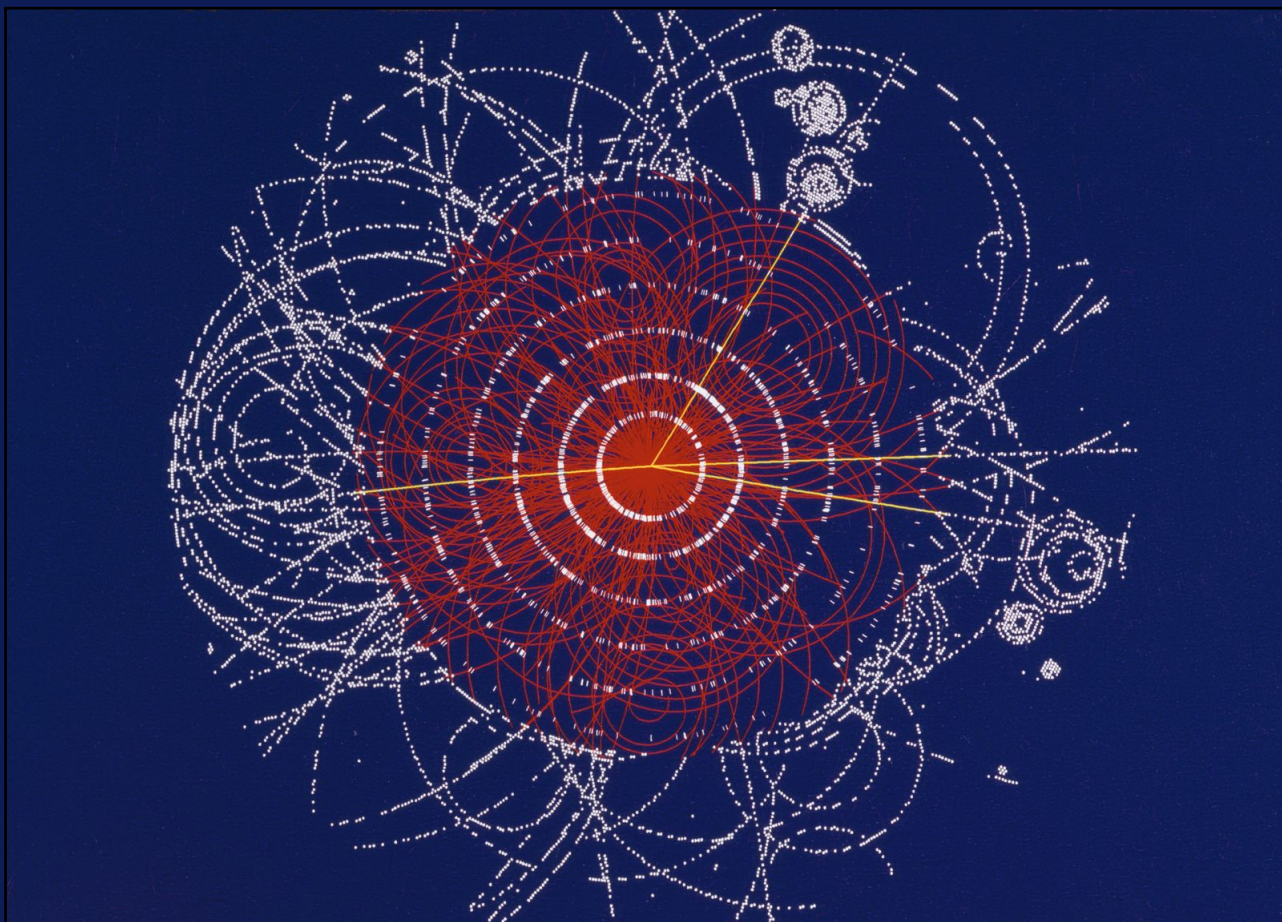


Quantum Field Theory

Leonardo Cerasi



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Notation

Conventions

In these notes, the Lorentz–Minkowski metric $\eta_{\mu\nu}$ has signature $(+, -, -, -)$ and Greek indices generally run over spacetime coordinates, while Latin indices are general \mathbb{N}_0 -indices defined in each context. Repeated indices are generally summed over, unless otherwise specified. The n -dimensional Levi–Civita symbol $\epsilon^{i_1 \dots i_n}$ is defined with the convention $\epsilon^{01 \dots n} = +1$.

All quantities are expressed in natural units, defined by $\hbar \equiv c \equiv 1$: in particular, the “dimensions” of a quantity generally refers to its mass dimension.

The normal ordering and the time ordering of an expression are respectively denoted by the normal-ordering operator \mathfrak{N} and by the time-ordering operator \mathfrak{T} .

Given $\alpha \in \mathbb{C}^{n \times n}$, with $n \in \mathbb{N}_0$, its complex conjugate is denoted as α^* , its transpose as α^\top and its Hermitian conjugate as $\alpha^\dagger := (\alpha^*)^\top$. Given a Dirac spinor $\Psi \in \mathbb{C}^4$, its Dirac dual (or Dirac adjoint) is defined as $\bar{\Psi} := \Psi^\dagger \gamma^0$.

The Landau symbol for a function $f : D \subseteq \mathbb{R} \rightarrow \mathbb{C}$ is defined by the condition $\exists M \in \mathbb{R} : |o(f(x))| \leq M |f(x)| \ \forall x \in D$.

Mathematical notation

The empty set is denoted by \emptyset and the power set of a set A by $\mathcal{P}(A) := \{B : B \subseteq A\}$. The counting numbers are $\mathbb{N} \equiv \{1, 2, 3, \dots\}$, and the natural numbers are defined by $\mathbb{N}_0 \equiv \{0\} \cup \mathbb{N}$. The imaginary unit is denoted by i and the unit quaternions by i, j, k , so that $\mathbb{C}(\mathbb{R}) = \text{span}(1, i)$ and $\mathbb{H}(\mathbb{R}) = \text{span}(1, i, j, k)$.

The permutation group of n objects, i.e. the n^{th} symmetric group, is denoted by S_n .

Given two \mathbb{K} -vector spaces V and W , with \mathbb{K} a generic field, the space of all \mathbb{K} -linear applications $f : V \rightarrow W$ is denoted by $\text{Hom}_{\mathbb{K}}(V, W)$: in particular, $\text{Hom}_{\mathbb{K}}(V) \equiv \text{End}(V)$. The subset of $\text{End}(V)$ of all automorphisms of V is the automorphism group $\text{Aut}(V)$, which is a group under composition of morphisms.

Given a manifold \mathcal{M} , the space of all smooth scalar functions on \mathcal{M} is denoted by $\mathcal{C}^\infty(\mathcal{M})$, the space of all vector fields on \mathcal{M} by $\mathfrak{X}(\mathcal{M})$, the space of all p -forms on \mathcal{M} by $\bigwedge^p(\mathcal{M})$ and the Grassmann algebra of \mathcal{M} by $\bigwedge(\mathcal{M}) := \bigoplus_{k=0}^n \bigwedge^k(\mathcal{M})$.

The exterior derivative is denoted by d , the partial derivative by ∂ , the nabla operator by $\nabla \equiv (\partial_1, \partial_2, \partial_3)$, the Laplacian operator by $\Delta \equiv \nabla^2$ and the D’Alembert operator by $\square \equiv \partial_0^2 - \Delta$.

A list of “important” Lie groups:

$\text{GL}(n, \mathbb{K}) := \{A \in \mathbb{K}^{n \times n} : \det A \neq 0\}$ general linear group (Lie group for $\mathbb{K} = \mathbb{R}, \mathbb{C}$)

$\text{SL}(n, \mathbb{K}) := \{A \in \text{GL}(n, \mathbb{K}) : \det A = 1\}$ special linear group (Lie group for $\mathbb{K} = \mathbb{R}, \mathbb{C}$)

$\text{O}(n) := \{A \in \mathbb{R}^{n \times n} : AA^\top = A^\top A = I_n\}$ orthogonal group

$\text{SO}(n) := \{A \in \text{O}(n) : \det A = 1\}$ special orthogonal group

$\text{U}(n) := \{A \in \mathbb{C}^{n \times n} : AA^\dagger = A^\dagger A = I_n\}$ unitary group

$\text{SU}(n) := \{A \in \text{U}(n) : \det A = 1\}$ special unitary group

Given a Lie group G , its associated Lie algebra is denoted by \mathfrak{g} .

Physical notation

The Lorentz group is denoted by $\text{SO}^+(1, 3)$, the Lorentz algebra by $\mathfrak{so}^+(1, 3)$, the Poincaré group by $\text{ISO}^+(1, 3)$, the Poincaré algebra by $\mathfrak{iso}^+(1, 3)$ and the Dirac algebra by $\mathfrak{cl}_{1,3}(\mathbb{C})$.

Hilbert and Fock spaces are generally denoted respectively by \mathcal{H} and \mathcal{F} , the Hamiltonian and the Lagrangian of a system by H and L , and the Hamiltonian density and the Lagrangian density of a field theory by $\mathcal{H} : H = \int d^3x H$ and $\mathcal{L} : L = \int d^3x L$. The action functional is defined as $\mathcal{S} := \int dt L = \int d^4x \mathcal{L}$.

The matrix element of a scattering process is generally denoted by \mathcal{M} , and its amplitude by $\mathcal{A} \equiv |\mathcal{M}|^2$.

The parity operator is denoted by \mathcal{P} , the charge-conjugation operator by \mathcal{C} and the time-reversal operator by \mathcal{T} .

For a general gauge theory, the covariant derivative is denoted by \mathcal{D}_μ , the connection (or gauge field) by \mathcal{A}_μ and the field-strength tensor by $\mathcal{F}_{\mu\nu}$. For the particular case of QED, these quantities are denoted respectively by D_μ , A_μ and $F_{\mu\nu}$.

The Feynman propagator in position and momentum space is denoted by $\Delta(x, y) \equiv \Delta(x - y)$ and $\tilde{\Delta}(p)$, the Dirac propagator by $S(x, y) \equiv S(x - y)$ and $\tilde{S}(p)$, and the photon propagator by $\Delta_{\mu\nu}(x, y) \equiv \Delta_{\mu\nu}(x - y)$ and $\tilde{\Delta}_{\mu\nu}(p)$.

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Part I

Field Theory

Classical Field Theory

§1.1 Continuous limit

§1.1.1 One-dimensional harmonic crystal

Consider a simple one-dimensional model of a crystal where atoms of mass $m \equiv 1$ lie at rest on the x -axis, with equilibrium positions labelled by $n \in \mathbb{N}$ and equally spaced by a distance a . Assuming these atoms are free to vibrate only in the x direction (longitudinal waves), and denoting the displacement of the atom at position n as η_n , one can write the Lagrangian for a **harmonic crystal** as:

$$L = \sum_n \left[\frac{1}{2} \dot{\eta}_n^2 - \frac{\lambda}{2} (\eta_n - \eta_{n-1})^2 \right] \quad (1.1)$$

where λ is the spring constant. From the Lagrange equations, the classical equations of motions are:

$$\ddot{\eta}_n = \lambda (\eta_{n+1} - 2\eta_n + \eta_{n-1}) \quad (1.2)$$

The solutions can be written as complex travelling waves:

$$\eta_n(t) = e^{i(kn - \omega t)} \quad (1.3)$$

where the dispersion relation is:

$$\omega^2 = 2\lambda (1 - \cos k) \approx \lambda k^2 \quad (1.4)$$

Therefore, in the long-wavelength limit $k \ll 1$ waves propagate with velocity $c = \sqrt{\lambda}$. To determine the normal modes, there need to be boundary conditions: imposing boundary conditions:

$$\eta_{n+N} = \eta_n \quad \implies \quad k_m = \frac{2\pi m}{N}, \quad m = 0, 1, \dots, N-1 \quad (1.5)$$

The normal-mode expansion can then be written as:

$$\eta(t) = \sum_{m=0}^{N-1} [\mathcal{A}_m e^{i(k_m n - \omega_m t)} + \mathcal{A}_m^* e^{-i(k_m n - \omega_m t)}] \quad (1.6)$$

where the complex conjugate is added to ensure that the total displacement is real. The momentum canonically-conjugated to the displacement is defined as:

$$\pi_n := \frac{\partial L}{\partial \dot{\eta}_n} = \dot{\eta}_n \quad (1.7)$$

In quantum mechanics, η_n and Π_n become operators with canonical commutator $[\hat{\eta}_j, \hat{\pi}_k] = i\hbar\delta_{jk}$. Implementing time evolution with the Heisenberg picture¹:

$$[\hat{\eta}_j(t), \hat{\pi}_k(t)] = i\hbar\delta_{jk} \quad (1.8)$$

The commutator of operators evaluated at different times requires solving the dynamics of the system. It is useful to introduce **annihilation** and **creation operators**² $\hat{a}(t)$ and $\hat{a}^\dagger(t)$, so that Eq. 1.6 becomes:

$$\hat{\eta}_n(t) = \sum_{m=0}^{N-1} \sqrt{\frac{\hbar}{2\omega_m}} \frac{1}{\sqrt{N}} [e^{i(k_m n - \omega_m t)} \hat{a}_m + e^{-i(k_m n - \omega_m t)} \hat{a}_m^\dagger] \quad (1.9)$$

where $[\hat{a}_j, \hat{a}_k^\dagger] = \delta_{jk}$ and the $N^{-1/2}$ ensures the normalization of normal modes. The proof of Eq. 1.8 follows from the finite Fourier series identity (sum of a geometric progression):

$$\sum_{m=0}^{N-1} e^{ik_m(n-n')} = N\delta_{nn'} \quad (1.10)$$

The Hamiltonian of the system can be written as:

$$\hat{H} = \sum_n \left[\frac{1}{2} \hat{\pi}_n^2 + \frac{\lambda}{2} (\hat{\eta}_n - \hat{\eta}_{n-1})^2 \right] = \sum_{m=0}^{N-1} \hbar\omega_m \left(\hat{a}_m^\dagger \hat{a}_m + \frac{1}{2} \right) \quad (1.11)$$

Generalizing the harmonic oscillator operator algebra (proven unique by Von Neumann), one can construct the Hilbert space for the harmonic crystal as:

$$\hat{a}_m |0\rangle \quad \forall m = 0, 1, \dots, N-1 \quad (1.12)$$

$$|n_0, n_1, \dots, n_{N-1}\rangle = \prod_{m=0}^{N-1} \frac{(\hat{a}_m^\dagger)^{n_m}}{\sqrt{n_m!}} |0\rangle \quad (1.13)$$

These are normalized eigenstates of Eq. 1.1 with energy eigenvalues:

$$E_0 = \frac{1}{2} \sum_{m=0}^{N-1} \hbar\omega_m \quad (1.14)$$

$$E_{n_0, n_1, \dots, n_{N-1}} = E_0 + \sum_{m=0}^{N-1} n_m \hbar\omega_m \quad (1.15)$$

This Hilbert space is called **Fock space** and the excited states **phonons**: these can be thought as “particles” and n_m is the number of phonons in the m^{th} normal mode.

¹Recall that $\hat{\mathcal{O}}(t) = e^{\frac{i}{\hbar}\hat{H}t}\hat{\mathcal{O}}(0)e^{-\frac{i}{\hbar}\hat{H}t}$ and $\frac{d\hat{\mathcal{O}}}{dt} = \frac{i}{\hbar}[\hat{H}, \hat{\mathcal{O}}]$.

²For a harmonic oscillator $\hat{H} = \frac{1}{2}\hat{p}^2 + \frac{1}{2}\omega^2\hat{x}^2$, so $\frac{d\hat{x}}{dt} = \hat{p}(t)$ and $\frac{d\hat{p}}{dt} = -\omega^2\hat{x}(t)$ and the solution can be written as:

$$\hat{x}(t) = \sqrt{\frac{\hbar}{2\omega}} [\hat{a}(t) + \hat{a}^\dagger(t)] \quad \hat{p}(t) = -i\omega\sqrt{\frac{\hbar}{2\omega}} [\hat{a}(t) - \hat{a}^\dagger(t)]$$

Inverting these expressions one finds $[\hat{a}(t), \hat{a}^\dagger(t)] = 1$ and $\hat{H} = \hbar\omega (\hat{a}^\dagger\hat{a} + \frac{1}{2})$. The time evolution $\hat{a}(t) = e^{-i\omega t}\hat{a}(0)$ ensures that \hat{H} is time-independent.

§1.1.2 One-dimensional harmonic string

Taking the continuum limit, the crystal becomes a string: to achieve this, one takes the limits $a \rightarrow 0$ and $N \rightarrow \infty$ while keeping the total length $R \equiv Na$ fixed. In this context, the displacement becomes a field $\eta(x, t)$ dependent on the continuous real coordinate $x \in [0, R]$, therefore:

$$(\eta_{n+1} - \eta_n)^2 \longrightarrow a^2 \left(\frac{\partial \eta}{\partial x} \right)^2 \quad \sum_n \longrightarrow \frac{1}{a} \int_0^R dx$$

Proposition 1.1.1 (Dirac delta)

In the continuous limit:

$$\frac{\delta_{nn'}}{a} \longrightarrow \delta(x - x') = \int_{\mathbb{R}} \frac{dk}{2\pi} e^{ik(x-x')}$$

Proof. By direct calculation:

$$a \sum_n f(an) \frac{\delta_{nm}}{a} = f(ma) \longrightarrow f(y) = \int_0^R dx f(x) \delta(x - y)$$

Recalling Eq. 1.10, since $k_m n = \frac{k_m}{a} na \rightarrow kx$, symmetrizing $k_m \in [-\pi, \pi]$ (instead of $[0, 2\pi]$) one finds:

$$\delta(x - x') \longleftarrow \frac{\delta_{nn'}}{a} = \frac{1}{Na} \sum_m e^{ik_m(n-n')} \longrightarrow \int_{\mathbb{R}} \frac{dk}{2\pi} e^{ik(x-x')}$$

where integration limits are $\pm \frac{\pi}{a} \rightarrow \pm \infty$. □

Proposition 1.1.2

The inverse Fourier transform of the Dirac delta reads:

$$\int_0^R dx e^{i(k-k')x} = 2\pi \delta(k - k') \quad (1.16)$$

Proof. Denoting the Fourier transform by the operator \mathfrak{F} , the thesis is a consequence of $\mathfrak{F}^{-1}\mathfrak{F} = \mathbb{1}$. □

By these relations, it can be seen that $\frac{dk}{2\pi}$ has the physical meaning of the number of normal modes per unit spatial volume with wavenumber between k and $k + dk$, while the interpretation of the divergent $\delta(0)$ varies: in x space, it is the reciprocal of the lattice spacing, i.e. the number of normal modes per unit spatial volume, but in k space $2\pi\delta(0)$ is the (hyper-)volume of the system.

In the continuous limit, the Lagrangian of the harmonic string becomes:

$$L = \int_0^R dx \left[\frac{1}{2} \rho_0 (\partial_t \eta)^2 - \frac{\kappa}{2} (\partial_x \eta)^2 \right]$$

where ρ_0 is the equilibrium mass density of the string. It is customary to absorb constants in the fields, thus, setting $\phi(x, t) \equiv \sqrt{\rho_0} \eta(x, t)$ and $\kappa = c^2 \rho_0$ and adding a pinning term $\propto \phi^2$, the

Lagrangian can be written as:

$$L = \int_0^R dx \left[\frac{1}{2} (\partial_t \phi)^2 - \frac{c^2}{2} (\partial_x \phi)^2 - \frac{m^2 c^4}{2} \phi^2 \right] \quad (1.17)$$

The classical equation of motion of this field yields:

$$\partial_t^2 \phi = c^2 \partial_x^2 \phi - m^2 c^4 \phi \quad (1.18)$$

The solutions of this wave equation can be written as:

$$\phi(x, t) = e^{i(kx - \omega_k t)} \quad (1.19)$$

with dispersion relation:

$$\omega_k^2 = c^2 k^2 + m^2 c^4 \quad (1.20)$$

To quantize this system, one needs to compute the Hamiltonian. The canonical momentum field is:

$$\Pi(x, t) := \frac{\partial L}{\partial(\partial_t \phi)} = \partial_t \phi(x, t) \quad (1.21)$$

The classical Hamiltonian can then be found as:

$$H = \int_0^R dx \left[\frac{1}{2} \Pi^2 + \frac{c^2}{2} (\partial_x \phi)^2 + \frac{m^2 c^4}{2} \phi^2 \right] \quad (1.22)$$

The quantum field is analogous to [Eq. 1.9](#):

$$\hat{\phi}(x, t) = \int_{\mathbb{R}} \frac{dk}{2\pi} \sqrt{\frac{\hbar}{2\omega_k}} \left[e^{i(kx - \omega_k t)} \hat{a}_k + e^{-i(kx - \omega_k t)} \hat{a}_k^\dagger \right] \quad (1.23)$$

with commutation relations:

$$[\hat{a}_k, \hat{a}_{k'}^\dagger] = 2\pi \delta(k - k') \quad (1.24)$$

$$[\hat{\phi}(x, t), \hat{\Pi}(x', t)] = i\hbar \delta(x - x') \quad (1.25)$$

The quantum Hamiltonian can be written as:

$$\hat{H} = \int_{\mathbb{R}} \frac{dk}{2\pi} \frac{1}{2} \hbar \omega_k \left(\hat{a}_k^\dagger \hat{a}_k + \hat{a}_k \hat{a}_k^\dagger \right) = E_0 + \int_{\mathbb{R}} \frac{dk}{2\pi} \hbar \omega_k \hat{a}_k^\dagger \hat{a}_k \quad (1.26)$$

The ground-state energy can be computed from [Eq. 1.14](#), defining $\text{Vol} := 2\pi \delta(k = 0)$:

$$E_0 = \text{Vol} \int_{\mathbb{R}} \frac{dk}{2\pi} \frac{1}{2} \hbar \omega_k \quad (1.27)$$

For a strictly continuous system there is no cut-off in the k integral, thus the zero-point energy diverges: however, this is not necessarily a problem, as often only changes in E_0 are relevant (and experimentally accessible), and in this case it is known as **Casimir energy**.

§1.2 Spacetime symmetries

§1.2.1 Lorentz group

Consider the group of linear transformations $x^\mu \mapsto \Lambda^\mu{}_\nu x^\nu$ on $\mathbb{R}^{1,3}$ which leave invariant the quantity $\eta_{\mu\nu} x^\mu x^\nu$, i.e. the orthogonal group $O(1, 3)$. The condition that $\Lambda^\mu{}_\nu$ must satisfy reads:

$$\eta_{\rho\sigma} = \eta_{\mu\nu} \Lambda^\mu{}_\rho \Lambda^\nu{}_\sigma \quad (1.28)$$

This implies that $\det \Lambda = \pm 1$: a transformation with $\det \Lambda = -1$ can always be written as the product of a transformation with $\det \Lambda = 1$ and a discrete transformation which reverses the sign of an odd number of coordinates, hence one further defines the special orthogonal group $SO(1, 3) := \{\Lambda \in O(1, 3) : \det \Lambda = 1\}$.

Writing explicitly the temporal component $1 = (\Lambda^0{}_0)^2 - (\Lambda^1{}_0)^2 - (\Lambda^2{}_0)^2 - (\Lambda^3{}_0)^2$, it is clear that $(\Lambda^0{}_0)^2 \geq 1$. Therefore, $O(1, 3)$ has two disconnected components: the orthochronous component with $\Lambda^0{}_0 \geq 1$ and the non-orthochronous component with $\Lambda^0{}_0 \leq -1$. Any non-orthochronous transformation can be written as the product of an orthochronous transformation and a discrete transformation which reverses the sign of the temporal component.

Definition 1.2.1 (Lorentz group)

The **Lorentz group** $SO^+(1, 3)$ is the orthochronous component of $SO(1, 3)$.

The discrete transformations are factored out of the Lorentz group: these are **parity** and **time reversal**, represented as $\mathcal{P}^\mu{}_\nu = \text{diag}(+1, -1, -1, -1)$ and $\mathcal{T}^\mu{}_\nu = \text{diag}(-1, +1, +1, +1)$. Applying these discrete transformations in all combinations ($\mathbb{1}$, \mathcal{P} , \mathcal{T} and \mathcal{PT}) one gets the four disconnected components of $O(1, 3)$, which are not simply connected. This means that Lorentz invariance does not include parity and time reversal invariance.

Considering an infinitesimal Lorentz transformation and applying Eq. 1.28:

$$\Lambda^\mu{}_\nu = \delta^\mu{}_\nu + \omega^\mu{}_\nu \quad \implies \quad \omega_{\mu\nu} = -\omega_{\nu\mu}$$

Anti-symmetry means that $\omega_{\mu\nu}$ has only 6 parameters, which define the Lorentz group: these can be identified by the 3 angles of spherical rotations in the (x, y) , (y, z) and (z, x) planes and the 3 rapidities of hyperbolic rotations in the (t, x) , (t, y) and (t, z) planes.

Proposition 1.2.1

The Lorentz group is a non-compact Lie group.

Proof. Spherical and hyperbolic rotations are continuous and differentiable functions respectively of angles and of rapidities, but while angles vary in $[0, 2\pi)$, rapidities vary in \mathbb{R} , hence the differentiable manifold associated to $SO^+(1, 3)$ is not compact. \square

§1.2.1.1 Lorentz algebra

The 6 parameters of the Lorentz group correspond to 6 generators of the associated Lorentz algebra. Labelling these generators as $J^{\mu\nu} : J^{\mu\nu} = -J^{\nu\mu}$, the generic element $\Lambda \in SO^+(1, 3)$ can be written as:

$$\Lambda = e^{-\frac{i}{2} \omega_{\mu\nu} J^{\mu\nu}} \quad (1.29)$$

The $\frac{1}{2}$ factor arises from each generator being counted twice (product of two anti-symmetric objects). Given an n -dimensional representation of $\text{SO}^+(1, 3)$, then $[J^{\mu\nu}]^i_j \in \mathbb{C}^{n \times n}$ and $[\Lambda]^i_j \in \mathbb{R}^{n \times n}$: for example, the $n = 1$ representation acts on **scalars**, which are invariant under Lorentz transformations, so $J^{\mu\nu} \equiv 0 \ \forall \mu, \nu = 0, \dots, 3$ and $\Lambda \equiv 1$.

4-vectors The $n = 4$ representation³ acts on **contravariant 4-vectors** v^μ , which transform according to $v^\mu \mapsto \Lambda^\mu_\nu v^\nu$, and **covariant 4-vectors** v_μ , which transform according to $v_\mu \mapsto \Lambda_\mu^\nu v_\nu$. In this representation, the generators are represented as $\mathbb{C}^{4 \times 4}$ matrices:

$$[J^{\mu\nu}]^\rho_\sigma = i(\eta^{\mu\rho}\delta^\nu_\sigma - \eta^{\nu\rho}\delta^\mu_\sigma) \quad (1.30)$$

This is an irreducible representation, and the associated Lie algebra $\mathfrak{so}^+(1, 3)$, called the **Lorentz algebra**, is:

$$[J^{\mu\nu}, J^{\sigma\rho}] = i(\eta^{\nu\rho}J^{\mu\sigma} - \eta^{\mu\rho}J^{\nu\sigma} - \eta^{\nu\sigma}J^{\mu\rho} + \eta^{\mu\sigma}J^{\nu\rho}) \quad (1.31)$$

It is convenient to rearrange the 6 independent components of $J^{\mu\nu}$ into two spatial vectors:

$$J^i := \frac{1}{2}\epsilon^{ijk}J^{jk} \quad K^i := J^{i0} \quad (1.32)$$

The $\mathfrak{so}^+(1, 3)$ algebra can then be rewritten as:

$$[J^i, J^j] = i\epsilon^{ijk}J^k \quad [J^i, K^j] = i\epsilon^{ijk}K^k \quad [K^i, K^j] = -i\epsilon^{ijk}J^k \quad (1.33)$$

The first equation defines a $\mathfrak{su}(2)$ sub-algebra, thus showing that J^i are the generators of angular momentum. Angles and rapidities are then defined as:

$$\theta^i := \frac{1}{2}\epsilon^{ijk}\omega^{jk} \quad \eta^i := \omega^{i0} \quad (1.34)$$

so that:

$$\Lambda = \exp[-i\boldsymbol{\theta} \cdot \mathbf{J} + i\boldsymbol{\eta} \cdot \mathbf{K}] \quad (1.35)$$

where $\boldsymbol{\theta}, \boldsymbol{\eta} \in \mathbb{R}^3$ and $\mathbf{J}, \mathbf{K} \in (\mathbb{R}^{4 \times 4})^3$. This definition reflects the *alias* interpretation: angles define counterclockwise rotations of vectors with respect to a fixed reference frame, while rapidities define boosts which increase velocities with respect to said frame.

§1.2.1.2 Tensor Representations

Given the action of $\text{SO}^+(1, 3)$ on covariant and contravariant vectors, it is clear to see that a generic (p, q) -tensor transforms as:

$$T^{\mu_1 \dots \mu_p}_{\nu_1 \dots \nu_q} \mapsto \Lambda^{\mu_1}_{\alpha_1} \dots \Lambda^{\mu_p}_{\alpha_p} \Lambda_{\nu_1}^{\beta_1} \dots \Lambda_{\nu_q}^{\beta_q} T^{\alpha_1 \dots \alpha_p}_{\beta_1 \dots \beta_q} \quad (1.36)$$

This shows that the representation of the Lorentz group which acts on (p, q) -tensors, although being of degree $n = 4^{p+q}$, is reducible into the direct product of $p + q$ 4-dimensional representations.

³To be precise, contravariant 4-vectors are elements of $V \equiv \mathbb{R}^{1,3}$, so they transform according to $v_\mu \mapsto \tilde{v}_\mu \equiv [\Lambda]^\mu_\nu v^\nu$, while covariant 4-vectors are elements of V^* , hence they require the transposed endomorphism (see Def. 1.4.2 in [7]) and inversely transform according to $\tilde{v}_\mu \mapsto v_\mu \equiv [\Lambda^\top]_\mu^\nu \tilde{v}_\nu := \tilde{v}_\nu [\Lambda]^\nu_\mu$. To invert this relation note that, from Eq. 1.28, $[\Lambda^{-1}]^\mu_\nu = [\Lambda]_\nu^\mu$, thus the transformation law is $v_\mu \mapsto \tilde{v}_\mu \equiv [\Lambda]_\mu^\nu v_\nu$. The identification $[\Lambda]^\mu_\nu \equiv \Lambda^\mu_\nu$ is an abuse of notation.

Moreover, consider the action of the Lorentz group on $(2, 0)$ -tensors: being $T^{\mu\nu} \mapsto \Lambda^\mu_\alpha \Lambda^\nu_\beta T^{\alpha\beta}$, if $T^{\mu\nu}$ is (anti-)symmetric it will remain so under a Lorentz transformation. Therefore, the 16-dimensional representation reduces to a 6-dimensional representation on anti-symmetric tensors and a 10-dimensional representation of symmetric tensors. Furthermore, the trace of a symmetric tensor is invariant, as $T \equiv \eta_{\mu\nu} T^{\mu\nu} \mapsto \eta_{\mu\nu} \Lambda^\mu_\alpha \Lambda^\nu_\beta T^{\alpha\beta} = T$, so the latter representation further reduces into a 9-dimensional representation on symmetric traceless tensors and a 1-dimensional representation on scalars. This means that:

$$\mathbf{4} \otimes \mathbf{4} = \mathbf{1} \oplus \mathbf{6} \oplus \mathbf{9} \quad (1.37)$$

These are irreducible representations which act on S , $A^{\mu\nu}$ and $\bar{S}^{\mu\nu} \equiv S^{\mu\nu} - \frac{1}{4}\eta^{\mu\nu}S$ respectively, with $A^{\mu\nu} \equiv \frac{1}{2}(T^{\mu\nu} - T^{\nu\mu})$ and $S^{\mu\nu} \equiv \frac{1}{2}(T^{\mu\nu} + T^{\nu\mu})$.

Decomposition under rotations Restricting the action of $\text{SO}^+(1, 3)$ to its subgroup $\text{SO}(3)$, tensors can be decomposed according to irreducible representations of $\text{SO}(3)$, which are labelled by the angular momentum $j \in \mathbb{N}_0$ and are of degree $n = 2j + 1$. Also recall the Clebsch–Gordan composition of angular momenta:

$$\mathbf{j}_1 \otimes \mathbf{j}_2 = \bigoplus_{j=|j_1-j_2|}^{j_1+j_2} \mathbf{j} \quad (1.38)$$

A Lorentz scalar α is a scalar under rotations too, so $\alpha \in \mathbf{0}$ (abuse of notation to denote that α transforms according to the $\mathbf{0}$ representation). A 4-vector v^μ is irreducible under the action of $\text{SO}^+(1, 3)$, but under $\text{SO}(3)$ it is decomposed into v^0 and \mathbf{v} , so $v^\mu \in \mathbf{0} \oplus \mathbf{1}$. A $(2, 0)$ -tensor then is:

$$\begin{aligned} T^{\mu\nu} \in (\mathbf{0} \oplus \mathbf{1}) \otimes (\mathbf{0} \oplus \mathbf{1}) &= (\mathbf{0} \otimes \mathbf{0}) \oplus (\mathbf{0} \otimes \mathbf{1}) \oplus (\mathbf{1} \otimes \mathbf{0}) \oplus (\mathbf{1} \otimes \mathbf{1}) \\ &= \mathbf{0} \oplus \mathbf{1} \oplus \mathbf{1} \oplus (\mathbf{0} \oplus \mathbf{1} \oplus \mathbf{2}) \end{aligned}$$

This is equivalent to Eq. 1.37: the trace is a scalar, so $S \in \mathbf{0}$, while the anti-symmetric part can be written as two spatial vectors A^{0i} and $\frac{1}{2}\epsilon^{ijk}A^{jk}$, so $A^{\mu\nu} \in \mathbf{1} \oplus \mathbf{1}$. The traceless symmetric part then decomposes as $\bar{S}^{\mu\nu} \in \mathbf{0} \oplus \mathbf{1} \oplus \mathbf{2}$ under spatial rotations.

Equivalently, $T^{\mu\nu}$ can be decomposed into $T^{00} \in (\mathbf{0} \otimes \mathbf{0})$, $T^{0i} \in (\mathbf{0} \otimes \mathbf{1})$, $T^{i0} \in (\mathbf{1} \otimes \mathbf{0})$ and $T^{ij} \in (\mathbf{1} \otimes \mathbf{1})$: the formers are a scalar and two spatial vectors associated to $\mathbf{0} \oplus \mathbf{1} \oplus \mathbf{1}$, while the latter can be decomposed into the trace, which is $\mathbf{0}$, the anti-symmetric part, which is $\mathbf{1}$ ($\epsilon^{ijk}A^{jk}$), and the traceless symmetric part, which is $\mathbf{2}$.

Example 1.2.1 (Gravitational waves)

Gravitational waves in **de Donder gauge** are described by a traceless symmetric matrix, therefore they have $j = 2$ (spin of the graviton).

There are two **invariant tensors** under $\text{SO}^+(1, 3)$: the metric $\eta_{\mu\nu}$, by Eq. 1.28, and the Levi–Civita symbol $\epsilon^{\mu\nu\sigma\rho}$:

$$\epsilon^{\mu\nu\sigma\rho} \mapsto \Lambda^\mu_\alpha \Lambda^\nu_\beta \Lambda^\sigma_\gamma \Lambda^\rho_\delta \epsilon^{\alpha\beta\gamma\delta} = (\det \Lambda) \epsilon^{\mu\nu\sigma\rho} = \epsilon^{\mu\nu\sigma\rho}$$

Note that, on the contrary, discrete transformations invert the handedness of spacetime, hence the Levi–Civita symbol changes sign: this hints to the facts that $\epsilon^{\mu\nu\rho\sigma}$ is a *tensor density*, since its transformation law is proportional to a power of the determinant of the coordinate transformation (this is true for general differentiable manifolds).

§1.2.1.3 Spinorial representations

The Lie algebras $\mathfrak{su}(2)$ and $\mathfrak{so}(3)$ both have three generators defined by:

$$[X_i, X_j] = i\epsilon^{ijk}X_k \quad (1.39)$$

which means that $SU(2)$ and $SO(3)$ are indistinguishable by infinitesimal transformations; however, they are globally different, as $SO(3)$ rotations are periodic by 2π , while $SU(2)$ rotations are periodic by 4π : in particular, it can be shown that $SO(3) \cong SU(2)/\mathbb{Z}_2$, i.e. $SU(2)$ is the universal cover of $SO(3)$. This means that $SU(2)$ representations can be labelled by $j \in \frac{1}{2}\mathbb{N}_0$, where half-integer spin representations are known as **spinorial representations**: they act on spinors, i.e. objects which change sign under rotations of 2π (thus not suitable to represent $SO(3)$).

Example 1.2.2 (Composition of two spin- $\frac{1}{2}$)

The $\frac{1}{2}$ representation of $SU(2)$ is a 2-dimensional representation where $J^i = \frac{\sigma^i}{2}$: the Pauli matrices satisfy $\sigma^i \sigma^j = \delta^{ij} + i\epsilon^{ijk}\sigma^k$, thus the $\mathfrak{su}(2)$ algebra is satisfied. Denoting the $m = \pm\frac{1}{2}$ states in the $\frac{1}{2}$ representation as $|\uparrow\rangle$ and $|\downarrow\rangle$, the Clebsch–Gordan decomposition $\frac{1}{2} \otimes \frac{1}{2} = \mathbf{0} \oplus \mathbf{1}$ yields the triplet ($j = 1$) $|\uparrow\uparrow\rangle, \frac{1}{\sqrt{2}}(|\uparrow\downarrow\rangle + |\downarrow\uparrow\rangle), |\downarrow\downarrow\rangle$ and the singlet ($j = 0$) $\frac{1}{\sqrt{2}}(|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle)$.

Proposition 1.2.2 (Lorentz algebra)

$$\mathfrak{so}^+(1, 3) \cong \mathfrak{su}(2) \oplus \mathfrak{su}(2) \quad (1.40)$$

Proof. Given the $\mathfrak{so}^+(1, 3)$ algebra in Eq. 1.32, it is possible to define:

$$\mathbf{J}_\pm := \frac{1}{2}(\mathbf{J} \pm i\mathbf{K}) \quad \implies \quad [\mathbf{J}_\pm^i, \mathbf{J}_\pm^j] = i\epsilon^{ijk}\mathbf{J}_\pm^k \quad [\mathbf{J}_\pm^i, \mathbf{J}_\mp^j] = 0 \quad (1.41)$$

These are two commuting $\mathfrak{su}(2)$ algebras, thus proving the thesis^a. \square

^aTo be precise, $\mathfrak{so}^+(1, 3)_\mathbb{C} \cong \mathfrak{su}(2)_\mathbb{C} \oplus \mathfrak{su}(2)_\mathbb{C}$, while $\mathfrak{so}^+(1, 3) \cong \mathfrak{sl}(2, \mathbb{C})$ since complexification yields the following isomorphisms (with Eq. 1.41):

$$\mathfrak{so}^+(1, 3) \hookrightarrow \mathfrak{so}^+(1, 3)_\mathbb{C} \cong \mathfrak{su}(2)_\mathbb{C} \oplus \mathfrak{su}(2)_\mathbb{C} \cong \mathfrak{sl}(2, \mathbb{C}) \oplus \mathfrak{sl}(2, \mathbb{C}) \cong \mathfrak{sl}(2, \mathbb{C})_\mathbb{C} \hookrightarrow \mathfrak{sl}(2, \mathbb{C})$$

Indeed, note that $SO^+(1, 3)$ is not globally equivalent to $SU(2) \times SU(2)$: instead, $SU(2) \times SU(2)/\mathbb{Z}_2 \cong SO(4)$, while the universal cover of $SO^+(1, 3)$ is $SL(2, \mathbb{C})$, as $SO^+(1, 3) \cong SL(2, \mathbb{C})/\mathbb{Z}_2$.

By Prop. 1.2.2, representations of $SO^+(1, 3)$ can be labelled by $(\mathbf{j}_-, \mathbf{j}_+) \in \frac{1}{2}\mathbb{N}_0 \times \frac{1}{2}\mathbb{N}_0$, with each index labelling a representation of $SU(2)$: as $\mathbf{J} = \mathbf{J}_+ + \mathbf{J}_-$, the $(\mathbf{j}_-, \mathbf{j}_+)$ representation contains states with all possible spins $|j_+ - j_-| \leq j \leq j_+ + j_-$, and it is a representation of degree $n = (2j_- + 1)(2j_+ + 1)$.

More formally, the $(\mathbf{j}_-, \mathbf{j}_+)$ representation is $\rho_{(j_-, j_+)} : \mathfrak{so}^+(1, 3) \rightarrow \mathfrak{gl}(V_{(j_-, j_+)})$, where $V_{(j_-, j_+)}$ is a $(2j_- + 1)(2j_+ + 1)$ -dimensional \mathbb{C} -vector space, defined up to similarity as (inverting Eq. 1.41):

$$\begin{aligned} \rho_{(j_-, j_+)}(J^i) &= J_{(j_+)}^i \otimes I_{(2j_-+1)} + I_{(2j_++1)} \otimes J_{(j_-)}^i \\ \rho_{(j_-, j_+)}(K^i) &= -i [J_{(j_+)}^i \otimes I_{(2j_++1)} - I_{(2j_++1)} \otimes J_{(j_-)}^i] \end{aligned} \quad (1.42)$$

where $\mathbf{J}_{(j)}$ is the $(2j + 1)$ -dimensional irreducible representation of $\mathfrak{su}(2)$ and \otimes is the tensor product (see Def. A.1.2). For example, the trivial representation is $(\mathbf{0}, \mathbf{0})$, as both $\mathbf{J}_\pm = \mathbf{0}$ and $\mathbf{J} = \mathbf{K} = \mathbf{0}$: this is the representation which acts on scalars.

Weyl representations $(\frac{1}{2}, 0)$ and $(0, \frac{1}{2})$ are 2-dimensional spinorial representations. These representations act on different spinors $(\psi_L)_\alpha, (\psi_R)_\beta \in \mathbb{C}^2$, with $\alpha, \beta = 1, 2$, which are respectively called **left-** and **right-handed Weyl spinors**.

In $(\frac{1}{2}, 0)$ the generators are $\mathbf{J}_- = \frac{\sigma}{2}$ and $\mathbf{J}_+ = 0$, while in $(0, \frac{1}{2})$ they are $\mathbf{J}_- = 0$ and $\mathbf{J}_+ = \frac{\sigma}{2}$, thus, by Eq. 1.42, $\mathbf{J}_L = \mathbf{J}_R = \frac{\sigma}{2}$ and $\mathbf{K}_L = -\mathbf{K}_R = i\frac{\sigma}{2}$, so that from Eq. 1.35:

$$\psi_L \mapsto \Lambda_L \psi_L = \exp \left[(-i\boldsymbol{\theta} - \boldsymbol{\eta}) \cdot \frac{\boldsymbol{\sigma}}{2} \right] \psi_L \quad (1.43)$$

$$\psi_R \mapsto \Lambda_R \psi_R = \exp \left[(-i\boldsymbol{\theta} + \boldsymbol{\eta}) \cdot \frac{\boldsymbol{\sigma}}{2} \right] \psi_R \quad (1.44)$$

Note that the generators K^i are not hermitian, as expected from Prop. A.2.3. Furthermore, as $\psi_{L/R} \in \mathbb{C}^2$, then $\Lambda_{L/R} \in \mathbb{C}^{2 \times 2}$ (in general the spinorial representations are complex representations).

Proposition 1.2.3

Given $\psi_L \in (\frac{1}{2}, 0)$ and $\psi_R \in (0, \frac{1}{2})$, then $\sigma^2 \psi_L^* \in (0, \frac{1}{2})$ and $\sigma^2 \psi_R^* \in (\frac{1}{2}, 0)$.

Proof. Recall that for Pauli matrices $\sigma^2 \sigma^i \sigma^2 = -(\sigma^i)^*$, so $\sigma^2 \Lambda_L^* \sigma^2 = \Lambda_R$ and:

$$\sigma^2 \psi_L^* \mapsto \sigma^2 (\Lambda_L \psi_L)^* = (\sigma^2 \Lambda_L^* \sigma^2) \sigma^2 \psi_L^* = \Lambda_R \sigma^2 \psi_L^* \implies \sigma^2 \psi_L^* \in (0, \frac{1}{2})$$

where $\sigma^2 \sigma^2 = I_2$ was used. The proof for $\sigma^2 \psi_R^*$ is analogous. \square

This allows to define **charge conjugation** on Weyl spinors:

$$\psi_L^c := i\sigma^2 \psi_L^* \quad \psi_R^c := -i\sigma^2 \psi_R^* \quad (1.45)$$

As of Prop. 1.2.3, charge conjugation transforms a left-handed Weyl spinor into a right-handed one and vice versa. Moreover, the i factor ensures that applying this operator twice yields the identity operator.

Dirac representation $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$ is a 4-dimensional complex representation: this representation acts on bispinors of the form $\Psi = ((\psi_L)_\alpha, (\xi_R)_\beta)^T \in \mathbb{C}^4$, called **Dirac spinors**, with transformation $\Lambda_D = \text{diag}(\Lambda_L, \Lambda_R) \in \mathbb{C}^{4 \times 4}$.

Definition 1.2.2 (Charge conjugation on Dirac spinors)

On Dirac spinors, the **charge conjugation operator** is defined as:

$$\Psi^c := \begin{pmatrix} -i\sigma^2 \xi_R^* \\ i\sigma^2 \psi_L^* \end{pmatrix} = -i \begin{bmatrix} 0 & \sigma^2 \\ -\sigma^2 & 0 \end{bmatrix} \Psi^* \quad (1.46)$$

This definition reflect the fact that charge conjugation changes the representation of the Weyl spinors (Prop. 1.2.3), i.e. the chirality of the components of the Dirac spinor.

Observation 1.2.1 (4-vector representation)

The other 4-dimensional representation of the Lorentz group is $(\frac{1}{2}, \frac{1}{2})$. To see that this is, in fact, the 4-vector representation, first consider the double-covering $\text{SO}^+(1, 3) \cong \text{SL}(2, \mathbb{C})/\mathbb{Z}_2$.

Lemma 1.2.1 (Double-covering of Lorentz group)

$$\mathrm{SO}^+(1, 3) \cong \mathrm{SL}(2, \mathbb{C})/\mathbb{Z}_2 \quad (1.47)$$

Proof. Define the set $\mathcal{H}(2) := \{H \in \mathbb{C}^{2 \times 2} : H^\dagger = H\}$ of all $\mathbb{C}^{2 \times 2}$ Hermitian matrices. Trivially $\mathcal{H}(2) \cong \mathbb{R}^{1,3}$, since:

$$\mathbb{R}^{1,3} \ni v^\mu = (v^0, v^1, v^2, v^3) \longleftrightarrow \begin{pmatrix} v^0 - v^3 & -v^1 + iv^2 \\ -v^1 - iv^2 & v^0 + v^3 \end{pmatrix} = V \in \mathcal{H}(2)$$

where the isomorphisms are $V = v_\mu \sigma^\mu$, $v^\mu = \frac{1}{2} \mathrm{tr}(\bar{\sigma}^\mu V)$, with $\sigma^\mu \equiv (I_2, \boldsymbol{\sigma})$, $\bar{\sigma}^\mu \equiv (I_2, -\boldsymbol{\sigma})$, since $\{I_2, \sigma^1, \sigma^2, \sigma^3\}$ is a basis of $\mathcal{H}(2)$.

Now, consider $M \in \mathrm{SL}(2, \mathbb{C})$ and define its linear action on $\mathcal{H}(2)$ as $V \mapsto MVM^\dagger$: then, $\mathrm{SL}(2, \mathbb{C})$ is an endomorphism group of $\mathcal{H}(2)$, since its action preserves Hermiticity, and $\det V \mapsto \det V$ as $\det M = 1$ by definition. But $\det V = v_\mu v^\mu$, hence the action of $\mathrm{SL}(2, \mathbb{C})$ leaves the Minkowski norm invariant: it is thus possible to define a group homomorphism $\Phi : \mathrm{SL}(2, \mathbb{C}) \rightarrow \mathrm{O}(1, 3)$. WTS $\Phi(\mathrm{SL}(2, \mathbb{C})) \subseteq \mathrm{SO}^+(1, 3)$: since $\mathrm{SL}(2, \mathbb{C})$ is simply connected (it is homeomorphic to $\mathbb{S}^3 \times \mathbb{R}^3$) and $I_2 \in \mathrm{SL}(2, \mathbb{C})$, then $\Phi(\mathrm{SL}(2, \mathbb{C}))$ must be connected, as Φ can be shown to be continuous, and it must contain the identity, but the only connected component of $\mathrm{O}(1, 3)$ which satisfies these conditions is $\mathrm{SO}^+(1, 3)$, hence $\Phi(\mathrm{SL}(2, \mathbb{C})) \subseteq \mathrm{SO}^+(1, 3)$. Then:

- surjectivity: to show that Φ surjective, recall that $\mathrm{SO}^+(1, 3)$ is generated by spatial (trigonometric) rotations and boosts (i.e. hyperbolic rotations). The former are clearly realized by Φ , since $\mathrm{SU}(2) < \mathrm{SL}(2, \mathbb{C})$ and $\Phi|_{\mathrm{SU}(2)}$ is the standard double-covering map for $\mathrm{SO}(3) \cong \mathrm{SU}(2)/\mathbb{Z}_2$. Then, by polar decomposition, any $M \in \mathrm{SL}(2, \mathbb{C})$ can be written as $M = UP$, with $U \in \mathrm{SU}(2)$ and $P = \exp(\mathbf{a} \cdot \boldsymbol{\sigma})$, $\mathbf{a} \in \mathbb{R}^3$: since P represents^a a boost along the direction $\hat{\mathbf{a}}$ of rapidity $2a$, all generators of $\mathrm{SO}^+(1, 3)$ are represented by Φ , i.e. $\mathrm{Im}\{\Phi\} = \mathrm{SO}^+(1, 3)$.
- $\ker \Phi = \{I_2, -I_2\}$: solving $\Phi(M) = \mathbf{1}_{\mathrm{SO}^+(1,3)}$ means that $MVM^\dagger = V \forall V \in \mathcal{H}(2)$. $V = I_2$ results in $M \in \mathrm{SU}(2) < \mathrm{SL}(2, \mathbb{C})$, and $V = \sigma_i$ that M commutes with all the Pauli matrices: this means that M commutes with all elements of $\mathcal{H}(2)$, hence it is a multiple of I_2 by Schur's Lemma (applicable since the Pauli matrices form the basis of an irreducible representation of $\mathrm{SU}(2)$), and $\det M = 1 \implies M = \pm I_2$.

Since $\{I_2, -I_2\} \cong \mathbb{Z}_2$, by the first isomorphism theorem $\mathrm{SO}^+(1, 3) \cong \mathrm{SL}(2, \mathbb{C})/\mathbb{Z}_2$. \square

^aFirst, note that, given a versor $\hat{\mathbf{n}} \in \mathbb{R}^3$, it is always possible to find a rotation $U_n \in \mathrm{SU}(2)$ such that $\hat{\mathbf{n}} \cdot \boldsymbol{\sigma} = U_n \sigma^3 U_n^\dagger$, hence $P = U_a \exp(a \sigma^3) U_a^\dagger$, thus WTS $\exp(a \sigma^3) = \mathrm{diag}(e^a, e^{-a})$ is a boost along the z -axis:

$$\begin{bmatrix} e^a & 0 \\ 0 & e^{-a} \end{bmatrix} \begin{pmatrix} v^0 - v^3 & -v^1 + iv^2 \\ -v^1 - iv^2 & v^0 + v^3 \end{pmatrix} \begin{bmatrix} e^a & 0 \\ 0 & e^{-a} \end{bmatrix} = \begin{pmatrix} e^{2a}(v^0 - v^3) & -v^1 + iv^2 \\ -v^1 - iv^2 & e^{-2a}(v^0 + v^3) \end{pmatrix}$$

i.e. $v^0 \pm v^3 \mapsto e^{\mp 2a}(v^0 \pm v^3)$ and $(v^1, v^2) \mapsto (v^1, v^2)$, which is equivalent to:

$$\begin{pmatrix} v^0 \\ v^3 \end{pmatrix} \mapsto \begin{bmatrix} \cosh \eta & -\sinh \eta \\ -\sinh \eta & \cosh \eta \end{bmatrix} \begin{pmatrix} v^0 \\ v^3 \end{pmatrix}$$

which is exactly the Lorentz transformation realizing a boost along the z -axis.

As of Lemma 1.2.1, given $\Lambda \in \text{SO}^+(1, 3)$, it is possible to associate to it $\Lambda_s \in \text{SL}(2, \mathbb{C})$ such that:

$$\Lambda^\mu{}_\nu \sigma^\nu = \Lambda_s \sigma^\mu \Lambda_s^\dagger \quad (1.48)$$

Recalling the proof of Prop. 1.2.2 $\mathfrak{sl}(2, \mathbb{C}) \cong \mathfrak{su}(2)_\mathbb{C} \oplus \mathfrak{su}(2)_\mathbb{C}$, hence:

$$\Lambda_s = \exp \left[(\mathbf{a} + i\mathbf{b}) \cdot \frac{\boldsymbol{\sigma}}{2} \right] \implies \Lambda_s^\dagger = \exp \left[(-\mathbf{a} + i\mathbf{b}) \cdot \frac{\boldsymbol{\sigma}}{2} \right]$$

where $\sigma^2 \boldsymbol{\sigma} \sigma^2 = -\boldsymbol{\sigma}^\top$ was used. Since $\exp(i\alpha \sigma^k) = I_2 \cos \alpha - i\sigma^k \sin \alpha$, the exponential is periodic of 2π in \mathbf{b} , so identify $\mathbf{a} \equiv -\boldsymbol{\eta}$ and $\mathbf{b} \equiv -\boldsymbol{\theta}$, i.e. $\Lambda_s \equiv \Lambda_L$ and $\Lambda_s^\dagger = \sigma^2 \Lambda_R^\top \sigma^2$: it is then clear that the representation of $\text{SO}^+(1, 3)$ acting on 4-vectors is $(\frac{1}{2}, \mathbf{0}) \otimes (\mathbf{0}, \frac{1}{2})$, with^a:

$$(\psi \otimes \xi) \equiv \psi \xi^\top \sigma^2 \implies \Lambda(\psi \otimes \xi) = \Lambda_s(\psi \otimes \xi) \Lambda_s^\dagger = (\Lambda_L \psi) \otimes (\Lambda_R \xi)$$

Note that the restriction of $(\psi \otimes \xi)$ being Hermitian correctly reduces the real dimension of the basis space from 8 to 4 (i.e. the basis space from $\mathbb{C}^2 \otimes \mathbb{C}^2$ to $\mathcal{H}(2)$). Now, given the canonical isomorphism $V \otimes W \cong W \otimes V$ for generic \mathbb{K} -vector spaces, the following equivalence relation holds⁴:

$$(\tfrac{1}{2}, \mathbf{0}) \otimes (\mathbf{0}, \tfrac{1}{2}) \cong (\tfrac{1}{2} \otimes \mathbf{0}, \mathbf{0} \otimes \tfrac{1}{2}) \cong (\tfrac{1}{2}, \tfrac{1}{2})$$

where Clebsch–Gordan decomposition was used. This shows that $(\frac{1}{2}, \frac{1}{2})$ is the representation of $\text{SO}^+(1, 3)$ acting on 4-vectors.

Moreover, by Clebsch–Gordan decomposition, $(\frac{1}{2}, \frac{1}{2}) \cong \mathbf{0} \oplus \mathbf{1}$, which is exactly the representation acting on 4-vectors ($v^0 \in \mathbf{0}$ and $\mathbf{v} \in \mathbf{1}$).

^aTo be precise, this is a representation equivalent to $(\frac{1}{2}, \mathbf{0}) \otimes (\mathbf{0}, \frac{1}{2})$, which would act on $[\psi \otimes \xi]_{(ij)} = \psi_i \xi_j$.

Parity Note that $\mathcal{P}\mathbf{K} = -\mathbf{K}$, as the velocity of the boost gets reversed, while $\mathcal{P}\mathbf{J} = \mathbf{J}$: this means that $\mathcal{P}\mathbf{J}_\pm = \mathbf{J}_\mp$, i.e. parity exchanges a $(\mathbf{j}_-, \mathbf{j}_+)$ representation into a $(\mathbf{j}_+, \mathbf{j}_-)$ representation. Therefore, a $(\mathbf{j}_-, \mathbf{j}_+)$ representation of $\text{SO}^+(1, 3)$ is a basis for the representation of the parity transformation if and only if $j_- = j_+$.

Example 1.2.3 (Parity on spinors)

While Weyl spinors (separately) are not a basis for a representation of the parity operator, Dirac spinors are.

§1.2.1.4 Dirac algebra

The Dirac algebra is $\mathfrak{cl}_{1,3}(\mathbb{C}) \cong \mathfrak{cl}_{1,3}(\mathbb{R}) \otimes \mathbb{C}$ (complexification). This Clifford algebra admits a matrix representation via $\gamma^\mu \in \mathbb{C}^{4 \times 4} \forall \mu = 0, 1, 2, 3$ such that:

$$\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu} I_4 \quad (1.49)$$

⁴In particular, a consequence is that, given two Lie algebras \mathfrak{g} and \mathfrak{h} , and given $\rho_1 : \mathfrak{g} \rightarrow V_1$, $\rho_2 : \mathfrak{g} \rightarrow V_2$, $\pi_1 : \mathfrak{h} \rightarrow W_1$ and $\pi_2 : \mathfrak{h} \rightarrow W_2$ their representations on \mathbb{C} -vector spaces, then $(\rho_1 \otimes \pi_1) \otimes (\rho_2 \otimes \pi_2) : (\mathfrak{g} \times \mathfrak{h}) \times (\mathfrak{g} \times \mathfrak{h}) \rightarrow \mathfrak{gl}((V_1 \otimes W_1) \otimes (V_2 \otimes W_2))$ and $(\rho_1 \otimes \rho_2) \otimes (\pi_1 \otimes \pi_2) : (\mathfrak{g} \times \mathfrak{g}) \times (\mathfrak{h} \times \mathfrak{h}) \rightarrow \mathfrak{gl}((V_1 \otimes V_2) \otimes (W_1 \otimes W_2))$ are equivalent.

The basis of this algebra is then given by:

I_4	1 matrix
γ^μ	4 matrices
$\gamma^{\mu\nu} \equiv \frac{1}{2}[\gamma^\mu, \gamma^\nu] \equiv \gamma^{[\mu}\gamma^{\nu]}$	6 matrices
$\gamma^{\mu\nu\rho} \equiv \gamma^{[\mu}\gamma^\nu\gamma^{\rho]} = i\epsilon^{\mu\nu\rho\sigma}\gamma_\sigma\gamma^5$	4 matrices
$\gamma^{\mu\nu\rho\sigma} \equiv \gamma^{[\mu}\gamma^\nu\gamma^\rho\gamma^{\sigma]} = -i\epsilon^{\mu\nu\rho\sigma}\gamma^5$	1 matrix

where $\gamma^5 := i\gamma^0\gamma^1\gamma^2\gamma^3$ is an additional gamma matrix. It is also convention to use $\sigma^{\mu\nu} \equiv i\gamma^{\mu\nu}$. Note that this matrix is not in the basis of $\mathfrak{cl}_{1,3}(\mathbb{C})$. It has the following properties:

$$(\gamma^5)^\dagger = \gamma^5 \quad (\gamma^5)^2 = I_4 \quad \{\gamma^5, \gamma^\mu\} = 0$$

Therefore, the Dirac algebra is a 16-dimensional unital associative \mathbb{C} -algebra, and its generic component is:

$$\Gamma = aI_4 + a_\mu\gamma^\mu + \frac{1}{2}a_{\mu\nu}\gamma^{\mu\nu} + \frac{1}{3!}a_{\mu\nu\rho}\gamma^{\mu\nu\rho} + \frac{1}{4!}a_{\mu\nu\rho\sigma}\gamma^{\mu\nu\rho\sigma} \quad (1.50)$$

where each coefficient is totally anti-symmetric.

The quadratic subspace $\mathfrak{cl}_{1,3}^{(2)}(\mathbb{C})$ of the Dirac algebra is of particular interest. By convention, define its 6 generators as $S^{\mu\nu} := \frac{i}{4}[\gamma^\mu, \gamma^\nu] \equiv \frac{1}{2}\sigma^{\mu\nu}$: it is straightforward to show that these generators satisfy Eq. 1.31, thus they give a 4-dimensional representation of $\mathfrak{so}^+(1, 3)$.

Proposition 1.2.4 (Dirac algebra and Dirac representation)

The generators of $\mathfrak{cl}_{1,3}^{(2)}(\mathbb{C})$ also give the $(\frac{1}{2}, \mathbf{0}) \oplus (\mathbf{0}, \frac{1}{2})$ representation of $\mathfrak{so}^+(1, 3)$.

Moreover, $\mathfrak{cl}_{1,3}^{(2)}(\mathbb{C}) \cong \text{Spin}(1, 3)$ (see §A.3.1.1). Given a spinor $\Psi \in \mathbb{C}^4$, this spin group acts as:

$$\Psi \mapsto \Lambda_{\frac{1}{2}}\Psi = \exp\left[-\frac{i}{2}\omega_{\mu\nu}S^{\mu\nu}\right]\Psi$$

Using the property $(\gamma^\mu)^\dagger = \gamma^0\gamma^\mu\gamma^0$, it is easy to see that $(S^{\mu\nu})^\dagger = -\gamma^0S^{\mu\nu}\gamma^0$, so, defining the **Dirac dual** $\bar{\Psi} := \Psi^\dagger\gamma^0$, this transforms via the inverse transformation:

$$\bar{\Psi} \mapsto \bar{\Psi}\Lambda_{\frac{1}{2}}^{-1}$$

Dirac bilinears The fact that Ψ and $\bar{\Psi}$ transform inversely allows to define objects with particular transformation relations.

Theorem 1.2.1 (Lorentz index of gamma matrices)

The gamma matrices transform under the $(\frac{1}{2}, \frac{1}{2})$ representation of $\text{SO}^+(1, 3)$ as:

$$\Lambda^\mu{}_\nu\gamma^\nu = \Lambda_{\frac{1}{2}}^{-1}\gamma^\mu\Lambda_{\frac{1}{2}} \quad (1.51)$$

Proof. Recalling the explicit form of the generators in Eq. 1.30:

$$[J^{\rho\sigma}]^\mu{}_\nu\gamma^\nu = i(\eta^{\rho\mu}\gamma^\sigma - \eta^{\sigma\mu}\gamma^\rho) = \frac{i}{2}(\gamma^\sigma\gamma^\rho\gamma^\mu + \gamma^\sigma\gamma^\mu\gamma^\rho - \gamma^\rho\gamma^\sigma\gamma^\mu - \gamma^\rho\gamma^\mu\gamma^\sigma)$$

As $[\gamma^\mu, \gamma^\nu] = \{\gamma^\mu, \gamma^\nu\} - 2\gamma^\nu\gamma^\mu = 2(\eta^{\mu\nu} - \gamma^\nu\gamma^\mu)$:

$$\begin{aligned} [\gamma^\mu, S^{\rho\sigma}] &= \frac{i}{4}[\gamma^\mu, \gamma^\rho\gamma^\sigma - \gamma^\sigma\gamma^\rho] = \frac{i}{4}(\gamma^\rho[\gamma^\mu, \gamma^\sigma] + [\gamma^\mu, \gamma^\rho]\gamma^\sigma - [\gamma^\mu, \gamma^\sigma]\gamma^\rho - \gamma^\sigma[\gamma^\mu, \gamma^\rho]) \\ &= \frac{i}{2}(-\gamma^\rho\gamma^\sigma\gamma^\mu - \gamma^\rho\gamma^\mu\gamma^\sigma + \gamma^\sigma\gamma^\mu\gamma^\rho + \gamma^\sigma\gamma^\rho\gamma^\mu) = [J^{\rho\sigma}]^\mu{}_\nu\gamma^\nu \end{aligned}$$

Expanding the thesis at first order:

$$\left[\mathbb{I}_4 - \frac{i}{2}\omega_{\rho\sigma}J^{\rho\sigma} \right]^\mu{}_\nu \gamma^\nu = \left[\mathbb{I}_4 + \frac{i}{2}\omega_{\rho\sigma}S^{\rho\sigma} \right] \gamma^\mu \left[\mathbb{I}_4 - \frac{i}{2}\omega_{\rho\sigma}S^{\rho\sigma} \right]$$

which becomes:

$$-[J^{\rho\sigma}]^\mu{}_\nu\gamma^\nu = S^{\rho\sigma}\gamma^\mu - \gamma^\mu S^{\rho\sigma}$$

This equation has been verified, so the thesis holds. \square

This result means that, although they are matrices, the μ in γ^μ is effectively a Lorentz index, thus allowing to construct dot products with 4-vectors; in this way, gamma matrices transform 4-vectors into operators on spinors⁵: $A^\mu \in (\frac{1}{2}, \frac{1}{2})$, but $\not{A} \equiv \gamma^\mu A_\mu \in (\frac{1}{2}, \mathbf{0}) \oplus (\mathbf{0}, \frac{1}{2})$.

A **Dirac bilinear** is an object $\bar{\Psi}\Gamma\Psi$, with $\Gamma \in \mathfrak{cl}_{1,3}(\mathbb{C})$. There are 5 basic bilinears:

$$\bar{\Psi}\Psi \quad \bar{\Psi}\gamma^\mu\Psi \quad \bar{\Psi}\gamma^5\Psi \quad \bar{\Psi}\gamma^5\gamma^\mu\Psi \quad \bar{\Psi}\sigma^{\mu\nu}\Psi$$

According to their transformation relations under the Lorentz group, these bilinears are respectively called scalar, vector, pseudo-scalar, pseudo-vector and tensor bilinear.

§1.2.1.5 Dual and adjoint representations

Recall that the adjoint representation of a Lie group is a representation acting on the associated Lie algebra: therefore, the adjoint representation of $\mathrm{SO}^+(1,3)$ is 6-dimensional and, by Prop. 1.2.2, it can be decomposed as $\mathbf{6} \cong (\mathbf{1}, \mathbf{0}) \oplus (\mathbf{0}, \mathbf{1})$.

Dual representations Consider a generic anti-symmetric tensor $A^{\mu\nu}$: its **dual tensor** is defined as $A_\star^{\mu\nu} := \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma}$. Every anti-symmetric tensor can then be decomposed as:

$$A_\pm^{\mu\nu} \equiv \frac{1}{2}(A^{\mu\nu} \pm iA_\star^{\mu\nu}) \quad (1.52)$$

which are respectively called the **self-dual** and **anti-self-dual part** of $A^{\mu\nu}$.

Lemma 1.2.2 ((Anti)-self-dual tensor)

An (anti)-self-dual tensor satisfies:

$$A^{\mu\nu} = \pm \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma} \quad (1.53)$$

⁵More precisely: spinors are defined as left-modules (even more precisely, as minimal left-ideals: see Def. A.3.3 and Def. A.3.4) over $\mathfrak{cl}_{1,3}(\mathbb{C})$, and $\not{A} = \gamma^\mu A_\mu \in \mathfrak{cl}_{1,3}(\mathbb{C})$ (specifically $\mathfrak{cl}_{1,3}^{(1)}(\mathbb{C})$, with notation from Eq. A.12).

Proof. Consider $A^{\mu\nu} = \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma}$:

$$A_+^{\mu\nu} = \frac{1}{2} \left(A^{\mu\nu} + \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma} \right) = \frac{1}{2} (i\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma}) = A^{\mu\nu} \quad A_-^{\mu\nu} = \frac{1}{2} \left(A^{\mu\nu} - \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma} \right) = 0$$

The other case is trivially verified. \square

Anti-symmetric tensors belong to the **6** representation, but this is a reducible representation. Indeed, the subspaces of self-dual and anti-self-dual tensors are invariant subspaces:

$$A^{\mu\nu} = \pm \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}A_{\rho\sigma} \implies \pm \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}(\Lambda_\rho^\alpha \Lambda_\sigma^\beta A_{\alpha\beta}) = -\frac{1}{4} \underbrace{\epsilon^{\mu\nu\rho\sigma}\epsilon_{\alpha\beta\gamma\delta}}_{\delta_{\alpha\beta\gamma\delta}^{\mu\nu\rho\sigma}} \Lambda_\rho^\alpha \Lambda_\sigma^\beta A^{\gamma\delta} = \Lambda^\mu_\alpha \Lambda^\nu_\beta A^{\alpha\beta}$$

Therefore, $\mathbf{6} = \mathbf{3} \oplus \bar{\mathbf{3}} \cong (\mathbf{1}, \mathbf{0}) \oplus (\mathbf{0}, \mathbf{1})$.

Adjoint representation To see explicitly that the adjoint representation of $\text{SO}^+(1, 3)$ is exactly **6**, recall that the adjoint representation is $\rho : \text{SO}^+(1, 3) \rightarrow \text{GL}(\mathfrak{so}^+(1, 3))$ and it acts on $\mathfrak{so}^+(1, 3)$ as $\rho(\Lambda)M = \Lambda M \Lambda^{-1}$, where $\Lambda \in \text{SO}^+(1, 3)$ and $M \in \mathfrak{so}^+(1, 3)$. As $[\Lambda^{-1}]^\mu_\nu = \Lambda_\nu^\mu$ by Eq. 1.28:

$$M^{\mu\nu} \mapsto \Lambda^\mu_\alpha M^{\alpha\beta} [\Lambda^{-1}]_\beta^\nu = \Lambda^\mu_\alpha \Lambda^\nu_\beta M^{\alpha\beta}$$

As $M \in \mathfrak{so}^+(1, 3)$ is anti-symmetric, this is the transformation relation of an anti-symmetric tensor: the adjoint representation is the $\mathbf{6} \cong (\mathbf{1}, \mathbf{0}) \oplus (\mathbf{0}, \mathbf{1})$.

Example 1.2.4 (Gauge theories)

In gauge theories (e.g. QED and QCD), the gauge potential A_μ transforms according to the fundamental representation $(\frac{1}{2}, \frac{1}{2})$, while the field strength $F_{\mu\nu}$ according to the adjoint representation $(\mathbf{1}, \mathbf{0}) \oplus (\mathbf{0}, \mathbf{1})$.

§1.2.1.6 Field representations

Given a field $\phi(x)$, under a Lorentz transformation $x^\mu \mapsto x'^\mu = \Lambda^\mu_\nu x^\nu$ it transforms as $\phi(x) \mapsto \phi'(x')$.

Scalar fields A scalar field transforms as:

$$\phi'(x') = \phi(x) \tag{1.54}$$

Consider an infinitesimal transformation $x'^\rho = x^\rho + \delta x^\rho$, with $\delta x^\rho = -\frac{i}{2}\omega_{\mu\nu} [J^{\mu\nu}]^\rho_\sigma x^\sigma$ as of Eq. 1.28. Then, by definition, $\delta\phi \equiv \phi'(x') - \phi(x) = 0$, which corresponds to the fact that the scalar representation of $\text{SO}^+(1, 3)$ is the trivial one ($J^{\mu\nu} = 0$).

However, one can consider the variation at fixed coordinate $\delta_0\phi \equiv \phi'(x) - \phi(x)$: while $\delta\phi$ studies only a single degree of freedom, as the point $P \in \mathbb{R}^{1,3}$ is kept constant and only $\phi(P)$ can vary (i.e. the base space is one-dimensional), $\delta_0\phi$ studies $\phi(P)$ with P varying over $\mathbb{R}^{1,3}$, thus the base space is now a space of functions, which is infinite-dimensional. Therefore, $\delta\phi$ provides a finite-dimensional representation of the generators, while $\delta_0\phi$ an infinite-dimensional one.

To explicit this representation:

$$\delta_0\phi = \phi'(x) - \phi(x) = \phi'(x' - \delta x) - \phi(x) = -\delta x^\rho \partial_\rho \phi = \frac{i}{2}\omega_{\mu\nu} [J^{\mu\nu}]^\rho_\sigma x^\sigma \partial_\rho \phi \equiv -\frac{i}{2}\omega_{\mu\nu} L^{\mu\nu} \phi$$

Recalling Eq. 1.30, the generators can be expressed as:

$$L^{\mu\nu} := i(x^\mu \partial^\nu - x^\nu \partial^\mu) \quad (1.55)$$

This is an infinite-dimensional representation, as it acts on the space of scalar fields. As $p^\mu = i\partial^\mu$ (with signature $(+, -, -, -)$), it is clear that $L^i \equiv \frac{1}{2}\epsilon^{ijk}L^{jk}$ is the orbital angular momentum.

Weyl fields A left-handed Weyl field transforms as:

$$\psi'_L(x') = \Lambda_L \psi_L(x) \quad (1.56)$$

with Λ_L defined in Eq. 1.43, and similarly for right-handed Weyl fields. The infinite-dimensional representation of the Lorentz generators determined by Weyl spinors can be found as:

$$\begin{aligned} \delta_0 \psi_L &\equiv \psi'_L(x) - \psi_L(x) = \psi'_L(x' - \delta x) - \psi_L(x) \\ &= \psi'_L(x') - \delta x^\rho \partial_\rho \psi_L(x) - \psi_L(x) = (\Lambda_L - I_2) \psi_L(x) - \delta x^\rho \partial_\rho \psi_L(x) \end{aligned}$$

The second term yields $L^{\mu\nu}$, while the first can be further elaborated by writing:

$$\Lambda_L = e^{-\frac{i}{2}\omega_{\mu\nu}S^{\mu\nu}} \quad (1.57)$$

Thus:

$$\delta_0 \psi_L = -\frac{i}{2}\omega_{\mu\nu}J^{\mu\nu}\psi_L$$

where the angular momentum separates into the orbital and the spin components:

$$J^{\mu\nu} = L^{\mu\nu} + S^{\mu\nu} \quad (1.58)$$

This separation is general: $L^{\mu\nu}$ is always expressed as in Eq. 1.55, while $S^{\mu\nu}$ depends on the specific representation. In the scalar representation $S^{\mu\nu} = 0$, while in the left/right-handed Weyl representation $S^{i0} = \pm i\frac{\sigma^i}{2}$.

Dirac fields A Dirac field transforms as:

$$\Psi'(x') = \Lambda_D \Psi(x) \quad (1.59)$$

where $\Lambda_D = \text{diag}(\Lambda_L, \Lambda_R)$. The infinite-dimensional representation of the Lorentz generators determined by Dirac spinors is:

$$\delta_0 \Psi \equiv \Psi'(x) - \Psi(x) = \Psi'(x') - \delta x^\rho \partial_\rho \Psi(x) - \Psi(x) = (\Lambda_D - I_4) \Psi(x) - \frac{i}{2}\omega_{\mu\nu}L^{\mu\nu}\Psi(x)$$

The first term can be rewritten defining:

$$\Lambda_D = e^{-\frac{i}{2}\omega_{\mu\nu}S^{\mu\nu}} \quad S^{0i} \equiv -\frac{i}{2} \begin{bmatrix} \sigma^i & 0 \\ 0 & -\sigma^i \end{bmatrix} \quad S^{ij} = -\frac{1}{2}\epsilon^{ijk} \begin{bmatrix} \sigma^k & 0 \\ 0 & \sigma^k \end{bmatrix} \equiv \frac{1}{2}\epsilon^{ijk}\Sigma^k$$

Therefore, a Dirac field transforms as:

$$\delta_0 \Psi = -\frac{i}{2}\omega_{\mu\nu}J^{\mu\nu}\Psi \quad (1.60)$$

with clear orbital and spin angular components:

$$J^{\mu\nu} = L^{\mu\nu} + S^{\mu\nu} \quad (1.61)$$

Vector fields A (contravariant) vector field transforms as:

$$V'^{\mu}(x') = \Lambda^{\mu}_{\nu} V^{\nu}(x) \quad (1.62)$$

A general vector field has a spin-0 and a spin-1 component, and it is acted on by the $(\frac{1}{2}, \frac{1}{2})$ representation.

§1.2.2 Poincaré group

Definition 1.2.3 (Poincaré group)

The **Poincaré group** is defined as $\text{ISO}^+(1, 3) := \text{T}^{1,3} \rtimes \text{SO}^+(1, 3)$, where $\text{T}^{1,3} \cong \mathbb{R}^{1,3}$ is the group of translations of $\mathbb{R}^{1,3}$, with multiplication $(a^{\mu}, \Lambda^{\mu}_{\nu}) \cdot (\bar{a}^{\mu}, \bar{\Lambda}^{\mu}_{\nu}) \equiv (a^{\mu} + \Lambda^{\mu}_{\nu} \bar{a}^{\nu}, \Lambda^{\mu}_{\rho} \bar{\Lambda}^{\rho}_{\nu})$.

Given a translation $x^{\mu} \mapsto x^{\mu} + a^{\mu}$, the associated group element can be written as:

$$T = e^{-i a_{\mu} P^{\mu}} \quad (1.63)$$

where P^{μ} is the 4-momentum operator. Clearly translations commute, and so do their generators; on the other hand, as \mathbf{P} is a vector under rotations, while P^0 (energy) a scalar, one has:

$$[J^i, P^j] = i \epsilon^{ijk} P^k \quad [J^i, P^0] = 0$$

These equations uniquely determine the **Poincaré algebra** $\mathfrak{iso}^+(1, 3)$:

$$\begin{aligned} [P^{\mu}, P^{\nu}] &= 0 \\ [J^{\mu\nu}, J^{\sigma\rho}] &= i(\eta^{\nu\rho} J^{\mu\sigma} - \eta^{\mu\rho} J^{\nu\sigma} - \eta^{\nu\sigma} J^{\mu\rho} + \eta^{\mu\sigma} J^{\nu\rho}) \\ [P^{\mu}, J^{\rho\sigma}] &= i(\eta^{\mu\rho} P^{\sigma} - \eta^{\mu\sigma} P^{\rho}) \end{aligned} \quad (1.64)$$

It's easy to check that $[K^i, P^0] = i P^i$, while $[J^i, P^0] = [P^i, P^0] = 0$: given that P^0 generates time translations, linear and angular momentum are conserved quantities, while \mathbf{K} is not.

§1.2.2.1 Field representations

Fields provide an infinite-dimensional representation of the Lorentz group as $J^{\mu\nu} = L^{\mu\nu} + S^{\mu\nu}$, where $S^{\mu\nu}$ does not depend on x^{μ} , but only on the spin of the field.

To represent P^{μ} on fields, their transformation law must be specified: all fields are required to be scalars under translations, independently of their spin. This means that, given a generic field $\phi(x)$, under a translation $x' = x + a$ it transforms as $\phi'(x') = \phi(x)$, so, under an infinitesimal translation $x' = x + \varepsilon$:

$$\begin{aligned} \delta_0 \phi &\equiv \phi'(x) - \phi(x) = \phi'(x' - \varepsilon) - \phi(x) = -\varepsilon^{\mu} \partial_{\mu} \phi(x) \\ &= e^{-i(-\varepsilon_{\mu}) P^{\mu}} \phi'(x') - \phi(x) = (e^{i\varepsilon_{\mu} P^{\mu}} - \mathbf{I}) \phi(x) = i\varepsilon_{\mu} P^{\mu} \phi(x) \end{aligned}$$

It is clear then that:

$$P^{\mu} = +i\partial^{\mu} \quad (1.65)$$

Explicitly, $P^0 = i\partial_t$ and $\mathbf{P} = -i\nabla$. It is trivial to check that these generators obey the Poincaré algebra.

§1.2.2.2 Particle representations

The Poincaré group can also be represented using the Hilbert space \mathcal{H} of a free particle as a basis. Denoting a generic state as $|\mathbf{p}, s\rangle \in \mathcal{H}$, where \mathbf{p} is the particle's momentum and s collectively labels all other quantum numbers, it is clear that \mathcal{H} is infinite-dimensional, as \mathbf{p} is a continuous unbounded variable.

Theorem 1.2.2 (Wigner's theorem)

On the Hilbert space of physical states, any symmetry transformation can be represented by a linear and unitary or anti-linear and anti-unitary operator.

By this theorem, Poincaré transformations can be represented by unitary matrices, i.e. \mathbf{J} , \mathbf{K} , \mathbf{P} and P^0 can be represented by hermitian operators. These representations can be labeled by Casimir operators, which for $\text{ISO}^+(1, 3)$ are easily found as $P_\mu P^\mu$ and $W_\mu W^\mu$, where W^μ is the **Pauli-Lubański operator**:

$$W^\mu := -\frac{1}{2}\epsilon^{\mu\nu\sigma\rho}J_{\nu\sigma}P_\rho \quad (1.66)$$

On single-particle states $P_\mu P^\mu = m^2$, while $W_\mu W^\mu$ can be conveniently computed in a particular frame (due to Lorentz invariance). If $m \neq 0$, this frame is the rest-frame of the particle:

$$W^\mu = -\frac{m}{2}\epsilon^{\mu\nu\sigma 0}J_{\nu\sigma} = \frac{m}{2}\delta^{\mu i}\epsilon^{ijk}J^{jk} = \delta^{\mu i}mJ^i$$

Therefore, on single-particle states of mass m and spin j , the Casimir operator takes the form:

$$W_\mu W^\mu = -m^2 j(j+1) \quad (1.67)$$

If $m = 0$, the rest-frame does not exist, but it is possible to choose a frame where $P^\mu = (\omega, 0, 0, \omega)$, where $W^0 = W^3 = \omega J^3$, $W^1 = \omega(J^1 - K^2)$ and $W^2 = \omega(J^2 + K^1)$, so that:

$$W_\mu W^\mu = -\omega^2 [(K^2 - J^1)^2 + (K^1 + J^2)^2] \quad (1.68)$$

It is clear that the $m \rightarrow 0$ limit is not trivial, and massive and massless representation need to be studied separately.

Massive representations Restricting to $m \in \mathbb{R}^+$ ($m^2 < 0$ states, called tachyons, are excluded), the massive representations are labeled by mass m and spin j : in fact, after a Lorentz transformation such that $P^\mu = (m, \mathbf{0})$, spatial rotations can still be performed, i.e. the subspace of single-particle states with momentum $P^\mu = (m, \mathbf{0})$ is still a basis for the representation of $\text{SU}(2)$ (as spinors must be included too). The group of transformations which leaves invariant a certain choice of P^μ is called the **little group**, so $\text{SU}(2)$ is the little group of massive single-particle states: massive representations are labelled by m and j , which means that massive particles of spin j have $2j + 1$ degrees of freedom.

Massless representations The little group for $P^\mu = (\omega, 0, 0, \omega)$ clearly is $\text{SO}(2)$, the group of rotations in the (x, y) plane generated by J^3 : as for any abelian group, its irreducible representations are one-dimensional, and they are labeled by the eigenvalue h of J^3 , which represents the angular momentum in the direction of propagation of the particle and is called **helicity**. Helicity can be shown to be quantized as $h \in \frac{1}{2}\mathbb{Z}_0$ (by topologic considerations on $\text{ISO}^+(1, 3) \equiv \mathbb{R}^4 \times \text{SL}(2, \mathbb{C})/\mathbb{Z}_2$).

As a consequence, massless particles only have one degree of freedom and are characterized by their helicity h . As $\text{SO}(2) \cong \text{U}(1)$, on a state of helicity h the little group is represented as:

$$U(\theta) = e^{-ih\theta} \quad (1.69)$$

Although massless particles with opposite helicities are logically two different species of particles, it can be written as $h = \hat{\mathbf{p}} \cdot \hat{\mathbf{J}}$ (unit vectors), so h is a pseudoscalar such that $h \mapsto -h$ under parity: this means that, if the interaction conserves parity, h and $-h$ must be symmetric.

Example 1.2.5 (Gauge bosons)

The electromagnetic and gravitational interactions conserve parity, thus photons and gravitons must be a basis for the representation of both $\text{ISO}^+(1, 3)$ and parity: photons can have $h = \pm 1$ (left- and right-handed), while gravitons have $h = \pm 2$.

Example 1.2.6 (Neutrinos)

Neutrinos only interact via the weak interaction, which does not conserve parity, and in fact the two states $h = \pm \frac{1}{2}$ are different particles: neutrinos have $h = -\frac{1}{2}$, while antineutrinos have $h = +\frac{1}{2}$.

§1.3 Classical equations of motion

Consider a **local field theory** of fields $\{\phi_i(x)\}_{i \in \mathcal{I}} \equiv \phi(x)$, where $x \in \mathbb{R}^{1,3}$ is a point in Minkowski spacetime. Its Lagrangian takes the form:

$$L = \int d^3x \mathcal{L}(\phi, \partial_\mu \phi) \quad (1.70)$$

where \mathcal{L} is the **Lagrangian density** of the theory (often referred to simply as the Lagrangian), which depends only on a finite number of derivatives. The action is then:

$$\mathcal{S} = \int dt L = \int d^4x \mathcal{L}(\phi, \partial_\mu \phi) \quad (1.71)$$

The integration is carried on the whole space-time, with usual boundary conditions that all fields decrease sufficiently fast at infinity; this also allows dropping all boundary terms.

Theorem 1.3.1 (Euler-Lagrange equations)

The **stationary action principle** $\delta\mathcal{S} = 0$ determines the classical equations of motion:

$$\frac{\partial \mathcal{L}}{\partial \phi_i} - \partial_\mu \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi_i)} = 0 \quad (1.72)$$

Proof. Varying Eq. 1.71:

$$\delta\mathcal{S} = \int d^4x \sum_{i \in \mathcal{I}} \left[\frac{\partial \mathcal{L}}{\partial \phi_i} \delta\phi_i + \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi_i)} \delta(\partial_\mu \phi_i) \right] = \int d^4x \sum_{i \in \mathcal{I}} \left[\frac{\partial \mathcal{L}}{\partial \phi_i} - \partial_\mu \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi_i)} \right] \delta\phi_i = 0$$

□

Corollary 1.3.1.1

Two Lagrangians which differ by a total divergence $\mathcal{L}' = \mathcal{L} + \partial_\mu K^\mu$ yield the same equations of motion.

Proof. This is a consequence of Stokes theorem:

$$\int_\Sigma d^4x \partial_\mu K^\mu = \int_{\partial\Sigma} dA n_\mu K^\mu$$

with n_μ normal vector to the $\partial\Sigma$ hypersurface.

□

From the Lagrangian, it is possible to define the conjugate momenta and the Hamiltonian density:

$$\Pi_i(x) := \frac{\partial \mathcal{L}}{\partial (\partial_0 \phi_i)} \quad (1.73)$$

$$\mathcal{H} = \sum_{i \in \mathcal{I}} \Pi_i(x) \partial_0 \phi_i(x) - \mathcal{L} \quad (1.74)$$

Unlike the Hamiltonian formalism, the Lagrangian formalism keeps Lorentz covariance explicit.

§1.3.1 Noether's theorem

Definition 1.3.1 (Infinitesimal transformation)

Given a field theory with fields $\{\phi_i\}_{i=1,\dots,k}$ and action $\mathcal{S}[\phi]$, an **infinitesimal transformation** parametrized by $\{\varepsilon^a\}_{a=1,\dots,N} : |\varepsilon^a| \ll 1$ is defined by two sets of functions $\{A_a^\mu(x)\}_{a=1,\dots,N}$ and $\{F_{i,a}(\phi, \partial\phi)\}_{i=1,\dots,k; a=1,\dots,N}$ such that:

$$\begin{aligned} x^\mu &\mapsto x'^\mu = x^\mu + \varepsilon^a A_a^\mu(x) \\ \phi_i(x) &\mapsto \phi'_i(x) = \phi_i(x) + \varepsilon^a F_{i,a}(\phi, \partial\phi) \end{aligned} \quad (1.75)$$

Definition 1.3.2 (Symmetry transformation)

An infinitesimal transformation is a **symmetry transformation** if it leaves $\mathcal{S}[\phi]$ invariant, regardless of ϕ being a solution of the equations of motion. It can further be classified as:

- **global symmetry**, if $\varepsilon^a \equiv \text{const.}$;
- **local symmetry**, if $\varepsilon^a = \varepsilon^a(x)$.

Symmetry transformations which leave spacetime unchanged, i.e. with $A_a^\mu(x) = 0$, are called **internal symmetries**.

Theorem 1.3.2 (Noether's theorem)

Given a global (but not local) symmetry parametrized by N generators, then there are N conserved currents $\{j_a^\mu(\phi)\}_{a=1,\dots,N}$ such that:

$$\partial_\mu j_a^\mu(\phi^{\text{cl}}) = 0 \quad (1.76)$$

where ϕ^{cl} is a classical solution of the equations of motion.

Proof. First, consider an infinitesimal transformation with slowly-varying parameters, that is $l |\partial_\mu \varepsilon^a| \ll |\varepsilon^a|$ (l characteristic scale of the field theory): being it not a local symmetry, $\delta\mathcal{S} \neq 0$ at $o(\varepsilon)$ and:

$$\mathcal{S}[\phi'] = \mathcal{S}[\phi] + \int d^4x [\varepsilon^a(x) K_a(\phi) - (\partial_\mu \varepsilon^a(x)) j_a^\mu(\phi) + o(\partial\partial\varepsilon)] + o(\varepsilon^2)$$

If $\varepsilon^a \equiv \text{const.}$ (global symmetry) then $\delta\mathcal{S}[\phi] = 0 \forall \phi$, therefore $K_a(\phi) = 0 \forall \phi$ (independent of ε). Assuming $\varepsilon^a(x) \rightarrow 0$ sufficiently fast as $x \rightarrow \infty$, then integration by parts yields:

$$\mathcal{S}[\phi'] = \mathcal{S}[\phi] + \int d^4x \varepsilon^a(x) \partial_\mu j_a^\mu(\phi) + o(\partial\partial\varepsilon) + o(\varepsilon^2)$$

This expression is independent of the choice of ϕ . Moreover, note that Eq. 1.75 can be rewritten as an internal transformation by setting:

$$\phi_i(x) \mapsto \phi'_i(x) = \phi_i(x - \varepsilon^a A_a) + \varepsilon^a F_{i,a} = \phi_i(x) + \varepsilon^a F_{i,a} - \varepsilon^a A_a^\mu \partial_\mu \phi_i \equiv \phi_i(x) + \delta\phi_i(x)$$

$\delta\phi_i(x)$ vanishes at infinity, therefore it is the kind of variation used to derive the equations of motion: choosing $\phi \equiv \phi^{\text{cl}}$ classical solution then implies $\delta\mathcal{S} = 0$ independently of ε , i.e. the thesis. \square

These are often called **Noether currents**, and the associated **Noether charges** are defined as:

$$Q_a := \int d^3x j_a^0(t, \mathbf{x}) \quad (1.77)$$

These are time-independent, as $\partial_0 Q_a = \int d^3x \partial_0 j_a^0 = - \int d^3x \partial_i j_a^i$: on all spacetime it vanishes by divergence theorem (fields vanish at infinity), but on a finite volume it yields a boundary term interpreted as the incoming and outgoing flux.

Proposition 1.3.1 (Noether currents)

The explicit expression of Noether currents is:

$$j_a^\mu = \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi_i)} [A_a^\nu(x) \partial_\nu \phi_i - F_{i,a}(\phi, \partial \phi)] - A_a^\mu(x) \mathcal{L} \quad (1.78)$$

Proof. Varying the action at $o(\partial \varepsilon)$:

$$\delta_\varepsilon \mathcal{S} = \delta_\varepsilon \int d^4x \mathcal{L} = \int \left[\delta_\varepsilon(d^4x) \mathcal{L} + d^4x \left(\frac{\partial \mathcal{L}}{\partial \phi_i} \delta_\varepsilon \phi_i + \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi_i)} \delta_\varepsilon(\partial_\mu \phi_i) \right) \right]$$

The Jacobian of Eq. 1.75 gives $d^4x \mapsto d^4x (1 + A_a^\mu \partial_\mu \varepsilon^a) + o(\varepsilon)$, while $\delta_\varepsilon \phi_i$ is not $o(\partial \varepsilon)$ and:

$$\delta_\varepsilon(\partial_\mu \phi_i) = \frac{\partial \phi'_i}{\partial x'^\mu} - \frac{\partial \phi_i}{\partial x^\mu} = \frac{\partial x^\nu}{\partial x'^\mu} \frac{\partial}{\partial x^\nu} (\phi_i + \varepsilon^a F_{i,a}) - \frac{\partial \phi_i}{\partial x^\mu} = -(\partial_\mu \varepsilon^a) (A_a^\nu \partial_\nu \phi_i - F_{i,a}) + o(\varepsilon)$$

The thesis follows from $\delta_\varepsilon \mathcal{S} = - \int d^4x (\partial_\mu \varepsilon^a) j_a^\mu + o(\partial \varepsilon) + o(\varepsilon^2)$. \square

If the considered infinitesimal transformation is not a global symmetry, then $\delta_\varepsilon \mathcal{S}$ has a non-vanishing $o(\varepsilon)$ term which gives rise to a quasi-conserved current:

$$\partial_\mu j_\mu^a = -(\delta_a \mathcal{L})_{\text{global}} \quad (1.79)$$

§1.3.1.1 Energy-momentum tensor

Consider spacetime translations: as all fields must be scalars under these transformations, they define a Noether current. In particular, translations have $A_\nu^\mu = \delta_\nu^\mu$ and $F_{i,\mu} = 0$ (the parameter index is a Lorentz index), so the conserved current is the **energy-momentum tensor** $j_\nu^\mu \equiv \theta^\mu_\nu$:

$$\theta^{\mu\nu} = \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi_i)} \partial^\nu \phi_i - \eta^{\mu\nu} \mathcal{L} \quad (1.80)$$

which is covariantly conserved on classical solutions of the equations of motion. The conserved Noether charge associated to the energy-momentum tensor is the **four-momentum**:

$$P^\mu := \int d^3x \theta^{0\mu} \quad (1.81)$$

The energy-momentum tensor so defined is not symmetric, however it can be made to via a tensor $A^{\rho\mu\nu}$ which is anti-symmetric w.r.t. (ρ, μ) : $T^{\mu\nu} \equiv \theta^{\mu\nu} + \partial_\rho A^{\rho\mu\nu}$ is physically equivalent from $\theta^{\mu\nu}$, as the second term is a vanishing spatial divergence in the definition of P^μ and is contracted to 0 in the conservation law.

§1.3.2 Scalar fields

§1.3.2.1 Real scalar fields

Consider a real scalar field ϕ : a non-trivial Poincaré-invariant action must contain $\partial_\mu \phi$ and must saturate each Lorentz index. For example:

$$\mathcal{S}[\phi] = \frac{1}{2} \int d^4x (\partial_\mu \phi \partial^\mu \phi - m^2 \phi^2) \quad (1.82)$$

The resulting equation of motion is the **Klein-Gordon equation**:

$$(\square + m^2) \phi = 0 \quad (1.83)$$

where $\square \equiv \partial_\mu \partial^\mu$. A plane wave $e^{\pm i p_\mu x^\mu}$ is a solution if $p^2 = m^2$, so the KG equation imposes the relativistic dispersion relation and m can be interpreted as the mass. As ϕ must be real, the general solution is a superposition of waves:

$$\phi(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} [a_{\mathbf{p}} e^{-i p_\mu x^\mu} + a_{\mathbf{p}}^* e^{i p_\mu x^\mu}]_{p^0=E_{\mathbf{p}}} \quad (1.84)$$

The positive energy solution has $E_{\mathbf{p}} = +\sqrt{\mathbf{p}^2 + m^2}$, but it contains both **positive** and **negative frequency modes** $e^{\mp i p_\mu x^\mu}$, while the $(2E_{\mathbf{p}})^{-1/2}$ factor is a convenient normalization of the $a_{\mathbf{p}}$ coefficients. The overall normalization of $\mathcal{S}[\phi]$ does not influence the equations of motion, however it is important for obtaining a positive-definite Hamiltonian. As the momentum conjugate to ϕ is $\Pi_\phi = \partial_0 \phi$, the Hamiltonian density is found as:

$$\mathcal{H} = \frac{1}{2} [\Pi_\phi^2 + (\nabla \phi)^2 + m^2 \phi^2] \quad (1.85)$$

The energy-momentum tensor is computed to be:

$$\theta^{\mu\nu} = \partial^\mu \phi \partial^\nu \phi - \eta^{\mu\nu} \mathcal{L} \quad (1.86)$$

It is trivial to see that $\theta^{00} = \mathcal{H}$: the Hamiltonian is the conserved charge related to the invariance under time translations.

To compute the conserved currents associated to Lorentz invariance, it is convenient to label the transformation parameters $\omega^{\mu\nu}$ by an anti-symmetric pair of Lorentz indices, so that **Eq. 1.75** become:

$$x^\mu \mapsto x'^\mu = x^\mu + \omega^\mu{}_\nu x^\nu = x^\mu + \frac{1}{2} \omega^{\rho\sigma} (\delta^\mu_\rho x_\sigma - \delta^\mu_\sigma x_\rho) \equiv x^\mu + \frac{1}{2} A^\mu_{(\rho\sigma)} \omega^{\rho\sigma}$$

As $F_{i,a} = 0$, from **Eq. 1.86** the conserved currents are:

$$j^{(\rho\sigma)\mu} = x^\rho \theta^{\mu\sigma} - x^\sigma \theta^{\mu\rho} \quad (1.87)$$

For spatial rotations, the conserved charge is:

$$M^{ij} = \int d^3x (x^i \theta^{0j} - x^j \theta^{0i}) = \int d^3x \partial_0 \phi (x^i \partial^j - x^j \partial^i) \phi = \frac{i}{2} \int d^3x [\phi L^{ij} (\partial_0 \phi) - (\partial_0 \phi) L^{ij} \phi]$$

where integration by parts was carried and L^{ij} is defined by **Eq. 1.55**. This can be generalized.

Definition 1.3.3 (Scalar product)

Given two real scalar fields ϕ_1 and ϕ_2 , their **scalar product** is defined as:

$$\langle \phi_1 | \phi_2 \rangle := \frac{i}{2} \int d^3x \phi_1 \overleftrightarrow{\partial}_0 \phi_2 \quad (1.88)$$

where $f \overleftrightarrow{\partial}_\mu g := f \partial_\mu g - (\partial_\mu f) g$.

Proposition 1.3.2

If ϕ_1 and ϕ_2 are KG solutions, then $\langle \phi_1 | \phi_2 \rangle$ is time-independent.

Proof. By the KG equation:

$$\partial_0 [\phi_1 \partial_0 \phi_2 - (\partial_0 \phi_1) \phi_2] = \phi_1 \partial_0^2 \phi_2 - (\partial_0^2 \phi_1) \phi_2 = \phi_1 \Delta \phi_2 - (\Delta \phi_1) \phi_2$$

which vanishes after integration by parts. \square

Note that this scalar product is not positive-definite.

Theorem 1.3.3 (Conserved charges)

Given a symmetry represented by a Lie group and a representation $L^{\mu\nu}$ of its generators as operators acting on fields, the value of the associated conserved charges on a solution ϕ of the equations of motion is:

$$M^{\mu\nu} = \langle \phi | L^{\mu\nu} | \phi \rangle \quad (1.89)$$

Example 1.3.1 (Four-momentum)

Applying Th. 1.3.3 to four-momentum $P^\mu = \langle \phi | i \partial^\mu | \phi \rangle$; for example, the $\mu = 0$ component is:

$$\begin{aligned} P^0 &= \langle \phi | i \partial^0 | \phi \rangle = \langle \phi | i \partial_0 | \phi \rangle = \frac{i}{2} \int d^3x [\phi (i \partial_0) \partial_0 \phi - (\partial_0 \phi) i \partial_0 \phi] \\ &= \frac{1}{2} \int d^3x [-\phi \partial_0^2 \phi + (\partial_0 \phi)^2] = \frac{1}{2} \int d^3x [-\phi (\Delta - m^2) \phi + (\partial_0 \phi)^2] \\ (\text{int. by parts}) &= \frac{1}{2} \int d^3x [(\nabla \phi)^2 + m^2 \phi^2 + (\partial_0 \phi)^2] = \frac{1}{2} \int d^3x \theta^{00} \end{aligned}$$

The free KG action can be generalized to a self-interacting real scalar field introducing a general potential:

$$\mathcal{S}[\phi] = \int d^4x \left[\frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) \right] \quad (1.90)$$

The quadratic term in $V(\phi)$ is the mass term, while higher-order terms describe the self-interaction.

§1.3.2.2 Complex scalar fields

Consider now two real scalar fields ϕ_1, ϕ_2 with the same mass m and combine them into a single complex scalar field $\phi \equiv \frac{1}{\sqrt{2}}(\phi_1 + i\phi_2)$. The KG action is the sum of the single actions and may

be written as:

$$\mathcal{S}[\phi] = \int d^4x (\partial_\mu \phi^* \partial^\mu \phi - m^2 \phi^* \phi) \quad (1.91)$$

Considering ϕ and ϕ^* as independent variables one obtains the KG equation, which yields the same mode expansion as Eq. 1.84:

$$\phi(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} [a_{\mathbf{p}} e^{-ip_\mu x^\mu} + b_{\mathbf{p}}^* e^{ip_\mu x^\mu}]_{p^0=E_{\mathbf{p}}} \quad (1.92)$$

Now $a_{\mathbf{p}}$ and $b_{\mathbf{p}}$ are independent, since there's no reality condition on ϕ .

Electric charge An interesting property of the complex KG field is the existence of a global U(1) symmetry of the action: $\mathcal{S}[\phi] \mapsto \mathcal{S}[\phi]$ under $\phi(x) \mapsto e^{i\theta} \phi(x)$. The associated Noether current can be computed from Eq. 1.78, using $\phi_i = (\phi, \phi^*)$, $A_a^\nu = 0$ and $F_{i,a} = (i, -i)$:

$$j_\mu = -i(\phi \partial_\mu \phi^* - \phi^* \partial_\mu \phi) = i\phi^* \overleftrightarrow{\partial}_\mu \phi \quad (1.93)$$

The conserved U(1) charge is then $Q_{U(1)} = i \int d^3x \phi^* \overleftrightarrow{\partial}_0 \phi = \langle \phi | \phi \rangle$, which is consistent with Th. 1.3.3 as the generator of U(1) is the identity operator.

§1.3.3 Spinor fields

§1.3.3.1 Weyl fields

Consider the theory of a left/right-handed Weyl field: recalling that $\psi_L^\dagger \sigma^\mu \psi_L$ and $\psi_R \bar{\sigma}^\mu \psi_R$ are 4-vectors, with $\sigma^\mu \equiv (1, \boldsymbol{\sigma})$ and $\bar{\sigma}^\mu \equiv (1, -\boldsymbol{\sigma})$, the kinetic term of the Lorentz-invariant Lagrangian can be written as:

$$\mathcal{L}_L = i\psi_L^\dagger \bar{\sigma}^\mu \partial_\mu \psi_L \quad \mathcal{L}_R = i\psi_R^\dagger \sigma^\mu \partial_\mu \psi_R \quad (1.94)$$

The i factor ensures the reality of the Lagrangian, as σ matrices are hermitian. The Lagrangian does not depend on $\partial_\mu \psi^*$, thus the Euler-Lagrange equations are:

$$(\partial_0 - \sigma^i \partial_i) \psi_L = 0 \quad (\partial_0 + \sigma^i \partial_i) \psi_R = 0 \quad (1.95)$$

These are known as **Weyl equations**: by $\sigma^i \sigma^j = \delta^{ij} + i\epsilon^{ijk} \sigma^k$, these equations imply two massless KG equations (assuming regular functions, so that $\partial_i \partial_j = \partial_j \partial_i$). Considering positive-energy plane-wave solutions of the form $\psi_L(x) = u_L \exp(-ip_\mu x^\mu) = u_L \exp(-iEt + i\mathbf{p} \cdot \mathbf{x})$, where u_L is a constant spinor, Weyl equations become:

$$\frac{\mathbf{p} \cdot \boldsymbol{\sigma}}{E} u_L = -u_L \quad \frac{\mathbf{p} \cdot \boldsymbol{\sigma}}{E} u_R = u_R$$

As these are massless fields $E = |\mathbf{p}|$, and since for $s = \frac{1}{2}$ fields $\mathbf{J} = \frac{\boldsymbol{\sigma}}{2}$ these equations become:

$$(\hat{\mathbf{p}} \cdot \mathbf{J}) u_L = -\frac{1}{2} u_L \quad (\hat{\mathbf{p}} \cdot \mathbf{J}) u_R = \frac{1}{2} u_R$$

These show that left/right-handed Weyl massless Weyl spinors have helicity $h = \mp \frac{1}{2}$, consistent with the fact that massless particles are helicity eigenstates.

The energy-momentum tensor can be computed from Eq. 1.80, noting that on classical solutions (by Weyl equations) the Lagrangian of the theory vanishes:

$$\theta^{\mu\nu} = i\psi_L^\dagger \bar{\sigma}^\mu \partial^\nu \psi_L \quad \theta^{\mu\nu} = i\psi_R^\dagger \sigma^\mu \partial^\nu \psi_R \quad (1.96)$$

Moreover, note that the Lagrangian is invariant under a global U(1) internal transformation with Noether currents and conserved charges:

$$\begin{aligned} j^\mu &= \psi_L^\dagger \bar{\sigma}^\mu \psi_L & j^\mu &= \psi_R^\dagger \sigma^\mu \psi_R \\ Q_{U(1)} &= \int d^3x \psi_L^\dagger \psi_L & Q_{U(1)} &= \int d^3x \psi_R^\dagger \psi_R \end{aligned}$$

Weyl Lagrangians are not invariant under parity, as $\psi_L \leftrightarrow \psi_R$.

Example 1.3.2 (Neutrinos)

Neutrinos are divided into three leptonic families: ν_e , ν_μ and ν_τ . Although they are massive $s = \frac{1}{2}$, in most contexts their mass can be neglected, thus they can be described by Weyl spinors: in particular, neutrinos by left-handed Weyl spinors and antineutrinos by right-handed Weyl spinors.

§1.3.3.2 Dirac fields

Consider a theory with both a left-handed and a right-handed Weyl spinor: one can construct two new Lorentz scalars, $\psi_L^\dagger \psi_R$ and $\psi_R^\dagger \psi_L$, as from Eq. 1.43-1.44 it is easy to check that $\Lambda_L^\dagger \Lambda_R = \Lambda_R^\dagger \Lambda_L = 1$. Two real combinations are $\psi_L^\dagger \psi_R + \psi_R^\dagger \psi_L$ and $i(\psi_L^\dagger \psi_R - \psi_R^\dagger \psi_L)$: under parity $\psi_L \leftrightarrow \psi_R$, so the former is a scalar and the latter a pseudoscalar. The Dirac Lagrangian then is:

$$\mathcal{L}_D = i\psi_L^\dagger \bar{\sigma}^\mu \partial_\mu \psi_L + i\psi_R^\dagger \sigma^\mu \partial_\mu \psi_R - m(\psi_L^\dagger \psi_R + \psi_R^\dagger \psi_L) \quad (1.97)$$

This Lagrangian is invariant under parity, as $\bar{\sigma}^\mu \partial_\mu \leftrightarrow \sigma^\mu \partial_\mu$ ($\partial_i \mapsto -\partial_i$). Considering ψ_i and ψ_i^* independent, the variation w.r.t. the latter yields:

$$i\bar{\sigma}^\mu \partial_\mu \psi_L = m\psi_R \quad i\sigma^\mu \partial_\mu \psi_R = m\psi_L \quad (1.98)$$

which is the **Dirac equation** in terms of Weyl spinors.

Proposition 1.3.3

The Dirac equation implies two massive Klein-Gordon equations.

Proof. Applying $i\sigma^\mu \partial_\mu$ to the first equation:

$$-\sigma^\mu \bar{\sigma}^\nu \partial_\mu \partial_\nu \psi_L = m i\sigma^\mu \partial_\mu \psi_R = m^2 \psi_L$$

Assuming that Ψ_L satisfies the hypotheses of Schwarz's Lemma, then $\partial_\mu \partial_\nu = \partial_\nu \partial_\mu$, and $\sigma^\mu \bar{\sigma}^\nu$ can be replaced by $\frac{1}{2}(\sigma^\nu \bar{\sigma}^\mu + \sigma^\mu \bar{\sigma}^\nu) = \frac{1}{2}(2\eta^{\mu\nu}) = \eta^{\mu\nu}$, yielding the thesis. \square

The m parameter can thus be interpreted as a mass, and now the two spinors are no longer helicity eigenstates. It is convenient to rewrite this theory in terms of a Dirac spinor, which defines the chiral representation:

$$\Psi = \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix}$$

The chiral representation of the Dirac algebra $\mathfrak{cl}_{1,3}(\mathbb{C})$ is defined as⁶:

$$\gamma^\mu \equiv \begin{bmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{bmatrix}$$

This allows to rewrite Dirac equation as:

$$(i\cancel{\partial} - m)\Psi = 0 \quad (1.100)$$

In the chiral representation, the Dirac adjoint simply is $\bar{\Psi} = (\psi_R^\dagger, \psi_L^\dagger)$, so the Dirac Lagrangian can be rewritten as:

$$\mathcal{L}_D = \bar{\Psi}(i\cancel{\partial} - m)\Psi \quad (1.101)$$

The γ^5 matrix in the chiral representation allows to define the projectors on left- and right-handed Weyl spinors:

$$\gamma^5 \equiv \begin{bmatrix} -I_2 & 0 \\ 0 & I_2 \end{bmatrix} \quad \Rightarrow \quad \frac{1 - \gamma^5}{2} \Psi = \begin{pmatrix} \psi_L \\ 0 \end{pmatrix} \quad \frac{1 + \gamma^5}{2} \Psi = \begin{pmatrix} 0 \\ \psi_R \end{pmatrix} \quad (1.102)$$

Proposition 1.3.4 (Unitary equivalence)

Given a constant $U \in \mathbb{C}^{2 \times 2} : U^\dagger = U^{-1}$, then under $\Psi \mapsto U\Psi$ the Lagrangian is invariant for $\gamma^\mu \mapsto U\gamma^\mu U^\dagger$.

Proof. Inserting $\Psi = U^\dagger \Psi'$ (and $\bar{\Psi} = \Psi'^\dagger U \gamma^0$) in the Dirac Lagrangian:

$$\mathcal{L}_D = \Psi'^\dagger U \gamma^0 (i\gamma^\mu \partial_\mu - m) U^\dagger \Psi' = \Psi'^\dagger U \gamma^0 U^\dagger (iU\gamma^\mu U^\dagger \partial_\mu - m) \Psi' = \bar{\Psi}' (i\gamma'^\mu \partial_\mu - m) \Psi'$$

where $\bar{\Psi}' \equiv \Psi'^\dagger \gamma^0$. The Lagrangian is thus unchanged. \square

It can be shown that the Dirac algebra is invariant under the transformation in [Prop. 1.3.4](#), thus it defines an equivalent representation of the algebra.

Proposition 1.3.5 (Lorentz invariance)

The Dirac equation is Lorentz invariant.

Proof. Recalling [Eq. 1.51](#):

$$\begin{aligned} (i\gamma^\mu \partial_\mu - m) \Psi(x) &\mapsto (i\gamma^\mu [\Lambda^{-1}]^\nu{}_\mu \partial_\nu - m) \Lambda_{\frac{1}{2}} \Psi(\Lambda^{-1}x) \\ &= \Lambda_{\frac{1}{2}} \Lambda_{\frac{1}{2}}^{-1} (i\gamma^\mu [\Lambda^{-1}]^\nu{}_\mu \partial_\nu - m) \Lambda_{\frac{1}{2}} \Psi(\Lambda^{-1}x) \\ &= \Lambda_{\frac{1}{2}} \left(i\Lambda_{\frac{1}{2}}^{-1} \gamma^\mu \Lambda_{\frac{1}{2}} [\Lambda^{-1}]^\nu{}_\mu \partial_\nu - m \right) \Psi(\Lambda^{-1}x) \\ &= \Lambda_{\frac{1}{2}} (i\Lambda^\mu{}_\sigma \gamma^\sigma [\Lambda^{-1}]^\nu{}_\mu \partial_\nu - m) \Psi(\Lambda^{-1}x) \\ &= \Lambda_{\frac{1}{2}} (i\gamma^\nu \partial_\nu - m) \Psi(\Lambda^{-1}x) \end{aligned}$$

⁶In terms of tensor products, the chiral representation of the Dirac algebra can be written as:

$$\gamma^0 = \sigma^1 \otimes I_2 \quad \gamma = i\sigma^2 \otimes \sigma \quad \gamma^5 = -\sigma^3 \otimes I_2 \quad (1.99)$$

Since 0 is a Lorentz scalar, this transformation law implies that $(i\not{\partial} - m) \Psi(y) = 0$, where $y = \Lambda^{-1}x$ are the coordinates in the new reference frame. This shows that the Dirac equation is Lorentz invariant. \square

Solutions The general solution to the Dirac equation is a superposition of plane waves, both with positive and negative-frequency modes respectively written as $\Psi(x) = u(p)e^{-ip_\mu x^\mu}$ and $\Psi(x) = v(p)e^{ip_\mu x^\mu}$, where $u(p)$ and $v(p)$ are Dirac spinors which, in the chiral representation, have both a left-handed and a right-handed Weyl spinor component. The Dirac equation then becomes:

$$(\not{p} - m) u(p) = 0 \quad (\not{p} + m) v(p) = 0 \quad (1.103)$$

Considering the massive positive-frequency solution in the rest frame, i.e. $p^\mu = (m, \mathbf{0})$, one finds $(\gamma^0 - 1)u(p) = 0$, which yields $u_L = u_R$: the KG equation imposes the mass-shell condition $p^2 = m^2$, but the Dirac equation, being a first-order equation, halves the number of independent degrees of freedom. A convenient normalization is:

$$u(p_0) = \sqrt{m} \begin{pmatrix} \xi \\ \xi \end{pmatrix}$$

where $\xi : \xi^\dagger \xi = 1$ is a two-component spinor which gives the spin orientation of the solution: $\xi = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$ for the spin-up solution and $\xi = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$ for the spin-down one.

The solution in a generic frame is obtained via a boost of the rest-frame solution, using transformation properties of left/right-handed spinors.

Proposition 1.3.6

The solution to the Dirac equation in a generic frame where $\mathbf{p} \parallel \mathbf{e}_z$ is given by:

$$u(p) = \begin{pmatrix} \begin{bmatrix} \sqrt{E-p^3} & 0 \\ 0 & \sqrt{E+p^3} \end{bmatrix} \xi \\ \begin{bmatrix} \sqrt{E+p^3} & 0 \\ 0 & \sqrt{E-p^3} \end{bmatrix} \xi \end{pmatrix} \quad (1.104)$$

Proof. Consider a boost along \mathbf{e}_z , parametrized by rapidity η , by Eq. 1.35:

$$\begin{pmatrix} E \\ p^3 \end{pmatrix} = \exp \left(\eta \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix} \right) \begin{pmatrix} m \\ 0 \end{pmatrix} = \left(\begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \cosh \eta + \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix} \sinh \eta \right) \begin{pmatrix} m \\ 0 \end{pmatrix} = \begin{pmatrix} m \cosh \eta \\ m \sinh \eta \end{pmatrix}$$

Thus, $E + p^3 = e^\eta m$ and $E - p^3 = e^{-\eta} m$. By Eq. 1.43-1.44, then:

$$\begin{aligned} u(p) &= \exp \left(-\frac{\eta}{2} \begin{bmatrix} \sigma_3 & 0 \\ 0 & -\sigma_3 \end{bmatrix} \right) \sqrt{m} \begin{pmatrix} \xi \\ \xi \end{pmatrix} = \left(\begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \cosh \frac{\eta}{2} - \begin{bmatrix} \sigma_3 & 0 \\ 0 & -\sigma_3 \end{bmatrix} \sinh \frac{\eta}{2} \right) \sqrt{m} \begin{pmatrix} \xi \\ \xi \end{pmatrix} \\ &= \begin{bmatrix} e^{\eta/2} \frac{1-\sigma^3}{2} + e^{-\eta/2} \frac{1+\sigma^3}{2} & 0 \\ 0 & e^{\eta/2} \frac{1+\sigma^3}{2} + e^{-\eta/2} \frac{1-\sigma^3}{2} \end{bmatrix} \sqrt{m} \begin{pmatrix} \xi \\ \xi \end{pmatrix} \\ &= \begin{bmatrix} \sqrt{E+p^3} \begin{bmatrix} 0 & 0 \\ 0 & 1 \end{bmatrix} + \sqrt{E-p^3} \begin{bmatrix} 1 & 0 \\ 0 & 0 \end{bmatrix} & 0 \\ 0 & \sqrt{E+p^3} \begin{bmatrix} 1 & 0 \\ 0 & 0 \end{bmatrix} + \sqrt{E-p^3} \begin{bmatrix} 0 & 0 \\ 0 & 1 \end{bmatrix} \end{bmatrix} \begin{pmatrix} \xi \\ \xi \end{pmatrix} \end{aligned}$$

which is the thesis. \square

This expression can be generalized for an arbitrary direction of \mathbf{p} :

$$u^s(p) = \begin{pmatrix} \sqrt{p_\mu \sigma^\mu} \xi^s \\ \sqrt{p_\mu \bar{\sigma}^\mu} \xi^s \end{pmatrix} \quad (1.105)$$

where the square root of a matrix is understood as taking the positive square root of its eigenvalues, and the polarization of ξ^s is made explicit.

It is convenient to work with specific spinors ξ : a useful choice are eigenvectors of σ^3 , in particular $\xi^1 = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$ (spin-up) and $\xi^2 = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$ (spin-down). In the ultra-relativistic (or massless) limit $p^\mu \rightarrow (E, 0, 0, E)$, so u^1 has only the right-handed component and u^2 only the left-handed one.

The negative-frequency solution is equivalent:

$$v^s(p) = \begin{pmatrix} \sqrt{p_\mu \sigma^\mu} \eta^s \\ -\sqrt{p_\mu \bar{\sigma}^\mu} \eta^s \end{pmatrix} \quad (1.106)$$

where the spin states are the same: $\eta^1 = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$ and $\eta^2 = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$.

Recalling the Dirac adjoint $\bar{u}^s(p) = u^{s\dagger}(p)\gamma^0$, $\bar{v}^s(p) = v^{s\dagger}(p)\gamma^0$ and the normalization choice $\xi^{r\dagger}\xi^s = \delta^{rs}$, $\eta^{r\dagger}\eta^s = \delta^{rs}$, several properties follow:

$$\bar{u}^r(p)u^s(p) = 2m\delta^{rs} \quad \bar{v}^r(p)v^s(p) = -2m\delta^{rs} \quad (1.107)$$

$$u^{r\dagger}(p)u^s(p) = 2E_{\mathbf{p}}\delta^{rs} \quad v^{r\dagger}(p)v^s(p) = 2E_{\mathbf{p}}\delta^{rs} \quad (1.108)$$

$$\bar{u}^r(p)v^s(p) = 0 \quad \bar{v}^r(p)u^s(p) = 0 \quad (1.109)$$

$$\bar{u}^r(p)\gamma^\mu u^s(p) = 2p^\mu\delta^{rs} \quad \bar{v}^r(p)\gamma^\mu v^s(p) = 2p^\mu\delta^{rs} \quad (1.110)$$

Note that $\bar{u}u \in \mathbb{C}$, while $u\bar{u} \in \mathbb{C}^{4 \times 4}$.

Proposition 1.3.7 (Spinor sums)

The sum over possible polarizations of a fermion yields:

$$\sum_s u^s(p)\bar{u}^s(p) = \not{p} + m \quad \sum_s v^s(p)\bar{v}^s(p) = \not{p} - m \quad (1.111)$$

Proof. Using $\sum_{s=1,2} \xi^s \xi^{s\dagger} = \mathbf{I}_2$:

$$\begin{aligned} \sum_s u^s(p)\bar{u}^s(p) &= \sum_s \begin{pmatrix} \sqrt{p_\mu \sigma^\mu} \xi^s \\ \sqrt{p_\mu \bar{\sigma}^\mu} \xi^s \end{pmatrix} (\xi^{s\dagger} \sqrt{p_\mu \bar{\sigma}^\mu} \quad \xi^{s\dagger} \sqrt{p_\mu \sigma^\mu}) \\ &= \begin{bmatrix} \sqrt{p_\mu \sigma^\mu} \sqrt{p_\mu \bar{\sigma}^\mu} & \sqrt{p_\mu \sigma^\mu} \sqrt{p_\mu \sigma^\mu} \\ \sqrt{p_\mu \bar{\sigma}^\mu} \sqrt{p_\mu \bar{\sigma}^\mu} & \sqrt{p_\mu \bar{\sigma}^\mu} \sqrt{p_\mu \sigma^\mu} \end{bmatrix} = \begin{bmatrix} m & p_\mu \sigma^\mu \\ p_\mu \bar{\sigma}^\mu & m \end{bmatrix} \end{aligned}$$

where the identity $(p_\mu \sigma^\mu)(p_\mu \bar{\sigma}^\mu) = p^2 = m^2$ was used. \square

Chiral symmetry Consider the massless Dirac Lagrangian. In this case, the action has two global internal symmetries:

$$\psi_L \mapsto e^{i\theta_L} \psi_L \quad \psi_R \mapsto e^{i\theta_R} \psi_R$$

Therefore, this theory is symmetric under $U(1) \times U(1)$. These can be rewritten as two distinct transformations of the Dirac spinor:

$$\Psi \mapsto e^{i\alpha} \Psi \quad (1.112)$$

$$\Psi \mapsto e^{i\beta\gamma^5} \Psi \quad (1.113)$$

where $\alpha \equiv \theta_L = \theta_R$ and $\beta \equiv -\theta_L = \theta_R$ respectively (to prove that the latter is a symmetry, from $\{\gamma^5, \gamma^\mu\} = 0$ it follows that $\gamma^\mu e^{i\beta\gamma^5} = e^{-i\beta\gamma^5} \gamma^\mu$). The former symmetry is called the vector $U(1)_V$ symmetry, with conserved Noether current the **vector current** $j_V^\mu := \bar{\Psi} \gamma^\mu \Psi$, while the latter is called the axial $U(1)_A$ symmetry, with conserved **axial current** $j_A^\mu := \bar{\Psi} \gamma^\mu \gamma^5 \Psi$.

While the vector symmetry is always a symmetry of the Dirac field, the axial symmetry is only for the massless Dirac field, as:

$$\partial_\mu j_A^\mu = 2im\bar{\Psi} \gamma^5 \Psi \quad (1.114)$$

Canonical Quantization

§2.1 Scalar fields

As for the quantization of a classical system in Quantum Mechanics, the quantization of a scalar field theory is performed promoting $\phi(t, \mathbf{x})$ and $\Pi(t, \mathbf{x})$ to hermitian operators in the Heisenberg picture and imposing the canonical equal-time commutation relation:

$$[\phi(t, \mathbf{x}), \Pi(t, \mathbf{y})] = i\delta^{(3)}(\mathbf{x} - \mathbf{y}) \quad (2.1)$$

while of course $[\phi(t, \mathbf{x}), \phi(t, \mathbf{y})] = [\Pi(t, \mathbf{x}), \Pi(t, \mathbf{y})] = 0$.

§2.1.1 Real scalar fields

A real scalar field is promoted to a real hermitian operator. In particular, by Eq. 1.84:

$$\phi(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} [a_{\mathbf{p}} e^{-ip_{\mu}x^{\mu}} + a_{\mathbf{p}}^{\dagger} e^{ip_{\mu}x^{\mu}}]_{p^0=E_{\mathbf{p}}} \quad (2.2)$$

In terms of creation and annihilation operators, the commutator Eq. 2.1 reads:

$$[a_{\mathbf{p}}, a_{\mathbf{q}}^{\dagger}] = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q}) \quad (2.3)$$

while $[a_{\mathbf{p}}, a_{\mathbf{q}}] = [a_{\mathbf{p}}^{\dagger}, a_{\mathbf{q}}^{\dagger}] = 0$. These can be regarded as the creation and annihilation operators of a collection of harmonic oscillators, one for each value of the momentum \mathbf{p} : the **Fock space** of the real scalar field can thus be constructed analogously to the Hilbert space of the harmonic oscillator.

Defining the **vacuum state** $|0\rangle : a_{\mathbf{p}} |0\rangle = 0 \ \forall \mathbf{p}$, suitably normalized as $\langle 0|0\rangle = 1$, the generic state of the Fock space is:

$$|\mathbf{p}_1, \dots, \mathbf{p}_n\rangle = \sqrt{2E_{\mathbf{p}_1}} \dots \sqrt{2E_{\mathbf{p}_n}} a_{\mathbf{p}_1}^{\dagger} \dots a_{\mathbf{p}_n}^{\dagger} |0\rangle \quad (2.4)$$

Proposition 2.1.1 (Normalization)

For one-particle states:

$$\langle \mathbf{p}_1 | \mathbf{p}_2 \rangle = 2E_{\mathbf{p}_1} (2\pi)^3 \delta^{(3)}(\mathbf{p}_1 - \mathbf{p}_2) \quad (2.5)$$

Proof. By Eq. 2.3, $\langle \mathbf{p}_1 | \mathbf{p}_2 \rangle = \sqrt{2E_{\mathbf{p}_1}} \sqrt{2E_{\mathbf{p}_2}} \langle 0 | a_{\mathbf{p}_1} a_{\mathbf{p}_2}^{\dagger} | 0 \rangle$, but $\langle 0 | a_{\mathbf{p}_1} a_{\mathbf{p}_2}^{\dagger} | 0 \rangle = \langle 0 | [a_{\mathbf{p}_1}, a_{\mathbf{p}_2}^{\dagger}] | 0 \rangle$, hence $\langle \mathbf{p}_1 | \mathbf{p}_2 \rangle = 2E_{\mathbf{p}_1} (2\pi)^3 \delta^{(3)}(\mathbf{p}_1 - \mathbf{p}_2)$. \square

Lemma 2.1.1

The combination $E_{\mathbf{p}}\delta^{(3)}(\mathbf{p} - \mathbf{q})$ is Lorentz invariant.

Proof. Recall that:

$$\delta(f(x) - f(x_0)) = \frac{1}{|f'(x_0)|} \delta(x - x_0) \quad (2.6)$$

A Lorentz boost along \mathbf{e}_i yields $p'_i = \gamma(p_i + \beta E)$ and $E' = \gamma(E + \beta p_i)$, so:

$$\delta^{(3)}(\mathbf{p} - \mathbf{q}) = \delta^{(3)}(\mathbf{p}' - \mathbf{q}') \frac{dp'_i}{dp_i} = \delta^{(3)}(\mathbf{p}' - \mathbf{q}') \gamma \left(1 + \beta \frac{dE}{dp_i} \right)$$

Note that $p^2 = E^2 - \mathbf{p}^2$ is a Lorentz invariant, thus $E dE = p_i dp_i$, so:

$$\delta^{(3)}(\mathbf{p} - \mathbf{q}) = \delta^{(3)}(\mathbf{p}' - \mathbf{q}') \frac{\gamma(E + \beta p_i)}{E} = \delta^{(3)}(\mathbf{p}' - \mathbf{q}') \frac{E'}{E}$$

Clearly, then, $E_{\mathbf{p}}\delta^{(3)}(\mathbf{p} - \mathbf{q})$ is Lorentz invariant. \square

This explains the choice of normalization.

Proposition 2.1.2 (KG Hamiltonian)

The Hamiltonian of a real scalar field can be written as:

$$H = \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} \left(a_{\mathbf{p}}^\dagger a_{\mathbf{p}} + \frac{1}{2} [a_{\mathbf{p}}, a_{\mathbf{p}}^\dagger] \right) \quad (2.7)$$

Proof. First of all, from Eq. 2.2:

$$\begin{aligned} \Pi(t, \mathbf{x}) &= \partial_0 \phi(t, \mathbf{x}) = -i \int \frac{d^3p}{(2\pi)^3} \sqrt{\frac{E_{\mathbf{p}}}{2}} \left[a_{\mathbf{p}} e^{-ip_\mu x^\mu} - a_{\mathbf{p}}^\dagger e^{ip_\mu x^\mu} \right]_{p^0=E_{\mathbf{p}}} \\ &= -i \int \frac{d^3p}{(2\pi)^3} \sqrt{\frac{E_{\mathbf{p}}}{2}} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} - a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] e^{i\mathbf{p} \cdot \mathbf{x}} \end{aligned}$$

$$\nabla \phi(t, \mathbf{x}) = i \int \frac{d^3p}{(2\pi)^3} \frac{\mathbf{p}}{\sqrt{2E_{\mathbf{p}}}} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} + a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] e^{i\mathbf{p} \cdot \mathbf{x}}$$

Inserting these expressions in Eq. 1.85:

$$\begin{aligned} \mathcal{H} &= \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} \left\{ -\frac{\sqrt{E_{\mathbf{p}}E_{\mathbf{q}}}}{2} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} - a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \left[a_{\mathbf{q}} e^{-iE_{\mathbf{q}}t} - a_{-\mathbf{q}}^\dagger e^{iE_{\mathbf{q}}t} \right] + \right. \\ &\quad \left. + \frac{-\mathbf{p} \cdot \mathbf{q} + m^2}{2\sqrt{E_{\mathbf{p}}E_{\mathbf{q}}}} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} + a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \left[a_{\mathbf{q}} e^{-iE_{\mathbf{q}}t} + a_{-\mathbf{q}}^\dagger e^{iE_{\mathbf{q}}t} \right] \right\} e^{i(\mathbf{p}+\mathbf{q}) \cdot \mathbf{x}} \end{aligned}$$

Recall the identity:

$$\int \frac{d^3x}{(2\pi)^3} e^{i(\mathbf{p}-\mathbf{q}) \cdot \mathbf{x}} = \delta^{(3)}(\mathbf{p} - \mathbf{q}) \quad (2.8)$$

Then (using $E_{-\mathbf{p}} = E_{\mathbf{p}}$):

$$\begin{aligned}
H &= \int d^3x \mathcal{H} \\
&= \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} \int d^3q \left\{ -\frac{\sqrt{E_{\mathbf{p}}E_{\mathbf{q}}}}{2} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} - a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \left[a_{\mathbf{q}} e^{-iE_{\mathbf{q}}t} - a_{-\mathbf{q}}^\dagger e^{iE_{\mathbf{q}}t} \right] + \right. \\
&\quad \left. + \frac{-\mathbf{p} \cdot \mathbf{q} + m^2}{2\sqrt{E_{\mathbf{p}}E_{\mathbf{q}}}} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} + a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \left[a_{\mathbf{q}} e^{-iE_{\mathbf{q}}t} + a_{-\mathbf{q}}^\dagger e^{iE_{\mathbf{q}}t} \right] \right\} \delta^{(3)}(\mathbf{p} + \mathbf{q}) \\
&= \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} \left\{ -\frac{E_{\mathbf{p}}}{2} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} - a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \left[a_{-\mathbf{p}} e^{-iE_{\mathbf{p}}t} - a_{\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] + \right. \\
&\quad \left. + \frac{\mathbf{p}^2 + m^2}{2E_{\mathbf{p}}} \left[a_{\mathbf{p}} e^{-iE_{\mathbf{p}}t} + a_{-\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \left[a_{-\mathbf{p}} e^{-iE_{\mathbf{p}}t} + a_{\mathbf{p}}^\dagger e^{iE_{\mathbf{p}}t} \right] \right\} \\
&= \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} \frac{E_{\mathbf{p}}}{2} \left\{ -\left[a_{\mathbf{p}} a_{-\mathbf{p}} e^{-2iE_{\mathbf{p}}t} - a_{-\mathbf{p}}^\dagger a_{-\mathbf{p}} - a_{\mathbf{p}} a_{\mathbf{p}}^\dagger + a_{-\mathbf{p}}^\dagger a_{\mathbf{p}}^\dagger e^{2iE_{\mathbf{p}}t} \right] \right. \\
&\quad \left. + \left[a_{\mathbf{p}} a_{-\mathbf{p}} e^{-2iE_{\mathbf{p}}t} + a_{-\mathbf{p}}^\dagger a_{-\mathbf{p}} + a_{\mathbf{p}} a_{\mathbf{p}}^\dagger + a_{-\mathbf{p}}^\dagger a_{\mathbf{p}}^\dagger e^{2iE_{\mathbf{p}}t} \right] \right\} \\
&= \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} \left[a_{-\mathbf{p}}^\dagger a_{-\mathbf{p}} + a_{\mathbf{p}} a_{\mathbf{p}}^\dagger \right] = \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} (a_{\mathbf{p}}^\dagger a_{\mathbf{p}} + a_{\mathbf{p}} a_{\mathbf{p}}^\dagger)
\end{aligned}$$

□

The second term in the Hamiltonian Eq. 2.7 is the sum of the zero-point energy of all oscillators and is proportional to $(2\pi)^3 \delta^{(3)}(0) \rightarrow V$, thus:

$$E_{\text{vac}} = \frac{V}{2} \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}}$$

This energy shows two divergences: the one coming from the infinite-volume limit (i.e. small momentum), regularized introducing an **infrared cutoff** in the form of a finite volume, and the one from the ultra-relativistic limit (i.e. large momentum), regularized introducing an **ultraviolet cutoff** in the form of a maximum momentum Λ . These divergences are retained in the expression for E_{vac} , as $E_{\text{vac}} \sim V$ and $E_{\text{vac}} \sim \Lambda^4$, but can be ignored (when ignoring gravity) since experiments are only sensitive to energy differences.

Discarding the zero-point energy, the Hamiltonian becomes:

$$H = \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} a_{\mathbf{p}}^\dagger a_{\mathbf{p}} \equiv \mathfrak{N} \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} (a_{\mathbf{p}}^\dagger a_{\mathbf{p}} + a_{\mathbf{p}} a_{\mathbf{p}}^\dagger) \quad (2.9)$$

where the **normal ordering operator** \mathfrak{N} was introduced, which acts by moving all creation operators to the left and all annihilation operators to the right (ex.: $\mathfrak{N}\{a_{\mathbf{p}} a_{\mathbf{p}}^\dagger\} = a_{\mathbf{p}}^\dagger a_{\mathbf{p}}$). It is now straightforward to compute the energy of a generic state in the Fock space, as $a_{\mathbf{p}}^\dagger a_{\mathbf{p}}$ is just a number operator:

$$H |\mathbf{p}_1, \dots, \mathbf{p}_n\rangle = (E_{\mathbf{p}_1} + \dots + E_{\mathbf{p}_n}) |\mathbf{p}_1, \dots, \mathbf{p}_n\rangle \quad (2.10)$$

Computing the spatial momentum from Eq. 1.81 as $P^i = \mathfrak{N} \int d^3x \theta^{0i} = \int d^3x \mathfrak{N}\{\partial_0 \phi \partial^i \phi\}$:

$$P^i = \int \frac{d^3p}{(2\pi)^3} p^i a_{\mathbf{p}}^\dagger a_{\mathbf{p}} \quad (2.11)$$

Therefore, the state $a_{\mathbf{p}}^\dagger |0\rangle$ can correctly be interpreted as a one-particle state with momentum \mathbf{p} , mass m and energy $E_{\mathbf{p}} = \sqrt{\mathbf{p}^2 + m^2}$. The generic state in the Fock space is a multiparticle state, with total energy and momentum the sum of the individual energies and momenta.

Finally, note that creation operators commute between themselves, hence multiparticle states are symmetric under exchange of pairs of particles, i.e. they obey the Bose-Einstein statistics: this agrees with the fact that quanta of a scalar field have no intrinsic spin. i.e. are spin-0 particles.

§2.1.1.1 Ladder operators

Recalling Def. 1.3.3, it is possible to relate the Fourier modes through the KG inner product, defined as:

$$\langle f, g \rangle := i \int d^3x f^*(x) \overleftrightarrow{\partial}_0 g(x) = i \int d^3x [f^*(x) \partial_0 g(x) - \partial_0 f^*(x) g(x)] \equiv i \int d^3x W_0[f^*, g](x) \quad (2.12)$$

where $W_0[f^*, g](x)$ is the *t*-Wrońskian¹ of $f^*(x), g(x)$. Although $W_0[f^*, g](x)$ is time-dependent, by Prop. 1.3.2, if $f^*(x), g(x)$ are solutions of the KG equation, then $\langle f^*, g \rangle$ is time-independent.

Proposition 2.1.3 (Orthonormality of Fourier modes)

Given positive- and negative-frequency Fourier modes:

$$f_{\mathbf{p}}(x) \equiv \frac{e^{-ip_{\mu}x^{\mu}}}{\sqrt{(2\pi)^3 2E_{\mathbf{p}}}} \quad f_{\mathbf{p}}^*(x) \equiv \frac{e^{ip_{\mu}x^{\mu}}}{\sqrt{(2\pi)^3 2E_{\mathbf{p}}}} \quad (2.14)$$

then:

$$\langle f_{\mathbf{p}}, f_{\mathbf{q}} \rangle = -\langle f_{\mathbf{p}}^*, f_{\mathbf{q}}^* \rangle = \delta^{(3)}(\mathbf{p} - \mathbf{q}) \quad \langle f_{\mathbf{p}}, f_{\mathbf{q}}^* \rangle = 0 \quad (2.15)$$

Proof. By direct calculation (using Eq. 2.8):

$$\begin{aligned} \langle f_{\mathbf{p}}, f_{\mathbf{q}} \rangle &= i \int d^3x f_{\mathbf{p}}^*(x) \overleftrightarrow{\partial}_0 f_{\mathbf{q}}(x) = i \int \frac{d^3x}{(2\pi)^3 \sqrt{4E_{\mathbf{p}}E_{\mathbf{q}}}} e^{i(\mathbf{q}-\mathbf{p})\cdot\mathbf{x}} \left(e^{iE_{\mathbf{p}}t} \overleftrightarrow{\partial}_0 e^{-iE_{\mathbf{q}}t} \right) \\ &= i \frac{\delta^{(3)}(\mathbf{q} - \mathbf{p})}{\sqrt{4E_{\mathbf{p}}E_{\mathbf{q}}}} (-iE_{\mathbf{q}} - iE_{\mathbf{p}}) e^{i(E_{\mathbf{p}}-E_{\mathbf{q}})t} = \delta^{(3)}(\mathbf{p} - \mathbf{q}) \end{aligned}$$

as $\delta(x - x_0)f(x) = \delta(x - x_0)f(x_0)$. Other products are equivalent. \square

Using the KG inner product it is possible to explicit the ladder operators.

¹Given $\{f_k(x)\}_{k=1,\dots,N} \subset \mathcal{C}^N(I, \mathbb{C})$, with $I \subset \mathbb{R}^n$ and $N \in \mathbb{N}$, their (*i*th-) Wrońskian is $W_i \in \mathcal{C}(I, \mathbb{C})$ defined as:

$$W[f_1, \dots, f_N](x) := \begin{vmatrix} f_1(x) & f_2(x) & \dots & f_N(x) \\ \partial_i f_1(x) & \partial_i f_2(x) & \dots & \partial_i f_N(x) \\ \vdots & \vdots & \ddots & \vdots \\ \partial_i^{N-1} f_1(x) & \partial_i^{N-1} f_2(x) & \dots & \partial_i^{N-1} f_N(x) \end{vmatrix} \quad (2.13)$$

If $\{f_k(x)\}_{k=1,\dots,N}$ are linearly dependent, so are the columns of the Wrońskian (as ∂_i is a linear operator), hence it vanishes: to show that a set of differentiable functions is linearly independent on a given set, it suffices to show that their Wrońskian does not vanish identically on said set (although it may vanish on a zero-measure subset).

Theorem 2.1.1 (Ladder operators)

The ladder operators for a real scalar field are:

$$\sqrt{2E_{\mathbf{p}}}a_{\mathbf{p}} = i \int d^3x e^{ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \phi(x) \quad \sqrt{2E_{\mathbf{p}}}a_{\mathbf{p}}^{\dagger} = -i \int d^3x e^{-ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \phi(x) \quad (2.16)$$

Proof. Eq. 2.2 can be cast in the form:

$$\phi(x) = \frac{1}{\sqrt{(2\pi)^3}} \int d^3p [f_{\mathbf{p}}(x)a_{\mathbf{p}} + f_{\mathbf{p}}^*(x)a_{\mathbf{p}}^{\dagger}] \implies \langle f_{\mathbf{p}}, \phi \rangle = \frac{a_{\mathbf{p}}}{\sqrt{(2\pi)^3}}$$

by Prop. 2.1.3. Then, inserting Eq. 2.14 yields the thesis ($a_{\mathbf{p}}^{\dagger}$ is analogous). \square

§2.1.2 Complex scalar fields

When considering a complex scalar field, Eq. 1.92 becomes:

$$\phi(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} [a_{\mathbf{p}} e^{-ip_{\mu}x^{\mu}} + b_{\mathbf{p}}^{\dagger} e^{ip_{\mu}x^{\mu}}]_{p^0=E_{\mathbf{p}}} \quad (2.17)$$

$$\phi^{\dagger}(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} [a_{\mathbf{p}}^{\dagger} e^{ip_{\mu}x^{\mu}} + b_{\mathbf{p}} e^{-ip_{\mu}x^{\mu}}]_{p^0=E_{\mathbf{p}}} \quad (2.18)$$

Now there are two independent sets of creation/annihilation operators, which obey the canonical commutation relation:

$$[a_{\mathbf{p}}, a_{\mathbf{q}}^{\dagger}] = [b_{\mathbf{p}}, b_{\mathbf{q}}^{\dagger}] = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q}) \quad (2.19)$$

with all other commutators vanishing. The Fock space is constructed by defining a vacuum state $|0\rangle : a_{\mathbf{p}}|0\rangle = b_{\mathbf{p}}|0\rangle = 0$ and then acting repeatedly with both creation operators. With normal ordering, one finds:

$$H = \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} (a_{\mathbf{p}}^{\dagger} a_{\mathbf{p}} + b_{\mathbf{p}}^{\dagger} b_{\mathbf{p}}) \quad (2.20)$$

$$P^i = \int \frac{d^3p}{(2\pi)^3} p^i (a_{\mathbf{p}}^{\dagger} a_{\mathbf{p}} + b_{\mathbf{p}}^{\dagger} b_{\mathbf{p}}) \quad (2.21)$$

The quanta of a complex scalar field are given by two different species of particles with the same mass.

Proposition 2.1.4 (U(1) charge)

The U(1) charge of the quantized complex scalar field is:

$$Q_{U(1)} = \int \frac{d^3p}{(2\pi)^3} (a_{\mathbf{p}}^{\dagger} a_{\mathbf{p}} - b_{\mathbf{p}}^{\dagger} b_{\mathbf{p}}) \quad (2.22)$$

Proof. By Eq. 1.93:

$$\begin{aligned}
Q_{U(1)} &= i \int d^3x \phi^\dagger \overleftrightarrow{\partial}_0 \phi = i \int d^3x \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_q}} \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_p}} \times \\
&\quad \times \left\{ [a_q^\dagger e^{iq_\mu x^\mu} + b_q e^{-iq_\mu x^\mu}] \partial_0 (a_p e^{-ip_\mu x^\mu} + b_p^\dagger e^{ip_\mu x^\mu}) + \right. \\
&\quad \left. - \partial_0 (a_q^\dagger e^{iq_\mu x^\mu} + b_q e^{-iq_\mu x^\mu}) [a_p e^{-ip_\mu x^\mu} + b_p^\dagger e^{ip_\mu x^\mu}] \right\} \\
&= \int d^3x \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_q}} \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_p}} \times \\
&\quad \times \left\{ E_p [a_q^\dagger e^{iq_\mu x^\mu} + b_q e^{-iq_\mu x^\mu}] [a_p e^{-ip_\mu x^\mu} - b_p^\dagger e^{ip_\mu x^\mu}] + \right. \\
&\quad \left. + E_q [a_q^\dagger e^{iq_\mu x^\mu} - b_q e^{-iq_\mu x^\mu}] [a_p e^{-ip_\mu x^\mu} + b_p^\dagger e^{ip_\mu x^\mu}] \right\} \\
&= \int d^3x \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_q}} \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_p}} \times \\
&\quad \times \left\{ E_p [a_q^\dagger e^{iE_q t} + b_{-q} e^{-iE_q t}] [a_p e^{-iE_p t} - b_{-p}^\dagger e^{iE_p t}] + \right. \\
&\quad \left. + E_q [a_q^\dagger e^{iE_q t} - b_{-q} e^{-iE_q t}] [a_p e^{-iE_p t} + b_{-p}^\dagger e^{iE_p t}] \right\} e^{i(\mathbf{p}-\mathbf{q}) \cdot \mathbf{x}} \\
&= \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} \left\{ [a_p^\dagger e^{iE_p t} + b_{-p} e^{-iE_p t}] [a_p e^{-iE_p t} - b_{-p}^\dagger e^{iE_p t}] + \right. \\
&\quad \left. + [a_p^\dagger e^{iE_p t} - b_{-p} e^{-iE_p t}] [a_p e^{-iE_p t} + b_{-p}^\dagger e^{iE_p t}] \right\} \\
&= \int \frac{d^3p}{(2\pi)^3} [a_p^\dagger a_p - b_{-p} b_{-p}^\dagger] = \int \frac{d^3p}{(2\pi)^3} (a_p^\dagger a_p - b_p b_p^\dagger)
\end{aligned}$$

Applying normal ordering yields the thesis. \square

While normal ordering was justified when considering the Hamiltonian on the grounds that the vacuum energy is unobservable, a charged vacuum would have observable effects; however when promoting ϕ to a quantum operator, the expression $\phi^\dagger \overleftrightarrow{\partial}_0 \phi$ presents an ordering ambiguity (ex.: $\phi^\dagger(\partial_0 \phi)$ or $(\partial_0 \phi)\phi^\dagger$), which is removed requiring the charge of the vacuum to vanish.

Being $a_p^\dagger a_p$ and $b_p^\dagger b_p$ number operators, the $U(1)$ charge is equal to the number of quanta created by a_p^\dagger minus the number of quanta created by b_p^\dagger , integrated over all momenta: in particular, $a_p^\dagger |0\rangle$ and $b_p^\dagger |0\rangle$ are both spin-zero particles of mass m and momentum \mathbf{p} , but they respectively have charges $Q_{U(1)} = +1$ and $Q_{U(1)} = -1$. This allows to properly interpret the negative-energy solutions of the KG equations: they are positive-energy particles with opposite $U(1)$ charge and are called **antiparticles**.

For a real scalar field, the reality condition reads $a_p = b_p$, thus it describes a field whose particle is its own antiparticle, and it is symmetric under any $U(1)$ symmetry.

§2.2 Spinor fields

A principle of QFT is the **spin-statistic theorem**: integer-spin fields are to be quantized imposing equal-time commutation relations, while half-integer-spin with equal-time anticommutation relations.

§2.2.1 Dirac fields

From the Dirac Lagrangian Eq. 1.101, the conjugate momentum to the Dirac field Ψ is computed as:

$$\Pi_\Psi = i\bar{\Psi}\gamma^0 = i\Psi^\dagger \quad (2.23)$$

Imposing the canonical anticommutation relation, according to the spin-statistic theorem:

$$\{\Psi_a(t, \mathbf{x}), \Psi_b^\dagger(t, \mathbf{y})\} = \delta^{(3)}(\mathbf{x} - \mathbf{y})\delta_{ab} \quad (2.24)$$

where $a, b = 1, 2, 3, 4$ are Dirac indices. Expanding the free Dirac field in plane waves:

$$\Psi(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [a_{\mathbf{p},s} u^s(p) e^{-ip_\mu x^\mu} + b_{\mathbf{p},s}^\dagger v^s(p) e^{ip_\mu x^\mu}]_{p^0=E_{\mathbf{p}}} \quad (2.25)$$

$$\bar{\Psi}(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [a_{\mathbf{p},s}^\dagger \bar{u}^s(p) e^{ip_\mu x^\mu} + b_{\mathbf{p},s} \bar{v}^s(p) e^{-ip_\mu x^\mu}]_{p^0=E_{\mathbf{p}}} \quad (2.26)$$

where the spinor wave functions $u^s(p), v^s(p)$ are given by Eq. 1.105-1.106. Translating Eq. 2.24 in terms of creation/annihilation operators:

$$\{a_{\mathbf{p},s}, a_{\mathbf{q},r}^\dagger\} = \{b_{\mathbf{p},s}, b_{\mathbf{q},r}^\dagger\} = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q}) \delta_{sr} \quad (2.27)$$

The Fock space is again constructed defining a vacuum state $|0\rangle : a_{\mathbf{p},s} |0\rangle = b_{\mathbf{p},s} |0\rangle = 0$ and then acting repeatedly on it with $a_{\mathbf{p},s}^\dagger, b_{\mathbf{p},s}^\dagger$:

$$\begin{aligned} & |(\mathbf{p}_1, s_1), \dots, (\mathbf{p}_n, s_n); (\mathbf{q}_1, r_1), \dots, (\mathbf{q}_m, r_m)\rangle \\ &= \sqrt{2E_{\mathbf{p}_1}} \dots \sqrt{2E_{\mathbf{p}_n}} \sqrt{2E_{\mathbf{q}_1}} \dots \sqrt{2E_{\mathbf{q}_m}} a_{\mathbf{p}_1, s_1}^\dagger \dots a_{\mathbf{p}_n, s_n}^\dagger b_{\mathbf{q}_1, r_1}^\dagger \dots b_{\mathbf{q}_m, r_m}^\dagger |0\rangle \end{aligned}$$

As these operators anticommute, states in this Fock space are antisymmetric under the exchange of particles, therefore spin- $\frac{1}{2}$ obey the Fermi-Dirac statistics (as of the spin-statistic theorem).

Proposition 2.2.1 (Dirac Hamiltonian)

The Hamiltoniana for a Dirac field Ψ is:

$$H = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} E_{\mathbf{p}} (a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} + b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s}) \quad (2.28)$$

Proof. By Eq. 2.23, the Hamiltonian density is:

$$\mathcal{H} = \Pi_\Psi \partial_0 \Psi - \mathcal{L}_D = i\Psi^\dagger \partial_0 \Psi - \bar{\Psi} (i\gamma^0 \partial_0 + i\gamma^i \partial_i - m) \Psi = \bar{\Psi} (-i\gamma^i \partial_i + m) \Psi$$

Therefore, using Eq. 2.25-2.26 and Eq. 2.8:

$$\begin{aligned}
 H &= \int d^3x \bar{\Psi} (-i\gamma^i \partial_i + m) \Psi = \int d^3x \bar{\Psi} (-i\boldsymbol{\gamma} \cdot \boldsymbol{\nabla} + m) \Psi \\
 &= \mathfrak{N} \int d^3x \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger \bar{u}^s(p) e^{ip_\mu x^\mu} + b_{\mathbf{p},s} \bar{v}^s(p) e^{-ip_\mu x^\mu}] \times \\
 &\quad \times (-i\boldsymbol{\gamma} \cdot \boldsymbol{\nabla} + m) [a_{\mathbf{q},r} u^r(q) e^{-iq_\mu x^\mu} + b_{\mathbf{q},r}^\dagger v^r(q) e^{iq_\mu x^\mu}] \\
 &= \mathfrak{N} \int d^3x \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger \bar{u}^s(p) e^{ip_\mu x^\mu} + b_{\mathbf{p},s} \bar{v}^s(p) e^{-ip_\mu x^\mu}] \times \\
 &\quad \times [(\boldsymbol{\gamma} \cdot \mathbf{q} + m) a_{\mathbf{q},r} u^r(q) e^{-iq_\mu x^\mu} + (-\boldsymbol{\gamma} \cdot \mathbf{q} + m) b_{\mathbf{q},r}^\dagger v^r(q) e^{iq_\mu x^\mu}]
 \end{aligned}$$

Using the Dirac equation in the form $(\not{p} - m)u(p) = (\not{p} + m)v(p) = 0$:

$$(\boldsymbol{\gamma} \cdot \mathbf{q} + m) u(q) = \gamma^0 E_{\mathbf{q}} u(q) \quad (-\boldsymbol{\gamma} \cdot \mathbf{q} + m) v(q) = -\gamma^0 E_{\mathbf{q}} v(q)$$

Therefore, omitting the constraint $p^0 = E_{\mathbf{p}}$ in the spinors' arguments and using Eq. 2.8:

$$\begin{aligned}
 H &= \mathfrak{N} \int d^3x \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger \bar{u}^s(\mathbf{p}) e^{iE_{\mathbf{p}}t} + b_{-\mathbf{p},s} \bar{v}^s(-\mathbf{p}) e^{-iE_{\mathbf{p}}t}] \times \\
 &\quad \times \gamma^0 E_{\mathbf{q}} [a_{\mathbf{q},r} u^r(\mathbf{q}) e^{-iE_{\mathbf{q}}t} - b_{-\mathbf{q},r}^\dagger v^r(-\mathbf{q}) e^{iE_{\mathbf{q}}t}] e^{i(\mathbf{q}-\mathbf{p}) \cdot \mathbf{x}} \\
 &= \mathfrak{N} \int \frac{d^3p}{2(2\pi)^3} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger \bar{u}^s(\mathbf{p}) e^{iE_{\mathbf{p}}t} + b_{-\mathbf{p},s} \bar{v}^s(-\mathbf{p}) e^{-iE_{\mathbf{p}}t}] \times \\
 &\quad \times \gamma^0 [a_{\mathbf{p},r} u^r(\mathbf{p}) e^{-iE_{\mathbf{p}}t} - b_{-\mathbf{p},s}^\dagger v^r(-\mathbf{p}) e^{iE_{\mathbf{p}}t}] \\
 &= \mathfrak{N} \int \frac{d^3p}{2(2\pi)^3} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger a_{\mathbf{p},r} u^{s\dagger}(\mathbf{p}) u^r(\mathbf{p}) + b_{-\mathbf{p},s} a_{\mathbf{p},r} v^{s\dagger}(-\mathbf{p}) u^r(\mathbf{p}) e^{-2iE_{\mathbf{p}}t} + \\
 &\quad - a_{\mathbf{p},s}^\dagger b_{-\mathbf{p},r}^\dagger u^{s\dagger}(\mathbf{p}) v^r(-\mathbf{p}) e^{2iE_{\mathbf{p}}t} - b_{-\mathbf{p},s} b_{-\mathbf{p},r}^\dagger v^{s\dagger}(-\mathbf{p}) v^r(-\mathbf{p})]
 \end{aligned}$$

Lemma 2.2.1

$$u^{s\dagger}(\mathbf{p}) v^r(-\mathbf{p}) = v^{s\dagger}(-\mathbf{p}) u^r(\mathbf{p}) = 0$$

Using Lemma 2.2.1, Eq. 1.108 and the antisymmetry $\mathfrak{N}\{b_{\mathbf{p},s} b_{\mathbf{p},s}^\dagger\} = -b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s}$:

$$H = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} E_{\mathbf{p}} \mathfrak{N}\{a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} - b_{\mathbf{p},s} b_{\mathbf{p},s}^\dagger\} = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} E_{\mathbf{p}} (a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} + b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s})$$

which completes the proof. \square

It can be seen that, using anticommutators, the Hamiltonian and its interpretation are analogous to that of the complex scalar field: if commutators were used, instead, one would get a final $-b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s}$ term, which is problematic as it yields an energy unbounded from below².

The momentum operator too is analogous to that of the complex scalar field, with the additional

²The spin-statistics theorem is implied by Lorentz invariance, positive energies and causality (see [1]).

spin degree of freedom.

Proposition 2.2.2 (Momentum operator)

The momentum operator of a Dirac field Ψ is:

$$\mathbf{P} = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} \mathbf{p} (a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} + b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s}) \quad (2.29)$$

Proof. By Eq. 1.80, the 0i-component of energy-momentum tensor of the Dirac Lagrangian is:

$$\theta^{0i} = \frac{\partial \mathcal{L}_D}{\partial(\partial_0 \Psi)} \partial^i \Psi = \bar{\Psi} i \gamma^0 \partial^i \Psi = -\Psi^\dagger i \partial_i \Psi$$

Thus, according to Eq. 1.81:

$$\begin{aligned} \mathbf{P} &= \mathfrak{N} \int d^3x \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger u^{s\dagger}(p) e^{ip_\mu x^\mu} + b_{\mathbf{p},s} v^{s\dagger}(p) e^{-ip_\mu x^\mu}] \times \\ &\quad \times (-i\nabla) [a_{\mathbf{q},r} u^r(q) e^{-iq_\mu x^\mu} + b_{\mathbf{q},r}^\dagger v^r(q) e^{iq_\mu x^\mu}] \\ &= \mathfrak{N} \int d^3x \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \int \frac{d^3q}{(2\pi)^3 \sqrt{2E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger u^{s\dagger}(p) e^{ip_\mu x^\mu} + b_{\mathbf{p},s} v^{s\dagger}(p) e^{-ip_\mu x^\mu}] \times \\ &\quad \times \mathbf{q} [a_{\mathbf{q},r} u^r(q) e^{-iq_\mu x^\mu} - b_{\mathbf{q},r}^\dagger v^r(q) e^{iq_\mu x^\mu}] \\ &= \mathfrak{N} \int \frac{d^3p}{(2\pi)^3} \frac{\mathbf{p}}{2E_{\mathbf{p}}} \sum_{s=1,2} \sum_{r=1,2} [a_{\mathbf{p},s}^\dagger u^{s\dagger}(\mathbf{p}) e^{iE_{\mathbf{p}}t} + b_{-\mathbf{p},s} v^{s\dagger}(-\mathbf{p}) e^{-iE_{\mathbf{p}}t}] \times \\ &\quad \times [a_{\mathbf{p},r} u^r(\mathbf{p}) e^{-iE_{\mathbf{p}}t} - b_{-\mathbf{p},r}^\dagger v^r(-\mathbf{p}) e^{iE_{\mathbf{p}}t}] \end{aligned}$$

Using Lemma 2.2.1, Eq. 1.108 and the antisymmetry $\mathfrak{N} b_{\mathbf{p},s} b_{\mathbf{p},s}^\dagger = -b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s}$:

$$\mathbf{P} = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} \mathbf{p} \mathfrak{N} \{a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} - b_{\mathbf{p},s} b_{\mathbf{p},s}^\dagger\} = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} \mathbf{p} (a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} + b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s})$$

which is the thesis. \square

Thus, both $a_{\mathbf{p},s}^\dagger$ and $b_{\mathbf{p},s}^\dagger$ create particles with energy $+E_{\mathbf{p}}$ and momentum \mathbf{p} : these are respectively called **fermions** and **antifermions**.

§2.2.1.1 Quantum numbers

Under a generic Lorentz transformation, the Dirac field Ψ transforms according to Eq. 1.60, i.e. $\Psi'(x) = e^{-\frac{i}{2} \omega_{\mu\nu} J^{\mu\nu}} \Psi(x)$, where the Lorentz generator $J^{\mu\nu}$ is defined in Eq. 1.61.

Proposition 2.2.3 (Angular momentum)

The conserved charge associated to the infinitesimal rotation Eq. 1.60 is:

$$\mathbf{J} = \int d^3x \Psi^\dagger (\mathbf{x} \times (-i\nabla) + \frac{1}{2} \boldsymbol{\Sigma}) \Psi \quad (2.30)$$

Proof. Consider a rotation of θ around the z -axis: it is described by $\omega_{12} = -\omega_{21} = \theta$, so, by Eq. 1.60:

$$\delta_0 \Psi = \theta (x^1 \partial^2 - x^2 \partial^1 - \tfrac{i}{2} \Sigma^3) \Psi = -\theta (x^1 \partial_2 - x^2 \partial_1 + \tfrac{i}{2} \Sigma^3) \Psi \equiv \theta \Delta \Psi$$

Using the same notation of Def. 1.3.1, $\epsilon \equiv \theta$ and $F = \Delta \Psi$, therefore, by Eq. 1.78, the temporal component of the conserved Noether current is (without negative sign, as it is mathematically equivalent):

$$j^0 = \frac{\partial \mathcal{L}_D}{\partial(\partial_0 \Psi)} \Delta \Psi = -i \Psi^\dagger ((\mathbf{x} \times \nabla)^3 + \tfrac{i}{2} \Sigma^3) \Psi$$

As the associated Noether charge is $J^3 = \int d^3x j^0$, this can be generalized to rotations around the x -axis and the y -axis, yielding:

$$\mathbf{J} = \int d^3x \Psi^\dagger (\mathbf{x} \times (-i\nabla) + \tfrac{1}{2} \Sigma) \Psi$$

which is the thesis. \square

For non-relativistic fermions, the first term gives the orbital angular momentum, while the second term gives the spin angular momentum. For relativistic fermions, this division is not straightforward.

To determine the spin of fermions, it is sufficient to consider them at rest. The spin along the z -axis is given by the S_z operator (at $t = 0$):

$$\begin{aligned} S_z &= \int d^3x \int \int \frac{d^3p d^3q}{(2\pi)^6 \sqrt{4E_{\mathbf{p}} E_{\mathbf{q}}}} e^{i(\mathbf{q}-\mathbf{p}) \cdot \mathbf{x}} \times \\ &\quad \times \sum_{t=1,2} \sum_{r=1,2} \left[a_{\mathbf{p},t}^\dagger u^{t\dagger}(\mathbf{p}) + b_{-\mathbf{p},t} v^{t\dagger}(-\mathbf{p}) \right] \frac{\Sigma^3}{2} \left[a_{\mathbf{q},r} u^r(\mathbf{q}) + b_{-\mathbf{q},r}^{\dagger} v^r(-\mathbf{q}) \right] \\ &= \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} \sum_{t=1,2} \sum_{r=1,2} \left[a_{\mathbf{p},t}^\dagger u^{t\dagger}(\mathbf{p}) + b_{-\mathbf{p},t} v^{t\dagger}(-\mathbf{p}) \right] \frac{\Sigma^3}{2} \left[a_{\mathbf{p},r} u^r(\mathbf{p}) + b_{-\mathbf{p},r}^{\dagger} v^r(-\mathbf{p}) \right] \end{aligned}$$

Noting that $S_z |0\rangle = 0$ (by definition of vacuum), then $S_z a_{\mathbf{0},s}^\dagger |0\rangle = \{S_z, a_{\mathbf{0},s}^\dagger\} |0\rangle$; the only non-zero terms are those proportional to $u^\dagger \Sigma^3 u$ and $v^\dagger \Sigma^3 v$, but the latter vanishes as $\{a, b\} = 0$, so the only term remaining is:

$$\{a_{\mathbf{p},t}^\dagger a_{\mathbf{p},r}, a_{\mathbf{0},s}^\dagger\} = a_{\mathbf{p},t}^\dagger \{a_{\mathbf{p},r}, a_{\mathbf{0},s}^\dagger\} = (2\pi)^3 \delta^{(3)}(\mathbf{p}) \delta_{rs} a_{\mathbf{p},t}^\dagger$$

The S_z operator thus acts as (recall Eq. 1.105 with $p^\mu = (m, 0, 0, 0)$):

$$S_z a_{\mathbf{0},s}^\dagger |0\rangle = \frac{1}{2E_{\mathbf{p}}} \sum_{t=1,2} u^{t\dagger}(\mathbf{0}) \frac{\Sigma^3}{2} u^s(\mathbf{0}) a_{\mathbf{0},t}^\dagger |0\rangle = \frac{1}{4} \left(2\xi^{1\dagger} \sigma^3 \xi^s a_{\mathbf{0},1}^\dagger + 2\xi^{2\dagger} \sigma^3 \xi^s a_{\mathbf{0},2}^\dagger \right) |0\rangle$$

Using $\sigma^3 \xi^1 = +\xi^1$ and $\sigma^3 \xi^2 = -\xi^2$, by $\xi^{r\dagger} \xi^s = \delta^{rs}$ one gets:

$$S_z a_{\mathbf{0},1}^\dagger |0\rangle = +\frac{1}{2} a_{\mathbf{0},1}^\dagger |0\rangle \quad S_z a_{\mathbf{0},2}^\dagger |0\rangle = -\frac{1}{2} a_{\mathbf{0},2}^\dagger |0\rangle$$

which means that fermions are spin- $\frac{1}{2}$ particles, with $a_{\mathbf{0},1}^\dagger$ creating $s = +\frac{1}{2}$ fermions and $a_{\mathbf{0},2}^\dagger$ creating $s = -\frac{1}{2}$ fermions. Conversely, it is equivalent to show that $b_{\mathbf{0},1}^\dagger$ creates $s = -\frac{1}{2}$

antifermions and $b_{0,2}^\dagger$ creates $s = +\frac{1}{2}$ antifermions (as b and b^\dagger are not in normal order in S_z , so there's an extra negative sign due to antisymmetry).

Another important conserved Noether charge of Dirac theory is that associated to the vector current $j_V^\mu = \bar{\Psi}\gamma^\mu\Psi$ (recall Eq. 1.112), which is (using $j_V^0 = \Psi^\dagger\Psi$):

$$Q_{U(1)_V} = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1,2} (a_{\mathbf{p},s}^\dagger a_{\mathbf{p},s} - b_{\mathbf{p},s}^\dagger b_{\mathbf{p},s}) \quad (2.31)$$

This means that fermions have $U(1)_V$ charge of $+1$, while antifermions of -1 .

state	S_z	$Q_{U(1)_V}$
$a_{\mathbf{p},1}^\dagger 0\rangle$	$+\frac{1}{2}$	$+1$
$a_{\mathbf{p},2}^\dagger 0\rangle$	$-\frac{1}{2}$	$+1$
$b_{\mathbf{p},1}^\dagger 0\rangle$	$-\frac{1}{2}$	-1
$b_{\mathbf{p},2}^\dagger 0\rangle$	$+\frac{1}{2}$	-1

Table 2.1: Quantum numbers for fermions in the Dirac theory.

§2.2.1.2 Ladder operators

As for the KG field, the ladder operators of the Dirac field can be made explicit too.

Proposition 2.2.4 (Ladder operators)

The ladder operators for a Dirac field are:

$$\sqrt{2E_{\mathbf{p}}} a_{\mathbf{p},s} = \int d^3x e^{ip_\mu x^\mu} \bar{u}^s(p) \gamma^0 \Psi(x) \quad \sqrt{2E_{\mathbf{p}}} a_{\mathbf{p},s}^\dagger = \int d^3x e^{-ip_\mu x^\mu} \bar{\Psi}(x) \gamma^0 u^s(p) \quad (2.32)$$

$$\sqrt{2E_{\mathbf{p}}} b_{\mathbf{p},s}^\dagger = \int d^3x e^{-ip_\mu x^\mu} \bar{v}^s(p) \gamma^0 \Psi(x) \quad \sqrt{2E_{\mathbf{p}}} b_{\mathbf{p},s} = \int d^3x e^{ip_\mu x^\mu} \bar{\Psi}(x) \gamma^0 v^s(p) \quad (2.33)$$

Proof. First, cast Eq. 2.25 in the form:

$$\Psi(x) = \int \frac{d^3k}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} \sum_{s=1,2} \left[a_{\mathbf{k},s} u^s(\mathbf{k}) e^{-iE_{\mathbf{k}}t} + b_{-\mathbf{k},s}^\dagger v^s(-\mathbf{k}) e^{iE_{\mathbf{k}}t} \right] e^{i\mathbf{k}\cdot\mathbf{x}}$$

Then:

$$\begin{aligned} \int d^3x e^{ip_\mu x^\mu} \Psi(x) &= \int \frac{d^3x d^3k}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} \sum_{s=1,2} \left[a_{\mathbf{k},s} u^s(\mathbf{k}) e^{i(E_{\mathbf{p}}-E_{\mathbf{k}})t} + b_{-\mathbf{k},s}^\dagger v^s(-\mathbf{k}) e^{i(E_{\mathbf{p}}+E_{\mathbf{k}})t} \right] e^{i(\mathbf{k}-\mathbf{p})\cdot\mathbf{x}} \\ &= \int \frac{d^3k}{\sqrt{2E_{\mathbf{k}}}} \delta^{(3)}(\mathbf{k}-\mathbf{p}) \sum_{s=1,2} \left[a_{\mathbf{k},s} u^s(\mathbf{k}) e^{i(E_{\mathbf{p}}-E_{\mathbf{k}})t} + b_{-\mathbf{k},s}^\dagger v^s(-\mathbf{k}) e^{i(E_{\mathbf{p}}+E_{\mathbf{k}})t} \right] \\ &= \frac{1}{\sqrt{2E_{\mathbf{p}}}} \sum_{r=1,2} \left[a_{\mathbf{p},r} u^r(\mathbf{p}) + b_{-\mathbf{p},r}^\dagger v^r(-\mathbf{p}) e^{2iE_{\mathbf{p}}t} \right] \end{aligned}$$

Lemma 2.2.2

$$\bar{u}^s(\mathbf{p})\gamma^0 v^r(-\mathbf{p}) = 0 \quad \bar{v}^s(\mathbf{p})\gamma^0 u^r(-\mathbf{p}) = 0 \quad (2.34)$$

Using Lemma 2.2.2 and Eq. 1.110:

$$\int d^3x e^{ip_\mu x^\mu} \bar{u}^s(p)\gamma^0 \Psi(x) = \sqrt{2E_{\mathbf{p}}} a_{\mathbf{p},s}$$

Analogously:

$$\int d^3x e^{-ip_\mu x^\mu} \bar{v}^s(p)\gamma^0 \Psi(x) = \sqrt{2E_{\mathbf{p}}} b_{\mathbf{p},s}^\dagger$$

To find the remaining ladder operators, either repeat the previous steps with $\bar{\Psi}(x)$ or use (as $(\gamma^\mu)^2 = \mathbb{I}_4$):

$$(\bar{u}^s \gamma^0 \Psi)^\dagger = (u^{s\dagger} \Psi)^\dagger = \Psi^\dagger u^s = \bar{\Psi} \gamma^0 u^s \quad (2.35)$$

This yields the thesis. \square

§2.2.2 Massless Weyl fields

The quantization of massless Weyl fields follows immediately from that of Dirac fields, and its useful to use Dirac notation:

$$\Psi_L \equiv \begin{pmatrix} \psi_L \\ 0 \end{pmatrix} \quad \Psi_R \equiv \begin{pmatrix} 0 \\ \psi_R \end{pmatrix}$$

Consider Ψ_L . As for Eq. 2.25:

$$\Psi_L(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [a_{\mathbf{p},s} u_L^s(p) e^{-ip_\mu x^\mu} + b_{\mathbf{p},s}^\dagger v_L^s(p) e^{ip_\mu x^\mu}]_{p^0=E_{\mathbf{p}}}$$

where the Dirac spinors now have the right-hand component vanishing, in the chiral representation. By Eq. 1.104-1.106, in the massless (ultra-relativistic) limit $p^\mu = (E, 0, 0, E)$, so spinors with $s = 1$ have only the right-handed component, while those with $s = 2$ only the left-handed one: therefore, only $s = 2$ spinors contribute to Ψ_L (and only $s = 1$ to Ψ_R).

By Tab. 2.1, it is clear that in this context $a_{\mathbf{p},2}^\dagger$ creates a particle with helicity $h = -\frac{1}{2}$, while $b_{\mathbf{p},2}^\dagger$ creates an antiparticle with $h = +\frac{1}{2}$: in general, a left-handed massless Weyl field describes particles with $h = -\frac{1}{2}$ and antiparticles with $h = +\frac{1}{2}$, while a right-handed one describes particles with $h = +\frac{1}{2}$ and antiparticles with $h = -\frac{1}{2}$.

§2.2.3 Discrete symmetries of fermionic fields

§2.2.3.1 Parity

Under parity $\mathcal{P}\mathbf{p} = -\mathbf{p}$ and $\mathcal{P}s = s$, so, for a generic particle of type a :

$$\mathcal{P} |a; \mathbf{p}, s\rangle = \eta_a |a; -\mathbf{p}, s\rangle \quad (2.36)$$

where η_a is a generic constant phase factor, since states in the Fock space which differ by a phase still represent the same physical state. As $\mathcal{P}^2 = \mathbb{1}$, this means that $\eta_a^2 = \pm 1$, as observables are built from an even number of fermionic ladder operators.

Non-Majorana fermions It is possible to prove that, for non-Majorana spin- $\frac{1}{2}$ fermions, it is possible to redefine \mathcal{P} so that $\eta_a = +1$, i.e. $\eta_a = \pm 1$.

Proposition 2.2.5

$$\mathcal{P}a_{\mathbf{p},s}^\dagger \mathcal{P} = \eta_a a_{-\mathbf{p},s}^\dagger \quad \mathcal{P}b_{\mathbf{p},s}^\dagger \mathcal{P} = \eta_b b_{-\mathbf{p},s}^\dagger \quad (2.37)$$

Proof. For a multiparticle state one must have:

$$\mathcal{P}a_{\mathbf{p},s}^\dagger b_{\mathbf{q},r}^\dagger |0\rangle = \eta_a \eta_b a_{-\mathbf{p},s}^\dagger b_{-\mathbf{q},r}^\dagger |0\rangle$$

Therefore:

$$\mathcal{P}a_{\mathbf{p},s}^\dagger = \eta_a a_{-\mathbf{p},s}^\dagger \mathcal{P} \quad \mathcal{P}b_{\mathbf{p},s}^\dagger = \eta_b b_{-\mathbf{p},s}^\dagger \mathcal{P}$$

Using $\mathcal{P}^2 = 1$ yields the thesis. \square

According to Wigner's theorem (Th. 1.2.2), this symmetry can be represented by a unitary operator, i.e. $\mathcal{P}^\dagger = \mathcal{P}^{-1}$, but $\mathcal{P}^2 = 1$, therefore $\mathcal{P}^\dagger = \mathcal{P}$. Then, Eq. 2.37 is valid for $a_{\mathbf{p},s}, b_{\mathbf{p},s}$ too, thus parity acts as:

$$\Psi(x) \mapsto \Psi'(x') = \mathcal{P}\Psi(x)\mathcal{P} \quad (2.38)$$

Explicitly:

$$\begin{aligned} \mathcal{P}\Psi(x)\mathcal{P} &= \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} \left[\eta_a a_{-\mathbf{p},s} u^s(p) e^{-ip_\mu x^\mu} + \eta_b b_{-\mathbf{p},s}^\dagger v^s(p) e^{ip_\mu x^\mu} \right]_{p^0=E_{\mathbf{p}}} \\ &= \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} \left[\eta_a a_{-\mathbf{p},s} u^s(p) e^{-iE_{\mathbf{p}}t + i\mathbf{p}\cdot\mathbf{x}} + \eta_b b_{-\mathbf{p},s}^\dagger v^s(p) e^{iE_{\mathbf{p}}t - i\mathbf{p}\cdot\mathbf{x}} \right]_{p^0=E_{\mathbf{p}}} \end{aligned}$$

Setting $\mathbf{p}' \equiv -\mathbf{p}, \mathbf{x}' \equiv -\mathbf{x}$ and noting that, by Eq. 1.104-1.106 $\mathbf{p} \mapsto -\mathbf{p}$ exchanges the left-handed and the right-handed components of the spinors, i.e. $u^s(p) \mapsto \gamma^0 u^s(p)$ and $v^s(p) \mapsto -\gamma^0 v^s(p)$:

$$\begin{aligned} \mathcal{P}\Psi(x)\mathcal{P} &= \int \frac{d^3p'}{(2\pi)^3 \sqrt{E_{\mathbf{p}'}}} \sum_{s=1,2} \left[\eta_a a_{\mathbf{p}',s} \gamma^0 u^s(p') e^{-iE_{\mathbf{p}'}t + i\mathbf{p}'\cdot\mathbf{x}'} - \eta_b b_{\mathbf{p}',s}^\dagger \gamma^0 v^s(p') e^{iE_{\mathbf{p}'}t - i\mathbf{p}'\cdot\mathbf{x}'} \right]_{p'^0=E_{\mathbf{p}'}} \\ &= \gamma^0 \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} \left[\eta_a a_{\mathbf{p},s} u^s(p) e^{-ip_\mu x^\mu} - \eta_b b_{\mathbf{p},s}^\dagger v^s(p) e^{ip_\mu x^\mu} \right]_{p^0=E_{\mathbf{p}}} \end{aligned}$$

Requiring that Ψ is a representation of parity, up to a phase, means $\eta_a = -\eta_b$, so that:

$$\mathcal{P}\Psi(t, \mathbf{x})\mathcal{P} = \eta_a \gamma^0 \Psi(t, -\mathbf{x}) \quad (2.39)$$

This, in the chiral representation, agrees with the fact that parity exchanges left-handed and right-handed Weyl spinors. The η_a factor cancels in any fermion bilinear involving only one type of particles; however, the relative phase factors of different particles can be observed: in particular, the opposite intrinsic parity of fermions and antifermions.

Spin-0 bosons As already noted, the scalar complex field is similar to the Dirac field, apart for the absence of spinors in the expansions of $\phi(x)$: in particular, this means that scalar fields have no relative negative signe between η_a and η_b , so a quantized complex scalar field gives a representation of parity if $\eta_a = \eta_b$, and the intrinsic parity of spin-0 particle and antiparticles is equal.

§2.2.3.2 Charge conjugation

Recall Eq. 1.46: charge conjugation acts on the classical Dirac field as $\Psi \mapsto -i\gamma^2\Psi^*$.

Definition 2.2.1 (Quantized charge conjugation)

The **charge conjugation operator** is defined as:

$$\mathcal{C}a_{\mathbf{p},s}\mathcal{C} = \eta_c b_{\mathbf{p},s} \quad \mathcal{C}b_{\mathbf{p},s}\mathcal{C} = \eta_c a_{\mathbf{p},s} \quad (2.40)$$

with $\eta_c = \pm 1$ for simplicity.

Thus, $\mathcal{C}^2 = \mathbb{1}$ too, and its physical interpretation is the exchange of particles and antiparticles, while leaving \mathbf{p} and s unchanged: as $a_{\mathbf{p},s}$ and $b_{\mathbf{p},s}$ create particles with opposite spin, this means that charge conjugation reverses the helicity of particles.

Lemma 2.2.3 (Charge conjugation on spinors)

$$\mathcal{C}u^s(p)\mathcal{C} = -i\gamma^2[v^s(p)]^* \quad \mathcal{C}v^s(p)\mathcal{C} = -i\gamma^2[u^s(p)]^* \quad (2.41)$$

Proposition 2.2.6 (Charge conjugation on Dirac fields)

Given a Dirac field $\Psi(x)$:

$$\mathcal{C}\Psi(x)\mathcal{C} = -i\eta_c\gamma^2[\Psi(x)]^* \quad (2.42)$$

Proof. Using Eq. 2.40-2.41:

$$\begin{aligned} \mathcal{C}\Psi(x)\mathcal{C} &= \eta_c \int \frac{d^3p}{(2\pi)^3\sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [b_{\mathbf{p},s}\mathcal{C}u^s(p)\mathcal{C}e^{-ip_{\mu}x^{\mu}} + a_{\mathbf{p},s}^{\dagger}\mathcal{C}v^s(p)\mathcal{C}e^{ip_{\mu}x^{\mu}}] \\ &= -i\eta_c\gamma^2 \int \frac{d^3p}{(2\pi)^3\sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [b_{\mathbf{p},s}[v^s(p)]^*e^{-ip_{\mu}x^{\mu}} + a_{\mathbf{p},s}^{\dagger}[u^s(p)]^*e^{ip_{\mu}x^{\mu}}] \\ &= -i\eta_c\gamma^2[\Psi(x)]^* \end{aligned}$$

which is the thesis. \square

As for parity, the transformation of quantized fields is analogous to that of classical fields, with an additional quantum phase factor which depends on the particle type.

Proposition 2.2.7 (Charge conjugation of the vector current)

$$\mathcal{C}(\bar{\Psi}\gamma^{\mu}\Psi)\mathcal{C} = -\bar{\Psi}\gamma^{\mu}\Psi \quad (2.43)$$

§2.2.3.3 Time reversal

Theorem 2.2.1 (Anti-unitary time reversal)

Time reversal cannot be implemented as a linear unitary operator.

Proof. Assume that the time reversal operator \mathcal{T} is linear and unitary; then, as it must be a symmetry of the free Dirac Lagrangian, $[\mathcal{T}, H] = 0$, so:

$$\begin{aligned}\mathcal{T}\Psi(t, \mathbf{x})\mathcal{T}|0\rangle &= \mathcal{T}e^{iHt}\Psi(\mathbf{x})e^{-iHt}\mathcal{T}|0\rangle = e^{iHt}\mathcal{T}\Psi(\mathbf{x})\mathcal{T}e^{-iHt}|0\rangle = e^{iHt}\mathcal{T}\Psi(\mathbf{x})\mathcal{T}|0\rangle \\ &= \Psi(-t, \mathbf{x})|0\rangle = e^{-iHt}\Psi(\mathbf{x})e^{iHt}|0\rangle = e^{-iHt}\Psi(\mathbf{x})|0\rangle\end{aligned}$$

assuming $H|0\rangle = 0$. But the first line is a sum of negative frequencies only, while the second line is a sum of positive frequencies only, which is absurd. \square

Time reversal is an example of operator which, according to Wigner's theorem (Th. 1.2.2), is represented as an anti-linear anti-unitary operator, i.e. an operator such that:

$$\langle a|\mathcal{T}^\dagger\mathcal{T}|b\rangle = \langle a|b\rangle^* \quad \mathcal{T}\lambda|a\rangle = \lambda^*\mathcal{T}|a\rangle \quad \forall |a\rangle, |b\rangle \in \mathcal{H}, \forall \lambda \in \mathbb{C}$$

In particular, anti-linearity (second property) solves the absurdum in the proof of Th. 2.2.1. Time reversal is expected to reverse the spin of particles, and this can be used to construct \mathcal{T} . First, define:

$$\xi^{-s} \equiv -i\sigma^2[\xi^s]^* \quad (2.44)$$

that is, $\xi^{-1} = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$ and $\xi^{-2} = \begin{pmatrix} -1 \\ 0 \end{pmatrix}$. This allows redefining $\eta^s \equiv \xi^{-s}$, and also $\xi^{-(-s)} = -\xi^s$.

Lemma 2.2.4 (Reversed spinors)

$$u^{-s}(-\mathbf{p}) = -\gamma^1\gamma^3[u^s(\mathbf{p})]^* \quad v^{-s}(-\mathbf{p}) = -\gamma^1\gamma^3[v^s(\mathbf{p})]^* \quad (2.45)$$

Proof. Defining $\tilde{p}^\mu \equiv (p^0, -\mathbf{p})$ and recalling that $\boldsymbol{\sigma}\boldsymbol{\sigma}^2 = -\boldsymbol{\sigma}^2\boldsymbol{\sigma}^*$:

$$u^{-s}(-\mathbf{p}) = \begin{pmatrix} \sqrt{\tilde{p}^\mu\sigma_\mu}(-i\sigma^2[\xi^s]^*) \\ \sqrt{\tilde{p}^\mu\bar{\sigma}_\mu}(-i\sigma^2[\xi^s]^*) \end{pmatrix} = \begin{pmatrix} -i\sigma^2\sqrt{p^\mu\sigma_\mu^*}[\xi^s]^* \\ -i\sigma^2\sqrt{p^\mu\bar{\sigma}_\mu^*}[\xi^s]^* \end{pmatrix} = -i \begin{pmatrix} \sigma^2 & 0 \\ 0 & \sigma^2 \end{pmatrix} [u^s(\mathbf{p})]^*$$

Noting that $i \text{diag}(\sigma^2, \sigma^2) = \gamma^1\gamma^3$ completes the proof. On the other hand:

$$\begin{aligned}v^{-s}(-\mathbf{p}) &= \begin{pmatrix} \sqrt{\tilde{p}^\mu\sigma_\mu}(-\xi^s) \\ -\sqrt{\tilde{p}^\mu\bar{\sigma}_\mu}(-\xi^s) \end{pmatrix} = \begin{pmatrix} \sigma^2\sqrt{p^\mu\sigma_\mu^*}\sigma^2(-\xi^s) \\ -\sigma^2\sqrt{p^\mu\bar{\sigma}_\mu^*}\sigma^2(-\xi^s) \end{pmatrix} \\ &= -i \begin{pmatrix} \sigma^2 & 0 \\ 0 & \sigma^2 \end{pmatrix} \begin{pmatrix} \sqrt{p^\mu\sigma_\mu^*}(-i\sigma^2[\xi^s]^*) \\ -\sqrt{p^\mu\bar{\sigma}_\mu^*}(-i\sigma^2[\xi^s]^*) \end{pmatrix} = -\gamma^1\gamma^3[v^s(-\mathbf{p})]^*\end{aligned}$$

\square

It is useful to define the reversed ladder operators:

$$a_{\mathbf{p},-s} \equiv (a_{\mathbf{p},2}, -a_{\mathbf{p},1}) \quad b_{\mathbf{p},-s} \equiv (b_{\mathbf{p},2}, -b_{\mathbf{p},1}) \quad (2.46)$$

Definition 2.2.2 (Time reversal)

The **time reversal operator** is defined as:

$$\mathcal{T}a_{\mathbf{p},s}\mathcal{T} = a_{-\mathbf{p},-s} \quad \mathcal{T}b_{\mathbf{p},s}\mathcal{T} = b_{-\mathbf{p},-s} \quad (2.47)$$

An additional overall phase is irrelevant. Moreover, as other discrete symmetries, $\mathcal{T}^2 = 1$.

	$\bar{\Psi}\Psi$	$\bar{\Psi}\gamma^5\Psi$	$\bar{\Psi}\gamma^\mu\Psi$	$\bar{\Psi}\gamma^\mu\gamma^5\Psi$	$\bar{\Psi}\sigma^{\mu\nu}\Psi$	∂_μ
\mathcal{C}	+1	+1	-1	+1	-1	+1
\mathcal{P}	+1	-1	$(-1)^\mu$	$-(-1)^\mu$	$(-1)^\mu(-1)^\nu$	$(-1)^\mu$
\mathcal{T}	+1	-1	$(-1)^\mu$	$(-1)^\mu$	$-(-1)^\mu(-1)^\nu$	$-(-1)^\mu$
\mathcal{CPT}	+1	+1	-1	-1	+1	-1

Table 2.2: Eigenvalues of fermion bilinears and derivative operator, with the shorthand $(-1)^\mu \equiv (+1, -1, -1, -1)$.

Proposition 2.2.8 (Time reversal on Dirac fields)

Given a Dirac field $\Psi(x)$:

$$\mathcal{T}\Psi(t, \mathbf{x})\mathcal{T} = -\gamma^1\gamma^3\Psi(-t, \mathbf{x}) \quad (2.48)$$

Proof. Using Lemma 2.2.4 and defining $\tilde{p}^\mu \equiv (p^0, -\mathbf{p})$, $\tilde{x}^\mu \equiv (-t, \mathbf{x})$, so that $p^\mu x_\mu = -\tilde{p}^\mu \tilde{x}_\mu$:

$$\begin{aligned}
\mathcal{T}\Psi(x)\mathcal{T} &= \int \frac{d^3p}{(2\pi)^3\sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} \mathcal{T} [a_{\mathbf{p},s}u^s(p)e^{-ip_\mu x^\mu} + b_{\mathbf{p},s}^\dagger v^s(p)e^{ip_\mu x^\mu}] \mathcal{T} \\
&= \int \frac{d^3p}{(2\pi)^3\sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [a_{-\mathbf{p},-s}[u^s(p)]^* e^{ip_\mu x^\mu} + b_{-\mathbf{p},-s}^\dagger [v^s(p)]^* e^{-ip_\mu x^\mu}] \\
&= -\gamma^3\gamma^1 \int \frac{d^3p}{(2\pi)^3\sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} [a_{\tilde{\mathbf{p}},-s}u^{-s}(\tilde{p})e^{-i\tilde{p}_\mu \tilde{x}^\mu} + b_{\tilde{\mathbf{p}},-s}^\dagger v^{-s}(\tilde{p})e^{i\tilde{p}_\mu \tilde{x}^\mu}] \\
&= \gamma^1\gamma^3 \int \frac{-d^3\tilde{p}}{(2\pi)^3\sqrt{2E_{\tilde{\mathbf{p}}}}} \sum_{s=1,2} [a_{\tilde{\mathbf{p}},s}u^s(\tilde{p})e^{-i\tilde{p}_\mu \tilde{x}^\mu} + b_{\tilde{\mathbf{p}},s}^\dagger v^s(\tilde{p})e^{i\tilde{p}_\mu \tilde{x}^\mu}] = -\gamma^1\gamma^3\Psi(\tilde{x})
\end{aligned}$$

□

§2.2.3.4 CPT symmetry

In order to study the invariance properties of fermionic Lagrangians, it is necessary to state the transformation relations of fermion bilinears (see Sec. 3.6 of [6] for details).

This shows that it is not possible to construct a Lorentz-invariant Lagrangian which violates CPT symmetry: this is an example of the **CPT theorem**, which states that, independently of the spin of the particle, a (local) Lorentz-invariant field theory with a hermitian Hamiltonian cannot violate CPT symmetry. A consequence of this theorem is that particles and antiparticles have exactly the same mass, which has been so far empirically confirmed.

§2.3 Electromagnetic field

§2.3.1 Maxwell theory

The electromagnetic field is described by a 4-vector A_μ , the **gauge potential**. From this, the field strength tensor is defined as:

$$F_{\mu\nu} := \partial_\mu A_\nu - \partial_\nu A_\mu \quad (2.49)$$

which is related to the electric and magnetic fields as $F^{0i} = -E^i$ and $F^{ij} = -\epsilon^{ijk} B^k$. The Lagrangian of the free electromagnetic field is:

$$\mathcal{L}_M = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \quad (2.50)$$

The associated equations of motion are:

$$\partial_\mu F^{\mu\nu} = 0 \quad (2.51)$$

Moreover, defining $\tilde{F}^{\mu\nu} \equiv \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} F_{\rho\sigma}$ (the Hodge dual), it is trivial to check that, by Schwarz lemma:

$$\partial_\mu \tilde{F}^{\mu\nu} = 0 \quad (2.52)$$

Eq. 2.51-2.52 are exactly Maxwell equations in the absence of sources: when written in terms of \mathbf{E} and \mathbf{B} , Eq. 2.51 gives the equations for $\nabla \cdot \mathbf{E}$ and $\nabla \times \mathbf{B}$, while Eq. 2.52 those for $\nabla \times \mathbf{E}$ and $\nabla \cdot \mathbf{B}$.

§2.3.1.1 Gauge invariance

A crucial local symmetry of the Maxwell Lagrangian is the symmetry under local gauge transformations like:

$$A_\mu(x) \mapsto A_\mu(x) - \partial_\mu \alpha(x) \quad (2.53)$$

with arbitrary $\alpha \in \mathcal{C}^\infty(\mathbb{R}^{1,3})$. Considering the free electromagnetic field, the global version of this transformation (that is, α independent of x) yields no conserved charge, as the associated Noether current vanishes identically.

Theorem 2.3.1 (Radiation gauge)

In the absence of sources, it is always possible to choose the **radiation gauge**:

$$A_0 = 0 \quad \nabla \cdot \mathbf{A} = 0 \quad (2.54)$$

Proof. Starting from a general gauge potential A_μ , the condition $A_0 = 0$ is achieved through:

$$A_\mu \mapsto A_\mu - \partial_\mu \int_{t_0}^t d\tau A_0(\tau, \mathbf{x})$$

Then, $A_0 = 0$ will remain unchanged if another gauge transformation with $\alpha(x) = \alpha(\mathbf{x})$ is performed. Consider:

$$\alpha(\mathbf{x}) = - \int_{\mathbb{R}^3} \frac{d^3 y}{4\pi |\mathbf{x} - \mathbf{y}|} \partial_i A^i(t, \mathbf{y})$$

which is independent of t since $E^i = -\partial_0 A^i$, as $A_0 = 0$, so $\partial_i E^i = 0$ implies $\partial_0 \partial_i A^i = 0$. Recall the identity:

$$\Delta_{\mathbf{x}} \frac{1}{|\mathbf{x} - \mathbf{y}|} = -4\pi \delta^{(3)}(\mathbf{x} - \mathbf{y}) \quad (2.55)$$

Thus:

$$\nabla \cdot \mathbf{A} \mapsto \nabla \cdot \mathbf{A} - \Delta_{\mathbf{x}} \alpha = \partial_i A^i(t, \mathbf{x}) - \partial_i A^i(t, \mathbf{x}) = 0$$

□

The radiation gauge clearly implies the **Lorentz gauge**:

$$\partial_\mu A^\mu = 0 \quad (2.56)$$

In this gauge, the equations of motions Eq. 2.51 become:

$$\square A^\mu = 0 \quad (2.57)$$

which are massless KG equations for each component of the gauge potential. Plane-wave solutions take the form:

$$A_\mu(x) = 2\Re\{\epsilon_\mu(k)e^{-ik_\mu x^\mu}\} \quad (2.58)$$

where $\epsilon_\mu(x)$ is the **polarization vector**. Then, Eq. 2.57 gives $k^2 = 0$, while the chosen radiation gauge implies $\epsilon_0 = 0$ and $\boldsymbol{\epsilon} \cdot \mathbf{k} = 0$: therefore, an electromagnetic wave has only two degrees of freedom, represented by a polarization vector $\boldsymbol{\epsilon}$ perpendicular to the direction of propagation.

Example 2.3.1 (Linear and circular polarization)

Given an electromagnetic wave with $\hat{\mathbf{k}} = \hat{\mathbf{e}}_z$, then there are two possible polarization directions perpendicular to $\hat{\mathbf{e}}_z$. A possible choice are **linear polarization vectors**:

$$\boldsymbol{\epsilon}_1 = (1, 0, 0) \quad \boldsymbol{\epsilon}_2 = (0, 1, 0) \quad (2.59)$$

Another choice is a linear combination of these vectors, the **circular polarization vectors**:

$$\boldsymbol{\epsilon}_R = \frac{\boldsymbol{\epsilon}_1 + i\boldsymbol{\epsilon}_2}{\sqrt{2}} \quad \boldsymbol{\epsilon}_L = \frac{\boldsymbol{\epsilon}_1 - i\boldsymbol{\epsilon}_2}{\sqrt{2}} \quad (2.60)$$

which are helicity eigenstates respectively with $h = +1$ and $h = -1$.

The advantage of the radiation gauge is that it exposes clearly the physical degrees of freedom of the electromagnetic field, while sacrificing explicit Lorentz covariance; on the other hand, the Lorentz gauge retains the explicit Lorentz covariance, at the cost of redundant degrees of freedom.

§2.3.1.2 Energy-momentum tensor

By Eq. 1.80, writing Eq. 2.50 explicitly in terms of A_μ , the energy-momentum tensor of the electromagnetic field is:

$$\theta^{\mu\nu} = -F^{\mu\rho}\partial^\nu A_\rho + \frac{1}{4}\eta^{\mu\nu}F^2 \quad (2.61)$$

with $F^2 \equiv F_{\mu\nu}F^{\mu\nu}$. To show the gauge-invariance of this tensor, recall Eq. 2.51:

$$\theta^{\mu\nu} \mapsto \theta^{\mu\nu} + F^{\mu\rho}\partial^\nu\partial_\rho\alpha \implies P^\mu \mapsto P^\mu + \int d^3x \partial_\rho(F^{0\rho}\partial^\mu\alpha) = P^\mu + \int d^3x \partial_i(F^{0i}\partial^\mu\alpha) = P^\mu$$

where the last term is a total spatial derivative, hence vanishing by divergence theorem provided that the field decreases sufficiently fast at infinity. To improve the energy-momentum tensor, add $\partial_\rho(F^{\mu\rho}A^\nu)$, which is covariantly conserved by itself and whose $\mu = 0$ component is a total spatial derivative, so to obtain:

$$T^{\mu\nu} = F^{\mu\rho}F_\rho{}^\nu + \frac{1}{4}\eta^{\mu\nu}F^2 \quad (2.62)$$

which is explicitly gauge-invariant and yields the usual expressions for the energy density $T^{00} = \frac{1}{2}(\mathbf{E}^2 + \mathbf{B}^2)$ and the momentum density $T^{0i} = (\mathbf{E} \times \mathbf{B})^i$.

In a general field theory, the observable quantities are the charges, not the currents: two Lagrangian densities which differ by a total 4-divergence are physically equivalent and give the same equations of motion, but the conserved currents obtained through Noether theorem are different, while the associated Noether charges are the same.

§2.3.1.3 Matter coupling

In the presence of an external current j^μ , Eq. 2.52 is not modified, as it is a consequence of the definition of $F^{\mu\nu}$ (assuming regular gauge fields), while Eq. 2.51 becomes:

$$\partial_\mu F^{\mu\nu} = j^\nu \quad (2.63)$$

By Schwarz lemma, this equation is consistent only if $\partial_\mu j^\mu = 0$. This can be understood in light of gauge invariance, considering the action:

$$\mathcal{S}_M = - \int d^4x \left[\frac{1}{4}F^2 + j^\mu A_\mu \right] \quad (2.64)$$

A gauge transformation $A_\mu \mapsto A_\mu - \partial_\mu \alpha$ implies $\mathcal{S}_M \mapsto \mathcal{S}_M + \int d^4x j^\mu \partial_\mu \alpha$: integrating by parts, it is clear that \mathcal{S}_M is gauge invariant only if $\partial_\mu j^\mu = 0$.

Dirac field The coupling of the electromagnetic field to the Dirac field is an example of the general procedure of writing a gauge-invariant action for a gauge theory. In particular, consider a theory with a global U(1) invariance, which is a symmetry of the free Dirac action by Eq. 1.112. Now generalize to a local U(1) symmetry:

$$\Psi(x) \mapsto e^{iq\alpha(x)}\Psi(x) \quad (2.65)$$

with $q \in \mathbb{R}$. This no longer is a symmetry of the Dirac action, however it can be combined with Eq. 2.53 defining the **covariant derivative**:

$$D_\mu \Psi := (\partial_\mu + iqA_\mu)\Psi \quad (2.66)$$

Proposition 2.3.1

$$D_\mu \Psi(x) \mapsto e^{iq\alpha(x)}D_\mu \Psi(x) \quad (2.67)$$

Proof. $D_\mu \Psi \mapsto [\partial_\mu + iq(A_\mu - \partial_\mu \alpha)]e^{iq\alpha}\Psi = e^{iq\alpha}[iq\partial_\mu \alpha + \partial_\mu + iqA_\mu - iq\partial_\mu \alpha]\Psi = e^{iq\alpha}D_\mu \Psi$ \square

The Lagrangian with a local U(1) symmetry is found replacing $\partial_\mu \mapsto D_\mu$ (**minimal coupling**): the global symmetry is gauged to a local symmetry, resulting in a gauge theory with gauge field A_μ . Applying this to Eq. 1.101:

$$\mathcal{L}_D = \bar{\Psi}(i\not{\partial} - m)\Psi - qA_\mu \bar{\Psi}\gamma^\mu\Psi \quad (2.68)$$

where $j_V^\mu := \bar{\Psi}\gamma^\mu\Psi$ is the Noether current associated to the global U(1) symmetry. The associated conserved charge then is:

$$Q = \int d^3x \bar{\Psi}\gamma^0\Psi = \int d^3x \Psi^\dagger\Psi \quad (2.69)$$

Complex scalar field A complex scalar field has a global U(1) symmetry $\phi \mapsto e^{iq\alpha}\phi$, thus the covariant derivative is identical to Eq. 2.67 and the gauged Lagrangian reads (recall Eq. 1.91):

$$\mathcal{L} = \partial_\mu\phi\partial^\mu\phi + iqA^\mu(\phi\partial_\mu\phi^* - \phi^*\partial_\mu\phi) + q^2|\phi|^2 A_\mu A^\mu - m^2\phi^*\phi \quad (2.70)$$

where $j_\mu := i\phi^*\overleftrightarrow{\partial}_\mu\phi$ is the Noether current associated to the global U(1) symmetry.

Higher order interaction terms Although a real scalar field cannot be coupled to the electromagnetic field through the minimal coupling (as the real condition imposes $q = 0$, i.e. a neutral field), interaction terms are possible via higher order terms, as $\mathcal{L}_{\text{int}} \sim \phi F_{\mu\nu}F^{\mu\nu}$ or $\mathcal{L}_{\text{int}} \sim \phi\epsilon_{\mu\nu\rho\sigma}F^{\mu\nu}F^{\rho\sigma}$.

The same is possible for the Dirac field too, for example with $\mathcal{L}_{\text{int}} \sim \bar{\Psi}\sigma^{\mu\nu}\Psi F_{\mu\nu}$: note, however, that these non-minimal couplings have dimensional coupling constants (with dimension of the inverse of a mass), which have a less fundamental significance than dimensionless coupling constant.

Example 2.3.2 (Neutral pions)

The neutral pion π^0 is described by a pseudoscalar field, thus its interaction with the electromagnetic field needs to be a pseudoscalar term, like $\epsilon_{\mu\nu\rho\sigma}F^{\mu\nu}F^{\rho\sigma}$ (this gives a good phenomenological description), as opposed to parity-invariant terms like $F_{\mu\nu}F^{\mu\nu}$.

§2.3.2 Quantization

Due to gauge symmetry, the gauge field gives a redundant physical description, therefore the quantization procedure can be carried in two different ways: fixing the gauge, thus working with only physical degrees of freedom but at the cost of losing explicit Lorentz invariance, or considering the whole A_μ , hence carrying spurious degrees of freedom.

§2.3.2.1 Quantization in the radiation gauge

As of Eq. 2.54, Maxwell equations Eq. 2.57 read $\square\mathbf{A} = \mathbf{0}$, with general classical solution:

$$\mathbf{A}(x) = \int \frac{d^3p}{(2\pi)^3\sqrt{2\omega_{\mathbf{p}}}} \sum_{\lambda=1,2} [\boldsymbol{\epsilon}(\mathbf{p}, \lambda)a_{\mathbf{p},\lambda}e^{-ip_\mu x^\mu} + \boldsymbol{\epsilon}^*(\mathbf{p}, \lambda)a_{\mathbf{p},\lambda}^*e^{ip_\mu x^\mu}]_{p^0=\omega_{\mathbf{p}}}$$

Inserting this expression into $\square\mathbf{A} = \mathbf{0}$ results in the mass-shell condition $p^2 = 0$, i.e. $\omega_{\mathbf{p}} = |\mathbf{p}|$. On the other hand, the gauge condition $\nabla \cdot \mathbf{A} = 0$ requires $\boldsymbol{\epsilon} \cdot \mathbf{p} = 0$: for each fixed \mathbf{p} , this

solution has two orthonormal solutions labelled by $\lambda = 1, 2$, which describe the two physical degrees of freedom of the electromagnetic field. The classical solution can be promoted to a hermitian operator as:

$$\mathbf{A}(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2\omega_{\mathbf{p}}}} \sum_{\lambda=1,2} \left[\boldsymbol{\epsilon}(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda} e^{-ip_\mu x^\mu} + \boldsymbol{\epsilon}^*(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda}^\dagger e^{ip_\mu x^\mu} \right]_{p^0=\omega_{\mathbf{p}}} \quad (2.71)$$

and imposing the canonical commutation relations:

$$[a_{\mathbf{p},\lambda}, a_{\mathbf{q},\lambda'}^\dagger] = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q}) \delta_{\lambda\lambda'} \quad (2.72)$$

In terms of commutators of A^i and conjugate momenta, consider that $\Pi_0 = 0$ (as $A_0 = 0$) and³:

$$\Pi^i = \frac{\delta}{\delta(\partial_0 A_i)} \left(-\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \right) = \frac{\delta}{\delta(\partial_0 A_i)} \left(-\frac{1}{2} F_{0i} F^{0i} \right) = -F^{0i} = E^i$$

Lemma 2.3.1

$$\frac{1}{2} \sum_{\lambda=1,2} [\epsilon^i(\mathbf{k}, \lambda) \epsilon^{*j}(\mathbf{k}, \lambda) + \epsilon^{*i}(-\mathbf{k}, \lambda) \epsilon^j(-\mathbf{k}, \lambda)] = \delta^{ij} - \frac{k^i k^j}{\mathbf{k}^2} \quad (2.73)$$

Proof. Trivially verified in a frame where $\mathbf{k} = (0, 0, k)$ choosing $\boldsymbol{\epsilon}(\mathbf{k}, 1) = (1, 0, 0)$ and $\boldsymbol{\epsilon}(\mathbf{k}, 2) = (0, 1, 0)$ and valid in any frame^a is ensured as both sides transform as tensor under rotations. \square

^aIn particular, the linear polarizations satisfy:

$$\sum_{\lambda=1,2} \epsilon^i(\mathbf{k}, \lambda) \epsilon^{*j}(\mathbf{k}, \lambda) = \delta^{ij} - \frac{k^i k^j}{\mathbf{k}^2} \quad (2.74)$$

Lemma 2.3.2

$$[A^i(t, \mathbf{x}), E^j(t, \mathbf{y})] = -i \int \frac{d^3k}{(2\pi)^3} e^{i\mathbf{k} \cdot (\mathbf{x} - \mathbf{y})} \left(\delta^{ij} - \frac{k^i k^j}{\mathbf{k}^2} \right) \quad (2.75)$$

The r.h.s. of this equation is similar to a Dirac delta and it is called **transverse Dirac delta**, defined in order to get $[\nabla \cdot \mathbf{A}(t, \mathbf{x}), \mathbf{E}(t, \mathbf{y})] = \mathbf{0}$ (as $\nabla \cdot \mathbf{A} = 0$), being the integrand proportional to:

$$k^i \left(\delta^{ij} - \frac{k^i k^j}{\mathbf{k}^2} \right) = k^j - k^j = 0$$

Fock space The standard construction of the Fock space proceeds defining the vacuum state $a_{\mathbf{p},\lambda} |0\rangle = 0$. The Hamiltonian and the linear momentum are then found as:

$$H = \frac{1}{2} \mathfrak{N} \int d^3x [\mathbf{E}^2 + \mathbf{B}^2] = \int \frac{d^3k}{(2\pi)^3} \sum_{\lambda=1,2} \omega_{\mathbf{k}} a_{\mathbf{k},\lambda}^\dagger a_{\mathbf{k},\lambda} \quad (2.76)$$

³Note that Π^i is the momentum conjugate to A_i , while $\Pi_i = -\Pi^i$ is the one conjugate to A^i .

$$\mathbf{P} = \mathfrak{N} \int d^3x \mathbf{E} \times \mathbf{B} = \int \frac{d^3k}{(2\pi)^3} \sum_{\lambda=1,2} \mathbf{k} a_{\mathbf{k},\lambda}^\dagger a_{\mathbf{k},\lambda} \quad (2.77)$$

Therefore, $a_{\mathbf{k},\lambda}^\dagger |0\rangle$ describes a massless particle with energy $\omega_{\mathbf{k}}$ and momentum \mathbf{k} . To study spin, it is necessary to switch to circular polarizations:

$$a_{\mathbf{k},\pm}^\dagger := \frac{1}{\sqrt{2}} \left(a_{\mathbf{k},1}^\dagger \pm i a_{\mathbf{k},2}^\dagger \right) \quad (2.78)$$

where $a_{\mathbf{k},1}, a_{\mathbf{k},2}$ are the linear polarizations $\boldsymbol{\epsilon}(\mathbf{k}, 1) = (1, 0, 0)$, $\boldsymbol{\epsilon}(\mathbf{k}, 2) = (0, 1, 0)$. Linear polarizations are not helicity eigenstates, while circular polarizations are:

$$\begin{aligned} S^3 a_{\mathbf{k},1}^\dagger |0\rangle &= +i a_{\mathbf{k},2}^\dagger |0\rangle & S^3 a_{\mathbf{k},+}^\dagger |0\rangle &= +a_{\mathbf{k},+}^\dagger |0\rangle \\ S^3 a_{\mathbf{k},2}^\dagger |0\rangle &= -i a_{\mathbf{k},1}^\dagger |0\rangle & S^3 a_{\mathbf{k},-}^\dagger |0\rangle &= -a_{\mathbf{k},-}^\dagger |0\rangle \end{aligned}$$

In conclusion, $a_{\mathbf{k},\pm}^\dagger |0\rangle$ describe massless particles with energy $\omega_{\mathbf{k}}$, momentum \mathbf{k} , spin 1 and helicity ± 1 : these are photons.

Defining angular momentum and boost generators too in terms of ladder operators, it is possible to show that Lorentz invariance is preserved by the quantization procedure, although not explicitly.

Discrete transformations It is possible to define parity and charge conjugation on photon states. The electric field is a true vector, while the magnetic field is a pseudovector, thus the gauge potential is a true vector: $\mathcal{P}\mathbf{A}(t, \mathbf{x}) = -\mathbf{A}(t, -\mathbf{x})$. In terms of photon states:

$$\mathcal{P} |\gamma; \mathbf{k}, \mathbf{s}\rangle = -|\gamma; -\mathbf{k}, \mathbf{s}\rangle \quad (2.79)$$

so the intrinsic parity of physical photon states is -1 .

As the fermionic current changes sign under charge conjugation (Eq. 2.43), it is a symmetry of the QED Lagrangian if $\mathcal{C}A^\mu\mathcal{C} = -A^\mu$, i.e. $\mathcal{C}a_{\mathbf{k},\lambda}^\dagger\mathcal{C} = -a_{\mathbf{k},\lambda}^\dagger$. As $\mathcal{C}|0\rangle = +|0\rangle$ and $\mathcal{C}^2 = \mathbb{1}$ by definition, then $\mathcal{C}a_{\mathbf{k},\lambda}^\dagger|0\rangle = \mathcal{C}a_{\mathbf{k},\lambda}\mathcal{C}\mathcal{C}|0\rangle = -a_{\mathbf{k},\lambda}^\dagger|0\rangle$, or:

$$\mathcal{C} |\gamma; \mathbf{k}, \mathbf{s}\rangle = -|\gamma; \mathbf{k}, \mathbf{s}\rangle \quad (2.80)$$

§2.3.2.2 Covariant quantization

The Maxwell Lagrangian Eq. 2.50 cannot be straightforwardly quantize, as Π^0 cannot be defined due to the absence of $\partial_0 A_0$ terms. The basic idea of the covariant quantization of the electromagnetic field, or Gupta-Bleuler quantization, is to start from a modified Lagrangian:

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} (\partial_\mu A^\mu)^2 \quad (2.81)$$

Conjugate momenta are then found to be:

$$\Pi^\mu = \frac{\partial \mathcal{L}}{\partial (\partial_0 A_\mu)} \implies \Pi^i = -F^{0i} = E^i \quad \Pi^0 = -\partial_\mu A^\mu$$

Canonical commutation relations now take the form:

$$[A^\mu(t, \mathbf{x}), \Pi^\nu(t, \mathbf{y})] = i\eta^{\mu\nu} \delta^{(3)}(\mathbf{x} - \mathbf{y}) \quad (2.82)$$

The metric ensures Lorentz covariance. The equations of motion are $\square A^\mu = 0$, thus the gauge field operators are:

$$A_\mu(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2\omega_{\mathbf{p}}}} \sum_{\lambda=0}^3 \left[\epsilon_\mu(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda} e^{-ip_\mu x^\mu} + \epsilon_\mu^*(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda}^\dagger e^{ip_\mu x^\mu} \right]_{p^0=\omega_{\mathbf{p}}} \quad (2.83)$$

$\square A_\mu = 0$ imposes $p^2 = 0$. Note that the modified Lagrangian is not gauge invariant, so there is no constraint on ϵ^μ ; in the frame $p^\mu = (p, 0, 0, p)$ a convenient choice of basis is $\epsilon^\mu(\mathbf{p}, \lambda) = \delta_\lambda^\mu$, hence only $\lambda = 1, 2$ satisfy $\epsilon_\mu p^\mu = 0$. Eq. 2.82 becomes:

$$[a_{\mathbf{p},\lambda}, a_{\mathbf{p},\lambda'}^\dagger] = -(2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q}) \eta_{\lambda\lambda'} \quad (2.84)$$

Note that the commutator for $\lambda = \lambda' = 0$ has a negative sign, which means that the norm is not positive defined on this Fock space (rendering impossible its interpretation as a probability):

$$\langle \mathbf{p}, \lambda | \mathbf{p}, \lambda \rangle = (2\omega_{\mathbf{p}}) \langle 0 | a_{\mathbf{p},\lambda} a_{\mathbf{p},\lambda}^\dagger | 0 \rangle = (2\omega_{\mathbf{p}}) \langle 0 | [a_{\mathbf{p},\lambda}, a_{\mathbf{p},\lambda}^\dagger] | 0 \rangle = -2\omega_{\mathbf{p}} V \eta_{\lambda\lambda} \quad (2.85)$$

(where $V \equiv (2\pi)^3 \delta^{(3)}(\mathbf{0})$) which is negative for $\lambda = 0$. However, the only physical states are those associated with transverse polarization vectors, thus these problematic states can be shown to be unphysical.

To recover the correct description of QED from the modified Lagrangian, it is necessary to restrict the Fock space; in particular, any two physical states must satisfy:

$$\langle \text{phys}' | \partial_\mu A^\mu | \text{phys} \rangle = 0 \quad (2.86)$$

Note that $\partial_\mu A^\mu$ can be decomposed into its positive- and negative-frequency parts $\partial_\mu A^\mu = (\partial_\mu A^\mu)^+ + (\partial_\mu A^\mu)^-$ as:

$$(\partial_\mu A^\mu)^+ \equiv -i \int \frac{d^3p}{(2\pi)^3 \sqrt{2\omega_{\mathbf{p}}}} \sum_{\lambda=0}^3 p_\mu \epsilon^\mu(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda} e^{-ip_\mu x^\mu} = (\partial_\mu A^\mu)^\dagger$$

Therefore, Eq. 2.86 is equivalent to:

$$(\partial_\mu A^\mu)^+ | \text{phys} \rangle = 0 \quad (2.87)$$

This is the definition of the physical subspace of the Fock space, as it preserves the linear structure of the physical Hilbert space. Now consider the most general superposition of polarization states with momentum \mathbf{k} , i.e. $|\mathbf{k}\rangle = \sum_{\lambda=0}^3 c_\lambda a_{\mathbf{k},\lambda}^\dagger |0\rangle$, and choose the frame $k^\mu = (k, 0, 0, k)$: in this frame the physical-state condition reads $c_0 + c_3 = 0$. Thus, all transverse states $\lambda = 1, 2$ are physical, while there is only one non-transverse state that remains:

$$|\phi\rangle \equiv (a_{\mathbf{k},0}^\dagger - a_{\mathbf{k},3}^\dagger) |0\rangle$$

The most general one-particle state of the physical subspace can then be written as $|\mathbf{k}_T\rangle + c |\phi\rangle$. However:

$$\langle \phi | \phi \rangle = \langle 0 | (a_{\mathbf{k},0} - a_{\mathbf{k},3}) (a_{\mathbf{k},0}^\dagger - a_{\mathbf{k},3}^\dagger) | 0 \rangle = \langle 0 | [a_{\mathbf{k},0}, a_{\mathbf{k},0}^\dagger] + [a_{\mathbf{k},3}, a_{\mathbf{k},3}^\dagger] | 0 \rangle = 0$$

This means that $|\phi\rangle$ is orthogonal to all physical states $|\varphi\rangle$. Moreover, only transverse states contribute to the energy and to the momentum, as $H, \mathbf{P} \sim -\eta^{\lambda\lambda'} a_{\mathbf{k},\lambda}^\dagger a_{\mathbf{k},\lambda'}$ and (as $c_0 + c_3 = 0$):

$$(a_{\mathbf{k},0} - a_{\mathbf{k},3}) |\psi\rangle = 0 \implies \langle \varphi' | -a_{\mathbf{k},0}^\dagger a_{\mathbf{k},0} + a_{\mathbf{k},3}^\dagger a_{\mathbf{k},3} | \varphi \rangle = \langle \varphi' | (-a_{\mathbf{k},0}^\dagger + a_{\mathbf{k},3}^\dagger) a_{\mathbf{k},3} | \varphi \rangle = 0$$

These facts mean that $|\mathbf{k}_T\rangle$ and $|\mathbf{k}_T\rangle + c |\phi\rangle$ are physically indistinguishable: photons can then be identified as the equivalence classes with respect to the equivalence relation $|\psi\rangle \sim |\psi\rangle + c |\phi\rangle$. This procedure has eliminated the spurious degrees of freedom, thus showing the equivalence between covariant quantization and gauge quantization.

§2.3.2.3 Ladder operators

As for the KG field and the Dirac field, the ladder operators of the electromagnetic field can be made explicit too.

Proposition 2.3.2 (Ladder operators)

The ladder operators for the electromagnetic field are (in radiation gauge):

$$\sqrt{2\omega_{\mathbf{k}}}a_{\mathbf{k},\lambda} = i\epsilon(\mathbf{k}, \lambda) \cdot \int d^3x e^{ik_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \mathbf{A}(x) \quad (2.88)$$

$$\sqrt{2\omega_{\mathbf{k}}}a_{\mathbf{k},\lambda}^{\dagger} = -i\epsilon^{*}(\mathbf{k}, \lambda) \cdot \int d^3x e^{-ik_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \mathbf{A}(x) \quad (2.89)$$

Proof. From Eq. 2.71:

$$\partial_0 \mathbf{A}(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2\omega_{\mathbf{p}}}} \sum_{\lambda=1,2} \left[-i\omega_{\mathbf{p}} \epsilon(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda} e^{-ip_{\mu}x^{\mu}} + i\omega_{\mathbf{p}} \epsilon^{*}(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda}^{\dagger} e^{ip_{\mu}x^{\mu}} \right]_{p^0=\omega_{\mathbf{p}}}$$

Therefore:

$$\begin{aligned} \int d^3x e^{ik_{\mu}x^{\mu}} \mathbf{A}(x) &= \int d^3x \int \frac{d^3p}{(2\pi)^3 \sqrt{2\omega_{\mathbf{p}}}} \sum_{\lambda=1,2} \left[\epsilon(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda} e^{i(k-p)_{\mu}x^{\mu}} + \epsilon^{*}(\mathbf{p}, \lambda) a_{\mathbf{p},\lambda}^{\dagger} e^{i(k+p)_{\mu}x^{\mu}} \right] \\ &= \frac{1}{\sqrt{2\omega_{\mathbf{k}}}} \sum_{\lambda=1,2} \left[\epsilon(\mathbf{k}, \lambda) a_{\mathbf{k},\lambda} + \epsilon^{*}(-\mathbf{k}, \lambda) a_{-\mathbf{k},\lambda}^{\dagger} e^{2i\omega_{\mathbf{k}}} \right] \\ \int d^3x e^{ik_{\mu}x^{\mu}} \partial_0 \mathbf{A}(x) &= \frac{i\omega_{\mathbf{k}}}{\sqrt{2\omega_{\mathbf{k}}}} \sum_{\lambda=1,2} \left[-\epsilon(\mathbf{k}, \lambda) a_{\mathbf{k},\lambda} + \epsilon^{*}(-\mathbf{k}, \lambda) a_{-\mathbf{k},\lambda}^{\dagger} e^{2i\omega_{\mathbf{k}}} \right] \end{aligned}$$

These can be combined as:

$$\int d^3x e^{ik_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \mathbf{A}(x) = \frac{i\omega_{\mathbf{k}}}{\sqrt{2\omega_{\mathbf{k}}}} \sum_{\lambda=1,2} [-2\epsilon(\mathbf{k}, \lambda) a_{\mathbf{k},\lambda}] = -i\sqrt{2\omega_{\mathbf{k}}} \sum_{\lambda=1,2} \epsilon(\mathbf{k}, \lambda) a_{\mathbf{k},\lambda}$$

In the radiation gauge $\epsilon(\mathbf{k}, \lambda) \cdot \epsilon(\mathbf{k}, \sigma) = \delta_{\lambda\sigma}$, hence Eq. 2.88. Analogously:

$$\int d^3x e^{-ik_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \mathbf{A}(x) = \frac{i\omega_{\mathbf{k}}}{\sqrt{2\omega_{\mathbf{k}}}} \sum_{\lambda=1,2} [2\epsilon^{*}(\mathbf{k}, \lambda) a_{\mathbf{k},\lambda}^{\dagger}] = i\sqrt{2\omega_{\mathbf{k}}} \sum_{\lambda=1,2} \epsilon^{*}(\mathbf{k}, \lambda) a_{\mathbf{k},\lambda}^{\dagger}$$

This completes the proof. □

Writing them in Lorentz-covariant form (using $\epsilon^0 = \epsilon^3 = 0$):

$$\sqrt{2\omega_{\mathbf{p}}}a_{\mathbf{p},\lambda} = i\epsilon^{\mu}(\mathbf{p}, \lambda) \int d^3x e^{ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 A_{\mu}(x) \quad (2.90)$$

$$\sqrt{2\omega_{\mathbf{p}}}a_{\mathbf{p},\lambda}^{\dagger} = -i\epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^3x e^{-ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 A_{\mu}(x) \quad (2.91)$$

Part II

Interacting Fields

Chapter 3

Perturbative Approach

In the case of free Hamiltonians/Lagrangians, free-particle states are eigenstates of the Hamiltonian and they cannot interact. To describe interactions between particles, non-linear terms must be included in \mathcal{H}/\mathcal{L} , which will couple Fourier modes (and particles that occupy them) to one another.

To preserve causality, these interaction terms must include only products of fields at the same spacetime point. Moreover, the discussion is restricted to terms which do not include derivatives of fields, so that $\mathcal{H}_{\text{int}} = -\mathcal{L}_{\text{int}}$.

Example 3.0.1 (ϕ^4 theory)

The so-called ϕ^4 theory is a scalar field theory with an interaction term of the kind:

$$\mathcal{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{1}{2} m^2 \phi^2 - \frac{\lambda}{4!} \phi^4 \quad (3.1)$$

for $\lambda \in \mathbb{R}$ a dimensionless coupling constant. This interaction describes the Higgs self-interaction in the Standard Model.

§3.1 Time evolution

Consider a system described by a Hamiltonian H . Two different ways of describing its time evolution are possible.

§3.1.1 Schrödinger picture

In the Schrödinger picture, states are treated as time-dependent, while operators are treated as time-independent. Consider an initial state $|\psi(0)\rangle \equiv |\psi\rangle$: its time-evolution is:

$$|\psi(t)\rangle = e^{-iHt} |\psi\rangle \quad (3.2)$$

In this picture, the amplitude for a process $|a(t_0)\rangle \rightarrow |b(t)\rangle$ is:

$$\mathcal{A} = \langle b | e^{-iH(t-t_0)} | a \rangle =: \langle b | S(t, t_0) | a \rangle \quad (3.3)$$

In the $t - t_0 \rightarrow \infty$ limit, S is known as **S -matrix**, and scattering amplitudes simply are its elements. This can be seen as an operator which realizes time evolution, as:

$$|\psi(t)\rangle = S(t) |\psi\rangle \quad (3.4)$$

where $S(t) \equiv S(t, 0)$.

Proposition 3.1.1 (Unitarity of time evolution)

$S(t, t_0)$ is a unitary operator on \mathcal{H} .

Proof. Consider a normalized initial state $|\psi\rangle : \langle\psi|\psi\rangle = 1$. Given a complete orthonormal eigenbasis $\{|n\rangle\}_{n \in \mathbb{N}}$ of \mathcal{H} , then it must be:

$$\sum_{n \in \mathbb{N}} |\langle n|S|\psi\rangle|^2 = 1 \quad (3.5)$$

This can also be written as:

$$\sum_{n \in \mathbb{N}} |\langle n|S|\psi\rangle|^2 = \sum_{n \in \mathbb{N}} \langle\psi|S^\dagger|n\rangle \langle n|S|\psi\rangle = \langle\psi|S^\dagger S|\psi\rangle \quad (3.6)$$

Being $|\psi\rangle$ arbitrary, it must be $S^\dagger S = S S^\dagger = \mathbb{1}_{\mathcal{H}}$. \square

The unitarity of the S -matrix expresses the conservation of probabilities. Moreover, it can be decomposed as:

$$S = \mathbb{I} + iT \quad (3.7)$$

where the **T -matrix** contains the information on the interactions, while the identity \mathbb{I} only represents a part of the incident wave-packet unaffected by interactions.

§3.1.2 Heisenberg picture

In the Heisenberg picture, states are treated as time-independent, while operators are treated as time-dependent. This picture is better suited for QFT, as fields (which are operators) depend both on \mathbf{x} and t , therefore the Heisenberg picture is natural in light of Lorentz covariance. Given a state $|\psi(t)\rangle$ in the Schrödinger picture, in the Heisenberg picture it is defined as:

$$|\psi, t\rangle := e^{iHt} |\psi(t)\rangle \quad (3.8)$$

which is manifestly time-independent. An operator \mathcal{O} in the Schrödinger picture instead becomes:

$$\mathcal{O}(t) := e^{iHt} \mathcal{O} e^{-iHt} \quad (3.9)$$

The label t in the definition of state is necessary to distinguish whose operator the state is eigenstate of: in fact, in general $\mathcal{O}(t) \neq \mathcal{O}(t')$ for $t \neq t'$, so they will have different eigenstates too.

In the Heisenberg picture, the S -matrix becomes:

$$\langle b|S(t, t_0)|a\rangle = \langle b, t|a, t_0\rangle \quad (3.10)$$

§3.2 Asymptotic theory

On a macroscopic scale, interaction times are extremely small. Therefore, it is convenient to make an assumption, the **adiabatic hypothesis**: while the interaction is described by $\mathcal{H} = \mathcal{H}_0 + \mathcal{H}_{\text{int}}$, where \mathcal{H}_0 is the free Hamiltonian, in the far past and the far future \mathcal{H}_{int} is adiabatically turned off, i.e. $\mathcal{H} \rightarrow \mathcal{H}_0$ as $t \rightarrow \pm\infty$. Moreover, as the out-going states can represent in-coming states for a successive process, the two Fock spaces¹ must be isomorphic, i.e. $\mathcal{F}_{\text{in}} \cong \mathcal{F}_{\text{out}}$: in particular, this implies the (physical) uniqueness of the vacuum, as $|0_{\text{in}}\rangle = |0_{\text{out}}\rangle \equiv |0\rangle$. This isomorphism is realized by the S -matrix $S \equiv S(+\infty, -\infty)$ (recall Eq. 3.3):

$$\langle \beta_{\text{out}} | \alpha_{\text{in}} \rangle = \langle \beta_{\text{in}} | S^\dagger | \alpha_{\text{in}} \rangle \quad \implies \quad |\beta_{\text{out}}\rangle = S |\beta_{\text{in}}\rangle \quad \implies \quad \phi_{\text{out}}(x) = S \phi_{\text{in}}(x) S^\dagger$$

where $\phi_{\text{in}}(x), \phi_{\text{out}}(x)$ are the free fields which generate $\mathcal{F}_{\text{in}}, \mathcal{F}_{\text{out}}$. Note that, to preserve covariance, the S -matrix must commute with Poincaré transformations. The adiabatic hypothesis asserts then:

$$\phi(x) \xrightarrow[t \rightarrow -\infty]{} \sqrt{Z} \phi_{\text{in}}(x) \quad \phi(x) \xrightarrow[t \rightarrow +\infty]{} \sqrt{Z} \phi_{\text{out}}(x) \quad (3.12)$$

where $Z \in \mathbb{R}$ is a renormalization factor. These limits must be understood in the weak sense, as they are not operator equations, but they are only valid for each matrix element separately².

§3.2.1 LSZ reduction formula

§3.2.1.1 Scalar fields

Consider a scattering process of a single species of neutral scalar particles. The ladder operators of a free real scalar field can be expressed via Eq. 2.16, which is time-independent, and they can be generalized to an interacting theory, in which case they are time-dependent (as Prop. 1.3.2 is only valid if $\phi(x)$ satisfies the free KG equation).

¹Formally, given a single-particle Hilbert space \mathcal{H} , the Fock space is defined as the following completion:

$$\mathcal{F}_\nu(\mathcal{H}) := \overline{\bigoplus_{n=0}^{\infty} S_\nu \mathcal{H}^{\otimes n}} \quad (3.11)$$

where S_ν is the operator which symmetrizes ($\nu = +$, for bosons) or antisymmetrizes ($\nu = -$, for fermions) the tensors it acts upon. In general:

$$\mathcal{F}_\nu(\mathcal{H}) = \mathbb{C} \oplus \mathcal{H} \oplus S_\nu(\mathcal{H} \otimes \mathcal{H}) \oplus S_\nu(\mathcal{H} \otimes \mathcal{H} \otimes \mathcal{H}) \oplus \dots$$

where $S_\nu \mathcal{H}^{\otimes n}$ consists of n -particle states (\mathbb{C} is the vacuum). A general state then is:

$$|\Psi\rangle_\nu = \sum_{n=0}^{\infty} |\Psi_n\rangle_\nu = a |0\rangle \oplus \sum_i a_i |\psi_i\rangle \oplus \sum_{i,j} a_{ij} |\psi_i, \psi_j\rangle_\nu \oplus \dots$$

The need for this infinite sum to converge in $\mathcal{F}_\nu(\mathcal{H})$ is solved by the completion, as it restricts the Fock space only to states with a finite inner-product-induced norm:

$$\| |\Psi\rangle_\nu \|^2 = \sum_{n=0}^{\infty} \langle \Psi_n | \Psi_n \rangle_\nu < \infty$$

²If that was not the case, then canonical quantization would imply $Z = 1$, i.e. that $\phi(x)$ is a free field. For details, see Sec. 5.1.2 of [5].

Proposition 3.2.1 (Time-dependent ladder operators)

Given an interacting scalar field theory $H = H_0 + H_{\text{int}}$, the time-evolution of generalized ladder operators is given by:

$$\frac{d}{dt}a_{\mathbf{p}}(t) = i[H_{\text{int}}, a_{\mathbf{p}}(t)] \quad (3.13)$$

Proof. Recall the Heisenberg equation for the time-evolution of an operator $\mathcal{O}(t)$ (in the Heisenberg representation):

$$\frac{d}{dt}\mathcal{O}(t) = i[H, \mathcal{O}] + \frac{\partial}{\partial t}\mathcal{O}(t) \quad (3.14)$$

Using Eq. 2.16, the explicit time-dependence of $a_{\mathbf{p}}(t)$ is solved as:

$$\frac{\partial}{\partial t}a_{\mathbf{p}}(t) = \frac{i}{\sqrt{2E_{\mathbf{p}}}} \frac{\partial}{\partial t} \int d^3x [\partial_0\phi(x) - iE_{\mathbf{p}}\phi(x)] e^{ip_{\mu}x^{\mu}}$$

Recall that for a general function $f = f(x_1, \dots, x_n, t)$:

$$\frac{df}{dt} = \sum_{i=1}^n \frac{\partial f}{\partial x_i} \frac{\partial x_i}{\partial t} + \frac{\partial f}{\partial t} \quad (3.15)$$

Therefore, treating $\phi(x)$ and $\partial_0\phi(x)$ as time-dependent variables:

$$\begin{aligned} \frac{\partial}{\partial t}a_{\mathbf{p}}(t) &= \frac{i}{\sqrt{2E_{\mathbf{p}}}} \left[\frac{d}{dt} - \partial_0\phi(x) \frac{\partial}{\partial \phi} - \partial_0^2\phi(x) \frac{\partial}{\partial (\partial_0\phi)} \right] \int d^3x [\partial_0\phi(x) - iE_{\mathbf{p}}\phi(x)] e^{ip_{\mu}x^{\mu}} \\ &= \frac{i}{\sqrt{2E_{\mathbf{p}}}} \int d^3x [\partial_0^2\phi(x) - iE_{\mathbf{p}}\partial_0\phi(x) + iE_{\mathbf{p}}\partial_0\phi(x) - (iE_{\mathbf{p}})^2\phi(x)] e^{ip_{\mu}x^{\mu}} + \\ &\quad - \frac{i}{\sqrt{2E_{\mathbf{p}}}} \int d^3x [-iE_{\mathbf{p}}\partial_0\phi(x) + \partial_0^2\phi(x)] e^{ip_{\mu}x^{\mu}} \\ &= \frac{i}{\sqrt{2E_{\mathbf{p}}}} \int d^3x [iE_{\mathbf{p}}\partial_0\phi(x) - (iE_{\mathbf{p}})^2\phi(x)] e^{ip_{\mu}x^{\mu}} = iE_{\mathbf{p}}a_{\mathbf{p}}(t) \end{aligned}$$

Next, by Eq. 2.7:

$$\begin{aligned} [H, a_{\mathbf{p}}(t)] &= [H_0, a_{\mathbf{p}}(t)] + [H_{\text{int}}, a_{\mathbf{p}}(t)] \\ &= \int \frac{d^3k}{(2\pi)^3} E_{\mathbf{k}} [a_{\mathbf{k}}^{\dagger}(t)a_{\mathbf{k}}(t), a_{\mathbf{p}}(t)] + [H_{\text{int}}, a_{\mathbf{p}}(t)] \\ &= - \int d^3k E_{\mathbf{k}} \delta^{(3)}(\mathbf{k} - \mathbf{p}) a_{\mathbf{k}}(t) + [H_{\text{int}}, a_{\mathbf{p}}(t)] = -E_{\mathbf{p}}a_{\mathbf{p}}(t) + [H_{\text{int}}, a_{\mathbf{p}}(t)] \end{aligned}$$

The first term cancels the partial derivative, yielding the thesis. \square

Moreover, in the adiabatic hypothesis it is possible to define **in-** and **out-ladder operators**:

$$a_{\mathbf{p}}^{\text{in,out}} \equiv \frac{1}{\sqrt{Z}} \lim_{t \rightarrow \mp\infty} a_{\mathbf{p}}(t) \quad (3.16)$$

In terms of the free in- and out-fields, they can be expressed as:

$$\sqrt{2E_{\mathbf{p}}} a_{\mathbf{p}}^{\text{in,out}\dagger} = -\frac{i}{\sqrt{Z}} \lim_{t \rightarrow \mp\infty} \int d^3x e^{-ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \phi(x) = -i \int d^3x e^{-ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0 \phi_{\text{in,out}}(x) \quad (3.17)$$

These operators respectively act on $\mathcal{F}_{\text{in,out}}$. Note that now these operators, although not explicitly time-dependent, depend implicitly on it as they are defined through the interacting field, i.e. they depend on the field at all times: the relation between in- and out-ladder operators is thus non-trivial.

Conversely, the in- and out-fields can be expressed in terms of in- and out-ladder operators:

$$\phi_{\text{in,out}}(x) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \left[a_{\mathbf{p}}^{\text{in,out}} e^{-ip_{\mu}x^{\mu}} + a_{\mathbf{p}}^{\text{in,out}\dagger} e^{ip_{\mu}x^{\mu}} \right]_{p^0=E_{\mathbf{p}}} \quad (3.18)$$

Proposition 3.2.2 (Behavior of ladder operators)

The in- and out-ladder operators are well-defined ladder operators on $\mathcal{F}_{\text{in,out}}$, i.e.:

$$\hat{P}_{\mu} a_{\mathbf{k}}^{\text{in,out}\dagger} |0_{\text{in}}\rangle = k_{\mu} a_{\mathbf{k}}^{\text{in,out}\dagger} |0_{\text{out}}\rangle \quad (3.19)$$

Proof. The thesis is equivalent to^a:

$$[\hat{P}_{\mu}, \phi_{\text{in,out}}(x)] = -i\partial_{\mu}\phi_{\text{in,out}}(x) \quad (3.20)$$

To show this, first note that Eq. 3.13 is equivalent to:

$$a_{\mathbf{k}}(t) = a_{\mathbf{k}}^{\text{in}} + i \int_{-\infty}^t d\tau [H_{\text{int}}, a_{\mathbf{k}}(\tau)] \equiv a_{\mathbf{k}}^{\text{int}} + i \int_{-\infty}^t d\tau \int \frac{d^3x}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} e^{ik_{\mu}x^{\mu}} \chi[\phi(\tau, \mathbf{x})]$$

where $\chi[\phi]$ is a certain functional of ϕ . Then:

$$\begin{aligned} \phi(x) &= \int \frac{d^3k}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} \left[a_{\mathbf{k}} e^{-ik_{\mu}x^{\mu}} + a_{\mathbf{k}}^{\dagger} e^{ik_{\mu}x^{\mu}} \right]_{k^0=E_{\mathbf{k}}} \\ &= \int \frac{d^3k}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} \left[a_{\mathbf{k}}^{\text{in}} e^{-ik_{\mu}x^{\mu}} + a_{\mathbf{k}}^{\text{int}\dagger} e^{ik_{\mu}x^{\mu}} \right]_{k^0=E_{\mathbf{k}}} + \\ &\quad + i \int \frac{d^3k}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} \int_{-\infty}^t d\tau \int \frac{d^3y}{(2\pi)^3 \sqrt{2E_{\mathbf{k}}}} [\chi[\phi(\tau, \mathbf{y})] e^{-ik_{\mu}(x^{\mu}-y^{\mu})} - \text{h.c.}]_{k^0=E_{\mathbf{k}}} \\ &\equiv \phi_{\text{in}}(x) + \int_{y^0 \in (-\infty, x^0]} d^4y G_{\text{S}}(x-y) \chi[\phi(y)] \end{aligned}$$

Therefore, at every time the interacting field can be seen as the in-field plus a source term. Now, consider that:

$$[\hat{P}_{\mu}, \phi(x)] = -i\partial_{\mu}\phi(x)$$

As the source terms vanishes as $t \rightarrow -\infty$ (with $x^0 \equiv t$), this equation holds for ϕ_{in} too, hence the thesis. The proof for ϕ_{out} is analogous. \square

^aRequires proof.

It is possible to make the relationship between in- and out-ladder operators explicit.

Lemma 3.2.1

$$\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p}}^{\text{in}\dagger} - a_{\mathbf{p}}^{\text{out}\dagger}) = \frac{i}{\sqrt{Z}} \int d^4x e^{-ip_{\mu}x^{\mu}} (\square + m^2)\phi(x) \quad (3.21)$$

$$\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p}}^{\text{out}} - a_{\mathbf{p}}^{\text{in}}) = \frac{i}{\sqrt{Z}} \int d^4x e^{ip_{\mu}x^{\mu}} (\square + m^2)\phi(x) \quad (3.22)$$

Proof. For any integrable $f(t, \mathbf{x})$ the following identity holds:

$$\left(\lim_{t \rightarrow +\infty} - \lim_{t \rightarrow -\infty} \right) \int d^3x f(t, \mathbf{x}) = \int_{-\infty}^{+\infty} dt \frac{\partial}{\partial t} \int d^3x f(t, \mathbf{x}) \quad (3.23)$$

Applying this to $f(t, \mathbf{x}) = -iZ^{-1/2}e^{-ik_{\mu}x^{\mu}}\overleftrightarrow{\partial}_0\phi(x)$ and using Eq. 3.17:

$$\begin{aligned} \sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p}}^{\text{out}\dagger} - a_{\mathbf{p}}^{\text{in}\dagger}) &= -\frac{i}{\sqrt{Z}} \int d^4x \partial_0 \left[e^{-ip_{\mu}x^{\mu}} \overleftrightarrow{\partial}_0\phi(x) \right] \\ &= -\frac{i}{\sqrt{Z}} \int d^4x \partial_0 \left[e^{-ip_{\mu}x^{\mu}} \partial_0\phi(x) - \phi(x) \partial_0 e^{-ip_{\mu}x^{\mu}} \right] \\ &= -\frac{i}{\sqrt{Z}} \int d^4x \left[e^{-ip_{\mu}x^{\mu}} \partial_0^2\phi(x) - \phi(x) \partial_0^2 e^{-ip_{\mu}x^{\mu}} \right] \\ &= -\frac{i}{\sqrt{Z}} \int d^4x \left[e^{-ip_{\mu}x^{\mu}} \partial_0^2\phi(x) - \phi(x) (\nabla^2 - m^2) e^{-ip_{\mu}x^{\mu}} \right] \end{aligned}$$

where the last line follows from $\partial_0^2 e^{-ip_{\mu}x^{\mu}} = -p_0^2 e^{-ip_{\mu}x^{\mu}} = (\mathbf{p}^2 - p^2) e^{-ip_{\mu}x^{\mu}} = (\nabla^2 - m^2) e^{-ip_{\mu}x^{\mu}}$. In order to perform integration by parts, note that initial and final particle states are understood to be convoluted to form wave-packets, which are localized in space, while $\phi(x)$ is not localized in time; hence, the second term can be integrated by parts twice, resulting in:

$$\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p}}^{\text{in}\dagger} - a_{\mathbf{p}}^{\text{out}\dagger}) = \frac{i}{\sqrt{Z}} \int d^4x e^{-ip_{\mu}x^{\mu}} (\partial_0^2 - \nabla^2 + m^2)\phi(x)$$

The other operator is found analogously. □

Note that if $\phi(x)$ were a free field, this operator would identically vanish: this is because the KG scalar product is time-independent in a free field theory, and so are ladder operators too. It is now possible to write the amplitude for a generic scattering process (i.e. S -matrix element). In particular, consider a process $|\mathbf{k}_1, \dots, \mathbf{k}_n; -\infty\rangle \rightarrow |\mathbf{p}_1, \dots, \mathbf{p}_m; +\infty\rangle$ with $\mathbf{k}_i \neq \mathbf{p}_j \forall i = 1, \dots, n, j = 1, \dots, m$: in this case, as there are no particles which remain unchanged by the interaction, the S -matrix reduces to the T -matrix.

Definition 3.2.1 (Bosonic chronological product)

Given two bosonic fields $f(x)$ and $g(x)$, their **chronological product** is defined as:

$$\begin{aligned} \mathfrak{T}\{f(x)g(y)\} &:= \begin{cases} g(y)f(x) & y^0 > x^0 \\ f(x)g(y) & y^0 < x^0 \end{cases} \\ &= \theta(y^0 - x^0)g(y)f(x) + \theta(x^0 - y^0)f(x)g(y) \end{aligned} \quad (3.24)$$

where $\theta(x)$ is the Heaviside distribution.

Theorem 3.2.1 (LSZ reduction formula for scalar fields)

The amplitude for the process $|\mathbf{k}_1, \dots, \mathbf{k}_n; -\infty\rangle \rightarrow |\mathbf{p}_1, \dots, \mathbf{p}_m; +\infty\rangle$, with $\mathbf{k}_i \neq \mathbf{p}_j \forall i = 1, \dots, n, j = 1, \dots, m$, is given by the **Lehmann–Symanzik–Zimmermann formula**:

$$\begin{aligned} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | \mathbf{k}_1, \dots, \mathbf{k}_n; -\infty \rangle &= \\ &= \prod_{i=1}^n \frac{k_i^2 - m^2}{i\sqrt{Z}} \int d^4x_i e^{-ik_{i\mu}x_i^\mu} \prod_{j=1}^m \frac{p_j^2 - m^2}{i\sqrt{Z}} \int d^4y_j e^{+ip_{j\mu}y_j^\mu} \times \\ &\quad \times \langle 0 | \mathfrak{T}\{\phi(x_1) \dots \phi(x_n) \phi(y_1) \dots \phi(y_m)\} | 0 \rangle \end{aligned} \quad (3.25)$$

Proof. Using Eq. 3.17 it is possible to extract a particle from the initial state:

$$\begin{aligned} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | \mathbf{k}_1, \dots, \mathbf{k}_n; -\infty \rangle &= \sqrt{2E_{\mathbf{k}_1}} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | a_{\mathbf{k}_1}^{\text{in}\dagger} | \mathbf{k}_2, \dots, \mathbf{k}_n; -\infty \rangle \\ &= \sqrt{2E_{\mathbf{k}_1}} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | a_{\mathbf{k}_1}^{\text{in}\dagger} - a_{\mathbf{k}_1}^{\text{out}\dagger} | \mathbf{k}_2, \dots, \mathbf{k}_n; -\infty \rangle \end{aligned}$$

as $a_{\mathbf{k}_1}^{\text{out}} | \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty \rangle = 0$ since $\mathbf{k}_i \neq \mathbf{p}_j$. Then, by Lemma 3.2.1:

$$\begin{aligned} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | \mathbf{k}_1, \dots, \mathbf{k}_n; -\infty \rangle &= \frac{i}{\sqrt{Z}} \int d^4x_1 e^{-ik_{1\mu}x_1^\mu} (\square_{x_1} + m^2) \times \\ &\quad \times \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | \phi(x_1) | \mathbf{k}_2, \dots, \mathbf{k}_n; -\infty \rangle \end{aligned}$$

Using the same argument (noting that $\mathfrak{T}\{a_{\mathbf{p}_1}^{\text{in}}\phi(x_1)\} = \phi(x_1)a_{\mathbf{p}_1}^{\text{in}}$, $\mathfrak{T}\{a_{\mathbf{p}_1}^{\text{out}}\phi(x_1)\} = a_{\mathbf{p}_1}^{\text{out}}\phi(x_1)$):

$$\begin{aligned} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | \phi(x_1) | \mathbf{k}_2, \dots, \mathbf{k}_n; -\infty \rangle &= \\ &= \sqrt{2E_{\mathbf{p}_1}} \langle \mathbf{p}_2, \dots, \mathbf{p}_m; +\infty | \mathfrak{T}\{(a_{\mathbf{p}_1}^{\text{out}} - a_{\mathbf{p}_1}^{\text{in}})\phi(x_1)\} | \mathbf{k}_2, \dots, \mathbf{k}_n; -\infty \rangle \\ &= \frac{i}{\sqrt{Z}} \int d^4y_1 e^{+ip_{1\mu}y_1^\mu} (\square_{y_1} + m^2) \langle \mathbf{p}_2, \dots, \mathbf{p}_m; +\infty | \mathfrak{T}\{\phi(y_1)\phi(x_1)\} | \mathbf{k}_2, \dots, \mathbf{k}_n; -\infty \rangle \end{aligned}$$

This procedure can be iterated^a, obtaining:

$$\begin{aligned} \langle \mathbf{p}_1, \dots, \mathbf{p}_m; +\infty | \mathbf{k}_1, \dots, \mathbf{k}_n; -\infty \rangle &= \prod_{i=1}^n \frac{i}{\sqrt{Z}} \int d^4x_i e^{-ik_{i\mu}x_i^\mu} \prod_{j=1}^m \frac{i}{\sqrt{Z}} \int d^4y_j e^{+ip_{j\mu}y_j^\mu} \times \\ &\quad \times (\square_{x_i} + m^2)(\square_{y_j} + m^2) \langle 0 | \mathfrak{T}\{\phi(x_1) \dots \phi(x_n) \phi(y_1) \dots \phi(y_m)\} | 0 \rangle \end{aligned}$$

It is possible to define the **N -point Green function** as:

$$G(x_1, \dots, x_N) \equiv \langle 0 | \mathfrak{T}\{\phi(x_1) \dots \phi(x_N)\} | 0 \rangle \quad (3.26)$$

In terms of its (4-dimensional) Fourier transform it reads:

$$G(x_1, \dots, x_N) = \prod_{i=1}^N \int \frac{d^4\xi_i}{(2\pi)^4} e^{i\xi_{i\mu}x_i^\mu} \tilde{G}(\xi_1, \dots, \xi_N)$$

Then:

$$(\square_{x_j} + m^2)G(x_1, \dots, x_N) = - \prod_{i=1}^N \int \frac{d^4\xi_i}{(2\pi)^4} e^{i\xi_{i\mu}x_i^\mu} (\xi_j^2 - m^2) \tilde{G}(\xi_1, \dots, \xi_N)$$

Substituting into the above expression:

$$\begin{aligned}
\prod_{i=1}^N \frac{i}{\sqrt{Z}} \int d^4 x_i e^{-i k_{i\mu} x_i^\mu} (\square_{x_i} + m^2) G(x_1, \dots, x_N) &= \\
&= - \prod_{i=1}^N \frac{i}{\sqrt{Z}} \int \frac{d^4 x_i d^4 \xi_i}{(2\pi)^4} e^{i(\xi_i - k_i)_\mu x_i^\mu} (\xi_i^2 - m^2) \tilde{G}(\xi_1, \dots, \xi_N) \\
&= \prod_{i=1}^N \frac{1}{i\sqrt{Z}} \int d^4 \xi_i \delta^{(4)}(\xi_i - k_i) (\xi_i^2 - m^2) \tilde{G}(\xi_1, \dots, \xi_N) \\
&= \prod_{i=1}^N \frac{k_i^2 - m^2}{i\sqrt{Z}} \tilde{G}(k_1, \dots, k_N) = \prod_{i=1}^N \frac{k_i^2 - m^2}{i\sqrt{Z}} \int d^4 x_i e^{-i k_{i\mu} x_i^\mu} G(x_1, \dots, x_N)
\end{aligned}$$

Using $G(x_1, \dots, x_n, y_1, \dots, y_m) = \langle 0 | \mathfrak{T} \{ \phi(x_1) \dots \phi(x_n) \phi(y_1) \dots \phi(y_m) \} | 0 \rangle$ yields the thesis. \square

^aTechnically, \square_{y_1} cannot be extracted from \mathfrak{T} , as it does not commute with the Heaviside distribution in its definition. However, the extraction of \square_{y_1} can be performed accounting for an additional term proportional to $\partial_0 \theta(y_1^0 - x_1^0) [\partial_0 \phi(y_1), \phi(x_1)] \sim \delta^{(4)}(y_1 - x_1)$, which does not alter the singular structure of the LSZ formula, thus leaving the residue and the resulting amplitude unchanged.

Going on mass-shell $p^2 \rightarrow m^2$, the Green function (or **correlation function**) can be shown to have singularities $\sim (p^2 - m^2)^{-1}$ for each particle involved: these poles are precisely cancelled by the factors in the LSZ reduction formula, thus leaving a finite result. In this sense, the amplitude of the process can be seen as the multi-pole residue of the Green function.

§3.2.1.2 Electromagnetic field

Consider now a scattering process of a species of massless spin-1 particles. In the adiabatic hypothesis (recall Eq. 2.90-2.91):

$$\sqrt{2\omega_{\mathbf{p}}} a_{\mathbf{p},\lambda}^{\text{in,out}\dagger} = -i\epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^3 x e^{-ip_\mu x^\mu} \overleftrightarrow{\partial}_0 A_\mu^{\text{in,out}}(x) \quad (3.27)$$

Lemma 3.2.2

$$\sqrt{2\omega_{\mathbf{p}}} (a_{\mathbf{p},\lambda}^{\text{in}\dagger} - a_{\mathbf{p},\lambda}^{\text{out}\dagger}) = \frac{i}{\sqrt{Z}} \epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^4 x e^{-ip_\mu x^\mu} \square A_\mu(x) \quad (3.28)$$

$$\sqrt{2\omega_{\mathbf{p}}} (a_{\mathbf{p},\lambda}^{\text{out}} - a_{\mathbf{p},\lambda}^{\text{in}}) = \frac{i}{\sqrt{Z}} \epsilon^\mu(\mathbf{p}, \lambda) \int d^4 x e^{-ip_\mu x^\mu} \square A_\mu(x) \quad (3.29)$$

Proof. Using Eq. 3.23:

$$\begin{aligned}
\sqrt{2\omega_{\mathbf{p}}} (a_{\mathbf{p},\lambda}^{\text{out}\dagger} - a_{\mathbf{p},\lambda}^{\text{in}\dagger}) &= -\frac{i}{\sqrt{Z}} \epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^4 x \partial_0 \left[e^{-ip_\mu x^\mu} \overleftrightarrow{\partial}_0 A_\mu(x) \right] \\
&= -\frac{i}{\sqrt{Z}} \epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^4 x \left[e^{-ip_\mu x^\mu} \partial_0^2 A_\mu(x) - A_\mu(x) \partial_0^2 e^{-ip_\mu x^\mu} \right] \\
&= -\frac{i}{\sqrt{Z}} \epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^4 x \left[e^{-ip_\mu x^\mu} \partial_0^2 A_\mu(x) - A_\mu(x) \nabla^2 e^{-ip_\mu x^\mu} \right]
\end{aligned}$$

where the last line follows from $\partial_0^2 e^{-ip_\mu x^\mu} = -p_0^2 e^{-ip_\mu x^\mu} = (\mathbf{p} - p^2) e^{-ip_\mu x^\mu} = \nabla^2 e^{-ip_\mu x^\mu}$. By the same observation made in the proof of Lemma 3.2.1, the second term is integrated by parts twice, resulting in:

$$\sqrt{2\omega_{\mathbf{p}}}(a_{\mathbf{p},\lambda}^{\text{in}\dagger} - a_{\mathbf{p},\lambda}^{\text{out}\dagger}) = \frac{i}{\sqrt{Z}} \epsilon^{\mu*}(\mathbf{p}, \lambda) \int d^4x e^{-ip_\mu x^\mu} (\partial_0^2 - \nabla^2) A_\mu(x)$$

The other operator follows analogously. \square

Note that, like for scalar fields, these operators vanish identically when dealing with free fields. Consider now a scattering process $|\mathbf{k}_1, \dots, \mathbf{k}_n; -\infty\rangle \rightarrow |\mathbf{p}_1, \dots, \mathbf{p}_m; +\infty\rangle$ in which no particles remain unchanged by the interaction, i.e. $\mathbf{k}_i \neq \mathbf{p}_j \forall i = 1, \dots, n, \forall j = 1, \dots, m$, so that the S -matrix reduces to the T -matrix. Then, as for Th. 3.2.1, particles can be extracted from the generic amplitude as:

$$\begin{aligned} & \langle \mathbf{k}_1, \dots, \mathbf{k}_n | iT | \mathbf{p}_1, \dots, \mathbf{p}_m \rangle \\ &= \prod_{i=1}^n \frac{k_i^2}{i\sqrt{Z}} \int d^4x_i e^{-ik_{i\mu} x_i^\mu} \sum_{\lambda_i=1,2} \epsilon^{\mu_i*}(\mathbf{k}_i, \lambda_i) \prod_{j=1}^m \int d^4y_j e^{+ip_{j\mu} y_j^\mu} \sum_{\sigma_j=1,2} \epsilon^{\nu_j}(\mathbf{p}_j, \sigma_j) \times \\ & \quad \times \langle 0 | \mathfrak{T} \{ A_{\mu_1}(x_1) \dots A_{\mu_n}(x_n) A_{\nu_1}(y_1) \dots A_{\nu_m}(y_m) \} | 0 \rangle \end{aligned} \quad (3.30)$$

§3.2.1.3 Spinor fields

Consider now a scattering process of a single species of spin- $\frac{1}{2}$ particles (and anti-particles). The adiabatic hypothesis now yields (recall Eq. 2.32-2.33):

$$\sqrt{2E_{\mathbf{p}}} a_{\mathbf{p},s}^{\text{in},\text{out}\dagger} = \frac{1}{\sqrt{Z}} \lim_{t \rightarrow \mp\infty} \int d^3x e^{-ip_\mu x^\mu} \bar{\Psi}(x) \gamma^0 u^s(p) = \int d^3x e^{-ip_\mu x^\mu} \bar{\Psi}_{\text{in},\text{out}}(x) \gamma^0 u^s(p) \quad (3.31)$$

$$\sqrt{2E_{\mathbf{p}}} b_{\mathbf{p},s}^{\text{in},\text{out}\dagger} = \frac{1}{\sqrt{Z}} \lim_{t \rightarrow \mp\infty} \int d^3x e^{ip_\mu x^\mu} \bar{v}^s(p) \gamma^0 \Psi(x) = \int d^3x e^{ip_\mu x^\mu} \bar{v}^s(p) \gamma^0 \Psi_{\text{in},\text{out}}(x) \quad (3.32)$$

Lemma 3.2.3

$$\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p},s}^{\text{in}\dagger} - a_{\mathbf{p},s}^{\text{out}\dagger}) = \frac{i}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) (i\overleftarrow{\not{\partial}} + m) u^s(p) e^{-ip_\mu x^\mu} \quad (3.33)$$

$$\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p},s}^{\text{out}} - a_{\mathbf{p},s}^{\text{in}}) = -\frac{i}{\sqrt{Z}} \int d^4x e^{ip_\mu x^\mu} \bar{u}^s(p) (i\overrightarrow{\not{\partial}} - m) \Psi(x) \quad (3.34)$$

$$\sqrt{2E_{\mathbf{p}}}(b_{\mathbf{p},s}^{\text{in}\dagger} - b_{\mathbf{p},s}^{\text{out}\dagger}) = \frac{i}{\sqrt{Z}} \int d^4x e^{-ip_\mu x^\mu} \bar{v}^s(p) (i\overrightarrow{\not{\partial}} - m) \Psi(x) \quad (3.35)$$

$$\sqrt{2E_{\mathbf{p}}}(b_{\mathbf{p},s}^{\text{out}} - b_{\mathbf{p},s}^{\text{in}}) = -\frac{i}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) (i\overleftarrow{\not{\partial}} + m) v^s(p) e^{ip_\mu x^\mu} \quad (3.36)$$

Proof. Using Eq. 3.23:

$$\begin{aligned}\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p},s}^{\text{in}\dagger} - a_{\mathbf{p},s}^{\text{out}\dagger}) &= -\frac{1}{\sqrt{Z}} \int d^4x \partial_0 [e^{-ip_\mu x^\mu} \bar{\Psi}(x) \gamma^0 u^s(p)] \\ &= -\frac{1}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) [-i\gamma^0 p_0 + \gamma^0 \overleftarrow{\partial}_0] u^s(p) e^{-ip_\mu x^\mu}\end{aligned}$$

By Eq. 1.103 $(\gamma^0 p_0 + \gamma^k p_k - m)u^s(p) = 0$, so:

$$\begin{aligned}\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p},s}^{\text{in}\dagger} - a_{\mathbf{p},s}^{\text{out}\dagger}) &= -\frac{1}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) [i\gamma^k p_k - im + \gamma^0 \overleftarrow{\partial}_0] u^s(p) e^{-ip_\mu x^\mu} \\ &= -\frac{1}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) [\gamma^0 \overleftarrow{\partial}_0 - \gamma^k \overrightarrow{\partial}_k - im] u^s(p) e^{-ip_\mu x^\mu}\end{aligned}$$

Assuming that initial and final particle states are convoluted to form wave-packets, integration by parts in spatial dimensions can be carried out, yielding:

$$\begin{aligned}\sqrt{2E_{\mathbf{p}}}(a_{\mathbf{p},s}^{\text{in}\dagger} - a_{\mathbf{p},s}^{\text{out}\dagger}) &= -\frac{1}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) [\gamma^0 \overleftarrow{\partial}_0 + \gamma^k \overleftarrow{\partial}_k - im] u^s(p) e^{-ip_\mu x^\mu} \\ &= \frac{i}{\sqrt{Z}} \int d^4x \bar{\Psi}(x) [i\overleftarrow{\not{\partial}} + m] u^s(p) e^{-ip_\mu x^\mu}\end{aligned}$$

Other operators are found analogously. □

Note that, like Eq. 3.21, these operators vanish identically when dealing with a free field.

Consider now a general scattering process involving both fermions and anti-fermions (quantities denoted by a tilde) $|\mathbf{k}_1, \dots, \mathbf{k}_{n_1}, \tilde{\mathbf{k}}_1, \dots, \tilde{\mathbf{k}}_{n_2}; -\infty\rangle \rightarrow |\mathbf{p}_1, \dots, \mathbf{p}_{m_1}, \tilde{\mathbf{p}}_1, \dots, \tilde{\mathbf{p}}_{m_2}; +\infty\rangle$ in which no particles remain unchanged by the interaction, i.e. $\mathbf{k}_i \neq \mathbf{p}_j, \mathbf{k}_k \neq \tilde{\mathbf{p}}_l \forall i = 1, \dots, n_1, j = 1, \dots, m_1, k = 1, \dots, n_2, l = 1, \dots, m_2$, so that the S -matrix reduces to the T -matrix. By the same reasoning of Th. 3.2.1, it is clear how to extract particles from this generic amplitude:

$$\langle \beta_{\text{fin}}; +\infty | \mathbf{k}, s; -\infty \rangle \mapsto \frac{i}{\sqrt{Z}} \int d^4x \sum_{a=1}^4 \langle \beta_{\text{fin}} | [\bar{\Psi}(x)]_a | 0 \rangle [(i\overleftarrow{\not{\partial}} + m)u^s(k)]_a e^{-ik_\mu x^\mu} \quad (3.37)$$

$$\langle \beta_{\text{fin}}; +\infty | \tilde{\mathbf{k}}, \tilde{s}; -\infty \rangle \mapsto \frac{i}{\sqrt{Z}} \int d^4\tilde{x} \sum_{a=1}^4 e^{-i\tilde{k}_\mu \tilde{x}^\mu} [\bar{v}^{\tilde{s}}(\tilde{k})(i\overrightarrow{\not{\partial}} - m)]_a \langle \beta_{\text{fin}} | [\Psi(\tilde{x})]_a | 0 \rangle \quad (3.38)$$

$$\langle \mathbf{p}, s; +\infty | \alpha_{\text{in}}; -\infty \rangle \mapsto -\frac{i}{\sqrt{Z}} \int d^4y \sum_{a=1}^4 e^{ip_\mu y^\mu} [\bar{u}^s(p)(i\overrightarrow{\not{\partial}} - m)]_a \langle 0 | [\Psi(y)]_a | \alpha_{\text{in}} \rangle \quad (3.39)$$

$$\langle \tilde{\mathbf{p}}, \tilde{s}; +\infty | \alpha_{\text{in}}; -\infty \rangle \mapsto -\frac{i}{\sqrt{Z}} \int d^4\tilde{y} \sum_{a=1}^4 \langle 0 | [\bar{\Psi}(\tilde{y})]_a | \alpha_{\text{in}} \rangle [(i\overleftarrow{\not{\partial}} + m)v^{\tilde{s}}(\tilde{p})]_a e^{i\tilde{p}_\mu \tilde{y}^\mu} \quad (3.40)$$

In general, the final “reduced” amplitude formula will contain a $(n_1 + n_2 + m_1 + m_2)$ -point Green function like (omitting spinor indices):

$$\langle 0 | \mathcal{T} \{ \bar{\Psi}(x_1) \dots \bar{\Psi}(x_{n_1}) \Psi(\tilde{x}_1) \dots \Psi(\tilde{x}_{n_2}) \Psi(y_1) \dots \Psi(y_{m_1}) \bar{\Psi}(\tilde{y}_1) \dots \bar{\Psi}(\tilde{y}_{m_2}) \} | 0 \rangle$$

§3.2.2 Correlation functions

With the LSZ reduction formulae, the problem of computing scattering amplitudes is reduced to that of computing correlation functions.

Consider a generic quantum field $\phi(x)$ described by a Hamiltonian $\mathcal{H} = \mathcal{H}_0 + \mathcal{H}_{\text{int}}$: the exact form of $\phi(x)$ is in general too difficult to obtain, as it satisfies a complicated non-linear equation of motion, so it cannot be written through an expansion in plane waves. To compute correlation functions, then, it is convenient to define a field related to $\phi(x)$:

$$\phi_I(t, \mathbf{x}) := e^{iH_0(t-t_0)} \phi(t_0, \mathbf{x}) e^{-iH_0(t-t_0)} \quad (3.41)$$

This field evolves with H_0 , i.e. is a free field, and it is called the **interaction picture** field. Being this a free field, it can be expanded as:

$$\phi_I(t, \mathbf{x}) = \int \frac{d^3p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \left[a_{\mathbf{p}} e^{-ip_\mu x^\mu} + a_{\mathbf{p}}^\dagger e^{ip_\mu x^\mu} \right]_{p^0=E_{\mathbf{p}}, x^0=t-t_0} \quad (3.42)$$

Proposition 3.2.3 (Interaction picture)

The transformation from the interaction picture to the Heisenberg picture is given by a unitary operator:

$$\phi(t, \mathbf{x}) = U^\dagger(t, t_0) \phi_I U(t, t_0) \quad (3.43)$$

with:

$$U(t, t_0) \equiv e^{iH_0(t-t_0)} e^{-iH(t-t_0)} \quad (3.44)$$

Proof. By direct calculation:

$$\begin{aligned} \phi(t, \mathbf{x}) &= e^{iH(t-t_0)} \phi(t_0, \mathbf{x}) e^{-iH(t-t_0)} \\ &= e^{iH(t-t_0)} e^{-iH_0(t-t_0)} \left[e^{iH_0(t-t_0)} \phi(t_0, \mathbf{x}) e^{-iH_0(t-t_0)} \right] e^{iH_0(t-t_0)} e^{-iH(t-t_0)} \\ &= e^{iH(t-t_0)} e^{-iH_0(t-t_0)} \phi_I(t, \mathbf{x}) e^{iH_0(t-t_0)} e^{-iH(t-t_0)} \equiv U^\dagger(t, t_0) \phi_I U(t, t_0) \end{aligned}$$

Unitarity is obvious, as H is a Hermitian operator. \square

Note that, since $[H_0, H_{\text{int}}] \neq 0$ in general, the two exponentials cannot be combined trivially (the Baker-Campbell-Hausdorff formula should be used instead).

Proposition 3.2.4 (Time evolution operator)

the time evolution operator is:

$$U(t, t_0) = \mathfrak{T} \exp \left[-i \int_{t_0}^t d\tau H_I(\tau) \right] \quad (3.45)$$

where:

$$H_I(t) \equiv e^{iH_0(t-t_0)} H_{\text{int}} e^{-iH_0(t-t_0)} \quad (3.46)$$

Proof. First of all, the time evolution operator solves the Schrödinger equation:

$$\begin{aligned} i\frac{\partial}{\partial t}U(t, t_0) &= e^{iH_0(t-t_0)}(H - H_0)e^{-iH(t-t_0)} = e^{iH_0(t-t_0)}H_{\text{int}}e^{-iH(t-t_0)} \\ &= e^{iH_0(t-t_0)}H_{\text{int}}e^{-iH_0(t-t_0)}e^{iH_0(t-t_0)}e^{-iH(t-t_0)} \equiv H_I(t)U(t, t_0) \end{aligned}$$

The solution to this equation which satisfies the initial condition $U(t_0, t_0) = 1$ is unique:

$$\begin{aligned} U(t, t_0) &= \mathfrak{T} \exp \left[-i \int_{t_0}^t d\tau H_I(\tau) \right] = \sum_{n=0}^{\infty} \frac{(-i)^n}{n!} \int_{[t_0, t]^n} dt_1 \dots dt_n \mathfrak{T}\{H_I(t_1) \dots H_I(t_n)\} \\ &= \sum_{n=0}^{\infty} (-i)^n \int_{t_0}^t dt_1 \int_{t_0}^{t_1} dt_2 \dots \int_{t_0}^{t_{n-1}} dt_n H_I(t_1) \dots H_I(t_n) \end{aligned}$$

It is in fact clear that:

$$\frac{\partial}{\partial t}U(t, t_0) = \sum_{n=0}^{\infty} (-i)^n \int_{t_0}^t dt_2 \dots \int_{t_0}^{t_{n-1}} dt_n H_I(t)H_I(t_2) \dots H_I(t_n) = -iH_I(t)U(t, t_0)$$

Hence, this is a solution to the above equation. \square

Lemma 3.2.4

$$U(t_1, t_0)U(t_0, t_2) = U(t_1, t_2) \quad (3.47)$$

Proof. Observe that:

$$i\frac{\partial}{\partial t}[U(t, t_0)U(t_0, t_2)] = \left[i\frac{\partial}{\partial t}U(t, t_0) \right] U(t_0, t_2) = H_I(t)U(t, t_0)U(t_0, t_2)$$

The boundary condition now reads $[U(t, t_0)U(t_0, t_2)]_{t=t_0} = U(t_0, t_2)$, and the unique solution for this boundary condition is:

$$U(t, t_0)U(t_0, t_2) = \mathfrak{T} \exp \left[-i \int_{t_0}^t d\tau H_I(\tau) \right]$$

which is precisely $U(t, t_2)$. \square

It is now possible to compute correlation functions.

Theorem 3.2.2 (Correlation functions)

Given a quantum field $\phi(x)$ described by a Hamiltonian $\mathcal{H} = \mathcal{H}_0 + \mathcal{H}_{\text{int}}$, then:

$$\langle 0 | \mathfrak{T} \{ \phi(x_1) \dots \phi(x_n) \} | 0 \rangle = \frac{\langle 0 | \mathfrak{T} \{ \phi_I(x_1) \dots \phi_I(x_n) \exp \left[-i \int d^4x \mathcal{H}_I \right] \} | 0 \rangle}{\langle 0 | \mathfrak{T} \{ \exp \left[-i \int d^4x \mathcal{H}_I \right] \} | 0 \rangle} \quad (3.48)$$

Proof. Consider WLOG $\{t_k\}_{k=1,\dots,n} : t_i < t_j$. Using Eq. 3.43 and Lemma 3.2.4:

$$\begin{aligned}\langle 0|\phi(x_1)\dots\phi(x_n)|0\rangle &= \langle 0|U^\dagger(t_1, t_0)\phi_I(x_1)U(t_1, t_0)U^\dagger(t_2, t_0)\dots U^\dagger(t_n, t_0)\phi_I(x_n)U(t_n, t_0)|0\rangle \\ &= \langle 0|U^\dagger(t_1, t_0)\phi_I(x_1)U(t_1, t_2)\dots U(t_{n-1}, t_n)\phi_I(x_n)U(t_n, t_0)|0\rangle\end{aligned}$$

Consider $t \in \mathbb{R} : t \gg t_1 > \dots > t_n \gg -t$, so that $U(t_n, t_0) = U(t_n, -t)U(-t, t_0)$ and $U^\dagger(t_1, t_0) = U^\dagger(t, t_0)U(t, t_1)$. Therefore:

$$\begin{aligned}\langle 0|\phi(x_1)\dots\phi(x_n)|0\rangle &= \langle 0|U^\dagger(t, t_0)[U(t, t_1)\phi_I(x_1)U(t_1, t_2)\phi_I(x_2)U(t_2, t_3)\dots \\ &\quad \dots U(t_{n-1}, t_n)\phi_I(x_n)U(t_n, -t)]U(-t, t_0)|0\rangle\end{aligned}$$

Note that this expression inside the square brackets is automatically time-ordered, so it may be rewritten as:

$$\begin{aligned}[\dots] &= \mathfrak{T}\{\phi_I(x_1)\dots\phi_I(x_n)U(t, t_1)U(t_1, t_2)\dots U(t_n, -t)\} \\ &= \mathfrak{T}\left\{\phi_I(x_1)\dots\phi_I(x_n)\exp\left[-i\int_{-t}^t d\tau H_I(\tau)\right]\right\}\end{aligned}$$

These considerations hold for arbitrary t_0 , so it can be set $t_0 = -t \rightarrow -\infty$. Then:

$$\langle 0|\mathfrak{T}\{\phi(x_1)\dots\phi(x_n)\}|0\rangle = \langle 0|U^\dagger(+\infty, -\infty)\mathfrak{T}\left\{\phi_I(x_1)\dots\phi_I(x_n)\exp\left[-i\int d^4x \mathcal{H}_I\right]\right\}|0\rangle$$

Note that $\langle 0|U^\dagger(\infty, -\infty)$ is the Hermitian conjugate of $U(\infty, -\infty)|0\rangle$, i.e. the state obtained evolving in time the vacuum state. As it was already assumed, the initial-state vacuum coincides physically to the final-state vacuum, i.e.:

$$U(\infty, -\infty)|0\rangle = e^{i\alpha}|0\rangle \quad \implies \quad e^{-i\alpha} = \left(\langle 0|\mathfrak{T}\left\{\exp\left[-i\int d^4x \mathcal{H}_I\right]\right\}|0\rangle\right)^{-1}$$

Finally:

$$\langle 0|\mathfrak{T}\{\phi(x_1)\dots\phi(x_n)\}|0\rangle = \frac{\langle 0|\mathfrak{T}\{\phi_I(x_1)\dots\phi_I(x_n)\exp[-i\int d^4x \mathcal{H}_I]\}|0\rangle}{\langle 0|\mathfrak{T}\{\exp[-i\int d^4x \mathcal{H}_I]\}|0\rangle}$$

which is the thesis. □

This allows to compute correlation functions from free fields, rather than interaction fields. Moreover, note that the functional dependence of \mathcal{H}_I in terms of $\phi_I(x)$ is the same as that of \mathcal{H}_{int} in terms of $\phi(x)$.

Example 3.2.1 (ϕ^4 potential)

Consider the ϕ^4 -interaction Hamiltonian $\mathcal{H}_{\text{int}} = \frac{\lambda}{4!}\phi^4$. Then:

$$\mathcal{H}_I = \frac{\lambda}{4!}e^{iH_0(t-t_0)}\phi^4e^{-iH_0(t-t_0)} = \frac{\lambda}{4!}\left[e^{iH_0(t-t_0)}\phi e^{-iH_0(t-t_0)}\right]^4 = \frac{\lambda}{4!}\phi_I^4$$

Note that Eq. 3.48 is naturally suited for perturbative evaluation with respect to coupling constants which appear in \mathcal{H}_{int} , thanks to the exponential function.

§3.3 Feynman diagrams

The problem of computing scattering amplitudes has been further reduced to that of computing n -point Green functions of free (interaction picture) fields.

§3.3.1 Propagators

§3.3.1.1 Feynman propagator

Theorem 3.3.1 (Feynman propagator)

Given a real scalar field $\phi(x)$, the **Feynman propagator** is computed as:

$$\Delta(x - y) := \langle 0 | \mathfrak{T} \{ \phi_I(x) \phi_I(y) \} | 0 \rangle = \int \frac{d^4 p}{(2\pi)^4} \frac{i}{p^2 - m^2 + i\epsilon} e^{-ip_\mu(x-y)^\mu} \quad (3.49)$$

with $\epsilon \rightarrow 0^+$.

Proof. First, decompose $\phi_I(x)$ into its positive- and negative-frequency parts (i.e. annihilation and creation parts):

$$\phi_I^+(x) \equiv \int \frac{d^3 p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} a_{\mathbf{p}} e^{-ip_\mu x^\mu} \quad \phi_I^-(x) \equiv \int \frac{d^3 p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} a_{\mathbf{p}}^\dagger e^{ip_\mu x^\mu}$$

Clearly $\phi_I^+(x) | 0 \rangle = 0$ and $\langle 0 | \phi_I^-(x) = 0$. Consider first the case $x^0 > y^0$:

$$\begin{aligned} \mathfrak{T} \{ \phi_I(x) \phi_I(y) \} &= \phi_I^+(x) \phi_I^+(y) + \phi_I^+(x) \phi_I^-(y) + \phi_I^-(x) \phi_I^+(y) + \phi_I^-(x) \phi_I^-(y) \\ &= \phi_I^+(x) \phi_I^+(y) + \phi_I^-(y) \phi_I^+(x) + \phi_I^-(x) \phi_I^+(y) + \phi_I^-(x) \phi_I^-(y) + [\phi_I^+(x), \phi_I^-(y)] \\ &= \mathfrak{N} \{ \phi_I(x) \phi_I(y) \} + [\phi^+(x), \phi^-(y)] \end{aligned}$$

Similarly, for $y^0 > x^0$ one has $\mathfrak{T} \{ \phi_I(x) \phi_I(y) \} = \mathfrak{N} \{ \phi_I(x) \phi_I(y) \} + [\phi^+(y), \phi^-(x)]$. Therefore, in general:

$$\begin{aligned} \mathfrak{T} \{ \phi_I(x) \phi_I(y) \} &= \mathfrak{N} \{ \phi_I(x) \phi_I(y) \} + \theta(x^0 - y^0) [\phi_I^+(x), \phi_I^-(y)] + \theta(y^0 - x^0) [\phi_I^+(y), \phi_I^-(x)] \\ &\equiv \mathfrak{N} \{ \phi_I(x) \phi_I(y) \} + \Delta(x - y) \end{aligned}$$

Now, observe that $\langle 0 | \mathfrak{N} \{ \phi_I(x) \phi_I(y) \} | 0 \rangle = 0$, as there is always either an annihilation operator acting on $| 0 \rangle$ or a creation operator acting on $\langle 0 |$. On the other hand, $\Delta(x - y) \in \mathbb{C}$, as $[a_{\mathbf{p}}, a_{\mathbf{q}}^\dagger] \in \mathbb{C}$, therefore $\langle 0 | \Delta(x - y) | 0 \rangle = \Delta(x - y) \langle 0 | 0 \rangle = \Delta(x - y)$. There only remains to compute the commutators:

$$\begin{aligned} \Delta(x - y) &= \int \frac{d^3 p d^3 q}{(2\pi)^6 \sqrt{4E_{\mathbf{p}} E_{\mathbf{q}}}} [\theta(x^0 - y^0) [a_{\mathbf{p}}, a_{\mathbf{q}}^\dagger] e^{-ip_\mu x^\mu + iq_\mu y^\mu} + \\ &\quad + \theta(y^0 - x^0) [a_{\mathbf{p}}, a_{\mathbf{q}}^\dagger] e^{-ip_\mu y^\mu + iq_\mu x^\mu}] \\ &= \int \frac{d^3 p d^3 q}{(2\pi)^3 \sqrt{4E_{\mathbf{p}} E_{\mathbf{q}}}} \delta^{(3)}(\mathbf{p} - \mathbf{q}) [\theta(x^0 - y^0) e^{-ip_\mu x^\mu + iq_\mu y^\mu} + \theta(y^0 - x^0) e^{-ip_\mu y^\mu + iq_\mu x^\mu}] \\ &= \int \frac{d^3 p}{(2\pi)^3 2E_{\mathbf{p}}} [\theta(x^0 - y^0) e^{-ip_\mu(x-y)^\mu} + \theta(y^0 - x^0) e^{ip_\mu(x-y)^\mu}] \end{aligned}$$

To show that this is equivalent to the thesis, note that Eq. 3.49 can be rewritten as:

$$\Delta(x - y) = \int \frac{d^3p}{(2\pi)^3} e^{i\mathbf{p} \cdot (\mathbf{x} - \mathbf{y})} \int_{\mathbb{R}} \frac{dp^0}{2\pi} \frac{i}{(p^0)^2 - E_{\mathbf{p}}^2 + i\epsilon} e^{-ip^0(x^0 - y^0)}$$

as $E_{\mathbf{p}} = +\sqrt{\mathbf{p}^2 + m^2}$. The integral in dp^0 can be computed in the complex p^0 -plane^a: the $i\epsilon$ -prescription slightly displaces the poles from the real axis^b, and in particular that in $p^0 = +E_{\mathbf{p}}$ is slightly below it while that in $p^0 = -E_{\mathbf{p}}$ is slightly above it. If $x^0 - y^0 > 0$, choose a semicircular (clockwise) contour in the lower-half plane, so that only the $p^0 = E_{\mathbf{p}}$ singularity is enclosed:

$$\Delta(x - y)|_{x^0 > y^0} = \int \frac{d^3p}{(2\pi)^3} e^{i\mathbf{p} \cdot (\mathbf{x} - \mathbf{y})} \frac{i}{2\pi} \left[(-2\pi i) \frac{e^{-iE_{\mathbf{p}}(x^0 - y^0)}}{2E_{\mathbf{p}}} \right] = \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} e^{-ip_{\mu}(x - y)^{\mu}}$$

If $x^0 - y^0 < 0$, choose a semicircular (counter-clockwise) contour in the upper-half plane:

$$\Delta(x - y)|_{y^0 > x^0} = \int \frac{d^3p}{(2\pi)^3} e^{i\mathbf{p} \cdot (\mathbf{x} - \mathbf{y})} \frac{i}{2\pi} \left[(2\pi i) \frac{e^{iE_{\mathbf{p}}(x^0 - y^0)}}{-2E_{\mathbf{p}}} \right] = \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} e^{ip_{\mu}(x - y)^{\mu}}$$

where in the last integral $\mathbf{p} \mapsto -\mathbf{p}$ was renamed. The proof is complete. \square

^aAs the integrand is $\sim (p^0)^{-2}$, the contribution of the curved part of a semicircular contour vanishes, only leaving the integral over the real axis.

^bThe poles are $p^0 = \pm \sqrt{E_{\mathbf{p}}^2 - i\epsilon} = \pm \left(E_{\mathbf{p}} - \frac{i\epsilon}{2E_{\mathbf{p}}} \right) + o(\epsilon^2)$.

The expression for the Feynman propagator in momentum space is trivially found:

$$\tilde{\Delta}(p) = \frac{i}{p^2 - m^2 + i\epsilon} \quad (3.50)$$

Moreover, the Feynman propagator is just a Green function for the KG operator:

$$(\square_x + m^2)\Delta(x - y) = \int \frac{d^4p}{(2\pi)^4} \frac{i}{p^2 - m^2 + i\epsilon} (-p^2 + m^2) e^{-ip_{\mu}(x - y)^{\mu}} = -i\delta^{(4)}(x - y)$$

§3.3.1.2 Dirac propagator

To explicitly compute the propagator for the Dirac field, it is necessary to include fermionic anti-symmetry when resolving time-ordering and normal-ordering, i.e. a factor of (-1) for each exchange of fermionic operators, e.g. $\mathfrak{T}\{\Psi(x_1)\Psi(x_2)\Psi(x_3)\Psi(x_4)\} = (-1)^3\Psi(x_3)\Psi(x_1)\Psi(x_4)\Psi(x_2)$ if $x_3^0 > x_1^0 > x_4^0 > x_2^0$, or $\mathfrak{N}\{a_{\mathbf{p}}a_{\mathbf{q}}a_{\mathbf{r}}^{\dagger}\} = (-1)^2a_{\mathbf{r}}^{\dagger}a_{\mathbf{p}}a_{\mathbf{q}} = (-1)^3a_{\mathbf{r}}^{\dagger}a_{\mathbf{q}}a_{\mathbf{p}}$ (equal expressions).

Definition 3.3.1 (Fermionic time-ordering)

Given a Dirac field $\Psi(x)$, the **time-ordering operator** acts as:

$$\mathfrak{T}\{\Psi(x)\bar{\Psi}(y)\} := \begin{cases} \Psi(x)\bar{\Psi}(y) & x^0 > y^0 \\ -\bar{\Psi}(y)\Psi(x) & y^0 > x^0 \end{cases} \quad (3.51)$$

Theorem 3.3.2 (Dirac propagator)

Given a Dirac field $\Psi(x)$, the **Dirac propagator** is computed as:

$$S(x - y) := \langle 0 | \mathfrak{T} \{ \Psi_I(x) \bar{\Psi}_I(y) \} | 0 \rangle = \int \frac{d^4 p}{(2\pi)^4} \frac{i(\not{p} + m)}{p^2 - m^2 + i\epsilon} e^{-ip_\mu(x-y)^\mu} \quad (3.52)$$

Proof. First, decompose $\Psi_I(x)$ and $\Psi_I(y)$ into positive- and negative-frequency parts:

$$\begin{aligned} \Psi_I^+(x) &= \int \frac{d^3 p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} a_{\mathbf{p},s} u^s(p) e^{-ip_\mu x^\mu} & \Psi_I^-(x) &= \int \frac{d^3 p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} b_{\mathbf{p},s}^\dagger v^s(p) e^{ip_\mu x^\mu} \\ \bar{\Psi}_I^+(x) &= \int \frac{d^3 p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} b_{\mathbf{p},s} \bar{v}^s(p) e^{-ip_\mu x^\mu} & \bar{\Psi}_I^-(x) &= \int \frac{d^3 p}{(2\pi)^3 \sqrt{2E_{\mathbf{p}}}} \sum_{s=1,2} a_{\mathbf{p},s}^\dagger \bar{u}^s(p) e^{ip_\mu x^\mu} \end{aligned}$$

Clearly $\Psi_I^+(x) | 0 \rangle = 0 = \bar{\Psi}_I^+(x) | 0 \rangle$ and $\langle 0 | \Psi_I^-(x) = 0 = \langle 0 | \bar{\Psi}_I^-(x)$. Consider first the case $x^0 > y^0$:

$$\begin{aligned} \mathfrak{T} \{ \Psi_I(x) \bar{\Psi}_I(y) \} &= \Psi_I^+(x) \bar{\Psi}_I^+(y) + \Psi_I^+(x) \bar{\Psi}_I^-(y) + \Psi_I^-(x) \bar{\Psi}_I^+(y) + \Psi_I^-(x) \bar{\Psi}_I^-(y) \\ &= \Psi_I^+(x) \bar{\Psi}_I^+(y) - \bar{\Psi}_I^-(y) \Psi_I^+(x) + \Psi_I^-(x) \bar{\Psi}_I^+(y) + \Psi_I^-(x) \bar{\Psi}_I^-(y) + \{ \Psi_I^+(x), \bar{\Psi}_I^-(y) \} \\ &\equiv \mathfrak{N} \{ \Psi_I(x) \bar{\Psi}_I(y) \} + \{ \Psi_I^+(x), \bar{\Psi}_I^-(y) \} \end{aligned}$$

Similarly, for $y^0 > x^0$:

$$\begin{aligned} \mathfrak{T} \{ \Psi_I(x) \bar{\Psi}_I(y) \} &= -\bar{\Psi}_I^+(y) \Psi_I^+(x) - \bar{\Psi}_I^-(y) \Psi_I^+(x) - \bar{\Psi}_I^+(y) \Psi_I^-(x) - \bar{\Psi}_I^-(y) \Psi_I^-(x) \\ &= \Psi_I^+(x) \bar{\Psi}_I^+(y) - \bar{\Psi}_I^-(y) \Psi_I^+(x) + \Psi_I^-(x) \bar{\Psi}_I^+(y) + \Psi_I^-(x) \bar{\Psi}_I^-(y) - \{ \bar{\Psi}_I^+(y), \Psi_I^-(x) \} \\ &\equiv \mathfrak{N} \{ \Psi_I(x) \bar{\Psi}_I(y) \} - \{ \bar{\Psi}_I^+(y), \Psi_I^-(x) \} \end{aligned}$$

Therefore, in general^a:

$$\begin{aligned} \mathfrak{T} \{ \Psi_I(x) \bar{\Psi}_I(y) \} &= \mathfrak{N} \{ \Psi_I(x) \bar{\Psi}_I(y) \} + \theta(x^0 - y^0) \{ \Psi_I^+(x), \bar{\Psi}_I^-(y) \} - \theta(y^0 - x^0) \{ \bar{\Psi}_I^+(y), \Psi_I^-(x) \} \\ &= \mathfrak{N} \{ \Psi_I(x) \bar{\Psi}_I(y) \} + S(x - y) \end{aligned}$$

By the same observations as in the previous proof $\langle 0 | \mathfrak{T} \{ \Psi_I(x) \bar{\Psi}_I(y) \} | 0 \rangle = S(x - y)$, so explicitly computing the anti-commutators:

$$\begin{aligned} S(x - y) &= \int \frac{d^3 p d^3 q}{(2\pi)^6 \sqrt{4E_{\mathbf{p}} E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} [\theta(x^0 - y^0) \{ a_{\mathbf{p},s}, a_{\mathbf{q},r}^\dagger \} u^s(p) \bar{u}^r(q) e^{-ip_\mu x^\mu + iq_\mu y^\mu} + \\ &\quad - \theta(y^0 - x^0) \{ b_{\mathbf{p},s}, b_{\mathbf{q},r}^\dagger \} v^r(q) \bar{v}^s(p) e^{-ip_\mu x^\mu + iq_\mu y^\mu}] \\ &= \int \frac{d^3 p d^3 q}{(2\pi)^3 \sqrt{4E_{\mathbf{p}} E_{\mathbf{q}}}} \sum_{s=1,2} \sum_{r=1,2} \delta^{(3)}(\mathbf{p} - \mathbf{q}) \delta_{sr} [\theta(x^0 - y^0) u^s(p) \bar{u}^r(q) e^{-ip_\mu x^\mu + iq_\mu y^\mu} + \\ &\quad - \theta(y^0 - x^0) v^r(q) \bar{v}^s(p) e^{-ip_\mu x^\mu + iq_\mu y^\mu}] \\ &= \int \frac{d^3 p}{(2\pi)^3 2E_{\mathbf{p}}} \sum_{s=1,2} [\theta(x^0 - y^0) u^s(p) \bar{u}^s(p) e^{-ip_\mu(x-y)^\mu} - \theta(y^0 - x^0) v^s(p) \bar{v}^s(p) e^{ip_\mu(x-y)^\mu}] \end{aligned}$$

Recalling Eq. 1.111:

$$S(x - y) = \int \frac{d^3 p}{(2\pi)^3 2E_{\mathbf{p}}} [\theta(x^0 - y^0) (\not{p} + m) e^{-ip_\mu(x-y)^\mu} - \theta(y^0 - x^0) (\not{p} - m) e^{ip_\mu(x-y)^\mu}]$$

By the same technique in the previous proof, to show that this expression is equal to the thesis, rewrite Eq. 3.52 as:

$$S(x-y) = \int \frac{d^3p}{(2\pi)^3} e^{i\mathbf{p}\cdot(\mathbf{x}-\mathbf{y})} \int \frac{dp^0}{2\pi} \frac{i(p^0\gamma^0 - \mathbf{p}\cdot\boldsymbol{\gamma} + m)}{(p^0)^2 - E_{\mathbf{p}}^2 + i\epsilon} e^{-ip^0(x^0-y^0)}$$

The $i\epsilon$ -prescription shifts the pole in $p^0 > 0$ below the real axis and that in $p^0 < 0$ above it. If $x^0 > y^0$, chose a semicircular (clockwise) contour in the lower-half plane, so that only the $p^0 = E_{\mathbf{p}}$ singularity is enclosed:

$$\begin{aligned} S(x-y)|_{x^0 > y^0} &= \int \frac{d^3p}{(2\pi)^3} e^{i\mathbf{p}\cdot(\mathbf{x}-\mathbf{y})} \frac{i}{2\pi} \left[(-2\pi i) \frac{E_{\mathbf{p}}\gamma^0 - \mathbf{p}\cdot\boldsymbol{\gamma} + m}{2E_{\mathbf{p}}} e^{-iE_{\mathbf{p}}(x^0-y^0)} \right] \\ &= \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} e^{ip_{\mu}\cdot(x-y)^{\mu}} (E_{\mathbf{p}}\gamma^0 - \mathbf{p}\cdot\boldsymbol{\gamma} + m) = \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} e^{ip_{\mu}\cdot(x-y)^{\mu}} (\not{p} + m) \end{aligned}$$

If instead $y^0 > x^0$, chose a semicircular (counter-clockwise) contour in the upper-half plane:

$$\begin{aligned} S(x-y)|_{x^0 < y^0} &= \int \frac{d^3p}{(2\pi)^3} e^{i\mathbf{p}\cdot(\mathbf{x}-\mathbf{y})} \frac{i}{2\pi} \left[(2\pi i) \frac{-E_{\mathbf{p}}\gamma^0 - \mathbf{p}\cdot\boldsymbol{\gamma} + m}{-2E_{\mathbf{p}}} e^{iE_{\mathbf{p}}(x^0-y^0)} \right] \\ &= \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} e^{ip_{\mu}\cdot(x-y)^{\mu}} (-E_{\mathbf{p}}\gamma^0 + \mathbf{p}\cdot\boldsymbol{\gamma} + m) \\ &= - \int \frac{d^3p}{(2\pi)^3 2E_{\mathbf{p}}} e^{ip_{\mu}\cdot(x-y)^{\mu}} (\not{p} - m) \end{aligned}$$

where in the second equality $\mathbf{p} \mapsto -\mathbf{p}$ was renamed. The proof is complete. \square

^aTo be precise, all these quantities should have spinor indices: indeed, $\Psi\bar{\Psi} \in \mathbb{C}^{4\times 4}$ while $\bar{\Psi}\Psi \in \mathbb{C}$, so the anti-commutator $\{\Psi, \bar{\Psi}\}$ is only defined with spinor indices, i.e. $\{\Psi_{\alpha}, \bar{\Psi}_{\beta}\}$. To shorten the notation, spinor indices are suppressed, and $\{\Psi, \bar{\Psi}\}, \{\bar{\Psi}, \Psi\}$ are treated as $\mathbb{C}^{4\times 4}$ matrices, always writing Dirac adjoints on the right in the explicit expressions.

Note that $S(x-y) \in \mathbb{C}^{4\times 4}$, with two Dirac indices: $S_{\alpha\beta}(x-y) \equiv \langle 0 | \mathfrak{T} \{ \Psi_{I,\alpha}(x) \bar{\Psi}_{I,\beta}(y) \} | 0 \rangle$. It is straightforward to see that the Feynman propagator in momentum space is:

$$\tilde{S}(x-y) = \frac{i(\not{p} + m)}{p^2 - m^2 + i\epsilon} \quad (3.53)$$

Lemma 3.3.1 (Dirac propagator in momentum space)

$$\tilde{S}(x-y) = \frac{i}{\not{p} - m} \quad (3.54)$$

Proof. $(\not{p} + m)(\not{p} - m) = \gamma^{\mu}\gamma^{\nu}p_{\mu}p_{\nu} - m^2 = \frac{1}{2}\{\gamma^{\mu}, \gamma^{\nu}\}p_{\mu}p_{\nu} - m^2 = \eta^{\mu\nu}p_{\mu}p_{\nu} - m^2 = p^2 - m^2. \quad \square$

This is interpreted as the inverse matrix $i(\not{p} - m)^{-1}$. As for the Feynman propagator, the Dirac propagator too is a Green function for the Dirac operator:

$$(i\not{\partial}_x - m)S(x-y) = \int \frac{d^4p}{(2\pi)^4} \frac{i}{\not{p} - m} (\gamma^{\mu}p_{\mu} - m) e^{-ip_{\mu}(x-y)^{\mu}} = i\delta^{(4)}(x-y)$$

§3.3.2 Wick's theorem

It is possible to simplify the computation of a general n -point Green function in terms of 2-point Green functions. In the following, the subscript I for the interaction-picture field is omitted.

§3.3.2.1 Bosonic fields

Definition 3.3.2 (Contraction)

Given a bosonic field $\phi(x)$, its **contraction** is defined as:

$$\overline{\phi(x)\phi(y)} := \begin{cases} [\phi^+(x), \phi^-(y)] & x^0 > y^0 \\ [\phi^+(y), \phi^-(x)] & y^0 > x^0 \end{cases} \quad (3.55)$$

Recalling the proof of Th. 3.3.1, it is clear that:

$$\mathfrak{T}\{\phi(x)\phi(y)\} = \mathfrak{N}\{\phi(x)\phi(y)\} + \overline{\phi(x)\phi(y)} \quad (3.56)$$

This relation can be generalized to n -point Green functions.

Theorem 3.3.3 (Wick's theorem)

Given a bosonic field $\phi(x)$:

$$\mathfrak{T}\{\phi(x_1) \dots \phi(x_n)\} = \mathfrak{N}\{\phi(x_1) \dots \phi(x_n) + \text{all possible contractions}\} \quad (3.57)$$

where both partially-contracted and fully-contracted terms are considered.

Proof. Use induction on n . Eq. 3.56 proves the $n = 2$ case, so assume the theorem true for generic $n - 1$ and consider the n case. WLOG $x_1^0 \geq \dots \geq x_n^0$ (if that's not the case, just relabel the points, as it does not affect Eq. 3.57), so, using the inductive step:

$$\begin{aligned} \mathfrak{T}\{\phi(x_1) \dots \phi(x_n)\} &= \phi(x_1) \mathfrak{T}\{\phi(x_2) \dots \phi(x_n)\} \\ &= (\phi^+(x_1) + \phi^-(x_1)) \mathfrak{N}\{\phi(x_2) \dots \phi(x_n) + \text{all contr. without } \phi(x_1)\} \end{aligned}$$

$\phi^-(x_1)$ is already in normal order (on the left), while $\phi^+(x_1)$ must be commuted with all other $\phi(x_i)$. For the term without contractions:

$$\begin{aligned} \phi^+(x_1) \mathfrak{N}\{\phi(x_2) \dots \phi(x_n)\} &= \mathfrak{N}\{\phi(x_2) \dots \phi(x_n)\} \phi^+(x_1) + [\phi^+(x_1), \mathfrak{N}\{\phi(x_2) \dots \phi(x_n)\}] \\ &= \mathfrak{N}\{\phi^+(x_1) \phi(x_2) \dots \phi(x_n)\} \\ &\quad + \mathfrak{N}\{\dots + \phi(x_2) \dots [\phi^+(x_2), \phi^-(x_i)] \dots \phi(x_n) + \dots\} \\ &= \mathfrak{N}\{\phi^+(x_1) \phi(x_2) \dots \phi(x_n) + \dots \overline{\phi^+(x_1) \phi(x_2) \dots \phi(x_i)} \dots \phi(x_n) + \dots\} \end{aligned}$$

Analogously, all possible contractions are obtained, proving the theorem. \square

Note that only fully-contracted terms contribute to the n -point Green function: every term with at least an uncontracted field vanishes when put between $\langle 0 | \cdot | 0 \rangle$, as there will be an annihilation operator acting on $|0\rangle$ or a creation operator acting on $\langle 0|$.

Example 3.3.1 (4-point Green function)

Applying Wick's theorem to a 4-point Green function yields:

$$\begin{aligned} \mathfrak{T}\{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)\} = & \mathfrak{N}\{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4) + \\ & \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} \\ & + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} \\ & + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} + \overbrace{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)} \end{aligned}$$

Only the last three terms contribute to the 4-point Green function. If $\phi(x)$ is a scalar field, then:

$$\begin{aligned} \langle 0 | \mathfrak{T}\{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)\} | 0 \rangle = & \Delta(x_1 - x_2)\Delta(x_3 - x_4) + \Delta(x_1 - x_3)\Delta(x_2 - x_4) + \\ & + \Delta(x_1 - x_4)\Delta(x_2 - x_3) \end{aligned}$$

This expression has immediate physical interpretation: indeed, $\Delta(x_1 - x_2)$ can be interpreted as the amplitude for the propagation of a particle from spacetime point x_1 to x_2 , while $\Delta(x_1 - x_2)\Delta(x_3 - x_4)$ is the amplitude for the propagation of a particle from x_1 to x_2 and one from x_3 to x_3 , without interacting with each other.

It is possible to associate intuitive graphs, called **Feynman diagrams**, to these kind of expressions; for example:

$$\Delta(x_1 - x_2)\Delta(x_3 - x_4) = \begin{array}{c} x_1 \bullet \text{-----} \bullet x_2 \\ x_3 \bullet \text{-----} \bullet x_4 \end{array}$$

Then, the 4-point Green function can be represented as:

$$\mathfrak{T}\{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)\} = \begin{array}{c} x_1 \bullet \text{-----} \bullet x_2 \\ x_3 \bullet \text{-----} \bullet x_4 \end{array} + \begin{array}{c} x_1 \bullet \\ | \\ x_3 \bullet \end{array} \begin{array}{c} x_2 \bullet \\ | \\ x_4 \bullet \end{array} + \begin{array}{c} x_1 \bullet \text{-----} \bullet x_2 \\ \diagdown \quad \diagup \\ x_3 \bullet \text{-----} \bullet x_4 \end{array}$$

Using this formalism to compute Eq. 3.48, when expanding the exponential in powers of \mathcal{H}_I , each term contains fields at the same spacetime point: this gives rise to less trivial Feynman diagrams.

§3.3.2.2 Fermionic fields**Definition 3.3.3** (Contraction)

Given a fermionic $\Psi(x)$, its **contraction** is defined as:

$$\overline{\Psi(x)}\bar{\Psi}(y) := \begin{cases} \{\Psi^+(x), \bar{\Psi}^-(y)\} & x^0 > y^0 \\ -\{\bar{\Psi}^+(y), \Psi^-(x)\} & y^0 > x^0 \end{cases} \quad (3.58)$$

$$\text{and } \overline{\Psi(x)}\Psi(y) = \bar{\Psi}(x)\bar{\Psi}(y) = 0.$$

By the proof of Th. 3.3.2, it is clear that:

$$\mathfrak{T}\{\Psi(x)\bar{\Psi}(y)\} = \mathfrak{N}\{\Psi(x)\bar{\Psi}(y)\} + \overline{\Psi(x)}\bar{\Psi}(y) \quad (3.59)$$

Wick's theorem remains basically unchanged for fermions.

Theorem 3.3.4 (Wick's theorem)

Given a fermionic field $\Psi(x)$:

$$\mathfrak{T}\{\Psi(x_1)\bar{\Psi}(x_2)\dots\} = \mathfrak{N}\{\Psi(x_1)\bar{\Psi}(x_2)\dots + \text{all possible contractions}\} \quad (3.60)$$

where both partially-contracted and fully-contracted terms are considered.

Proof. Inserting everything inside $\mathfrak{N}\{\cdot\}$ ensures that the correct signs are considered due to fermionic anti-commutation, therefore the proof is analogous to that of Th. 3.3.3. \square

As for the bosonic counterparts, only fully-contracted terms contribute to the n -point Green function, as each partially-contracted term will have either an annihilation operator acting on $|0\rangle$ or a creation operator acting on $\langle 0|$ (or both).

§3.3.2.3 Scattering in $\lambda\phi^4$ -theory

Consider a real scalar field theory with $\mathcal{H}_I = \frac{\lambda}{4!}\phi^4$ and consider a scattering process with two initial particles with momenta $\mathbf{k}_1, \mathbf{k}_2$ and two final particles with momenta $\mathbf{p}_1, \mathbf{p}_2$. Eq. 3.25 and Eq. 3.48 give the amplitude:

$$\begin{aligned} \langle \mathbf{p}_1, \mathbf{p}_2 | iT | \mathbf{k}_1, \mathbf{k}_2 \rangle &= \frac{k_1^2 - m^2}{i\sqrt{Z}} \frac{k_2^2 - m^2}{i\sqrt{Z}} \frac{p_1^2 - m^2}{i\sqrt{Z}} \frac{p_2^2 - m^2}{i\sqrt{Z}} \int \prod_{i=1}^4 d^4x_i e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - p_{3\mu}x_3^\mu - p_{4\mu}x_4^\mu)} \times \\ &\quad \times \frac{\langle 0 | \mathfrak{T}\{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)\exp[-i\frac{\lambda}{4!}\int d^4x\phi(x)]\} | 0 \rangle}{\langle 0 | \mathfrak{T}\{\exp[-i\frac{\lambda}{4!}\int d^4x\phi(x)]\} | 0 \rangle} \end{aligned}$$

To compute this expression up to $o(\lambda)$, consider that for $\lambda\phi^4$ -theory $Z = 1 + o(\lambda^2)$, thus it is safe to set $Z \equiv 1$.

The $o(\lambda^0)$ term is obtained setting $\lambda = 0$, i.e. no coupling, but in the absence of coupling there is no scattering either and the amplitude must be trivial. Recalling Ex. 3.3.1:

$$\begin{aligned} &\int d^4x_1 d^4x_2 d^4x_3 d^4x_4 e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - p_{3\mu}x_3^\mu - p_{4\mu}x_4^\mu)} \langle 0 | \mathfrak{T}\{\phi(x_1)\phi(x_2)\phi(x_3)\phi(x_4)\} | 0 \rangle \\ &= \int d^4x_1 d^4x_2 d^4x_3 d^4x_4 e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - p_{3\mu}x_3^\mu - p_{4\mu}x_4^\mu)} [\Delta(x_1 - x_2)\Delta(x_3 - x_4) + \dots] \\ &= \left[\int d^4x d^4X e^{i(p_1 + p_2)_\mu X^\mu + i(p_1 - p_2)_\mu x^\mu / 2} \Delta(x) \right] \left[\int d^4y d^4Y e^{-i(k_1 + k_2)_\mu Y^\mu - i(k_1 - k_2)_\mu y^\mu / 2} \Delta(y) \right] + \dots \\ &= (2\pi)^4 \delta^{(4)}(p_1 + p_2) \frac{i}{\left(\frac{p_1 - p_2}{2}\right)^2 - m^2} (2\pi)^4 \delta^{(4)}(k_1 + k_2) \frac{i}{\left(\frac{k_1 - k_2}{2}\right)^2 - m^2} + \dots \end{aligned}$$

with $x = x_1 - x_2$, $2X = x_1 + x_2$ and $y = x_3 - x_4$, $2Y = x_3 + x_4$. Then, using $f(x)\delta(x - x_0) = f(x_0)\delta(x - x_0)$:

$$\begin{aligned} & \int d^4x_1 d^4x_2 d^4x_3 d^4x_4 e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - p_{3\mu}x_3^\mu - p_{4\mu}x_4^\mu)} \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \} | 0 \rangle \\ &= (2\pi)^4 \delta^{(4)}(p_1 + p_2) (2\pi)^4 \delta^{(4)}(k_1 + k_2) \frac{i}{p_1^2 - m^2} \frac{i}{k_1^2 - m^2} + \dots \end{aligned}$$

Note that the $i\epsilon$ term must be retained while p is an integration variable, as it gives the prescription of going around the poles, but it can be dropped when p is the momentum on an external leg, i.e. fixed. This expression only has two poles, therefore:

$$\langle \mathbf{p}_1, \mathbf{p}_2 | iT | \mathbf{k}_1, \mathbf{k}_2 \rangle = -(2\pi)^8 \delta^{(4)}(p_1 + p_2) \delta^{(4)}(k_1 + k_2) (p_2^2 - m^2)(k_2^2 - m^2) + \dots + o(\lambda)$$

These terms vanish when going on mass-shell, confirming that at $o(\lambda^0)$ there is no interaction, i.e. no contribution to the scattering amplitude. This is a general feature of $n \rightarrow m$ scattering amplitudes: *disconnected graphs do not contribute to the amplitude*, as they do not provide enough pole factors to cancel those in the LSZ formula.

The first non-trivial contribution is that at $o(\lambda)$. The numerator of Eq. 3.48 is:

$$\int \prod_{i=1}^4 d^4x_i e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - p_{3\mu}x_3^\mu - p_{4\mu}x_4^\mu)} \left(-i \frac{\lambda}{4!} \right) \int d^4x \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \phi^4(x) \} | 0 \rangle$$

The only non-vanishing contribution is the fully-contracted term:

$$\langle 0 | \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \phi(x) \phi(x) \phi(x) \phi(x) | 0 \rangle = \begin{array}{c} x_1 \quad \quad \quad x_2 \\ \quad \quad \quad \diagdown \quad \diagup \\ \quad \quad \quad x \\ \quad \quad \quad \diagup \quad \diagdown \\ x_3 \quad \quad \quad x_4 \end{array}$$

as it is the only connected diagram between the possible ones. Note that this diagram corresponds to $4!$ equal contractions (only one is shown), therefore the numerator becomes:

$$\begin{aligned} & -i\lambda \int d^4x \prod_{i=1}^4 d^4x_i e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - p_{3\mu}x_3^\mu - p_{4\mu}x_4^\mu)} \Delta(x_1 - x) \Delta(x_2 - x) \Delta(x_3 - x) \Delta(x_4 - x) \\ &= -i\lambda \tilde{\Delta}(p_1) \tilde{\Delta}(p_2) \tilde{\Delta}(k_1) \tilde{\Delta}(k_2) \int d^4x e^{i(p_1 + p_2 - k_1 - k_2)_\mu x^\mu} \\ &= -i\lambda (2\pi)^4 \delta^{(4)}(p_1 + p_2 - k_1 - k_2) \frac{i}{p_1^2 - m^2} \frac{i}{p_2^2 - m^2} \frac{i}{k_1^2 - m^2} \frac{i}{k_2^2 - m^2} \end{aligned}$$

There remains to evaluate the denominator: this term gives only **vacuum-to-vacuum transitions**, i.e. diagrams with no external legs. To account for these, note that each connected diagram which contributes to the numerator can be “dressed” with all possible vacuum-to-vacuum transitions; indeed, to be precise:

$$\begin{aligned} & \langle 0 | \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \phi(x) \phi(x) \phi(x) \phi(x) | 0 \rangle = \\ &= \begin{array}{c} x_1 \quad \quad \quad x_2 \\ \quad \quad \quad \diagdown \quad \diagup \\ \quad \quad \quad x \\ \quad \quad \quad \diagup \quad \diagdown \\ x_3 \quad \quad \quad x_4 \end{array} + \begin{array}{c} \text{bubble} \\ x_1 \quad \quad \quad x_2 \\ \quad \quad \quad \diagdown \quad \diagup \\ \quad \quad \quad x \\ \quad \quad \quad \diagup \quad \diagdown \\ x_3 \quad \quad \quad x_4 \end{array} + \begin{array}{c} \text{two bubbles} \\ x_1 \quad \quad \quad x_2 \\ \quad \quad \quad \diagdown \quad \diagup \\ \quad \quad \quad x \\ \quad \quad \quad \diagup \quad \diagdown \\ x_3 \quad \quad \quad x_4 \end{array} + \dots \end{aligned}$$

$$= \begin{array}{c} x_1 \bullet \\ \quad \diagdown \quad \diagup \\ \quad \quad x \\ \quad \diagup \quad \diagdown \\ x_3 \bullet \quad x_4 \bullet \end{array} \times \left[1 + \text{diagram with two circles sharing a vertex} + \text{diagram with three circles sharing a vertex} + \dots \right]$$

This is nothing but the perturbative expansion of the denominator of Eq. 3.48, therefore they cancel out leaving only the already computed diagram³. Thus, the transition amplitude at $o(\lambda)$ is:

$$\langle \mathbf{p}_1, \mathbf{p}_2 | iT | \mathbf{k}_1, \mathbf{k}_2 \rangle = -i\lambda(2\pi)^4 \delta^{(4)}(p_1 + p_2 - k_1 - k_2) \quad (3.61)$$

§3.3.2.4 Scattering in $\lambda\phi^3$ -theory

To illustrate the case of internal lines, it is useful to consider a real scalar field theory with $\mathcal{H}_I = \frac{\lambda}{3!}\phi^3$: in this case, 3 lines meet at each vertex, instead of 4. The $2 \rightarrow 2$ scattering amplitude is:

$$\begin{aligned} \langle \mathbf{p}_1, \mathbf{p}_2 | iT | \mathbf{k}_1, \mathbf{k}_2 \rangle &= \prod_{i=1,2} \frac{p_i^2 - m^2}{i\sqrt{Z}} \prod_{j=1,2} \frac{k_j^2 - m^2}{i\sqrt{Z}} \int \prod_{i=1}^4 d^4 x_i e^{i(p_{1\mu} x_1^\mu + p_{2\mu} x_2^\mu - k_{1\mu} x_3^\mu - k_{4\mu} x_4^\mu)} \times \\ &\quad \times \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \exp \left[-i \frac{\lambda}{3!} \int d^4 x \phi^3(x) \right] \} | 0 \rangle \end{aligned}$$

where the denominator was omitted by implicitly considering only connected graphs with 4 external legs. This amplitude does not have a $o(\lambda^0)$ contribution (as it only has disconnected diagrams) nor an $o(\lambda)$ contribution (as the product of an odd number of fields cannot be fully contracted). At $o(\lambda^2)$ the integral becomes:

$$\begin{aligned} \langle \mathbf{p}_1, \mathbf{p}_2 | iT | \mathbf{k}_1, \mathbf{k}_2 \rangle &= \frac{p_1^2 - m^2}{i\sqrt{Z}} \frac{p_2^2 - m^2}{i\sqrt{Z}} \frac{k_1^2 - m^2}{i\sqrt{Z}} \frac{k_2^2 - m^2}{i\sqrt{Z}} \int \prod_{i=1}^4 d^4 x_i e^{i(p_{1\mu} x_1^\mu + p_{2\mu} x_2^\mu - k_{1\mu} x_3^\mu - k_{4\mu} x_4^\mu)} \times \\ &\quad \times \frac{1}{2!} \left(-i \frac{\lambda}{3!} \right)^2 \int d^4 x d^4 y \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \phi^3(x) \phi^3(y) \} | 0 \rangle \end{aligned}$$

In order to obtain a fully connected diagrams, there needs to be a contraction $\overline{\phi(x)}\phi(y)$ between the two vertices, creating an internal line. Note that the $1/2!$ factor accounts for the symmetry

³This too is a general statement. Consider a generic process with possible connected diagrams $\{C_i\}_{i \in \mathcal{I}}$ and possible vacuum-to-vacuum diagrams $\{V_j\}_{j \in \mathcal{J}}$. Then, as shown above, the contribution of a connected diagram C_i is really:

$$C_i \times \sum_{\{n_j\} \subset \mathbb{N}_0} \prod_{j \in \mathcal{J}} \frac{1}{n_j!} V_j^{n_j}$$

with $\{n_j\}_{j \in \mathcal{J}} \subset \mathbb{N}_0$ (the symmetry factor is due to the swap of identical diagrams). The sum of all non-vanishing diagrams thus is:

$$\sum (\text{diagrams}) = \sum_{i \in \mathcal{I}} C_i \times \sum_{\{n_j\} \subset \mathbb{N}_0} \prod_{j \in \mathcal{J}} \frac{1}{n_j!} V_j^{n_j} = \sum_{i \in \mathcal{I}} C_i \times \prod_{j \in \mathcal{J}} \sum_{n=1}^{\infty} \frac{1}{n!} V_j^n = \sum_{i \in \mathcal{I}} C_i \times \prod_{j \in \mathcal{J}} \exp V_j = \sum_{i \in \mathcal{I}} C_i \times \exp \sum_{j \in \mathcal{J}} V_j$$

which can be recast in the form:

$$\sum (\text{diagrams}) = \sum (\text{connected}) \times \exp \sum (\text{vacuum-to-vacuum}) \quad (3.62)$$

The exponential exactly cancels $\langle 0 | \mathfrak{T} \{ \exp [-i \int d^4 x \mathcal{H}_I] \} | 0 \rangle$, leaving only connected diagrams in the calculation.

$x \leftrightarrow y$, while $1/3!$ accounts for sets of equivalent contractions. Examples of possible diagrams are:

$$\begin{array}{c}
 \begin{array}{ccc}
 x_1 & & x_3 \\
 & \searrow & \nearrow \\
 & x & y \\
 & \nearrow & \searrow \\
 x_2 & & x_4
 \end{array}
 \end{array}
 = \Delta(x_1 - x)\Delta(x_2 - x)\Delta(x - y)\Delta(y - x_3)\Delta(y - x_4)$$

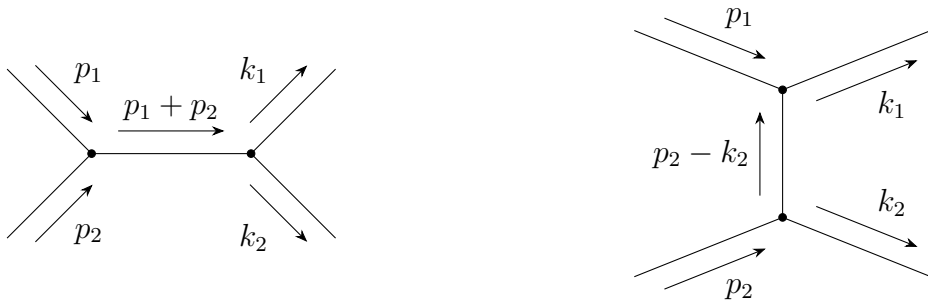
$$\begin{array}{c}
 \begin{array}{ccc}
 x_1 & & x_3 \\
 & \searrow & \nearrow \\
 & y & \\
 & \nearrow & \searrow \\
 x_2 & & x_4
 \end{array}
 \end{array}
 = \Delta(x_1 - y)\Delta(x_2 - x)\Delta(x - y)\Delta(y - x_3)\Delta(x - x_4)$$

These diagrams are inequivalent, as they correspond to different sets of contractions. Computing for example the first one:

$$\begin{aligned}
 & (-i\lambda)^2 \int \prod_{i=1}^4 d^4 x_i e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - k_{1\mu}x_3^\mu - k_{4\mu}x_4^\mu)} \times \\
 & \quad \times \int d^4 x d^4 y \Delta(x_1 - x)\Delta(x_2 - x)\Delta(x - y)\Delta(y - x_3)\Delta(y - x_4) \\
 & = (-i\lambda)^2 \tilde{\Delta}(p_1)\tilde{\Delta}(p_2)\tilde{\Delta}(k_1)\tilde{\Delta}(k_2) \int d^4 x d^4 y e^{i(p_1+p_2)_\mu x^\mu - i(k_1+k_2)_\mu y^\mu} \Delta(x - y) \\
 & = (-i\lambda)^2 \tilde{\Delta}(p_1)\tilde{\Delta}(p_2)\tilde{\Delta}(k_1)\tilde{\Delta}(k_2) \int d^4 X d^4 Y e^{i(p_1+p_2-k_1-k_2)_\mu Y^\mu + i(p_1+p_2+k_1+k_2)_\mu X^\mu/2} \Delta(X) \\
 & = (-i\lambda)^2 (2\pi)^4 \delta^{(4)}(p_1 + p_2 - k_1 - k_2) \tilde{\Delta}(p_1)\tilde{\Delta}(p_2)\Delta(p_1 + p_2)\tilde{\Delta}(k_1)\tilde{\Delta}(k_2)
 \end{aligned}$$

with $X = x - y$, $2Y = x + y$. As expected, momentum-space propagators associated to external legs are canceled by pole terms in the LSZ formula: now there remains only the $\tilde{\Delta}(p_1 + p_2)$ factor associated to the internal line.

It is useful to work in momentum space, rather than coordinate space. For example, in momentum space the above diagrams become:



§3.3.2.5 General considerations

The example calculations carried for $\lambda\phi^3$ - and $\lambda\phi^4$ -theory hint to the general technique to compute general **tree diagrams**, i.e. connected diagrams without internal loops. For a scalar field theory, in momentum space:

1. draw all possible topologically-inequivalent connected graphs corresponding to the given initial and final states and multiply them by their symmetry factor;
2. each external leg is associated to a pole factor in the LSZ reduction formula, hence they can be omitted directly (“leg amputation”);
3. as there is always a momentum-conserving δ -function, it is convenient to define the matrix element \mathcal{M} for the process $|\text{in}\rangle \rightarrow |\text{out}\rangle$ as:

$$\langle \mathbf{p}_1, \dots, \mathbf{p}_m | iT | \mathbf{k}_1, \dots, \mathbf{k}_n \rangle \equiv (2\pi)^4 \delta^{(4)} \left(\sum_{i=1}^m p_i - \sum_{j=1}^n k_j \right) i\mathcal{M} \quad (3.63)$$

4. energy-momentum conservation must be imposed separately at each vertex;
5. each vertex is associated to a $-i\lambda$ factor, where λ is the coupling constant of the interaction;
6. check if that symmetry factors cancel fully or partially with the $1/n!$ factor at $o(\lambda^n)$ and possible numerical factors in the definition of \mathcal{H}_I .

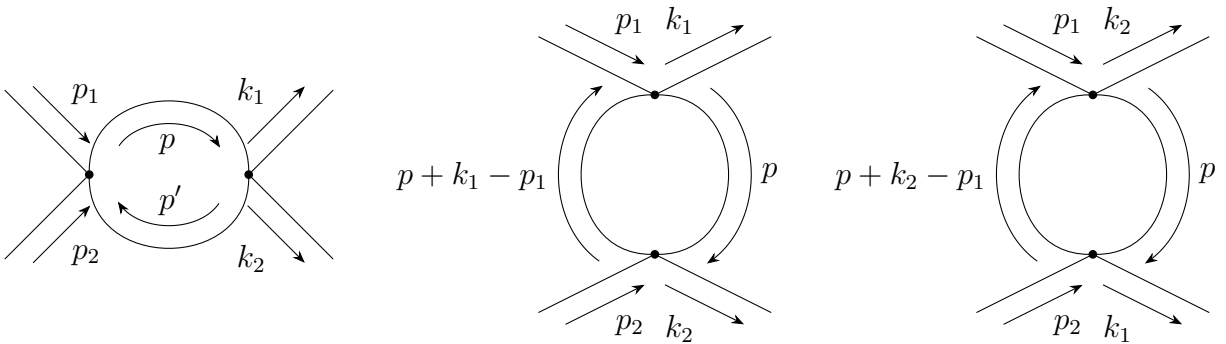
The fact that the considered diagrams are topologically inequivalent needs to be stressed out. Moreover, symmetry factors need to be computed carefully, as illustrated in §4.2.1.3.

§3.3.3 Loop diagrams

To illustrate how loop diagrams are treated, consider the $o(\lambda^2)$ corrections to the $2 \rightarrow 2$ scattering in $\lambda\phi^4$ -theory:

$$\begin{aligned} & \frac{1}{2!} \left(-i\frac{\lambda}{4!} \right)^2 \int \prod_{i=1}^4 d^4 x_i e^{i(p_{1\mu}x_1^\mu + p_{2\mu}x_2^\mu - k_{1\mu}x_3^\mu - k_{4\mu}x_4^\mu)} \times \\ & \times \int d^4 x d^4 y \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \phi^4(x) \phi^4(y) \} | 0 \rangle \end{aligned}$$

There are three possible inequivalent full contractions, associated to the connected diagrams:



with $p' = p_1 + p_2 - p$. It is clear that energy conservation does not completely fix momenta of internal lines, and in particular of loops. Consider for example the first diagram: it is obtained contracting $\phi(x_1)$ and $\phi(x_2)$ with two $\phi(x)$ ($4 \cdot 3$ ways), $\phi(x_3)$ and $\phi(x_4)$ with two $\phi(y)$ ($4 \cdot 3$ ways)

and two $\phi(x)$ with two $\phi(y)$ (2 ways), so, considering another 2 factor for $x \leftrightarrow y$ symmetry, this diagram has symmetry number $S = (4!)^2$. The resulting contribution is:

$$\begin{aligned} & \frac{(-i\lambda)^2}{2} \int \prod_{i=1}^4 d^4 x_i e^{i(p_{1\mu} x_1^\mu + p_{2\mu} x_2^\mu - k_{1\mu} x_3^\mu - k_{4\mu} x_4^\mu)} \int d^4 x d^4 y \Delta(x_1 - x) \Delta(x_2 - x) \Delta(x - y) \Delta(y - x) \\ &= \frac{(-i\lambda)^2}{2} \tilde{\Delta}(p_1) \tilde{\Delta}(p_2) \tilde{\Delta}(k_1) \tilde{\Delta}(k_2) \int d^4 x d^4 y e^{i(p_1 + p_2)_\mu x^\mu - i(k_1 + k_2)_\mu y^\mu} \Delta^2(x - y) \\ &= \frac{(-i\lambda)^2}{2} \tilde{\Delta}(p_1) \tilde{\Delta}(p_2) \tilde{\Delta}(k_1) \tilde{\Delta}(k_2) (2\pi)^4 \delta^{(4)}(p_1 + p_2 - k_1 - k_2) \int \frac{d^4 k}{(2\pi)^4} \tilde{\Delta}(k) \tilde{\Delta}(p_1 + p_2 - k) \end{aligned}$$

Therefore, the additional feature of loop integrals is the integration over unfixed momenta, assigning the right propagators to each internal line. It is possible to define the **loop-amplitude** for $\lambda\phi^4$ -theory:

$$\mathcal{A}(p) := \frac{(-i\lambda)^2}{2} \int \frac{d^4 k}{(2\pi)^4} \frac{i}{k^2 - m^2 + i\epsilon} \frac{i}{(p - k)^2 - m^2 + i\epsilon} \quad (3.64)$$

so that the $2 \rightarrow 2$ scattering amplitude at one-loop level (i.e. $\mathcal{O}(\lambda^2)$) is:

$$i\mathcal{M}_{2 \rightarrow 2} = -i\lambda + \mathcal{A}(p_1 + p_2) + \mathcal{A}(p_1 - k_1) + \mathcal{A}(p_1 - k_2) \quad (3.65)$$

§3.3.3.1 Divergences

The integral in Eq. 3.64 diverges at large k , which is an example of **UV divergence**. To study this divergence, first of all note that $(p - k)^2 \rightarrow k^2$ as $k \rightarrow \infty$, so that it is safe to set $p = 0$. Due to the $i\epsilon$ -prescription, in the complex k^0 -plane the pole in $k^0 > 0$ is below the real axis, while that in $k^0 < 0$ is above it: the integration path can then be rotated counterclockwise from the real axis to the imaginary axis by $k^0 \mapsto ik^0$, which is called a **Wick rotation**. The amplitude then becomes:

$$\mathcal{A}(0) = \frac{(-i\lambda)^2}{2} \int \frac{d^4 k}{(2\pi)^4} \left(\frac{i}{k^2 - m^2 + i\epsilon} \right)^2 \mapsto i \frac{\lambda^2}{2} \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k^2 + m^2)^2}$$

where $k^2 = k_0^2 + \mathbf{k}^2$ is now a Euclidean momentum. To regulate the UV divergence, a cutoff $\Lambda : k^2 < \Lambda^2$ is introduced, so that:

$$\begin{aligned} \mathcal{A}(0) &= i \frac{\lambda^2}{2} \int \frac{d^4 k}{(2\pi)^4} \frac{1}{k^4} + \text{finite terms} = i \frac{\lambda^2}{2} \frac{1}{(2\pi)^4} 2\pi^2 \int^\Lambda \frac{dk}{k} + \text{finite terms} \\ &= \frac{i\lambda^2}{16\pi^2} \log \Lambda + \text{finite terms} \end{aligned}$$

where $2\pi^2$ is the solid angle in \mathbb{R}^4 . It can be shown that, for general p , the singular part of $\mathcal{A}(p)$ does not depend on p ; hence, it can be written:

$$i\mathcal{M}_{2 \rightarrow 2} = -i\lambda + i\lambda^2 (\beta_0 \log \Lambda + \text{finite terms})$$

where:

$$\beta_0 \equiv \frac{3}{16\pi^2} \quad (3.66)$$

Another example of loop divergence in $\lambda\phi^4$ -theory is that associated to the 2-point Green function at $\mathcal{O}(\lambda^1)$:

$$\langle 0 | \mathcal{T} \{ \phi(x_1) \phi(x_2) \phi^4(x) \} | 0 \rangle = 4 \cdot 3 \cdot \overbrace{\phi(x_1) \phi(x_2) \phi(x) \phi(x) \phi(x) \phi(x)} = 4 \cdot 3 \times \text{tadpole diagram}$$

This diagram is typically called a **tadpole**. Its contribution can be computed in momentum performing a Wick rotation:

$$-iB \equiv \frac{-i\lambda}{2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2 + m^2} = -i \frac{\lambda}{32\pi^2} \left(\Lambda^2 - m^2 \log \frac{\Lambda^2 + m^2}{m^2} \right) \quad (3.67)$$

Note that the tadpole is independent of p , and it has both a quadratic and logarithmic UV divergence.

Lemma 3.3.2 (Tadpole resummation)

In $\lambda\phi^4$ -theory, for a $|\mathbf{p}\rangle \rightarrow |\mathbf{p}\rangle$ scattering, tadpole diagrams collectively contribute as:

$$= \frac{i}{p^2 - m^2 - B}$$

Proof. The represented diagram series reduces to a geometric series in momentum space:

$$\begin{aligned} & \tilde{\Delta}(p) + \tilde{\Delta}(p)(-iB)\tilde{\Delta}(p) + \tilde{\Delta}(p)(-iB)\tilde{\Delta}(p)(-iB)\tilde{\Delta}(p) + \dots \\ &= \tilde{\Delta}(p) \sum_{n=0}^{\infty} \left[-iB\tilde{\Delta}(p) \right]^n = \tilde{\Delta} \frac{1}{1 + iB\tilde{\Delta}(p)} = \frac{i}{p^2 - m^2} \left[1 - \frac{B}{p^2 - m^2} \right]^{-1} = \frac{i}{p^2 - m^2 - B} \end{aligned}$$

which is the thesis. \square

This shows that accounting for tadpole diagrams is equivalent to shifting the mass $m^2 \mapsto m^2 + B$.

§3.3.4 Feynman rules

To compute amplitudes up to a desired order, it is not necessary to explicitly expand the exponential each time and then perform Wick contractions, as these steps are summarized by a simple set of rules.

§3.3.4.1 $\lambda\phi^n$ -theory

Consider a scalar field theory with Hamiltonian:

$$\mathcal{H} = \mathcal{H}_{\text{KG}} + \frac{\lambda}{n!} \phi^n \quad (3.68)$$

The Feynman rules in momentum space for this theory are:

1. for each internal line: $\bullet \xrightarrow{p} \bullet = \tilde{\Delta}(p)$;
2. for each external leg: $\bullet \text{---} = 1$;
3. for each vertex: $\bullet = -i\lambda$;
4. impose momentum conservation at each vertex;
5. integrate over each undetermined loop momentum p with measure $\frac{d^4 p}{(2\pi)^4}$;
6. divide by the symmetry factor of the diagram.

The problem of computing scattering amplitude is thus reduced to that of drawing every inequivalent fully-connected Feynman diagram (modulo void-to-void diagrams). These rules allow to compute the matrix element \mathcal{M} , which is related to the T -matrix element by Eq. 3.63 and to the S -matrix element by $S = I + iT$.

§3.3.4.2 Yukawa theory

Consider a field theory with a real scalar field $\phi(x)$ and a Dirac field $\Psi(x)$. The **Yukawa Hamiltonian** is:

$$\mathcal{H} = \mathcal{H}_{\text{KG}} + \mathcal{H}_{\text{D}} + g\bar{\Psi}\Psi\phi \quad (3.69)$$

where g is a dimensionless coupling constant. This is a simplified model of QED. Note that the form of \mathcal{H}_{int} allows interaction vertices with two fermion lines and one scalar line. The Feynman rules in momentum space for this theory are:

1. propagators:

$$(a) \text{ scalar: } \bullet \xrightarrow{q} \bullet = \frac{i}{q^2 - m_\phi^2 + i\epsilon} = \tilde{\Delta}(q);$$

$$(b) \text{ fermion: } \bullet \xrightarrow{p} \bullet = \frac{i(\not{p} + m)}{p^2 - m^2 + i\epsilon} = \tilde{S}(p);$$

$$2. \text{ vertex: } \begin{array}{c} \nearrow \\ \bullet \\ \searrow \end{array} \text{---} = -ig;$$

3. external legs:

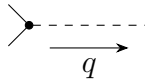
- (a) initial state:

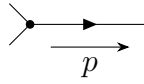
$$i. \text{ scalar: } \begin{array}{c} \nearrow \\ \bullet \\ \searrow \end{array} \text{---} \xleftarrow{q} = 1;$$

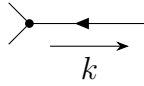
$$ii. \text{ fermion: } \begin{array}{c} \nearrow \\ \bullet \\ \searrow \end{array} \text{---} \xleftarrow{p} = u^s(p);$$

$$iii. \text{ anti-fermion: } \begin{array}{c} \nearrow \\ \bullet \\ \searrow \end{array} \text{---} \xleftarrow{k} = \bar{v}^s(k);$$

(b) final state:

i. scalar:  = 1;

ii. fermion:  = $\bar{u}^s(p)$;

iii. anti-fermion:  = $v^s(k)$;

4. impose momentum conservation at each vertex;

5. integrate over each undetermined loop momentum p with measure $\frac{d^4p}{(2\pi)^4}$ (each closed fermionic loop determines an additional (-1) factor due to anti-commutation);

6. divide by the symmetry factor of the diagram.

Note that for the Yukawa interaction $g\bar{\Psi}\Psi\phi$ the symmetry factor of each diagram is $S = 1$, as the three field in \mathcal{H}_{int} cannot substitute for one another in contractions. Moreover, Dirac indices are always contracted together along fermionic lines.

It can be shown that an interaction mediated by a scalar field is always attractive, both for ff , $f\bar{f}$ and $\bar{f}\bar{f}$ scattering.

§3.3.4.3 Quantum Electrodynamics

Consider a field theory with a Dirac field $\Psi(x)$ and a vector (gauge) field $A_\mu(x)$. The **QED interaction Hamiltonian** is:

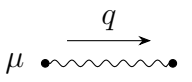
$$\mathcal{H}_{\text{int}} = e\bar{\Psi}\gamma^\mu\Psi A_\mu \quad (3.70)$$

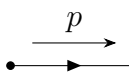
This interaction term determines interaction vertices with two fermionic lines and one vector line. The photon propagator is derived in §4.2.2:

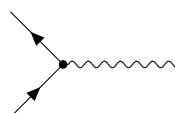
$$\Delta_{\mu\nu}(x-y) := \langle 0 | \mathfrak{T} \{ A_\mu(x) A_\nu(y) \} | 0 \rangle \implies \tilde{\Delta}_{\mu\nu}(k) = \frac{-i}{k^2 + i\epsilon} \eta_{\mu\nu} \quad (3.71)$$

Note that the spatial components $A_i(x)$ have the same propagator as a massless real scalar field, while $A_0(x)$ has an additional negative sign. The Feynman rules in momentum space for QED are:

1. propagators:

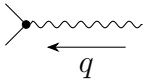
(a) vector:  = $\tilde{\Delta}_{\mu\nu}(q)$;

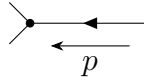
(b) fermion:  = $\tilde{S}(p)$;

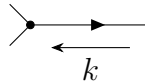
2. vertex:  = $-ie\gamma^\mu$;

3. external legs:

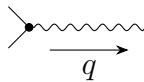
(a) initial state:

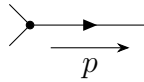
i. vector:  $\mu = \epsilon_\mu(q);$

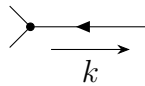
ii. fermion:  $= u^s(p);$

iii. anti-fermion:  $= \bar{v}^s(k);$

(b) final state:

i. scalar:  $\mu = \epsilon_\mu^*(q);$

ii. fermion:  $= \bar{u}^s(p);$

iii. anti-fermion:  $= v^s(k);$

4. impose momentum conservation at each vertex;

5. integrate over each undetermined loop momentum p with measure $\frac{d^4p}{(2\pi)^4}$ (each closed fermionic loop determines an additional (-1) factor due to anti-commutation);

6. divide by the symmetry factor of the diagram.

Of course, when considering fermions with charge Q (in units of $|e|$), substitute $e \mapsto Q|e|$ in the above formulae.

It can be shown that the interaction mediated by a vector field is attractive for $f\bar{f}$ scattering and repulsive for $f f$ and $\bar{f}\bar{f}$ scattering.

Path Integral Quantization

§4.1 Path integrals

In contrast to canonical quantization, in which fields are promoted to operators, in path-integral quantization they remain ordinary functions. While the former allows for a more direct understanding of the notion of particle (thanks to ladder operators), the latter has the advantage of not being intrinsically perturbative, thus better describing theories with non-perturbative effects.

§4.1.1 Path integral in Quantum Mechanics

Consider a general (bosonic) quantum system, with Hermitian operators $\{\hat{q}_a, \hat{p}_a\}_{a=1,\dots,n}$ which satisfy the canonical commutation relations:

$$[\hat{q}_a, \hat{p}_b] = i\delta_{ab} \quad (4.1)$$

Eigenstates of these operators form two improperly-normalized complete orthonormal sets of eigenstates, with scalar product:

$$\langle q|p\rangle = \prod_{a=1}^n \frac{1}{\sqrt{2\pi}} e^{iq_a p_a} \equiv (2\pi)^{-n/2} e^{iq \cdot p} \quad (4.2)$$

with $q \equiv \{q_1, \dots, q_n\}, p \equiv \{p_1, \dots, p_n\}$. The switch from the Schrödinger picture to the Heisenberg one is carried by the Hamiltonian $\hat{H}(\hat{q}, \hat{p})$ of the system, which is assumed to be normal-ordered with all q_a on the left of all p_a : in this picture, all of the above relations remain valid for equal-time states.

Theorem 4.1.1 (Hamiltonian path integral)

The amplitude of $|q_i, t_i\rangle \rightarrow |q_f, t_f\rangle$ is a functional integral:

$$\langle q_f, t_f | q_i, t_i \rangle = \int_{q(t_i)=q_i}^{q(t_f)=q_f} \prod_{a=1}^n dq_a \frac{dp_a}{2\pi} \exp i \int_{t_i}^{t_f} dt \left[\sum_{b=1}^n \dot{q}_b(t) p_b(t) - H(q(t), p(t)) \right] \quad (4.3)$$

Proof. Consider a partition $\{t_m\}_{m=0,\dots,N} : t_m < t_j \forall m < j$ of the interval $[t_i, t_f]$, so that $t_0 \equiv t_i$ and $t_N \equiv t_f$: for simplicity, take them equally spaced, i.e. $t_m = t_0 + m\epsilon$, with

$N\epsilon = t_f - t_i$. The amplitude for $|q_i, t_i\rangle \rightarrow |q_f, t_f\rangle$ can thus be computed as:

$$\langle q_f, t_f | q_i, t_i \rangle = \int_{\mathbb{R}^{n(N-1)}} \prod_{m=0}^{N-1} dq_m \langle q_{m+1}, t_{m+1} | q_m, t_m \rangle$$

where the completeness relation was repeatedly used. Switching to the Schrödinger picture:

$$\begin{aligned} \langle q_{m+1}, t_{m+1} | q_m, t_m \rangle &= \langle q_{m+1} | e^{-i\hat{H}(t_{m+1}-t_m)} | q_m \rangle \\ &= \langle q_{m+1} | e^{-i\hat{H}\epsilon} | q_m \rangle = \int_{\mathbb{R}^n} dp_m \langle q_{m+1} | p_m \rangle \langle p_m | e^{-i\hat{H}\epsilon} | q_m \rangle \\ &= (2\pi)^{-n/2} \int_{\mathbb{R}^n} dp_m e^{iq_{m+1} \cdot p_m} \langle p_m | e^{-i\hat{H}(\hat{q}, \hat{p})\epsilon} | q_m \rangle \\ &= \frac{1}{(2\pi)^n} \int_{\mathbb{R}^n} dp_m e^{i(q_{m+1}-q_m) \cdot p_m} e^{-iH(q_m, p_m)\epsilon} \end{aligned}$$

The amplitude is then rewritten as:

$$\begin{aligned} \langle q_f, t_f | q_i, t_i \rangle &= \frac{1}{(2\pi)^n} \int_{\mathbb{R}^{2n(N+1)}} \prod_{m=0}^{N-1} dq_m dp_m e^{i[(q_{m+1}-q_m) \cdot p_m - H(q_m, p_m)\epsilon]} \\ &= \frac{1}{(2\pi)^n} \int_{\mathbb{R}^{2n(N+1)}} \prod_{m=0}^{N-1} dq_m dp_m \exp \sum_{k=1}^{N-1} i\epsilon \left[\frac{q_{k+1} - q_k}{\epsilon} \cdot p_k - H(q_k, p_k) \right] \end{aligned}$$

Consider now the $\epsilon \rightarrow 0$ limit, i.e. $N \rightarrow \infty$: the integration is then performed on an infinite number of variables, hence becoming a functional integral. The integration is carried on the set of functions $\{q(t) : \mathbb{R}^n \rightarrow \mathbb{R}^n : q(t_i) = q_i \wedge q(t_f) = q_f\}$; moreover, note that, in the $\epsilon \rightarrow 0$ limit, the sum becomes a Riemann integral:

$$\langle q_f, t_f | q_i, t_i \rangle = \frac{1}{(2\pi)^n} \int_{q(t_i)=q_i}^{q(t_f)=q_f} dq dp \exp i \int_{t_i}^{t_f} dt [\dot{q}(t) \cdot p(t) - H(q(t), p(t))]$$

which is precisely the thesis. \square

Proposition 4.1.1 (Lagrangian path integral)

For a p -quadratic Hamiltonian, the path integral can be written as:

$$\langle q_f, t_f | q_i, t_i \rangle = \int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] e^{\frac{i}{\hbar} \mathcal{S}[q, \dot{q}]} \quad (4.4)$$

where $\mathcal{S}[q, \dot{q}]$ is the action functional of the system and $[dq]$ is a suitably-normalized integration measure.

Proof. Consider a system with a p -quadratic Hamiltonian $H(q, p) = \frac{p^2}{2} + V(q)$ (however, the thesis holds for a general quadratic form of p), so that the previous proof can be rewritten

with a Gaussian integral (Eq. A.38):

$$\begin{aligned}
\langle q_{m+1}, t_{m+1} | q_m, t_m \rangle &= \frac{1}{(2\pi)^n} e^{-i\epsilon V(q_m)} \int_{\mathbb{R}^n} dp_m e^{i(q_{m+1}-q_m) \cdot p_m - \frac{i\epsilon}{2} p_m^2} \\
&= \frac{1}{(2\pi)^n} e^{-i\epsilon V(q_m)} \prod_{a=1}^n \int_{\mathbb{R}} dp_{m,a} e^{i(q_{m+1,a}-q_{m,a}) p_{m,a} - \frac{i\epsilon}{2} p_{m,a}^2} \\
&= \frac{1}{(2\pi)^n} e^{-i\epsilon V(q_m)} \prod_{a=1}^n \sqrt{\frac{2\pi}{i\epsilon}} e^{-\frac{1}{2i\epsilon} (q_{m+1,a}-q_{m,a})^2} \\
&= \left(\frac{1}{2\pi i\epsilon} \right)^{n/2} e^{i\epsilon \left[\frac{1}{2} \left(\frac{q_{m+1}-q_m}{\epsilon} \right)^2 - V(q_m) \right]}
\end{aligned}$$

Therefore:

$$\langle q_f, t_f | q_i, t_i \rangle = \left(\frac{1}{2\pi i\epsilon} \right)^{nN/2} \int_{\mathbb{R}^{n(N-1)}} \prod_{m=0}^{N-1} dq_m \exp \sum_{k=0}^{N-1} i\epsilon \left[\frac{1}{2} \left(\frac{q_{k+1}-q_k}{\epsilon} \right)^2 - V(q_k) \right]$$

In the $\epsilon \rightarrow 0$ limit, with the same considerations as above and absorbing the normalization in the integration measure, this becomes:

$$\langle q_f, t_f | q_i, t_i \rangle = \int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] \exp \left[i \int_{t_i}^{t_f} dt L(q(t), \dot{q}(t)) \right] =: \int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] e^{iS[q, \dot{q}]}$$

Reinstating explicitly \hbar yields the thesis. \square

This formulation elucidates the classical limit $\hbar \rightarrow 0$: while quantistically a particle explores all possible trajectories, which are then weighted by a phase e^{iS} , classically only trajectories which extremize the action are relevant, as those which are far from the extrema get extremely-oscillating phases under small deformations. Thus, the least-action principle has been recovered. The path integral is not only useful for computing amplitudes, but it is also used to express matrix elements on the coordinate-eigenbasis.

Lemma 4.1.1 (Time-ordered products)

Given $\{t_k\}_{k=1,\dots,N} \subset [t_i, t_f]$ and operators $\{\hat{\mathcal{O}}_k(\hat{q}(t))\}_{k=1,\dots,N} \equiv \{\hat{\mathcal{O}}_k(t)\}_{k=1,\dots,N}$, $N \in \mathbb{N}$, then:

$$\int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] \mathcal{O}_1(t_1) \dots \mathcal{O}_N(t_N) e^{iS} = \langle q_f, t_f | \mathfrak{T} \{ \hat{\mathcal{O}}_1(t_1) \dots \hat{\mathcal{O}}_N(t_N) \} | q_i, t_i \rangle \quad (4.5)$$

where \mathfrak{T} is the time-ordered product, which is defined as:

$$\mathfrak{T} \{ f(x) g(y) \} := \begin{cases} f(x) g(y) & x^0 > y^0 \\ g(y) f(x) & y^0 > x^0 \end{cases} \quad (4.6)$$

Proof. First, note that:

$$\int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] = \int_{\mathbb{R}^n} d\bar{q} \int_{q(t_i)=q_i}^{q(t)=\bar{q}} [dq] \int_{q(t)=\bar{q}}^{q(t_f)=q_f} [dq]$$

so that (due to the additivity of the action integral):

$$\begin{aligned}
 \int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] \mathcal{O}_k(t) e^{iS} &= \int_{\mathbb{R}^n} d\bar{q} \underbrace{\int_{q(t_i)=q_i}^{q(t)=\bar{q}} [dq] \exp \left[i \int_{t_i}^t dt L \right]}_{\langle \bar{q}, t | q_i, t_i \rangle} \mathcal{O}_k(\bar{q}) \underbrace{\int_{q(t)=\bar{q}}^{q(t_f)=q_f} [dq] \exp \left[i \int_t^{t_f} dt L \right]}_{\langle q_f, t_f | \bar{q}, t \rangle} \\
 &= \int_{\mathbb{R}^n} d\bar{q} \langle q_f, t_f | \bar{q}, t \rangle \mathcal{O}_k(\bar{q}) \langle \bar{q}, t | q_i, t_i \rangle = \int_{\mathbb{R}^n} d\bar{q} \langle q_f, t_f | \hat{\mathcal{O}}_k(t) | \bar{q}, t \rangle \langle \bar{q}, t | q_i, t_i \rangle \\
 &= \langle q_f, t_f | \hat{\mathcal{O}}_k(t) | q_i, t_i \rangle
 \end{aligned}$$

This links the matrix element of the operator $\hat{\mathcal{O}}_k(t)$ to the path integral of the function $\mathcal{O}_k(t)$. Consider now that, being $\mathcal{O}_k(t) \in \mathbb{C}$, i.e. commuting:

$$\int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] \mathcal{O}_1(t_1) \dots \mathcal{O}_N(t_N) e^{iS} = \int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] \mathcal{O}_{\pi(1)}(t_{\pi(1)}) \dots \mathcal{O}_{\pi(N)}(t_{\pi(N)}) e^{iS}$$

where $\pi \in S_N : t_{\pi(1)} \geq \dots \geq t_{\pi(N)}$ is a time-ordering permutation. Applying the above result:

$$\int_{q(t_i)=q_i}^{q(t_f)=q_f} [dq] \mathcal{O}_1(t_1) \dots \mathcal{O}_N(t_N) e^{iS} = \langle q_f, t_f | \hat{\mathcal{O}}_{\pi(1)}(t_{\pi(1)}) \dots \hat{\mathcal{O}}_{\pi(N)}(t_{\pi(N)}) | q_i, t_i \rangle$$

But $\hat{\mathcal{O}}_{\pi(1)}(t_{\pi(1)}) \dots \hat{\mathcal{O}}_{\pi(N)}(t_{\pi(N)}) =: \mathfrak{T}\{\hat{\mathcal{O}}_1(t_1) \dots \hat{\mathcal{O}}_N(t_N)\}$, thus completing the proof. \square

§4.1.2 Path integral in Quantum Field Theory

To generalize the path-integral formalism to QFT, consider the index a as running over points \mathbf{x} in space and over a spin/helicity (i.e. species/helicity) index m :

$$\hat{q}_a(t) \longrightarrow \hat{q}_m(t, \mathbf{x}) \quad |q, t\rangle \longrightarrow |q(\mathbf{x}), t\rangle$$

The generalizations of Eq. 4.3-4.4 is then straightforward:

$$\begin{aligned}
 \langle q_f(\mathbf{x}), t_f | q_i(\mathbf{x}), t_i \rangle &= \int_{q(t_i, \mathbf{x})=q_i(\mathbf{x})}^{q(t_f, \mathbf{x})=q_f(\mathbf{x})} \prod_m dq(\mathbf{x}) \frac{dp(\mathbf{x})}{2\pi} \times \\
 &\times \exp i \int_{t_i}^{t_f} dt \int_{\mathbb{R}^3} d^3x \left[\sum_m \dot{q}_m(t, \mathbf{x}) p_m(t, \mathbf{x}) - \mathcal{H}[q(t, \mathbf{x}), p(t, \mathbf{x})] \right] \quad (4.7)
 \end{aligned}$$

$$\langle q_f(\mathbf{x}), t_f | q_i(\mathbf{x}), t_i \rangle = \int_{q(t_i, \mathbf{x})=q_i(\mathbf{x})}^{q(t_f, \mathbf{x})=q_f(\mathbf{x})} \mathcal{D}q \exp i \int_{t_i}^{t_f} dt \int_{\mathbb{R}^3} d^3x \mathcal{L}[q(t, \mathbf{x}), \dot{q}(t, \mathbf{x})] \quad (4.8)$$

where $\mathcal{D}q$ generalizes the measure $[dq]$. It is best to express the path integral not for coordinate eigenstates, but for experimentally-observed states: these are states at $t_i \rightarrow -\infty, t_f \rightarrow +\infty$ with a definite number of particles (of various types).

Theorem 4.1.2 ((Bosonic) Path integral)

The amplitude for the (bosonic) process $|\alpha_{\text{in}}\rangle \rightarrow |\beta_{\text{out}}\rangle$ reads:

$$\langle \beta_{\text{out}} | \alpha_{\text{in}} \rangle = \int \mathcal{D}q \exp \left[i \int_{\mathbb{R}^4} dt d^3x \mathcal{L}[q(t, \mathbf{x}), \dot{q}(t, \mathbf{x})] \right] \langle \beta_{\text{out}} | q(\mathbf{x}), +\infty \rangle \langle q(\mathbf{x}), -\infty | \alpha_{\text{in}} \rangle \quad (4.9)$$

where the functional integral is now unconstrained.

Proof. Using Eq. 4.8:

$$\begin{aligned} \langle \beta_{\text{out}} | \alpha_{\text{in}} \rangle &= \int dq_i(\mathbf{x}) dq_f(\mathbf{x}) \langle \beta_{\text{out}} | q_f(\mathbf{x}), +\infty \rangle \langle q_f(\mathbf{x}), +\infty | q_i(\mathbf{x}), -\infty \rangle \langle q_i(\mathbf{x}), -\infty | \alpha_{\text{in}} \rangle \\ &= \int dq_i(\mathbf{x}) dq_f(\mathbf{x}) \int_{q(-\infty, \mathbf{x})=q_i(\mathbf{x})}^{q(+\infty, \mathbf{x})=q_f(\mathbf{x})} \mathcal{D}q \exp \left[i \int_{\mathbb{R}^4} dt d^3x \mathcal{L}[q(t, \mathbf{x}), \dot{q}(t, \mathbf{x})] \right] \times \\ &\quad \times \langle \beta_{\text{out}} | q_f(\mathbf{x}), +\infty \rangle \langle \alpha_{\text{in}} | q_i(\mathbf{x}), -\infty \rangle \end{aligned}$$

This is equivalent to having an unconstrained functional integral, hence the thesis. \square

This result is easily translated into Hamiltonian form. The expression for the wavefunctions which appear in Eq. 4.9 depends on the nature of the quantum field $q_m(t, \mathbf{x})$ considered (which, in this case, is constrained to be bosonic as of Eq. 4.1).

Note that Lemma 4.1.1 still holds: it allows to compute n -point Green functions via path integrals.

Proposition 4.1.2 (Correlation functions)

Given a (bosonic) quantum field $q(x)$, described by a Lagrangian \mathcal{L} , the n -point Green function can be computed as:

$$\langle 0 | \mathfrak{T} \{ \hat{q}(x_1) \dots \hat{q}(x_n) \} | 0 \rangle = \frac{\int \mathcal{D}q q(x_1) \dots q(x_n) e^{iS}}{\int \mathcal{D}q e^{iS}} \quad (4.10)$$

where \hat{q} denotes the field in the Heisenberg picture and the integrals are performed over field configurations $q(x) : q(\pm\infty, \mathbf{x}) = 0$.

Proof. By Lemma 4.1.1:

$$\int_{q(-\infty, \mathbf{x})=0}^{q(+\infty, \mathbf{x})=0} \mathcal{D}q q(x_1) \dots q(x_n) e^{iS} = \langle 0, +\infty | \mathfrak{T} \{ \hat{q}(x_1) \dots \hat{q}(x_n) \} | 0, -\infty \rangle$$

As the final-state vacuum is assumed to physically coincide with the initial-state vacuum, then:

$$|0, +\infty\rangle = e^{i\alpha} |0, -\infty\rangle \implies e^{-i\alpha} = \langle 0, +\infty | 0, -\infty \rangle$$

Inserting $\langle 0, +\infty | = \langle 0, -\infty | e^{-i\alpha}$ in the above equation and noting that:

$$\langle 0, +\infty | 0, -\infty \rangle = \int_{q(-\infty, \mathbf{x})=0}^{q(+\infty, \mathbf{x})=0} \mathcal{D}q e^{iS} \quad (4.11)$$

yields the thesis ($|0, -\infty\rangle \equiv |0\rangle$). \square

This is the equivalent of Eq. 3.48 evaluated with path integrals: in particular, it is clear that the denominator Eq. 4.11 gives the vacuum-to-vacuum transition amplitudes, “undressing” the numerator of vacuum-to-vacuum disconnected diagrams. Moreover, note how the normalization of $\mathcal{D}q$ carries no physical significance, as it simplifies in the fraction.

§4.2 Functional quantization of fields

§4.2.1 Scalar fields

First of all, consider a free scalar field theory:

$$G(x_1, \dots, x_n) = \frac{\int \mathcal{D}\phi \phi(x_1) \dots \phi(x_n) e^{i\mathcal{S}}}{\int \mathcal{D}\phi e^{i\mathcal{S}}} \quad \mathcal{S} = \frac{1}{2} \int d^4x (\partial_\mu \phi \partial^\mu \phi - m^2 \phi^2)$$

An integral so defined is not assured to converge, as the oscillating factor $e^{i\mathcal{S}}$ is not necessarily sufficient to provide convergence over large fluctuations, i.e. over field configurations with large action value. To ensure convergence, it is conventional to add a small convergence factor to the free Lagrangian:

$$\mathcal{L}_0 = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{1}{2} m^2 \phi^2 + \frac{i\epsilon}{2} \phi^2 \quad (4.12)$$

with $\epsilon \rightarrow 0^+$. This renders $G(x_1, \dots, x_n)$ a Gaussian integral, ensuring its convergence. The convergence factor can be viewed as an imaginary displacement of the mass $m^2 \mapsto m^2 - i\epsilon$: this is precisely the $i\epsilon$ -prescription for the Feynman propagator.

Definition 4.2.1 (Scalar generating functional)

Given a scalar field theory with Lagrangian \mathcal{L} , the **generating functional** is defined as:

$$Z[J] := \int \mathcal{D}\phi \exp \left[i \int d^4x (\mathcal{L} + J(x)\phi(x)) \right] \quad (4.13)$$

where $J(x)$ is a scalar field.

The generating functional is a functional of the **source field** $J(x)$, and its functional derivatives¹ are the correlation functions of the theory.

Lemma 4.2.1 (Correlation functions from generating functionals)

$$G(x_1, \dots, x_n) = \frac{1}{Z[0]} \left[\left(-i \frac{\delta}{\delta J(x_1)} \right) \dots \left(-i \frac{\delta}{\delta J(x_n)} \right) Z[J] \right]_{J=0} \quad (4.16)$$

¹The *functional derivative* is a formal manipulation which, given $J : \mathbb{R}^n \rightarrow \mathbb{R}$, is defined by:

$$\frac{\delta}{\delta J(x)} J(y) \equiv \delta^{(n)}(x - y) \quad (4.14)$$

It is clear, then, that given $f : \mathbb{R}^n \rightarrow \mathbb{C}^m$:

$$\frac{\delta}{\delta J(x)} \int_{\mathbb{R}^n} d^n y J(y) f(y) = f(x) \quad (4.15)$$

It also follows the usual derivative rules.

Proof. Using Eq. 4.15:

$$\begin{aligned} \frac{1}{Z[0]} \left[\left(-i \frac{\delta}{\delta J(x_1)} \right) \cdots \left(-i \frac{\delta}{\delta J(x_n)} \right) Z[J] \right]_{J=0} \\ = \frac{1}{Z[0]} \int d^4x \left[\left(-i \frac{\delta}{\delta J(x_1)} \right) \cdots \left(-i \frac{\delta}{\delta J(x_n)} \right) e^{i \int d^4y J(y) \phi(y)} \right]_{J=0} e^{iS} \\ = \frac{1}{Z[0]} \int d^4x \phi(x_1) \cdots \phi(x_n) e^{iS} \end{aligned}$$

which is precisely Eq. 4.10. \square

§4.2.1.1 Free theory

The generating functional for a free scalar theory can be written explicitly.

Proposition 4.2.1 (Free generating functional)

The generating functional for a free KG field ϕ is:

$$Z_0[J] = Z_0[0] \exp \left[-\frac{1}{2} \int d^4x d^4y J(x) \Delta(x-y) J(y) \right] \quad (4.17)$$

where $\Delta(x-y)$ is the Feynman propagator.

Proof. Integrating by parts the kinetic term of the action:

$$\int d^4x [\mathcal{L}_0 + J\phi] = \int d^4x \left[-\frac{1}{2} \phi(x) (\square + m^2 - i\epsilon) \phi(x) + J(x) \phi(x) \right] \equiv -\frac{1}{2} \langle \phi | \hat{A} | \phi \rangle + \langle J | \phi \rangle$$

(with abuse of notation) where the shifted KG operator $\hat{A} \equiv \square + m^2 - i\epsilon$ was introduced (Hermitian as $\epsilon \rightarrow 0^+$). Introducing a momentum-space basis:

$$|k\rangle : \langle x | k \rangle = e^{-ik_\mu x^\mu} \quad \int_{\mathbb{R}^4} \frac{d^4k}{(2\pi)^4} |k\rangle \langle k| = \mathbb{1} \quad (4.18)$$

it is possible to switch to momentum space through a 4D Fourier transform^a:

$$\phi(x) = \int \frac{d^4k}{(2\pi)^4} \langle x | k \rangle \langle k | \phi \rangle = \int \frac{d^4k}{(2\pi)^4} \phi(k) e^{-ik_\mu x^\mu} \quad \hat{A} |k\rangle = (-k^2 + m^2 - i\epsilon) |k\rangle$$

It is then possible to rewrite the integral above in momentum space:

$$\int d^4x [\mathcal{L}_0 + J\phi] = \int \frac{d^4k}{(2\pi)^4} \left[\frac{1}{2} \phi^*(k) (k^2 - m^2 + i\epsilon) \phi(k) + J^*(k) \phi(k) \right]$$

Note that $\phi^*(k) = \phi(-k)$, $J^*(k) = J(-k)$ and $\langle J | x \rangle = \langle x | J \rangle$. To further simplify the integral, define a shifted field:

$$|\chi\rangle = |\phi\rangle - \hat{A}^{-1} |J\rangle$$

The Jacobian of this transformation is trivially 1, thus:

$$\begin{aligned} \int d^4x [\mathcal{L}_0 + J\phi] &= -\frac{1}{2} (\langle \chi | + \langle J | \hat{A}^{-1}) \hat{A} (|\chi\rangle + \hat{A}^{-1} |J\rangle) + \langle J | \chi \rangle + \langle J | \hat{A}^{-1} |J\rangle \\ &= -\frac{1}{2} \langle \chi | \hat{A} | \chi \rangle + \frac{1}{2} \langle J | \hat{A}^{-1} |J\rangle \end{aligned}$$

As the second term is independent of the free field, it is then clear that:

$$Z_0[J] = \int \mathcal{D}\chi \exp \left[-\frac{i}{2} \langle \chi | \hat{A} | \chi \rangle + \frac{i}{2} \langle J | \hat{A}^{-1} | J \rangle \right] = Z_0[0] \exp \left[\frac{i}{2} \langle J | \hat{A}^{-1} | J \rangle \right]$$

Recall that the Green's function of the KG operator is the Feynman propagator Eq. 3.49, thus:

$$\hat{A}^{-1} J(x) = i \int d^4 y \Delta(x-y) J(y)$$

as $\hat{A} \Delta(x-y) = -i \delta^{(4)}(x-y)$. Then:

$$\exp \left[\frac{i}{2} \langle J | \hat{A}^{-1} | J \rangle \right] = \exp \left[\frac{i}{2} \int d^4 x J(x) \hat{A}^{-1} J(x) \right] = \exp \left[-\frac{1}{2} \int d^4 x d^4 y J(x) \Delta(x-y) J(y) \right]$$

This completes the proof. \square

^aThe Fourier transform is on all spacetime dimensions, as all possible paths contribute to the path integral, not just the solutions to the equations of motion.

Therefore, the propagator is a bi-local distribution. It is possible to express the integral in momentum space too, using $\hat{A}^{-1} |k\rangle = (-k^2 + m^2 - i\epsilon)^{-1} |k\rangle = i\tilde{\Delta}(k) |k\rangle$:

$$Z_0[J] = Z_0[0] \exp \left[-\frac{1}{2} \int \frac{d^4 k}{(2\pi)^4} J^*(k) \tilde{\Delta}(k) J(k) \right] \quad (4.19)$$

Using Lemma A.4.2, it is possible to make $Z[0]$ explicit:

$$Z_0[0] = \int \mathcal{D}\phi \exp \left[-\frac{i}{2} \langle \phi | \hat{A} | \phi \rangle \right] = \mathcal{N}(\det \hat{A})^{-1/2} \quad (4.20)$$

where $\det \hat{A}$ is known as a **functional determinant**. Although it has a physical meaning, as $Z_0[0] = \langle 0|0\rangle$ represents the vacuum-to-vacuum amplitude, the functional determinant cancels in the expression for the correlation functions (“undressing” of vacuum-to-vacuum diagrams).

§4.2.1.2 Interacting theory

Consider now a scalar field theory with interaction Lagrangian $\mathcal{L}_{\text{int}} = V[\phi]$. Lemma 4.2.1 is independent on whether the theory is free or not, so the generating integral needs only to be generalized in a way which suits the formalism of correlation functions and Wick's theorem.

Proposition 4.2.2 (Interacting generating functional)

The generating functional for an interacting KG field ϕ with $\mathcal{L}_{\text{int}} = -V[\phi]$ can be written as:

$$Z[J] = \sum_{n \in \mathbb{N}_0} \frac{(-i)^n}{n!} \left(\int d^4 x V \left[-i \frac{\delta}{\delta J(x)} \right] \right)^n Z_0[J] \quad (4.21)$$

where $Z_0[J]$ is the generating functional for the free theory.

Proof. Manipulating Def. 4.2.1:

$$\begin{aligned}
Z[J] &= \int \mathcal{D}\phi e^{iS + i \int d^4y J(y)\phi(y)} \\
&= \int \mathcal{D}\phi e^{i \int d^4y (\mathcal{L}_0 + V[\phi] + J(y)\phi(y))} \\
&= \int \mathcal{D}\phi e^{i \int d^4y (\mathcal{L}_0 + J(y)\phi(y))} \sum_{n=0}^{\infty} c_n \phi^n \\
&= \int \mathcal{D}\phi \sum_{n=0}^{\infty} c_n \left(-i \frac{\delta}{\delta J} \right)^n e^{i \int d^4x (\mathcal{L}_0 + J\phi)} \\
&= \int \mathcal{D}\phi e^{-i \int d^4x V[-i \frac{\delta}{\delta J(x)}]} e^{i \int d^4y (\mathcal{L}_0 + J(y)\phi(y))} \\
&= e^{-i \int d^4x V[-i \frac{\delta}{\delta J(x)}]} Z_0[J]
\end{aligned}$$

Expanding the first exponential yields the thesis. \square

Hence, in order to compute the correlation functions for the interacting theory, one only needs the generating functional of the free theory:

$$\begin{aligned}
&\langle 0 | \mathfrak{T} \{ \phi(x_1) \dots \phi(x_n) \} | 0 \rangle \\
&= \frac{1}{Z[0]} \left[\left(-i \frac{\delta}{\delta J(x_1)} \right) \dots \left(-i \frac{\delta}{\delta J(x_n)} \right) \sum_{k \in \mathbb{N}} \frac{(-i)^k}{k!} \left(\int d^4x V \left[-i \frac{\delta}{\delta J(x)} \right] \right)^k Z_0[J] \right]_{J=0} \quad (4.22)
\end{aligned}$$

This expression is analogous to Wick's theorem: due to the final condition $J = 0$, only terms which do not contain any $J(x)$ contribute to the correlation function, and these are precisely the fully-contracted terms. Moreover, as of Eq. 4.17, only terms with an even number of functional derivatives are non-vanishing, as an odd number of them results in terms with at least one $J(x)$: this allows the interaction potential to shape the interaction vertex.

Example 4.2.1 (Free scalar field theory)

Computing the 2-point Green function in a free scalar field theory is trivial with Eq. 4.17:

$$\begin{aligned}
\langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \} | 0 \rangle &= \frac{1}{Z_0[0]} \left[\left(-i \frac{\delta}{\delta J(x_1)} \right) \left(-i \frac{\delta}{\delta J(x_2)} \right) Z_0[J] \right]_{J=0} \\
&= \left(-i \frac{\delta}{\delta J(x_1)} \right) \left(-i \frac{\delta}{\delta J(x_2)} \right) \exp \left[-\frac{1}{2} \int d^4x d^4y J(x) \Delta(x-y) J(y) \right] \Big|_{J=0} \\
&= \frac{1}{Z_0[0]} \left(-i \frac{\delta}{\delta J(x_1)} \right) \left[i \int d^4y \Delta(x_2 - y) J(y) \right] Z_0[J] \Big|_{J=0} \\
&= \frac{1}{Z_0[0]} \left[\Delta(x_2 - x_1) + \right. \\
&\quad \left. - \int d^4y \Delta(x_2 - y) J(y) \int d^4y \Delta(x_1 - y) J(y) \right] Z_0[J] \Big|_{J=0} = \Delta(x_2 - x_1)
\end{aligned}$$

This expression can be used to compute the 3-point Green function:

$$\begin{aligned}
& \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \} | 0 \rangle \\
&= \frac{1}{Z_0[0]} \left(-i \frac{\delta}{\delta J(x_1)} \right) \left[\Delta(x_3 - x_2) - \int d^4 y \Delta(x_2 - y) J(y) \int d^4 y \Delta(x_3 - y) J(y) \right] Z_0[J] \Big|_{J=0} \\
&= \frac{1}{Z_0[0]} \left[i \Delta(x_3 - x_2) \int d^4 y \Delta(x_1 - y) J(y) + i \Delta(x_2 - x_1) \int d^4 y \Delta(x_3 - y) J(y) \right. \\
&\quad \left. + i \Delta(x_3 - x_1) \int d^4 y \Delta(x_2 - y) J(y) + \right. \\
&\quad \left. - i \int d^4 y \Delta(x_1 - y) J(y) \int d^4 y \Delta(x_2 - y) J(y) \int d^4 y \Delta(x_3 - y) J(y) \right] Z_0[J] \Big|_{J=0} = 0
\end{aligned}$$

The first non-trivial application of Wick's theorem is the 4-point Green function:

$$\begin{aligned}
& \langle 0 | \mathfrak{T} \{ \phi(x_1) \phi(x_2) \phi(x_3) \phi(x_4) \} | 0 \rangle \\
&= \frac{1}{Z_0[0]} \left(-i \frac{\delta}{\delta J(x_1)} \right) \left[i \Delta(x_4 - x_3) \int d^4 y \Delta(x_2 - y) J(y) + \right. \\
&\quad \left. + i \Delta(x_3 - x_2) \int d^4 y \Delta(x_4 - y) J(y) + i \Delta(x_4 - x_2) \int d^4 y \Delta(x_3 - y) J(y) + \right. \\
&\quad \left. - i \int d^4 y \Delta(x_2 - y) J(y) \int d^4 y \Delta(x_3 - y) J(y) \int d^4 y \Delta(x_4 - y) J(y) \right] Z_0[J] \Big|_{J=0} \\
&= \frac{1}{Z_0[0]} \left[\Delta(x_4 - x_3) \Delta(x_2 - x_1) + \Delta(x_3 - x_2) \Delta(x_4 - x_1) + \right. \\
&\quad \left. + \Delta(x_4 - x_2) \Delta(x_3 - x_1) + (\propto J^2) + (\propto J^4) \right] Z_0[J] \Big|_{J=0} \\
&= \Delta(x_4 - x_3) \Delta(x_2 - x_1) + \Delta(x_3 - x_2) \Delta(x_4 - x_1) + \Delta(x_4 - x_2) \Delta(x_3 - x_1)
\end{aligned}$$

§4.2.1.3 Symmetry factors

Consider Eq. 4.21 for $\lambda\phi^n$ -theory:

$$\begin{aligned}
Z[J] &= Z_0[0] \sum_{v,p \in \mathbb{N}_0} \frac{1}{v!} \left[\int d^4 x \frac{-i\lambda}{n!} \left(-i \frac{\delta}{\delta J(x)} \right)^n \right]^v \frac{1}{p!} \left[-\frac{1}{2} \int d^4 x d^4 y J(x) \Delta(x - y) J(y) \right]^p \\
&= Z_0[0] \sum_{v,p \in \mathbb{N}_0} \frac{1}{v! p! (n!)^v 2^p} \left[-i\lambda \int d^4 x \left(-i \frac{\delta}{\delta J(x)} \right)^n \right]^v \left[-\int d^4 x d^4 y J(x) \Delta(x - y) J(y) \right]^p
\end{aligned}$$

Each term in this expansion is represented by a Feynman diagram with v vertices and p propagators. This expression helps clarifying how the symmetry factor of a diagram is computed: first of all, note that there are $n!$ ways to rearrange functional derivatives in each v -term, amounting to a total of $(n!)^v$ equivalent ways for each term of the expansion, paired to $v!$ ways to rearrange the vertices; then, there are also $p!$ ways to rearrange propagators, with 2^p equivalent choices of sides (in total for p propagators).

Although it would seem that the combinatorial factor exactly cancels the prefactor of each term in the above expansion, it needs to be noted that not all these rearrangements are actual distinct (but equivalent) diagrams. The actual combinatorial factor of the diagram is the

symmetry factor, which can be computed as:

$$S = 2^\beta g \prod_k (k!)^{\alpha_k} \quad (4.23)$$

where β is the number of lines which connects a vertex to itself, g is the number of permutations of vertices which leave the diagram unchanged, and α_k is the number of couples of vertices connected by k identical lines.

§4.2.2 Electromagnetic field

With the functional formalism, it is possible to prove Eq. 3.71. The action of the electromagnetic field (massless spin-1 field) can be expressed as:

$$\begin{aligned} \mathcal{S} &= \int d^4x \left[-\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \right] = -\frac{1}{4} \int d^4x (\partial_\mu A_\nu - \partial_\nu A_\mu) (\partial^\mu A^\nu - \partial^\nu A^\mu) \\ &= -\frac{1}{2} \int d^4x (\partial_\mu A_\nu \partial^\mu A^\nu - \partial^\mu A_\nu \partial^\nu A_\mu) = \frac{1}{2} \int d^4x (A_\nu \square A^\nu - A_\nu \partial^\mu \partial^\nu A_\mu) \\ &= \frac{1}{2} \int d^4x \int \frac{d^4k d^4p}{(2\pi)^8} A_\nu(k) (e^{-ik_\alpha x^\alpha} \square e^{-ip_\alpha x^\alpha} \eta^{\mu\nu} - e^{-ik_\alpha x^\alpha} \partial^\mu \partial^\nu e^{-ip_\alpha x^\alpha}) A_\mu(p) \\ &= \frac{1}{2} \int d^4x \int \frac{d^4k d^4p}{(2\pi)^8} e^{-i(k+p)_\alpha x^\alpha} A_\nu(k) (-p^2 \eta^{\mu\nu} + p^\mu p^\nu) A_\mu(p) \\ &= \frac{1}{2} \int \frac{d^4k}{(2\pi)^4} A_\mu(-k) (-k^2 \eta^{\mu\nu} + k^\mu k^\nu) A_\nu(k) \end{aligned}$$

This kinetic term is orthogonal to k_μ , i.e. the action vanishes for all gauge fields of the form $A_\mu(k) = k_\mu \alpha(k)$: this is equivalent to saying that the Green's equation for this kinetic operator has no solution, as the latter is singular and therefore non-invertible. This problem is due to gauge invariance: indeed, $F_{\mu\nu}$ is invariant under (Eq. 2.53):

$$A_\mu(x) \mapsto A_\mu(x) + \partial_\mu \alpha(x) \quad (4.24)$$

and the non-invertibility arises from modes gauge-equivalent to $A_\mu(x) \equiv 0$ (obtained with the gauge transformation $\partial_\mu \alpha(x) = -A_\mu(x)$), i.e. non-physical degrees of freedom.

It is possible to isolate and quantize physical degrees of freedom, thanks to a technique due to Faddeev and Popov (introduced in [3]), which yields (see Sec. 9.4 of [6] for a detailed derivation):

$$\tilde{\Delta}_{\mu\nu}(k) = \frac{-i}{k^2 + i\epsilon} \left[\eta^{\mu\nu} - (1 - \xi) \frac{k^\mu k^\nu}{k^2} \right] \quad (4.25)$$

where $\xi \in \mathbb{R}$ is a gauge-fixing parameter. The most useful choices are $\xi = 0$ (Landau gauge), $\xi = 1$ (Feynman gauge) and $\xi = 3$ (Yennie gauge): the Feynman gauge $\xi = 1$ is adopted from now on. In position space:

$$\Delta_{\mu\nu}(x - y) = \int \frac{d^4k}{(2\pi)^4} e^{-ik_\alpha x^\alpha} \frac{-i\eta^{\mu\nu}}{k^2 + i\epsilon} \quad (4.26)$$

The formalism of the generating functional can be extended to massless spin-1 fields too.

Definition 4.2.2 (Generating functional for the electromagnetic field)

Given a gauge field A_μ described by a Lagrangian \mathcal{L} , the **generating functional** is defined as:

$$Z[J] := \int \mathcal{D}A \exp \left[i \int d^4x (\mathcal{L} + J^\mu(x) A_\mu(x)) \right] \quad (4.27)$$

where $J^\mu(x)$ is a 4-vector field.

Lemma 4.2.2 (Correlation functions from generating functionals)

$$\langle 0 | \mathfrak{T} \{ A_{\mu_1}(x_1) \dots A_{\mu_n}(x_n) \} | 0 \rangle = \frac{1}{Z[0]} \left[\left(-i \frac{\delta}{\delta J^{\mu_1}(x_1)} \right) \dots \left(-i \frac{\delta}{\delta J^{\mu_n}(x_n)} \right) Z[J] \right]_{J=0} \quad (4.28)$$

Proof. The proof is analogous to that of Lemma 4.2.1, using $\frac{\delta}{\delta J^\mu(x)} J^\nu(y) = \delta^\nu_\mu \delta^{(4)}(x - y)$. \square

Proposition 4.2.3 (Generating functional and the propagator)

$$Z[J] = Z[0] \exp \left[\frac{1}{2} \int d^4x d^4y J^\mu(x) \Delta_{\mu\nu}(x - y) J^\nu(y) \right] \quad (4.29)$$

Proof. The proof is analogous to that of Prop. 4.2.1. The extra negative sign is due to $\Delta_{\mu\nu}(x - y)$ solving:

$$(\square g^{\mu\nu} - \partial^\mu \partial^\nu) \Delta_{\mu\nu}(x - y) = i \delta^{(4)}(x - y) \quad (4.30)$$

Note that the source term is $i \delta^{(4)}(x - y)$ and not $-i \delta^{(4)}(x - y)$ (as for the KG operator). \square

§4.2.3 Spinor fields

To describe fermions in the functional formalism, it is necessary to define them through Grassmann fields (see §A.3.2.1) of the form:

$$\Psi(x) = \sum_i \theta_i \Psi_i(x)$$

where θ_i are Grassmann numbers and $\Psi_i(x)$ are a basis of Dirac spinor fields. This allows to extend the path-integral formalism to fermionic (anti-commuting) fields.

Example 4.2.2 (2-point Green function)

Using Eq. 4.10, it is possible to recover the Dirac propagator:

$$\langle 0 | \mathfrak{T} \{ \Psi(x_1) \bar{\Psi}(x_2) \} | 0 \rangle = \frac{\int \mathcal{D}\bar{\Psi} \mathcal{D}\Psi \Psi(x_1) \bar{\Psi}(x_2) \exp \left[i \int d^4x \bar{\Psi} (i \not{\partial} - m + i\epsilon) \Psi \right]}{\int \mathcal{D}\bar{\Psi} \mathcal{D}\Psi \exp \left[i \int d^4x \bar{\Psi} (i \not{\partial} - m + i\epsilon) \Psi \right]}$$

where convergence is assured by the $i\epsilon$ -prescription. Using Eq. A.37 for the numerator and Eq. A.36 for the denominator (note that $\bar{\Psi}$ is used, instead of Ψ^* , as they are unitarily equivalent), the functional determinant of the $i \not{\partial} - m + i\epsilon$ operator cancels out, leaving the matrix element of the inverse operator $(i \not{\partial} - m + i\epsilon)$ (more precisely, $[-i (i \not{\partial} - m + i\epsilon)]^{-1}$). To compute this matrix element in momentum space, consider the representation of the

fermionic field in it:

$$\Psi(x) = \int \frac{d^4k}{(2\pi)^4} \Psi(k) e^{-ik_\mu x^\mu} \quad \bar{\Psi}(x) = \int \frac{d^4k}{(2\pi)^4} \bar{\Psi}(k) e^{ik_\mu x^\mu} \quad (4.31)$$

Therefore, the matrix element is:

$$\langle 0 | \mathcal{T} \{ \Psi(x_1) \bar{\Psi}(x_2) \} | 0 \rangle = \int \frac{d^4k}{(2\pi)^4} \frac{1}{-i(\not{k} - m + i\epsilon)} e^{-ik_\mu(x_1 - x_2)^\mu} = S(x_1 - x_2)$$

which is precisely the Dirac propagator (Eq. 3.52).

To derive the Feynman rules for a fermionic theory, it is possible to define a generating functional.

Definition 4.2.3 (Fermionic generating functional)

Given a theory \mathcal{L} with a fermionic field Ψ , the **generating functional** is defined as:

$$Z[\eta, \bar{\eta}] = \int \mathcal{D}\bar{\Psi} \mathcal{D}\Psi \exp \left[i \int d^4x (\mathcal{L} + \bar{\eta}(x) \Psi(x) + \bar{\Psi}(x) \eta(x)) \right] \quad (4.32)$$

where $\eta(x)$ is a Grassmann field.

Proposition 4.2.4 (Free generating functional)

The generating functional for a free fermionic field Ψ is:

$$Z_0[\bar{\eta}, \eta] = Z_0[0, 0] \exp \left[- \int d^4x d^4y \bar{\eta}(x) S(x - y) \eta(y) \right] \quad (4.33)$$

Proof. Rewriting the exponent of Eq. 4.32 in momentum space:

$$\begin{aligned} \int d^4x (\mathcal{L}_0 + \bar{\eta} \Psi + \bar{\Psi} \eta) &= \int d^4x \int \frac{d^4k d^4p}{(2\pi)^8} \bar{\Psi}(k) e^{ik_\mu x^\mu} (i\not{\partial} - m + i\epsilon) \Psi(p) e^{-ip_\mu x^\mu} + \\ &\quad + \int d^4x \int \frac{d^4k d^4p}{(2\pi)^8} \bar{\eta}(k) e^{ik_\mu x^\mu} \Psi(p) e^{-ip_\mu x^\mu} \\ &\quad + \int d^4x \int \frac{d^4k d^4p}{(2\pi)^8} \bar{\Psi}(k) e^{ik_\mu x^\mu} \eta(p) e^{-ip_\mu x^\mu} \\ &= \int \frac{d^4k}{(2\pi)^4} [\bar{\Psi}(k) (\not{k} - m + i\epsilon) \Psi(k) + \bar{\eta}(k) \Psi(k) + \bar{\Psi}(k) \eta(k)] \end{aligned}$$

Now, shift the field as:

$$\Psi \mapsto \Xi = \Psi + (i\not{\partial} - m + i\epsilon)^{-1} \eta \equiv \Psi + \hat{A}^{-1} \eta$$

The integral can then be rewritten as (recall the Hermitianity of the Dirac operator):

$$\begin{aligned} \int d^4x (\mathcal{L}_0 + \bar{\eta}\Psi + \bar{\Psi}\eta) &= \int \frac{d^4k}{(2\pi)^4} \left[(\bar{\Xi} - \bar{\eta}\hat{A}^{-1}) \hat{A} (\Xi - \hat{A}^{-1}\eta) + \bar{\eta}\Xi + \bar{\Xi}\eta - 2\bar{\eta}\hat{A}^{-1}\eta \right] \\ &= \int \frac{d^4k}{(2\pi)^4} [\bar{\Xi}\hat{A}\Xi - \bar{\eta}\hat{A}^{-1}\eta] \\ &= \int d^4x \bar{\Xi}(i\not{\partial} - m + i\epsilon)\Xi - \int \frac{d^4k}{(2\pi)^4} \bar{\eta}(k) \frac{1}{\not{k} - m + i\epsilon} \eta(k) \end{aligned}$$

Therefore, the generating functional can be expressed as (using the shift-invariance of Grassmann integrals):

$$\begin{aligned} Z_0[\bar{\eta}, \eta] &= \int \mathcal{D}\bar{\Xi} \mathcal{D}\Xi \exp \left[i \int d^4x \bar{\Xi}(i\not{\partial} - m + i\epsilon)\Xi \right] \exp \left[- \int \frac{d^4k}{(2\pi)^4} \bar{\eta}(k) \frac{i}{\not{k} - m + i\epsilon} \eta(k) \right] \\ &= Z_0[0, 0] \exp \left[- \int d^4x d^4y \bar{\eta}(x) S(x - y) \eta(y) \right] \end{aligned}$$

which is the thesis. \square

As for the free scalar field, using Eq. A.36 makes $Z_0[0]$ an explicit functional determinant of the kinetic operator:

$$Z_0[0, 0] = \mathcal{N} \det \hat{A} \quad (4.34)$$

To compute correlation functions from the generating functional, note that for Grassmann numbers:

$$\frac{d}{d\eta} \theta \eta = -\frac{d}{d\eta} \eta \theta = -\theta \quad (4.35)$$

Lemma 4.2.3 (Correlation functions from generating functionals)

$$\langle 0 | \mathfrak{T} \{ \Psi(x_1) \bar{\Psi}(x_2) \dots \} | 0 \rangle = \frac{1}{Z[0, 0]} \left[\left(-i \frac{\delta}{\delta \bar{\eta}(x_1)} \right) \left(i \frac{\delta}{\delta \eta(x_2)} \right) \dots Z[\bar{\eta}, \eta] \right]_{\bar{\eta}, \eta=0} \quad (4.36)$$

Proof. Analogous to the proof of Lemma 4.2.1 with Eq. 4.35. \square

Finally, note that the anti-commuting nature of Grassmann numbers makes Eq. 3.51 trivial.

Example 4.2.3 (QED Feynman rules)

Consider the QED Lagrangian:

$$\mathcal{L}_{\text{QED}} = \bar{\Psi}(i\not{D} - m)\Psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} \quad (4.37)$$

Expanding the covariant derivative $D_\mu = \partial_\mu + ieA_\mu$:

$$\mathcal{L}_{\text{QED}} = \bar{\Psi}(i\not{\partial} - m)\Psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - e \bar{\Psi} \gamma^\mu \Psi A_\mu$$

It is then clear that this theory describes spin- $\frac{1}{2}$ (anti-)fermions which propagate with $\tilde{S}(p)$ and massless spin-1 gauge bosons which propagate with $\tilde{\Delta}_{\mu\nu}(k)$. The interaction vertex involves a fermion, an anti-fermion and a gauge boson, and it acts as $-ie \int d^4x \gamma^\mu$.

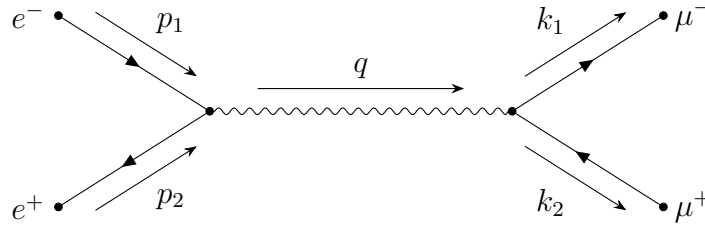
Amplitudes in QED

§5.1 Tree-level amplitudes

§5.1.1 $e^+e^- \rightarrow \mu^+\mu^-$ scattering

The $e^+e^- \rightarrow \mu^+\mu^-$ scattering is the simplest process in QED and one of the most important in high-energy physics, as it is used to calibrate e^+e^- colliders.

Recalling the Feynman rules for QED from §3.3.4.3, the amplitude can be found from the only Feynman diagram of the process¹:



with $q = p_1 + p_2 = k_1 + k_2$ by momentum conservation. By Eq. 3.63, then:

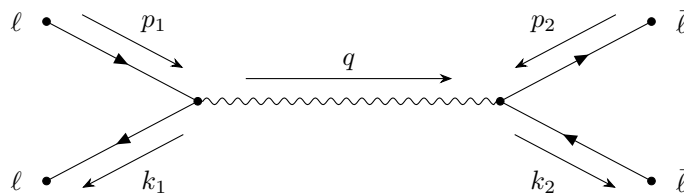
$$i\mathcal{M} = \bar{u}^{s_1}(k_1)(-ie\gamma^\mu)v^{s_2}(k_2)\frac{-i\eta_{\mu\nu}}{q^2 + i\epsilon}\bar{v}^{r_2}(p_2)(-ie\gamma^\nu)u^{r_1}(p_1) = \frac{ie^2}{q^2}\bar{u}^{s_1}(k_1)\gamma^\mu v^{s_2}(k_2)\bar{v}^{r_2}(p_2)\gamma_\mu u^{r_1}(p_1)$$

Lemma 5.1.1 (Bi-spinor products)

$$(\bar{u}\gamma^\mu v)^* = \bar{v}\gamma^\mu u \quad (5.1)$$

Proof. $(\bar{u}\gamma^\mu v)^* = v^\dagger(\gamma^\mu)^\dagger\bar{u}^\dagger = v^\dagger(\gamma^\mu)^\dagger(u^\dagger\gamma^0)^\dagger = v^\dagger(\gamma^\mu)^\dagger\gamma^0 u = v^\dagger\gamma^0\gamma^\mu u = \bar{v}\gamma^\mu u.$ \square

¹Note that, if $\ell\bar{\ell} \rightarrow \ell\bar{\ell}$, with $\ell = e^-, \mu^-$, was considered instead, there would be a second Feynman diagram to be computed:



with $q = p_1 - k_1 = k_2 - p_2$.

With this lemma, it is easy to see that (as $\bar{u}\gamma^\mu v \in \mathbb{C}$):

$$\begin{aligned} |\mathcal{M}_{r_1, r_2, s_1, s_2}|^2 &= \frac{e^4}{q^4} \bar{u}^{s_1}(k_1) \gamma^\mu v^{s_2}(k_2) \bar{v}^{r_2}(p_2) \gamma_\mu u^{r_1}(p_1) \bar{v}^{s_2}(k_2) \gamma^\nu u^{s_1}(k_1) \bar{u}^{r_1}(p_1) \gamma_\nu v^{r_2}(p_2) \\ &= \frac{e^4}{q^4} \bar{u}^{s_1}(k_1) \gamma^\mu v^{s_2}(k_2) \bar{v}^{s_2}(k_2) \gamma^\nu u^{s_1}(k_1) \bar{v}^{r_2}(p_2) \gamma_\mu u^{r_1}(p_1) \bar{u}^{r_1}(p_1) \gamma_\nu v^{r_2}(p_2) \end{aligned}$$

This matrix element has explicit dependence on spin states: although not impossible, it is difficult to experimentally retain control over them, both in the initial- and final-state. For this reason, it is convenient to consider a matrix element which is averaged over initial spin states and summed over final spin states (as usually the detector does not distinguish them):

$$|\mathcal{M}|^2 = \frac{1}{2} \sum_{r_1=1,2} \frac{1}{2} \sum_{r_2=1,2} \sum_{s_1} \sum_{s_2=1,2} |\mathcal{M}_{r_1, r_2, s_1, s_2}|^2 \quad (5.2)$$

Recalling Eq. 1.111 and writing spinor indices explicitly, the first half of the matrix element is:

$$\begin{aligned} \frac{e^4}{4q^4} \sum_{s_1, s_2=1,2} \sum_{a,b,c,d=1}^4 [\bar{u}^{s_1}(k_1)]_a \gamma_{ab}^\mu [v^{s_2}(k_2)]_b [\bar{v}^{s_2}(k_2)]_c \gamma_{cd}^\nu [u^{s_1}(k_1)]_d \\ = \frac{e^4}{4q^4} \sum_{a,b,c,d=1}^4 [k_1 + m_\mu]_{da} \gamma_{ab}^\mu [k_2 - m_\mu]_{bc} \gamma_{cd}^\nu = \frac{e^4}{4q^4} \text{tr}\{(\not{k}_1 + m_\mu) \gamma^\mu (\not{k}_2 - m_\mu) \gamma^\nu\} \end{aligned}$$

The unpolarized matrix element then is:

$$|\mathcal{M}|^2 = \frac{e^4}{4q^4} \text{tr}\{(\not{k}_1 + m_\mu) \gamma^\mu (\not{k}_2 - m_\mu) \gamma^\nu\} \text{tr}\{(\not{p}_1 + m_e) \gamma_\mu (\not{p}_2 - m_e) \gamma_\nu\} \quad (5.3)$$

This approach is general: any QED amplitude involving external fermions can be simplified into traces of products of gamma matrices, when squared and summed/averaged over spin states, hence eliminating explicit dependence on spinors.

§5.1.1.1 Traces

Due to their importance in QED, traces of gamma matrices should be studied systematically.

Proposition 5.1.1 (Odd number of gamma matrices)

For any $n \in \mathbb{N}$ odd:

$$\text{tr}\{\gamma^{\mu_1} \dots \gamma^{\mu_n}\} = 0 \quad (5.4)$$

Proof. As $\{\gamma^\mu, \gamma^5\} = 0$ and given the cyclic property of the trace:

$$\begin{aligned} \text{tr}\{\gamma^{\mu_1} \dots \gamma^{\mu_n}\} &= \text{tr}\{\gamma^5 \gamma^5 \gamma^{\mu_1} \dots \gamma^{\mu_n}\} = -\text{tr}\{\gamma^5 \gamma^{\mu_1} \gamma^5 \dots \gamma^{\mu_n}\} \\ &= (-1)^n \text{tr}\{\gamma^5 \gamma^{\mu_1} \dots \gamma^{\mu_n} \gamma^5\} = (-1)^n \text{tr}\{\gamma^5 \gamma^5 \gamma^{\mu_1} \dots \gamma^{\mu_n}\} = (-1)^n \text{tr}\{\gamma^{\mu_1} \dots \gamma^{\mu_n}\} \end{aligned}$$

Therefore, the trace vanishes for any odd $n \in \mathbb{N}$. \square

Proposition 5.1.2 (Even number of gamma matrices)

$$\text{tr}\{\gamma^\mu \gamma^\nu\} = 4\eta^{\mu\nu} \quad (5.5)$$

$$\text{tr}\{\gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma\} = 4(\eta^{\mu\nu} \eta^{\rho\sigma} - \eta^{\mu\rho} \eta^{\nu\sigma} + \eta^{\mu\sigma} \eta^{\nu\rho}) \quad (5.6)$$

Proof. Using $\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu} \mathbf{I}_4$:

$$\begin{aligned} \text{tr}\{\gamma^\mu \gamma^\nu\} &= \text{tr}\{2\eta^{\mu\nu} \mathbf{I}_4 - \gamma^\nu \gamma^\mu\} = 8\eta^{\mu\nu} - \text{tr}\{\gamma^\mu \gamma^\nu\} \\ \text{tr}\{\gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma\} &= \text{tr}\{2\eta^{\mu\nu} \gamma^\rho \gamma^\sigma - \gamma^\nu \gamma^\mu \gamma^\rho \gamma^\sigma\} \\ &= \text{tr}\{2\eta^{\mu\nu} \gamma^\rho \gamma^\sigma - \gamma^\nu 2\eta^{\mu\rho} \gamma^\sigma + \gamma^\nu \gamma^\rho 2\eta^{\mu\sigma} - \gamma^\nu \gamma^\rho \gamma^\sigma \gamma^\mu\} \\ &= 4(\eta^{\mu\nu} \eta^{\rho\sigma} - \eta^{\mu\rho} \eta^{\nu\sigma} + \eta^{\mu\sigma} \eta^{\nu\rho}) - \text{tr}\{\gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma\} \end{aligned}$$

which is the thesis. \square

In general, the trace of the product of n gamma matrices can be reduced to traces of products of $n - 2$ gamma matrices. Traces involving $\gamma^5 := i\gamma^0 \gamma^1 \gamma^2 \gamma^3$ are of interest too.

Proposition 5.1.3 (Odd number of gamma matrices)

For any $n \in \mathbb{N}$ odd:

$$\text{tr}\{\gamma^{\mu_1} \dots \gamma^{\mu_n} \gamma^5\} = 0 \quad (5.7)$$

Proof. As in the proof of Prop. 5.1.1:

$$\text{tr}\{\gamma^{\mu_1} \dots \gamma^{\mu_n} \gamma^5\} = (-1)^n \text{tr}\{\gamma^5 \gamma^{\mu_1} \dots \gamma^{\mu_n} \gamma^5 \gamma^5\} = (-1)^n \text{tr}\{\gamma^{\mu_1} \dots \gamma^{\mu_n} \gamma^5\}$$

which vanishes for any odd $n \in \mathbb{N}$. \square

Proposition 5.1.4 (Even number of gamma matrices)

$$\text{tr} \gamma^5 = \text{tr}\{\gamma^\mu \gamma^\nu \gamma^5\} = 0 \quad (5.8)$$

$$\text{tr}\{\gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma \gamma^5\} = -4i\epsilon^{\mu\nu\rho\sigma} \quad (5.9)$$

Proof. Recalling that $\{\gamma^5, \gamma^\mu\} = 0$:

$$\text{tr} \gamma^5 = \text{tr}\{\gamma^0 \gamma^0 \gamma^5\} = -\text{tr}\{\gamma^0 \gamma^5 \gamma^0\} = -\text{tr}\{\gamma^0 \gamma^0 \gamma^5\} = -\text{tr} \gamma^5$$

Analogously, given $\alpha \neq \mu, \nu$ (so that $\{\gamma^\mu, \gamma^\alpha\} = \{\gamma^\nu, \gamma^\alpha\} = \{\gamma^5, \gamma^\alpha\} = 0$):

$$\text{tr}\{\gamma^\mu \gamma^\nu \gamma^5\} = \text{tr}\{\gamma^0 \gamma^0 \gamma^\mu \gamma^\nu \gamma^5\} = (-1)^3 \text{tr}\{\gamma^0 \gamma^\mu \gamma^\nu \gamma^5 \gamma^0\} = -\text{tr}\{\gamma^\mu \gamma^\nu \gamma^5\}$$

Consider now $\text{tr}\{\gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma \gamma^5\}$: the above reasoning still applies if any couple of equal indices is present, therefore the trace must be proportional to $\epsilon^{\mu\nu\rho\sigma}$, which agrees with the anti-commutation relation of different gamma matrices. Computing $\text{tr}\{\gamma^0 \gamma^1 \gamma^2 \gamma^3 \gamma^5\}$ gives the correct proportionality constant $-4i$. \square

It is useful to recall an identity for the Levi-Civita symbol.

Lemma 5.1.2 (Levi–Civita symbol)

Given $n \in \mathbb{N}$ and $\{i_k\}_{k=1,\dots,n}, \{j_k\}_{k=1,\dots,n} \subset \{1, \dots, n\}$, then:

$$\epsilon_{i_1 \dots i_k i_{k+1} \dots i_n} \epsilon^{i_1 \dots i_k j_{k+1} \dots j_n} = n! \delta_{i_{k+1} i_n}^{j_{k+1} \dots j_n} \quad (5.10)$$

where the **generalized Kronecker delta** is defined as:

$$\delta_{\nu_1 \dots \nu_p}^{\mu_1 \dots \mu_p} := \begin{vmatrix} \delta_{\nu_1}^{\mu_1} & \dots & \delta_{\nu_1}^{\mu_p} \\ \vdots & \ddots & \vdots \\ \delta_{\nu_p}^{\mu_1} & \dots & \delta_{\nu_p}^{\mu_p} \end{vmatrix} \quad (5.11)$$

Gamma matrices' contractions are important too.

Proposition 5.1.5 (Contractions)

$$\gamma^\mu \gamma_\mu = 4 \quad (5.12)$$

$$\gamma^\mu \gamma^\nu \gamma_\mu = -2\gamma^\nu \quad (5.13)$$

$$\gamma^\mu \gamma^\nu \gamma^\rho \gamma_\mu = 4\eta^{\nu\rho} \quad (5.14)$$

$$\gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma \gamma_\mu = -2\gamma^\nu \gamma^\rho \gamma^\sigma \quad (5.15)$$

§5.1.1.2 Unpolarized matrix element

It is now possible to evaluate the matrix element in Eq. 5.3. Explicitly, the traces are:

$$\begin{aligned} \text{tr}\{(\not{k}_1 + m_\mu)\gamma^\mu(\not{k}_2 - m_\mu)\gamma^\nu\} &= \text{tr}\{k_{1\alpha}k_{2\beta}\gamma^\alpha\gamma^\mu\gamma^\beta\gamma^\nu + m_\mu k_{2\beta}\gamma^\mu\gamma^\beta\gamma^\nu - k_{1\alpha}m_\mu\gamma^\alpha\gamma^\mu\gamma^\nu - m_\mu^2\gamma^\mu\gamma^\nu\} \\ &= k_{1\alpha}k_{2\beta}4(\eta^{\alpha\mu}\eta^{\beta\nu} - \eta^{\alpha\beta}\eta^{\mu\nu} + \eta^{\alpha\nu}\eta^{\mu\beta}) - m_\mu^2 4\eta^{\mu\nu} \\ &= 4[k_1^\mu k_2^\nu + k_1^\nu k_2^\mu - \eta^{\mu\nu}(k_1 \cdot k_2 + m_\mu^2)] \end{aligned}$$

where $u \cdot v \equiv u^\mu v_\mu$. The other trace is analogous, so that:

$$\begin{aligned} |\mathcal{M}|^2 &= \frac{4e^4}{q^4} [k_1^\mu k_2^\nu + k_1^\nu k_2^\mu - \eta^{\mu\nu}(k_1 \cdot k_2 + m_\mu^2)] [p_{1\mu}p_{2\mu} + p_{1\nu}p_{2\nu} - \eta_{\mu\nu}(p_1 \cdot p_2 + m_e^2)] \\ &= \frac{4e^4}{q^4} [(k_1 \cdot p_1)(k_2 \cdot p_2) + (k_1 \cdot p_2)(k_2 \cdot p_1) - (p_1 \cdot p_2)(k_1 \cdot k_2) - m_\mu^2(p_1 \cdot p_2) + \\ &\quad (k_1 \cdot p_2)(k_2 \cdot p_1) + (k_1 \cdot p_1)(k_2 \cdot p_2) - (p_1 \cdot p_2)(k_1 \cdot k_2) - m_\mu^2(p_1 \cdot p_2) + \\ &\quad - 2(k_1 \cdot k_2)(p_1 \cdot p_2) - 2m_e^2(k_1 \cdot k_2) + \\ &\quad + 4(k_1 \cdot k_2)(p_1 \cdot p_2) + 4m_e^2(k_1 \cdot k_2) + 4m_\mu^2(p_1 \cdot p_2) + 4m_\mu^2 m_e^2] \\ &= \frac{8e^4}{q^4} [(k_1 \cdot p_1)(k_2 \cdot p_2) + (k_1 \cdot p_2)(k_2 \cdot p_1) + m_\mu^2(p_1 \cdot p_2) + m_e^2(k_1 \cdot k_2) + 2m_\mu^2 m_e^2] \end{aligned}$$

As $m_e \sim 10^{-4}m_\mu$, it is possible to neglect terms which contain m_e , thus the unpolarized matrix element is:

$$|\mathcal{M}|^2 = \frac{8e^4}{q^4} [(k_1 \cdot p_1)(k_2 \cdot p_2) + (k_1 \cdot p_2)(k_2 \cdot p_1) + m_\mu^2(p_1 \cdot p_2)] \quad (5.16)$$

Note that, for a generic $2 \rightarrow 2$ process, the unpolarized matrix element depends on 4 momenta, i.e. 16 variables. However, momentum conservation reduces this number by 4, as $k_1 + k_2 =$

$p_1 + p_2$ (4 equations), and mass-shell conditions on these 4 momenta reduces it further by 4. This leaves only 8 variables, but 6 of them are fixed by Lorentz invariance² (from 6 independent generators of $SO^+(1,3)$): the unpolarized matrix element hence only depends on 2 independent variables.

Definition 5.1.1 (Mandelstam variables)

Given a $2 \rightarrow 2$ scattering such that $i_1(k_1) + i_2(k_2) \rightarrow f_1(p_1) + f_2(p_2)$, then the **Mandelstam variables** are defined as:

$$s := (p_1 + p_2)^2 = (k_1 + k_2)^2 \quad (5.17)$$

$$t := (k_1 - p_1)^2 = (k_1 - p_2)^2 \quad (5.18)$$

$$u := (k_1 - p_2)^2 = (k_2 - p_1)^2 \quad (5.19)$$

Proposition 5.1.6

Only two Mandelstam variables are independent, as:

$$s + t + u = m_{i_1}^2 + m_{i_2}^2 + m_{f_1}^2 + m_{f_2}^2 \quad (5.20)$$

Proof. By direct calculation, using momentum conservation:

$$\begin{aligned} s + t + u &= p_1^2 + p_2^2 + 2p_1 \cdot p_2 + k_1^2 + p_1^2 - 2k_1 \cdot p_1 + k_2^2 + p_1^2 - 2k_2 \cdot p_1 \\ &= 3p_1^2 + p_2^2 + k_1^2 + k_2^2 - 2p_1 \cdot (k_1 + k_2 - p_2) = p_1^2 + p_2^2 + k_1^2 + k_2^2 \end{aligned}$$

Mass-shell conditions yield the thesis. \square

It is possible to parametrize the 2 degrees of freedom of $e^+e^- \rightarrow \mu^+\mu^-$ scattering with two Mandelstam variables, which for this process can be expressed as:

$$s = 2(k_1 \cdot k_2 + m_\mu^2) = 2(p_1 \cdot p_2 + m_e^2)$$

$$t = -2k_1 \cdot p_1 + m_e^2 + m_\mu^2 = -2k_2 \cdot p_2 + m_e^2 + m_\mu^2$$

$$u = -2k_2 \cdot p_1 + m_e^2 + m_\mu^2 = -2k_1 \cdot p_2 + m_e^2 + m_\mu^2$$

Neglecting terms with m_e , it is possible to write:

$$|\mathcal{M}|^2 = \frac{8e^4}{s^2} \left[\frac{t - m_\mu^2}{2} \frac{t - m_\mu^2}{2} + \frac{u - m_\mu^2}{2} \frac{u - m_\mu^2}{2} + m_\mu^2 \frac{s}{2} \right]$$

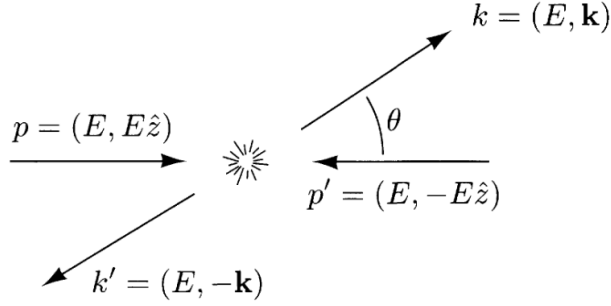
which is:

$$|\mathcal{M}|^2 = \frac{2e^4}{s^2} [(t - m_\mu^2)^2 + (u - m_\mu^2)^2 + 2m_\mu^2 s] \quad (5.21)$$

To compare this theoretical prediction with experimental data, however, it is better to choose an energy and an angle as the 2 independent variables. To do so, it is necessary to choose a reference frame: a natural choice is the **center-of-mass frame**, defined as the frame in which $\mathbf{p}_1 = -\mathbf{p}_2$. In this RF, momenta are parametrized as (see Fig. 5.1):

$$p_1 = (E, E\hat{\mathbf{e}}_z) \quad p_2 = (E, -E\hat{\mathbf{e}}_z) \quad k_1 = (E, \mathbf{k}) \quad k_2 = (E, -\mathbf{k})$$

²Note that momentum conservation is a consequence of translational invariance, therefore it can be united to Lorentz invariance to get Poincaré invariance, and indeed $ISO^+(1,3)$ has $6 + 4 = 10$ independent generators.

Figure 5.1: Center-of-mass RF for a $2 \rightarrow 2$ scattering process.

with:

$$|\mathbf{k}| = \sqrt{E^2 - m_\mu^2} \quad \mathbf{k} \cdot \hat{\mathbf{e}}_z = |\mathbf{k}| \cos \theta$$

Note that this is the most general expression for $2 \rightarrow 2$ scattering with $m_{i_1} = m_{i_2}$ and $m_{f_1} = m_{f_2}$. In this frame, the Mandelstam variables are:

$$s = 4E^2 \equiv E_{\text{cm}}^2$$

$$t = -2(E^2 - E|\mathbf{k}| \cos \theta) + m_e^2 + m_\mu^2 \approx -2E(E - |\mathbf{k}| \cos \theta) + m_\mu^2$$

$$u = -2(E^2 + E|\mathbf{k}| \cos \theta) + m_e^2 + m_\mu^2 \approx -2E(E + |\mathbf{k}| \cos \theta) + m_\mu^2$$

Inserting into Eq. 5.21:

$$\begin{aligned} |\mathcal{M}|^2 &= \frac{2e^4}{16E^4} [4E^2(E - |\mathbf{k}| \cos \theta)^2 + 4E^2(E + |\mathbf{k}| \cos \theta)^2 + 8m_\mu^2 E^2] \\ &= \frac{e^4}{2E^2} [(E - |\mathbf{k}| \cos \theta)^2 + (E + |\mathbf{k}| \cos \theta)^2 + 2m_\mu^2] \\ &= e^4 \left[1 + \frac{|\mathbf{k}|^2}{E^2} \cos^2 \theta + \frac{m_\mu^2}{E^2} \right] = e^4 \left[1 + \frac{m_\mu^2}{E^2} + \left(1 - \frac{m_\mu^2}{E^2} \right) \cos^2 \theta \right] \end{aligned}$$

§5.1.1.3 Unpolarized cross-section

The differential cross-section can be computed using Eq. B.2. In the case of $e^+e^- \rightarrow \mu^+\mu^-$ scattering, the flux factor is:

$$\Phi \equiv 4E_{\mathbf{p}_1} E_{\mathbf{p}_2} \left| \frac{\mathbf{p}_1}{E_{\mathbf{p}_1}} - \frac{\mathbf{p}_2}{E_{\mathbf{p}_2}} \right| = 4E^2 \left| \frac{E\hat{\mathbf{e}}_z}{E} - \frac{-E\hat{\mathbf{e}}_z}{E} \right| = 8E^2$$

On the other hand, the invariant phase-space measure is:

$$d\Pi_2 = \frac{d^3k_1}{(2\pi)^3 2E_{\mathbf{k}_1}} \frac{d^3k_2}{(2\pi)^3 2E_{\mathbf{k}_2}} (2\pi)^4 \delta^{(4)}(p_1 + p_2 - k_1 - k_2)$$

Note that $p_1 + p_2 = (2E, \mathbf{0})$, therefore $\mathbf{k}_1 = -\mathbf{k}_2$, i.e. $E_{\mathbf{k}_1} = E_{\mathbf{k}_2}$, so:

$$d\Pi_2 = \frac{d^3k_1}{(2\pi)^2 (2E_{\mathbf{k}_1})^2} \delta(2E - 2E_{\mathbf{k}_1}) = \frac{\mathbf{k}_1^2 dk_1 d\Omega_2}{16\pi^2 E_{\mathbf{k}_1}^2} \delta(2E_{\mathbf{k}_1} - 2E)$$

where $d\Omega_2$ is the measure on the 2-sphere³.

Lemma 5.1.3 (Scaling property)

Given a differentiable $f : I \subset \mathbb{R} \rightarrow \mathbb{R}$ with roots $\{x_i\}_{i \in \mathcal{I}} \subset I : f'(x_i) \neq 0 \ \forall i \in \mathcal{I}$, then:

$$\delta(f(x)) = \sum_{i \in \mathcal{I}} \frac{\delta(x_i)}{|f'(x_i)|} \quad (5.25)$$

Using Lemma 5.1.3, then:

$$\begin{aligned} d\Pi_2 &= \frac{\mathbf{k}_1^2 dk_1 d\Omega_2}{16\pi^2 E_{\mathbf{k}_1}^2} \delta(2\sqrt{m_\mu^2 + \mathbf{k}_1^2} - 2E) = \frac{\mathbf{k}_1^2 dk_1 d\Omega_2}{16\pi^2 E_{\mathbf{k}_1}^2} \frac{\delta(|\mathbf{k}_1| - \sqrt{E^2 - m_\mu^2})}{2 \frac{|\mathbf{k}_1|}{\sqrt{m_\mu^2 + \mathbf{k}_1^2}}} \\ &= \frac{|\mathbf{k}_1| dk_1 d\Omega_2}{32\pi^2 E} \delta(|\mathbf{k}_1| - \sqrt{E^2 - m_\mu^2}) \end{aligned}$$

Integrating over $|\mathbf{k}_1|$ and inserting into Eq. B.2 yields a general formula for $2 \rightarrow 2$ scattering:

$$\frac{d\sigma}{d\Omega_2} = \frac{1}{256\pi^2 E^2} \sqrt{1 - \frac{m_\mu^2}{E^2}} |\mathcal{M}|^2 \quad (5.26)$$

Inserting the unpolarized matrix element:

$$\frac{d\sigma}{d\Omega_2} = \frac{e^4}{256\pi^2 E^2} \sqrt{1 - \frac{m_\mu^2}{E^2}} \left[1 + \frac{m_\mu^2}{E^2} + \left(1 - \frac{m_\mu^2}{E^2} \right) \cos^2 \theta \right] \quad (5.27)$$

Proposition 5.1.7 (Total cross-section)

³In general, \mathbb{S}^n can be described as a surface embedded in \mathbb{R}^{n+1} , defined by:

$$\sum_{i=1}^{n+1} x_i^2 = 1 \quad (5.22)$$

Therefore, \mathbb{S}^n can be parametrized by *hyperspherical coordinates* $\{\varphi_1, \dots, \varphi_{n-1}, \varphi_n\} \subset [0, \pi]^{n-1} \times [0, 2\pi)$ as:

$$\begin{aligned} x_1 &= \cos \varphi_1 \\ x_2 &= \sin \varphi_1 \cos \varphi_2 \\ x_3 &= \sin \varphi_1 \sin \varphi_2 \cos \varphi_3 \\ &\vdots \\ x_{n-1} &= \sin \varphi_1 \dots \sin \varphi_{n-2} \cos \varphi_{n-1} \\ x_n &= \sin \varphi_1 \dots \sin \varphi_{n-2} \sin \varphi_{n-1} \cos \varphi_n \\ x_{n+1} &= \sin \varphi_1 \dots \sin \varphi_{n-2} \sin \varphi_{n-1} \sin \varphi_n \end{aligned} \quad (5.23)$$

It is then possible to define the *hyperspherical measure* on \mathbb{S}^n :

$$d\Omega_n \equiv \sin^{n-1} \varphi_1 \sin^{n-2} \varphi_2 \dots \sin^2 \varphi_{n-2} \sin \varphi_{n-1} d\varphi_1 d\varphi_2 \dots d\varphi_{n-2} d\varphi_{n-1} d\varphi_n \quad (5.24)$$

The **total cross-section** for $e^+e^- \rightarrow \mu^+\mu^-$ scattering is:

$$\sigma_{\text{tot}} = \frac{\pi\alpha^2}{3E^2} \sqrt{1 - \frac{m_\mu^2}{E^2}} \left(1 + \frac{1}{2} \frac{m_\mu^2}{E^2}\right) \quad (5.28)$$

Proof. Integrating Eq. 5.27:

$$\begin{aligned} \sigma_{\text{tot}} &:= \int d\Omega_2 \frac{d\sigma}{d\Omega_2} = \frac{\alpha^2}{16E^2} \sqrt{1 - \frac{m_\mu^2}{E^2}} \int_0^{2\pi} d\varphi \int_{-1}^1 d\cos\theta \left[1 + \frac{m_\mu^2}{E^2} + \left(1 - \frac{m_\mu^2}{E^2}\right) \cos^2\theta\right] \\ &= \frac{\alpha^2}{16E^2} \sqrt{1 - \frac{m_\mu^2}{E^2}} 2\pi \left[2 \left(1 + \frac{m_\mu^2}{E^2}\right) + \frac{2}{3} \left(1 - \frac{m_\mu^2}{E^2}\right)\right] = \frac{\pi\alpha^2}{3E^2} \sqrt{1 - \frac{m_\mu^2}{E^2}} \left(1 + \frac{1}{2} \frac{m_\mu^2}{E^2}\right) \end{aligned}$$

where $\alpha^2 \equiv \frac{e^2}{4\pi}$ was used. \square

It is clear that $\sigma_{\text{tot}} \propto \frac{\alpha^2}{s}$: the $\alpha^2 (1 + o(\alpha^2))$ is understood from the structure of the interaction vertex, while the s^{-1} dependence can be derived from dimensional analysis and requiring σ_{tot} to be a scalar.

§5.1.1.4 Helicity structure

Proposition 5.1.8 (Helicity eigenstates)

In the high-energy limit, i.e. $m \approx 0$, chirality eigenstates are helicity eigenstates too:

$$\hat{h} = \frac{\gamma^5}{2} \quad (5.29)$$

Proof. By Eq. 1.103, in the high-energy limit $\not{p}u = \not{p}v = 0$, therefore (as $\gamma^0\gamma^0 = I_4$):

$$\gamma^5\gamma^0\boldsymbol{\gamma} \cdot \mathbf{p}u = E\gamma^5u \quad \gamma^5\gamma^0\boldsymbol{\gamma} \cdot \mathbf{p}v = E\gamma^5v$$

As $m = 0$, then $\mathbf{p} = E\hat{\mathbf{p}}$. Moreover, in the chiral representation:

$$\gamma^5\gamma^0\gamma^i = \begin{bmatrix} -I_2 & 0 \\ 0 & I_2 \end{bmatrix} \begin{bmatrix} 0 & I_2 \\ I_2 & 0 \end{bmatrix} \begin{bmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{bmatrix} = \begin{bmatrix} \sigma^i & 0 \\ 0 & \sigma^i \end{bmatrix} = \Sigma^i$$

Recalling that the spin operator is $\mathbf{S} = \frac{\boldsymbol{\Sigma}}{2}$ and that the helicity operator is $\hat{h} = \mathbf{S} \cdot \hat{\mathbf{p}}$, then by the above equations the helicity operator is proportional to the chirality operator, yielding the thesis. \square

Recall Eq. 1.105-1.106. In the high-energy limit $m = 0$, i.e. $\mathbf{p} = E\hat{\mathbf{p}}$, so that:

$$u(p) = \sqrt{E} \begin{pmatrix} \sqrt{I_2 - \hat{\mathbf{p}} \cdot \boldsymbol{\sigma} \xi} \\ \sqrt{I_2 + \hat{\mathbf{p}} \cdot \boldsymbol{\sigma} \xi} \end{pmatrix} \quad v(p) = \sqrt{E} \begin{pmatrix} \sqrt{I_2 - \hat{\mathbf{p}} \cdot \boldsymbol{\sigma} \eta} \\ -\sqrt{I_2 + \hat{\mathbf{p}} \cdot \boldsymbol{\sigma} \eta} \end{pmatrix}$$

Consider chirality eigenstates $\xi_\pm : \hat{\mathbf{p}} \cdot \boldsymbol{\sigma} \xi_\pm = \pm \xi_\pm$ and $\eta_\pm : \hat{\mathbf{p}} \cdot \boldsymbol{\sigma} \eta_\pm = \mp \eta_\pm$, so that:

$$u_+(p) = \sqrt{2E} \begin{pmatrix} 0 \\ \xi_+ \end{pmatrix} \quad u_-(p) = \sqrt{2E} \begin{pmatrix} \xi_- \\ 0 \end{pmatrix} \quad v_+(p) = \sqrt{2E} \begin{pmatrix} \eta_+ \\ 0 \end{pmatrix} \quad v_-(p) = \sqrt{2E} \begin{pmatrix} 0 \\ -\eta_- \end{pmatrix}$$

As $\gamma^5 u_\pm = \pm u$ and $\gamma^5 v_\pm = \mp v_\pm$, these are clearly chirality eigenstates⁴, hence they are helicity eigenstates too:

$$\hat{h}u_\pm = \pm \frac{1}{2}u_\pm \quad \hat{h}v_\pm = \mp \frac{1}{2}v_\pm \quad (5.30)$$

Given the structure of the interaction vertex $e\bar{\Psi}\gamma^\mu\Psi$, the f and \bar{f} which interact at a vertex have the same chirality:

$$\bar{\Psi}\gamma^\mu\Psi' = (\psi_L^\dagger \quad \psi_R^\dagger) \begin{bmatrix} 0 & I_2 \\ I_2 & 0 \end{bmatrix} \begin{bmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{bmatrix} \begin{pmatrix} \psi_L' \\ \psi_R' \end{pmatrix} = \begin{pmatrix} \psi_L^\dagger \bar{\sigma}^\mu \psi_L' \\ \psi_R^\dagger \sigma^\mu \psi_R' \end{pmatrix}$$

which vanishes if f and \bar{f} have opposite chirality. This means that f and \bar{f} have opposite helicities, as per Eq. 5.30, i.e. parallel spins. In the case of $e^+e^- \rightarrow \mu^+\mu^-$ scattering, in the center-of-mass frame the e^+e^- pair must then annihilate in a photon with spin projection ± 1 along the $\hat{\mathbf{e}}_z$ axis, while the $\mu^+\mu^-$ pair must be created by a photon with spin projection ± 1 along the $\hat{\mathbf{k}} = (\sin\theta, 0, \cos\theta)$ axis ($\hat{\mathbf{e}}_y$ chosen such that $\hat{\mathbf{e}}_y \cdot \mathbf{k} = 0$). Recalling Eq. 2.60, there are two possible choices for the polarization vector of each photon:

$$\epsilon_{\text{in}}^\mu = \frac{1}{\sqrt{2}}(0, 1, \pm i, 0) \quad \epsilon_{\text{out}}^\mu = \frac{1}{\sqrt{2}}(0, \cos\theta, \pm i, -\sin\theta)$$

where \pm denotes a right/left-handed photon (i.e. $h = \pm 1$). The matrix element is proportional to the overlap between these polarization vectors, which is:

$$\epsilon_{\text{in}}^\mu \epsilon_{\text{out}\mu} = \frac{1}{2}(\pm_{\text{in}} \pm_{\text{out}} 1 - \cos\theta)$$

The amplitude is then proportional to the sum of the four possible square moduli:

$$|\mathcal{M}|^2 \propto \sum |\epsilon_{\text{in}}^\mu \epsilon_{\text{out}\mu}|^2 \propto (1 + \cos\theta)^2 + (1 - \cos\theta)^2 = (1 + \cos^2\theta) \quad (5.31)$$

Therefore, the angular dependence of the cross-section too is fixed a priori by symmetry considerations.

⁴Recall Lemma 2.2.3: as $\{\gamma^5, \gamma^\mu\} = 0$, chirality anti-commutes with charge conjugation, therefore antiferms have flipped chirality.

Part III

Renormalization

Chapter 6

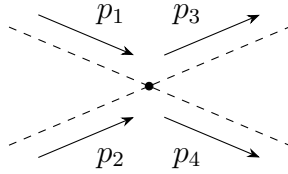
1-Loop Renormalization

As already noted in §3.3.3.1, the computation of Next-to-Leading Order (NLO) corrections often produces diverging amplitudes. However, this seems to be a consequence of Heisenberg's principle, implying that a quantum field theory cannot be valid at all energy scales at once.

§6.1 Renormalization of $\lambda\phi^4$ -theory

§6.1.1 1-Loop diagrams

Consider $\lambda\phi^4$ -theory and compute the scattering cross-section for $\phi\phi \rightarrow \phi\phi$ at LO:



$$\Rightarrow d\sigma^{\text{LO}} = \frac{1}{\Phi} \frac{d^3p_3}{(2\pi)^3 2E_{p_3}} \frac{d^3p_4}{(2\pi)^3 2E_{p_4}} (2\pi)^4 \delta^{(4)}(p_1 + p_2 - p_3 - p_4) \lambda^2$$

In the CM frame:

$$p_1 = (E, 0, 0, p) \quad p_2 = (E, 0, 0, -p) \quad p_3 = (E, p \sin \theta, 0, p \cos \theta) \quad p_4 = (E, -p \sin \theta, 0, -p \cos \theta)$$

$$s = 4E^2 \quad t = -2p^2(1 + \cos \theta) \quad u = -2p^2(1 - \cos \theta)$$

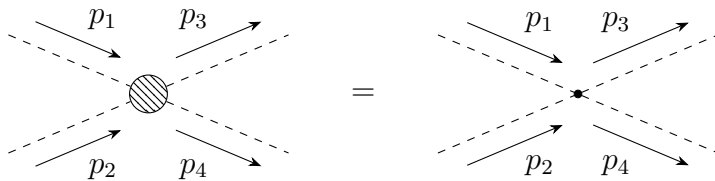
with flux factor $\Phi = 8pE$. Then the cross-section becomes:

$$d\sigma^{\text{LO}} = \frac{\lambda^2}{8pE} \frac{d^3p_3}{(2\pi)^2 (2E_{p_3})^2} \delta(2E - 2\sqrt{p_3^2 + m^2}) = \frac{\lambda^2}{8pE} \frac{p_3 dp_3 d\Omega_2}{(2\pi)^2 4E_{p_3}} \delta(E_{p_3} - E) = \frac{\lambda^2}{8pE} \frac{p}{16\pi E} d\cos \theta$$

Therefore, there's no angular dependence:

$$\frac{d\sigma^{\text{LO}}}{d\cos \theta} = \frac{\lambda^2}{128\pi} \frac{1}{E^2} \quad (6.1)$$

Now, consider the NLO corrections to this cross-section, i.e. the 1-loop corrections:



$$= \text{Feynman diagram with four external lines (p1, p2, p3, p4) meeting at a central vertex, with a shaded circle representing a loop.} + o(\lambda^2) \equiv -i\lambda F(s, t)$$

where $F(s, t)$ is called **form factor**.

Proposition 6.1.1 (Form factor)

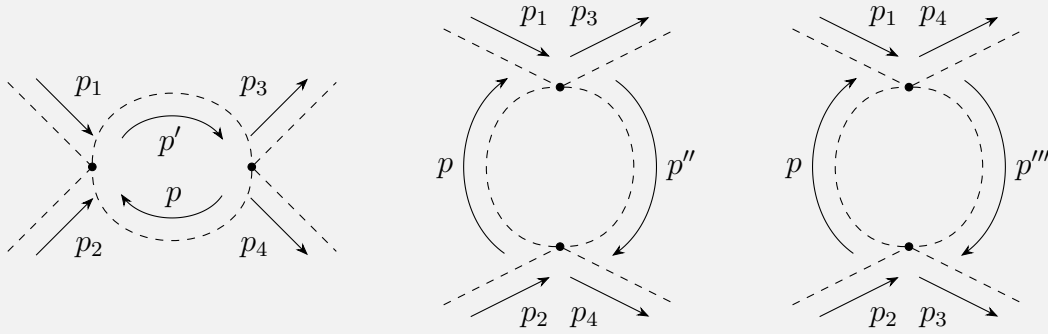
In $\lambda\phi^4$ -theory, the form factor at NLO is:

$$F(s, t) = 1 - i\lambda (V(s) + V(t) + V(u)) + o(\lambda^2) \quad (6.2)$$

with:

$$V(p^2) := \frac{1}{2} \int \frac{d^4k}{(2\pi)^4} \frac{i}{k^2 - m^2} \frac{i}{(p+k)^2 - m^2} \quad (6.3)$$

Proof. The NLO correction to the scattering amplitude is given by three loop diagrams:



with $p' = p + p_1 + p_2$, $p'' = p + p_1 - p_3$ and $p''' = p + p_1 - p_4$. These correspond to loop diagrams in the s , t and u channels, therefore, recalling Eq. 3.64, the thesis is trivially found. \square

Lemma 6.1.1 (Feynman parameters)

For $\{A_i\}_{i=1,\dots,n} \subseteq \mathbb{C}$ and $\{m_i\}_{i=1,\dots,n} \subseteq \mathbb{R}^+$, with $n \in \mathbb{N}$:

$$\frac{1}{A_1^{m_1} \dots A_n^{m_n}} = \int_{[0,1]^n} dx_1 \dots dx_n \delta(\sum_{i=1}^n x_i - 1) \frac{\prod_{i=1}^n x_i^{m_i-1}}{(\sum_{i=1}^n x_i A_i)^{m_1+\dots+m_n}} \frac{\Gamma(m_1 + \dots + m_n)}{\Gamma(m_1) \dots \Gamma(m_n)} \quad (6.4)$$

This allows to rewrite $V(p^2)$:

$$V(p^2) = -\frac{1}{2} \int \frac{d^4k}{(2\pi)^4} \int_0^1 dx [k^2 - m^2 + x(2k \cdot p + p^2)]^{-2} \equiv -\frac{1}{2} \int_0^1 dx \int \frac{d^4\ell}{(2\pi)^4} \frac{1}{(\ell^2 - M^2)^2}$$

having defined:

$$\ell \equiv k + xp \quad M^2 \equiv m^2 - x(1-x)p^2 \quad (6.5)$$

Lemma 6.1.2

$$I_d(n) \equiv \int \frac{d^d k}{(2\pi)^d} \frac{1}{(k^2 - M^2 + i\epsilon)^n} = \frac{i(-1)^n}{(4\pi)^{d/2}} \frac{\Gamma(n - \frac{d}{2})}{\Gamma(n)} \left(\frac{1}{M^2} \right)^{n-d/2} \quad (6.6)$$

Proof. The poles of the integrand function are $k_0 = \pm [\sqrt{M^2 + \mathbf{k}^2} - i\epsilon]$, therefore the integration path $I = \mathbb{R}$ can be rotated counterclockwise by $\frac{\pi}{2}$ without overlapping with any pole. The resulting Wick rotation $k_0 \mapsto ik_0$, i.e. $k \mapsto k_E$ (see Obs. 6.1.1), transforms the d -dimensional integral with Minkowskian metric into a d -dimensional integral with (negative)

Euclidean metric:

$$I_d(n) = i(-1)^n \int \frac{d^d k_E}{(2\pi)^d} \frac{1}{(k_E^2 + M^2 - i\epsilon)^n} = i(-1)^n \int_{\mathbb{S}^{d-1}} \frac{d\Omega_{d-1}}{(2\pi)^d} \int_{\mathbb{R}^+} dk_E \frac{k_E^{d-1}}{(k_E^2 + M^2)^n}$$

It can be shown that:

$$\int_{\mathbb{R}^+} dk_E \frac{k_E^{d-1}}{(k_E^2 + M^2)^n} = \frac{1}{2} \frac{\Gamma(n - \frac{d}{2}) \Gamma(\frac{d}{2})}{\Gamma(n)} \left(\frac{1}{M^2} \right)^{n - \frac{d}{2}} \quad (6.7)$$

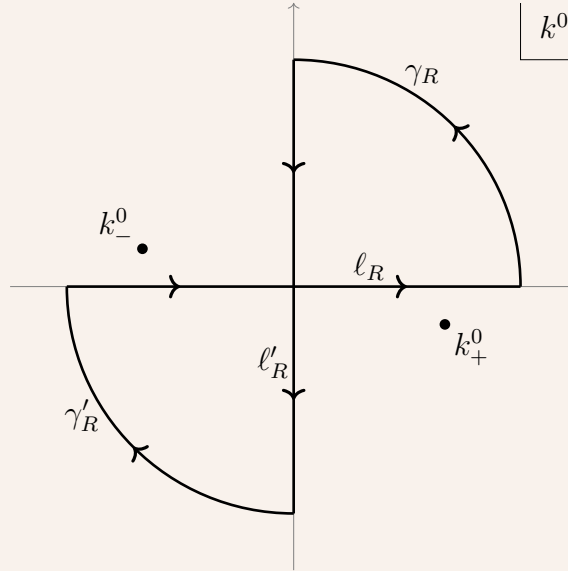
Inserting Eq. A.41 yields the thesis. \square

Observation 6.1.1 (Wick rotations)

Due to Lorentz invariance, the general form a d -dimensional loop integral is:

$$I = \int d^d k f(k^2)$$

However, these kinds of integrals cannot be treated with the standard techniques of Analysis, since the measure $d^d k$ is defined on the Lorentzian manifold $\mathbb{R}^{1,d-1}$ with Minkowskian metric $\eta_{\mu\nu} = (+1, -1, \dots, -1)$. In order to treat them in a standard way, consider first the integral $I_0 \equiv \int_{\mathbb{R}} dk^0 f(k^2)$ on the temporal component: by extending k^0 to the whole complex plane \mathbb{C} , this integral turns into the integral of a meromorphic^a function $f : \mathbb{C} \rightarrow \mathbb{C}$ over the path $\ell_\infty \equiv \mathbb{R}$ (identifying the curve with its support). In particular, consider the case of massive particles only, so that the poles of $f(k^2)$ have non-vanishing real parts: due to the general form $k^2 - m^2 + i\epsilon$ of the denominators of propagators, the poles come in couples of the form $k_\pm^0 = \pm \sqrt{m^2 + \mathbf{k}^2} \mp i\epsilon$. For example, the case of a single propagator is drawn as:



Considering the paths drawn in this figure, it is clear that, in the case of massive particle, the poles of $f(k^2)$ will always be outside of the total closed path γ ; hence, by Cauchy's residue theorem:

$$\oint_{\gamma} dk^0 f(k^2) = 0$$

The integrals over the arcs γ_R and γ'_R vanish for $R \rightarrow \infty$, hence this equation reduces to:

$$\int_{-\infty}^{\infty} dk^0 f(k^2) = \int_{-i\infty}^{i\infty} dk^0 f(k^2)$$

Then, substituting $k_E^0 = ik^0$, $\mathbf{k}_E = \mathbf{k}$ in the second integral, i.e. the so-called **Wick rotation**, the main result of this trick is found:

$$\int_{\mathbb{R}} dk^0 f(k^2) = i \int_{\mathbb{R}} dk_E^0 f(-k_E^2) \implies \int d^d k f(k^2) = i \int d^d k_E f(-k_E^2)$$

where the Euclidean momentum is now defined on the Euclidean manifold \mathbb{R}^d with Euclidean metric $\delta_{\mu\nu} = (+1, +1, \dots, +1)$. The Wick rotation thus turns a Minkowskian integral into a Euclidean integral, which can be treated with the standard techniques of Analysis.

^aGiven an open set $D \subset \mathbb{C}$, then $f : D \rightarrow \mathbb{C}$ is *meromorphic* if it is holomorphic on $D - P$, where $P \subset D$ is a set of isolated points called *poles*. Recall that a function $f : D \rightarrow \mathbb{C}$ is *holomorphic* on D if it is complex differentiable at every point in D .

Eq. 6.6 diverges for $(d, n) = (4, 2)$, as $\Gamma(z) = \frac{1}{z} - \gamma_E + o(z)$, therefore a cutoff needs to be introduced to regularize the integral:

$$\begin{aligned} V(p^2) &= \lim_{\Lambda \rightarrow \infty} \frac{-i}{2} \int_0^1 dx \int_0^\Lambda \frac{dk_E}{8\pi^2} \frac{k_E^3}{(k_E^2 + M^2)^2} = \lim_{\Lambda \rightarrow \infty} \frac{-i}{2} \int_0^1 dx \int_{M^2}^{\Lambda^2 + M^2} \frac{d\rho}{16\pi^2} \frac{\rho - M^2}{\rho^2} \\ &= \lim_{\Lambda \rightarrow \infty} \frac{-i}{32\pi^2} \int_0^1 dx \left[\log \left(1 + \frac{\Lambda^2}{M^2} \right) - \frac{\Lambda^2}{\Lambda^2 + M^2} \right] \\ &= \lim_{\Lambda \rightarrow \infty} \frac{-i}{32\pi^2} \int_0^1 dx \left[\log \frac{\Lambda^2}{M^2} - 1 + o\left(\frac{M^2}{\Lambda^2}\right) \right] \end{aligned}$$

The form factor can then be rewritten as:

$$F(s, t) = 1 + \frac{\lambda}{32\pi^2} \left[3 + \lim_{\Lambda \rightarrow \infty} \int_0^1 dx \left(\log \frac{M^2(s)}{\Lambda^2} + \log \frac{M^2(t)}{\Lambda^2} + \log \frac{M^2(u)}{\Lambda^2} \right) \right] + o(\lambda^2) \quad (6.8)$$

Now, define the threshold values $s_0 \equiv 4m^2$, $t_0 = u_0 \equiv 0$. The **physical coupling constant** λ_{phys} is defined as:

$$\left. \frac{d\sigma}{d \cos \theta} \right|_{s_0, t_0} = \frac{\lambda_{\text{phys}}^2}{128\pi m^2} \quad (6.9)$$

which is linked to the **bare coupling constant** $\lambda_{\text{bare}} \equiv \lambda$ by:

$$\lambda_{\text{bare}} F(s_0, t_0) = \lambda_{\text{phys}} \quad (6.10)$$

Proposition 6.1.2 (Bare and physical coupling)

The NLO form factor can be expressed as:

$$\lambda F(s, t) = \lambda_{\text{phys}} \left[1 - \frac{\lambda_{\text{phys}}}{32\pi^2} \int_0^1 dx \left(\log \frac{M^2(s)}{M^2(s_0)} + \log \frac{M^2(t)}{M^2(t_0)} + \log \frac{M^2(u)}{M^2(u_0)} \right) \right] + o(\lambda_{\text{phys}}^3) \quad (6.11)$$

Proof. By Eq. 6.10:

$$\lambda_{\text{phys}} = \lambda(1 - c_0\lambda) + o(\lambda^3) \iff \lambda = \lambda_{\text{phys}}(1 + c_0\lambda_{\text{phys}}) + o(\lambda_{\text{phys}}^3)$$

Therefore:

$$\lambda(1 - c\lambda) = \lambda_{\text{phys}}(1 + c_0\lambda_{\text{phys}})(1 - c\lambda_{\text{phys}}) + o(\lambda_{\text{phys}}^3) = \lambda_{\text{phys}}(1 - (c - c_0)\lambda_{\text{phys}}) + o(\lambda_{\text{phys}}^3)$$

Inserting Eq. 6.8 yields the thesis. \square

Expressing λ_{bare} as a function of λ_{phys} eliminates the divergence regularized by the cutoff, which is absorbed into λ_{phys} , the observable coupling constant. Note that the choice of s_0 and t_0 to define λ_{phys} is purely conventional: a more useful choice is imposing:

$$M^2(s_0) = M^2(t_0) = M^2(u_0) = \mu^2 \quad (6.12)$$

It is important to stress the difference between Λ and μ : Λ is a **regularization scale**, introduced to regularize diverging loop integrals, while μ is a **renormalization point**, introduced to eliminate the Λ -dependence¹. Both parameters are arbitrary, hence physical observables cannot depend on any of them: this means that the explicit μ -dependence of $F(s, t)$ must cancel the implicit μ -dependence of λ , so that λ_{phys} does not depend on it.

The infinities which cancel expressing $F(s, t)$ in terms of λ_{phys} are explicit when expressing $\lambda_{\text{bare}} = \lambda_{\text{bare}}(\lambda_{\text{phys}})$. In fact:

$$\lambda_{\text{phys}} = \lambda_{\text{bare}} \left[1 + \frac{3\lambda_{\text{bare}}}{32\pi^2} \left(1 + \log \frac{\mu^2}{\Lambda^2} \right) \right] + o(\lambda_{\text{bare}}^3)$$

Thus, choosing $\mu^2 \propto \Lambda^2$, it is possible to express λ_{phys} as a limit of λ_{bare} , and eventually $\lambda_{\text{phys}} = \lambda_{\text{bare}}$: the bare coupling constant is the physical coupling at the scale $\Lambda \rightarrow \infty$ (infinitely-small distances), i.e. including the fluctuations at all scales. Due to Heisenberg's principle, then, λ_{bare} is expected to diverge: excitations of the fields in momentum space are defined as normal coordinates of the system in position space, i.e. $p \sim \partial_x \phi \sim \frac{1}{\Delta x}$, which diverges as $\Delta x \rightarrow 0$ (as $\Lambda \rightarrow \infty$).

§6.1.2 Renormalized perturbation theory

The process of renormalization, which allows to obtain finite results from diverging amplitudes, can be implemented in various ways. That of introducing a UV cut-off Λ , expressing the bare observables in terms of the physical observables at some scale μ and then setting $\Lambda \rightarrow \infty$ is rather inefficient at higher orders in perturbation theory: thus, a more efficient regularization scheme must be adopted. The main regulators which are typically used are:

- cut-off regulator: a parameter $\Lambda \in \mathbb{R}^+$ is introduced to cut off the UV regime, i.e. imposing $k_E^2 \leq \Lambda^2$, so that the domain of UV divergences is removed entirely from the integration. However, this regulator effectively cuts off the infinite number of multi-particle states that can be produced at high energies (see Fig. 6.2), thus violating unitarity (and gauge symmetry, too);

¹The regularization scale Λ is an unphysical artifact introduced to isolate UV divergences, it depends on the specific regularization scheme adopted (UV cut-off, dimensional regularization, etc.) and it disappears in renormalized results, while the renormalization point μ is the physical reference scale which defines renormalized physical parameters.

- lattice regularization: to avoid the high-momentum regime, a discretization parameter $a \in \mathbb{R}^+$ is introduced, so that the position space is discretized as $x^\mu = n^\mu a$, with $n^\mu \in \mathbb{Z}^d$, thus reducing the momentum space to the first Brillouin zone $|k^\mu| \leq \frac{2\pi}{a} \forall \mu = 0, 1, \dots, d-1$. Clearly, this regularization scheme is analogous to the UV cut-off, therefore it carries the same problems of violation of unitarity and gauge symmetry, with the addition of violation of Lorentz symmetry too (it is not possible to define trigonometric and hyperbolic rotations in a discretized space);
- convergent propagators: UV divergences can be cured by making propagators more convergent with additional higher-order terms in the denominator. For example, adding a term like $+\frac{1}{\Lambda^2}\phi\Box\phi$ to the Lagrangian for a scalar field, the Feynman propagator becomes $i\tilde{\Delta}^{-1}(p) = p^2 + \frac{1}{\Lambda^2}p^4 - m^2$, which is more convergent than the usual propagator. This higher-order terms, however, may break the symmetries of the original Lagrangians;
- Pauli-Villars regulator: the UV behavior of loop integrals is modulated by the introduction of auxiliary particles, which are then eliminated by taking their masses to infinity. This regulator preserves gauge invariance in abelian gauge theories, but it is not gauge covariant in non-abelian gauge theories;
- dimensional regularization: in this scheme, the regulator is the number of dimensions, and in fact it preserves both unitarity, gauge symmetry and Lorentz invariance.

Observation 6.1.2 (Dimensional regularization)

In general, the dimensional regularization scheme consists in the analytic continuation of loop momenta to $d = 4 - 2\epsilon$, with $\epsilon \in \mathbb{C} : \Re \epsilon < 0$. This procedure turns loop integrals into meromorphic functions of $\epsilon \in \mathbb{C}$, allowing for the isolation of divergences as poles in ϵ .

The dimensional regularization prescription leaves freedom in choosing the dimensionality of external momenta, as well as the number of polarizations of both external and internal particles, thus allowing for the definition of different regularization schemes: the **conventional dimensional regularization** (CDR), in which all momenta and polarization are analytically continued to d dimensions, and the **'t Hooft–Veltman scheme** (HV), in which only internal momenta and polarizations are.

When considering non-chiral gauge theories like QED and QCD, CDR is the most natural choice, as the main difference between CDR and HV is the treatment of purely 4-dimensional objects, i.e. γ^5 and $\epsilon^{\mu\nu\sigma\rho}$. In particular, in CDR both the Dirac algebra and Lorentz indices are analytically continued to d dimensions, leading to a mathematical inconsistency when $d \notin \mathbb{N}$. In 4-dimensional Minkowski space $\mathbb{R}^{1,3}$ with metric signature $\eta = (+, -, -, -)$, the Dirac algebra is defined as $\mathfrak{cl}_{1,3}(\mathbb{C}) \cong \mathfrak{cl}_{1,3}(\mathbb{R}) \otimes \mathbb{C}$ and it admits a matrix representation with generators $\{\gamma^\mu\}_{\mu=0,1,2,3} \subset \mathbb{C}^{4 \times 4}$ such that:

$$\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu}I_4 \quad (6.13)$$

As in CDR Lorentz indices too are d -dimensional, the Dirac algebra becomes $\mathfrak{cl}_{1,d-1}(\mathbb{C}) \simeq \mathfrak{cl}_{1,d-1}(\mathbb{R}) \otimes \mathbb{C}$, where usual Minkowski space $\mathbb{R}^{1,3}$ is replaced by the pseudo-Euclidean space $\mathbb{R}^{1,d-1}$ with metric signature $\eta = (+, -, \dots, -)$. This structure, however, is ill-defined, as for $d \in \mathbb{N}$ then $\dim_{\mathbb{C}} \mathfrak{cl}_{1,d-1}(\mathbb{C}) = 2^d$, but an algebra of dimension $2^d : d \notin \mathbb{N}$ is meaningless. To solve this issue, in CDR spinor indices remain 4-dimensional: this is realized considering a matrix representation generated by $\{\gamma^\mu\}_{\mu=0,\dots,d-1} \subset \mathbb{C}^{4 \times 4}$ and imposing the formal relations **Eq. 6.13**. Consistency is achieved with the analytical continuation of trace identities: indeed, traces obtained by recursively applying **Eq. 6.13** (like $\text{tr}\{\gamma^\mu\gamma^\nu\} = 4\eta^{\mu\nu}$) are still valid in d -

dimensions, as the only dependence on dimension comes from contractions such as $\eta^\mu_\mu = d$. A fatal inconsistency of CDR arises when considering γ^5 . In $\mathfrak{cl}_{1,3}(\mathbb{C})$, this matrix is defined as $\gamma^5 := \frac{i}{4!}\epsilon_{\mu\nu\rho\sigma}\gamma^\mu\gamma^\nu\gamma^\rho\gamma^\sigma = i\gamma^0\gamma^1\gamma^2\gamma^3$ and has the property $\{\gamma^5, \gamma^\mu\} = 0 \ \forall \mu = 0, 1, 2, 3$, which allows to prove the following identity:

$$\text{tr}\{\gamma^5\gamma^\mu\gamma^\nu\gamma^\rho\gamma^\sigma\} = -4i\epsilon^{\mu\nu\rho\sigma} \quad (6.14)$$

This construction cannot be generalized consistently to $d \notin \mathbb{N}$. To show this, assume a d -dimensional generalization $\gamma^5 \in \mathbb{C}^{4 \times 4} : \{\gamma^5, \gamma^\mu\} = 0 \ \forall \mu = 0, \dots, d-1$, so that $\gamma^5\gamma^{\mu_1} \dots \gamma^{\mu_n} = (-1)^n \gamma^{\mu_1} \dots \gamma^{\mu_n} \gamma^5$. By the cyclicity of the trace, then:

$$(1 - (-1)^n) \text{tr}\{\gamma^5\gamma^{\mu_1} \dots \gamma^{\mu_n}\} = 0 \quad (6.15)$$

Consider $n = d$: clearly $1 - (-1)^d = 1 - e^{i\pi d} \neq 0$ for $d \notin \mathbb{N}$, so $\text{tr}\{\gamma^5\gamma^{\mu_1} \dots \gamma^{\mu_d}\} = 0$. This is an open contradiction to Eq. 6.14, as analytic continuation should continuously preserve the top product of the algebra as $\epsilon \rightarrow 0$.

This contradiction is the explicit manifestation of a more profound topological issue of analytically continuing the number of dimensions: the Levi-Civita symbol in $d = 4$ is linked² to the Grassmann algebra $\bigwedge(\mathbb{R}^{1,3})$, and in particular to its top-form, but $\bigwedge^k(\mathbb{R}^d)$ is only defined for $d \in \mathbb{N}$, so the top exterior subspace $\bigwedge^d(\mathbb{R}^{1,d-1})$ is meaningless for $d \notin \mathbb{N}$ and the Levi-Civita symbol cannot be analytically continued to $d = 4 - 2\epsilon$ dimensions.

In the following, the CDR scheme is adopted, since no chiral interactions are considered.

Proposition 6.1.3 (Dimensional analysis)

In $d = 4$:

$$[\phi] = [M] \quad [\lambda] = 1 \quad (6.16)$$

Proof. In natural units $\hbar = c = 1$, each physical quantity has the dimension of a power of mass, as:

$$[\hbar] = [E][T] = [M][T] \equiv 1 \quad [c] = [L][T]^{-1} \equiv 1$$

where $E = m$ ($c^2 = 1$) was used. Therefore, $[L] = [T] = [M]^{-1}$. In this system, the action must be a dimensionless quantity:

$$\mathcal{S} = \int d^4x \mathcal{L} \quad \implies \quad 1 = [L]^4[\mathcal{L}] = [M]^{-4}[\mathcal{L}] \quad \implies \quad [\mathcal{L}] = [M]^4$$

As $[\partial_\mu] = [L]^{-1} = [M]$, then, it is clear that $[\phi] = [M]$ and, being $\mathcal{L}_{\text{int}} = -\frac{\lambda}{4!}\phi^4$, $[\lambda] = 1$. \square

Proposition 6.1.4 (Dimensionality in CDR)

In $d = 4 - 2\epsilon$:

$$[\phi] = [M]^{1-\epsilon} \quad [\lambda] = [M]^{2\epsilon} \quad (6.17)$$

²Given an n -dimensional vector space $V(\mathbb{K})$, its *Grassmann algebra* (or exterior algebra) is the \mathbb{N}_0 -graded algebra $\bigwedge(V) = \bigoplus_{k=0}^n \bigwedge^k(V)$ of k -forms. It can be shown that $\dim_{\mathbb{K}} \bigwedge^k(V) = \binom{n}{k}$, hence the top exterior subspace $\bigwedge^n(V)$ is 1-dimensional: indeed, given a basis $\{e_i\}_{i=1,\dots,n} \subset V$, it is $\bigwedge^n(V) = \langle e_1 \wedge \dots \wedge e_n \rangle$, and the Levi-Civita symbol is normalized so that $\epsilon^{1\dots n}$ has the same sign of $e_1 \wedge \dots \wedge e_n$. See §A.3.2.

Proof. In the dimensional regularization scheme the differential becomes $d^4x \mapsto d^d x$, therefore the Lagrangian must have:

$$[\mathcal{L}] = [M]^d = [M]^{4-2\epsilon}$$

$[\partial_\mu] = [M]$ remains unchanged, therefore $[\phi] = [M]^{1-\epsilon}$ and, as a consequence, $[\lambda] = [M]^{2\epsilon}$. \square

To explicit the dimensionality of the (bare) coupling constant, a regularization scale μ_{reg} is introduced, so that:

$$\lambda = \mu_{\text{reg}}^{2\epsilon} \lambda_{\text{bare}} \quad (6.18)$$

Proposition 6.1.5

In $d = 4 - 2\epsilon$:

$$-i\lambda V(p^2) = -\frac{\lambda_{\text{bare}}}{32\pi^2} \int_0^1 dx \left[\frac{1}{\epsilon} - \gamma_E + \ln 4\pi - \log \frac{M^2}{\mu_{\text{reg}}^2} + o(\epsilon) \right] \quad (6.19)$$

where $\gamma_E \simeq 0.5772$ is the Euler–Mascheroni constant.

Proof. From Eq. 6.3 and using Eq. 6.6:

$$\begin{aligned} -i\lambda V(p^2) &= \frac{i\lambda}{2} \int_0^1 dx \int \frac{d^d \ell}{(2\pi)^d} \frac{1}{(\ell^2 - M^2)^2} = \frac{i\lambda_{\text{bare}}}{2} \mu_{\text{reg}}^{2\epsilon} \int_0^1 dx \frac{i}{(4\pi)^{2-\epsilon}} \frac{\Gamma(\epsilon)}{2} (M^2)^{-\epsilon} \\ &= -\frac{\lambda_{\text{bare}}}{32\pi^2} \int_0^1 dx (4\pi)^\epsilon \Gamma(\epsilon) \left(\frac{M^2}{\mu_{\text{reg}}^2} \right)^{-\epsilon} \end{aligned}$$

Now, the Laurent series in ϵ for this expression is:

$$\begin{aligned} -i\lambda V(p^2) &= -\frac{\lambda_{\text{bare}}}{32\pi^2} \int_0^1 dx (1 + \epsilon \log 4\pi + o(\epsilon)) \left(\frac{1}{\epsilon} - \gamma_E + o(\epsilon) \right) \left(1 - \epsilon \log \frac{M^2}{\mu_{\text{reg}}^2} + o(\epsilon) \right) \\ &= -\frac{\lambda_{\text{bare}}}{32\pi^2} \int_0^1 dx \left[\frac{1}{\epsilon} - \gamma_E + \log 4\pi - \log \frac{M^2}{\mu_{\text{reg}}^2} + o(\epsilon) \right] \end{aligned}$$

which is the thesis. \square

The dimensionally-regularized form factor, then, takes the form:

$$F(s, t) = 1 - \frac{\lambda_{\text{bare}}}{32\pi^2} \left[3 \left(\frac{1}{\epsilon} - \gamma_E + \log 4\pi \right) - \int_0^1 dx \left(\log \frac{M^2(s)}{\mu_{\text{reg}}^2} + \log \frac{M^2(t)}{\mu_{\text{reg}}^2} + \log \frac{M^2(u)}{\mu_{\text{reg}}^2} \right) \right] + o(\epsilon)$$

It would then be possible to proceed with the renormalization of the coupling constant specifying a cinematic (s_0, t_0) and setting $\lambda_{\text{phys}} = \lambda_{\text{bare}} F(s_0, t_0)$. However, given the arbitrariness of the renormalization point μ_{ren} , it is algebraically convenient to choose an unphysical cinematic which simplifies the expression of λ_{phys} .

Definition 6.1.1 (Minimal subtraction scheme)

The **minimal subtraction scheme** (MS) is a renormalization scheme with renormalization

point such that:

$$\lambda_{\overline{\text{MS}}} = \lambda_{\text{bare}} \left[1 - \frac{3\lambda_{\text{bare}}}{32\pi^2} \frac{1}{\epsilon} \right] + o(\lambda_{\text{bare}}^3) \quad (6.20)$$

When working with dimensional regularization, it is more useful to adopt the **modified MS scheme** ($\overline{\text{MS}}$), which includes the universal constant $-\gamma_E + \log 4\pi$ in the expression for the physical coupling constant:

$$\lambda_{\overline{\text{MS}}} = \lambda_{\text{bare}} \left[1 - \frac{3\lambda_{\text{bare}}}{32\pi^2} \left(\frac{1}{\epsilon} - \gamma_E + \log 4\pi \right) \right] + o(\lambda_{\text{bare}}^3) \quad (6.21)$$

Theorem 6.1.1 (Scales in the $\overline{\text{MS}}$ scheme)

In the $\overline{\text{MS}}$ scheme:

$$\mu_{\text{reg}} = \mu_{\text{ren}} \quad (6.22)$$

Proof. As of Eq. 6.21, using the same reasoning as in the proof of Prop. 6.1.2:

$$\lambda_{\text{bare}} = \lambda_{\overline{\text{MS}}} \left[1 + \frac{3\lambda_{\overline{\text{MS}}}}{32\pi^2} \left(\frac{1}{\epsilon} - \gamma_E + \log 4\pi \right) \right] + o(\lambda_{\overline{\text{MS}}}^3) \quad (6.23)$$

Moreover:

$$\begin{aligned} \lambda_{\text{bare}} F(s, t) &= \lambda_{\text{bare}} \left[1 - \frac{3\lambda_{\text{bare}}}{32\pi^2} \left(\frac{1}{\epsilon} - \gamma_E + \log 4\pi \right) \right] \\ &\quad + \frac{\lambda_{\text{bare}}^2}{32\pi^2} \int_0^1 dx \left(\log \frac{M^2(s)}{\mu_{\text{reg}}^2} + \log \frac{M^2(t)}{\mu_{\text{reg}}^2} + \log \frac{M^2(u)}{\mu_{\text{reg}}^2} \right) + o(\lambda_{\text{bare}}^3) \end{aligned}$$

Then, noting that $\lambda_{\text{bare}}^2 = \lambda_{\overline{\text{MS}}}^2 + o(\lambda_{\overline{\text{MS}}}^3)$:

$$\lambda_{\text{bare}} F(s, t) = \lambda_{\overline{\text{MS}}} \left[1 + \frac{\lambda_{\overline{\text{MS}}}}{32\pi^2} \int_0^1 dx \left(\log \frac{M^2(s)}{\mu_{\text{reg}}^2} + \log \frac{M^2(t)}{\mu_{\text{reg}}^2} + \log \frac{M^2(u)}{\mu_{\text{reg}}^2} \right) \right] + o(\lambda_{\overline{\text{MS}}}^3)$$

Comparing this expression to Eq. 6.11, it is clear that μ_{reg} is also the renormalization point of the scheme. \square

Hence, in the $\overline{\text{MS}}$ scheme $\mu_{\text{reg}} = \mu_{\text{ren}} \equiv \mu$.

§6.1.2.1 Lagrangian counterterms

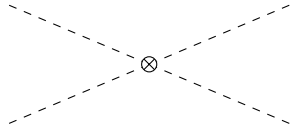
Define $Z_\lambda : \lambda_{\text{bare}} = Z_\lambda \lambda_{\overline{\text{MS}}} \equiv \lambda_{\overline{\text{MS}}} + \delta_\lambda$, so that:

$$Z_\lambda = 1 + \frac{3\lambda_{\overline{\text{MS}}}}{32\pi^2} \left(\frac{1}{\epsilon} - \gamma_E + \log 4\pi \right) + o(\lambda_{\overline{\text{MS}}}^2) \quad \delta_\lambda = \frac{3\lambda_{\overline{\text{MS}}}^2}{32\pi^2} \left(\frac{1}{\epsilon} - \gamma_E + \log 4\pi \right) + o(\lambda_{\overline{\text{MS}}}^3)$$

Inserting into the Lagrangian:

$$\begin{aligned} \mathcal{L} &= \frac{1}{2} (\partial_\mu \phi \partial^\mu \phi - m^2 \phi^2) - \frac{\lambda}{4!} \phi^4 \\ &= \frac{1}{2} (\partial_\mu \phi \partial^\mu \phi - m^2 \phi^2) - \frac{\lambda_{\overline{\text{MS}}}}{4!} \phi^4 - \frac{\delta_\lambda}{4!} \phi^4 + o(\lambda_{\overline{\text{MS}}}^3) \end{aligned}$$

It is then clear that renormalizing the coupling constant is equivalent to adding a counterterm to the Lagrangian, which contains the original divergence. This counterterm induces a new interaction vertex:



$$= -i\delta_\lambda \quad (6.24)$$

This process is general: divergences are cured order by order in perturbation theory by adding each time new counterterms to the Lagrangian: hence the name of renormalized perturbation theory³. At each order, each new counterterm is fixed by an appropriate physical condition.

§6.1.3 Renormalizability

The problem of renormalization has been reduced to that of classifying all possible counterterms to the Lagrangian, where each counterterm arises from the normalization of a different class of diverging Feynman diagrams.

Definition 6.1.2 (Superficial degree of divergence)

Given a scalar field theory in d dimensions, the **superficial degree of divergence** of a Feynman diagram with L loops and P_s scalar propagators is:

$$D := dL - 2P_s \quad (6.25)$$

The superficial degree of divergence gives an idea of how badly the considered diagram diverges: indeed, each loop corresponds to an integration with measure $d^d k$, which has dimension $[M]^d$, and each propagator corresponds to a factor of $\tilde{\Delta}(p) \sim (p^2 - m^2)^{-1}$, which has dimension $[M]^{-2}$, so D is the net power of momentum present in the amplitude:

$$\mathcal{M} \sim \int \frac{dk}{k^{1-D}}$$

Introducing a UV cut-off Λ , the divergence will be different for different values of D :

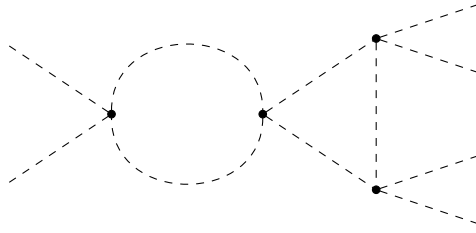
- $D > 0$: $\mathcal{M} \sim \Lambda^D$, hence the diagram has a UV divergence which must be renormalized;
- $D = 0$: $\mathcal{M} \sim \log \Lambda$, hence the diagram has a logarithmic UV divergence;
- $D < 0$: $\mathcal{M} \sim \Lambda^{-|D|}$, hence the diagram is UV finite.

This, however, is a naive analysis: in fact, the diagram may have diverging sub-diagrams which make the actual degree of divergence worse, or it may have symmetries which make divergences partially cancel, leaving a better degree of divergence. As an example, a trivial diagram with no loops and no propagators has $D = 0$, but it is not a diverging diagram.

The problem of diverging sub-diagrams requires a more careful analysis. Consider the following

³The name may be confusing, as the counterterms still contain divergences.

diagram in 4-dimensional $\lambda\phi^4$ -theory:



$$D = 4 \cdot 2 - 2 \cdot 5 = -2$$

Although $D < 0$, the loop integral diverges: however, in renormalized perturbation theory this divergence is cured by the counterterm Eq. 6.24. This is a general property: apart from superficially divergent diagrams, the only other divergences arise from diverging sub-diagrams of superficially finite diagrams, but these divergences are cured by the counterterms already added to the Lagrangian. Hence, the renormalization procedure correctly cures all UV divergences.

Theorem 6.1.2 (Bogolyubov–Parasyuk theorem)

For any renormalized quantum field theory, the perturbative expansion of matrix elements of the S -matrix is UV-finite.

Proof. The proof is found in [2, 4]. □

§6.1.3.1 $\lambda\phi^n$ -theory

Lemma 6.1.3 (Loop number)

Given a diagram with P propagators and V vertices, then:

$$L = P - V + 1 \quad (6.26)$$

Proof. Each loop is associated to an unconstrained momentum: each propagator corresponds to a momentum, while each vertex to a δ which constrains the momenta of the diagram, but there's also an overall momentum-conserving δ which makes one of the vertex-associated δ redundant, hence $L = P - (V - 1)$, which is the thesis. □

This result is valid in general for any field theory, and P represents the total number of propagators (even of different kinds).

Considering the specific case of $\lambda\phi^n$ -theory, instead, there's another relation which can simplify the computation of the superficial degrees of freedom: as each vertex has n scalar lines, then:

$$nV = 2P_s + N_s \quad (6.27)$$

where N_s is the number of external scalar lines. This expression takes into account the fact that each propagator links two vertices.

Theorem 6.1.3 (Superficial degree of divergence)

In $\lambda\phi^n$ -theory in d dimensions, the **superficial degree of divergence** is:

$$D = d + \left(n \frac{d-2}{2} - d \right) V - \frac{d-2}{2} N \quad (6.28)$$

Proof. By Eq. 6.27 $P = \frac{1}{2}(nV - N)$, thus by Eq. 6.26:

$$L = \left(\frac{n}{2} - 1\right) V - \frac{N}{2} + 1$$

Recalling Eq. 6.25:

$$D = d \left[\left(\frac{n}{2} - 1\right) V - \frac{N}{2} + 1 \right] - nV + N = d + \left(n \frac{d-2}{2} - d\right) V - \frac{d-2}{2} N$$

which is the thesis. \square

Example 6.1.1 ($\lambda\phi^4$ -theory)

The case of $\lambda\phi^4$ -theory in 4 dimensions is a particularly simple one, as:

$$D = 4 - N \tag{6.29}$$

Hence, the only classes of diverging diagrams are vacuum-to-vacuum diagrams, which diverge as $\sim \Lambda^4$, the $1 \rightarrow 1$ diagrams (i.e. the 2-point Green function of the theory and all its possible loop corrections), which diverges as $\sim \Lambda^2$, and the $2 \rightarrow 2$ scattering, which diverges as $\sim \log \Lambda$ (agreeing with Eq. 6.8). In this case there is only a finite number of classes of diverging diagrams.

Th. 6.1.3 can be derived also with dimensional analysis alone. In d -dimensional $\lambda\phi^n$ -theory, it is trivial to show that:

$$[\phi] = [M]^{\frac{d-2}{2}} \quad [\lambda] = [M]^{d-n\frac{d-2}{2}}$$

The dimensionality of λ is the dimensionality of a diagram with a single vertex, i.e. with $N = n$ amputated legs, and, as diagrams with the same N have the same dimensionality, of any diagram with $N = n$: generalizing, the dimensionality of a diagram with N amputated legs should be $d - N\frac{d-2}{2}$. On the other hand, if the diagram has V vertices, by definition it should have dimension $[\lambda]^V [M]^D$, thus:

$$\left(d - n\frac{d-2}{2}\right) V + D = d - N\frac{d-2}{2}$$

which is exactly Eq. 6.28. It is then clear that D decreases as V increases (i.e. going to higher orders in perturbation theory) if and only if the coupling constant has positive mass dimension, leaving only three possible classes theories:

- super-renormalizable theories: the coupling constant has strictly positive mass dimension, hence there is only a finite number of diverging diagrams.
- renormalizable theories: the coupling constant is dimensionless, hence there is only a finite number of classes of diverging diagrams (although their absolute number may be infinite);
- non-renormalizable theories: the coupling constant has strictly negative mass dimension, hence there is an infinite number of classes of diverging diagrams.

Note that these are general properties, not restricted to real scalar field theories.

Example 6.1.2 ($d = 4$ theories)

Consider $\lambda\phi^n$ -theory in $d = 4$ dimensions. The coupling constant in $\mathcal{L}_{\text{int}} = -\frac{1}{n!}\lambda\phi^n$ has dimension $4 - n$, hence:

- $\lambda\phi^3$ -theory is super-renormalizable: $D = 4 - V - N$, thus the only diverging diagrams are the 2-point Green function (loop corrections have $D < 0$) and the single interaction vertex, while all other diagrams have a finite amplitude;
- $\lambda\phi^4$ -theory is renormalizable: $D = 4 - N$, so the finitely-many classes of diverging diagrams are those listed in Ex. 6.1.1;
- $\lambda\phi^n$ -theory with $n \geq 5$ is non-renormalizable: for example, in $\lambda\phi^5$ -theory $D = 4 + V - N$, hence increasing the number of vertices increases the superficial degree of divergence, i.e. $\forall N \in \mathbb{N}_{\geq 2} \exists V \in \mathbb{N} : D > 0$.

The conceptual distinction between renormalizable and non-renormalizable theories is then clear: renormalizable theories can be renormalized with a finite set of counterterms, which can be fixed with a finite number of physical conditions (i.e. comparison with experiments), while non-renormalizable theories require an infinite set of counterterms, hence having no predictive power (as infinite parameters must be fixed by comparison with experiments).

Although at first glance non-renormalizability may seem problematic, in reality it is not. Consider a non-renormalizable theory with coupling constant $g \propto M^{-k}$, with $k \in \mathbb{R}^+$ and some mass scale M , and assume that the theory has already been renormalized up to $(N - 1)$ -point amplitudes, with $N \in \mathbb{N}_{\geq 2}$: then, an N -point amplitude at typical energy scale E can be expressed as:

$$\mathcal{A}_N(E) = \mathcal{A}_N^{(0)}(E) \left[1 + c_1 \frac{E^k}{M^k} + c_2 \frac{E^{2k}}{M^{2k}} + \cdots + c_n \frac{E^{nk}}{M^{nk}} \right]$$

where c_1, \dots, c_{n-1} are finite constant already fixed by physical conditions. As c_n should be fixed too by a new physical condition, the theory would not be predictive even at order n : however, at low energy $E \ll M$, this lack of predictivity due to c_n is negligible, as the quantity $(E/M)^{nk}$ can be assumed to be below the sensitivity of the experimental apparatus taken into consideration. Therefore, it is clear that non-renormalizable theories are acceptable low-energy theories: they gain predictive power by limiting the set of counterterms to be fixed to those necessary at the desired energy scale, as the ignorance about higher-order divergences is beyond the required accuracy. At $E \sim M$ the above series expansion diverges, signaling a lack of meaning of the theory: non-renormalizable theories have a built-in scale which provides their limit of validity, while renormalizable theories are in principle valid at all energy scales.

Example 6.1.3 (Fermi theory)

Consider the Fermi theory for β^- decay, which describes the decay $n^0 \rightarrow p^+ + e^- + \bar{\nu}_e$ as a 4-point vertex with interaction Lagrangian:

$$\mathcal{L}_{\text{F,int}} = -G_{\text{F}} \bar{\Psi}_n \Psi_p \bar{\Psi}_\nu \Psi_e \quad (6.30)$$

Since $[\Psi] = \frac{d-1}{2}$, then $[G_{\text{F}}] = 2 - d$, which makes the Fermi theory clearly non-renormalizable in $d = 4$. This, however, hints to the fact that this theory is only a low-energy approximation (i.e. an “effective” field theory) of the theory of the weak interaction, whose gauge bosons (W^\pm and Z^0) are massive and therefore short-lived, making the interaction a short-range one (and indeed approximated by a 4-point vertex).

Since the β^- decay is mediated by a W^- boson, the mass scale of the coupling is expected to be M_W , which is exactly the case:

$$G_F = \frac{\sqrt{2}}{8} \frac{g^2}{M_W^2} = 1.1663787(6) \cdot 10^{-5} \text{ GeV}^{-2}$$

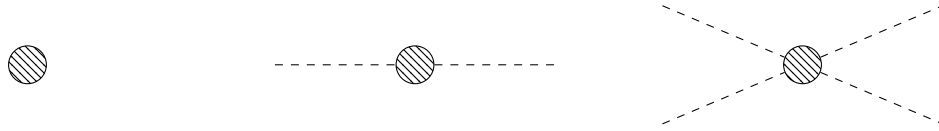
where g is the coupling constant of the weak interaction.

§6.1.3.2 $\lambda\phi^4$ -theory

Consider $\lambda\phi^4$ -theory in $d = 4$ dimensions:

$$\mathcal{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{1}{2} m_0^2 \phi^2 - \frac{\lambda_0}{4!} \phi^4 \quad (6.31)$$

where m_0, λ_0 are the bare mass and coupling constant. By Eq. 6.28, the superficial degree of divergence of a diagram depends only on the number of external legs, as $D = 4 - N$; moreover, as the theory is symmetric under $\phi \rightarrow -\phi$, all amplitudes with an odd number of external legs vanish. The only classes of divergent diagrams then are:



The first class is that of vacuum-to-vacuum diagrams, which are unobservable vacuum energy shifts as they are all accounted for in the definition of correlation functions (see §3.3.2.3). The second class of diagrams are the loop corrections to the 2-point Green function and, having $D = 2$, has the general diverging structure $\sim c_1 \Lambda^2 + c_2 p^2 \log \Lambda$, where p is the initial-state four-momentum. The third class of diagrams are the loop corrections to the $2 \rightarrow 2$ scattering amplitude, with $D = 0$ and diverging structure $\sim c_3 \log \Lambda$ (agreeing with §6.1.1). Ignoring the vacuum-to-vacuum diagrams, there are three infinite constants that need to be absorbed into three unobservable parameters of the theory: the bare mass, the bare coupling constant and the field-strength.

Correlation functions To express the identity operator on the Fock space \mathcal{F} of the theory, it is necessary to choose a complete basis of \mathcal{F} . Consider an eigenstate of the Hamiltonian $H |\lambda_0\rangle = m_\lambda |\lambda_0\rangle$ such that $\mathbf{P} |\lambda_0\rangle = \mathbf{0}$ (i.e. $m_\lambda^2 = E_0^2$): as per Lorentz invariance, all possible boosts of $|\lambda_0\rangle$ are still eigenstates of H , and they have all the possible three-momenta. Conversely, any eigenstate of H with definite three-momentum can be expressed as the boost of some zero-momentum eigenstate $|\lambda_0\rangle$. Denoting with $|\lambda_{\mathbf{p}}\rangle$ the eigenstate of H with three-momentum \mathbf{p} obtained as a boost of $|\lambda_0\rangle$ and defining $E_{\mathbf{p}}(\lambda) \equiv \sqrt{\mathbf{p}^2 + m_\lambda^2}$, where $m_\lambda : H |\lambda_0\rangle = m_\lambda |\lambda_0\rangle$, then the completeness relation on \mathcal{F} can be written as⁴:

$$\mathbf{1}_{\mathcal{F}} = |0\rangle \langle 0| + \sum_{\lambda} \int \frac{d^3 p}{(2\pi)^3} \frac{1}{2E_{\mathbf{p}}(\lambda)} |\lambda_{\mathbf{p}}\rangle \langle \lambda_{\mathbf{p}}| \quad (6.32)$$

where $|0\rangle$ is the vacuum state of \mathcal{F} .

⁴Recall that for the one-particle Hilbert space \mathcal{H} with relativistic normalization (Eq. 2.5):

$$\mathbf{1}_{\mathcal{H}} = \int \frac{d^3 p}{(2\pi)^3} \frac{1}{2E_{\mathbf{p}}} |\mathbf{p}\rangle \langle \mathbf{p}| \quad (6.33)$$

Proposition 6.1.6 (Correlation functions)

The correlation function for a real scalar field theory can be written as (with $x^0 > y^0$):

$$\langle 0 | \mathfrak{T} \{ \phi(x) \phi(y) \} | 0 \rangle = \sum_{\lambda} \int \frac{d^4 p}{(2\pi)^4} \frac{i}{p^2 - m_{\lambda}^2 + i\epsilon} e^{-ip_{\mu}(x-y)^{\mu}} |\langle 0 | \phi(0) | \lambda_0 \rangle|^2 \quad (6.34)$$

Proof. Inserting the identity in the 2-point Green function, as $x^0 > y^0$:

$$\langle 0 | \phi(x) \phi(y) | 0 \rangle = \langle 0 | \phi(x) | 0 \rangle \langle 0 | \phi(y) | 0 \rangle + \sum_{\lambda} \int \frac{d^3 p}{(2\pi)^3} \frac{1}{2E_{\mathbf{p}}(\lambda)} \langle 0 | \phi(x) | \lambda_{\mathbf{p}} \rangle \langle \lambda_{\mathbf{p}} | \phi(y) | 0 \rangle$$

The first term is constant and can be neglected, as it usually vanishes by symmetry or Lorentz invariance. Manipulating the matrix elements:

$$\begin{aligned} \langle 0 | \phi(x) | \lambda_{\mathbf{p}} \rangle &= \langle 0 | e^{ix_{\mu} P^{\mu}} \phi(0) e^{-ix_{\mu} P^{\mu}} | \lambda_{\mathbf{p}} \rangle = e^0 \langle 0 | \phi(0) | \lambda_{\mathbf{p}} \rangle e^{-ip_{\mu} x^{\mu}} \Big|_{p^0=E_{\mathbf{p}}} \\ &= \langle 0 | U^{-1} U \phi(0) U^{-1} U | \lambda_{\mathbf{p}} \rangle e^{-ip_{\mu} x^{\mu}} \Big|_{p^0=E_{\mathbf{p}}} = \langle 0 | \phi(0) | \lambda_0 \rangle e^{-ip_{\mu} x^{\mu}} \Big|_{p^0=E_{\mathbf{p}}} \end{aligned}$$

where U is a unitary operator implementing a Lorentz boost $\mathbf{p} \mapsto \mathbf{0}$, using that $U \phi(0) U^{-1} = \phi(0)$ (this is valid for a spinless field, while higher-spin fields have non-trivial transformation properties). Introducing an integration in dp^0 as in the proof of Th. 3.3.1 yields the thesis. \square

The case $x^0 < y^0$ is analogous, and a common expression can be formulated.

Definition 6.1.3 (Källén–Lehmann representation)

The **Källén–Lehmann spectral representation** of the correlation function of a real scalar field theory is defined as:

$$\langle 0 | \mathfrak{T} \{ \phi(x) \phi(y) \} | 0 \rangle = \int_0^{\infty} \frac{dM^2}{2\pi} \rho(M^2) \Delta(x-y; M^2) \quad (6.35)$$

where $\rho(M^2)$ is a positive **spectral density** function:

$$\rho(M^2) := \sum_{\lambda} 2\pi \delta(M^2 - m_{\lambda}^2) |\langle 0 | \phi(0) | \lambda_0 \rangle|^2 \quad (6.36)$$

As shown in Fig. 6.2, the one-particle states contribute to the spectral density with an isolated δ :

$$\rho(M^2) = 2\pi \delta(M^2 - m^2) Z + (\text{terms with } M^2 \gtrsim (2m)^2) \quad (6.37)$$

where Z is a number coming from the squared matrix element and it is called **field-strength renormalization** (the same factor appearing in the LSZ reduction formula). Note that m is the physical mass of the particle, as opposed to the bare mass m_0 appearing in the Lagrangian. The Fourier transform of the 2-point Green function then is:

$$\tilde{\Delta}(p) = \frac{iZ}{p^2 - m^2 + i\epsilon} + \int_{\sim 4m^2}^{\infty} \frac{dM^2}{2\pi} \rho(M^2) \frac{i}{p^2 - M^2 + i\epsilon} \quad (6.38)$$

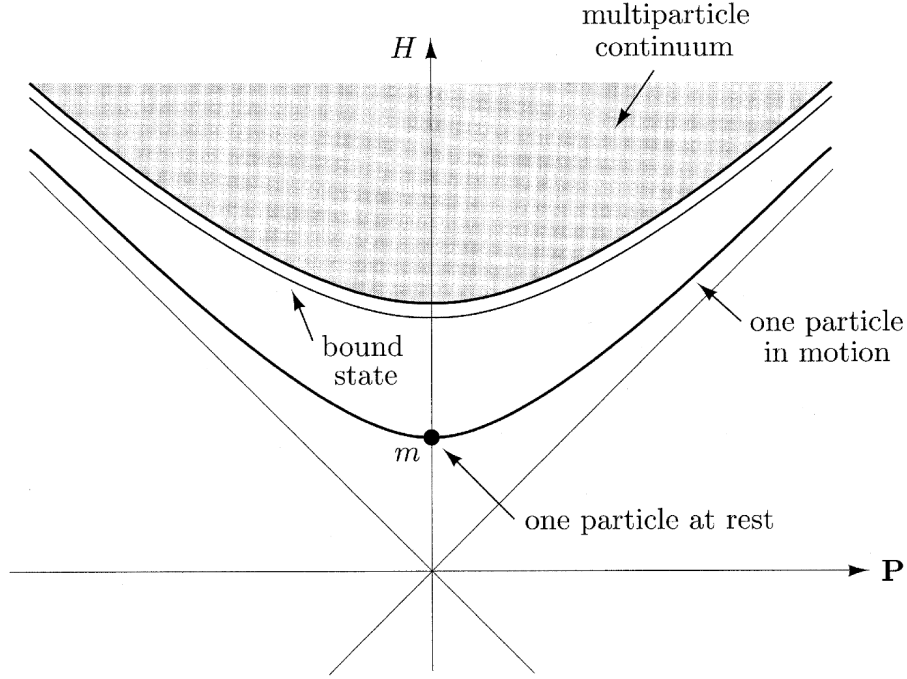


Figure 6.1: Hyperboloids in energy-momentum space formed by the eigenvalues of $P^\mu = (H, \mathbf{P})$.

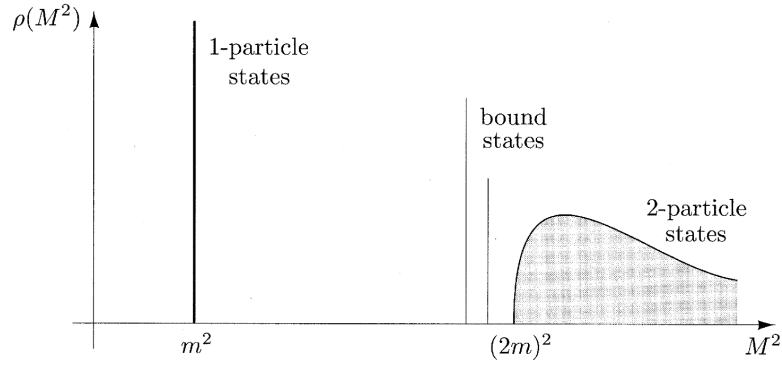


Figure 6.2: Spectral density function for a typical interacting field theory.

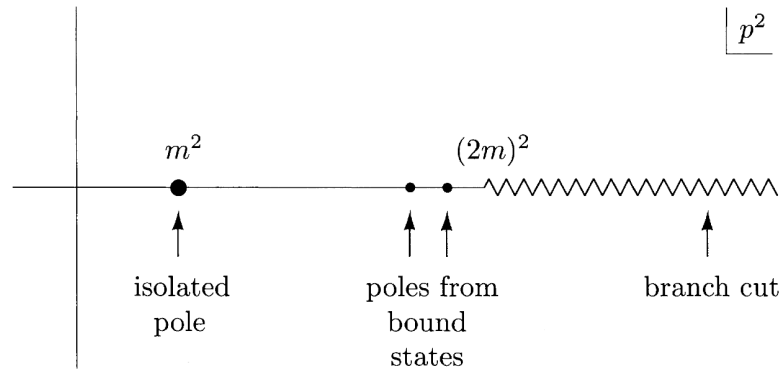


Figure 6.3: Analytic structure in the complex p^2 -plane of $\tilde{\Delta}(p)$.

The pole structure of $\tilde{\Delta}(p)$ is shown in Fig. 6.3: the first term gives an isolated simple pole in $p^2 = m^2$, while the second term contributes a branch cut beginning at $p^2 = (2m)^2$ (any two-particle bound states will appear as an additional δ in $\rho(M^2)$, i.e. an additional isolated pole below the branch cut). Eq. 6.38 is a direct generalization of Eq. 3.50: the former contains the

field-strength renormalization $Z = |\langle \lambda_0 | \phi(0) | 0 \rangle|^2$, i.e. the probability for $\phi(0)$ to create a given state from vacuum, while the latter has it implicitly, since in the free theory $|\langle \mathbf{p} | \phi(0) | 0 \rangle|^2 = 1$; moreover, Eq. 3.50 is the amplitude of a single particle propagating from x to y , as in the free theory $\phi(0)$ can only create a single particle from vacuum, while Eq. 6.38 contains contributions from multiparticle intermediate states with a continuous mass spectrum.

Counterterms Eq. 6.38 has a single term which is singular at $p^2 = m^2$, and to eliminate the awkward residue Z it is possible to set:

$$\phi \equiv \sqrt{Z} \phi_r \quad (6.39)$$

This is the renormalized field (recall that it is equivalent to renormalize multiplicatively or additively).

Proposition 6.1.7 (Renormalized $\lambda\phi^4$ -theory)

The renormalized Lagrangian of $\lambda\phi^4$ -theory takes the form:

$$\mathcal{L} = \frac{1}{2} \partial_\mu \phi_r \partial^\mu \phi_r - \frac{1}{2} m^2 \phi_r^2 - \frac{\lambda}{4!} \phi_r^4 + \frac{1}{2} \delta_Z \partial_\mu \phi_r \partial^\mu \phi_r - \frac{1}{2} \delta_m \phi_r^2 - \frac{\delta_\lambda}{4!} \phi_r^4 \quad (6.40)$$

with:

$$\delta_Z \equiv Z - 1 \quad \delta_m \equiv m_0^2 Z - m^2 \quad \delta_\lambda \equiv \lambda_0 Z^2 - \lambda \quad (6.41)$$

Proof. Inserting Eq. 6.39 into the Lagrangian of $\lambda\phi^4$ -theory:

$$\mathcal{L} = \frac{1}{2} Z \partial_\mu \phi_r \partial^\mu \phi_r - \frac{1}{2} m_0^2 Z \phi_r^2 - \frac{\lambda_0}{4!} \phi_r^4$$

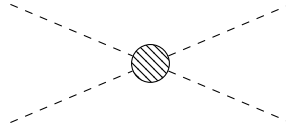
The isolation of counterterms is then trivial. □

This Lagrangian is composed of the standard $\lambda\phi^4$ -theory Lagrangian, with renormalized field, mass and coupling constant, plus three counterterms: this is precisely the number of infinite constants appearing in diverging amplitudes of the theory. The Feynman rules for the renormalized theory are:

$$\begin{aligned} \text{---} \xrightarrow{p} \text{---} &= \frac{i}{p^2 - m^2 + i\epsilon} & \text{---} \xrightarrow{p} \otimes \text{---} &= i(p^2 \delta_Z - \delta_m) \\ \text{---} \times \text{---} &= -i\lambda & \text{---} \otimes \text{---} &= -i\delta_\lambda \end{aligned}$$

To make the definition of the counterterms effective, the physical conditions which fix Z , m and λ must be set, in order to set their value via comparison with experiments. The three **renormalization conditions** are:

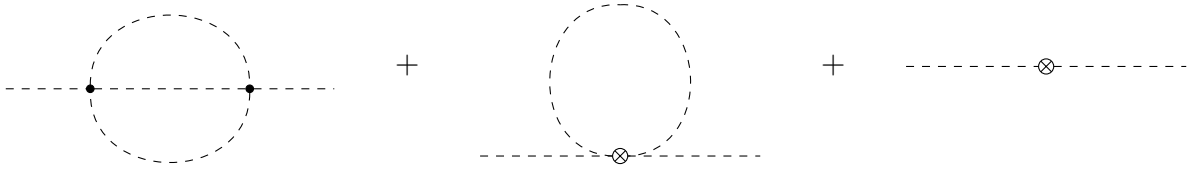
$$\text{---} \xrightarrow{p} \text{---} \text{---} \text{---} = \frac{i}{p^2 - m^2} + (\text{terms regular at } p^2 = m^2)$$



$$= -i\lambda \quad \text{at} \quad s = 4m^2, \quad t = u = 0$$

These physical conditions fix the pole and the residue of the 2-point Green function and the zero-momentum $2 \rightarrow 2$ scattering amplitude: once loop integrals have been regularized, the δ_Z , δ_m and δ_λ are adjusted (renormalized) to satisfy the renormalization conditions and, after this adjustment, the amplitudes should be finite and independent of the regulator. This marks the successful renormalization of the theory⁵.

At $o(\lambda)$ order only δ_λ is non-trivial, as δ_Z and δ_m come into play only at $o(\lambda^2)$ with diagrams:



This means that:

$$\delta_\lambda = 1 + o(\lambda) \quad \delta_Z = 1 + o(\lambda^2) \quad \delta_m = 1 + o(\lambda^2) \quad (6.42)$$

agreeing with §6.1.2.1. Tadpoles produced by the physical vertex (not the counterterm) should be considered too, but, as per Lemma 3.3.2, they only contribute as a shift between m and m_0 .

⁵Note that in the $\overline{\text{MS}}$ scheme the physical conditions are replaced by subtraction of poles: this means that the $\overline{\text{MS}}$ mass $m_{\overline{\text{MS}}}$ is not the pole of the 2-point Green function, hence it is not physically observable.

§6.2 Renormalization of QED

Recall the QED Lagrangian:

$$\mathcal{L}_{\text{QED}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \bar{\Psi}(\mathrm{i}\not{\partial} - m)\Psi - e\bar{\Psi}\not{A}\Psi \quad (6.43)$$

which is implicitly assumed to be normal-ordered. The interaction between the fermion field Ψ and the gauge field A_μ is expressed by the interaction Lagrangian $\mathcal{L}_{\text{int}} = -e\bar{\Psi}\not{A}\Psi$, which depends on the coupling constant e .

Observation 6.2.1 (Renormalizability of QED)

Consider the QED Lagrangian Eq. 6.43 in d dimensions. Then, imposing $[\mathcal{L}_{\text{QED}}] = d$ and $[m] = [\not{\partial}] = 1$, it is trivial to see that:

$$[A_\mu] = \frac{d-2}{2} \quad [\Psi] = \frac{d-1}{2} \quad [e] = \frac{4-d}{2} \quad (6.44)$$

which hints to the fact that QED is renormalizable only for $d \leq 4$. To verify this, compute the superficial degree of divergence of a diagram with L loops, P_b bosonic propagators and P_f fermionic propagator:

$$D = dL - 2P_b - P_f$$

Recalling Lemma 6.1.3, the loop number can be expressed as $L = P_b + P_f - V + 1$, so:

$$D = d - dV + (d-2)P_b + (d-1)P_f$$

Finally, noting that each vertex has 2 fermionic lines and 1 bosonic line, i.e. $2V = 2P_f + N_b$ and $V = 2P_b + N_b$, the following relation is found:

$$D = d + \frac{d-4}{2}V - \frac{d-2}{2}N_b - \frac{d-1}{2}N_f \quad (6.45)$$

Clearly, then, if $d > 4$ there is an infinite number of classes of diverging diagrams, since D increases as V increases. On the contrary, for $d = 4$ QED is renormalizable, as there are only finitely-many classes of diverging diagrams (those with $N_b + \frac{3}{2}N_f \leq 4$), while for $d < 4$ it is super-renormalizable, as there are only finitely-many diverging diagrams.

By Eq. 3.7, the S -matrix element for a generic scattering process $|i\rangle \rightarrow |f\rangle$ is $S_{if} = \delta_{if} + iT_{if}$. Assuming a non-trivial scattering, the significant quantity is the T -matrix element, which can be computed from the LSZ formula and is found to be (Eq. 3.63):

$$iT_{if} = (2\pi)^4 \delta^{(4)}(p_f - p_i) i\mathcal{M}_{if} \quad (6.46)$$

where the matrix element \mathcal{M}_{if} receives contributions from all the diagrams arising from the possible contractions determined by Wick's theorem (Th. 3.3.3 and Th. 3.3.4). In particular, due to Eq. 3.48, this contributions arrange themselves in a power series of the coupling constant, hence the general matrix element can be written as:

$$\mathcal{M}_{if} = \sum_{n \in \mathbb{N}} \mathcal{M}_{if}^{(n)} \quad (6.47)$$

with $\mathcal{M}_{if}^{(n)} \propto e^n$. Note that the power series is over \mathbb{N} and not \mathbb{N}_0 , since diagrams without interactions are disconnected (except for propagators) and do not contribute to S_{if} .

Example 6.2.1 (Corrections to the photon propagator)

Consider the propagator of the gauge field A_μ . In the free theory, it is simply given by Eq. 4.25, while in an interacting theory it becomes:

$$\Delta_{\mu\nu}(x-y) \equiv \langle 0 | \mathfrak{T} \{ A_\mu(x) A_\nu(y) \} | 0 \rangle = \sum_{n=0}^{\infty} \frac{i^n}{n!} \langle 0 | \mathfrak{T} \{ A_{I\mu}(x) A_{I\nu}(y) \mathcal{L}_I(x_1) \dots \mathcal{L}_I(x_n) \} | 0 \rangle$$

where Eq. 3.48 was used (with the conventional normalization to ignore vacuum-to-vacuum diagrams). For ease of notation, in the following all fields are understood to be in the interaction picture. Specializing the computation for QED:

$$\Delta_{\mu\nu}(x-y) = \sum_{n=0}^{\infty} \frac{(-ie)^n}{n!} \int dx_1 \dots dx_n \langle 0 | \mathfrak{T} \{ A_\mu(x) A_\nu(y) j^{\alpha_1}(x_1) A_{\alpha_1}(x_1) \dots j^{\alpha_n}(x_n) A_{\alpha_n}(x_n) \} | 0 \rangle$$

where the fermionic vector current is defined as $j^\mu := \bar{\Psi} \gamma^\mu \Psi$. Now, it is possible to apply Wick's theorem: e.g. up to $o(e^2)$ order:

$$\Delta_{\mu\nu}(x-y) = \mu \text{---} \text{---} \text{---} \nu + \mu \text{---} \text{---} \text{---} \text{---} \text{---} \nu + o(e^4)$$

since, due to the form of the interaction vertex $\mathcal{L}_{\text{int}} \sim A_\mu$, only terms with $n \in \mathbb{N}_0$ even give non-vanishing contributions. Therefore, all power series in QED can be expressed either in terms of the coupling constant e or of the structure constant $\alpha \equiv \frac{e^2}{4\pi}$: indeed, it is useful to define the following notation for $n \in \mathbb{N}_0$ (highest number of loops in the corrections):

$$\Delta_{\mu\nu}(x-y) \equiv \Delta_{\mu\nu}^{[n]}(x-y) + o(\alpha^{n+1}) \quad (6.48)$$

Note that the power series for the propagator can be rearranged as follows:

$$\Delta_{\mu\nu} = \text{---} \text{---} \text{---} + \text{---} \text{---} \text{---} + \text{---} \text{---} \text{---} + \dots \quad (6.49)$$

$$\text{---} \text{---} \text{---} = \text{---} \text{---} \text{---} + \text{---} \text{---} \text{---} + \text{---} \text{---} \text{---} + \dots \quad (6.50)$$

Hence, each blob is itself a power series in α .

From Eq. 6.45 it is possible to state all relevant classes of diverging diagrams in $d = 4$ QED:

$$\text{---} \text{---} \text{---} \sim \Lambda^2 \quad \text{---} \text{---} \text{---} \sim \Lambda \quad \text{---} \text{---} \text{---} \sim \log \Lambda$$

These are respectively the corrections to the photon propagator (collectively denoted as vacuum polarization), to the electron propagator (denoted as electron self-energy) and to the electron-photon interaction vertex.

Therefore, since there are three classes of diverging diagram, the QED Lagrangian is expected to be renormalized with three counterterms.

§6.2.1 Diverging amplitudes

§6.2.1.1 Interaction vertex

First, consider the electron-photon vertex, which at 1-loop order can be written as:

$$\mu \xrightarrow{q} \text{vertex} = \mu \xrightarrow{q} \text{vertex} + \mu \xrightarrow{q} \text{vertex} \equiv -ie\Gamma_\mu^{[1]}(q)$$

where the dependence on q only is to be proved. Note that vacuum polarization and electron self-energy diagrams have been omitted: these are indeed only corrections to the photon and electron propagators, and are not relevant to determine the vertex counterterm.

Lemma 6.2.1 (Generalized vertex function)

The electron-photon vertex function can be expressed as:

$$\Gamma^\mu(q) = \gamma^\mu F_1(q^2) + \frac{i\sigma^{\mu\nu}}{2m} q_\nu F_2(q^2) \quad (6.51)$$

where $F_1(q^2)$ is the **Dirac form factor** and $F_2(q^2)$ is the **Pauli form factor**.

Proof. Denoting the incoming and outgoing fermionic momenta with p and p' , in general Γ^μ is expressed in terms of γ^μ , since $\Gamma_\mu^{[0]} = \gamma_\mu$, of the 4-vectors p and p' and of the constants e and m . This is an exhaustive list, since no other objects appear in the Feynman rules for the diagrams involved^a.

The form of Γ^μ is further constrained by Lorentz invariance. Indeed, it must transform as a 4-vector (in the sense of γ^μ , recall Th. 1.2.1), hence it can only be a linear combination of γ^μ , p^μ and p'^μ , or, more conveniently, $p'^\mu + p^\mu$ and $p'^\mu - p^\mu$:

$$\Gamma^\mu = A\gamma^\mu + B(p' + p)^\mu + C(p' - p)^\mu$$

A priori, the coefficients could contain \not{p} and \not{p}' , but $\not{p}u(p) = mu(p)$ and $\bar{u}(p')\not{p}' = m\bar{u}(p')$ (note that Γ^μ only has sense between Dirac spinors), hence it is possible to restrain the form of the coefficients to only contain ordinary scalars: the only non-trivial scalar available is $q^2 = 2m^2 - 2p \cdot p'$, hence the coefficients can only be functions of q^2 .

Now, make use of the Ward identity $q_\mu \Gamma^\mu = 0$. The second term trivially vanishes, as does the first one since $\bar{u}(p')(\not{p}' - \not{p})u(p) = \bar{u}(p')(m - m)u(p) = 0$: the only remaining term is the third one, which does not identically vanish, therefore $C \equiv 0$.

Finally, it is customary to rewrite the B coefficient using the Gordon identity:

$$\bar{u}(p')\gamma^\mu u(p) = \bar{u}(p') \left[\frac{p'^\mu + p^\mu}{2m} + \frac{i\sigma^{\mu\nu}}{2m} q_\nu \right] u(p) \quad (6.52)$$

so that the vertex function is expressed in the same form as in the thesis. \square

^aIn principle $\epsilon^{\mu\nu\rho\sigma}$ (or equivalently γ^5) could be present too, but it is forbidden by the fact that QED is a parity-conserving theory.

The vertex function, and in particular the form factors, contains all the information on how an electron is influenced by an electromagnetic field. In presence of an external classical field $A_\mu^{\text{cl}}(x)$, the QED Lagrangian is modified with an additional interaction term $-eA_\mu^{\text{cl}}j^\mu$, and the S -matrix element for the scattering of an electron from this field is:

$$i\mathcal{M}(2\pi)\delta(p^0 - p'^0) = -ie\bar{u}(p')\Gamma^\mu(q)u(p)\tilde{A}_\mu^{\text{cl}}(q)$$

where $\tilde{A}_\mu^{\text{cl}}(q)$ is the Fourier transform of $A_\mu^{\text{cl}}(x)$. This expression is particularly useful to compute the Coulomb scattering amplitude of a non-relativistic electron from static fields.

Electrostatic field Suppose $A_\mu^{\text{cl}}(x) = (\phi(\mathbf{x}), \mathbf{0})$. Then $\tilde{A}_\mu^{\text{cl}}(q) = ((2\pi)\delta(q^0)\tilde{\phi}(\mathbf{q}), \mathbf{0})$, and the matrix element becomes:

$$i\mathcal{M} = -ie\bar{u}(p')\Gamma^0(q)u(p)\tilde{\phi}(\mathbf{q})$$

Moreover, assume that $\phi(\mathbf{x})$ is slowly varying over a large (even macroscopic) region, so that $\tilde{\phi}(\mathbf{q})$ is peaked at $\mathbf{q} = \mathbf{0}$. This allows to take the limit $\mathbf{q} \rightarrow \mathbf{0}$ (i.e. $q \rightarrow 0$, since $q^0 = 0$) in the spinor matrix element, meaning that only the Dirac form factor contributes:

$$\bar{u}(p')\gamma^0 u(p) = u^\dagger(p')u(p) \approx 2E_p \xi'^\dagger \xi \approx 2m \xi'^\dagger \xi$$

where Eq. 1.108 in the limit $p' - p \rightarrow 0$ was used. Thus, the matrix element is expressed as:

$$i\mathcal{M} \approx -ieF_1(0)\tilde{\phi}(\mathbf{q}) \cdot 2m\xi'^\dagger \xi$$

which is the Born approximation⁶ for the non-relativistic scattering from a potential of the form:

$$V(\mathbf{x}) = -eF_1(0)\phi(\mathbf{x}) \quad (6.53)$$

This means that $F_1(0)$ is the electric charge of the electron, in units of $-e$: hence, it must be $F_1(0) = 1$. Note that $\Gamma^{[0]}(q) = \gamma^\mu$, i.e. at leading order already $F_1^{[0]}(q^2) = 1$, thus radiative corrections to the electric form factor must vanish at $q^2 = 0$.

Magnetostatic field Suppose $A_\mu^{\text{cl}}(x) = (0, A_i^{\text{cl}}(\mathbf{x}))$. Then $\tilde{A}_\mu^{\text{cl}}(q) = (0, (2\pi)\delta(q^0)\tilde{A}_i^{\text{cl}}(\mathbf{q}))$, and the matrix element becomes:

$$i\mathcal{M} = -ie\bar{u}(p')\Gamma^i(q)u(p)\tilde{A}_i^{\text{cl}}(\mathbf{q}) = -ie\bar{u}(p') \left[\gamma^i F_1(q^2) + \frac{i\sigma^{i\nu} q_\nu F_2(q^2)}{2m} \right] u(p)\tilde{A}_i^{\text{cl}}(\mathbf{q})$$

In the non-relativistic limit, keeping only terms which are linear in momenta:

$$u(p) \approx \sqrt{m} \left(\begin{pmatrix} 1 - \frac{\mathbf{p} \cdot \boldsymbol{\sigma}}{2m} \\ 1 + \frac{\mathbf{p} \cdot \boldsymbol{\sigma}}{2m} \end{pmatrix} \xi \right) \implies \bar{u}(p')\gamma^i u(p) \approx 2m\xi'^\dagger \left[\frac{\mathbf{p}' \cdot \boldsymbol{\sigma}}{2m} \sigma^i + \sigma^i \frac{\mathbf{p} \cdot \boldsymbol{\sigma}}{2m} \right] \xi$$

This expression for the electric term can be further expanded using $\sigma^i \sigma^j = \delta^{ij} \mathbf{I}_2 + i\epsilon^{ijk} \sigma^k$, which result in a spin-independent term proportional to $(p' + p)^i \mathbf{I}_2$ and a spin-dependent term

⁶In general, the Born approximation for the scattering of an electron from a static potential $V(\mathbf{x})$ consists in modelling the incoming and outgoing electron asymptotic states as plane waves of defined momentum, so that the matrix element reads (with $\mathbf{q} \equiv \mathbf{p}_f - \mathbf{p}_i$):

$$\mathcal{M} = \langle \mathbf{p}_f, s_f | V(\mathbf{x}) | \mathbf{p}_i, s_i \rangle = \bar{u}^{s_f}(p_f) \int d^3x e^{i(\mathbf{p}_f - \mathbf{p}_i) \cdot \mathbf{x}} V(\mathbf{x}) u^{s_i}(p_i) = \bar{u}^{s_f}(p_f) \tilde{V}(\mathbf{q}) u^{s_i}(p_i) \quad (6.54)$$

proportional to $\epsilon^{ijk}(p' - p)^j \sigma^k$. The first term contributes to the usual kinetic operator from non-relativistic quantum mechanics and can be ignored in this treatment, hence only the spin-dependent term is retained:

$$\bar{u}(p') \gamma^i u(p) \approx 2m \xi'^{\dagger} \left[\frac{-i}{2m} \epsilon^{ijk} q^j \sigma^k \right] \xi$$

The magnetic term is already linear in momentum, hence it can be approximated with only the leading term in the expansion for the spinors:

$$\bar{u}(p') \frac{i\sigma^{i\nu}}{2m} q_{\nu} u(p) \approx m \begin{pmatrix} \xi'^{\dagger} & \xi'^{\dagger} \end{pmatrix} \frac{i}{2m} \epsilon^{ijk} \begin{bmatrix} \sigma^k & 0 \\ 0 & \sigma^k \end{bmatrix} q_j \begin{pmatrix} \xi \\ \xi \end{pmatrix} = 2m \xi'^{\dagger} \left[\frac{-i}{2m} \epsilon^{ijk} q^j \sigma^k \right] \xi$$

Thus, the linear spin-dependent contribution to the vertex function reads, as $q \rightarrow 0$:

$$i\mathcal{M} \approx -2ime \xi'^{\dagger} \frac{-i}{2m} \epsilon^{ijk} q^j \sigma^k [F_1(0) + F_2(0)] \xi \tilde{A}_i^{\text{cl}}(\mathbf{q}) = 2ime \xi'^{\dagger} \frac{\sigma^k}{2m} [F_1(0) + F_2(0)] \xi \tilde{B}^k(\mathbf{q})$$

where $\tilde{B}^k(\mathbf{q}) \equiv i\epsilon^{ijk} q^j \tilde{A}_{\text{cl}}^k(\mathbf{q})$ is the Fourier transform of the magnetic field generated by the classical vector potential $\mathbf{A}^{\text{cl}}(\mathbf{x})$. Again, this matrix element can be interpreted as the Born approximation for the non-relativistic scattering of an electron from a potential of the form:

$$V(\mathbf{x}) = -\boldsymbol{\mu} \cdot \mathbf{B}(\mathbf{x}) \quad (6.55)$$

with the electron magnetic moment defined as:

$$\boldsymbol{\mu} = -\frac{e}{m} [F_1(0) + F_2(0)] \frac{\boldsymbol{\sigma}}{2} \equiv -g \frac{e}{2m} \mathbf{S} \quad (6.56)$$

where $\mathbf{S} \equiv \frac{\boldsymbol{\sigma}}{2}$. The **Landé factor** is the found to be:

$$g = 2[F_1(0) + F_2(0)] = 2 + 2F_2(0) \quad (6.57)$$

Since $F_2^{[0]}(0) = 0$, at leading order $g^{[0]} = 2$. To compute the first-order correction to the Landé factor, it is necessary to compute the magnetic form factor at leading order.

Proposition 6.2.1 (Form factors)

At 1-loop order in dimensional regularization, the Dirac and Pauli form factors are:

$$F_1^{[1]}(q^2) = 1 + \frac{\alpha}{2\pi} (4\pi)^{\epsilon} \Gamma(\epsilon) \int_0^1 dx \int_0^1 dy \int_0^1 dz \delta(x+y+z-1) \mathcal{K}(\epsilon, x, y, z, q^2, m^2) \Delta^{-\epsilon} \quad (6.58)$$

$$F_2^{[1]}(q^2) = \frac{\alpha}{\pi} \int_0^1 dx \int_0^1 dy \int_0^1 dz \delta(x+y+z-1) \frac{m^2 z(1-z)}{m^2(1-z)^2 - xyq^2} \quad (6.59)$$

with $\Delta \equiv m^2(1-z)^2 - xyq^2$ and the ϵ -finite function \mathcal{K} defined as:

$$\mathcal{K}(\epsilon, x, y, z, q^2, m^2) \equiv \left[1 - \epsilon + 2 \frac{\Gamma(1-\epsilon)}{\Gamma(\epsilon)} \frac{(1-x)(1-y)q^2 + (1-4z+z^2)m^2}{m^2(1-z)^2 - xyq^2} \right] \quad (6.60)$$

Proof. Copy proof from OneNote

□

Note that, as long as Eq. 6.63 is valid (which is just the Ward identity $k_\mu \Pi^{\mu\nu}(k) = 0$), then $\Pi(k^2)$ is regular on the light-cone $k^2 = 0$, which means that the exact photon propagator has a simple pole at $k^2 = 0$ and the photon remains massless at all orders in perturbation theory. On the contrary, if $\Pi^{\mu\nu}(k)$ contained a term like $M^2 \eta^{\mu\nu}$ without a compensating $k^\mu k^\nu$ term, then the photon mass would be shifted to M .

Proposition 6.2.2 (Vacuum polarization)

At 1-loop order in dimensional regularization, the **vacuum polarization** is:

$$\Pi^{[1]}(q^2) = -\frac{2\alpha}{\pi} (4\pi)^\epsilon \Gamma(\epsilon) \int_0^1 dx x(1-x) \Delta^{-\epsilon} \quad (6.64)$$

with $\Delta \equiv m^2 - x(1-x)q^2$.

Proof. At 1-loop order, Eq. 6.50 reduces to:

Hence, recalling that fermionic loops determine an additional (-1) factor:

$$\begin{aligned} i\Pi_{\mu\nu}^{[1]}(q^2) &= (-ie)^2 (-1) \int \frac{d^4 k}{(2\pi)^4} \frac{\text{tr}\{\gamma_\mu i(\not{k} + m) \gamma_\nu i(\not{k} + \not{q} + m)\}}{(k^2 - m^2 + i\epsilon)((k+q)^2 - m^2 + i\epsilon)} \\ &= -4e^2 \int \frac{d^4 k}{(2\pi)^4} \frac{2k_\mu k_\nu + k_\mu q_\nu + q_\mu k_\nu - \eta_{\mu\nu}(k^2 + k \cdot q - m^2)}{(k^2 - m^2 + i\epsilon)(k^2 + 2k \cdot q + q^2 - m^2 + i\epsilon)} \end{aligned}$$

Introducing Feynman parameters (Lemma 6.1.1):

$$\begin{aligned} \Pi_{\mu\nu}^{[1]}(q^2) &= 4ie^2 \int_0^1 dx \int \frac{d^4 \ell}{(2\pi)^4} [2\ell_\mu \ell_\nu + (1-2x)(\ell_\mu q_\nu + \ell_\nu q_\mu) - 2x(1-x)q^\mu q^\nu + \\ &\quad - \eta_{\mu\nu}(\ell^2 + (1-2x)\ell \cdot q - x(1-x)q^2 - m^2)] (\ell^2 - \Delta + i\epsilon)^{-2} \end{aligned}$$

where the integration has been shifted to $\ell = k + xq$, with $\Delta \equiv m^2 - x(1-x)q^2$. Note that both the integration measure and the denominator are even under $\ell \rightarrow -\ell$, hence all terms with an odd numerator vanish (i.e. the linear terms in ℓ , in this case). Then, the integral can be broken up into two contributions:

$$\begin{aligned} \Pi_{\mu\nu}^{[1]}(q^2) &= 4ie^2 \int_0^1 dx \left[\int \frac{d^4 \ell}{(2\pi)^4} \frac{2\ell_\mu \ell_\nu - \eta_{\mu\nu} \ell^2}{(\ell^2 - \Delta + i\epsilon)^2} + \right. \\ &\quad \left. + (\eta_{\mu\nu}(m^2 + x(1-x)q^2) - 2x(1-x)q^\mu q^\nu) \int \frac{d^4 \ell}{(2\pi)^4} \frac{1}{(\ell^2 - \Delta + i\epsilon)^2} \right] \end{aligned}$$

The second integral is in the form of Eq. 6.6, while the first one requires additional treatment.

Lemma 6.2.3

$$\int \frac{d^d k}{(2\pi)^d} \frac{k^{2m}}{(k^2 - \Delta + i\epsilon)^n} = \frac{i(-1)^{n+m}}{(4\pi)^{d/2}} \frac{\Gamma(n - \frac{d}{2} - m) \Gamma(\frac{d}{2} + m)}{\Gamma(\frac{d}{2}) \Gamma(n)} \Delta^{-n+\frac{d}{2}+m} \quad (6.65)$$

Proof. Performing a Wick rotation:

$$\int \frac{d^d k}{(2\pi)^d} \frac{k^{2m}}{(k^2 - \Delta + i\epsilon)^n} = i(-1)^{n+m} \int \frac{d^d \ell_E}{(2\pi)^d} \frac{\ell_E^{2m}}{(\ell_E^2 + \Delta)^n}$$

Parametrizing the Euclidean domain of integration \mathbb{R}^d with spherical coordinates and using Eq. A.41:

$$\int \frac{d^d k}{(2\pi)^d} \frac{k^{2m}}{(k^2 - \Delta + i\epsilon)^n} = \frac{2i(-1)^{n+m}}{(4\pi)^{d/2} \Gamma(\frac{d}{2})} \int_{\mathbb{R}^+} d\ell \frac{\ell^{2m+d-1}}{(\ell^2 + \Delta)^n}$$

Recalling Eq. 6.7 yields the thesis. \square

Now, recall the following identity for tensorial momentum integrals, valid for a general integrable $f(k^2)$:

$$\int d^d k k^\mu k^\nu f(k^2) = \frac{1}{d} \eta^{\mu\nu} \int d^d k k^2 f(k^2) \quad (6.66)$$

which is a consequence of Lorentz invariance. With this identity, the numerator of the first integral becomes $-\eta_{\mu\nu}(1 - 2/d)\ell^2$, so that, by Eq. 6.65:

$$\Pi_{\mu\nu}^{[1]}(q^2) = 8e^2(q_\mu q_\nu - \eta_{\mu\nu} q^2) \int_0^1 dx x(1-x) \frac{\Gamma(2 - \frac{d}{2})}{(4\pi)^{d/2}} \Delta^{d/2-2} \equiv (\eta_{\mu\nu} q^2 - q_\mu q_\nu) \Pi(q^2)$$

where the transverse projector has been factored out. Then:

$$\Pi^{[1]}(q^2) = -8e^2 \int_0^1 dx x(1-x) \frac{\Gamma(2 - \frac{d}{2})}{(4\pi)^{d/2}} \Delta^{d/2-2}$$

Inserting $d = 4 - 2\epsilon$ yields the thesis. \square

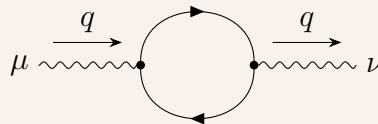
The Laurent series in ϵ of the vacuum polarization then reads:

$$\Pi^{[1]}(q^2) = -\frac{\alpha}{3\pi} \frac{1}{\epsilon} + o(\epsilon^0) \quad (6.67)$$

This clearly shows that the pole term at 1-loop order does not depend on the energy scale q^2 , which is a feature of all loop orders⁸.

Observation 6.2.2 (Vacuum polarization and charge screening)

Consider corrections to the photon propagator like:



The physical interpretation of these corrections is that the vacuum is filled with virtual electric dipoles, hence the name of **vacuum polarization**. In particular, in presence of an

⁸This is to be expected: in fact, the divergences encapsulated in the ϵ -poles are UV divergences, and in the UV limit two energy scales q^2 and μ^2 become equal, since they both $q^2, \mu^2 \rightarrow \infty$.

electric charge, e.g. a proton, these virtual dipoles orient themselves according to said charge, thus resulting in a screening of the Coulomb potential: this phenomenon is known as *charge screening*.

§6.2.1.3 Electron self-energy

Finally, take the corrections to the electron propagator. They are analogous to those in Eq. 6.49, but in this case the structure of geometric series is manifest, and they can be resummed in the exact electron propagator as:

$$\tilde{S}(p) = \frac{i}{\not{p} - m - \Sigma(p)} \quad (6.68)$$

Proposition 6.2.3 (Electron self-energy)

At 1-loop order in dimensional regularization, the **electron self-energy** is:

$$\Sigma^{[1]}(p) = -\frac{\alpha}{2\pi}(4\pi)^\epsilon \Gamma(\epsilon) \int_0^1 dx ((1-x)\not{p} - 2m) \Delta^{-\epsilon} \quad (6.69)$$

with $\Delta \equiv xm^2 - x(1-x)p^2$.

Proof. At 1-loop order, the only correction to the electron propagator is:

$$\text{Diagram} \equiv -i\Sigma^{[1]}(p)$$

With the Feynman rules for QED:

$$\begin{aligned} -i\Sigma^{[1]}(p) &= (-ie)^2 \int \frac{d^4k}{(2\pi)^4} \frac{\gamma^\mu i(\not{k} + \not{p} + m) \gamma_\mu (-i)}{((k+p)^2 - m^2 + i\epsilon)(k^2 + i\epsilon)} \\ &= 2e^2 \int \frac{d^4k}{(2\pi)^4} \frac{\not{k} + \not{p} - 2m}{(k^2 + 2k \cdot p + p^2 - m^2 + i\epsilon)(k^2 + i\epsilon)} \end{aligned}$$

Introducing Feynman parameters (Lemma 6.1.1):

$$\Sigma^{[1]}(p) = 2ie^2 \int_0^1 dx \int \frac{d^4\ell}{(2\pi)^4} \frac{\ell + (1-x)\not{p} - 2m}{(\ell^2 - \Delta + i\epsilon)^2}$$

where the integration has been shifted to $\ell \equiv k + xq$, with $\Delta \equiv xm^2 - x(1-x)p^2$. Since both the integration measure and the denominator are even under $\ell \rightarrow -\ell$, the first term of the numerator vanishes, leaving only:

$$\Sigma^{[1]}(p) = 2ie^2 \int_0^1 dx ((1-x)\not{p} - 2m) \int \frac{d^4\ell}{(2\pi)^4} \frac{1}{(\ell^2 - \Delta + i\epsilon)^2}$$

The loop integral can now be evaluated with Eq. 6.6, and the thesis is trivially found. \square

Like the electric form factor, the electron self-energy too presents an IR divergence on the light-cone $k^2 = 0$, which here is manifest as a divergence of the integral over the Feynman parameter x , and it can again be regulated introducing a small photon mass μ . The IR divergence,

however, does not affect the UV-divergent part (as it is to be expected), which can be isolated with a Laurent series in ϵ :

$$\Sigma^{[1]}(\not{p}) = -\frac{1}{\epsilon} \frac{\alpha}{4\pi} (\not{p} - 4m) + o(\epsilon^0) \quad (6.70)$$

This completes the computation of diverging diagrams in QED: the electron-photon vertex in Eq. 6.61, the vacuum polarization in Eq. 6.67 and the electron self-energy in Eq. 6.70.

§6.2.2 Renormalization conditions

At this point, the UV divergences of QED have been regularized, i.e. expressed in terms of a regulator (as poles in ϵ , in this case). There remains to fix renormalization conditions in order to absorb these divergences into unobservable bare parameters. To do so, define the renormalized (i.e. physical) field-strengths, mass and charge as:

$$\Psi = \sqrt{Z_2} \Psi_r \quad A^\mu = \sqrt{Z_3} A_r^\mu \quad m = Z_m m_r \quad e = \frac{Z_1}{Z_2 \sqrt{Z_3}} e_r \quad (6.71)$$

so that the bare Lagrangian Eq. 6.43 can be rewritten as in terms of finite renormalized quantities and infinite renormalization coefficients:

$$\mathcal{L}_{\text{QED}} = -\frac{1}{4} Z_3 F_r^{\mu\nu} F_{\mu\nu}^r + Z_2 \bar{\Psi}_r (i\not{\partial} - Z_m m_r) \Psi_r - Z_1 e_r \bar{\Psi}_r \not{A}_r \Psi_r \quad (6.72)$$

It is convenient to work with renormalization counterterms, which are defined from the renormalization coefficients as:

$$Z_1 = 1 + \delta_1 \quad Z_2 = 1 + \delta_2 \quad Z_3 = 1 + \delta_3 \quad Z_2 Z_m = 1 + \frac{\delta_m}{m_r} \quad (6.73)$$

Writing the Lagrangian counterterms explicitly:

$$\begin{aligned} \mathcal{L}_{\text{QED}} = & -\frac{1}{4} F_r^{\mu\nu} F_{\mu\nu}^r + \bar{\Psi}_r (i\not{\partial} - m_r) \Psi_r - e_r \bar{\Psi}_r \not{A}_r \Psi_r + \\ & + \frac{1}{2} \delta_3 A_r^\mu (\eta_{\mu\nu} \square - \partial_\mu \partial_\nu) A_r^\nu + \bar{\Psi}_r (i\delta_2 \not{\partial} - \delta_m) \Psi_r - \delta_1 e_r \bar{\Psi}_r \not{A}_r \Psi_r \end{aligned} \quad (6.74)$$

The Feynman rules for this Lagrangian are trivially found. Assuming always to work in Feynman gauge:

$$\begin{aligned} \mu \xrightarrow{k} \nu &= \frac{-i\eta_{\mu\nu}}{k^2 + i\epsilon} & \xrightarrow{p} &= \frac{i}{\not{p} - m_r + i\epsilon} & \mu \xrightarrow{k} \begin{array}{c} \nearrow \\ \searrow \end{array} &= -ie_r \gamma^\mu \\ \mu \xrightarrow{k} \otimes \nu &= -i(\eta_{\mu\nu} k^2 - k_\mu k_\nu) \delta_3 & \xrightarrow{p} \otimes &= i(\delta_2 \not{p} - \delta_m) \\ \mu \xrightarrow{k} \otimes \begin{array}{c} \nearrow \\ \searrow \end{array} &= -ie_r \gamma^\mu \delta_1 \end{aligned}$$

Now, these counterterms must be chosen in order to cancel the divergences of the three classes of diverging diagrams, i.e. the photon and electron propagators and the interaction vertex.

§6.2.2.1 Electron propagator

The geometric series to compute the exact electron propagator is still valid, however now the blob correction contains 1PI diagrams with both “normal” vertices and “counter”-vertices, i.e. it has the form:

Then, it is useful to generalize the electron self-energy introducing:

$$-i\Sigma_r(\not{p}) \equiv i(\delta_2 \not{p} - \delta_m) - i\Sigma(\not{p}) \quad (6.75)$$

so that the exact electron propagator is written as:

$$\tilde{S}(\not{p}) = \frac{i}{\not{p} - m_r - \Sigma_r(\not{p})} \quad (6.76)$$

Now, the relevant renormalization conditions imposed on the electron propagator are on its pole and the related residue:

$$\tilde{S}^{-1}(m_p) = 0 \quad i \frac{\partial}{\partial \not{p}} \tilde{S}^{-1}(m_p) = 1 \quad (6.77)$$

where the pole mass m_p is determined by the subtraction scheme adopted, i.e. by how the finite parts of the counterterms are defined.

On-shell scheme In the on-shell scheme, the pole mass is taken to be the physical mass of the electron, i.e. $m_p = m_r$. Therefore, the above conditions translate to:

$$\Sigma_r(m_r) = 0 \quad \Sigma'_r(m_r) = 0 \quad (6.78)$$

which allow to define both δ_m and δ_2 . At 1-loop order, recalling Eq. 6.70, the divergent part of the counterterms is found solving the above system:

$$\delta_2^{[1]} = -\frac{1}{\epsilon} \frac{\alpha_r}{4\pi} + o(\alpha_r \epsilon^0) \quad \delta_m^{[1]} = -\frac{1}{\epsilon} \frac{\alpha_r}{\pi} + o(\alpha_r \epsilon^0) \quad (6.79)$$

From Eq. 6.73, noting that $(1 + \delta_2)^{[1]} = 1 - \delta_2^{[1]}$, the renormalization coefficients are found to be:

$$Z_2^{[1]} = 1 - \frac{1}{\epsilon} \frac{\alpha_r}{4\pi} + o(\alpha_r \epsilon^0) \quad Z_m^{[1]} = 1 - \frac{1}{\epsilon} \frac{3\alpha_r}{4\pi} + o(\alpha_r \epsilon^0) \quad (6.80)$$

Minimal subtraction scheme In the MS scheme, the counterterms only contain a divergent part, hence, by Eq. 6.75:

$$\delta_2^{[1]} = -\frac{1}{\epsilon} \frac{\alpha_r}{4\pi} \quad \delta_m^{[1]} = -\frac{1}{\epsilon} \frac{\alpha_r}{\pi} \quad (6.81)$$

This explicitly shows that the divergent part of the counterterms is scheme-independent. The renormalization coefficients in the MS scheme thus differ from those in the on-shell scheme only by a finite quantity. Recall that it is also possible to define a modified minimal subtraction scheme ($\overline{\text{MS}}$), which additionally includes in the counterterms the universal factor $\log(4\pi e^{-\gamma_E})$ present in dimensionally-regularized calculations.

Note that in this case the pole mass $m_p = m_{\text{MS}}$ is different from the physical mass m_r of the electron, but their difference is finite and given by the $o(\alpha_r \epsilon^0)$ terms in Eq. 6.79.

§6.2.2.2 Photon propagator

The exact photon propagator too is modified by the presence of counterterms. In particular, the blob correction in this case has the form:

$$\mu \text{ --- } \text{---} \otimes \text{---} \nu + \mu \text{ --- } \text{---} \text{---} \text{---} \nu$$

Defining a generalized vacuum polarization as:

$$i\Pi_r^{\mu\nu}(k) = -i(\eta^{\mu\nu}k^2 + k^\mu k^\nu)\delta_3 + i\Pi^{\mu\nu}(k) \quad \Pi_r(k^2) = \Pi(k^2) - \delta_3 \quad (6.82)$$

the exact photon propagator can be written as:

$$\tilde{\Delta}_{\mu\nu}(k) = \frac{-i}{k^2} \eta_{\mu\nu} \frac{1}{1 - \Pi_r(k^2)} + \frac{i}{k^2} \frac{k_\mu k_\nu}{k^2} \frac{\Pi_r(k^2)}{1 - \Pi_r(k^2)} \quad (6.83)$$

In this case, the propagator already has a pole at $k^2 = 0$, hence it is sufficient to fix its residue to $-i\eta_{\mu\nu}$, i.e.:

$$\Pi_r(0) = 0 \quad (6.84)$$

Thus, $\delta_3 = \Pi(0)$, which at 1-loop can be computed from Eq. 6.64. The divergent part is:

$$\delta_3^{[1]} = -\frac{1}{\epsilon} \frac{\alpha_r}{3\pi} + o(\alpha_r \epsilon^0) \quad Z_3^{[1]} = 1 - \frac{1}{\epsilon} \frac{\alpha_r}{3\pi} + o(\alpha_r \epsilon^0) \quad (6.85)$$

§6.2.2.3 Interaction vertex

Given the definition of the renormalization coefficient for e_r in Eq. 6.71, the corrections to the electron and the photon propagator are factored out of the vertex function, hence the corrections to the interaction vertex are only given by:

$$\mu \text{ --- } \text{---} \otimes \text{---} + \mu \text{ --- } \text{---} \text{---} \text{---}$$

The resulting generalized vertex function is:

$$\Gamma_r^\mu(k) = \Gamma^\mu(k) + \gamma^\mu \delta_1 \quad (6.86)$$

It is conventional to define e_r as the electron charge measured at great distance (i.e. at $k \rightarrow 0$), hence the renormalization condition is:

$$\Gamma_r^\mu(0) = \gamma^\mu \quad (6.87)$$

Recalling Eq. 6.51, the counterterm is correctly determined by the electric form factor only (since the magnetic form factor is not divergent). At 1-loop order, from Eq. 6.61:

$$\delta_1^{[1]} = -\frac{1}{\epsilon} \frac{\alpha_r}{4\pi} + o(\alpha_r \epsilon^0) \quad Z_1 = 1 - \frac{1}{\epsilon} \frac{\alpha_r}{4\pi} + o(\alpha_r \epsilon^0) \quad (6.88)$$

§6.2.2.4 Gauge symmetry and renormalization

Comparing Eq. 6.80 and Eq. 6.88, it may seem a coincidence that the divergent parts of Z_1 and Z_2 are equal. However, this is a consequence of the gauge symmetry which defines the QED, i.e. the U(1) vector symmetry.

Lemma 6.2.4 (Ward–Takahashi identity)

Defining $\Gamma^\mu(p+k, p) \equiv \Gamma^\mu(k)$ in order to denote the incoming momentum in the fermionic current (not necessarily on-shell), the vertex function can be written as:

$$-ik_\mu \Gamma^\mu(p+k, p) = \tilde{S}^{-1}(\not{p} + \not{k}) - \tilde{S}^{-1}(\not{p}) \quad (6.89)$$

Proof. See §7.5 of [6]. □

Note that the Ward–Takahashi identity is a direct consequence of the conservation of the Noether current $j^\mu = \bar{\Psi}\gamma^\mu\Psi$ associated to the gauge symmetry of QED, i.e. $\partial_\mu j^\mu = 0$.

Proposition 6.2.4

At all loop orders:

$$Z_1 = Z_2 \quad (6.90)$$

Proof. Apply the Ward–Takahashi identity to the bare QED Lagrangian:

$$k_\mu \Gamma^\mu(p+k, p) = \not{k} - \Sigma(\not{p} + \not{k}) + \Sigma(\not{p})$$

Taking the limit of all particles going on-shell, i.e. $k \rightarrow 0$ and $p \rightarrow m_r$, and recalling that for on-shell electrons the vertex function is given by Eq. 6.51, this equation becomes:

$$F_1(0) = 1 - \Sigma'(m_r)$$

However, from the renormalization conditions⁹Eq. 6.78,6.87 it follows that, at all loop orders:

$$F_1(0) = 1 - \delta_1 \quad \Sigma'(m_r) = \delta_2$$

Hence, $\delta_1 = \delta_2$ at all loop orders, which is the thesis. □

From this equivalence, it follows that the charge is truly renormalized as $e = Z_3^{-1/2}e_r$. This has an important consequence when considering QED with multiple fermionic species (e.g. muons): the renormalized electric charge has the same universal value for all the species, since it only depends on the renormalization coefficient for the photon field.

§6.2.3 Renormalization group

Consider the Coulomb scattering $e^-e^- \rightarrow e^-e^-$, which at tree-level has the matrix element (in the t -channel):

$$\begin{array}{c} \text{Diagram: Two incoming fermion lines (arrows pointing right) meet at a vertex, exchange a photon (wavy line) with momentum q, and then split into two outgoing fermion lines (arrows pointing right).} \end{array} \equiv i\mathcal{M}_0(q^2) \sim \frac{e^2}{q^2}$$

⁹Indeed, the definitions of the renormalization coefficients given in Eq. 6.71 are equivalent to defining:

$$\Gamma^\mu(0) = Z_1^{-1}\gamma^\mu \quad \Sigma(\not{p}) \sim (1 - Z_2^{-1})\not{p} \quad (6.91)$$

where e is the bare coupling constant. Accounting for loop corrections to the photon propagator, i.e. inserting Eq. 6.62, results into the observed matrix element:

$$i\mathcal{M}(q^2) \sim \frac{e^2}{q^2} \frac{1}{1 - \Pi(q^2)}$$

This allows to define an observed “effective” electron charge $e_{\text{eff}}^2(q^2)$, which depends on the energy scale q^2 at which the experiment is performed:

$$e_{\text{eff}}^2(q^2) := \frac{e^2}{1 - \Pi(q^2)} \quad (6.92)$$

Recalling that $e_r^2 = Z_3 e^2 = e^2/(1 - \Pi(0))$, it is clear that the renormalized charge is defined as $e_r \equiv e_{\text{eff}}(0)$, i.e. as the electron charge observed at great distance, as already noted when imposing the renormalization condition for the electron-photon vertex: this is the effect of vacuum polarization, where the physical charge of the electron $e_r = 1.602176634 \cdot 10^{-19} \text{ C}$ is affected by a progressively stronger screening effect when probed at smaller distances, i.e. at larger momenta. Indeed, at 1-loop order:

$$\alpha_{\text{eff}}^{[1]}(q^2) = \alpha_r \frac{1 - \Pi^{[1]}(0)}{1 - \Pi^{[1]}(q^2)} = \frac{\alpha_r}{1 - \Pi^{[1]}(q^2) + \Pi^{[2]}(0)}$$

with $\alpha_r = e_r^2/4\pi \simeq 1/137.036$ and:

$$\Pi^{[1]}(q^2) - \Pi^{[1]}(0) = \frac{2\alpha^{[1]}}{\pi} \int_0^1 dx x(1-x) \log \left[1 - x(1-x) \frac{q^2}{m^2} \right]$$

In the relativistic limit¹⁰ $-q^2 \gg m^2$ this expression becomes, using $\alpha^{[1]} = \alpha_r$:

$$\Pi^{[1]}(q^2) - \Pi^{[1]}(0) = \frac{\alpha}{3\pi} \left[\log \frac{-q^2}{m^2} - \frac{5}{3} + o\left(\frac{m^2}{q^2}\right) \right]$$

which clearly shows that $\alpha_{\text{eff}}(q^2)$ becomes larger as $|q^2|$ increases, as predicted by the charge screening due to vacuum polarization.

Moreover, note that $\Pi(k^2) \propto e^2$, hence, defining $\pi(k^2) \equiv \Pi(k^2)/e^2$, it is possible to invert this relation to find the bare electric charge in terms of the observed one:

$$e^2 = \frac{e_{\text{eff}}^2(q^2)}{1 - e_{\text{eff}}^2(q^2)\pi(q^2)} \quad (6.93)$$

Now, consider two measurements of e_{eff} at two different scales q^2 and μ^2 . From Eq. 6.92:

$$\frac{1}{e_{\text{eff}}^2(q^2)} - \frac{1}{e_{\text{eff}}^2(\mu^2)} = \pi(\mu^2) - \pi(q^2)$$

Since the divergent part of the vacuum polarization is independent of the energy scale (Eq. 6.67), this expression is correctly finite, as it is to be expected from physically observable quantities.

¹⁰The two channels that contribute to the Coulomb scattering are the t -channel and the u -channel. In the CM frame, since $p_1 = (E_{\mathbf{p}}, \mathbf{p})$, $p_2 = (E_{\mathbf{p}}, -\mathbf{p})$, $k_1 = (E_{\mathbf{p}}, \mathbf{p}')$ and $k_2 = (E_{\mathbf{p}}, -\mathbf{p}')$, with $\|\mathbf{p}'\| = \|\mathbf{p}\|$, defining the scattering angle $\theta : \cos \theta = \hat{\mathbf{p}}' \cdot \hat{\mathbf{p}}$, the Mandelstam variables are:

$$t = -2\mathbf{p}^2(1 - \cos \theta) \leq 0 \quad u = -2\mathbf{p}^2(1 + \cos \theta) \leq 0$$

where $t = 0$ only for forward scattering ($\theta = 0$) and $u = 0$ only for backward scattering ($\theta = \pi$). In the relativistic regime $\mathbf{p}^2 \gg m^2$, hence $-t, -u \gg m^2$.

Lemma 6.2.5

In the relativistic regime $-q^2, -\mu^2 \gg m^2$:

$$\pi^{[1]}(\mu^2) - \pi^{[1]}(q^2) \simeq \frac{1}{12\pi^2} \log \frac{\mu^2}{q^2}$$

Proof. From Eq. 6.64:

$$\begin{aligned} \pi^{[1]}(\mu^2) - \pi^{[1]}(q^2) &= -\frac{1}{2\pi^2} \int_0^1 dx x(1-x) [-\log \Delta(\mu^2) + \log \Delta(q^2) + o(\epsilon)] \\ &= \frac{1}{2\pi^2} \int_0^1 dx x(1-x) \log \frac{1-x(1-x)\frac{\mu^2}{m^2}}{1-x(1-x)\frac{q^2}{m^2}} \simeq \frac{1}{2\pi^2} \int_0^1 dx x(1-x) \log \frac{\mu^2}{q^2} \end{aligned}$$

Performing the integration yields the thesis. \square

With this relation, it is possible to relate directly the measurements of the electric charge at different energy scales:

$$\alpha_{\text{eff}}^{[1]}(q^2) = \frac{\alpha_{\text{eff}}^{[1]}(\mu^2)}{1 - \frac{\alpha_{\text{eff}}^{[1]}(\mu^2)}{3\pi} \log \frac{q^2}{\mu^2}} \quad (6.94)$$

This dependence of the physical coupling on the energy scale is known as the **running of the coupling**, and it is effectively a prediction of the renormalization of the theory (which has been verified experimentally).

The above equation completely determines the dependence of the coupling on the energy scale (at 1-loop order). Switching the notation to $q \mapsto \mu$ and $\mu \mapsto \mu_0$ and differentiating with respect to $\log \mu^2$:

$$\frac{d}{d \log \mu^2} \frac{1}{\alpha_{\text{eff}}^{[1]}(\mu^2)} = -\frac{1}{3\pi} \implies \frac{d\alpha_{\text{eff}}^{[1]}}{d \log \mu^2} = \alpha_{\text{eff}}^{[1]} \frac{4}{3} \frac{\alpha_{\text{eff}}^{[1]}}{4\pi}$$

Generalizing to higher loop orders, the **renormalization group** equation is found:

$$\frac{d\alpha_{\text{eff}}(\mu^2)}{d \log \mu^2} = -\beta(\alpha_{\text{eff}}(\mu^2)) \quad (6.95)$$

where the **Callan–Symanzik β -function** is defined as:

$$\beta(\alpha) \equiv \alpha \sum_{n \in N_0} \beta_n \left(\frac{\alpha}{4\pi} \right)^{n+1} \quad (6.96)$$

Clearly, for QED $\beta_0 = -\frac{4}{3}$, which means that the coupling gets stronger as the energy increases. On the contrary, for theories with $\beta_0 > 0$ (like QCD), the coupling is weaker in the UV and stronger in the IR, a phenomenon known as *asymptotic freedom*.

Furthermore, viewing masses as additional couplings, it is possible to define a massive β -function, so that the scale-dependence of the observed mass can be written as:

$$\frac{dm_{\text{eff}}(\mu^2)}{d \log \mu^2} = -\beta_m(\alpha_{\text{eff}}(\mu^2), m_{\text{eff}}(\mu^2)) \quad (6.97)$$

Now, it is possible to study the renormalization of n -point Green functions. In particular, consider a connected 1PI Green function with k photon lines and n electron lines:

$$\tilde{G}^{(k,n)}(\{p_i\}, \alpha, m, \Lambda) \equiv \langle 0 | \mathfrak{T} \{ A(p_1) \dots A(p_k) \Psi(p_{k+1}) \dots \Psi(p_{k+n}) \} | 0 \rangle$$

where Lorentz indices and Dirac adjoints have been suppressed for brevity. Then, renormalizing the field-strengths:

$$\tilde{G}^{(k,n)}(\{p_i\}, \alpha, m, \Lambda) = Z_3^{k/2} Z_2^{n/2} \tilde{G}_r^{(k,n)}(\{p_i\}, \alpha_r, m_r, \mu)$$

where the regulator Λ is substituted by the renormalization point μ . Note that \tilde{G}_r , α_r and m_r are finite, while \tilde{G} , α , m , Z_2 and Z_3 are divergent with respect to the regulator (i.e. for $\Lambda \rightarrow \infty$ with a UV cut-off or $\epsilon \rightarrow 0$ in dimensional regularization): in particular, \tilde{G} , α and m only depend on the regulator Λ and not on the renormalization point μ , hence taking the derivative with respect to $\log \mu^2$ yields:

$$0 = \frac{k}{2} Z_3^{k/2} \frac{1}{Z_3} \frac{\partial Z_3}{\partial \log \mu^2} Z_2^{n/2} \tilde{G}_r^{(k,n)} + Z_3^{k/2} \frac{n}{2} \frac{1}{Z_2} \frac{\partial Z_2}{\partial \log \mu^2} \tilde{G}_r^{(k,n)} + Z_3^{k/2} Z_2^{n/2} \frac{d\tilde{G}_r^{(k,n)}}{d \log \mu^2}$$

Rewriting the first two terms and expanding the total derivative, this equation becomes:

$$\left[\frac{\partial}{\partial \log \mu^2} - \beta(\alpha_r) \frac{\partial}{\partial \alpha_r} - \beta_m(\alpha_r, m_r) \frac{\partial}{\partial m_r} + k\gamma_\gamma(\alpha_r) + n\gamma_e(\alpha_r, m_r) \right] \tilde{G}_r^{(k,n)}(\{p_i\}, \alpha_r, m_r, \mu) \quad (6.98)$$

which is known as the **Callan–Symanzik equation**. The two factors γ_γ and γ_e are known respectively as the photon and the electron *anomalous dimensions*; in general, the anomalous dimension of a field ϕ_i is defined as:

$$\gamma_i(\{\lambda_j\}) := \frac{\partial \log \sqrt{Z_i}}{\partial \log \mu^2} \quad (6.99)$$

where Z_i is the renormalization coefficient of the field strength $\phi_i = Z_i \phi_i^r$. Note that the anomalous dimension is independent of the regulator, contrary to Z_i , as it depends only on the various couplings of the theory which affect ϕ_i .

Example 6.2.2 (Evolution of a massless renormalizable theory)

Consider a quantum field theory with a single massless quantum field ϕ and a single interaction vertex with dimensionless coupling α , and suppose to conduct two experiments to evaluate the n -point Green function at two different energy scales, i.e. with momenta $\{p_i\}_{i=1,\dots,n}$ and $\{e^s p_i\}_{i=1,\dots,n}$, with $s \in \mathbb{R}$ a rescaling parameter. Then, defining the dimension $[\tilde{G}_r^{(n)}(\{p_i\}, \alpha_r, \mu)] \equiv d_n$, the n -point Green function classically rescales as:

$$\tilde{G}_r^{(n)}(\{e^s p_i\}, \alpha_r, \mu) = e^{d_n s} \tilde{G}_r^{(n)}(\{p_i\}, \alpha_r, e^{-s} \mu)$$

There remains to evaluate the scaling of the Green function due to the renormalization point. To do so, consider the Callan–Symanzik equation for this massless theory:

$$\left[\frac{\partial}{\partial \log \mu^2} - \beta(\alpha_r) \frac{\partial}{\partial \alpha_r} + n\gamma(\alpha_r) \right] \tilde{G}_r^{(n)}(\{p_i\}, \alpha_r, \mu) = 0$$

Now, parametrize^a this equation in the (μ, α_r) -plane so that:

$$\frac{d \log \mu^2(s)}{ds} = 1 \quad \frac{d \alpha_r(s)}{ds} = -\beta(\alpha_r(s))$$

with initial conditions $\mu(0) = \bar{\mu}$, $\alpha_r(0) = \bar{\alpha}_r$. The Callan–Symanzik equation then becomes:

$$\frac{d}{ds} \tilde{G}_r^{(n)}(s) = -n\gamma(\alpha_r(s)) \tilde{G}_r^{(n)}(s)$$

with $\tilde{G}_r^{(n)}(s) \equiv \tilde{G}_r^{(n)}(\{p_i\}, \alpha_r(s), \mu(s))$. This ODE is trivially solved as:

$$\tilde{G}_r^{(n)}(s) = \tilde{G}_r^{(n)}(0) \exp \left[-n \int_0^s dt \gamma(\alpha_r(t)) \right]$$

Since the characteristic equation for $\mu(s)$ has solution $\mu(s) = e^{s/2} \bar{\mu}$, the total scaling of the n -point Green function is:

$$\tilde{G}_r^{(n)}(\{e^s p_i\}, \alpha_r, \mu) = \tilde{G}_r^{(n)}(\{p_i\}, \alpha_r, \mu) \exp \left[d_n s - n \int_0^{-2s} dt \gamma(\alpha_r(t)) \right] \quad (6.100)$$

This clarifies the nomenclature “anomalous dimension”.

^aThis is the *method of characteristics* to solve quasi-linear first-order PDEs.

Intuitively, the Callan–Symanzik equation expresses the fact that physical observables cannot depend on the renormalization point μ , which is an arbitrary energy scale introduced in the calculations to eliminate the explicit dependence on the regulator. Indeed, it constraints the dependence of renormalized quantities on μ in such a way that when the physical observables (i.e. S -matrix elements) are expressed in terms of them, if the renormalization point varies, the renormalized quantities vary too but the physical observables do not.

Part IV

Gauge Theories

Non-Abelian Gauge Theories

In §2.3 it was shown how, from the gauge invariance Eq. 2.53 of Electrodynamics, it was necessary to introduce a covariant derivative Eq. 2.67, ultimately leading to the QED Lagrangian. The same process can be generalized to arbitrary groups of gauge transformations (**gauge groups**, for short), thus defining a field (gauge) theory starting from its symmetry properties.

§7.1 Yang-Mills Lagrangian

Consider $n \in \mathbb{N}$ fermionic fields $\{\psi_k(x)\}_{k=1,\dots,n}$ and an n -spinor $\Psi(x)$ defined as:

$$\Psi(x) = \begin{pmatrix} \psi_1(x) \\ \vdots \\ \psi_n(x) \end{pmatrix} \quad (7.1)$$

As a gauge group, consider a d -dimensional Lie group G : WLOG, be G a simply-connected Lie group, so that each element can be expressed via the exponential map, and be it compact too, so that its representations are unitary. Then, consider $\{T^a\}_{a=1,\dots,d}$ an n -dimensional representation of the associated Lie algebra \mathfrak{g} , so that the action of G on Ψ can be expressed as:

$$\Psi(x) \mapsto V(x)\Psi(x) \quad V(x) = \exp[i\theta_a(x)T^a] \equiv e^{i\alpha(x)} \quad (7.2)$$

where the Lie parameters $\{\theta_a(x)\}_{a=1,\dots,d} \subset \mathcal{C}^\infty(\mathbb{R}^{1,3})$ define a local gauge transformation. The aim is to define a Lagrangian which is invariant under this transformation, i.e. the Lagrangian of a (local) gauge theory.

Simple terms invariant under global phase rotations, like the fermion mass term $m\bar{\Psi}\Psi$, are of course invariant under Eq. 7.2 too, but derivatives need a careful treatment: indeed, the limit-definition of a derivative involves fields at different spacetime points, which have different transformations according to Eq. 7.2. In order to define a derivative of Ψ , it is necessary to introduce a factor to subtract values of $\Psi(x)$ in a meaningful way, so consider $U(y, x) \in U(n)$: $U(x, x) = 1$ and which transforms under the action of G as:

$$U(y, x) \mapsto e^{i\alpha(y)}U(y, x)e^{-i\alpha(x)} \quad (7.3)$$

In general, $U(y, x) \mapsto V(y)U(y, x)V^\dagger(x)$. It is clear that $U(y, x)\Psi(x)$ and $\Psi(y)$ have the same transformation law, so they can be meaningfully subtracted.

Definition 7.1.1 (Covariant derivative)

Given $n^\mu \in \mathbb{R}^{1,3}$, the **covariant derivative** of a fermionic field $\Psi(x)$ along n^μ is defined as:

$$n^\mu \mathcal{D}_\mu \Psi(x) := \lim_{\varepsilon \rightarrow 0} \frac{1}{\varepsilon} [\Psi(x + \varepsilon n) - U(x + \varepsilon n, x) \Psi(x)] \quad (7.4)$$

where $U(y, x)$ is defined in Eq. 7.3.

To make this definition explicit, it is necessary to get an expression of $U(y, x)$ for points at infinitesimal distance. Given the unitarity of $U(y, x)$, it can be expressed through the generators $\{T^a\}_{a=1, \dots, d}$ as:

$$U(x + \varepsilon n, x) = I_n + i g \varepsilon n^\mu \mathcal{A}_\mu^a(x) T_a + o(\varepsilon^2) \quad (7.5)$$

where $g \in \mathbb{R}$ is a constant. The new vector field $\mathcal{A}_\mu^a(x)$ (actually, d different vector fields) is a **connection**, and it allows to express the covariant derivative as (directly from Eq. 7.4):

$$\mathcal{D}_\mu = \partial_\mu - i g \mathcal{A}_\mu^a T_a \quad (7.6)$$

Proposition 7.1.1

The covariant derivative $\mathcal{D}_\mu \Psi$ transforms as Ψ .

Proof. From Eq. 7.3-7.5 (recalling that ∂_μ is anti-Hermitian):

$$\begin{aligned} I_n + i g \varepsilon n^\mu \mathcal{A}_\mu^a T_a &\mapsto V(x + \varepsilon n) (I_n + i g \varepsilon n^\mu \mathcal{A}_\mu^a T_a) V^\dagger(x) \\ &= [(1 + \varepsilon n^\mu \partial_\mu) V(x)] V^\dagger(x) + V(x) (i g \varepsilon n^\mu \mathcal{A}_\mu^a T_a) V^\dagger(x) + o(\varepsilon^2) \\ &= I_n - \varepsilon n^\mu V(x) \partial_\mu V^\dagger(x) + V(x) (i g \varepsilon n^\mu \mathcal{A}_\mu^a T_a) V^\dagger(x) + o(\varepsilon^2) \end{aligned}$$

Hence, the connection transforms as:

$$\mathcal{A}_\mu^a(x) T_a \mapsto V(x) \left[\mathcal{A}_\mu^a(x) T_a + \frac{i}{g} \partial_\mu \right] V^\dagger(x)$$

The derivative $\partial_\mu V^\dagger(x)$ is non-trivial to compute, as G is in general non-Abelian, hence the exponent does not necessarily commute with its derivative. At $o(\theta)$:

$$\begin{aligned} \mathcal{A}_\mu^a(x) T_a &\mapsto (I_n + i \theta^b(x) T_b + o(\theta^2)) \left[\mathcal{A}_\mu^a(x) T_a + \frac{i}{g} \partial_\mu \right] (I_n - i \theta^c(x) T_c + o(\alpha^2)) \\ &= (I_n + i \theta^b(x) T_b + o(\theta^2)) \left[\mathcal{A}_\mu^a(x) T_a - i \mathcal{A}_\mu^a(x) \theta^c(x) T_a T_c + \frac{1}{g} \partial_\mu \theta^c(x) T_c + o(\theta^2) \right] \\ &= \mathcal{A}_\mu^a(x) T_a - i \mathcal{A}_\mu^a(x) \theta^c(x) T_a T_c + i \theta^b(x) \mathcal{A}_\mu^a(x) T_b T_a + \frac{1}{g} \partial_\mu \theta^c(x) T_c + o(\theta^2) \\ &= \mathcal{A}_\mu^a(x) T_a + f^{abc} \mathcal{A}_\mu^a(x) \theta^b(x) T_c + \frac{1}{g} \partial_\mu \theta^a(x) T_a + o(\theta^2) \end{aligned}$$

Then:

$$\begin{aligned} \mathcal{D}_\mu \Psi(x) &\mapsto [\partial_\mu - i g \mathcal{A}_\mu^a(x) T_a - i g f^{abc} \mathcal{A}_\mu^a(x) \theta^b(x) T_c - i \partial_\mu \theta^a(x) T_a] (I_n + i \theta^a(x) T_a) \Psi(x) + o(\theta^2) \\ &= [\partial_\mu + i \theta^a T_a \partial_\mu + i \partial_\mu \theta^a T_a - i g \mathcal{A}_\mu^a T_a + g \mathcal{A}_\mu^a \theta^b T_a T_b - i g f^{abc} \mathcal{A}_\mu^a \theta^b T_c - i \partial_\mu \theta^a T_a + o(\theta^2)] \Psi \\ &= [\partial_\mu + i \theta^a T_a \partial_\mu - i g \mathcal{A}_\mu^a T_a + g \mathcal{A}_\mu^a \theta^b T_a T_b - i g f^{abc} \mathcal{A}_\mu^a \theta^b T_c + o(\theta^2)] \Psi \end{aligned}$$

Recognizing $T_a T_b - i f^{abc} T_c = T_b T_a$ allows to write:

$$\begin{aligned} \mathcal{D}_\mu \Psi(x) &\mapsto [\partial_\mu + i\theta^a(x) T_a \partial_\mu - ig \mathcal{A}_\mu^a(x) T_a + g\theta^b(x) T_b \mathcal{A}_\mu^a(x) T_a + o(\theta^2)] \Psi(x) \\ &= [I_n + i\theta^a(x) T_a + o(\theta^2)] (\partial_\mu - ig \mathcal{A}_\mu^a(x) T_a) \Psi(x) = V(x) \mathcal{D}_\mu \Psi(x) \end{aligned}$$

which is the thesis. \square

The gauge-invariant Lagrangian can thus be built using covariant derivatives (minimal coupling prescription), but there needs to be included a kinetic term for the connection, i.e. a gauge-invariant term depending on $\mathcal{A}_\mu^a(x)$ only.

Lemma 7.1.1 (Field-strength tensor)

The commutator of covariant derivatives reads:

$$[\mathcal{D}_\mu, \mathcal{D}_\nu] = -ig \mathcal{F}_{\mu\nu}^a T_a \quad (7.7)$$

with the **field-strength tensor** defined as:

$$\mathcal{F}_{\mu\nu}^a := \partial_\mu \mathcal{A}_\nu^a - \partial_\nu \mathcal{A}_\mu^a + g f^{abc} \mathcal{A}_\mu^b \mathcal{A}_\nu^c \quad (7.8)$$

Proof. By direct computation:

$$\begin{aligned} [\mathcal{D}_\mu, \mathcal{D}_\nu] &= [\partial_\mu, \partial_\nu] - ig [\mathcal{A}_\mu^a, \partial_\nu] T_a - ig [\partial_\mu, \mathcal{A}_\nu^a] T_a - g^2 \mathcal{A}_\mu^b \mathcal{A}_\nu^c [T_b, T_c] \\ &= -ig (\mathcal{A}_\mu^a \partial_\nu - \partial_\nu \mathcal{A}_\mu^a - \mathcal{A}_\mu^a \partial_\nu) T_a - ig (\mathcal{A}_\nu^a \partial_\mu - \partial_\mu \mathcal{A}_\nu^a - \mathcal{A}_\nu^a \partial_\mu) T_a - ig^2 f^{bca} \mathcal{A}_\mu^b \mathcal{A}_\nu^c T_c \\ &= -ig (\partial_\mu \mathcal{A}_\nu^a - \partial_\nu \mathcal{A}_\mu^a + g f^{bca} \mathcal{A}_\mu^b \mathcal{A}_\nu^c) T_a \end{aligned}$$

Using $f^{bca} = f^{abc}$, as it is always possible to choose generators such that the structure constants are completely antisymmetric¹, yields the thesis. \square

Note that the field-strength tensor is not itself a gauge-invariant quantity, as really there are d different field-strength tensors. However, it is straightforward to construct gauge-invariant combinations of $\mathcal{F}_{\mu\nu}^a$.

Theorem 7.1.1 (Gauge invariance)

Any globally-symmetric function of Ψ , $\mathcal{F}_{\mu\nu}^a$ and their covariant derivatives is also locally-symmetric, i.e. gauge invariant.

Lemma 7.1.2 (Kinetic term)

The following term is gauge-invariant:

$$\text{tr}\{(\mathcal{F}_{\mu\nu}^a T_a)^2\} = 2 \mathcal{F}_{\mu\nu}^a \mathcal{F}_a^{\mu\nu} \quad (7.9)$$

This allows defining the simplest non-Abelian gauge theory, **Yang-Mills theory**:

$$\mathcal{L}_{\text{YM}} = -\frac{1}{4} \mathcal{F}_{\mu\nu}^a \mathcal{F}_a^{\mu\nu} \quad (7.10)$$

¹Requires proof.

Appendices

Appendix A

Mathematical Reference

§A.1 Multilinear algebra

Definition A.1.1 (Tensor product of vectors)

Given two \mathbb{K} -vector spaces V, W with bases $\{\mathbf{e}_i\}_{i=1,\dots,n}, \{\mathbf{f}_j\}_{j=1,\dots,m}$, the **tensor product** $V \otimes W$ is defined as:

$$\mathbf{v} \otimes \mathbf{w} = \sum_{i=1}^n \sum_{j=1}^m v^i w^j \mathbf{e}_i \otimes \mathbf{f}_j \quad \Longleftrightarrow \quad \begin{pmatrix} v_1 \\ \vdots \\ v_n \end{pmatrix} \otimes \begin{pmatrix} w_1 \\ \vdots \\ w_m \end{pmatrix} = \begin{pmatrix} v_1 w_1 \\ \vdots \\ v_1 w_m \\ \vdots \\ v_n w_1 \\ \vdots \\ v_n w_m \end{pmatrix} \quad (\text{A.1})$$

Definition A.1.2 (Tensor product of applications)

Given \mathbb{K} -linear maps $f : V \rightarrow V'$ and $g : W \rightarrow W'$ represented by $F \equiv [f_{i,j}] \in \mathbb{K}^{n' \times n}$ and $G \equiv [g_{i,j}] \in \mathbb{K}^{m' \times m}$, the **tensor product** $f \otimes g : V \otimes W \rightarrow V' \otimes W'$ is defined as $F \otimes G \in \mathbb{K}^{(n'm') \times (nm)}$:

$$[F \otimes G] = \begin{bmatrix} f_{1,1}g_{1,1} & \cdots & f_{1,1}g_{1,m} & & f_{1,n}g_{1,1} & \cdots & f_{1,n}g_{1,m} \\ \vdots & \ddots & \vdots & \cdots & \vdots & \ddots & \vdots \\ f_{1,1}g_{m',1} & \cdots & f_{1,1}g_{m',m} & & f_{1,n}g_{m',1} & \cdots & f_{1,n}g_{m',m} \\ & & \vdots & \ddots & & & \vdots \\ f_{n',1}g_{1,1} & \cdots & f_{n',1}g_{1,m} & & f_{n',n}g_{1,1} & \cdots & f_{n',n}g_{1,m} \\ \vdots & \ddots & \vdots & \cdots & \vdots & \ddots & \vdots \\ f_{n',1}g_{m',1} & \cdots & f_{n',1}g_{m',m} & & f_{n',n}g_{m',1} & \cdots & f_{n',n}g_{m',m} \end{bmatrix}$$

It is clear that $(F \otimes G)(\mathbf{v} \otimes \mathbf{w}) = (F\mathbf{v}) \otimes (G\mathbf{w})$.

§A.2 Lie groups

Definition A.2.1 (Lie groups)

A **Lie group** is a group whose elements depend in a continuous and differentiable way on a set of real parameters $\{\theta_a\}_{a=1,\dots,d} \subset \mathbb{R}^d$.

A Lie group can be seen both as a group and as a d -dimensional differentiable manifold (with coordinates θ_a). WLOG it is always possible to choose $g(0, \dots, 0) = e$.

Definition A.2.2 (Representations)

Given a group G and a vector space $V(\mathbb{K})$, a **representation** of G on V is a homomorphism $\rho : G \rightarrow \text{GL}(V)$.

A representation ρ which is a isomorphism is called **faithful**. As $\text{GL}(V) \cong \mathbb{K}^{n \times n}$, with $n \equiv \dim_{\mathbb{K}} V$, it is usual to represent G as matrices acting on elements of V , i.e. $\rho : G \rightarrow \mathbb{K}^{n \times n}$.

Theorem A.2.1

Given a Lie group G and $g \in G$ connected with the identity, a representation of degree n on $V(\mathbb{C})$ is:

$$\rho(g(\theta)) = e^{i\theta_a T^a} \quad (\text{A.2})$$

where $\{T^a\}_{a=1,\dots,d} \subset \mathbb{C}^{n \times n}$ are the **generators** of G on V .

Definition A.2.3 (Lie algebras)

Given a Lie group G with generators $\{T^a\}_{a=1,\dots,d} \subset \mathbb{C}^{n \times n}$ on $V(\mathbb{C})$, its **Lie algebra** is:

$$[T^a, T^b] = i f^{ab}_c T^c \quad (\text{A.3})$$

where f^{ab}_c are called the **structure constants**.

The sum over repeated indices is understood.

Proposition A.2.1

The Lie algebra of a Lie group is independent of the representation.

Proposition A.2.2

Any d -dimensional abelian Lie algebra is isomorphic to the direct sum of d one-dimensional Lie algebras.

As a consequence, all irreducible representations of an abelian Lie group are of degree $n = 1$.

Definition A.2.4 (Casimir operators)

Given a Lie group with generators $\{T^a\}_{a=1,\dots,d} \subset \mathbb{C}^{n \times n}$ on $V(\mathbb{C})$, a **Casimir operator** is an operator which commutes with each generator.

Given an irreducible representation on V , Casimir operators are operators proportional to $\mathbb{1}_V$, and the proportionality constants can be used to label the representation: they correspond to conserved physical quantities.

Proposition A.2.3

A non-compact group cannot have finite unitary representations, except for those with trivial non-compact generators.

This means that the non-compact component of a group cannot be represented with unitary operators of finite dimension.

§A.2.1 Adjoint representation

To define the adjoint representation of a Lie group, it is necessary to give more formal definitions.

Definition A.2.5 (Lie algebras)

An n -dimensional \mathbb{K} -**Lie algebra** \mathfrak{g} is an n -dimensional \mathbb{K} -vector space equipped with a bilinear map $[\cdot, \cdot] : \mathfrak{g} \times \mathfrak{g} \rightarrow \mathfrak{g}$ such that:

1. $[X, Y] = -[Y, X] \quad \forall X, Y \in \mathfrak{g}$;
2. $[X, [Y, Z]] + [Y, [Z, X]] + [Z, [X, Y]] = 0 \quad \forall X, Y, Z \in \mathfrak{g}$ (Jacobi identity).

A Lie algebra \mathfrak{g} is **commutative** if $[X, Y] = 0 \quad \forall X, Y \in \mathfrak{g}$. Note the link with Lie groups (to be formalized later): given a Lie group G , denoting as \mathcal{G} the associated differentiable manifold, then the Lie algebra \mathfrak{g} associated to G is $\mathfrak{g} := T_e \mathcal{G}$ (tangent space at the identity).

Example A.2.1 (\mathbb{R}^3 as a Lie algebra)

Let $\mathfrak{g} = \mathbb{R}^3$ and $[\cdot, \cdot] : \mathbb{R}^3 \times \mathbb{R}^3 \ni (\mathbf{x}, \mathbf{y}) \mapsto [\mathbf{x}, \mathbf{y}] \equiv \mathbf{x} \times \mathbf{y} \in \mathbb{R}^3$. Then \mathfrak{g} is a Lie algebra.

Definition A.2.6 (Lie algebra morphisms)

Let $\mathfrak{g}, \mathfrak{h}$ be Lie algebras. A linear map $\varphi : \mathfrak{g} \rightarrow \mathfrak{h}$ is a **Lie algebra homomorphism** if:

$$\varphi([X, Y]) = [\varphi(X), \varphi(Y)] \quad \forall X, Y \in \mathfrak{g}$$

If φ is bijective, then it is a **Lie algebra isomorphism**.

A Lie algebra isomorphism $\varphi : \mathfrak{g} \rightarrow \mathfrak{g}$ is a **Lie algebra automorphism** $\varphi \in \text{Aut } \mathfrak{g}$, where $\text{Aut } \mathfrak{g}$ is the **automorphism group** of \mathfrak{g} (a group under the composition of morphisms).

Definition A.2.7 (Adjoint map (pt. 1))

Let \mathfrak{g} be a Lie group and, given $X \in \mathfrak{g}$, define $\text{ad}_X : \mathfrak{g} \rightarrow \mathfrak{g} : Y \mapsto \text{ad}_X(Y) := [X, Y]$. Then the **adjoint map** on \mathfrak{g} is the map $\text{ad} : \mathfrak{g} \rightarrow \text{End } \mathfrak{g} : X \mapsto \text{ad}_X$.

By Jacobi identity, the adjoint map is a **derivation** of the Lie bracket, as:

$$\text{ad}_X([Y, Z]) = [\text{ad}_X(Y), Z] + [Y, \text{ad}_X(Z)]$$

Proposition A.2.4

Given a Lie algebra \mathfrak{g} , the adjoint map $\text{ad} : \mathfrak{g} \rightarrow \text{End } \mathfrak{g}$ is a Lie algebra homomorphism.

Proof. Note that, by Jacobi identity:

$$\text{ad}_{[X,Y]}(Z) = [[X,Y], Z] = [X, [Y, Z]] + [Y, [Z, X]]$$

Moreover:

$$[\text{ad}_X, \text{ad}_Y](Z) = [X, [Y, Z]] - [Y, [X, Z]] = [X, [Y, Z]] + [Y, [Z, X]]$$

Therefore, $\text{ad}_{[X,Y]} = [\text{ad}_X, \text{ad}_Y]$, which is the thesis. \square

Let \mathfrak{g} be an n -dimensional \mathbb{K} -Lie algebra and $\{X_i\}_{i=1,\dots,n} \subset \mathfrak{g}$ a basis of \mathfrak{g} . Then there are unique constants $c_{ijk} \in \mathbb{K}$:

$$[X_i, X_j] = \sum_{k=1}^n c_{ijk} X_k$$

known as **structure constants**.

Theorem A.2.2 (Lie algebras from Lie groups)

Let $G \subset \text{GL}(n, \mathbb{C})$ be a Lie group. Then $\mathfrak{g} = \{X \in \mathbb{C}^{n \times n} : e^{tX} \in G \ \forall t \in \mathbb{R}\}$ is a Lie algebra.

Even if G is a complex Lie group, its Lie algebra can still be real.

Theorem A.2.3 (Induced Lie algebra homomorphism)

Let G, H be Lie groups, with Lie algebra $\mathfrak{g}, \mathfrak{h}$, and let $\varphi : G \rightarrow H$ be a Lie group homomorphism. Then there exists a unique \mathbb{R} -linear map $\Phi : \mathfrak{g} \rightarrow \mathfrak{h}$ such that:

$$\varphi(e^X) = e^{\Phi(X)} \ \forall X \in \mathfrak{g}$$

The map Φ has additional properties:

1. $\Phi(gXg^{-1}) = \varphi(g)\Phi(X)\varphi(g)^{-1} \ \forall X \in \mathfrak{g}, \forall g \in G$;
2. $\Phi([X, Y]) = [\Phi(X), \Phi(Y)] \ \forall X, Y \in \mathfrak{g}$ (Lie algebra homomorphism);
3. $\Phi(X) = \left. \frac{d}{dt} \right|_{t=0} \varphi(e^{tX}) \ \forall X \in \mathfrak{g}$.

To phrase Th. A.2.3 in the language of manifolds, Φ is the **derivative** of φ at the identity: $\Phi = d\varphi_e$.

Definition A.2.8 (Adjoint map (pt. 2))

Let G be a Lie group, with Lie algebra \mathfrak{g} . The **adjoint map** of $g \in G$ is the linear map $\text{Ad}_g : \mathfrak{g} \rightarrow \mathfrak{g} : \text{Ad}_g(X) := gXg^{-1}$.

As Ad_g is clearly invertible, with $\text{Ad}_g^{-1} = \text{Ad}_{g^{-1}}$, then $\text{Ad}_g \in \text{GL}(\mathfrak{g}) \ \forall g \in G$. Furthermore, it is clear that $\text{Ad}_g([X, Y]) = [\text{Ad}_g(X), \text{Ad}_g(Y)] \ \forall X, Y \in \mathfrak{g}, \forall g \in G$, therefore each adjoint map is a Lie algebra homomorphism.

Proposition A.2.5 (Adjoint representation)

Let G be a Lie group, with Lie algebra \mathfrak{g} . Then the map $\text{Ad} : G \rightarrow \text{GL}(\mathfrak{g}) : g \mapsto \text{Ad}_g$ is a homomorphism.

Recalling Def. A.2.2, $\text{Ad} : G \rightarrow \text{GL}(\mathfrak{g})$ is a representation of G on \mathfrak{g} , called the **adjoint representation**.

As $\text{GL}(\mathfrak{g}) \cong \text{GL}(n, \mathbb{K})$ (with $n \equiv \dim_{\mathbb{K}} \mathfrak{g}$ and $\mathbb{K} = \mathbb{R}$ or \mathbb{C}), it can be viewed as a Lie group itself, and its Lie algebra is $\mathfrak{gl}(\mathfrak{g})$. Thus, $\text{Ad} : G \rightarrow \text{GL}(\mathfrak{g})$ is a Lie group homomorphism (as it can be shown to be continuous).

Proposition A.2.6 (Adjoint maps)

Let G be a Lie group, with Lie algebra \mathfrak{g} . Then, given the Lie group homomorphism $\text{Ad} : G \rightarrow \text{GL}(\mathfrak{g})$, the induced Lie algebra homomorphism is $\text{ad} : \mathfrak{g} \rightarrow \mathfrak{gl}(\mathfrak{g})$ such that $\text{ad}_X(Y) = [X, Y]$.

Proof. By Th. A.2.3, the Lie algebra homomorphism induced by $\varphi \equiv \text{Ad}$ is:

$$\Phi(X) \equiv \text{ad}_X = \left. \frac{d}{dt} \right|_{t=0} \text{Ad}_{e^{tX}}$$

Hence:

$$\text{ad}_X(Y) = \left. \frac{d}{dt} \right|_{t=0} e^{tX} Y e^{-tX} = [X, Y]$$

□

This result links the two adjoint maps in Def. A.2.7-A.2.8.

§A.2.2 $\text{SU}(n)$ Lie group

The $\text{SU}(n)$ group is the group of unitary transformations of n -dimensional complex vectors. Its (faithful) fundamental representation thus is:

$$\text{SU}(n) = \{U \in \mathbb{C}^{n \times n} : UU^\dagger = U^\dagger U = I_n \wedge \det U = +1\}$$

The generators of $\text{SU}(n)$ can be found setting $U = \exp(i\theta_a T^a) = I_n + i\theta_a T^a$ and using $U^\dagger U = I_n$:

$$T^a = T^{a\dagger} \tag{A.4}$$

Moreover, by the Jacobi formula $(\det A(t)) \frac{d}{dt} (\det A(t)) = \text{tr}(A(t)^{-1} \frac{d}{dt} A(t))$ evaluated at $t = 0$:

$$\text{tr} T^a = 0 \tag{A.5}$$

The traceless condition can be generalized to all semi-simple Lie algebras. Therefore, the generators of $\text{SU}(n)$ are $\mathbb{C}^{n \times n}$ hermitian traceless matrices: the dimension of $\mathfrak{su}(n)$ then is $n^2 - 1$.

The adjoint representation can be given by representing the generators of the Lie group (i.e. the basis of the Lie algebra) with the structure constants of the Lie algebra:

$$(T_{\text{ad}}^b)_{ac} = if_{ab}^c \tag{A.6}$$

Proposition A.2.7 (Structure constants)

The structure constants of a Lie algebra satisfy the Lie algebra.

Proof. As $[T^a, T^b] = if^ab_c T^c$, the Jacoby identity becomes (recalling that f^ab_c is totally antisymmetric):

$$\begin{aligned} [[T^a, T^b], T^c] + [[T^b, T^c], T^a] + [[T^c, T^a], T^b] &= 0 \\ \iff f^{ab}_d f^{dc}_e + f^{bc}_d f^{da}_e + f^{ca}_d f^{db}_e &= 0 \end{aligned}$$

The condition $([T^a_{\text{ad}}, T^c_{\text{ad}}])_{be} = if^{ac}_d (T^d_{\text{ad}})_{be}$ then gives:

$$f^{ba}_d f^{dc}_e - f^{bc}_d f^{da}_e = f^{ac}_d f^{bd}_e \iff f^{ab}_d f^{dc}_e + f^{bc}_d f^{da}_e + f^{ca}_d f^{db}_e = 0$$

These two expressions are equal, hence the thesis. \square

Moreover, since the structure constant are real, the adjoint representation is always a real representation: the adjoint representation of $\text{SU}(n)$ has degree $n^2 - 1$.

Representation are labelled by their Casimir operators. For any simple Lie algebra, given a representation \mathbf{r} , a Casimir operator is defined as:

$$T_{\mathbf{r}}^a T_{\mathbf{r}}^a = C_2(\mathbf{r}) \mathbf{I}_{n_{\mathbf{r}}} \quad (\text{A.7})$$

This is called the **quadratic Casimir operator**, as it is associated to $T^2 \equiv T^a T^a$ (a Casimir operator since $[T^b, T^2] = if^{ba}_c \{T^c, T^a\} = 0$ by antisymmetry).

Proposition A.2.8 (Quadratic Casimir operator)

For the fundamental and the adjoint representations \mathbf{n} and \mathbf{g} of $\text{SU}(n)$, the quadratic Casimir operator is:

$$C_2(\mathbf{n}) = T_{\text{R}} \frac{n^2 - 1}{n} \quad C_2(\mathbf{g}) = 2T_{\text{R}} n \quad (\text{A.8})$$

where T_{R} is the normalization factor defined as:

$$\text{tr}(T_{\mathbf{n}}^a T_{\mathbf{n}}^b) = T_{\text{R}} \delta^{ab} \quad (\text{A.9})$$

Proof. It is always possible to chose generators such that Eq. A.9 holds^a, hence, contracting it with δ^{ab} (with $a, b = 1, \dots, n^2 - 1$, as they label the basis of $\mathfrak{su}(n)$) and recalling Eq. A.7:

$$C_2(\mathbf{n})n = T_{\text{R}}(n^2 - 1)$$

To compute the Casimir operator for the adjoint representation, first consider the decomposition of the direct product of two representations:

$$\mathbf{r}_1 \otimes \mathbf{r}_2 = \bigoplus_i \mathbf{r}_i$$

In this representation $T_{\mathbf{r}_1 \otimes \mathbf{r}_2}^a = T_{\mathbf{r}_1}^a \otimes \mathbf{1}_{\mathbf{r}_2} + \mathbf{1}_{\mathbf{r}_1} \otimes T_{\mathbf{r}_2}^a$, and it acts on tensor objects Ξ_{pq} whose first index transforms according to \mathbf{r}_1 and the second index according to \mathbf{r}_2 . Recalling that $\text{tr} T^a = 0$:

$$\begin{aligned} \text{tr}(T_{\mathbf{r}_1 \otimes \mathbf{r}_2}^a)^2 &= \text{tr}((T_{\mathbf{r}_1}^a)^2 \otimes \mathbf{1}_{\mathbf{r}_2} + 2T_{\mathbf{r}_1}^a \otimes T_{\mathbf{r}_2}^a + \mathbf{1}_{\mathbf{r}_1} \otimes (T_{\mathbf{r}_2}^a)^2) \\ &= \text{tr}(C_2(\mathbf{r}_1) \mathbf{1}_{\mathbf{r}_1} \otimes \mathbf{1}_{\mathbf{r}_2}) + \text{tr}(C_2(\mathbf{r}_2) \mathbf{1}_{\mathbf{r}_1} \otimes \mathbf{1}_{\mathbf{r}_2}) = (C_2(\mathbf{r}_1) + C_2(\mathbf{r}_2))n_{\mathbf{r}_1}n_{\mathbf{r}_2} \end{aligned}$$

However, by the decomposition above:

$$\text{tr}(T_{\mathbf{r}_1 \otimes \mathbf{r}_2}^a)^2 = \sum_i C_2(\mathbf{r}_i) n_{\mathbf{r}_i}$$

Consider $\mathbf{n} \otimes \mathbf{n}^*$, where \mathbf{n}^* is the complex conjugate of the fundamental representation (for complex representations, \mathbf{r} and \mathbf{r}^* are generally inequivalent representations): then Ξ_{pq} contains a term proportional to the invariant δ_{pq} , while the other $n^2 - 1$ independent components transform as a general $n \times n$ traceless tensor, i.e. under the adjoint representation of $\text{SU}(n)$ (as of Eq. A.4-A.5), thus $\mathbf{n} \otimes \mathbf{n}^* = \mathbf{1} \oplus \mathbf{g}$ and the above identity becomes:

$$(C_2(\mathbf{1}) + C_2(\mathbf{g}))(n^2 - 1) = (C_2(\mathbf{n}) + C_2(\mathbf{n}^*))n^2$$

Using $C_2(\mathbf{1}) = 0$ (as all generators are trivially zero) and $C_2(\mathbf{n}^*) = C_2(\mathbf{n})$:

$$C_2(\mathbf{g})(n^2 - 1) = 2T_R \frac{n^2 - 1}{n} n^2$$

which completes the proof. □

^aRequires proof.

§A.2.2.1 $\text{SU}(2)$ Lie group

The fundamental representation of $\text{SU}(2)$ is $T_2^a = \frac{\sigma^a}{2}$, while $\mathfrak{su}(2)$ is defined by commutators $[T_2^a, T_2^b] = i\epsilon^{abc}T_2^c$ (as $\sigma^a\sigma^b = \delta^{ab}\mathbf{I}_2 + i\epsilon^{abc}\sigma^c$). The adjoint representation then is:

$$(T_{\mathbf{g}}^a)_{ij} = i\epsilon^{iaj} \tag{A.10}$$

Explicitly:

$$T_{\mathbf{g}}^1 = \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{bmatrix} \quad T_{\mathbf{g}}^2 = \begin{bmatrix} 0 & 0 & i \\ 0 & 0 & 0 \\ -i & 0 & 0 \end{bmatrix} \quad T_{\mathbf{g}}^3 = \begin{bmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{bmatrix}$$

As an aside, these are exactly the generators of the fundamental representation of $\text{SO}(3)$: this is due to the adjoint map of $\text{SU}(2)$ being also the double-covering map on $\text{SO}(3) \cong \text{SU}(2)/\mathbb{Z}_2$.

§A.3 Algebras

Definition A.3.1 (Associative algebras)

An n -dimensional **associative \mathbb{K} -algebra** \mathcal{A} is an n -dimensional vector space $V(\mathbb{K})$ equipped with a bilinear map $\mathcal{A} \times \mathcal{A} \ni (a, b) \mapsto ab \in \mathcal{A}$ such that:

1. $(ab)c = a(bc) \quad \forall a, b, c \in \mathcal{A}$ (associativity);
2. $a(b + c) = ab + ac \wedge (a + b)c = ac + bc \quad \forall a, b, c \in \mathcal{A}$;
3. $\lambda(ab) = (\lambda a)b = a(\lambda b) \quad \forall \lambda \in \mathbb{K}, \forall a, b \in \mathcal{A}$.

An algebra is said to be **unital** if $\exists 1 \in \mathcal{A} : 1a = a1 = a \quad \forall a \in \mathcal{A}$, called **identity element**.

Definition A.3.2 (Algebra morphisms)

Given two associative \mathbb{K} -algebras \mathcal{A}, \mathcal{B} , a \mathbb{K} -linear map $\varphi : \mathcal{A} \rightarrow \mathcal{B}$ is an **algebra morphism** if:

$$\varphi(ab) = \varphi(a)\varphi(b) \quad \forall a, b \in \mathcal{A}$$

If $\varphi(1_{\mathcal{A}}) = 1_{\mathcal{B}}$, then φ is a **unital morphism**. An algebra morphism $\varphi : \mathcal{A} \rightarrow \mathcal{A}$ is an **endomorphism**, and if $\varphi^2 = \mathbb{1}_{\mathcal{A}}$ it is an **involution**.

Proposition A.3.1 (Even subalgebras)

Given a unital associative \mathbb{K} -algebra \mathcal{A} and an involution $\varphi \in \text{End } \mathcal{A}$, then:

$$\mathcal{A} = \mathcal{A}^+ \oplus \mathcal{A}^-$$

where, defining $\pi := \frac{1}{2}(\mathbb{1}_{\mathcal{A}} + \varphi)$:

$$\mathcal{A}^+ := \pi(\mathcal{A}) = \{a \in \mathcal{A} : \varphi(a) = a\}$$

$$\mathcal{A}^- := (\mathbb{1}_{\mathcal{A}} - \pi)(\mathcal{A}) = \{a \in \mathcal{A} : \varphi(a) = -a\}$$

As $\mathcal{A}^+\mathcal{A}^+, \mathcal{A}^-\mathcal{A}^- \subset \mathcal{A}^+$ and $\mathcal{A}^+\mathcal{A}^-, \mathcal{A}^-\mathcal{A}^+ \subset \mathcal{A}^-$, then \mathcal{A}^+ is a subalgebra of \mathcal{A} , the **even subalgebra**.

Definition A.3.3 (Modules)

Given a unital associative \mathbb{K} -algebra \mathcal{A} , a **right-module** M is a \mathbb{K} -vector space equipped with a \mathbb{K} -bilinear multiplication map $\mathcal{A} \times M \ni (a, m) \mapsto am \in M$ such that:

1. $(ab)m = a(bm) \quad \forall a, b \in \mathcal{A}, \forall m \in M$;
2. $1_{\mathcal{A}}m = m \quad \forall m \in M$.

Definition A.3.4 (Ideals)

Given a unital associative \mathbb{K} -algebra \mathcal{A} , a **right-ideal** is a subalgebra $\mathcal{I} \subset \mathcal{A}$ such that:

1. $aj \in \mathcal{I} \quad \forall a \in \mathcal{A}, \forall j \in \mathcal{I}$.

A right-ideal is said minimal if it is non-trivial and does not contain any non-trivial sub-right-ideal.

The analogous left- definitions are clear. A left-ideal of \mathcal{A} can be viewed as a right-module on \mathcal{A} .

§A.3.1 Clifford algebras

Definition A.3.5 (Clifford algebras)

Given an n -dimensional vector space $V(\mathbb{K})$ with a quadratic form q , associated linear form¹ ω and orthogonal basis $\{e_i\}_{i=1,\dots,n}$, and a unital associative \mathbb{K} -algebra \mathcal{A} , a **Clifford mapping** is an injective \mathbb{K} -linear map $\rho : V \rightarrow \mathcal{A} : 1 \notin \rho(V) \wedge \rho(x)^2 = -q(x)1 \ \forall x \in V$. If $\rho(V)$ generates \mathcal{A} , then (\mathcal{A}, ρ) is a **Clifford algebra** for (V, q) , and is denoted by $\mathfrak{cl}(V)$.

Lemma A.3.1

$$\{\rho(x), \rho(y)\} = -2\omega(x, y)1 \quad \forall x, y \in V$$

Proof. By direct computation:

$$\rho(x)\rho(y) + \rho(y)\rho(x) = \rho(x+y)^2 - \rho(x)^2 - \rho(y)^2 = -(q(x+y) - q(x) - q(y))1 = -2\omega(x, y)1$$

which is the thesis. \square

Setting for ease of reading $\rho(x) \equiv x$, it is clear that $x \perp y \implies xy = -yx$.

More intuitively, the Clifford algebra $\mathfrak{cl}(V)$ can be seen as the associative algebra generated by V setting $xy = -\omega(x, y)1 \ \forall x, y \in V$, so that:

$$\{x, y\} = 2\omega(x, y)1 \quad \forall x, y \in V \tag{A.11}$$

In general, $\mathfrak{cl}_{m,n}(\mathbb{R})$ denotes the Clifford algebra associated to $\mathbb{R}^{m,n}$ with quadratic form:

$$q(x_1, \dots, x_m, y_1, \dots, y_n) = \sum_{i=1}^m x_i^2 - \sum_{j=1}^n y_j^2$$

Example A.3.1 (Complex numbers)

Via Clifford algebras, $\mathbb{C} \cong \mathfrak{cl}_{0,1}(\mathbb{R})$. In fact, $\mathbb{R}^{0,1}$ has orthonormal basis $\{e_1\}$ such that $q(e_1) = -1$, i.e. $e_1^2 = -1$, so elements of the Clifford algebra are generated by $\{1, e_1\}$: identifying $e_1 \equiv i$ gives the desired isomorphism.

Example A.3.2 (Quaternions)

$\mathbb{H} \cong \mathfrak{cl}_{0,2}(\mathbb{R})$. Indeed, $\mathbb{R}^{0,2}$ has orthonormal basis $\{e_1, e_2\} : e_1^2 = e_2^2 = -1$; moreover, as $e_1 \perp e_2$, then $e_1 e_2 = -e_2 e_1$ by Eq. A.11, so elements of $\mathfrak{cl}_{0,2}(\mathbb{R})$ are generated by $\{1, e_1, e_2, e_1 e_2\}$:

¹Given a vector space $V(\mathbb{K})$, a quadratic form is a map $q : V \rightarrow \mathbb{K}$ such that $q(\lambda x) = \lambda^2 q(x) \ \forall \lambda \in \mathbb{K}, \forall x \in V$, and the associated bilinear form is a \mathbb{K} -bilinear map $\omega : V \times V \rightarrow \mathbb{K}$ such that $\omega(x, y) = \frac{1}{2}(q(x+y) - q(x) - q(y))$, which is manifestly symmetric and $q(x) = \omega(x, x)$.

setting $e_1 \equiv i$, $e_2 \equiv j$ and $e_1 e_2 = k$ yields the result.

§A.3.1.1 Spin groups

Given an n -dimensional \mathbb{K} -vector space, the Clifford algebra $\mathfrak{cl}(V)$ is finite dimensional and is naturally \mathbb{N}_0 -graded² as:

$$\mathfrak{cl}(V) = \bigoplus_{i=0}^n \mathfrak{cl}^{(i)}(V) \quad (\text{A.12})$$

where $\mathfrak{cl}^{(0)}(V) = \mathbb{K}$, $\mathfrak{cl}^{(1)}(V) = V$ and $\mathfrak{cl}^{(2)}(V) \equiv \text{Spin}(V)$ is the **spin group** of V . The spin group is a Lie group and, via its natural action on V , can be shown to be $\text{Spin}(V) \cong \mathfrak{so}(V)$.

§A.3.2 Grassmann algebras

For the definition of the tensor product, recall Eq. A.1.

Definition A.3.6 (Tensor algebras)

Given a vector space $V(\mathbb{K})$, for all $k \in \mathbb{N}$ define the k^{th} **tensor power**:

$$T^k V := V^{\otimes k} \equiv \underbrace{V \otimes \cdots \otimes V}_{k \text{ times}}$$

By convention $T^0 V \equiv \mathbb{K}$. The **tensor algebra** (or free algebra) of V is then defined as:

$$T(V) := \bigoplus_{k=0}^{\infty} T^k V \quad (\text{A.13})$$

Given $k \in \mathbb{N}$, $T^k V$ is a \mathbb{K} -vector consisting of all tensors on V of order $(k, 0)$ (hence, $T^0 V \equiv \mathbb{K}$ must be seen as a 1D \mathbb{K} -vector space). Expanding the direct sum, the tensor algebra is:

$$T(V) = \mathbb{K} \oplus V \oplus (V \otimes V) \oplus (V \otimes V \otimes V) \oplus \dots \quad (\text{A.14})$$

and it is a unital associative algebra with multiplication determined by the canonical isomorphism $T^k V \otimes T^\ell V \rightarrow T^{k+\ell} V$: in particular, $T(V)$ is an \mathbb{N}_0 -graded associative algebra.

Definition A.3.7 (Grassmann algebras)

Given a vector space $V(\mathbb{K})$ and the two-sided ideal \mathcal{I} generated by^a $\{x^{\otimes k} : x \in V, k \in \mathbb{N}_{\geq 2}\}$, the **Grassman algebra** (or exterior algebra) of V is defined as the quotient algebra:

$$\bigwedge(V) := T(V)/\mathcal{I} \quad (\text{A.16})$$

The **exterior product** of two elements $\alpha, \beta \in \bigwedge(V)$ is defined as:

$$\alpha \wedge \beta \equiv \alpha \otimes \beta \mod \mathcal{I} \quad (\text{A.17})$$

²Given an index set \mathcal{I} and a \mathbb{K} -vector space, the latter is \mathcal{I} -graded if there exists a family of subspaces $\{V_i\}_{i \in \mathcal{I}}$ of V such that:

$$V = \bigoplus_{i \in \mathcal{I}} V_i$$

^aGiven a unital associative algebra \mathcal{A} and a subset $X \subset \mathcal{A}$, the two-sided ideal \mathcal{I} of \mathcal{A} generated by X is:

$$\mathcal{I} := \left\{ \sum_{i=1}^k a_i x_i b_i : a_i, b_i \in \mathcal{A}, x_i \in X \right\} \quad (\text{A.15})$$

The quotient algebra is found analogously to the quotient group. Note that the exterior product is an **alternating product**, as by definition $\omega \wedge \omega = 0 \forall \omega \in \Lambda(V)$. In general, by definition $x_1 \wedge \cdots \wedge x_k = 0$ if $x_i = x_j \in V$ for some $i \neq j \in [1, k]$.

Proposition A.3.2 (Anticommutativity)

Given a vector space $V(\mathbb{K})$, a k -ple $\{x_i\}_{i=1,\dots,k}$, $k \in \mathbb{N}$, and $\sigma \in S^k$, then:

$$x_{\sigma(1)} \wedge \cdots \wedge x_{\sigma(k)} = \text{sgn } \sigma x_1 \wedge \cdots \wedge x_k \quad (\text{A.18})$$

Proof. For $k = 2$:

$$0 = (x + y) \wedge (x + y) = x \wedge x + x \wedge y + y \wedge x + y \wedge y = x \wedge y + y \wedge x$$

from which $x \wedge y = -y \wedge x$. □

The Grassmann algebra is \mathbb{N}_0 -graded too. Defining the k^{th} **exterior power** $\bigwedge^k(V)$ as the \mathbb{K} -vector subspace of $\bigwedge(V)$ spanned by elements of the form $x_1 \wedge \cdots \wedge x_k$, with $x_i \in V$, then:

$$\bigwedge(V) = \bigoplus_{k=0}^{\infty} \bigwedge^k(V) \quad (\text{A.19})$$

If $\omega \in \bigwedge^k(V)$ can be expressed as the exterior product of k elements of V , then ω is said to be **decomposable**.

Theorem A.3.1 (Dimension of the Grassmann algebra)

Given an n -dimensional vector space $V(\mathbb{K})$, then $\bigwedge^k(V) = \{0\} \forall k > \dim_{\mathbb{K}} V$, i.e.:

$$\bigwedge(V) = \bigoplus_{k=0}^n \bigwedge^k(V) \quad (\text{A.20})$$

Proof. Given a basis $\{e_i\}_{i=1,\dots,n}$ of V , WTS $\{e_{i_1} \wedge \cdots \wedge e_{i_k}\}_{1 \leq i_1 < \cdots < i_k \leq n}$ is a basis of $\bigwedge^k(V)$. Consider $x_1 \wedge \cdots \wedge x_k \in \bigwedge^k(V)$, with $\{x_i\}_{i=1,\dots,k} \subset V$. Then, denoting the component of x_i along e_j as x_i^j , it is clear that:

$$x_1 \wedge \cdots \wedge x_k = \sum_{i_1=1}^n \cdots \sum_{i_k=1}^n x_1^{i_1} \cdots x_k^{i_k} e_{i_1} \wedge \cdots \wedge e_{i_k}$$

Note that every term with $e_{i_\ell} = e_{i_m}$ for some $\ell \neq m \in [1, k]$ vanishes, while the remaining terms can be reordered so that $1 \leq i_1 \leq \cdots \leq i_k \leq n$: this shows that $\{e_{i_1} \wedge \cdots \wedge e_{i_k}\}_{1 \leq i_1 < \cdots < i_k \leq n}$ is a basis of $\bigwedge^k(V)$. It is then clear that if $k > n$, then there will be at least two equal basis vector in every basis element: therefore, by the alternating nature of the

exterior product, $\bigwedge^k(V) = 0$. □

It can be shown that $\dim_{\mathbb{K}} \bigwedge^k(V) = \binom{n}{k}$, so $\dim_{\mathbb{K}} \bigwedge(V) = 2^n$. The graded structure of the Grassmann algebra is ensured by:

$$\bigwedge^k(V) \wedge \bigwedge^p(V) \subset \bigwedge^{k+p}(V) \quad (\text{A.21})$$

This is an anticommutative grading, as:

$$\alpha \in \bigwedge^k(V), \beta \in \bigwedge^p(V) \implies \alpha \wedge \beta = (-1)^{kp} \beta \wedge \alpha \quad (\text{A.22})$$

§A.3.2.1 Grassmann numbers

Given an n -dimensional vector field $V(\mathbb{C})$ (possibly $n = \infty$), the elements of $\bigwedge(V)$ are called **Grassmann numbers** (or supernumbers). The general form of a Grassmann number is:

$$z = c_0 + \sum_{k=1}^n \sum_{1 \leq i_1 < \dots < i_k \leq n} c_{i_1 \dots i_k} \theta_{i_1} \dots \theta_{i_k} \equiv z_B + z_S \quad (\text{A.23})$$

where $c_{i_1 \dots i_k} \in \mathbb{C}$ are completely-antisymmetric tensors of rank $(k, 0)$ and $\theta_i \theta_j \equiv \theta_i \wedge \theta_j$ (basis of the Grassmann algebra). z_B is called the **body** and z_S the **soul** of the supernumber z .

Proposition A.3.3 (Soul)

If $n < \infty$, then the soul of a supernumber is nilpotent:

$$z_S^{n+1} = 0 \quad (\text{A.24})$$

Lemma A.3.2 (Equations)

If $n < \infty$, then:

$$\theta_i z = 0 \quad \forall i = 1, \dots, n \implies z = c \theta_1 \dots \theta_n, \quad c \in \mathbb{C} \quad (\text{A.25})$$

If $n = \infty$, instead:

$$\theta_a z = 0 \quad \forall a \in \mathbb{N} \implies z = 0 \quad (\text{A.26})$$

In analogy to Hermitian conjugation, an **involution** (or conjugation) is defined for supernumbers:

$$(\theta_i \theta_j)^* = \theta_j \theta_i = -\theta_i \theta_j \quad (\text{A.27})$$

known as the deWitt convention. Moreover, note that products of an odd number of Grassmann variables anticommute with each other: these are called **a-numbers**. On the other hand, products of an even number of Grassmann variables commute with each other (and with every Grassmann number): these are called **c-numbers**. This decomposition induces a Z_2 -grading of the algebra, showing that Grassmann algebras are supercommutative algebras³. Note that c-numbers form a subalgebra of $\bigwedge(V)$, while a-numbers do not (they are only a subspace).

³A *superalgebra* is a Z_2 -graded associative algebra $\mathcal{A} = \mathcal{A}_0 \oplus \mathcal{A}_1$ with bilinear multiplication $\mathcal{A} \times \mathcal{A} \rightarrow \mathcal{A}$ such that $\mathcal{A}_i \mathcal{A}_j \subset \mathcal{A}_{i+j}$, where indices read modulo 2. Elements of each \mathcal{A}_i are called homogeneous.

A *supercommutative algebra* is a superalgebra such that any two homogeneous elements x, y satisfy $yx = (-1)^{|x||y|}xy$, with $|x|$ denoting the grade of x , i.e. $x \in \mathcal{A}_i \implies |x| = i \pmod{2}$.

§A.3.2.2 Dual numbers

The case $n = 1$ is of particular interest: the elements of the exterior algebra of a 1-dimensional complex vector space are called **dual numbers**.

As for Grassmann variables $\theta\eta = -\epsilon\theta$, so that θ^2 , every analytic function reduces to a linear one on dual numbers (through its Taylor series): $f(\theta) = A + B\theta$, with $A, B \in \mathbb{C}$, is the most general function on dual numbers. Clearly, derivation is defined as:

$$\frac{d}{d\theta} (A + B\theta) \equiv B \quad (\text{A.28})$$

On the other hand, integration (only considered on the whole domain of θ) is less trivial and can be defined imposing two conditions:

1. linearity: $\int d\theta [af(\theta) + bg(\theta)] = a \int d\theta f(\theta) + b \int d\theta g(\theta) \quad \forall a, b \in \mathbb{C};$
2. shift-invariance: $\int d\theta f(\theta) = \int d\theta f(\theta + \eta).$

The only linear function satisfying the second constraint is the constant function, thus integration is conventionally defined as:

$$\int d\theta (A + B\theta) \equiv B \quad (\text{A.29})$$

Therefore, on Grassmann numbers, integration and derivation are the same thing: indeed, if it were $I = D^{-1}$, as $D^2 = 0$, then $0 = ID^2 = D$, which is absurd.

Integration can be extended to complex integration. Parting from deWitt's convention, define:

$$\theta = \frac{\theta_1 + i\theta_2}{\sqrt{2}} \quad \theta^* = \frac{\theta_1 - i\theta_2}{\sqrt{2}} \quad (\text{A.30})$$

These satisfy (simple check):

$$(\theta\eta)^* = \eta^*\theta^* \quad (\text{A.31})$$

θ and θ^* are distinct Grassmann numbers, therefore $d\theta_1 d\theta_2 = \frac{1}{i} d\theta^* d\theta = -\frac{1}{i} d\theta d\theta^*$. The order ambiguity of integrating over multiple Grassmann numbers is solved setting the convention of performing the innermost integral first:

$$\int d\theta d\eta \eta\theta = +1 \quad (\text{A.32})$$

Lemma A.3.3 (Gaussian integral)

Given $b \in \mathbb{C}$:

$$\int d\theta^* d\theta e^{-\theta^* b \theta} = b \quad (\text{A.33})$$

Proof. $\int d\theta^* d\theta e^{-\theta^* b \theta} = \int d\theta^* d\theta (1 - \theta^* b \theta) = \int d\theta^* d\theta (1 + b \theta \theta^*) = b.$ □

Introducing an extra factor of $\theta\theta^*$ in the integrand yields an extra b^{-1} factor (as in standard Gaussian integrals):

$$\int d\theta^* d\theta \theta\theta^* e^{-\theta^* b \theta} = 1 \quad (\text{A.34})$$

Proposition A.3.4 (Unitary invariance)

Given $U \in U(n)$ and setting $\theta'_i = \sum_{j=1}^n U_{ij} \theta_j$, then:

$$\int \prod_{i=1}^n d\theta_i^* \theta_i f(\theta) = \int \prod_{i=1}^n d\theta_i'^* \theta_i' f(\theta') \quad (\text{A.35})$$

Proof. By the Prop. A.3.2:

$$\begin{aligned} \prod_{i=1}^n \theta'_i &= \sum_{i_1, \dots, i_n=1}^n \frac{1}{n!} \epsilon^{i_1 \dots i_n} \theta'_{i_1} \dots \theta'_{i_n} = \sum_{i_1, \dots, i_n=1}^n \sum_{j_1, \dots, j_n=1}^n \frac{1}{n!} \epsilon^{i_1 \dots i_n} U_{i_1 j_1} \dots U_{i_n j_n} \theta_{j_1} \dots \theta_{j_n} \\ &= \sum_{i_1, \dots, i_n=1}^n \sum_{j_1, \dots, j_n=1}^n \frac{1}{n!} \epsilon^{i_1 \dots i_n} \epsilon^{j_1 \dots j_n} U_{i_1 j_1} \dots U_{i_n j_n} \prod_{k=1}^n \theta_k = \det U \prod_{i=1}^n \theta_i \end{aligned}$$

Equivalently:

$$\prod_{i=1}^n \theta_i'^* = (\det U)^* \prod_{i=1}^n \theta_i^*$$

As $U \in U(n)$, $(\det U) (\det U)^* = 1$, so:

$$\prod_{i=1}^n \theta_i'^* \theta_i' = \prod_{i=1}^n \theta_i^* \theta_i$$

The same holds for the measure. As the only term of $f(\theta)$ which survives the integration has exactly one factor of each θ_i^* and one of each θ_i , the integral remains unchanged. \square

With unitary invariance, it is possible to extend complex integration to the multi-dimensional case.

Lemma A.3.4 (Gaussian integral)

Given $B \in \mathbb{C}^{n \times n}$ Hermitian:

$$\int \prod_{i=1}^n d\theta_i^* d\theta_i e^{-\sum_{j,k=1}^n \theta_j^* B_{jk} \theta_k} = \det B \quad (\text{A.36})$$

Proof. Being B Hermitian, it can be diagonalized by a unitary transformation, so, given its eigenvalues $\{b_i\}_{i=1, \dots, n} \subset \mathbb{C}$:

$$\int \prod_{i=1}^n d\theta_i^* d\theta_i e^{-\sum_{j,k=1}^n \theta_j^* B_{jk} \theta_k} = \int \prod_{i=1}^n d\theta_i'^* d\theta_i' e^{-\sum_{j=1}^n \theta_j'^* b_j \theta_j'} = \prod_{i=1}^n \int d\theta_i^* d\theta_i e^{-\theta_i^* b_i \theta_i} = \prod_{i=1}^n b_i$$

By Binet's theorem, this is precisely $\det B$. \square

With extra factors, the standard Gaussian behavior is recovered:

$$\int \prod_{i=1}^n d\theta_i^* d\theta_i \theta_p \theta_q e^{-\sum_{j,k=1}^n \theta_j^* B_{jk} \theta_k} = (\det B) [B^{-1}]_{pq} \quad (\text{A.37})$$

§A.4 Gaussian integrals

Lemma A.4.1

Given $\alpha, \beta \in \mathbb{C} : \Re \alpha \geq 0$:

$$\int_{\mathbb{R}} dx e^{-\frac{1}{2}\alpha x^2 + \beta x} = \sqrt{\frac{2\pi}{\alpha}} e^{\frac{\beta^2}{2\alpha}} \quad (\text{A.38})$$

Proposition A.4.1 (Generalized Gaussian integral)

Given a non-singular matrix $A \in \mathbb{R}^{n \times n}$ and a vector $\mathbf{J} \in \mathbb{R}^n$:

$$\int_{\mathbb{R}^n} d^n x e^{-\frac{1}{2}\mathbf{x}^\top A \mathbf{x} + \mathbf{J} \cdot \mathbf{x}} = \sqrt{\frac{(2\pi)^n}{\det A}} e^{\frac{1}{2}\mathbf{J}^\top A^{-1} \mathbf{J}} \quad (\text{A.39})$$

Lemma A.4.2

Given a non-singular matrix $A \in \mathbb{R}^{n \times n}$ and a vector $\mathbf{J} \in \mathbb{R}^n$:

$$\int_{\mathbb{R}^n} d^n x e^{-\frac{i}{2}\mathbf{x}^\top A \mathbf{x} + i\mathbf{J} \cdot \mathbf{x}} = \sqrt{\frac{(2\pi i)^n}{\det A}} e^{-\frac{i}{2}\mathbf{J}^\top A^{-1} \mathbf{J}} \quad (\text{A.40})$$

Proposition A.4.2 (Solid angle in \mathbb{R}^n)

$$\Omega_{n-1} \equiv \int_{\mathbb{S}^{n-1}} d\Omega_{n-1} = \frac{2\pi^{\frac{n}{2}}}{\Gamma(\frac{n}{2})} \quad (\text{A.41})$$

Proof. Consider the n -dimensional Gaussian integral:

$$\begin{aligned} \int_{\mathbb{R}^n} d^n x e^{-\sum_{i=1}^n x_i^2} &= \pi^{n/2} \\ &= \int_{\mathbb{S}^{n-1}} d\Omega_{n-1} \int_{\mathbb{R}^+} dr r^{n-1} e^{-r^2} = \Omega_{n-1} \int_{\mathbb{R}^+} \frac{d\rho}{2} \rho^{\frac{n}{2}-1} e^{-\rho} =: \Omega_{n-1} \frac{\Gamma(\frac{n}{2})}{2} \end{aligned}$$

where $\rho \equiv r^2$. □

Appendix B

Scattering Processes

§B.1 Cross-sections and S -matrix

The cross-section generalizes the notion of geometric area of an object. Consider a flux of particles (particles per unit area per unit time):

$$j = \frac{dn}{dS dt}$$

Then, the number of incident particles diffused by a disk of section S in a time Δt is:

$$N = j S \Delta t$$

The (classical) **cross-section** of the disk is then defined as:

$$\sigma := \frac{N}{j \Delta t} \quad (\text{B.1})$$

From a quantum point of view, N is the probability of the transition $|\text{in}\rangle \rightarrow |\text{out}\rangle$ between the initial and final states, thus:

$$d\sigma = \frac{1}{j \Delta t} |\langle \text{out}(f) | \text{in} \rangle|^2 df$$

where f is some quantity which the final state depends on.

§B.1.1 Classical and quantum definition

§B.1.2 Phase-space integration

Proposition B.1.1 (Scattering cross-section)

Given a $2 \rightarrow n$ scattering process with well-defined initial momenta p_A and p_B , then the **differential cross-section** is:

$$d\sigma = \frac{1}{2E_A 2E_B |\mathbf{v}_A - \mathbf{v}_B|} \prod_{k=1}^n \frac{d^3 p_k}{(2\pi)^3 2E_k} |\mathcal{M}(\mathcal{AB} \rightarrow \{f\})|^2 (2\pi)^4 \delta^{(4)}(p_A + p_B - \sum_{i=1}^n p_i) \quad (\text{B.2})$$

where $\mathcal{M}(\mathcal{AB} \rightarrow \{f\})$ is the matrix element of the scattering process and $\mathbf{v}_k \equiv \frac{\mathbf{p}_k}{E_k}$ is the velocity of the k^{th} particle.

This cross-section is differential in $3n - 4$ variables, as the momentum-conserving $\delta^{(4)}$ allows to perform 4 integrations. The integration measure is defined as the **invariant n -body phase space**:

$$\int d\Pi_n \equiv \prod_{k=1}^n \int \frac{d^3 p_k}{(2\pi)^3 2E_k} (2\pi)^4 \delta^{(4)}(P - \sum_{i=1}^n p_i) \quad (\text{B.3})$$

with P the total initial 4-momentum.

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