \).

That all the properties in (1) and (2) remain valid for general bounded, measurable functions can be shown by uniform approximation by step functions. Concerning the integrals in (2), note that the measures \( E_u \) and \( E_{u,v} \) are all finite and a uniformly convergent sequence of functions \( (f_n)_n \) is uniformly bounded in \( n \). This allows to use the dominated convergence theorem. \( \square \)

The following properties of operators obtained by integrating a spectral measure can be easily deduced from the properties of spectral integrals.

**Corollary 24.2.16.** Let \( E : \mathcal{A} \to L(\mathcal{H}) \) be a spectral measure and \( f \in \mathcal{L}^\infty(\Omega) \). Then the following holds:

1. \( T_E(f) \) is normal.
2. \( T_E(f) \) is selfadjoint if and only if \( f \) is real-valued \( (\overline{f} = f) \) almost everywhere.
3. \( T_E(f) \) is unitary if and only if \( |f| = 1 \) almost everywhere.

**Definition 24.2.17.** Let \( X \) be a topological space and \( E : \mathcal{B}(X) \to L(\mathcal{H}) \) a spectral measure on the Borel-\( \sigma \)-algebra of \( X \). Then the support of \( E \) is defined as the complement of the largest open set of measure 0.

Clearly
\[
\int_X f(x) \, dE(x) = \int_{\text{supp}(E)} f(x) \, dE(x).
\]
Therefore one may in fact extend the map $T_E$ further to those measurable functions which are bounded on $\text{supp}(E)$, i.e. $\mathcal{L}^\infty(\text{supp}(E), E)$.

**Theorem 24.2.18.** Let $E: B(\mathbb{C}) \to L(\mathcal{H})$ be a spectral measure on the Borel-$\sigma$-algebra of $\mathbb{C}$. Assume that $\text{supp} E \subseteq \mathbb{C}$ is compact, and set $T := T_E(\text{id})$. Then $\text{supp} E = \sigma(T)$.

**Proof.** First let $\lambda \in \text{supp} E$. Then for any $\epsilon > 0$ we have $E(B_\epsilon(\lambda)) \neq 0$. For any $u \in \text{Ran}(E(B_\epsilon(\lambda))) \neq \{0\}$ we have

$$E_{u,u}(B_\epsilon(\lambda)) = \langle E(B_\epsilon(\lambda))u, u \rangle = \langle u, u \rangle = \|u\|^2 = \|E_{u,u}\|,$$

thus $\text{supp} E_{u,u} \subseteq B_\epsilon(\lambda)$. It follows that

$$\|(T - \lambda)u\|^2 = \int_{B_\epsilon(\lambda)} |z - \lambda|^2 dE_{u,u}(z) \leq \epsilon^2 \cdot \|u\|^2.$$

Since $\epsilon > 0$ was arbitrary, it follows that $(T - \lambda)|_{E(B_\epsilon(\lambda))} = 0$. But then $T - \lambda$ cannot be invertible, hence $\lambda \in \sigma(T)$.

Conversely, let $\lambda \notin \text{supp}(E)$. Since $\text{supp}(E) \subseteq \mathbb{C}$ is closed, there exists $\epsilon > 0$ such that $B_\epsilon(\lambda) \subseteq \text{supp}(E)^c$, hence $E(B_\epsilon(\lambda)) = 0$. So $B_\epsilon(\lambda)$ has $E$-measure zero and therefore the function $f: \mathbb{C} \to \mathbb{C}$ with $f(z) = (z - \lambda)^{-1}$ for $z \neq \lambda$ and $f(\lambda) = 0$ is essentially bounded with respect to $E$. Thus we may form the operator $T_E(f) \in L(\mathcal{H})$. But clearly $T_E(f)$ is an inverse to $T - \lambda = T_E(z - \lambda)$ by multiplicity of $T_E$ and $T_E(1) = \text{id}_\mathcal{H}$. So $\lambda \notin \sigma(T)$. \hfill $\Box$

**24.2.5. Proof of the Spectral Theorem for self-adjoint bounded operators.** In this section we give a proof of the Spectral Theorem [24.2.1] which states that any selfadjoint operator $A \in L(\mathcal{H})$ determines a unique spectral measure $E$ on $\sigma(A)$ such that $A = \int_{\sigma(A)} \lambda dE(\lambda)$.

The basic idea of the followig approach is to use a “duality argument”, meaning that instead of directly constructing a spectral measure we construct the map which is supposed to be the integral with respect to this measure. Such “integrals”, which have to be operators, will be defined as functions of the given operator $T$. If we manage to define this map for a sufficiently large class of functions containing characteristic functions, then the desired measure $E$ may be reconstructed from the integral by the identity $E(A) = \mathds{1}_A(T)(= \int_{\sigma(T)} \mathds{1}_A(x) dE(x))$.

The starting point is to define $f(T)$ when $f$ is a polynomial function in the usual way.

**Definition 24.2.19.** Let $T \in L(\mathcal{H})$ be given. For a polynomial $p(x) = \sum_{j=0}^n a_j x^j$ with coefficients $a_0, \ldots, a_n \in \mathbb{C}$ we define the operator $p(T) \in L(\mathcal{H})$ by $p(T) := \sum_{j=0}^n a_j T^n$.

Next we want to extend this definition to continuous functions defined on $\sigma(T)$ in the case where $T$ is selfadjoint. This will be done by approximating continuous functions by polynomials. The following statement is crucial for this.

**Theorem 24.2.20.** Let $T \in L(\mathcal{H})$ and let $p(x) = \sum_{j=0}^n a_j x^j$ be a complex polynomial. Then $\sigma(p(T)) = p(\sigma(T))$. Moreover, if $T$ is selfadjoint, then also $\|p(T)\|_\infty = \|p(\sigma(T))\|_\infty$.

**Proof.** For a constant polynomial $p = a_0$ both $\sigma(p(T))$ and $p(\sigma(T))$ simply equal $\{a_0\}$. So assume in the following that $p$ is nonconstant.

First let $\lambda \in \sigma(T)$ be given. We need to show that $p(\lambda) \in \sigma(p(T))$, i.e., that $p(\lambda)\text{id} - p(T)$ is not invertible. Using polynomial division we may write $p(\lambda) - p(x) = (\lambda - x)q(x)$ for some other polynomial $q$ since the left-hand side vanishes for $x = \lambda$. It follows that $p(\lambda)\text{id} - p(T) = (\lambda\text{id} - T)q(T)$. Since $\lambda \in \sigma(T)$, the operator $\lambda\text{id} - T$ is not invertible, hence neither can be $p(\lambda)\text{id} - p(T)$. This shows $p(\lambda) \in \sigma(p(T))$.

Conversely, let $\lambda \in \sigma(p(T))$. We need to show that $\lambda \in p(\sigma(T))$. Using the fundamental theorem of algebra, we write $\lambda - p(x) = a(z_1 - x) \cdots (z_n - x)$ for certain $a, z_1, \ldots, z_n \in \mathbb{C}$. It follows that

$$\lambda\text{id} - p(T) = a(z_1\text{id} - T) \cdots (z_n\text{id} - T).$$

Since $\lambda \in \sigma(p(T))$, the left-hand side is non-invertible. Therefore at least one of the factors on the right-hand side is also non-invertible. Without loss $z_1\text{id} - T$ is not invertible, thus $z_1 \in \sigma(T)$.

Further to those measurable functions which are bounded on $\text{supp}(E)$, i.e. $\mathcal{L}^\infty(\text{supp}(E), E)$.
Since on the other hand \( z_1 \) is (by definition) a zero of the polynomial \( \lambda - p(z) \) it follows that 
\[ \lambda = p(z_1) \in p(\sigma(T)). \]

Now assume moreover that \( T \) is selfadjoint. Then \( p(T)^* = \overline{p(T)} \), where \( \overline{p} \) is the polynomial with complex-conjugated coefficients, i.e., \( \overline{p}(x) = \sum_{j=0}^n \overline{a_j} x^j \) (nothing happens to the argument \( x \)). It follows that \( p(T)p(T)^* = |p|^2(T) \), where \( |p|^2 = p\overline{p} \) is again a polynomial (since in \( \overline{p} \) only the coefficients are complex conjugated). Hence
\[
\|p(T)\|_\infty^2 = \|p(T)p(T)^*\|_\infty = \|\|p|^2(T)\|_\infty.
\]

The norm equality \( \|p(T)\|_\infty = \|p\|_{\sigma(T)}\|_\infty \) for \( T \) selfadjoint and \( p \) a polynomial now allows to define \( f(T) \) for continuous functions \( f \) via uniform approximation of \( f \) by polynomials. Such an approximation is possible by the Stone-Weierstraß theorem (note that \( \sigma(T) \subseteq \mathbb{C} \) is compact since \( T \) is bounded).

**Theorem 24.2.21 (Continuous functional calculus).** Let \( T \in L(H) \) be selfadjoint. The map assigning to each complex polynomial \( p \) the operator \( p(T) \) as defined in Definition 24.2.17 has a unique continuous extension to a map \( C(\sigma(T)) \to L(H) \), \( f \mapsto f(A) \) with respect to the supremum norm on \( C(\sigma(T)) \) and the operator norm on \( L(H) \). Moreover, the following properties are satisfied:

1. For all \( f, g \in C(\sigma(T)) \) and \( \alpha, \beta \in \mathbb{C} \) it holds that
   \[
   (\alpha f + \beta g)(T) = \alpha f(T) + \beta f(T),
   \]
   \[
   (fg)(T) = f(T)g(T),
   \]
   \[
   f(T)^* = \overline{f(T)},
   \]
   \[
   [f(T), g(T)] = 0.
   \]

   In particular, the operator \( f(A) \) is selfadjoint if and only if \( f \) is real-valued.

2. For every \( f \in C(\sigma(T)) \) with \( f \geq 0 \) also \( f(T) \geq 0 \).

3. If \( Au = \lambda u \) for \( u \in H \) and \( \lambda \in \mathbb{C} \), then \( f(A)u = f(\lambda)u \) for every \( f \in C(\sigma(T)) \).

4. For every \( f \in C(\sigma(T)) \) one has \( \sigma(f(A)) = f(\sigma(A)) \).

5. For every \( f \in C(\sigma(T)) \) one has \( \|f(A)\| = \|f\|_\infty \).

**Proof.** Let \( f \in C(\sigma(T)) \) be given. Since \( \sigma(T) \subseteq \mathbb{R} \) is compact, by the Stone-Weierstrass theorem (see for instance [Rudin, Principles of Mathematical Analysis, Ch. 7]) there exists a sequence \( (p_n)_n \) of polynomials such that \( \|p_n - f\|_\infty \to 0 \). In particular \( (p_n)_n \) is a Cauchy sequence with respect to \( \|\cdot\|_\infty \) on \( C(\sigma(T)) \). By Theorem 24.2.20 we have
\[
\|p_n(T) - p_m(T)\|_\infty = \|p_n - p_m\|\sigma(T)\|_\infty = \|p_n\|_{\sigma(T)} - p_m\|_{\sigma(T)}\|_\infty.
\]
Therefore also \( (p_n(T))_n \) is a Cauchy sequence in \( L(H) \). Since \( L(H) \) is complete, there exists a limit and we set \( f(T) := \lim_{n \to \infty} p_n(T) \). It is not difficult to check that this limit is independent of the choice of the sequence \( (p_n)_n \) with \( p_n \to f \) in \( C(\sigma(T)) \). Moreover, it follows immediately that \( \|f(T)\|_\infty = \esssup|f| \) holds, which proves (5). This also shows continuity of the map \( f \mapsto f(T) \), from which uniqueness follows as well.

The properties in (1) and (3) are obviously satisfied for polynomial functions, and are verified for other continuous functions by approximation by polynomials. The claim about selfadjointness follows from 24.2.27.

Concerning (2), let \( f \in C(\sigma(T)) \) with \( f \geq 0 \) be given. Define \( g := \sqrt{f} \in C(\sigma(T)) \). Then \( g^2 = f \) and so also \( g(A)^2 = f(A) \) by 24.2.20. It follows that for \( u \in H \) we have
\[
\langle f(A)u, u \rangle = \langle g(A)^2 u, u \rangle = \langle g(A)u, g(A)u \rangle = \|g(A)u\|^2 \geq 0.
\]
Note here that \( g(A) \) is selfadjoint by (1) since \( g \) is real-valued.

It remains to show (4). To this end, let \( f \in C(\sigma(T)) \) be given.

We prove \( \sigma(f(T)) \subseteq f(\sigma(T)) \) by contradiction. Assume that \( \lambda \notin f(\sigma(T)) \). Then we may define \( g \in C(\sigma(T)) \) by \( g(x) := (\lambda - f(x))^{-1} \). By (24.2.22) it follows that \( g(T) \) is an inverse to \( \lambda \text{id} - f(T) \). Hence \( \lambda \notin \sigma(f(T)) \).

To prove \( \sigma(f(T)) \subseteq \sigma(f(T)) \), let \( \lambda \in \sigma(T) \) be given. Suppose that \( \lambda \notin \sigma(f(T)) \). Again define \( g \in C(\sigma(T)) \) by \( g(x) := \lambda - f(x) \), then by assumption \( g(T) = \lambda \text{id} - f(T) \) is invertible. Since \( g(\lambda) = 0 \) it is possible to choose a sequence \( (\varphi_n)_n \) of functions \( \varphi_n \in C(\sigma(T)) \) with \( \| \varphi_n \|_\infty = 1 \) for all \( n \in \mathbb{N} \) but whose supports shrink towards \( \{ \lambda \} \) sufficiently fast such that \( \| \varphi_n g \|_\infty \to 0 \). Using invertibility of \( g(T) \) it now follows that

\[
1 = \| \varphi_n \|_\infty = \| \varphi_n(T) \|_\infty = \| \varphi_n(T) g(T) g(T)^{-1} \|_\infty \leq \| \varphi_n(T) g(T) \|_\infty \cdot \| g(T)^{-1} \|_\infty \to 0.
\]

This is a contradiction, thus \( \lambda \in \sigma(f(T)) \) must hold.

With the continuous functional calculus in hand we can now construct the spectral measure of a selfadjoint operator by using “duality arguments”. Let \( T \in L(\mathcal{H}) \) be selfadjoint. By the previous theorem we have a continuous linear map

\[
C(\sigma(T)) \to L(\mathcal{H}), \quad f \mapsto f(T) .
\]

We want to show that this map is in fact given by integration with respect to a spectral measure (which is to be found). To find this measure we will make use of the Riesz representation theorem which states that any continuous (or positive) linear map from \( L(\mathcal{H}) \) to \( \mathbb{C} \) is given by integration with respect to some complex (or positive) measure. More precisely, the following holds.

**Theorem 24.2.22 (Riesz representation theorem for the dual of \( C(X) \)).**

Let \( X \) be a compact Hausdorff space (e.g., \( X \subseteq \mathbb{R}^n \) compact) and let \( I : C(X) \to \mathbb{C} \) be a linear map.

1. If \( I \) is continuous with respect to the supremum norm on \( C(X) \), then there exists a unique complex measure \( \mu \) on the Borel-\( \sigma \)-algebra of \( X \) such that

\[
\forall f \in C(X) : \quad I(f) = \int_K f(x) \, d\mu(x) . \tag{24.2.29}
\]

2. If \( I \) is positive, i.e., \( I(f) \geq 0 \) for all \( f \in C(X) \) with \( f \geq 0 \), then there exists a unique (positive) measure \( \mu \) on the Borel-\( \sigma \)-algebra of \( X \) such that (24.2.29) holds.

**Idea of the proof.** The idea is that one would like to define the measure \( \mu \) by setting \( \mu(A) = I(1_A) \). However, this is not possible since characteristic functions are not continuous (for most sets). Therefore one proceeds by approximating characteristic functions by continuous functions. Verifying that this really works involves a number of technical details which the interested reader can find in [[Folland, Real Analysis, Ch.7]] for instance. \( \square \)

We cannot directly apply this result to the map (*) since this maps takes values in \( L(\mathcal{H}) \) instead of \( \mathbb{C} \). However, the Riesz representation theorem can be applied to “expectation values” of (*), i.e., for fixed \( u, v \in \mathcal{H} \) it can be applied to the linear function

\[
C(\sigma(T)) \to \mathbb{C}, \quad f \mapsto \langle f(T)u, v \rangle .
\]

Note here that this function is continuous since

\[
| \langle f(T)u, v \rangle | \leq \| f(T)u \| \cdot \| v \| \leq \| f(T) \|_\infty \cdot \| u \| \cdot \| v \| = \| f \|_\infty \cdot \| u \| \cdot \| v \| . \tag{24.2.30}
\]

Applying the Riesz representation theorem (note again that \( \sigma(T) \) is compact since \( T \) is bounded) we get a complex measure \( E_{u,v} \), which depends on \( u \) and \( v \), such that

\[
\forall f \in C(\sigma(T)) : \quad \langle f(T)u, v \rangle = \int_{\sigma(T)} f(x) \, dE_{u,v}(x) . \tag{24.2.30}
\]
All these complex measures $E_{u,v}$, where $u$ and $v$ range over all of $\mathcal{H}$ can finally be combined into one spectral measure $E$ by the Riesz representation theorem for jointly continuous bilinear maps on Hilbert spaces (see Section 24.3). Namely, for fixed $A \in \sigma(T)$ consider the map

$$B_A : \mathcal{H} \times \mathcal{H} \to \mathbb{C}, \quad (u, v) \mapsto E_{u,v}(A) = \int_{\sigma(A)} \mathbb{1}_A \, dE_{u,v}.$$  

It follows immediately from (24.2.30) that the map $B_A$ is sesquilinear. Moreover, for any $u, v \in \mathcal{H}$ we have

$$|B_A(u, v)| = \left| \int_{\sigma(A)} \mathbb{1}_A(x) \, dE_{u,v}(x) \right| \leq \|\mathbb{1}_A\|_{\infty} \cdot \|E_{u,v}\| = \|u\| \cdot \|v\|,$$

where $\|E_{u,v}\|$ denotes the total variation of the complex measure $E_{u,v}$ (see Section 24.3). This shows that $B_A$ is jointly continuous. By the Riesz representation theorem for jointly continuous bilinear maps on Hilbert spaces (see Section 24.3) there thus exists $E(U) \in L(\mathcal{H})$ such that

$$\forall u, v \in \mathcal{H} : \langle E(A)u, v \rangle = B_A(u, v) = \int_{\sigma(T)} \mathbb{1}_A \, dE_{u,v}. \tag{24.2.31}$$

We now show that this define a spectral measure from which $T$, and also the whole continuous functional calculus of $T$ can be recovered by integration.

**Theorem 24.2.23 (Spectral theorem for selfadjoint operators).** Let $T \in \mathcal{H}$ be selfadjoint. The map $E : B(\sigma(T)) \to L(\mathcal{H})$ constructed above is a spectral measure. Moreover, for every $f \in C(\sigma(A))$ one has $f(T) = \int_{\sigma(T)} f(x) \, dE(x)$, were the left-hand side is defined as in Theorem 24.2.21 and the right-hand side as in Theorem 24.2.14. In particular we have $A = \int_{\sigma(A)} x \, dE(x)$.

**Proof.** The proof in fact becomes simpler if we first extend the functional calculus to all bounded, measurable functions, and only afterwards verify that $E$ is a spectral measure.

To define $f(T)$ also for $f \in L^\infty(\sigma(T))$, similarly as in the construction of $E$ we start with the function $B_f : \mathcal{H} \times \mathcal{H} \to \mathbb{C}$ defined by $B_f(u, v) := \int_{\sigma(T)} f(x) \, dE_{u,v}(x).$ Just as the map $B_A$ above, also $B_f$ is sesquilinear and jointly continuous. So by the Riesz representation theorem for jointly continuous sesquilinear maps in Hilbert spaces there exists a unique operator $f(T) \in L(\mathcal{H})$ such that $(f(T)u, v) = B_f(u, v)$ for all $u, v \in \mathcal{H}$.

Note that in particular $\mathbb{1}_A(T) = E(A)$ holds for any $A \in B(\sigma(T))$. That $E$ maps into the orthogonal (selfadjoint) projections will follow easily from this once we have established that this new functional calculus is still multiplicative and involutive.

Let us first show that multiplicativity remains to hold. Fix $u, v \in \mathcal{H}$. Due to multiplicativity of the continuous functional calculus, for any $f, g \in C(\sigma(T))$ we have

$$\int_{\sigma(T)} f(x)g(x) \, dE_{u,v}(x) = \langle (f(T)g(T))u, v \rangle = \langle f(T)g(T)u, v \rangle = \int_{\sigma(T)} f(x) \, dE_{g(T)u,v}(x).$$

Thinking of $g \in C(\sigma(T))$ as being fixed, this shows that integration with respect to the two measures $g \, dE_{u,v}$ and $E_{f(T)u,v}$ coincides on $C(\sigma(T))$. From the uniqueness statement in Riesz representation theorem for the dual of $C(\sigma(T))$ it follows that actually these two measures coincide. Therefore it follows that for any $f \in L^\infty(\Omega)$ we have

$$\int_{\sigma(T)} f(x)g(x) \, dE_{u,v}(x) = \int_{\sigma(T)} f(x) \, dE_{g(T)u,v}(x).$$

\(^3\)For $f$ continuous, by (24.2.30) this definition of $f(T)$ coincides with the previous one of Theorem 24.2.21 by (24.2.30).
The right-hand side can be rewritten as
\[
\int_{\sigma(T)} f(x) \, dE_{g(T)u,v}(x) = \langle f(T)g(T)u, v \rangle = \langle g(T)u, f(T)^*v \rangle = \int_{\sigma(T)} g(x) \, dE_{u,f(T)^*v}(x).
\]

Together with the previous identity, we see that
\[
\int_{\sigma(T)} f(x)g(x) \, dE_{u,v}(x) = \int_{\sigma(T)} g(x) \, dE_{u,f(T)^*v}(x) \quad (*)
\]
holds for all \( g \in C(\sigma(T)) \) and \( f \in L^\infty(\sigma(T)) \). Analogously as before, it follows from the Riesz representation theorem for the dual of \( C(\sigma(T)) \) that the measures \( f \, dE_{u,v} \) and \( dE_{u,f(T)^*v} \) coincide. Therefore \((*)\) remains valid also for \( g \in L^\infty(\sigma(T)) \). Thus, in summary, we have seen that for any \( f, g \in L^\infty(\sigma(T)) \) we have
\[
\langle (fg)(T)u, v \rangle = \int_{\sigma(T)} f(x)g(x) \, dE_{u,v}(x) \overset{(*)}{=} \int_{\sigma(T)} g(x) \, dE_{u,f(T)^*v}(x)
\]
\[
= \langle g(T)u, f(T)^*v \rangle = \langle f(T)g(T)u, v \rangle.
\]

Since also \( u, v \in \mathcal{H} \) were arbitrary, this shows multiplicativity. As an immediate consequence, for any \( A \in B(\sigma(T)) \) and \( u, v \in \mathcal{H} \) we have
\[
\langle E(A)^2u, v \rangle = \langle 1_A(T)1_A(T)u, v \rangle = \langle (1_A)^2(T)u, v \rangle \overset{T^2=1_A}{=} \langle 1_A(T)u, v \rangle = \langle E(A)u, v \rangle.
\]

This shows that \( E(A) \) is a projection. To show that \( E(A) \) is an orthogonal projection, we have to show that \( E(A) \) is selfadjoint. To this end, note that for any \( u, v \in \mathcal{H} \) and \( f \in C(\sigma(T)) \) we have
\[
\int_{\sigma(T)} f(x) \, dE_{u,v}(x) = \langle f(T)u, v \rangle = \langle v, f(T)^*u \rangle = \langle f(T)^*v, u \rangle = \langle f(T)v, u \rangle
\]
\[
= \int_{\sigma(T)} f(x) \, dE_{v,u}(x) = \int_{\sigma(T)} f(x) \, d\overline{E_{v,u}(x)}.
\]

As before, since this holds for all \( f \in C(\sigma(T)) \), it follows that \( E_{u,v} = \overline{E_{v,u}} \). But this precisely means that for any \( A \in B(\sigma(T)) \) we have
\[
\langle E(A)u, v \rangle = E_{u,v}(A) = \overline{E_{v,u}(A)} = \langle E(A)v, u \rangle = \langle u, E(A)v \rangle.
\]

This shows that \( E(A)^* = E(A) \). We conclude that \( E \) indeed maps into the orthogonal projections of \( \mathcal{H} \).

By a similar computation as above (Exercise 27.39), using that \( \overline{E_{u,v}} = E_{v,u} \) holds for any \( u, v \in \mathcal{H} \), it follows that \( f(T)^* = \overline{f(T)} \) holds for any \( f \in L^\infty(\sigma(T)) \).

That \( E(\sigma(T)) = \text{id}_\mathcal{H} \) holds follows since \( E(\sigma(T)) = 1_{\sigma(T)}(T) \) and since the constant function 1 is mapped to \( \text{id}_\mathcal{H} \) already by the continuous functional calculus.

It remains to show \( \sigma \)-additivity. To this end, let \( A_1, A_2, \ldots, \sigma(T) \) be disjoint and measurable. Set \( A := \bigcup_{i=1}^n A_i \). Further set \( f_n := \sum_{i=1}^n 1_{A_i} \), then \( f_n \to 1_A \) pointwise and \( \|f_n\|_\infty \leq 1 \) for all \( n \in \mathbb{N} \). Now let \( u \in \mathcal{H} \) be given. To show that \( \sum_{i=1}^n E(A_i)u \) converges to \( E(A)u \), by Exercise ?? it suffices to show that \( \sum_{i=1}^n \langle E(A_i)u, v \rangle \to \langle E(A)u, v \rangle \) for all \( v \in \mathcal{H} \) and \( \sum_{i=1}^n \|E(A_i)u\| \to \|E(A)u\| \). Concerning weak convergence, for every \( v \in \mathcal{H} \) we have by dominated convergence \((\|f_n\|_\infty \leq 1)\)
\[
\langle \sum_{i=1}^n E(A_i)u, v \rangle = \int_{\sigma(T)} f_n(x) \, dE_{u,v}(x) \xrightarrow{\text{dom. conv.}} \int_{\sigma(T)} 1_A(x) \, dE_{u,v}(x) = \langle E(A)u, v \rangle.
\]
Concerning convergence of the norms, we have
\[
\left\| \sum_{i=1}^{n} E(A_i)u \right\|^2 = \|f_n(A)u\|^2 = \langle f_n(A), f_n(A)u \rangle = \langle f_n(A)^* f_n(A)u, u \rangle = \langle |f_n|^2(A)u, u \rangle.
\]
This concludes the proof that \( E \) is a spectral measure. The other claims have already been shown.

We restate the possibility to define bounded, measurable functions of a selfadjoint operator in a definition. For continuous functions, respectively polynomials, this definition coincides with the previous ones.

**Definition 24.2.24** (\( \mathcal{L}^\infty \)-functional calculus for selfadjoint operators). Let \( T \in L(\mathcal{H}) \) be selfadjoint, and let \( E : B(\sigma(T)) \to L(\mathcal{H}) \) be the spectral measure of \( T \) as in Theorem 24.2.23. For any \( f \in \mathcal{L}^\infty(\sigma(T)) \) we define the operator \( f(T) \in L(\mathcal{H}) \) by
\[
f(T) = T_E(f) = \int_{\sigma(T)} f(x) \, dE(x),
\]
with the integral defined as in Theorem 24.2.15.

The properties of this functional calculus, most of which have already been shown, are summarized in the following theorem.

**Theorem 24.2.25** (Properties of the functional calculus). Let \( T \in L(\mathcal{H}) \) be selfadjoint.

1. For all \( f, g \in \mathcal{L}^\infty(\sigma(T)) \) and \( \alpha, \beta \in \mathbb{C} \) it holds that
\[
(\alpha f + \beta g)(T) = \alpha f(T) + \beta f(T),
\]
\[
(fg)(T) = f(T)g(T),
\]
\[
(f(T))^* = f(T)^*,
\]
\[
[f(T), g(T)] = 0.
\]
In particular, the operator \( f(T) \) is selfadjoint if and only if \( f \) is real-valued.

2. For every \( f \in \mathcal{L}^\infty(\sigma(T)) \) with \( f \geq 0 \) also \( f(T) \geq 0 \).

3. If \( Tu = \lambda u \) for \( u \in \mathcal{H} \) and \( \lambda \in \mathbb{C} \), then \( f(T)u = f(\lambda)u \) for every \( f \in \mathcal{L}^\infty(\sigma(T)) \).

4. For every \( f \in \mathcal{L}^\infty(\sigma(T)) \) one has \( \sigma(f(T)) \subseteq f(\sigma(T)) \).

5. For every \( f \in \mathcal{L}^\infty(\sigma(T)) \) one has \( \|f(T)\|_{\infty} = \text{ess sup}|f| \).

**Proof.** The first part has already been shown in the proof of the previous theorem. The second part follows from multiplicativity precisely as for continuous functions. The third part can be shown by the same trick that was used several times in the proof of the previous theorem. The fourth part can again be proven precisely as for continuous functions. The last estimate can be directly established using the definition of \( f(T) \).

That the fourth statement has to be changed compared to the continuous functional calculus can be expected since one may change measurable function \( f \) on a set of measure zero (with respect to \( E \)) without changing the operator \( f(T) \), also see Exercise 27.51.
24.2. The spectral theorem for normal and several commuting operators. The spectral theorem can be extended to normal operators by noting that similar to a complex number, any normal operator \( T \in L(H) \) can be written as \( T = A + iB \) for operators \( A, B \in L(H) \) which are selfadjoint ("real") and commute. Simply set
\[
A := \frac{T + T^*}{2}, \quad B := \frac{T - T^*}{2i},
\]
then obviously \( T = A + iB \) holds. It is easy to verify that \( A \) and \( B \) are selfadjoint. They commute since \( T \) is normal. The idea for constructing a spectral measure for \( T \) is now now that similarly to how the product measure \( dx \, dy \) is the Lebesgue measure on \( \mathbb{R}^2 \), the spectral measure of \( T \) should be the product of the spectral measures of \( A \) and \( B \).

Therefore we now explain in which sense a product of spectral measures can indeed be defined. This works nicely, if the different spectral measures commute. In the following explanation we focus on the product of two commuting spectral measures, but everything can be extended in a straightforward sense to any finite number of commuting spectral measures.

**Theorem 24.2.26.** Let \( X_1, \ldots, X_n \) be complete, separable metric spaces, \( H \) a Hilbert space, and let \( E_1: B(X_1) \rightarrow L(H), \ldots, E_n: B(X_n) \rightarrow L(H) \) be spectral measures. Suppose that \( E_1, \ldots, E_n \) are pairwise commuting, i.e., \( E_i(A), E_j(B) = 0 \) holds for all \( i, j \in \{1, \ldots, n\} \) and \( A \in B(X_i), B \in B(X_j) \). Then there exists a unique spectral measure \( E: B(X_1 \times \cdots \times X_n) \rightarrow L(H) \) satisfying
\[
\forall A_1 \in B(X_1), \ldots, A_n \in B(X_n): \quad E(A_1 \times \cdots \times A_n) = E_1(A_1) \cdots E_n(A_n).
\]
(Note that due to the commutation assumption the ordering of the operators \( E_1(A_1), \ldots, E_n(A_n) \) on the right-hand side plays no role.)

We will not prove this theorem, although the proof is not very hard, but would take us too far into measure-theoretic arguments. However, let us make some brief remarks. The idea is of course to take (24.2.38) as definition for \( E \) on the subset of \( B(X_1 \times X_2) \) which consists of all product sets. By "standard" measure-theoretic arguments one can extend \( E \) to the product-\( \sigma \)-algebra of \( B(X_1) \) and \( B(X_2) \) (which is the \( \sigma \)-algebra generated by all products). The separability assumption gives that this product-\( \sigma \)-algebra coincides with \( B(X_1 \times X_2) \), whereas in general it might be smaller (see [Folland, Real Analysis, Sec.1.2]). For details we refer to [Birman, Chapter 5, Section 2][].

24.2.7. The joint spectral theorem for commuting selfadjoint operators. In this section we show that given several selfadjoint operators which commute with each other, there exists a "joint spectral measure" on the product of the spectra of the operators, from which all of these operators as well as functions of several of them can be reproduced via integration.

**Theorem 24.2.27** (Joint spectral theorem). Let \( T_1, \ldots, T_n \in L(H) \) be selfadjoint, pairwise commuting operators, and denote by \( E_1, \ldots, E_n \) their respective spectral measures. There exists a unique spectral measure \( E \) on \( B(\sigma(T_1) \times \cdots \times \sigma(T_n)) \) such that \( E(A_1 \times \cdots \times A_n) = E_1(A_1) \cdots E_n(A_n) \) holds for all \( A_1 \in B(\sigma(T_1)), \ldots, A_n \in B(\sigma(T_n)) \). Moreover, for any \( f_1 \in L^\infty(\sigma(T_1)), \ldots, f_n \in L^\infty(\sigma(T_n)) \) one has
\[
\int f_1(x_1) \cdots f_n(x_n) \, dE(x_1, \ldots, x_n) = f_1(T_1) \cdots f_n(T_n).
\]
In particular,
\[
A_1 = \int x_1 \, dE(x_1, \ldots, x_n), \ldots, A_n = \int x_n \, dE(x_1, \ldots, x_n).
\]

The spectral measure \( E \) is also the unique spectral measure on \( B(\sigma(T_1) \times \cdots \times \sigma(T_n)) \) with the property that the last formulas hold.

**Proof.** The proof for existence of such a measure is similar to the scalar case and can be found in []. The idea is to approximate general measurable subsets of the product by products of measurable sets. The identity about the integral of a product of functions is easily verified for
simple functions, the proof for general bounded, measurable functions works can afterwards be done by approximation.

One calls the support of the joint spectral measure of operators $T_1, \ldots, T_n$ their joint spectrum, and denotes it by $\sigma(T_1, \ldots, T_n)$. In general, $\sigma(T_1, \ldots, T_n)$ may only be a subset of $\sigma(T_1) \times \cdots \times \sigma(T_n)$.

24.2.8. The spectral theorem for normal operators.

**Theorem 24.2.28** (Spectral theorem for normal operators). Let $T \in L(H)$ be a normal operator. Then there exists a unique spectral measure $E$ on $\mathcal{B}(\sigma(T))$ such that

$$T = \int_{\sigma(T)} \lambda \, dE(\lambda).$$

**Proof.** Define

$$A := \frac{T + T^*}{2}, \quad B := \frac{T - T^*}{2i}.$$ 

Then $T = A + iB$ and $A$ and $B$ are both selfadjoint. Moreover, due to $T^*T = TT^*$, the operators $A$ and $B$ commute. Now first take the joint spectral measure. Probably one shows that its support is precisely $\sigma(T)$.

24.2.9. Unbounded Selfadjoint Operators.

- dense domains
- explain difference between symmetric and selfadjoint operators
- state the spectral theorem for unbounded operators

Since this Hamiltonian is unbounded, we need to specify a suitable domain of definition $\mathcal{D}(H)$. For simplicity, we here take the smooth and square integrable functions, $\mathcal{D}(H) = L^2(\mathbb{R}^3) \cap C^\infty(\mathbb{R}^3)$. Then $\mathcal{H}$ satisfies the requirement (1.1.6) in the sense that

$$\langle H \phi | \psi \rangle_{\mathcal{H}} = \langle \phi | H \psi \rangle_{\mathcal{H}} \quad \text{for all } \phi, \psi \in \mathcal{D}(H).$$

In mathematical terms, the Hamiltonian is a symmetric operator (selfadjointness is a stronger statement; we postpone the details to Section ?? below).
A Brief Overview of Microlocal Analysis

In the previous chapter we constructed the fermionic signature operator
\[ S_m : \mathcal{H}_m \to \mathcal{H}_m , \]
being a bounded and symmetric operator. We now briefly recall why this construction is useful. Basically, there are two ways to proceed:

- Construction of a corresponding causal fermion system:
  The subspace \( \mathcal{H} \) is introduced typically as the negative spectral subspace of the fermionic signature operator,
  \[ \mathcal{H} := \chi_{(-\infty,0)}(\mathcal{S}_m)(\mathcal{H}_m) \subset \mathcal{H}_m . \]
  We saw that in the Minkowski vacuum, \( \mathcal{H} \) corresponds precisely of the subspace of all negative-frequency solutions of the Dirac equation. In the presence of an external potential, \( \mathcal{H} \) can be regarded as a generalization of the Dirac sea picture to the time-dependent setting.

  Now we can proceed as described in Section 5.4: For \( 0 < \varepsilon < \varepsilon_{\text{max}} \) we again introduce regularization operators \( \mathcal{R}_\varepsilon \)
  \[ \mathcal{R}_\varepsilon : \mathcal{H} \to C^0(\mathcal{M}, S_M) \]
  and define the local correlation operators \( F^\varepsilon(x) \in \mathbb{L}(\mathcal{H}) \) by
  \[ \langle u | F^\varepsilon(x) v \rangle_{\mathcal{H}} = -\langle (\mathcal{R}_\varepsilon u)(x) | (\mathcal{R}_\varepsilon v)(x) \rangle \quad \text{for all } u, v \in \mathcal{H} . \]
  Taking the push-forward measure,
  \[ \rho = (F^\varepsilon)_* \mu \]
  we obtain a causal fermion system \((\mathcal{H}, F, \rho)\).

  Now one can analyze the structures of the obtained causal fermion systems and relate them to the causal structure of Minkowski space, the potentials contained in \( \mathcal{B} \), etc. One can also introduce particles and/or anti-particles by “occupying additional state” or by “creating holes in the Dirac sea.” The most important and difficult task is to analyze the causal action for such systems in the limit \( \varepsilon \searrow 0 \). For which potentials and particle/anti-particle configurations does one have a critical point of the causal action? This question is analyzed in [48].

- Alternatively, one can prefer not to introduce the regularization at this stage. This is a good procedure as long as the interaction is purely by an external field. As soon as one wants to describe the full interaction (no matter if using the causal action principle or in the framework of standard quantum field theory), an ultraviolet regularization is essential, giving rise to the subtle question of what happens in the limit \( \varepsilon \searrow 0 \).

Let us outline the procedure without regularization. Then the right object to consider is the unregularized kernel of the fermionic projector \( P(x, y) \) defined as follows. Recall that in Section 16.2 the causal fundamental solution was defined as an operator to the solution space

\[ k_m := \frac{1}{2\pi i} (s' - s^\wedge) : C^\infty_c(\mathcal{M}, S\mathcal{M}) \to \mathcal{H}_m \cap C^\infty_w(\mathcal{M}, S\mathcal{M}) . \]
Multiplying by the projection operator to the negative spectral subspace of the fermionic signature operator gives the unregularized fermionic projector $P$,

$$P = -\chi(-\infty,0)(S) k_m : C_0^\infty(M,SM) \to \mathcal{H}_m \cap C_0^\infty(M,SM).$$

Next, one represents this operator as an integral operator with a distributional kernel:

**Theorem 25.0.1.** There is a unique distribution $P \in \mathcal{D}'(M \times M)$ such that for all $\phi, \psi \in C_0^\infty(M,SM)$,

$$<\phi|P|\psi> = P(\phi \otimes \psi).$$

The proof uses the Schwartz kernel theorem; see [76, Theorem 3.12].

Next, the connection to quantum field theory is made by the following theorem, which is based on an abstract construction due to Araki [1] (see [75, Theorem 1.4] or [48, Theorem 1.5.1]):

**Theorem 25.0.2.** There are fermionic field operators $\hat{\Psi}_\alpha(x)$ and $\hat{\Psi}_\beta(y)^*$ together with a ground state $|0>$ with the following properties:

(a) The canonical anti-commutation relations hold:

$$\{\hat{\Psi}_\alpha(x),\hat{\Psi}_\beta(y)^*\} = k_m(x,y)_\beta^\alpha, \quad \{\hat{\Psi}_\alpha(x),\hat{\Psi}_\beta(y)\} = 0 = \{\hat{\Psi}_\alpha(x)^*,\hat{\Psi}_\beta(y)^*\}.$$

(b) The two-point function is given by

$$<0|\hat{\Psi}_\alpha(x)\hat{\Psi}_\beta(y)^*|0> = -P(x,y)_\beta^\alpha.$$

We again point out that this theorem defines a quasi-free (i.e. non-interacting) quantum field theory. Introducing an interaction is a difficult task. A first step is to introduce Wick ordering and time-ordered products. It turns out that in order to make mathematical sense of these operations, the two-point distribution must be of Hadamard form. If this is the case, one can even define the Clearly, here we cannot enter the constructions of perturbative quantum field theory (the interested reader is referred for example to [122]). But the Hadamard property can be understood simply in the context of distribution theory, as we now outline.
Clearly, causal fermion systems were developed for describing spacetimes. Also, the resulting structures are closely tied to structure in spacetime. Nevertheless, by slightly generalizing the definition, it also applies to the Euclidean and Riemannian setting:

**Definition 26.0.1.** Given a complex Hilbert space \((\mathcal{H}, \langle \cdot, \cdot \rangle_{\mathcal{H}})\) and parameters \(p, q \in \mathbb{N}_0\) with \(p \leq q\), we let \(\mathcal{F} \subset L(\mathcal{H})\) be the set of all self-adjoint operators on \(\mathcal{H}\) of finite rank, which (counting with multiplicities) have at most \(p\) positive and at most \(q\) negative eigenvalues. On \(\mathcal{F}\) we are given a positive measure \(\rho\) (defined on a \(\sigma\)-algebra of subsets of \(\mathcal{F}\)), the so-called universal measure. We refer to \((\mathcal{H}, \mathcal{F}, \rho)\) as a topological fermion system of spin signature \((p, q)\).

In the case \(p = q\), we call \((\mathcal{H}, \mathcal{F}, \rho)\) a causal fermion system of spin dimension \(n := p\). If \(p = 0\), we call \((\mathcal{H}, \mathcal{F}, \rho)\) a Riemannian fermion system of spin dimension \(n := q\).

This definition was first given in \([58]\), where also many examples and applications are worked out.

For a Riemannian fermion system, the causal structure of Definition 5.5.1 becomes trivial. Namely, in this case the operator \(-y\) is positive semi-definite, so that its square root \(\sqrt{-y}\) is well-defined as a positive semi-definite operator. Using that the spectrum is invariant under cyclic permutations, it follows that

\[
x y = -x \sqrt{-y} \sqrt{-y}\text{ is isospectral to } \sqrt{-y} (-x) \sqrt{-y},
\]

and the last operator product is obviously positive semi-definite. As a consequence, the eigenvalues of \(xy\) are all real and non-negative. As a consequence, the points \(x\) and \(y\) are necessarily timelike separated, except in the trivial case that all eigenvalues of \(xy\) coincide. In view of this shortcoming, it is not clear whether the causal action principle is sensible for Riemannian fermion systems. It is conceivable that other variational principles are of mathematical interest, but this has not yet been explored.
Overview of the Continuum Limit

The continuum limit is a limit which allows to deduce quantum theory, general relativity
and quantum field theory from the theory of causal fermion systems. In this chapter, we give an
overview about this limit. A thorough explanation of all details can be found in the simultaneously appearing book [48].

Among other things, in this overview, we intend to address worries about the legitimacy of
this limit. Therefore, we have carefully chosen the presentation such that the reader can clearly
identify what is “put into” the analysis, and what comes out. Furthermore, this chapter only
builds on the concepts introduced in Ch. ??, no further knowledge is required at this point.

27.1. General Idea

The continuum limit allows to connect the theory of causal fermion systems to quantum the-
tory, general relativity and quantum field theory. However, we want to make it very clear from the
beginning that the continuum limit does not merely consist of changing the value of one constant
of the theory. Indeed, the theory, as outlined, e.g., in Ch. ??, does not posses constants on the
very fundamental level. Constants, such as the charge or mass of particles, only appear when
specifying the minimizers of the theory. To see what this means consider again the example of
section ??, where we have constructed the causal fermion system representing Minkowski vacuum.
If, in this case, we change the mass of the Dirac sea particles $\psi_k$, from $m$ to $m'$, and carry out
the same construction as in Sec. ??, we obtain a different causal fermion system. (This is so since
different mass implies different solutions of the Dirac eq. and therefore different local correlation
matrices in general.) This means that if one changes the constants characterzing a causal fermion
system one really changes the causal fermion system itself. And in general, in cases where the
constants are not merely the masses of Dirac sea particles, it is not clear that a change of constant
will yields a causal fermion system which is again a minimizer of the causal action principle.

This shows that in the context of this theory, it is natural that the continuum limit is indeed a
limit of causal fermion systems. This means that one considers a family of causal fermion systems
$(\mathcal{H}_\tau, \mathcal{F}_\tau, \rho_\tau)_{\tau \in (-\varepsilon, \varepsilon)}$ and the limit

$(\mathcal{H}_\tau, \mathcal{F}_\tau, \rho_\tau) \to (\mathcal{H}_0, \mathcal{F}_0, \rho_0), \quad (27.1.1)$

where $(\mathcal{H}_0, \mathcal{F}_0, \rho_0)$ is a causal fermion system which is in some sense “distinguished”. In practise,
both the number $n$ and $f$ in the definition of causal fermion systems (Def. ??), and thus $\mathcal{H}$ and
$\mathcal{F}$, are fixed. Therefore, Eq. (??) reduces to

$\rho_\tau \to \rho_0. \quad (27.1.2)$

In this sense, the continuum limit is a limit.

At this point, the reader probably wonders how something as concrete as quantum theory or
general relativity can “come out” of something of the sort of Eq. (27.1.2). What lies behind this is
a fact which we appeared already in Ch. ?? and Ch. ??: Causal fermion systems can be described
with objects analogous to objects of contemporary physical theories, e.g. wave-functions, tangent spaces, connections, etc., \( c.f. \) sections ??, ?? and ?? . Therefore, the limit (27.1.2) is at the same time a limit of all those objects:

\[
\psi_k \tau \rightarrow \psi_0
\]

\[
\nabla_\tau \rightarrow \nabla_0
\]

\[
...
\]

The essential question in the continuum limit is thus \textit{which equations relate those objects} and the central statement is that \textit{in the limit (27.1.2), those are the equations of contemporary physics}. Taking a limit is necessary since the “distinguished” causal fermion systems \((\mathcal{H}_0, \mathcal{F}_0, \rho_0)\) have a very singular behavior, \( c.f. \) below, and thus the limit corresponds to a regularization of \((\mathcal{H}_0, \mathcal{F}_0, \rho_0)\).

This is the general concept behind the continuum limit, which we will explain in more detail in this chapter. Note that the name \textit{continuum limit} is of historical origin: The concepts which led to the theory of causal fermion systems were first built out of discrete spacetimes, and the continuum limit (27.1.2) was indeed a limit which consisted out of a family of such discrete systems which approximated a continuous spacetime. This is not so any more, both the theory of causal fermion systems and the continuum limit are of course fully compatible with continuous spacetimes.

### 27.2. The Fermionic Projector

Ferm. Projektor: Evt. bekannt aus Kapt. ... projeziert auf besetzte zustände... deswegen wichtigstes objekt in der analyse der theorie... auch unabhängig von TCFS: wichtiges objekt: zeichnet eindeutiges vakuum aus, erfüllt die Hadamardbedingung (...) deswegen heißt das buch auch TCFS and Ferm. Proj.

auch im kont-limis: wichtige rolle erklären, warum...

The fermionic projector, which was the central object of the investigations of Chs. ??, ?? and ??, also plays a key role in the continuum limit. This is so since it projects upon all occupied wave-functions, i.e. all wave-functions which are “really there”. (\textit{C.f.} the explanation of this on p. ??) - bra-ket notation - link zu wellenfunktionen -

The general idea of the continuum limit is the following. We have seen in Ch. ?? that given a causal fermion system, the fermionic projector \( P \) can be introduced, and we have seen that it projects upon all occupied wave functions. Furthermore, we have seen in Ch. ?? that the physical situation of Minkowski space (and Dirac “particles” thereon) can be reexpressed as a causal fermion system, and in Ch. ?? that the fermionic projector of the causal fermion system describing Minkowski space takes the form

\[
P(x, y) = \frac{1}{(2\pi)^4} \int \delta(k^2 - m^2) \theta(-k^0) e^{-i(x-y)k} d^4k.
\]  

(27.2.1)

It is easy to check that \( P \) satisfies the vacuum Dirac equation

\[
(i\partial_x - m)P(x, y) = 0,
\]  

(27.2.2)

The physical picture is that this causal fermion system can be described as though Dirac particles were propagating in vacuum, i.e. without external potential.

The central question behind the continuum limit is whether there is also a causal fermion system which can be described as though there were Dirac particles propagating in an external potential. Of course, since one is only interested in “physical” causal fermion systems, according to the postulates of the theory, this causal fermion system should be a solution of the causal variational principle. In mathematical terms, the question is therefore: Is there a causal fermion
system which solves the causal variational principle and whose fermionic projector \( \tilde{P} \) satisfies the Dirac equation with external potential \( \mathcal{B} \):

\[
(i\partial_x + \mathcal{B} - m)\tilde{P}(x, y) = 0. \tag{27.2.3}
\]

The essential conceptual point is that the requirement of it being a solution of the causal variational principle leads to constraint equations which involve \( \mathcal{B} \) and the fermionic projector. Only if those equations are fulfilled can the causal fermion system be a minimizer of the causal variational principle. The analysis, which we sketch in some more detail in this chapter, shows that those equations turn out to be the equations known from contemporary physics, e.g. the Yang-Mills equation coupled to the Dirac current.

The external potential can a priori have any form. It could be, e.g., a vector potential, i.e. \( \mathcal{B}(x) = \mathcal{A}(x) \), where \( \mathcal{A} \) is a 4-vector on Minkowski space. Within this chapter, we will always assume that \( \mathcal{B} \) is of the form

\[
\mathcal{B} = \chi_L \mathcal{A}_L + \chi_R \mathcal{A}_R. \tag{27.2.4}
\]

where \( \chi_L \) and \( \chi_R \) are as described on p. ?? and \( \mathcal{A}_L, \mathcal{A}_R \) are called left-handed and right-handed component, respectively. In order to obtain the limit for the gravitational dynamics, \( \mathcal{B} \) can be chosen as a differential operator, c.f. [48].

### 27.3. Overview of the Methods

In order to plug the causal fermion system corresponding to \( \tilde{P} \) into the causal variational principle, we first need to express \( \tilde{P} \) explicitly for a given potential \( \mathcal{B} \). We have already seen a non-perturbative construction thereof in Ch. ??, but for the purpose of the continuum limit it is more helpful to do so perturbatively. Once this is done, we reformulate the causal variational principle in terms of the eigenvalues of the quantity \( A_{xy} = \tilde{P}(x,y)\tilde{P}(y,x) \), and insert the result of the perturbative construction.

#### 27.3.1. Perturbative construction of \( \tilde{P} \).

The perturbative construction of \( \tilde{P} \) is carried out in two steps:

**Step (1):** Construct the Greens functions of the Dirac equation with external potential \( \mathcal{B} \), perturbatively.

**Step (2):** Carry the result of step (1) over to \( \tilde{P} \).

**Step (1):** We already know the advanced and retarded Greens functions of the vacuum Dirac equation from Ch. ??: Those are the functions

\[
s_{m}^{\vee} (x,y) = \int \frac{d^4 k}{(2\pi)^4} s_{m}^{\vee} (k) e^{-ik(x-y)}, \tag{27.3.1}
\]

with

\[
s_{m}^{\vee} (k) = \lim_{\epsilon \to 0^{-}} \frac{k + m}{k^2 - m^2 - i\epsilon k^0}\quad \text{and}\quad s_{m}^{\wedge} (k) = \lim_{\epsilon \to 0^{+}} \frac{k + m}{k^2 - m^2 + i\epsilon k^0}. \tag{27.3.2}
\]

For convenience, we include the mass \( m \) of the Dirac equation into the external potential. This means that we consider the case where the vacuum Dirac equation is massless, and we have as Dirac equation with external potential

\[
(i\partial_x + B) \tilde{P}(x, y) = 0. \tag{27.3.3}
\]
where $B = B + m$. The advanced and retarded Green’s functions thereof can be obtained from $s^\vee_m$ and $s^\wedge_m$ perturbatively by

$$
\tilde{s}^\vee_m = \sum_{n=0}^{\infty} (-s^\vee_m B)^n s^\vee_m , \quad \tilde{s}^\wedge_m = \sum_{n=0}^{\infty} (-s^\wedge_m B)^n s^\wedge_m .
$$

This perturbation expansion is called \textit{causal perturbation expansion}.

In order to be able to analyze and utilize the terms $(-s^\vee/\wedge_m B^k s^\vee/\wedge_m)$ of the perturbation expansion, we introduce a further expansion ...

\textbf{Step (2)}:

In order to answer this question, we calculate the fermionic projector $\tilde{P}$ perturbatively in powers of $B$ (still without specifying the dependence of $B$ on $x$) and then plug $\tilde{P}$ into the causal variational principle, or to be more precise into the Euler-Lagrange equations of the causal variational principle. What one finds then is that the causal fermion system which $\tilde{P}$ belongs to can only by a minimizer of the causal variational principle (and hence, can “realizable in nature” (c.f. Sec.??)) if the external potential $B$ and the wave functions $\psi$ are related in a particular way. If, e.g., we choose $B$ as in (??), we find that the causal fermion system is a minimizer only if $A_L, A_R$ and $\psi$ are related by

$$
(i\partial + \gamma^5 A_a - m)\Psi = 0 , \quad C_0 (\partial^k A^l_a - \Box A^k_a) - C_2 A^k_a = 12\pi^2 \nabla \gamma^5 \gamma^k \Psi ,
$$

where $A_a = A_L + A_R$. Here, $C_0$ and $C_2$ are constants which are explained in Sec. ?? below. Those are the Dirac and Yang-Mills equation, i.e. the equations which govern the interaction of fermions and the electromagnetic field.

This shows that indeed the right physical equations appear in the continuum limit for this potential. In the remainder of this chapter we explain in more detail how the analysis is performed and how the result is obtained. For the treatment of other potentials and the deduction of the Einstein equations we refer to [48].

We now explain in more detail how the analysis works. Mathematically, the central point of the analysis is that the fermionic projector $P(x, y)$ diverges on the light-cone, i.e. if $(x - y)^2 = 0$.

In order to cope with this, one initially expands not only in powers of $B$ but also in powers of $\frac{1}{\xi}$, according to the following definition.

---

1This can be understood from the fact that if $(x - y)^2 = 0$, the exponential term vanishes and therefore the integral diverges.
27.5. THE CONTINUUM LIMIT OF CAUSAL FERMION SYSTEMS

We now mention a few directions which cannot be covered in the lectures and give a brief outlook.

27.4. Analysis of the Causal Action Principle

The causal action principle was introduced at the beginning of these lectures (see Section ??). But we never came back to it. Indeed, entering the analysis of this variational principle would have been a separate topic. Here we can only mention a few works in this direction. In the papers \[41, 43\], the existence of minimizers is proved in various situation. In simple terms, these results show that the causal action principle makes mathematical sense.

The important question is how the minimizers look like. In \[25\] a few simple numerical examples are discussed. In \[40\] an effect of spontaneous symmetry breaking elaborated. This can be interpreted as a manifestation of a more general effect of “spontaneous structure formation” (see \[57\] Section 3). In \[79\] \[13\] the Euler-Lagrange equations corresponding to the causal action principle are worked out. The analysis in \[13\] reveals a tendency of minimizers to be discrete. For a physical discussion of this effect we refer to \[57\] Section 4. Finally, the paper \[56\] is devoted to the question how an initial value problem can be posed for causal variational principles, and whether it has a unique solution.

Except for these few results, nothing is known on the structure of the minimizers. It is an interesting open problem to get a better understanding of the structure of minimizers of causal variational principles. It is planned to study the causal action principle in more advanced explicit examples in the near future.

27.5. The Continuum Limit of Causal Fermion Systems

It was the main motivation early on to get an alternative description of quantum field theory. This connection can indeed be made by considering the so-called continuum limit. Similar as in Section ??, one introduces an ultraviolet regularization on a microscopic scale \(\varepsilon\). But then one analyzes the Euler-Lagrange equations corresponding to the causal action principle in the limit as \(\varepsilon \downarrow 0\) when the regularization is removed. A subtle point is that the results do depend on how the regularization is introduced. In order to analyze this effect in detail, one considers a large class of regularizations ("method of variable regularization"). It turns out that the structure of the effective equations obtained in the continuum limit are indeed regularization independent; only the coupling constants and the bosonic masses depend on the details of the regularizations.

Due to these complications, it does not seem suitable to enter the details of the continuum limit analysis in a lecture series. But we the interested reader to the survey article \[44\] and the references therein. In the more recent paper \[45\], the continuum limit is worked out for a system involving neutrinos, giving rise to an effective interaction described by a left-handed, massive SU(2) gauge field and a gravitational field. The next step is to work out a model involving leptons and quarks \[46\], which should combine gravity with all the interactions of the standard model. The plan is to publish the three papers \[42, 45, 46\] together with a nice general introduction as a book.

We finally point out that it is an important open problem to work out the detailed expansion of the effective interaction in terms of Feynman diagrams (including all loop diagrams) and to compare it to the standard expansion of perturbative quantum field theory. This is a major project for the future.
Exercise 27.1. (Lorentz transformation of spinors) In the lecture, you have been introduced to the Lorentz transformation, to the Dirac equation and to spinors. If you are not trained as a physicist, you might wonder about the physical meaning of those quantities, and about the relation between them. The goal of this exercise is to help you understand this better.

1. First of all, explain briefly what Lorentz transformations are used for in physics. Make an example connected to you driving your car to university.
2. Electrons can be described by the Dirac equation, i.e., they are spinors. What happens to spinors when you make a Lorentz transformation, i.e., when you look at them from a moving reference frame? Why does anything have to happen at all?
3. In the lecture, you have seen the following lemma:
   For any orthochronous proper Lorentz transformation \( \Lambda \) there is a unitary matrix \( U(\Lambda) \) such that
   \[
   U(\Lambda) \Lambda^l_j \gamma_j U(\Lambda)^{-1} = \gamma^l .
   \]
   Explain its meaning and give a proof.

   Hints: Since \( \Lambda \) is orthochronous and proper, you can write it in the form \( \Lambda = \exp(\lambda) \), where \( \lambda \) denotes matrix \( (\Lambda^l_j)_{l,j} \). Take a look the family \( \Lambda(t) = \exp(\lambda t) \) of Lorentz transformations and differentiate it at \( t = 0 \) to obtain
   \[
   \lambda^l_j g_{lk} = -g_{jm} \lambda^m_k ,
   \]
   where \( g_{jm} \) is the Lorentz metric. (Use that Lorentz transformation leave the Lorentz metric invariant.)
   Use this equation to evaluate that the matrix
   \[
   u := \frac{1}{4} \lambda^l_k \gamma^k
   \]
   is anti-selfadjoint (with respect to the spin scalar product) and thus that the matrix \( U(t) := \exp(tu) \) is unitary. Finally, take a look at the family of matrices
   \[
   A(t) := U(t) \Lambda(t)^\dagger \gamma^j U(t)^{-1} .
   \]
   Use the commutation relations \( [\gamma^l \gamma^k, \gamma^j] = 2 (\gamma^l \delta^j_k - \delta^j_l \gamma^k) \) to evaluate \( \frac{d}{dt} A(t) \), and prove the lemma by comparing \( A(0) \) and \( A(1) \).

Exercise 27.2. (Getting familiar with the Dirac equation) In this exercise you will familiarize yourself with some of the objects in the context of the Dirac equation, which will be used throughout this lecture.

1. Take a look at the Dirac equation with external field, \((i \gamma^k (\partial_k - i A_k) - m)\psi = 0\). Multiply it with the operator \((i \gamma^l (\partial_l - i A_l) + m)\) and compare the result to the Klein-Gordon equation with external field. Explain why the additional term describes the coupling of spin to the electromagnetic field.
2. Check that in the Dirac representation, the Gamma matrices are self-adjoint w.r.t. the spin scalar product. In light of exercise 1, do you have an idea why this self-adjointness is important?

Please limit all text answers during this tutorial to 2-3 sentences (quarter of a page) per question.
(3) Take a look at the pseudoscalar matrix \( \rho = i\gamma^0\gamma^1\gamma^2\gamma^3 \). Evaluate whether it is self-adjoint and whether it commutes with the Gamma matrices. What is it used for?

(4) In order to define an operator whose square is the wave operator, Dirac introduced the Gamma matrices and their anti-commutation relations. (They form a so-called Clifford algebra.) As you know, Dirac introduced them as \( 4 \times 4 \)-matrices. Why so? What is the minimal dimension of a representation of this algebra? (Please prove your answer.)

(5) Show that the Gamma matrices satisfy the relation \( [\gamma^i, \gamma^j] = 2(\gamma^i g^{kj} - \delta^i_j \gamma^k) \) which you have (probably) used in exercise 1.

Exercise 27.3. Compute the support of the \( \delta \)-distribution at the origin.

Exercise 27.4. Show that the support of a distribution on \( \Omega \) is a closed subset of \( \Omega \).

Exercise 27.5. Show that for \( f \in L^1_{\text{loc}}(\Omega) \), the support of the regular distribution \( T_f \) coincides with the essential support of \( f \), which is defined by

\[
\text{esssup} f := \Omega \setminus \bigcup \{ U \subset \Omega \text{ open: } f|_U = 0 \text{ almost everywhere} \}.
\]

Exercise 27.6. Let \( T \in \mathcal{D}'(\Omega) \) and \( \chi \in C^\infty(\Omega) \) with \( \chi|_{\text{supp} T} = 1 \). Show that \( \chi T = T \).

Exercise 27.7. Let \( T \in \mathcal{D}'(\Omega) \) and \( \alpha \in \mathbb{N}^n \). Show that \( \text{supp}(\partial^\alpha T) \subset \text{supp}(T) \). Find a simple example where the inclusion is strict.

Exercise 27.8. Let \( T \in \mathcal{D}'(\Omega) \). Show that \( \text{supp} T \subset \Omega \) is compact if and only if \( T \) extends to a continuous linear functional on \( \mathcal{E}(\Omega) \).

Hint: Recall the characterization of \( \mathcal{E} \)-continuity given in Exercise 27.69.

Exercise 27.9. Let \( f \in C(\mathbb{R}) \subset L^1_{\text{loc}}(\mathbb{R}) \) be given by \( f(x) = |x| \). Compute \( \text{singsupp} T_f \).

Exercise 27.10. Show that the singular support of a distribution is always closed and contained in the support of the distribution.

Exercise 27.11. Let \( T \in \mathcal{D}'(\Omega) \) and \( \alpha \in \mathbb{N}^n \). Show that \( \text{singsupp}(T) = \text{singsupp}(\partial^\alpha T) \).

Exercise 27.12. Let \( T \in \mathcal{D}'(\Omega) \) and \( f \in C^\infty(\Omega) \). Show that \( \text{singsupp}(T) \subset \text{singsupp}(f T) \). Give a simple example where equality fails.

Exercise 27.13. Let \( S, T \in \mathcal{D}'(\Omega) \). Show that \( \text{singsupp}(S + T) \subset \text{singsupp} S + \text{singsupp} T \). Give a simple example where equality fails.

Exercise 27.14. Compute the singular support of regular distribution \( T_H \in \mathcal{D}'(\mathbb{R}) \) where \( H \in L^1_{\text{loc}}(\mathbb{R}) \) is the Heaviside function (cf. Example 24.1.21).

Exercise 27.15. Compute the singular support of the Cauchy principal value of \( \frac{1}{x} \) (cf. Exercise ??).

Exercise 27.16. Let \( T \in \mathcal{D}'(\Omega) \). Show that \( \text{singsupp} T = \emptyset \) if and only if \( T = T_f \) for a smooth function \( f \in C^\infty(\Omega) \).

Exercise 27.17. Give (or look up) a proof of Theorem 24.1.37

Hint: The proof of Theorem 24.1.10 might give you some idea for estimating the integrals.

Exercise 27.18. Show that for \( T \in \mathcal{D}'(\mathbb{R}^n) \) and \( \psi \in C^\infty_0(\mathbb{R}^n) \) one has

\[
\text{supp}(T \ast \psi) \subset \text{supp} T + \text{supp} \psi.
\]

Exercise 27.19. Let \( f, g \in C_0(\mathbb{R}^n) \) with \( \text{supp}(f g) \subset [-R, R]^n \) for some \( R > 0 \). For any \( N \in \mathbb{N} \), decompose \( [-R, R]^n \) into \( N^n \) cubes of same size and let \( x_1, \ldots, x_{N^n} \) denote their middle points. Show that the functions \( x \mapsto \sum_{j=1}^N f(x - x_j)g(x_j) \left( \frac{R}{N} \right)^n \) converges uniformly (in \( x \)) to \( x \mapsto (f \ast g)(x) \) as \( N \to \infty \).
Exercise 27.20. Let \( T \in D'(\mathbb{R}^n) \). Show that \( \varphi \mapsto \langle \tilde{T}, \varphi \rangle := \langle T, \varphi \rangle \) defines a distribution \( \tilde{T} \in D'(\mathbb{R}^n) \). Show further that \( \tilde{T}_f = T_f \) for any \( f \in C^\infty(\mathbb{R}^n) \).

Exercise 27.21. Let \( S \in D'(\mathbb{R}^n) \) and \( T \in \mathcal{E}'(\mathbb{R}^n) \). Show that \( \text{Exercise 27.25} \) really defines a distribution, i.e., a (sequentially) continuous linear functional on \( D(\mathbb{R}^n) \).

Exercise 27.22. Use the convolution theorem to solve the Cauchy problem for the heat equation

\[
\begin{align*}
\partial_t u(t, x) &= \Delta u(t, x) \\
u(0, x) &= u_0(x)
\end{align*}
\]

Assume that \( u_0 \in C^\infty(\mathbb{R}^n) \). Why does this only work for positive times?

Exercise 27.23. Let \( P(\xi) = \sum_{|\alpha| \leq m} c_\alpha \xi^\alpha \) be some polynomial in \( \xi \). We write \( P(D) \) for the differential operator arising from \( P \) by formally replacing every \( \xi \) by \( i\partial \). How would \( P(\xi) \) look like in order that \( P(D) = \Delta \)? Consider the first-order initial value problem

\[
\begin{align*}
\partial_t u(t, x) &= P(D)u(t, x) \\
u(0, x) &= u_0(x)
\end{align*}
\]

Using the convolution theorem formally, i.e. without worrying about any convergence or differentiability issues, find a representation formula for the solution. Can you find some conditions on \( P(\xi) \) such that this formula makes rigorous sense for \( u_0 \in \mathcal{S}(\mathbb{R}^n) \) and yields a true solution?

Exercise 27.24. Use the convolution theorem to solve the Cauchy problem for the wave equation

\[
\begin{align*}
\Box u &= 0 \\
u(0, x) &= u_0(x) \\
\partial_t u(0, x) &= u_1(x)
\end{align*}
\]

Here \( \Box = -\partial_t^2 + \Delta \) is the wave operator. You do not need to compute the inverse Fourier transforms explicitly this time.

Exercise 27.25. For any \( n \in \mathbb{N} \) define \( G_n : \mathbb{R}^n \setminus \{0\} \to \mathbb{R} \) by

\[
G_1(x) = \begin{cases} 
  x & x > 0 \\
  0 & x \leq 0
\end{cases}, \quad G_2(x) = \frac{1}{2\pi} \ln(|x|) \quad \text{and} \quad G_n(x) = -\frac{|x|^{2-n}}{(n-2)c_n} \quad \text{for} \quad n \geq 3.
\]

Interpret \( G_n \) as distribution (how?) and show that \( \Delta_n G_n = \delta \) in the sense of distributions, where \( \Delta_n = \sum_{j=1}^n \partial_j^2 \) denotes the Laplace operator and \( \delta \) the \( \delta \)-distribution at the origin.

Exercise 27.26. For \( k > 0 \), define \( G_H : \mathbb{R}^3 \setminus \{0\} \to \mathbb{C} \) by

\[
G_H(x) = \frac{1}{4\pi|x|} e^{ik|x|}.
\]

Interpret \( G_H \) as distribution (how?) and show that \((\Delta_3 + k^2)G_H = -\delta \), where \( \Delta_3 = \partial_1^2 + \partial_2^2 + \partial_3^2 \) is the Laplace operator and \( \delta \) the \( \delta \)-distribution at the origin. Show further that \( G_H \to G_3 \) for \( k \to \infty \) in the sense of distributions, where \( G_3 \) is the Green’s function of the Laplace equation in three dimensions from Exercise 27.25.

Exercise 27.27. For some parameter \( \lambda > 0 \) define \( G_W : \mathbb{R}^n \times (\mathbb{R} \setminus \{0\}) \to \mathbb{R} \) by

\[
G_W(x, t) := \frac{1}{(4\pi\lambda t)^{\frac{n}{2}}} e^{-\frac{|x|^2}{4\lambda t}} \Theta(t).
\]

Interpret \( G_W \) as distribution on \( \mathbb{R}^{n+1} \) and show that \((\partial_t - \lambda \Delta)G_W = \delta \). Compute the singular support of \( G_W \).
Exercise 27.28. For some constant $c \in \mathbb{R}$ define $G_S : \mathbb{R}^n \times (\mathbb{R} \setminus \{0\}) \to \mathbb{C}$ by

$$G_S(x,t) = \frac{c}{t^2} e^{-\frac{|x|^2}{t^2}} \Theta(t).$$

1. Show that $G_S$ is locally integrable if and only if $n = 1$.
2. Show that

$$\langle G_S, \phi \rangle := \lim_{r \to 0} \int_0^\infty \int_{\mathbb{R}^n} \frac{c}{t^2} e^{-\frac{|x|^2}{t^2}} \phi(x,t) \, dx \, dt$$

defines a distribution.
3. Determine $c$ such that $(i\partial_t + \Delta)G_S = \delta$ in the sense of distributions.
4. Compute the singular support of $G_S$.

Exercise 27.29. Define $I : \mathbb{R}^2 \setminus \{0\} \to \mathbb{C}$ by

$$I(x,y) = \frac{1}{\pi(x+iy)}.$$

Explain how $I$ may be understood as distribution on $\mathbb{R}^2$. Show that $\frac{1}{2}(\partial_x + i\partial_y)I = \delta$ holds in the sense of distributions.

Exercise 27.30. Show that

$$\langle G^\pm, \phi \rangle := \int_{\mathbb{R}^3} \frac{1}{4\pi|x|} \phi(\pm|x|, x) \, d^3x$$

defines distribution $G^\pm \in \mathcal{D}'(\mathbb{R}^3)$ (pay attention to dimensions here). Compute the support and singular support of $G^\pm$. Show that $(\partial_t^2 - \Delta)G^\pm = \delta$ holds in the sense of distributions.

Exercise 27.31 (Spectral measure of a hermitean matrix). Verify that the function defined by (24.2.2) satisfies the properties of a spectral measure. Then verify that for this measure, the integral (24.2.1) coincides with (24.2.2).

Exercise 27.32. Let $A : \mathcal{H} \to \mathcal{H}$ be linear. Show that the following are equivalent:

1. $A$ is continuous.
2. $A$ is continuous in $0$.
3. $\exists C > 0 : \forall x \in \mathcal{H} : \|Ax\| \leq C \cdot \|x\|.$

Exercise 27.33. Prove Theorem 24.2.5.

Exercise 27.34. Construct a bounded operator $A$ on $\ell^2$ which is injective but not surjective. Conclude that not every element of $\sigma(A)$ is an eigenvalue of $A$.

Exercise 27.35 (Neumann series). Let $A \in L(\mathcal{H})$ be given. Show that if $\sum_{k=0}^\infty A^k$ converges with respect to the operator norm on $L(\mathcal{H})$, then the limit is the inverse to $\text{id}_\mathcal{H} - A$. Moreover, show that for $\|A\| < 1$ the series always converges.

Exercise 27.36. Let $A \in L(\mathcal{H})$ be given. Show that the resolvent $R(\lambda) = (\lambda \text{id}_\mathcal{H} - A)$ can be expressed as a power series in $\lambda$ around any point $\lambda_0 \in \rho(A)$.

Hint: Use the Neumann series (cf. Exercise 27.35).

Exercise 27.37. Fill in the details in the proof of Theorem 24.2.7.

Exercise 27.38. Let $A, B \in L(\mathcal{H})$ and $\alpha, \beta \in \mathbb{C}$.

1. Show that $(\alpha A + \beta B)^* = \bar{\alpha} A^* + \bar{\beta} B^*$. Moreover, show that $(A^*)^* = A$.
2. Show that $\|A^*\|_\infty = \|A\|_\infty$ and $\|AA^*\|_\infty = \|A\|_\infty^2$.

Exercise 27.39. Let $\mathcal{H}$ be a complex Hilbert space and $A \in L(\mathcal{H})$. Show that $A$ is selfadjoint if and only if $\langle Ax, x \rangle \in \mathbb{R}$ holds for all $x \in \mathcal{H}$.

Exercise 27.40. Prove Proposition 24.2.9.
Exercise 27.41. Let $E : \mathcal{A} \to L(H)$ be a spectral measure. Verify that $E_u$ and $E_{u,v}$ defined by (24.2.11) and (24.2.12) are positive, respectively complex measures for any given $u, v \in H$. Moreover, verify the following properties:

1. $4E_{u,v} = E_{u+v} + iE_{u+i-v} - E_{u-v} - iE_{u-i-v}$
2. $E_{u,v} = \frac{E_u E_v}{E_{u+v}}$
3. $|E_{u,v}(A)|^2 \leq E_u(A) E_v(A)$
4. If $A_1, A_2, \ldots$ are pairwise disjoint, then $\sum_{n=1}^{\infty} |E_{u,v}(A_n)| \leq \sqrt{E_u(A)} \sqrt{E_v(A)}$

Exercise 27.42. Let $(\Omega, \mathcal{A})$ be a measurable space and let $E : \mathcal{A} \to L(\mathbb{C})$ be a spectral measure. Show that $E$ is a multiple of the Dirac measure at some point.

Exercise 27.43. Prove Theorem (24.2.12).

Exercise 27.44. Deduce uniqueness of the map $T_E$ in Theorem (24.2.15) from part (2) and the Riesz representation theorem for Hilbert spaces.

Exercise 27.45. Let $E : \mathcal{A} \to L(H)$ be a spectral measure and $T := E(x) \in L(H)$. For a polynomial $p(x) = \sum_{j=0}^{n} a_n x^n$ with coefficients $a_0, \ldots, a_n \in \mathbb{C}$ define the operator $p(T) \in L(H)$ by $p(T) := \sum_{j=0}^{n} a_n T^n$. Show that $p(T) = T_E(p) = \int_{\Omega} p(x) \, dE(x)$.

Exercise 27.46. Prove Corollary (24.2.16).

Exercise 27.47. Let $T \in L(H)$ be selfadjoint, and let $p$ be a polynomial with real coefficients. Show that $p(T)$ is selfadjoint again.

Exercise 27.48. Carry out the details of those steps in the proof of Theorem (24.2.21) where above it was simply stated that they work by “approximating by polynomials”.

Exercise 27.49. Show that the measurable functional calculus is also involutive, i.e., show that $f(T)^* = \overline{f(T)}$ holds for any $f \in \mathcal{L}(\sigma(T))$.

Exercise 27.50. Fill in the details of the previous proof.

Exercise 27.51. Give an example of a selfadjoint operator $T$ and a bounded, measurable function $f : \sigma(T) \to \mathbb{C}$ with $\sigma(f(T)) \neq f(\sigma(T))$.

Exercise 27.52. Give an example of a selfadjoint operator $T$ and a bounded, measurable function $f : \sigma(T) \to \mathbb{C}$ with $\|f(T)\|_{\infty} \neq \|f\|_{\infty}$, where $\|f\|_{\infty}$ denotes the (actual) supremum (not the essential supremum).

Exercise 27.53. Let $T \in L(H)$ be selfadjoint and let $E : B(\sigma(T)) \to L(H)$ be the corresponding spectral measure. Show that for any $B \in L(H)$ one has $[T, B] = 0$ if and only if $[E(U), B] = 0$ for all $U \in B(\sigma(T))$. Show moreover that in this case also $[f(T), B] = 0$ for all bounded measurable functions $f : B(\sigma(T)) \to \mathbb{C}$.

Exercise 27.54 (Existence of smooth functions with compact support). Show that the function $f : \mathbb{R}^n \to \mathbb{R}$ defined by

\[
    f(x) := \begin{cases} 
        e^{\frac{-1-|x|^2}{1-|x|^2}} & |x| < 1 \\
        0 & |x| \geq 1 
    \end{cases}
\]

is smooth and has compact support.

Hint: First deal with the case $n = 1$. Show smoothness in this case by induction on the order of differentiation.

Exercise 27.55. Let $f \in \mathcal{E}(\mathbb{R}^n)$. For any $n \in \mathbb{N}$ define $f_n \in \mathcal{E}(\mathbb{R}^n)$ by $f_n(x) := f(\frac{x}{n})$. Show that $(f_n)_n$ $\mathcal{E}$-converges to a constant function. If additionally $f \in \mathcal{D}(\mathbb{R}^n)$ and $f \neq 0$ hold, show that $f_n \in \mathcal{D}(\mathbb{R}^n)$ holds as well but this sequence does not converge in $\mathcal{D}(\mathbb{R}^n)$.

Exercise 27.56. Show that the following are equivalent ways of saying that $u \in C^\infty(\mathbb{R}^n)$ is a Schwartz function:
(1) \( \forall \alpha, \beta \in \mathbb{N}^n : \sup_{x \in \mathbb{R}^n} |(x)^\alpha \partial^\beta u(x)| < \infty. \)
(2) \( \forall \alpha, \beta \in \mathbb{N}^n : \sup_{x \in \mathbb{R}^n} |\partial^\alpha((x)^\beta u(x))| < \infty. \)
(3) \( \forall \alpha, \beta \in \mathbb{N}^n \exists C > 0 \forall x \in \mathbb{R}^n : |(x)^\alpha \partial^\beta u(x)| < C. \)
(4) \( \forall \alpha, \beta \in \mathbb{N}^n \exists C > 0 \forall x \in \mathbb{R}^n : |\partial^\alpha((x)^\beta u(x))| < C. \)
(5) \( \forall \alpha, \beta \in \mathbb{N}^n \exists C > 0 \forall x \in \mathbb{R}^n : |\partial^\alpha u(x)| < \frac{C}{(x^\alpha)^\beta}. \)

Here \( (x) := (1 + |x|^2)^{\frac{1}{2}} \) for any \( x \in \mathbb{R}^n. \)

**Exercise 27.57.** Carry out the details left out in Example 24.1.3.

**Exercise 27.58.** Prove Lemma 24.1.3 for the case \( p = \infty. \)

**Exercise 27.59.** (Moderately growing functions). A function \( f \in C^\infty(\mathbb{R}^n) \) is said to be moderately growing if the following holds:

\[
\forall \alpha \in \mathbb{N}^n \exists N \in \mathbb{N} : \sup_{x \in \mathbb{R}^n} |(x)^{-N} \partial^\alpha f(x)| < \infty. \tag{27.5.2}
\]

We denote by \( \mathcal{O}_M(\mathbb{R}^n) \) the set of moderately growing functions. Show that \( \mathcal{O}_M(\mathbb{R}^n) \) is a vector space which contains all polynomials, all Schwartz functions, and all smooth functions with compact support. Is a moderately growing function also a Schwartz function? Give an example of a smooth function which does not belong to \( \mathcal{O}_M(\mathbb{R}^n). \)

**Exercise 27.60.** Show that \( T : \mathcal{E}(\mathbb{R}) \to \mathbb{C}, T(u) := u(0) + u'(1) \) is a distribution.

**Exercise 27.61.** Show that \( I : \mathcal{S}(\mathbb{R}) \to \mathbb{R}, I(u) = \int_{\mathbb{R}} u(x) \, dx \) defines a tempered distribution. Why does not I define a distribution in \( \mathcal{E}'(\mathbb{R}). \) Show that for any \( [a, b] \subset \mathbb{R}, \) the assignment \( u \mapsto \int_a^b u(x) \, dx \) defines an element of \( \mathcal{E}'(\mathbb{R}). \)

(Remark: The general case follows by approximation.)

**Exercise 27.62.** (Fundamental calculus in variations of continuous case). Show that any continuous function on \( \Omega \) is locally integrable. Suppose now that \( T_f = T_g \) for two continuous functions \( f, g \in C(\Omega). \) Show that \( f = g. \)

**Exercise 27.63.** Show also that \( T_f \in \mathcal{E}'(\Omega) \) if \( \text{supp } f \) is compact, and \( T_f \in \mathcal{S}'(\Omega) \) if \( f \in \mathcal{O}_M(\mathbb{R}^n) \) is of moderate growth.

**Exercise 27.64.** (Integrability of inverse powers). Let \( \alpha > 0 \) and let \( f : \mathbb{R}^n \to \mathbb{R} \) be defined by \( f(x) = |x|^{-\alpha} \) for \( x \neq 0 \) and \( f(0) = 0. \) Find out for which \( \alpha > 0 \) one has \( f \in L^1_{\text{loc}}(\mathbb{R}^n). \)

**Exercise 27.65.** (Logarithmic poles are integrable). Consider the function \( f : \mathbb{R} \to \mathbb{R} \) defined as \( f(x) = \log |x| \) for \( x \neq 0 \) and \( f(0) = 0. \) Show that \( f \in L^1_{\text{loc}}(\mathbb{R}). \)

**Exercise 27.66.** (An identity for the Cauchy principal value of \( \frac{1}{x} \)). Show that for any \( u \in \mathcal{D}(\mathbb{R}) \) and any even function \( \varphi \in L^1(\mathbb{R}) \cap C^1(\mathbb{R}) \) with \( \varphi(0) = 1 \) one has

\[
(\text{v.p.} \left( \frac{1}{x} \right), u) = \int_{\mathbb{R}} \frac{u(x) - u(0)\varphi(x)}{x} \, dx,
\]

where the right-hand side is to be understood as an honest integral (i.e., without any further manipulations). First explain why this integral exists in the first place.

**Exercise 27.67.** Fill in the details in example 24.1.11. Show furthermore that the \( \delta \)-distribution and also any regular distribution is a distribution of the type described in this example.

**Exercise 27.68.** For a smooth function \( u \in C^\infty(\mathbb{R}), \) denote by \( u^{(k)} \in C^\infty(\mathbb{R}) \) its \( k \)-the derivative. Show that

\[
T(u) := \sum_{k=0}^{\infty} u^{(k)}(k) \quad \forall u \in \mathcal{D}(\mathbb{R}) \tag{27.5.3}
\]

defines a distribution \( T \in \mathcal{D}'(\mathbb{R}). \) Is \( T \) also in \( \mathcal{S}'(\mathbb{R}) \) or \( \mathcal{E}'(\mathbb{R})? \)
Exercise 27.69. Show that a linear operator \( A : \mathcal{E}(\Omega) \to \mathbb{C} \) is sequentially continuous if and only if it is sequentially continuous at 0 if and only if there exist a compact subset \( K \subset \Omega \) and a constant \( C > 0 \), \( m \in \mathbb{N} \) such that
\[
\forall u \in \mathcal{E}(\Omega) : |Au| \leq C \sum_{|\alpha| \leq m} \| \partial^\alpha u \|_{L^\infty(K)} .
\] (27.5.4)

Exercise 27.70. Show that a linear operator \( A : \mathcal{S}(\mathbb{R}^n) \to \mathbb{C} \) is sequentially continuous if and only if it is sequentially continuous at 0 if and only if there exist \( m \in \mathbb{N} \) and a constant \( C > 0 \) such that
\[
\forall u \in \mathcal{S}(\mathbb{R}^n) : |Au| \leq C \sum_{|\alpha|,|\beta| \leq m} \| x^\alpha \partial^\beta u \|_{L^\infty(\mathbb{R}^n)} .
\] (27.5.5)

Exercise 27.71 (Model \( \delta \)-sequences). Let \( \psi \in \mathcal{C}^\infty(\mathbb{R}^n) \cap L^1(\mathbb{R}^n) \) with \( \int_{\mathbb{R}^n} \psi(x) \, dx = 1 \). For any \( k \in \mathbb{N} \), define \( \psi_k \in \mathcal{C}^\infty(\mathbb{R}^n) \) by
\[
\psi_k(x) := k^n \psi(kx) .
\] (27.5.6)

Show that \( (\psi_k)_k \) is a \( \delta \)-sequence. Moreover, show that if \( \text{supp} \psi \) is compact, then \( (\psi_k)_k \) is a strict \( \delta \)-sequence.

Remark: \( \delta \)-sequences of this type are called model \( \delta \)-sequences or mollifiers. If additionally \( \psi \geq 0 \), one speaks of a positive mollifier and if \( \psi(x) = \varphi(|x|) \) for some \( \varphi \in \mathcal{C}^\infty_0(\mathbb{R}^+) \), one speaks of a symmetric mollifier.

Exercise 27.72. Show that the following sequences are all delta sequences on \( \mathbb{R} \).

1. The Gaussians \( \psi_k(x) = ke^{-\pi k^2 x^2} \).
2. The Lorentz curves \( \psi_k(x) = \frac{1}{\pi(1+k^2 x^2)} \).
3. The Fejer kernel \( \psi_k(x) = \frac{k}{\pi} \left( \frac{\sin kx}{kx} \right)^2 \).

You may skip showing the normalization (this is where the constants come from).

Exercise 27.73. Show sequential continuity also in those cases of Theorem 24.1.19 which were not explicitly treated in the proof.

Exercise 27.74 (Oscillatory integrals). Let \( a, \phi \in \mathcal{C}^\infty(\mathbb{R}^n_x \times \mathbb{R}^N_\xi) \) satisfy the following assumption:

A) \( \phi \) is a phase function, i.e. \( \phi \) is real-valued and satisfies \( \phi(x, \lambda \xi) = \lambda \phi(x, \xi) \) for all \( \lambda > 0 \) and all \( (x, \xi) \in \mathbb{R}^n_x \times \mathbb{R}^N_\xi \) (positive homogeneity of degree one in \( \xi \)).

B) \( a \) is a symbol of degree \( m > 0 \) for some \( m \in \mathbb{N} \), i.e. for any \( \alpha \in \mathbb{N}^n \) and any \( \beta \in \mathbb{N}^N \) there exists a constant \( C_{\alpha, \beta} > 0 \) such that
\[
| \partial_\xi^\beta \partial_\xi^\alpha a(x, \xi) | \leq C_{\alpha, \beta} (1 + |\xi|)^{m-|\alpha|} \quad \forall (x, \xi) \in \mathbb{R}^n_x \times \mathbb{R}^N_\xi .
\] (27.5.7)

C) \( \phi \) has no critical points on the support of \( a \).

For such functions, consider the so-called oscillatory integrals
\[
u(x) = \int e^{i\phi(x, \xi)} a(x, \xi) \, d\xi ,
\] (27.5.8)
which in general is only well-defined if \( m < -n + 1 \). (Why?) Show that \( u \) can however always be interpreted as a tempered distribution in the following way:

1. Show that for every \( \epsilon > 0 \) the integral
\[
u_\epsilon(x) = \int e^{i\phi(x, \xi)} a(x, \xi) e^{-\epsilon \frac{|\xi|^2}{2}} \, d\xi
\] (27.5.9)
is well-defined and defines a bounded, smooth function \( u_\epsilon \in \mathcal{C}^\infty(\mathbb{R}^n_x) \).

2. Show that \( u_\epsilon \) interpreted as (regular) tempered distribution converges in \( \mathcal{S}' \) as \( \epsilon \to 0 \).
Hint: Notice that $e^{i\phi(x,\xi)} = (i\partial_x \phi(x,\xi))^{-1} \partial_x e^{i\phi(x,\xi)}$. Use this and partial integration to show that for $\varphi \in S(\mathbb{R}^n)$ one can write

$$\int u(x)\varphi(x)\,dx = \int_{\mathbb{R}^n} e^{-\epsilon|x|^2} \left( \int_{\mathbb{R}^n} e^{i\phi(x,\xi)} \mathcal{L}(a(x,\xi)\varphi(x))\,dx \right)\,d\xi,$$

where $\mathcal{L}$ is some differential operator. Notice that one can choose operators $\mathcal{L}$ of arbitrary high order. Use this and the symbol property of $\phi$ to show that in such a way one can obtain a nicely decaying integrand such that the limit $\epsilon \to 0$ exists.

Exercise 27.75. Show that (24.1.14) and (24.1.15) do in fact define distributions again. Cover the different cases $\mathcal{D}', \mathcal{E}'$ and $S'$.

Exercise 27.76. Show that (24.1.15) is also compatible with viewing smooth functions as distributions.

Exercise 27.77. Compute the distributional derivative of the absolute value function $|\cdot| \in L^1_{\text{loc}}(\mathbb{R}) \subset \mathcal{D}'(\mathbb{R})$. Show that it is again a regular distribution.

Exercise 27.78. For an interval $[a, b] \subset \mathbb{R}$, let $\mathbf{1}_{[a,b]} \in L^1_{\text{loc}}(\mathbb{R})$ be the corresponding characteristic function. Compute its derivative as a distribution.

Exercise 27.79. By Exercise 27.67, $\log |x|$ is locally integrable on $\mathbb{R}$. Show that its distributional derivative is the Cauchy principal value $\mathbf{v.p.}(\frac{1}{x})$, cf. Exercise 15.

Exercise 27.80. Show that $x \cdot \mathbf{v.p.}(\frac{1}{x}) = 1$, where $\mathbf{v.p.}(\frac{1}{x})$ denotes the Cauchy principal value distribution (cf. exercise 15).

Exercise 27.81 (Iterative application of derivatives and multiplication). Let $T \in \mathcal{E}'(\Omega)$, $\alpha, \beta \in \mathbb{N}^n$ and $f, g \in C^\infty(\Omega)$. Show that $\partial^\alpha (\partial^\beta T) = \partial^{\alpha+\beta} T$ and $f g T = (f g) T$.

Exercise 27.82 (Leibniz rule for distributions). Either let $f \in C^\infty(\Omega)$ and $T \in \mathcal{D}'(\Omega)$ or $T \in \mathcal{E}'(\Omega)$, or let $f \in \mathcal{O}_M(\mathbb{R}^n)$ and $T \in S'(\mathbb{R}^n$). Show that the Leibniz rule

$$\partial^\alpha (f T) = \sum_{\beta \leq \alpha} \binom{\alpha}{\beta} (\partial^{\alpha - \beta} f)(\partial^\beta T)$$

(27.5.10)

holds for any $\alpha \in \mathbb{N}^n$.

Exercise 27.83. Let $\alpha \in \mathbb{N}^n$, $f \in C^\infty(\Omega)$ and $A = \sum_{|\alpha| \leq m} a_\alpha \partial^\alpha$. Show that the maps

$$u \mapsto \partial^\alpha u, \quad u \mapsto f u, \quad u \mapsto A u$$

are linear and sequentially continuous on all introduced test function and distribution spaces. In the tempered case, take $\Omega = \mathbb{R}^n$ and assume that $f$ and the coefficients of $A$ are all of moderate growth.

Using the sequential continuity on the test function spaces, give another (very quick) argument why (24.1.14) and (24.1.15) do indeed define distributions again.

Exercise 27.84. For $n \in \mathbb{N}$ let $G_n : \mathbb{R}^n \to \mathbb{R}$ be defined by $G_n(x) := x\Theta(x)$ with $\Theta$ being the Heaviside function, $G_2(x) := \frac{1}{2} \ln|x|$ and $G_n(x) := -\|x\|^{n-2} n c_n^{-1}$ for $n \geq 3$, where $c_n = |S^{n-1}| = (2\pi)^{\frac{n}{2}} \Gamma\left(\frac{n}{2}\right)$ denotes the volume of the $(n-1)$-dimensional unit sphere. To be complete, set $G_n(0) := 0$. Show that $G_n$ is locally integrable on $\mathbb{R}^n$ for every $n \in \mathbb{N}$ and is a Green’s function for the Laplace operator at $x_0 = 0$. Use $G_n$ to also find Green’s functions for the Laplace operator at an arbitrary point $x_0 \in \mathbb{R}^n$.

Exercise 27.85. For $k > 0$, define $G_H : \mathbb{R}^3 \to \mathbb{C}$ by $G^H_k(x) = -\frac{\exp(ik|x|)}{4\pi|x|^2}$, and $G_H(0) = 0$. Show that $G_H$ is locally integrable and a Green’s function for the so-called Helmholtz operator $H_k := \Delta + k^2$. Further show that $\lim_{k \to 0} G^H_k = G_3$ in $\mathcal{D}'(\mathbb{R}^3)$, where $G_3 \in \mathcal{D}'(\mathbb{R}^3)$ denotes the Green’s function of the Laplace operator defined in exercise 27.83.
Exercise 27.86. Define $G_W : \mathbb{R}^n \times \mathbb{R} \to \mathbb{R}$ by $G_W(x,t) := (4\pi t)^{-\frac{n}{2}} \exp(-\frac{|x|^2}{4t})\Theta(t)$, where $\Theta$ is the Heaviside function ($G_W(x,0) := 0$). Show that $G_W$ is locally integrable on $\mathbb{R}^n \times \mathbb{R}$ and a Green’s function of the heat operator $W = \partial_t - \Delta$.

Exercise 27.87. Define $G_S : \mathbb{R}^n \times \mathbb{R} \to \mathbb{R}$ by $G_S(x,t) := c_n t^{-\frac{n}{2}} \exp(-\frac{|x|^2}{4t})\Theta(t)$, where $\Theta$ is the Heaviside function and $c_n$ a constant ($G_S(x,0) := 0$). Show that $G_S$ is locally integrable on $\mathbb{R}^n \times \mathbb{R}$ if and only if $n = 1$. Show that

$$\langle \hat{G}_S, u \rangle := \lim_{\epsilon \to 0} \int_0^\infty \int_{\mathbb{R}^n} G_S(x,t)u(x,t) \, dx \, dt$$

defines a distribution $\hat{G}_S \in \mathcal{D}'(\mathbb{R}^n \times \mathbb{R})$. Show further that this distribution is a Green’s function of the Schrödinger operator $S := i\partial_t + \Delta$ if $c_n$ is chosen appropriately (how?).

Exercise 27.88. Define $G_{\overline{\sigma}} : \mathbb{R}^2 \to \mathbb{C}$ by $G_{\overline{\sigma}}(x,y) := \frac{1}{\pi(x+y)^2}$ and $G_{\overline{\sigma}}(0) = 0$. Show that $G_{\overline{\sigma}}$ is locally integrable on $\mathbb{R}^2$ and a Green’s function of the operator $\overline{\sigma} := \frac{1}{2}(\partial_x + i\partial_y)$.

Hint: By Green’s theorem one has $\int_U \overline{\sigma}u(x,y) \, dx \, dy = 1 \int_{\partial U} u(x,y)(dx + i \, dy)$.

Exercise 27.89. Define distributions $G_{\Box}^\pm \in \mathcal{D}'(\mathbb{R} \times \mathbb{R}^3)$ by

$$\langle G_{\Box}^\pm, u \rangle := \int_{\mathbb{R}^3} \frac{1}{4\pi|x|}u(\pm|x|,x) \, d^3x.$$

Show that this indeed defines distributions and show that $G_{\Box}^\pm$ are Green’s function of the wave operator $\Box := \partial_t^2 - \Delta$.

Exercise 27.90 (Fourier transform of a Gaussian). Compute the Fourier transform of the Gaussian $f(x) = \exp(-\frac{1}{2}|x|^2)$ by first reducing to $n = 1$ and then deriving a well-known ordinary differential equation satisfied by $\hat{f}(\xi)$.

Exercise 27.91. Prove the properties of the Fourier transform on $S(\mathbb{R}^n)$ as stated in Theorem 24.1.20. Use the rapid decay of Schwartz functions to estimate integrals. Concerning the Fourier inversion, plug in the functions $g_\epsilon(\xi) = e^{i\epsilon x}e^{-|\xi|^2}$ for $\epsilon > 0$ into (24.1.21) and examine what happens to both sides in the limit $\epsilon \to 0$. The Fourier transform of $g_\epsilon$ can be computed using exercise 27.90 by completing squares.

Exercise 27.92. Prove Theorem 24.1.28 about the Fourier transform of tempered distributions.

Exercise 27.93. Compute the Fourier transform of the $\delta$-distribution at a general point $x_0 \in \mathbb{R}^n$ and show that it is a regular distribution.

Exercise 27.94. Compute the distributional Fourier transform of the (bounded but not integrable) functions $e^{ix}$, $\sin(x)$ and $\cos(x)$.

Exercise 27.95. For $\sigma \in \mathbb{R}$ given, compute the distributional Fourier transform of the (bounded but not integrable) “complex Gaussian function” $u(x) = \exp(-i\sigma |x|^2)$.

Exercise 27.96. Compute the distributional Fourier transform of the Heaviside function $\theta$, cf. exercise 24.1.21 using the following trick: First compute the Fourier transform of $f_\epsilon(x) := \theta(x)e^{-\epsilon x}$ for $\epsilon > 0$ and then examine the limit $\epsilon \to 0$.

Exercise 27.97. Compute the Fourier transform of the Cauchy principal value distribution, cf. example 24.1.15.

Exercise 27.98. Compute the distributional Fourier transform of the the logarithm (which is locally integrable and of moderate growth at infinity).

Exercise 27.99. Let $T \in \mathcal{S}'(\mathbb{R}^n)$. 

To that purpose, proceed as follows: (where operator products are defined as in the lecture). Use the uniqueness of solutions of the
**\( y \)**
show that the function
\[
 s \quad \text{and check why this construction is unique. To this end, make the ansatz}
\]
causal Green's functions of the Dirac equation (**344 REMAINING EXERCISES**
\[
(\psi + \mathcal{B} - m) \psi(x) = 0 \quad \text{in an external field,}
\]
\[
\hat{s}^\vee_m = \sum_{n=0}^{\infty} (-s_m^\vee \mathcal{B})^n s^\vee_m , \quad \hat{s}^\wedge_m = \sum_{n=0}^{\infty} (-s_m^\wedge \mathcal{B})^n s^\wedge_m .
\]
and check why this construction is unique. To this end, make the ansatz
\[
\hat{s}^\vee_m = \sum_{n=0}^{\infty} s^{(n)} ,
\]
where **s**\(^{(n)}\) is exactly of order **B**\(^n\) and show that \( \forall n > 1 \)
\[
(i\phi - m) s^{(n)} = -\mathcal{B} s^{(n-1)}
\]
(where operator products are defined as in the lecture). Use the uniqueness of solutions of the homogeneous Dirac equation to show that, if we require
\[
\text{supp } \hat{s}^\vee_m(x, \cdot) \subset J^\vee_\gamma(x),
\]
then **s**\(^{(0)}\) = \( \hat{s}^\vee_m \) and \( s^{(n)} = -s_m^\vee \mathcal{B} s^{(n-1)} \) hold uniquely.
EXERCISE 27.104. Recall the definition of the Hilbert space $\mathcal{H}_t$ and of the time evolution operator $U^{t,t'} : \mathcal{H}_t \to \mathcal{H}_{t'}$ from the lecture. Perform a partial Fourier transformation to show that

$$ U^{t,t'}(\vec{k}) = \int_{-\infty}^{\infty} (\vec{k} + m) \gamma^0 \delta(k^2 - m^2) |_{k=(\omega, \vec{k})} e^{i\omega(t-t')} \, d\omega. $$

Carry out the $\omega$ Integration to obtain

$$ U^{t',t}(\vec{k}) = \sum_{\pm} \Pi_{\pm}(\vec{k}) e^{\mp i\omega(t-t')}, $$

where $\Pi_{\pm}(\vec{k})$ are projection operators and

$$ \omega(\vec{k}) = \sqrt{|\vec{k}|^2 + m^2} \quad \text{and} \quad k_{\pm} = (\pm\omega(\vec{k}), \vec{k}). $$

Use this to conclude that $U^{t,t'}$ is unitary. (Note that it is part of the exercise to show that the $\Pi_{\pm}(\vec{k})$ are projection operators.)

EXERCISE 27.105. From the lecture you know the fermionic projector of the vacuum in momentum space

$$ P_m(k) = (\vec{k} + m) \delta(k^2 - m^2) \Theta(-k^0). $$

In this exercise, you will calculate its (inverse) Fourier transform,

$$ P_m(x,y) = \int \frac{d^d k}{(2\pi)^d} P_m(k) e^{-i k \cdot (x-y)}, $$

which is the (kernel of) the fermionic projector of the vacuum in “position space”. To do so,

1. calculate the inverse Fourier transform of $\delta(k^2 - m^2) \Theta(-k^0)$,

$$ T_{m^2}(x,y) := \int \frac{d^d k}{(2\pi)^d} \delta(k^2 - m^2) \Theta(-k^0) e^{-i k \cdot (x-y)} $$

using Bessel functions.

Hint: Use spherical coordinates for $\xi = (y-x)$ and $k$ in “spatial directions” (i.e., use 3-dim. spherical coordinates and keep the $t / k^0$-component). Carry out the integral over the angles (as compared to time- and radial coordinates) first. Then, carry out the $k^0$ integration and look up the resulting formula in a textbook or website which contains information about Fourier transforms of Bessel functions (possibly you might have to rearrange the integral s.t. you integrate $(-\infty, \infty)$ rather than $(0, \infty)$). It might prove helpful to set $r=0$ and $t=0$ before integrating, respectively.

If you have trouble revisit GRIPS at the end of the week to see if more hints for this exercise have been added.

2. Expand the resulting Bessel functions to obtain the formula for $T_{m^2}$ which you have seen in the lecture.

3. Argue why $P_m(x,y) = (i\partial_x + m) T_{m^2}$ and calculate $P_m(x,y)$ by applying the derivative.

Ps: This exercise is not so easy. Don’t get your hair teared ;)

EXERCISE 27.106. (Causal Green’s functions of the Dirac equation) In the lecture you were introduced to the Dirac equation and its the causal Green’s functions

$$ s^\vee_m(k) = \lim_{\epsilon \searrow 0} \frac{\vec{k} + m}{k^2 - m^2 - i\epsilon k^0}, \quad \text{and} \quad s^\wedge_m(k) = \lim_{\epsilon \searrow 0} \frac{\vec{k} + m}{k^2 - m^2 + i\epsilon k^0}, $$

as well as to the Feynman propagator

$$ s^F_m(k) := \lim_{\epsilon \searrow 0} \frac{\vec{k} + m}{k^2 - m^2 + i\epsilon}. \quad (27.5.12) $$
(1) Argue that these limits exist as distributions, and show that if one takes the inverse Fourier transform of $s_m^\gamma(x)$ and $\hat{s}_m^\gamma(k)$, 

$$s_m(x,y) = \int \frac{d^4k}{(2\pi)^4} s_m(k) e^{ik(x-y)},$$

one obtains

$$\text{supp } s_m^\gamma(x,.) \subset J_x^\gamma$$

$$\text{supp } \hat{s}_m^\gamma(x,.) \subset J_x^\gamma,$$

i.e. $s_m^\gamma/\hat{s}_m^\gamma(x,y)$ are causal. Give a brief interpretation of what this means.

(2) Show that the Fourier transform of the Feynman propagator is

$$s_m^\gamma(x,y) = \frac{i}{(2\pi)^4} \int \mathbb{R}^3 (k + m) e^{-ik(x-y)} \bigg|_{k=(\epsilon(t) \omega(\vec{p}), \vec{p})} d\mu_{\vec{p}}$$

where $t := (y - x)^0$, $\omega(\vec{p}) := \sqrt{|\vec{p}|^2 + m^2}$ and $d\mu_{\vec{p}} := \frac{d^3p}{2\pi|\vec{p}|}$ ($\epsilon$ denotes the sign function $\epsilon(x) = 1$ for $x \geq 0$ and $\epsilon(x) = -1$ otherwise).

Try to interpret the result: Which kind of frequencies appear for positive times, which kind for negative times? What could this mean physically?

Hints: ad i) You can show this by proving first that $s_m^\gamma(x - y)$ is 0 if $(x^0 - y^0) > 0$ and then invoking Lorentz invariance to show that $\text{supp } s_m^\gamma(x,.) \subset J_x^\gamma$. (The latter argument can be given in words.) But feel free to look for a different solution! Use residues to calculate the integrals!

**Exercise 27.107. (The Dyson series (5+1 points))** In this exercise, we will study the so-called Dyson series and the ordered exponential. We will take a look at how both appear in quantum field theory. In exercise 4.3 you will then see how this appears in the context of the light cone expansion of the fermionic projector.

In quantum field theory, one is given some Hamiltonian $H$ (a linear though, unfortunately, not bounded operator on some Hilbert (Fock) space $\mathcal{H}$). It splits up into a “free” part $H_0$ and an interaction part $H_I$:

$$H = H_0 + H_I.$$  

What one is interested in is the time evolution operator belonging to $H$, i.e. $U(t) = e^{-iHt}$. Usually, the time evolution operator belonging to $H_0$, i.e. $U_0(t) = e^{-iH_0t}$, has been understood to a sufficient extend, and we therefore assume that it is known. What remains to be understood is the “effect” of $H_I$. To this end, one studies the operator $V(t) := U_0(-t)U(t)$.

(1) Show that $V(t)$ satisfies the differential equation

$$\frac{dV(t)}{dt} = -iH_I(t)V(t),$$

where $H_I(t) = U_0(-t)H_1U_0(t)$. (Note that $H_0$ and $H_I$ may not commute.)

(2) Re-express this differential equation as an integral equation taking the form $V(t) = \mathbb{1} + \int_0^t \ldots$

(3) On the right-hand-side of your integral equation, there is a factor $V(\tau)$ (in case you denoted your integration variable by $\tau$). You can now find an expression for $V(t)$ up to a given order in $H_I$ by plugging in the integral equation into itself several times. Write down the expression up to second order in $H_I$.

(4) Generalize this to arbitrary order in $H_I$, i.e. show that $V(t) \doteq \mathbb{1} + \sum_{n=1}^{\infty} V_n(t)$ for a suitable $V_n(t)$ of the form

$$V_n(t) = \frac{1}{n!} \int_0^t \int_0^{\tau_n} \ldots \int_0^{\tau_2} \ldots d\tau_1 \ldots d\tau_{n-1} d\tau_n,$$

where $\doteq$ denotes that this is a formal expansion. This is the so-called Dyson series, which is in the heart of QFT.
**REMAINING EXERCISES**

**Bonus question:** Does the series converge in some topology if \( H_1 \) is a bounded operator? (Prove your answer.)

(5) Let \( t_1, \ldots, t_n \) be distinct elements of \( \mathbb{R} \). The time-ordered product of operators \( A_1, \ldots, A_n \) is defined as

\[
T[A_1(t_1) \cdots A_n(t_n)] := A_{i_1}(t_{i_1}) \cdots A_{i_n}(t_{i_n}) \ \text{for} \ t_{i_1} > t_{i_2} > \ldots > t_{i_n}.
\]

(This means just that one orders the product \( A_1(t_1) \cdots A_n(t_n) \) s.t. operators evaluated at highest times are on the left.) Show that

\[
V(t) = T \exp \left[ \frac{i}{\hbar} \int_0^t H_1(\tau) d\tau \right],
\]

where the exponential is defined via its series expansion. This is the definition time-ordered exponential.

(6) Show that it is exactly the ordered exponential which you have defined in the lecture up to peculiarities in the choice of integration borders.

**Exercise 27.108.** (Understanding Theorem 4.3) Lemma 4.2 of the lecture (script) gives you the means to calculate operator products of the form \( S^{(l)} \cdot V S^{(k)} \), where \( V \) is a multiplication operator such as \( V(x) = \gamma^k A_k(x) = \mathbb{A}(x) \). In the last exercise, you have extended this lemma to hold also for the case \( S^{(l)} \cdot V S^{(-1)} \). This is good since now we can also calculate operator products of the form \( sV_s \), where \( s \) is one of the Dirac Green's functions, \( s = s^\wedge \) or \( s = s^\vee \). In this exercise, you will see how this works.

(1) Remind yourself: Why do we want to calculate operator products of the form \( sV_s \) to a given order on the light cone? What's the big goal?

(2) Calculate the operator product \( -s\bar{A}s \) to leading order on the light cone. Show that your result is of the form (4.32) (Theorem 4.3).

Instructions: Use Eq. (4.12) of the script to rewrite the factors \( s \) in the operator product in terms of \( S \). Consider only terms including \( S^{(l)} \) with \( l \leq 0 \) and argue why it is sufficient to do so. Use the result of the last exercise to obtain (with abbreviations as introduced in the lecture (script))

\[
(-s\bar{A}s)(x, y) = i \int_x^y (y - x)^j A_j(z_1) dz_1 s(x, y) + \mathcal{O}((x - y)^{-2}).
\]

(3) Use this result to obtain \( s\bar{A}s\bar{A}s \) to leading order. Generalize this to obtain \( (-s\bar{A})^n s \) to leading order. Show that all terms are of the form (4.32). Hint: Look at the iteration-example (4.25) given in the lecture and do it analogously: Redefine \( A^k(z) \) to include also the terms which you already iterated.

(4) Use this to show that

\[
\tilde{s}(x, y) = P\exp \left( -i \int_x^y (y - x)_{j} A^j(z) dz \right) s(x, y) + \mathcal{O}((x - y)^{-2}).
\]

What does this mean for you? Why does the ordered exponential (and thus the Dyson series) appear here? Do you think that this is a hint that quantum field theory might be a first or second order effective theory of something more fundamental or do you not think so?

(5) **Bonus question:** In QFT, what researchers do is to calculate the leading terms of the Dyson series, i.e. e.g. \( V_1(t), V_2(t) \) and possibly even \( V_3(t) \). This is called “perturbative QFT”. Why is the name “non-perturbative” justified for the things which you have learned so far? Hint: Think about which quantity the word “perturbative” refers to and use the result of question (iv) for your answer.
EXERCISE 27.109. (Mass and spatial normalization of the fermionic projector) In the lecture you have seen that the fermionic projector can be defined with mass- and spatial normalization. Also, in exercise 3.2 you have seen that this indeed makes a difference if one considers the fermionic projector with external potential. In this exercise, you will learn about a symmetry between the two normalizations which was only discovered very recently.

Look for the article “Perturbative Description of the Fermionic Projector: Normalization, Causality and Furry’s Theorem” by Felix Finster and Jürgen Tolksdorf on http://arxiv.org. Read section 3.4 carefully and carry out all the intermediate steps which were not carried out explicitly in the publication.

EXERCISE 27.110. (Mass and spatial normalization of the fermionic projector) In the lecture you have seen that the fermionic projector can be defined with mass- and spatial normalization. Also, in exercise 3.2 you have seen that this indeed makes a difference if one considers the fermionic projector with external potential. In this exercise, you will learn about a symmetry between the two normalizations which was only discovered very recently.

Recall the following definitions from the lecture:

\[ (A|_t B)(x, z) := 2\pi \int_{\mathbb{R}^3} A(x, (t, \vec{y})) \gamma^0 B((t, \vec{y}), z) \, d^3 y, \]  
\[ A \cdot B = C \quad \text{stands for} \quad A_m B_{m'} = \delta(m - m') C_m. \]  

(1) Prove the following equations:

\[ k|_t k = k_m = k_m|_t p_m = p_m|_t k_m \]  

(27.5.16)

EXERCISE 27.111. (Yang-Mills tensor) One of the aims of the lecture is to show how everyday physical equations appear in the continuum limit. In this exercise, we will see a first hint for that. To this end, consider the following terms in the light cone expansion of \( s(\chi_R A) \):s:

\[ \chi_L (\gamma^\sigma R \tilde{A} s) (x, y) \quad \text{phase-free} \quad \chi_L S^{(0)}(x, y) \xi^i \int_x^y dz \, [0, 1 | 0] (\partial A_i)(z) \]  
\[ - \chi_L S^{(0)}(x, y) \int_x^y dz \, [0, 0 | 0] \tilde{A}(z) + \chi_L S^{(0)}(x, y) A(x) + \ldots, \]  

(27.5.17)

where \( \chi_L \) is one of the chiral projectors and abbreviations were used as in the lecture.

(1) Show that each of those terms is phase-free as defined in the lecture, and explain why no calculation is necessary to write down Eq. (27.5.17) to a given order on the light cone.

(2) Show that the three terms are equal to just one term which contains the Yang-Mills tensor \( F_{ik} = \partial_i A_k - \partial_k A_i \).

(3) Explain what one has to do in order to obtain the contributions to the light cone expansion of the residual fermionic projector \( \tilde{P}_{\text{res}} \) from an equation such as Eq. (27.5.17). Transform the three terms accordingly.

EXERCISE 27.112. (Signature operator) Prove the following proposition, which shows that the causal fundamental solution is the signature operator of the dual pairing \(<.|.>\).

Proposition 5.1: For any \( \psi_m \in H_m \) and \( \phi \in C^\infty_0(M, SM) \),

\[ (\psi_m | k_m \phi) = <\psi_m|\phi>. \]  

(27.5.18)

Hint: Proceed as follows:

(1) First, assume that \( \psi_m \in C^\infty_0(M, SM) \). Chose Cauchy surfaces \( N_+ \) and \( N_- \) which lie the future and past of supp \( \phi \), respectively. Show that then

\[ (\psi_m | k_m \phi) = i \int_{\Omega} \nabla_j <\psi_m| \gamma^j \tau_m^\Lambda \phi> x \, d\mu(x), \]  

(27.5.19)

where \( \Omega \) is the spacetime region between the two surfaces.
(2) Use that \( \psi_m \) satisfies the Dirac equation to show that
\[
(\psi_m \big| k_m \phi) = \int_M <\psi_m|\phi> \, d\mu(x) \tag{27.5.20}
\]

(3) In order to extend the result to general \( \psi_m \in \mathcal{H}_m \), let \( \psi_m^{(n)} \in \mathcal{H}_m \cap C_\infty(M, SM) \) be a sequence which converges in \( \mathcal{H}_m \) to \( \psi_m \). Show that \( (\psi_m^{(n)} \big| k_m \phi) \to (\psi_m \big| k_m \phi) \) and that \( <\psi_m^{(n)}|\phi> \to <\psi_m|\phi> \). (For the last step, use that \( \phi \) is compactly supported and Fubini's theorem. Also, for the whole exercise, do not forget about our good friends Mr. Cauchy and Mr. Schwarz.)

Exercise 27.113. (Regularization operator) Prove that the family \( (\mathfrak{R}_\epsilon)_\epsilon > 0 \) of operators
\[
(\mathfrak{R}_\epsilon \psi)(t, \vec{x}) := (U^{t, t_0}(\eta_\epsilon \ast \psi | t_0 = t)) (\vec{x}) \tag{27.5.21}
\]
is a family of regularization operators, where \( U^{t, t_0} \) is the time evolution operator defined in the lecture,
\[
\eta_\epsilon(\vec{x}) = e^{-\frac{|\vec{x}|^2}{\epsilon^2}} \frac{1}{\epsilon^3}
\]
and
\[
(\eta_\epsilon \ast \psi)(\vec{x}) := \int \eta_\epsilon(\vec{x} - \vec{y}) \psi(\vec{y}) \, d^3 y.
\]
You may use the following inequality:
\[
\| (U^{t, t_0} \phi)(t, \vec{x}) \|_{C^1} \leq C \| \phi \|_{W^{2, 2}(\mathbb{R}^3)} \quad \forall \phi \in \mathcal{H}_m, \tag{27.5.22}
\]
where \( W^{k,p} \) denotes the Sobolev spaces (sometimes also denoted by \( H^{k,p} \)).

Hint: Within this exercise, do not think of \( \mathcal{H}_m \) as representing the identification of the \( \mathcal{H}_t \) which were introduced in the context of the time evolution operators.
Note that on the right hand side of Eq. (5.2) in the lecture, the norm denotes the norm on \( \mathcal{H}_m \). For the proof of the latter, use Eq. (27.5.22), a Fourier transformation in spatial directions, Plancherel’s theorem and the behavior of convolutions under the Fourier transform.

Exercise 27.114. (Bonus exercise) Prove inequality \( (27.5.22) \).

Exercise 27.115. (Causal fermion systems) Let \( (\mathcal{H} = l^2(\mathbb{N}), \langle \cdot | \cdot \rangle) \) be the Hilbert space of square integrable valued sequences, i.e.,
\[
\langle u|v \rangle := \sum_{l=1}^\infty \bar{u}_l \, v_l.
\]
Let \( x_k \) be an operator on \( \mathcal{H} \) which is defined as
\[
(x_k u)_k = -u_{k+1}
\]
\[
(x_k u)_{k+1} = -u_k
\]
\[
(x_k u)_l = 0 \quad \forall \, l \in \mathbb{N} \setminus \{k, k+1\},
\]
i.e. for \( u = (u_1, u_2, ..., u_k, u_{k+1}, ....) \),
\[
x_k u := \begin{cases} 0, \ldots, 0, & -u_{k+1}, -u_k, 0, \ldots \end{cases} \quad \text{\( k-1 \) entries}
\]

(1) Construct a causal fermion system \( (\mathcal{H}, \mathcal{F}, \rho) \) where \( \mathcal{H} \) is as above and where \( \text{supp} \, \rho = \{x_k\}_{k \in \mathbb{N}} \).

Hint: Use a discrete counting measure, i.e. a sum of Dirac measures.

(2) Derive the following objects for this causal fermion system and analyze them:
\( (\text{a}) \) Spacetime \( M \)
(b) Spin spaces \( (S_x, \prec, \succ, \sim x) \)
(c) Wave functions \( \psi \)
(d) Krein space \( \mathcal{K} \) (including topology!)

(3) Answer the following questions: Is the map \( \Psi : \mathcal{H} \to \mathcal{K} \) well-defined and continuous? Is it an embedding? Are \( \mathcal{H} \) and \( \mathcal{K} \) isomorphic?

(4) Try to give a physical interpretation of what this causal fermion system could represent.

**Exercise 27.116. (Eigenvalues of \( xy \))** Let \((\mathcal{H}, \mathcal{F}, \rho)\) be a causal fermion system. Show that for \( x, y \in \mathcal{F} \), the operator product \( xy \) has the same non-trivial eigenvalues as the closed chain \( P(x, y)P(y, x) \), where \( P(x, y) \) is defined in Eq. (5.5).

Hint: Remember that you cannot use the characteristic polynomial on infinite dimensional vector spaces! Therefore, argue why, for this purpose, you can restrict \( xy \) to \( S_y \). Prove the identity \( \det(BC - \lambda I) = \det(CB - \lambda I) \) and use it.

**Exercise 27.117. (The fermionic projector in bra-ket notation)** Familiarize yourself with the bra-ket notation (e.g. read the English Wikipedia article on this subject). Afterwards, prove that
\[
P(x, y) = -\sum_{k \in \mathbb{N}} |\psi^u_k(x)\rangle \langle \psi^u_k(y)|.
\]

**Exercise 27.118. (\( \gamma \)-matrices) Prove the following identities:**
\[
\begin{align*}
(1) \quad & \gamma^5 = \frac{1}{4\pi} \varepsilon_{\mu
u\alpha\beta} \gamma^\mu \gamma^\nu \gamma^\alpha \gamma^\beta \\
(2) \quad & \gamma^\mu \gamma^\nu = 4I_4 \\
(3) \quad & \gamma^\mu \gamma^\nu \gamma^\mu = -2\gamma^\nu \\
(4) \quad & \gamma^\mu \gamma^\nu \gamma^\alpha \gamma^\beta = 4\eta^{\rho\sigma} I_4 \\
(5) \quad & \gamma^\mu \gamma^\nu \gamma^\alpha \gamma^\beta = -2\gamma^\sigma \gamma^\Gamma \gamma^\nu \\
(6) \quad & \text{tr}(\gamma^\mu) = 0 \\
(7) \quad & \text{trace of any product of an odd number of } \gamma^\mu \text{ is zero} \\
(8) \quad & \text{trace of } \gamma^5 \text{ times a product of an odd number of } \gamma^\mu \text{ is still zero} \\
(9) \quad & \text{tr}(\gamma^\mu \gamma^\nu) = 4\delta^{\mu\nu} \\
(10) \quad & \text{tr}(\gamma^{\mu_1} \ldots \gamma^{\mu_n}) = \text{tr}(\gamma^{\mu_1} \ldots \gamma^{\mu_1})
\end{align*}
\]

**Exercise 27.119. (Eigenvalues of the closed chain)** Prove Lemma 7.1 of the lecture, i.e. prove that the characteristic polynomial of \( A^\varepsilon_{xy} \) has the roots
\[
\lambda_{\pm} = g\gamma + h\gamma \pm \sqrt{(g\gamma)^2 - g^2 \gamma^2 + (g\gamma + h\gamma)^2}.
\]

Show that they either form a complex conjugate pair or else they are both real and have the same sign.

Hint: From the lecture you know that you can write the (regularized) fermionic projector as
\[
P^\varepsilon(x, y) = g_\varepsilon(x, y) \gamma^5 + h\varepsilon(x, y) , P^\varepsilon(y, x) = g_\varepsilon(x, y) \gamma^5 + h\varepsilon(x, y).
\]

With this, calculate the matrix \( A^\varepsilon_{xy} =: A \) and decompose it into the form
\[
A = A_1 + A_2 + \mu,
\]
where
\[
A_1 = \frac{1}{2} [g\gamma, g\gamma], \quad A_2 = h\gamma + g\gamma.
\]

Show that the matrices \( A_1 \) and \( A_2 \) anti-commute and use this to show that
\[
(A - \mu)^2 = (g\gamma)^2 - g^2 \gamma^2 + (g\gamma + h\gamma)^2.
\]

Use this polynomial equation in \( A \) to find the zeros of the characteristic polynomial of \( A \). In order to show that if they are real, they have the same sign, calculate the product \( \lambda_+, \lambda_- \) and show that
\[
\lambda_+ \lambda_- = (g^2 - h^2)(\gamma^2 - \gamma^2) = |(g^2 - h^2)|^2 \geq 0.
\]
Exercise 27.120. (Understanding the \(i\varepsilon\) - Regularization) Note: The next exercise sheet will contain a “reloaded” version of this exercise with more explanations. But go ahead and try it already if you like.

Prove the following statements from the lecture:

1. 
\[
T^{(0)}_{[p]} \rightarrow -\frac{1}{8\pi^3} \frac{1}{2r} \left( t - r - i\varepsilon \right), \quad T^{(-1)}_{[p]} \rightarrow -\frac{2}{r} \frac{\partial}{\partial t} T^{(0)}_{[p]} = -\frac{1}{8\pi^3} \frac{1}{r^2} \frac{1}{(t - r - i\varepsilon)^2},
\]

2. 
\[
A_{\varepsilon}^{x y} = \frac{1}{256\pi^6} \frac{t^2 - r^2 + \varepsilon^2}{|t - r + i\varepsilon|^4} - \frac{1}{256\pi^6} \frac{i\varepsilon[\gamma_0, \bar{\xi}]}{|t - r + i\varepsilon|^4}.
\]

3. The bilinear contribution to \(A_{\varepsilon}^{x y}\) makes the eigenvalues \(\lambda_{\pm}\) complex.

4. The leading divergence of the integral (7.11) in \(\varepsilon\) can be described by the scaling
\[
T^{(n)} \sim (\varepsilon|\xi|)^{n-1} \quad \text{and} \quad dt \sim \varepsilon.
\]

5. The scaling of the other divergent contributions is obtained by multiplying with powers of
\[
\varepsilon |\xi|.
\]

6. 
\[
\frac{1}{4} \text{Tr} \left( (\xi - i\varepsilon\gamma_0)(\xi + i\varepsilon\gamma_0) \right) = \frac{1}{2} (\xi - i\varepsilon\gamma_0)^2 + \frac{1}{2} (\xi + i\varepsilon\gamma_0)^2 + (\text{higher orders in } \varepsilon/|\xi|)
\]

Here is a list of formulae from the lecture which you might find helpful in proving this exercise. (Of course, you can find all of them at various places in the script.)

\[
T_{m^2}(x, y) := \int \frac{d^4k}{(2\pi)^4} \delta(k^2 - m^2) \Theta(-k^0) e^{-ik(x-y)}
\]

\[
T_{m^2}(x, y) = -\frac{1}{8\pi^3} \left( \frac{\text{PP}}{(y-x)^2} + i\pi\delta((y-x)^2) \epsilon((y-x)^0) \right)
\]

\[
+ \frac{m^2}{32\pi^3} \sum_{j=0}^{\infty} \frac{(-1)^j}{j! (j+1)!} \frac{(m^2(y-x)^2)^j}{4^j}
\]

\[
\times \left( \log|m^2(y-x)^2| + c_j + i\pi \Theta((y-x)^2) \epsilon((y-x)^0) \right)
\]

\[
T^{(i)} = \left( \frac{d}{da} \right)^i T_{a=0}
\]

\[
P(x, y) = (i\phi - m) T_{m^2}(x, y)
\]

\[
\hat{P}^\varepsilon(k) = (\hat{k} + m) \delta(k^2 - m^2) \Theta(-k^0) \exp(\varepsilon k^0).
\]

\[
\lim_{\varepsilon \to 0} \log(a\xi^2 - i\varepsilon\xi^0) + i\pi = \log|a\xi^2| + i\pi \Theta(\xi^2) \epsilon(\xi^0)
\]

Exercise 27.121. (Understanding the \(i\varepsilon\) - Regularization Reloaded (1))
(1) Look up the formula for the Greensfunction of the Klein-Gordon-Operator in position space, \( T_n(x,y) \), in the script \((a = m^2)\). Drop the log(a)-Term (this corresponds to the “regularization” introduced in Ch. 4) and the smooth contributions (since we are presently only interested in the behavior of the singular components). Then, calculate \( T^{(n)} \) for \( n \leq 0 \).

(2) Explain the idea behind the \( \varepsilon \)-regularization and explain why it amounts to setting \( t \to t - \varepsilon \), where \( \xi^0 \equiv t \).

(3) Plug in the \( \varepsilon \)-regularization into \( T^{(0)} \). Use that the only goal we presently have is to evaluate any composite expressions of \( T \)'s weakly on the light-cone away from the origin, which means that \( t \approx |\xi| =: r \) to show that

\[
T^{(0)}_{[0]} \to -\frac{1}{8\pi^3} \, \frac{1}{2r(t-r-\varepsilon)^2}.
\]

(4) Use the defining formula for \( T^{(-1)} \),

\[
\xi_k T^{(-1)}(x,y) := 2 \frac{\partial}{\partial x^k} T^{(0)}(x,y)
\]

for \( k = 0 \) to obtain

\[
T^{(-1)} \to T^{(-1)}_{[0]} = -\frac{1}{8\pi^3 r^2} \, \frac{1}{(t-r-\varepsilon)^2}.
\]

(Note that Eq. (27.5.30) follows from the corresponding equation for \( S \) due to the residual argument.)

(5) Given the Greens-functions for the Klein-Gordon operator, \( T \), you know how to construct the fermionic operator. Do so for the regularized and massless case, i.e., set

\[
P^\varepsilon(x,y) = \frac{i}{2} \gamma^\varepsilon T^{(-1)}_{[0]}.
\]

With this, show that

\[
A^\varepsilon_{xy} := P^\varepsilon(x,y) P^\varepsilon(y,x) = \frac{1}{256\pi^6 r^4} \frac{(t^2-r^2+\varepsilon^2)^2}{|t-r+i\varepsilon|^4} - \frac{1}{256\pi^6 r^4} \, \frac{i\varepsilon [\gamma^0, \xi]}{|t-r+i\varepsilon|^4}.
\]

(6) Show that if \( x \) is on the light cone of \( y \) (or vice versa), which means that \( t = r \), then for every \( \varepsilon > 0 \), the last term of Eq. (27.5.31) makes the eigenvalues \( \lambda_\pm \) of \( A_{xy} \) complex. What does this mean physically?

Hint: In the last exercise session, we have discussed that if one seeks the eigenvalues of a matrix containing \( \gamma^k \)'s, it is a good idea to check if the matrix fulfills a nice polynomial equation by squaring it. Try this here, it works.

\[
(\varepsilon [\gamma^0, \xi])^2 = -4\varepsilon^2 |\xi|^2 < 0
\]

**Exercise 27.122.** (Understanding the \( \varepsilon \)-Regularization Reloaded (2)) Next, we want to understand the scaling behavior of \( T^{(n)}_{[n]} \) if the regularization is removed. (We need to know this in order to evaluate which terms to keep and which terms to “throw away” when analyzing the continuum limit.) To this end, we take a look at simple fractions which are evaluated weakly on the light cone,

\[
\int_{|\xi| - \varepsilon}^{|\xi| + \varepsilon} dt \, \eta(t, \xi) \, \frac{\prod_{\alpha=1}^{n} T^{(a_\alpha)}_o \prod_{\beta=1}^{n} T^{(b_\beta)}_o}{\prod_{\gamma=1}^{n} T^{(c_\gamma)}_o \prod_{\delta=1}^{n} T^{(d_\delta)}_o}
\]

and investigate their divergence with \( \varepsilon \to 0 \).

**Assumption:** \( \eta(t) \) “varies on the scale \( l_{\text{macro}} \)”, i.e.,

\[
\frac{\partial}{\partial t} \eta(t') \approx \frac{1}{l_{\text{macro}}} \eta(t')
\]
when \( t' \approx r \).

Proceed as follows:

1. First of all, use 8.2(i) to show that \( T(n) \sim \left( (t - i\varepsilon - r)(t - i\varepsilon + r) \right)^{n-1} \), where \( \sim \) denotes proportionality. (Think about why you can drop the \( \log(\xi) \)-Term.)

2. Assume that the support of \( \eta(t) \) is contained in the interval \( (r - \varepsilon, r + \varepsilon) \), where \( r = |\xi| \).

\[
(27.5.32) \sim \int_{-\infty}^{\infty} \eta(t) \left( (t - i\varepsilon - r)(t - i\varepsilon + r) \right)^N \left( (t + i\varepsilon - r)(t + i\varepsilon + r) \right)^M \left( (t + i\varepsilon - r)(t + i\varepsilon + r) \right)^P \left( (t - i\varepsilon - r)(t - i\varepsilon + r) \right)^Q \, dt
\]

for suitable \( N, P, Q, M \in \mathbb{R}_0^+ \).

(3) Assume that \( N - P < 0 \). The integrand of the last equation has a pole at \( t = r + i\varepsilon \). Close the integration contour (utilizing the support-condition of \( \eta(t) \)) and evaluate the integral with the residue theorem. (Assume that \( \eta \) is only supported around this pole, i.e. neglect the other poles.)

4. The residuum which you get contains a factor

\[
\frac{d^{N+P-1}}{dt^{N+P-1}} \left[ \eta(t) \left( (t - i\varepsilon - r)(t - i\varepsilon + r) \right)^N \left( (t + i\varepsilon - r)(t + i\varepsilon + r) \right)^M \left( (t + i\varepsilon - r)(t + i\varepsilon + r) \right)^P \left( (t - i\varepsilon - r)(t - i\varepsilon + r) \right)^Q \right] \bigg|_{t=r+i\varepsilon}.
\]

Apply one \( t \)-derivative to the term in brackets and evaluate it at \( t = r + i\varepsilon \). Evaluate how the four different terms which you obtain behave when \( \varepsilon \to 0 \). (Assume \( M - Q < 0 \).)

With this, prove that

\[
(27.5.34) \sim r^{N-P+M+Q} \eta(r) \varepsilon^{N-P+M+Q+1} \left( 1 + \mathcal{O}\left( \frac{\varepsilon}{r} \right) \right) \left( 1 + \mathcal{O}\left( \frac{\varepsilon}{l^{\text{macro}}} \right) \right)
\]

(5) Finally, use part (ii) to show that if one associates one of the \( \varepsilon \) of Eq. (27.5.34) to \( dt \),

\[
dt \sim \varepsilon,
\]

then one has indeed

\[
T(n) \sim (\varepsilon|\xi|)^{n-1} \left( 1 + \mathcal{O}\left( \frac{\varepsilon}{r} \right) \right) \left( 1 + \mathcal{O}\left( \frac{\varepsilon}{l^{\text{macro}}} \right) \right)
\]

as stated in the lecture.

**Exercise 27.123. (Spectral projectors of the regularized vacuum)** Show that

\[
F_\pm = \frac{1}{2} \left( 1 \pm \frac{|\xi|}{z - \xi} \right) + \xi (\deg \leq 0) + (\deg < 0).
\]

\[
(27.5.35)
\]

and

\[
F_\pm P(x, y) = \left\{ \begin{array}{cc} 0 & \text{for "+"} \\ \xi \xi^{-1} T_{[0]}^{-1} & \text{for "-"} + (\deg < 2) \end{array} \right. \mathcal{O} \left( \frac{\varepsilon}{l^{\text{macro}}} \right)
\]

\[
(27.5.36)
\]

**Exercise 27.124. (Resolvent formalism)** In the lecture, we have defined

\[
G_k = \frac{1}{2\pi i} \oint_{\Gamma_k} (z - A)^{-1} \, dz
\]

where \( \Gamma_k \) is a small contour around the \( k \)-th eigenvalue of the matrix \( A \) which does not enclose any other eigenvalue. If you are not familiar with the resolvent formalism, have a look e.g. into Kato, “Perturbation Theory for Linear Operators”, p. 36-40.

1. Determine the spectrum of the following matrices and the corresponding operators \( G_k \):

\[
(a) \quad A := \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}
\]

\[
(b) \quad B := \begin{pmatrix} \lambda_1 & 0 \\ 0 & \lambda_2 \end{pmatrix}
\]

(2) Prove the resolvent identity.

(3) Show that \( G_k G_l = \delta_{k,l} G_l \). Hint: For the case \( k = l \), use the fact that the two integration contours \( \Gamma_k \) and \( \Gamma_l \) lie outside one another. For the case \( k = l \) use that you can deform the contour continuously (so long as it does not “swipe over another pole”) and deform \( \Gamma_k \) and \( \Gamma_l \) such that one encloses the other.

**Exercise 27.125.** (Understanding the contraction rules) This exercise intends to make you familiar with the contraction rules in the formalism of the continuum limit. The goal is to prove the contraction rules for the \( \varepsilon \)-regularization.

(1) Explain what the general idea is behind the method of variable regularization and behind the formalism of the continuum limit.

In order to prove the contraction rules for the \( \varepsilon \)-regularization, we first need to know explicitly the functions \( \xi_{[p]}^{(n)} \), \( \bar{\xi}_{[p]}^{(n)} \) and \( T_{[p]}^{(n)} \). (Set \( T_{[p]}^{(n)} = 0 \).) Since we are working in the \( \varepsilon \)-regularization, we define

\[
\xi_{[p]}^{(n)} := \xi|_{z^0 \to z^0 - \varepsilon} \quad \forall n, p
\]

(2) Prove that

\[
T_{[p]}^{(0)} = -\frac{1}{8\pi^3} \frac{1}{2r(t + i\varepsilon - r)} \left( 1 + O\left( \frac{\bar{z}}{r} \right) \right) \quad (27.5.38)
\]

\[
T_{[p]}^{(n)} = c_n \left( 2r(t + i\varepsilon - r) \right)^{n-1} \left( \log(2r(t + i\varepsilon - r)) + d_n \right) \left( 1 + O\left( \frac{\bar{z}}{r} \right) \right) \quad \forall n \geq 1 \quad (27.5.39)
\]

where \( c_n, d_n \in \mathbb{C} \).

(3) Use the equation

\[
\frac{\partial}{\partial x_k} T^{(l)} = \frac{1}{2} \xi_k T^{(l-1)} \quad (27.5.40)
\]

to fix all coefficients \( c_n, d_n \). Hint: Only consider the case \( k = 0 \) and use again that \( t = r + O(\varepsilon) \). Mind the hidden minus sign (\( \xi = (y - x) \)).

We can now prove the contraction rules

\[
(\xi_{[p]}^{(n)})^j (\xi_{[p']}^{(n')})_j = \frac{1}{2} \left( \xi_{[p]}^{(n)} \varepsilon_{[p']}^{(n')} + \varepsilon_{[p]}^{(n)} \xi_{[p']}^{(n')} \right) \quad (27.5.41)
\]

\[
(\xi_{[p]}^{(n)})^j (\bar{\xi}_{[p']}^{(n')})_j = \frac{1}{2} \left( \bar{\xi}_{[p]}^{(n)} \varepsilon_{[p']}^{(n')} + \varepsilon_{[p]}^{(n)} \bar{\xi}_{[p']}^{(n')} \right) \quad (27.5.42)
\]

\[
\bar{\xi}_{[p]}^{(n)} T_{[p]}^{(n)} = -4 \left( n T_{[p]}^{(n+1)} + T_{[p]}^{(n+2)} \right). \quad (27.5.43)
\]

(4) In the context of the \( \varepsilon \)-regularization, it suffices to define \( z_{[p]}^{(n)} := (\xi_{[p]}^{(n)})^j (\xi_{[p]}^{(n)})_j \). Prove that with this definition, eqs. (27.5.41) and (27.5.42) hold up to \( O(\varepsilon^2) \).

(5) Prove that up to smooth contributions we have for the \( T_{[p]}^{(n)} \) derived in part (iii)

\[
(\xi_{[p]}^{(n)})^2 T_{[p]}^{(n)} = -4 n T_{[p]}^{(n+1)},
\]

and show that this corresponds exactly to eq. (27.5.43).

**Exercise 27.126.** (Bra-ket notation II) Prove that in the context of the Bra-ket notation, for \( u, v \in \mathcal{H} \),

\[
\text{Tr} \left( |v\rangle \langle u| A \right) = \text{Tr} \left( \langle u| A v \rangle \right) = \langle u| A v \rangle.
\]

Hint: Use an orthonormal basis of \( \mathcal{H} \).

**Exercise 27.127.** (Double null spinor frame) Given the \( \varepsilon \) regularization, calculate explicitly the double null spinor frame \( \left( f_L^+, \xi_L^+, f_R^+, \bar{\xi}_L^+ \right) \) at a point \( x \) with respect to a point \( y \). (Up to \( O(\varepsilon) \).)

Verify that \( \xi f_L^+ = f_R^+ \).

Hint: Remember that the double null spinor frame is composed of a particular set of eigenvectors.
of the matrix $A_{x,y}$ (hence the term “w.r.t. a point $y$”). Assume that the spatial component of $\xi := x - y$ is normalized and points into the $z$-direction, $\hat{\xi} = (0,0,1)$. In order to determine $\gamma^c$, calculate the matrix $\chi^c F_s$ explicitly and determine its eigenvector. (And remember that the $\xi$ in Eq. (8.9) in the script is really $\gamma^{(n)}$.)

**Exercise 27.128.** (Matrix elements in the double null spinor frame) In $T_{[0]}^{(-1)}$ be an arbitrary regularization of $T^{(-1)}$ and let
\[ P_0(x,y) := \frac{i}{2} \chi^c T_{[0]}^{(-1)} . \] Determine the matrix elements $P_0(x,y)^{+}_{\pm \pm}$ and $P_0(x,y)^{+ \pm}_{\pm}$ (denoted by $F^\pm_\pm(P_0(x,y))$ and $F^\pm_\mp(P_0(x,y))$ in the script). Simplify the expressions as far as possible.

**Hint:** Use the cyclic property of the trace, the anti-commutation relations of the gamma matrices and the contraction rules.

**Exercise 27.129.** (Gauge potentials) In this exercise, you will see how the gauge freedom of electrodynamics is contained in causal fermion systems.

1. Show that if $A_\mu$ is a solution of the homogeneous Maxwell equations on Minkowski space, so is $A'_\mu := A_\mu + \partial_\mu \Lambda$, where $\Lambda$ is a real-valued function on Minkowski space.
2. The Dirac equation coupled to electromagnetism, $(i \partial + A - m) \psi = 0$, is not invariant under the transformation $A'_\mu := A_\mu + \partial_\mu \Lambda$, hence this is not a gauge transformation. Show that if one suitably transforms $\psi$ together with $A$,
   \[ A_\mu(x) \rightarrow A'_\mu(x), \]
   \[ \psi(x) \rightarrow U(x) \psi \]
   this is a gauge transformation. (Determine $U(x)$.)
3. Use exercise 6.3 to show that in the context of causal fermion systems, the transformation $P(x,y) \rightarrow U(x)P(x,y)U(y)^*$.
4. Show that the transformation leaves the causal action $S$ invariant. This means that if a causal fermion system can be described by the vector potential $A_\mu$, there is another causal fermion system which is a minimizer of the causal variational principle which can be described by the vector potential $A'_\mu$.

**Exercise 27.130.** (Perturbation of the eigenvalues of the closed chain) Assume that you want to study if a causal fermion system can be described by Dirac spinors $\psi$ coupled to some quantity $\mathcal{X}$. If $\mathcal{X} = \gamma^5 A$, you know already from the lecture or talks that the causal fermion system can only be a minimizer of the causal action principle (and hence “physically admissible”) if $A$ and $\psi$ are related by the Dirac equation coupled to the Yang-Mills equation.

Now assume that you wanted to study this for some (possibly very exotic) $\mathcal{X}$ and that you find that the corresponding perturbation of the fermionic projector to first order $\delta P$ (sometimes also denoted by $\Delta P$ in the script) is
\[ \delta P(x,y) = \chi^c \psi, \]
where $\chi^c$ is the left-handed chiral projector and $\psi$ is some arbitrary vector in Minkowski space.

1. Calculate the corresponding variation of the eigenvalues of the closed chain to first order in the variation, $\delta \lambda^{x,y}_k$, up to leading order in the degree. Which eigenvalues change, which not?
2. Assume that $\psi = \xi + \bar{\xi}$ and calculate $\delta \lambda^{x,y}_k$ explicitly (i.e. your result should not contain any Gamma matrices and traces thereof any more).
(3) Assume that you have derived the more general perturbation
\[ \delta P(x,y) = \chi_L \psi_L + \chi_R \psi_R. \] (27.5.48)
where \( \nu_L \) and \( \nu_R \) are vectors in Minkowski space. Based on your results in part (i), answer the following question: For which vectors \( \nu_L \) and \( \nu_R \) does
\[ |\lambda^{xy}_k| = |\lambda^{xy}_l| \quad \forall k, l \]
hold?

(4) What does this mean for potential \( \chi \) which you (hypothetically) derived from equation (27.5.48)?

Hint: Use equation (8.35) from the script, and remember that the \( F^{xy}_k \) are given by \( \chi_{L/R} F_{\pm} \).
Remaining Exercises from Online Course 2021

EXERCISE 27.131. (The regularized fermionic projector in Minkowski space) Consider the kernel of the fermionic projector regularized in momentum space by a convergence-generating factor $e^{-\varepsilon |k^0|}$, i.e.

$$P^\varepsilon(x,y) = \int_{\mathbb{R}^4} \frac{d^4k}{(2\pi)^4} \left( \frac{k + m}{\Theta(-k^0)} \delta(k^2 - m^2) e^{-ik(x-y)} e^{-\varepsilon |k^0|} \right). \tag{27.5.49}$$

(1) Show that $P^\varepsilon(x,y)$ can be written as $\rho^\varepsilon + \beta^\varepsilon$, for $\rho^\varepsilon, \beta^\varepsilon$ smooth functions of $\xi = y - x$.

(2) Compute $P^\varepsilon(x,x)$. Is this matrix invertible? How does it scale in $\varepsilon$?

For $\xi$ spacelike or timelike, i.e. away from the lightcone, the limit $\varepsilon \searrow 0$ of (27.5.49) is well-defined. More precisely, it can be shown that $\rho^\varepsilon \to \alpha \xi$ and $\beta^\varepsilon \to \beta$ pointwise, for $\alpha, \beta$ smooth complex functions. Find smooth real functions $a, b$ such that

$$\lim_{\varepsilon \to 0} A^\varepsilon_{xy} = a/\xi + b. \tag{27.5.50}$$

How do the eigenvalues of (27.5.50) look like? Discuss them in relation to the notion of causality in the setting of causal fermion systems.

EXERCISE 27.132. (Understanding the connection between causal structure and closed chain) Let $x, y \in \mathbb{R}^4$ be timelike separated vectors and assume that $\xi := y - x$ is normalized to $\eta(\xi, \xi) = 1$. As explained in Exercise 5.1, the limit $\varepsilon \searrow 0$ of the closed chain $A^\varepsilon_{xy}$ takes the form $A = a/\xi + b$. Consider the matrices

$$F^\pm := \frac{1}{2} \left( I \pm \xi \right) \in \text{M}(4, \mathbb{C}).$$

Prove the following statements.

(1) $F^\pm$ have rank two and map to eigenspaces of $A$. What are the corresponding eigenvalues?

(2) $F^\pm$ are idempotent and symmetric with respect to the spin inner product $\prec \cdot, \cdot \succ$ on $\mathbb{C}^4$.

(3) The image of the matrices $F^\pm$ is positive or negative definite.

(4) The image of $F_+$ is orthogonal to that of $F_-$ (with respect to the spin inner product)

The result of this exercise can be summarized by saying that the $F^\pm$ are the spectral projection operators of $A$.

EXERCISE 27.133. (Spin spaces in Minkowski space - part 1) Let $\mathcal{H}^-_m$ denote the Hilbert space of negative-energy solutions of the Dirac equation as introduced in the lecture. By means of a convergence-generating factor as in Exercise 5.1 it is possible to define a bounded regularization operator

$$\mathcal{R}_\varepsilon : \mathcal{H}^-_m \to \mathcal{H}^-_m \cap C^\infty(\mathbb{R}^4, \mathbb{C}^4),$$

which can be proved to be injective. As you know from the lecture, this allows us to define local correlation operators $F^\varepsilon(x)$ on $\mathcal{H}^-_m$ via

$$\langle u | F^\varepsilon(x) v \rangle := -\langle \mathcal{R}_\varepsilon u(x), \mathcal{R}_\varepsilon v(x) \rangle. \tag{27.5.51}$$

This gives rise to a causal fermion system, called the regularized Dirac sea vacuum.

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Conclude that $S$ distinguishes a direction of time where $T_A, B$.

In the example of a causal fermion system describing a Dirac sea vacuum in Minkowski space, $P(A, B)$ be a Hilbert basis of $\mathcal{H}^\infty_\alpha$. More precisely, for $(2)$ Referring to Exercise 5.1, deduce that, for spacelike separation, $\epsilon = 0.$

Conclude that $P^\epsilon(x, y) \in \mathcal{H}^\infty_\alpha$ and $P^\epsilon(x, y)$ are linearly independent.

Exercise 27.135. (Time direction in Minkowski space) A causal fermion systems distinguishes a direction of time through the anti-symmetric functional

$$C : M \times M \ni (A, B) \mapsto i \text{tr}(BA \pi_B \pi_A - A B \pi_A \pi_B) \in \mathbb{R}. \quad (27.5.52)$$

A point $B \in M$ is said to lie in the future (past) of $A \in M$ whenever $C(A, B)$ is positive (negative). In the example of a causal fermion system describing a Dirac sea vacuum in Minkowski space, the functional $\mathcal{C}$ can be equivalently represented on $\mathbb{R}^4$ by (see Lemma 1.2.11 in the blue book)

$$C : \mathbb{R}^4 \times \mathbb{R}^4 \ni (x, y) \mapsto i \text{Tr}_{\mathbb{C}^4} (P^\epsilon(x, y) P^\epsilon(y, x) - P^\epsilon(x, y), A^\epsilon_{xy})$$

Exercise 27.136. (The time-direction functional in Minkowski space (6 points)) Away from the light-cone the kernel of the fermionic projector $P^\epsilon$ converges to a smooth function $P$. More precisely, for $\xi \in \mathbb{R}^3 \setminus L_0$,

$$P^\epsilon(\xi) \to P(\xi) = (i\partial - m)T_{m2}(\xi), \quad T_{m2}(\xi) := \int \frac{d^4k}{(2\pi)^4} \delta(k^2 - m^2) \Theta(-k^0)e^{-ik\cdot\xi}.$$
A causal fermion system distinguishes a direction of time by means of the anti-symmetric real functional
\[ C : M \times M \ni (A, B) \mapsto i \text{tr}(BA \pi_B \pi_A - AB \pi_A \pi_B) \in \mathbb{R}. \]

From Exercise 5.1-(ii) we know that \( P^\varepsilon(0) \) is invertible. Let us define \( \nu := P^\varepsilon(0)^{-1} \in \text{Mat}(4, \mathbb{C}). \)

Using the identifications as in Exercise 5.3-4 \((x \equiv F^\varepsilon(x), S_x \cong \mathbb{C}^4)\), prove the following identities (up to global constants)
1. \( \pi_x y x \pi_y \pi_x S_x = P^\varepsilon(x, y) P^\varepsilon(y, x) P^\varepsilon(x, y) \nu P^\varepsilon(y, x) \nu. \)
2. \( C(x, y) = i \text{Tr}_{C^4}(P^\varepsilon(x, y) \nu P^\varepsilon(y, x) \nu[A, \pi_y]). \)

Let \( x, y \) be spacelike separated. Using (2) and (ii), what can you infer about the size of the functional \( C(x, y) \) in the limit \( \varepsilon \searrow 0? \) Hint: Discuss the commutator in (ii). The scaling in \( \varepsilon \) from Exercise 5.1-(ii) may be useful.

**Exercise 27.137. (Closedness of the local correlation function (6 points))** The goal of this exercise is to show that the local correlation operators, as defined in Exercise 5.3, realizes a one-to-one topological identification of Minkowski space with a closed subset of \( \mathcal{F} \). Let us define the local correlation function by
\[ F^\varepsilon : \mathbb{R}^4 \to \mathcal{F}, \quad \langle u, F^\varepsilon(x) v \rangle := -\langle \mathfrak{R}_x u(x), \mathfrak{R}_y v \rangle. \quad (27.5.53) \]

Thanks to the translation invariance of the Dirac sea, it can be proved that all the \( F^\varepsilon(x) \) are unitarily equivalent, in particular they have the same norm.

1. **Continuity:** The regularization operator as in Exercise 5.3 can be chosen to fulfill:
   a. There is \( C > 0 \) such that \( |\mathfrak{R} u(x)| \leq C ||u|| \) for all \( x \in \mathbb{R}^4 \)
   b. For all \( x \in \mathbb{R}^4 \) and \( \delta > 0 \) there is \( r > 0 \) such that \( |\mathfrak{R} u(x) - \mathfrak{R} u(y)| \leq \delta ||u|| \) for all \( u \in \mathcal{H}_m \) and all \( y \in B_r(x) \).

   Use these properties to show that \( F^\varepsilon \) is continuous in the operator topology.

2. **Injectivity:** Let \( F^\varepsilon(x) = F^\varepsilon(y) \). We need to show that \( x = y \). Referring to Exercise 1.2, consider the elements \( u_n(p) \in \mathcal{H}_m \) whose regularization reads \((\varepsilon_4 = (0, 0, 0, 1))\)
\[ (\mathfrak{R}_x u_n(p))(z) = \int_{\mathbb{R}^3} (p_-(k) \varepsilon_4) h_n(k - p) e^{-\varepsilon \omega(k)} e^{i(\omega(k) t + k \cdot z)} d^3 k, \]
where \( h_n \) is a Dirac delta sequence. Apply definition \((27.5.53)\) to the vectors \( u_n^{(0)}, u_n^{(p)} \) and take the limit \( n \to \infty \). How can the arbitrariness of \( p \) be exploited in order to infer that \( x = y? \) Motivate your answer. Hint: Note that \( p_-(p) \) depends continuously on \( p \).

3. **Closedness:** The final step consists in proving that the local correlation function is closed, i.e. it maps closed sets to closed sets. In particular, it follows that \( F^\varepsilon(\mathbb{R}^4) \) is closed and that the inverse \( (F^\varepsilon)^{-1}|_{F^\varepsilon(\mathbb{R}^4)} \) is continuous. The identification is then complete.

   a. Let \( K \subset \mathcal{F} \) be compact and let \( \{x_n\} \subset H := (F^\varepsilon)^{-1}(K) \). By compactness of \( K \) there exists a subsequence \( \{y_n\} \subset H \) such that \( F^\varepsilon(y_n) \to A \in K \). Show that \( A \) is self-adjoint and different from zero. Hint: How could the comment after \((27.5.53)\) be exploited?

   b. Show that the subspace of \( \mathcal{H}_m \) of solutions of the form
\[ u_\varphi(x) := \int_{\mathbb{R}^3} d^3 k (p_-(k) \varphi(k)) e^{i(\omega(k) t + k \cdot x)}, \quad \text{with} \ \varphi \in S(\mathbb{R}^4, \mathbb{C}^4), \]
is dense. Deduce that there exists at least one \( \varphi \in S(\mathbb{R}^4, \mathbb{C}^4) \) such that \( \langle u_\varphi, A u_\varphi \rangle \neq 0 \).

   c. Convince yourself that \( e^{-\varepsilon \omega} p_\varphi \in S(\mathbb{R}^3, \mathbb{C}^4) \) for any \( \varphi \in S(\mathbb{R}^3, \mathbb{C}^4) \).
(d) It can be proved that the solutions of the Dirac equation with initial data in \( S(\mathbb{R}^3, \mathbb{C}^4) \) decay polynomially in both space and time direction. Use (2) and (3) to show that the sequence \( \{ y_n \} \) cannot be unbounded.

(e) Conclude that \( \{ x_n \} \) has a converging subsequence in \( H \).

Exercise 27.138. (On the differentiable manifold structure of regular points (4 points)) Let \( \mathcal{H} \) be a Hilbert space of finite dimension \( N \). The set \( \mathcal{F}^{reg} \) of regular points can be endowed with a differentiable structure. Precisely, let \( x \in \mathcal{F}^{reg} \). Choosing a Hilbert basis and using a block matrix representation in \( \mathcal{H} = S_x \oplus S_x^\perp \cong \mathbb{C}^{2n} \oplus \mathbb{C}^{N-2n} \), the operator \( x \) can be rewritten as

\[
 x = \begin{pmatrix} X & 0 \\ 0 & 0 \end{pmatrix} \in \text{Mat}(N-2n, \mathbb{C}), \quad X \in \text{Symm}(2n, \mathbb{C}) := \{ A \in \text{Mat}(2n, \mathbb{C}), \ A^\dagger = A \}, \quad (27.5.54)
\]

where \(^\dagger\) refers to the Euclidean scalar product. With this identification in mind, we now define

\[
 \Phi : (\text{Symm}(2n, \mathbb{C}) \oplus \text{L}(\mathbb{C}^{2n}, \mathbb{C}^{N-2n})) \cap B_{\varepsilon}(0) \to \mathcal{F}^{reg}
\]

\[
 (A, B) \mapsto \begin{pmatrix} X + A & B \\ B^\dagger & (X + A)^{-1} B \end{pmatrix}
\]

Prove the following statements.

(1) For sufficiently small \( \varepsilon \), \( \Phi \) is well-defined, continuous and injective.

(2) For sufficiently small \( \delta \), \( B_{\delta}(x) \subset \text{im} \Phi \) and the restriction \( \Phi^{-1}(B_{\delta}(x)) \to B_{\delta}(x) \) is homeomorphism.

\( \text{Hint: With the help of a unitary operator } U \text{ diagonalize any } y \in \mathcal{F}^{reg} \text{ as in } (27.5.54). \)

\( \text{Exploit this to show that } y \text{ must take the form } (27.5.55), \text{ if its distance from } x \text{ is sufficiently small.} \)

(3) Could one repeat the exercise if the Euclidean inner product is replaced by the canonical spin inner product on \( \mathbb{C}^{2n} \) ? Motivate your answer.

Exercise 27.139. (A distinguished Riemannian metric on the regular points (2 points)) Let \( \mathcal{H} \) be a Hilbert space of finite dimension \( N \) and let \( n < N/2 \) denote the spin dimension. The set of regular points \( \mathcal{F}^{reg} \) inherits from the ambient space \( L(\mathcal{H}) \cong \mathbb{R}^{2(N \times N)} \) a smooth manifold structure of dimension \( 4n(N-n) \). For any \( x \in \mathcal{F}^{reg} \) we introduce the bilinear form

\[
 h_x : T_x \mathcal{F}^{reg} \times T_x \mathcal{F}^{reg} \ni (u, v) \mapsto \text{tr}(uv) \in \mathbb{R}. \quad (27.5.56)
\]

The goal of this exercise is to show that \( h \) is a well-defined Riemannian metric on \( \mathcal{F}^{reg} \).

(1) Show that every \( u \in T_x \mathcal{F}^{reg} \subset L(\mathcal{H}) \) is self-adjoint.

(2) Show that \( (27.5.56) \) is a positive-definite inner product which depends smoothly on \( x \).

Exercise 27.140. (Variations of the universal measure and the physical wave functions) Let \( (\mathcal{H}, \mathcal{F}, \rho) \) be a causal fermion system with \( M := \text{supp } \rho \). Let \( F_\tau : M \to \mathcal{F} \) with \( \tau \in (-\delta, \delta) \) be a family of continuous functions which satisfy the following conditions,

\[
 F_0 = I \quad \text{and} \quad F_\tau |_{M \cap K} \equiv I \quad \text{with } K \subset M \text{ compact.}
\]

Moreover, we assume that \( F_\tau \) is differentiable in \( \tau \) and that all points of \( K \) are regular. This gives rise to a variation of the universal measure by defining \( \rho_\tau := (F_\tau)_* \rho \). This exercise explains how such a variation can be realized by a variation of the physical wave functions.

(1) Fix \( x \in K \). Show that by decreasing \( \delta \), one can arrange that the operators \( F_\tau(x) \) have maximal rank \( 2n \) for all \( \tau \in (-\delta, \delta) \)

(2) We introduce the spin spaces

\[
 S^x_\tau := \text{Im} \ F_\tau(x), \quad \langle \cdot, | \cdot \rangle^x_\tau := -\langle \cdot | F_\tau(x) \cdot \rangle_\mathcal{H} | S^x_\tau \times S^x_\tau.
\]

Construct a family of isometries (w.r.t. the spin inner products)

\[
 V_\tau(x) : (S_x, \langle \cdot, | \cdot \rangle_x) \to (S^x_\tau, \langle \cdot, | \cdot \rangle^x_\tau)
\]
which is differentiable in \( \tau \). \textit{Hint: For example, one can work with the orthogonal projections in } \( \mathcal{K} \) \textit{and take the polar decomposition with respect to the spin scalar products.}

(3) Consider the variation of the wave evaluation operator given by
\[
\Psi_\tau(x) := V_\tau(x)^{-1} \pi_{F_\tau(x)} : \mathcal{K} \to S_x.
\]
Show that \( F_\tau(x) = -\Psi_\tau(x)^*\Psi_\tau(x) \).

(4) So far, the point \( x \in K \) was fixed. We now extend the construction so that \( x \) can be varied. Use a compactness argument to show that there is \( \delta > 0 \) such that the operators \( F_\tau(x) \) have maximal rank \( 2n \) for all \( \tau \in (-\delta, \delta) \) and all \( x \in K \). Show that the mappings \( V_\tau(x) \) can be introduced such that they depend continuously in \( x \) and are differentiable in \( \tau \).

\[\text{Exercise 27.141. (Unitary variations in the Krein space)}\]

Let \( (\mathcal{K}, \mathcal{F}, \rho) \) be a causal fermion system with corresponding Krein space \( (\mathcal{K}, \langle \cdot, \cdot \rangle) \). Assume that \( \psi^u \in \mathcal{K} \) for all \( u \in \mathcal{K} \), so that the evaluation operator \( \Psi : \mathcal{K} \to \mathcal{K} \) is well-defined. Let \( U_\tau \) be a continuous family of unitary operators on \( \mathcal{K} \). Setting
\[
\Psi_\tau := U_\tau \circ \Psi : \mathcal{K} \to \mathcal{K},
\]
we obtain a corresponding variation of the physical wave functions, which gives rise to a variation of the universal measure by setting
\[
\rho_\tau := (F_\tau)_* \rho \quad \text{with} \quad F_\tau(x) := -\Psi_\tau(x)^*\Psi_\tau(x).
\]
Show that the volume and the trace constraints are satisfied by these variations.

\[\text{Exercise 27.142. (Jets in Minkowski Space (6 points))}\]

Let \( \mathcal{K}_m \) the Hilbert space of negative-energy solutions of the Dirac equation. We introduce a \textit{regularization operator} by inserting a cutoff in momentum space. More precisely, let \( \chi \) denote the characteristic function of the unit ball in \( \mathbb{R}^3 \). Then, on the dense subset \( \mathcal{D} \subset \mathcal{K}_m \) of smooth solutions \( u \) of the form (cf. Exercise 1.2)
\[
u(x) := \int_{\mathbb{R}^3} d^3k p_-(k) \varphi(k) e^{i(\omega(k)t+k \cdot x)}, \quad \text{with} \quad \varphi \in \mathcal{S}(\mathbb{R}^3, \mathbb{C}^4),
\] \[(27.5.57)\]
we define \( \mathcal{R}_\varepsilon u \in \mathcal{K}_m \) by cutting off the momenta larger than \( \varepsilon^{-1} \), i.e.
\[
(\mathcal{R}_\varepsilon u)(x) := \int_{\mathbb{R}^3} d^3k \chi(\varepsilon k) p_-(k) \varphi(k) e^{i(\omega(k)t+k \cdot x)}.
\]
(i) Show that \( \mathcal{R}_\varepsilon : \mathcal{D} \to \mathcal{K}_m \) is continuous, hence it can be uniquely extended to all of \( \mathcal{K}_m \).
Show that such extension is self-adjoint and idempotent.
(ii) Define \( \mathcal{H}^\varepsilon := \text{Im} \mathcal{R}_\varepsilon \) and \( F^\varepsilon : \mathbb{R}^4 \to \mathcal{F} \) as in Exercise 6.2-(1). Show that \( \text{Im} F^\varepsilon(x) \subset \mathcal{H}^\varepsilon \).

For any \( a \in \mathbb{R}^4 \) we define the \textit{translation operators}
\[
U_a : \mathcal{K}_m \to \mathcal{K}_m, \quad (U_a u)(x) := u(x + a).
\]
(iii) Working in momentum space with solutions of the form \[(27.5.57)\] show that \( U \) is well-defined and that it gives rise to a strongly-continuous (i.e. \( U_a u \to u \) as \( a \to 0 \) for all \( u \)) unitary representation of the translation group \( \mathbb{R}^4 \).

(iv) Show that \( U_a(\mathcal{H}^\varepsilon) \subset \mathcal{H}^\varepsilon \) and that the self-adjoint generators of \( U \) are bounded on \( \mathcal{H}^\varepsilon \), i.e. show that
\[
\mathcal{H}^\varepsilon \ni u \mapsto \frac{d}{ds}\bigg|_0 U_{se_j} u \in \mathcal{H}^\varepsilon
\]
is a well-defined bounded operator for any \( j = 0, 1, 2, 3 \). \textit{Hint: Work on the dense subset } \( \mathcal{R}_\varepsilon(\mathcal{D}) \subset \mathcal{H}^\varepsilon \).

(v) Show that \( [U_a, \mathcal{R}_\varepsilon] = 0 \) and \( U_a^* F^\varepsilon(x) U_a = F^\varepsilon(x + a) \).
there exists a unique vector field $X$ with initial condition $\omega_n$.

These transformation can be described infinitesimally by jets $u_j = (0, u_j)$. Determine $u_j$ and show that they solve the weak Euler-Lagrange equations.

**Exercise 27.143. (In preparation for Moser Theorem - part 1)** Let $M$ be a connected, oriented and closed manifold of dimension $n$ and $\omega \in \Omega^n M$ a volume form. It can be proved that

$$\int_M \omega = 0 \quad \text{if and only if there exists } \eta \in \Omega^{n-1} M \text{ such that } \omega = d\eta.$$ 

One direction follows immediately from Stokes Theorem. The goal of this exercise is to understand the idea behind the proof of the other direction by working in a very simple setting. More precisely, you are asked to prove the following statement: Let $n = 1, 2$ and $\omega \in \Omega^n \mathbb{R}^n$ a differentiable form with compact support contained in $(0,1)^n$. Show that there exists a differential form $\eta \in \Omega^{n-1} \mathbb{R}^n$, again with compact support in $(0,1)^n$, such that $\omega = d\eta$.

**Hint:** The case $n = 1$ is immediate. The case $n = 2$ is a bit trickier. The proof boils down to finding a smooth function $F \in C^\infty((0,1)^2, \mathbb{R})$ such that $\partial_1 F^2 - \partial_2 F^1 = \omega(\partial_1, \partial_2)$. Try to reduce the problem to the one-dimensional case by treating one variable as a parameter. Taking an average over $(0,1)$ may be useful.

**Exercise 27.144. (In preparation for Moser Theorem - part 2)** Let $M$ be a $n$-dimensional manifold and $\omega \in \Omega^n M$ a volume form. It can be proved that for any $\eta \in \Omega^{n-1} M$ there exists a unique vector field $X \in TM$ such that $\iota_X \omega = \eta$. Once again, the goal of this exercise is to understand the proof by working in very specific setting. More precisely, you are asked to prove the statement in the case $n = 2$.

**Hint:** How could the Fréchet-Riesz representation theorem be exploited here? Differentiability may be proved by working with local trivializations.

**Exercise 27.145. (In preparation for Moser theorem - part 3)** Let $M$ be an $n$-dimensional compact manifold. A time-dependent vector field is a differentiable function $X : \mathbb{R} \times M \to TM$ such that $X(t,p) \in T_p M$ for all $t \in \mathbb{R}$ and $p \in M$. Its flow is the unique differentiable mapping $\alpha : \mathbb{R} \times M \to M$ with the property that $\alpha(t,\cdot)p$ is the integral curve of $X$ with initial condition $p$, i.e. $\alpha(0, p) = p$ and $\alpha(t, p) = X(t, \alpha(t, p))$ for all $t \in \mathbb{R}$. We use the simpler notation $X_t := X(t, \cdot)$ and $\alpha_t := \alpha(t, \cdot)$.

1. It can be proved that, for any differential form $\omega \in \Omega^k M$,

$$\frac{d}{ds} \bigg|_{t} (\alpha^*_s \omega) = \alpha^*_t (\mathcal{L}_X \omega), \quad (27.5.58)$$

where $\mathcal{L}_\xi$ is the Lie derivative along the vector field $\xi$ and $^*$ denotes the pull-back. Check identity (27.5.58) for $\omega$ of the form $f, dh$ and $f dh$, where $f, h \in C^\infty(M)$.

**Hint:** How does $X_t$ act on $f$? Look up the definition and the properties of the pull-back and the Lie derivative. In particular, remember that $d$ commute with $^*$, $\mathcal{L}_\xi$ and $d/ds$.

2. Let $d\rho \in \Omega^n M$ be a volume form. Assume the existence of two differentiable families

$$\Phi : \mathbb{R} \to \text{Diff}(M) \quad \text{and} \quad \varphi : \mathbb{R} \to C^\infty(M), \quad \text{with } \Phi_0 = \text{Id}_M \quad \text{and} \quad \varphi_0 \equiv 1,$$

such that

$$d\rho = (\Phi_t)^* (\varphi_t \ d\rho) \quad \text{for all } t \in \mathbb{R}. \quad (27.5.59)$$

We now introduce the time-dependent vector field

$$X(t, p) := \frac{d}{ds} \bigg|_{t} \Phi_s (\varphi_t^{-1}(p)) \in T_p M.$$ 

Show that $\Phi$ is the flow of $X$ and that the couple $(-\varphi_0, X_0)$ defines an inner solution.
Hint: Differentiate (27.5.59) through using the chain rule. What is the connection between the divergence and the Lie derivative?

Exercise 27.146. (On inner solutions) Let \( d\rho \in \Omega^n M \) be a volume form and consider differentiable one-parameter families \( h_t \in C^\infty(M, \mathbb{R}^+) \), \( \Phi_t \in \text{Diff}(M) \), with the property that
\[
\dot{h}_0 = 0, \quad \Phi_0 = \text{Id}_M, \quad h_0 d\rho = (\Phi_t)^*(h_t d\rho) \quad \text{for all } t \in \mathbb{R}.
\] (27.5.60)

Rewrite (27.5.60) as in Exercise 8.4-(ii) and show that the corresponding inner solution depends only on \( h_0 \) and \( X_0 \).

Exercise 27.147. (Moser’s theorem) Let \( M \) be an connected, oriented and closed manifold of dimension \( n \). Let \( \omega_0, \omega_1 \in \Omega^n M \) be volume forms with the property that
\[
\int_M \omega_0 = \int_M \omega_1.
\]
Find a diffeomorphism \( \Phi \in \text{Diff}(M) \) such that \( \omega_0 = \Phi^* \omega_1 \).

Hint: First, apply Exercise 8.2. Find then a simple differentiable curve \( \omega_s \) of volume forms connecting \( \omega_0 \) and \( \omega_1 \) and apply Exercise 8.3, obtaining in this way a time-dependent vector field. Using Exercise 8.4 show that the pull-back of \( \omega_s \) through the flow is constant. What is then \( \Phi \)?

Exercise 27.148. (Nonlinear surface layer integral) We now use Moser’s theorem to prove conservation of the nonlinear surface layer integral in the compact setting. To this aim, let \( \mathcal{F} \) be a smooth manifold and \( \mathcal{L} \in C^\infty(\mathcal{F} \times \mathcal{F}, \mathbb{R}^+) \). Furthermore, let \( M_0, M_1 \) be two connected, oriented and closed submanifolds of \( \mathcal{F} \) of dimension \( n \), and let \( F \in \text{Diff}(M_0, M_1) \). Finally, let \( d\rho_0 \in \Omega^n M_0 \) and \( d\rho_1 \in \Omega^n M_1 \) be two volume forms. The nonlinear surface layer integral is defined for any relatively compact \( \Omega \subset M_0 \) by
\[
\gamma^\Omega_F(\rho_0, \rho_1) := \int_{F(\Omega)} d\rho_1(x) \int_{M_0 \setminus \Omega} d\rho_0(y) \mathcal{L}(x, y) - \int_{\Omega} d\rho_0(x) \int_{M_1 \setminus F(\Omega)} d\rho_1(y) \mathcal{L}(x, y),
\]
where \( \rho_0 \) and \( \rho_1 \) are the measures induced by the volume forms on the corresponding manifolds.

The goal of this exercise is the following: Show that \( F \) can be modified in a way that \( \gamma^\Omega_F(\rho_0, \rho_1) = 0 \) for all relatively compact \( \Omega \subset M_0 \). To this aim, let us introduce the volume forms
\[
d\nu_0(x) := \left( \int_{M_1} \mathcal{L}(x, y) d\rho_1(y) \right) d\rho_0(x) \quad \text{and} \quad d\nu_1(x) := \left( \int_{M_0} \mathcal{L}(x, y) d\rho_0(y) \right) d\rho_1(x)
\]

(1) Using the symmetry of the Lagrangian \( \mathcal{L}(x, y) = \mathcal{L}(y, x) \) show that, for all \( \Omega \),
\[
\gamma^\Omega_F(\rho_0, \rho_1) = \int_{\Omega} (F^* d\nu_1 - d\nu_0).
\]

(2) Using Exercise 9.1 show that there is \( \Phi \in \text{Diff}(M_0) \) such that \( (F \circ \Phi^{-1})^* d\nu_1 = d\nu_0 \).

Hint: Note that \( M_0 \) and \( M_1 \) are different manifolds! How do you apply Exercise 9.1? Points (i) and (ii) together give the claim: By modifying \( F \) to \( F \circ \Phi^{-1} \) we obtain a conservation law for the nonlinear surface layer integral.
APPENDIX A

The Spin Coefficients

In this appendix, we verify by explicit computation that the matrices $E_j$ containing the spin coefficients as given by (4.2.24),

$$E_j = \frac{i}{2} \Gamma \partial_j \Gamma - \frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n + \frac{i}{8} \text{Tr} \left( \Gamma G_j \nabla_m G^m \right) \Gamma,$$  

(A.0.1)

have the following behavior under gauge transformations:

$$E_j \rightarrow UE_j U^{-1} \quad \text{for U(1) gauge transformations} \quad (A.0.2)$$

$$E_j \rightarrow UE_j U^{-1} + iU(\partial_j U^{-1}) \quad \text{for SU(2, 2) gauge transformations} \quad (A.0.3)$$

Under U(1) gauge transformations, all the terms in (A.0.1) remain unchanged because $U$ and its partial derivatives commute with $\Gamma$ as well as with the $G^j$. Therefore, the relation (A.0.2) is obvious. Thus it remains to consider SU(2, 2) gauge transformations. Our goal is to verify (A.0.3) soll for a fixed spacetime point $p$.

We decompose the gauge transformation $U$ as $U = U_2 U_1$ with

$$U_1(x) = U(p)$$

$$U_2(x) = U(x) U^{-1}(p).$$

Being constant, the first transformation clearly satisfies the transformation law (A.0.3). Therefore, it suffices to consider a gauge transformation $U$ with $U(p) = 1$. Then (A.0.3) can be written as

$$\tilde{E}_j = E_j - i \partial_j U.$$  

(A.0.4)

We now compute the transformation law of each of the summands in (A.0.1) after each other:

(1) $\frac{i}{2} \Gamma \partial_j \Gamma$
   
   (i) odd transformations:
   
   $$\frac{i}{2} \tilde{\Gamma} \partial_j \tilde{\Gamma} = \frac{i}{2} \Gamma (\partial_j \Gamma + [\partial_j U, \Gamma])$$
   
   $$= \frac{i}{2} \Gamma \partial_j \Gamma + \frac{i}{2} \Gamma \left( (\partial_j U) \Gamma - \Gamma (\partial_j U) \right)$$
   
   $$= \frac{i}{2} \Gamma \partial_j \Gamma - i \Gamma \partial_j U$$
   
   $$= \frac{i}{2} \Gamma \partial_j \Gamma - i \partial_j U$$

   (ii) even transformations:
   
   $$\frac{i}{2} \tilde{\Gamma} \partial_j \tilde{\Gamma} = \frac{i}{2} \Gamma \left( \partial_j \Gamma + [\partial_j U, \Gamma] \right) = \frac{i}{2} \Gamma \partial_j \Gamma$$

Thus for odd transformations, we get the correct transformation law, whereas for even transformations the desired term $i \partial_j U$ is still missing.

(2) $-\frac{i}{16} \text{Tr}(G^m \nabla_j G^n) G_m G_n$
(i) odd transformations:

\[ -\frac{i}{16} \text{Tr} \left( \tilde{G}^m \nabla_j \tilde{G}^n \right) \tilde{G}_m \tilde{G}_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - \frac{i}{16} \text{Tr} \left( G^m \left[ \partial_j U, G^n \right] \right) G_m G_n \quad (A.0.5) \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n , \quad (A.0.6) \]

where we used that \( G^m, G^n \) and \( \partial_j U \) are odd, implying that the trace in the last summand in (A.0.5) vanishes.

(ii) \( \partial_j U = i\sigma_{kl} \) for arbitrary indices \( k,l \):

\[ -\frac{i}{16} \text{Tr} \left( \tilde{G}^m \nabla_j \tilde{G}^n \right) \tilde{G}_m \tilde{G}_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - \frac{i}{16} \text{Tr} \left( G^m \left[ i\sigma_{kl}, G^n \right] \right) G_m G_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - \frac{i}{16} \text{Tr} \left( \left[ G^n, G^m \right] i\sigma_{kl} \right) G_m G_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - \frac{i}{8} \text{Tr} \left( \sigma^{mn} i\sigma_{kl} \right) \sigma_{mn} \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - i\partial_j U \]

(iii) \( \partial_j U = \Gamma \):

\[ -\frac{i}{16} \text{Tr} \left( \tilde{G}^m \nabla_j \tilde{G}^n \right) \tilde{G}_m \tilde{G}_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - \frac{i}{16} \text{Tr} \left( G^m \left[ \Gamma, G^n \right] \right) G_m G_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n - \frac{i}{16} \text{Tr} \left( \left[ G^n, G^m \right] \right) G_m G_n \]

\[ = -\frac{i}{16} \text{Tr} \left( G^m \nabla_j G^n \right) G_m G_n , \]

because \( \text{Tr} \left( \Gamma \sigma^{mn} \right) = 0 \) for all \( m,n \).

Thus we get the correct transformation law for bilinear transformations \( \partial_j U = i\sigma_{kl} \).

(3) \[ \frac{i}{8} \text{Tr} \left( \Gamma G_j \nabla_m G^n \right) \Gamma \]

\[ = \frac{i}{8} \text{Tr} \left( \tilde{\nabla} G_j \nabla_m \tilde{G}^n \right) \tilde{\Gamma} + \frac{i}{8} \text{Tr} \left( \Gamma G_j \left[ \partial_m U, G^n \right] \right) \Gamma \]

\[ = \frac{i}{8} \text{Tr} \left( \Gamma G_j \nabla_m G^n \right) \Gamma + \frac{i}{8} \text{Tr} \left( \partial_m U \left[ G^n, G_j \right] \right) \Gamma \]

\[ = \frac{i}{8} \text{Tr} \left( \Gamma G_j \nabla_m G^n \right) \Gamma - \frac{i}{4} \text{Tr} \left( \partial_m U \Gamma \delta_j^m \right) \Gamma \]

\[ = \frac{i}{8} \text{Tr} \left( \Gamma G_j \nabla_m G^n \right) \Gamma - \frac{i}{4} \text{Tr} \left( \left( \partial_j U \right) \Gamma \right) \Gamma , \]

where we used the relations

\[ [G^j, \Gamma G^k] = -\Gamma \left\{ G^j, G^k \right\} \]

as well as the anti-commutation relations for Dirac matrices. We again distinguish different cases:

(i) \( \partial_j U \) is odd or \( \partial_j U = i\sigma_{kl} \):

\[ \frac{i}{8} \text{Tr} \left( \tilde{\nabla} G_j \nabla_m \tilde{G}^n \right) \tilde{\Gamma} = \frac{i}{8} \text{Tr} \left( \Gamma G_j \nabla_m G^n \right) \Gamma \]
(ii) $\partial_j U = \Gamma$:

$$\frac{i}{8} \text{Tr} \left( \bar{\Gamma} \, \bar{G}_j \, \nabla_m \bar{G}^m \right) \bar{\Gamma} = \frac{i}{8} \text{Tr} \left( \Gamma \, G_j \, \nabla_m G^m \right) \Gamma - i \partial_j U$$

Hence we get the correct transformation law if $\partial_j U = \Gamma$.

Adding all the terms gives the desired transformation law \[A.0.4\].
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