Quantum Field Theory

June 18, 2019

Contents

1 Classical field theory

1.1 Field theory in continuum

Euler-Lagrange-equation

$$\partial_\mu \left(\frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \right) - \frac{\partial \mathcal{L}}{\partial \phi} = 0 \tag{1.1.1}$$

momentum density

$$\pi(x) = \frac{\partial \mathcal{L}}{\partial \dot{\phi}(\boldsymbol{x})} \tag{1.1.2}$$

Hamiltonian density

$$\mathcal{H}(\phi(\boldsymbol{x}), \pi(\boldsymbol{x})) = \pi(\boldsymbol{x})\dot{\phi}(\boldsymbol{x}) - \mathcal{L}(\phi, \partial_\mu \phi) \tag{1.1.3}$$

1.2 Noether Theorem

If a Lagrangian field theory has an infinitisimal symmetry, then there is an associated current j^μ, which is conserved.

$$\partial_\mu j^\mu = 0 \tag{1.2.1}$$

$$j^\mu = \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \Delta\phi - X^m u \tag{1.2.2}$$

Energy-momentum tensor (stress-energy tensor)

Asymmetric version

$$\Theta^\mu_\nu = \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \partial_\nu \phi - \delta^\mu_\nu \mathcal{L} \tag{1.2.3}$$

General version

$$T^{\mu\nu} = \Theta^{\mu\nu} + \partial_\lambda f^{\mu\nu\lambda} \tag{1.2.4}$$

with $f^{\lambda\mu\nu} = -f^{\mu\lambda\nu}$ or $\partial_\mu \partial \nu f^{\lambda\mu\nu} = 0$

2 Klein-Gordon theory

(Real) Lagrangian density

$$\mathcal{L} = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 \tag{2.0.1}$$

Quantization

$$\begin{aligned} [\phi(\boldsymbol{x}), \phi(\boldsymbol{x}')] = [\pi(\boldsymbol{x}), \pi(\boldsymbol{x}')] = 0 \\ [\phi(\boldsymbol{x}), \pi(\boldsymbol{x}')] = i\delta^{(3)}(\boldsymbol{x} - \boldsymbol{x}') \end{aligned} \tag{2.0.2}$$

Decomposition into Fourier modes

$$\phi(\boldsymbol{x}) = \int \frac{\mathrm{d}^3 p}{(2\pi)^3}\frac{1}{\sqrt{2E_p}}\left(a_{\boldsymbol{p}}e^{i\boldsymbol{p}\cdot\boldsymbol{x}} + a^\dagger_{\boldsymbol{p}}e^{-i\boldsymbol{p}\cdot\boldsymbol{x}}\right) \tag{2.0.3}$$

$$\pi(\boldsymbol{x}) = \int \frac{\mathrm{d}^3 p}{(2\pi)^3}(-i)\sqrt{\frac{E_p}{2}}\left(a_{\boldsymbol{p}}e^{i\boldsymbol{p}\cdot\boldsymbol{x}} - a^\dagger_{\boldsymbol{p}}e^{-i\boldsymbol{p}\cdot\boldsymbol{x}}\right) \tag{2.0.4}$$

thus the commutation relations for ladder operators:

$$\left[a_{\boldsymbol{p}}, a_{\boldsymbol{p}'}\right] = \left[a^\dagger_{\boldsymbol{p}}, a^\dagger_{\boldsymbol{p}'}\right] = 0 \tag{2.0.5}$$

$$\left[a_{\boldsymbol{p}}, a^\dagger_{\boldsymbol{p}'}\right] = (2\pi)^3\delta^{(3)}(\boldsymbol{p} - \boldsymbol{p}') \tag{2.0.6}$$

Hamiltonian in terms of ladder operator

$$H = \int \frac{\mathrm{d}^3 p}{(2\pi)^3}E_p\left(a_{\boldsymbol{p}}a^\dagger_{\boldsymbol{p}} + \frac{1}{2}\left[a_{\boldsymbol{p}}, a^\dagger_{\boldsymbol{p}}\right]\right) \tag{2.0.7}$$

Normlisation it's also lorentz-invariante

$$\langle p|q\rangle = 2E_p(2\pi)^3\delta^{(3)}(\boldsymbol{p} - \boldsymbol{q}) \tag{2.0.8}$$

2.1 Heisenberg-picture fields

Heisenberg-picture

$$|\psi_H\rangle = e^{iHt}\,|\psi_s(t)\rangle \tag{2.1.1}$$

$$O_H(t) = e^{iHt}O_S\,e^{-iHt} \tag{2.1.2}$$

Field operator

$$\phi(x) = \phi(\boldsymbol{x}, t) = \int \frac{\mathrm{d}^3 p}{(2\pi)^3}\frac{1}{\sqrt{2E_p}}\left(a_p e^{ipx} + a^\dagger_p e^{-ipx}\right) \tag{2.1.3}$$

2.2 Commutations and propogators

Commutations

$$[\phi(x), \phi(y)] = D(x-y) - D(y-x) \begin{cases} = 0 & \text{if } (x-y) \text{ is space-like} \\ \neq 0 & \text{otherwise} \end{cases} \tag{2.2.1}$$

$$D(x-y) = \int \frac{\mathrm{d}^3 p}{(2\pi)^3} \frac{1}{2E_p} e^{-ip(x-y)} \tag{2.2.2}$$

Propogator

$$\langle 0|\phi(x)\phi(y)|0\rangle = D(x-y) \tag{2.2.3}$$

Feynman propagator

$$\begin{aligned} D_F(x-y) &= \langle 0|T\phi(x)\phi(y)|0\rangle \\ &= \Theta(x^0 - y^0)D(x-y) + \Theta(y^0 - x^0)D(y-x) \end{aligned} \tag{2.2.4}$$

$$D_F(x-y) = \int \frac{\mathrm{d}^4 p}{(2\pi)^4} \frac{i}{p^2 - m^2 + i\epsilon} e^{-ip(x-y)} \tag{2.2.5}$$

3 Quantization of the Dirac field

3.1 Dirac equation

$$\left(i\gamma^\mu \partial_\mu - m\right)\phi(x) = 0 \tag{3.1.1}$$

Standard representation (Dirac's)

$$\gamma_0 = \begin{pmatrix} \mathbb{1}_2 & 0 \\ 0 & -\mathbb{1}_2 \end{pmatrix} \quad \boldsymbol{\gamma} = \begin{pmatrix} 0 & \boldsymbol{\sigma} \\ -\boldsymbol{\sigma} & 0 \end{pmatrix} \tag{3.1.2}$$

Lorentz transformation

$$\Lambda = \exp\left(\frac{1}{2}\omega_{\mu\nu} M^{\mu\nu}\right) \tag{3.1.3}$$

with ω set of parameters and M the generator of Lie algebra.

Spinor representation

$$S^{\rho\sigma} = \frac{1}{4}\left[\gamma^\rho, \gamma^\sigma\right] = \frac{1}{2i}\sigma^{\rho\sigma} \tag{3.1.4}$$

$$\tag{3.1.5}$$

Spinor transformation

$$S(\Lambda) = \exp\left(\frac{1}{2}\omega_{\mu\nu} S^{\mu\nu}\right) \tag{3.1.6}$$

$$\psi'_a(x) = S_{ab}(\Lambda)\psi_b(\Lambda^{-1}x) \tag{3.1.7}$$

adjoint spinor

$$\bar{\psi} = \psi^\dagger \gamma^0 \tag{3.1.8}$$

Fifth gamma matrix

$$\gamma^5 := i\gamma^0\gamma^1\gamma^2\gamma^3 \tag{3.1.9}$$

$$\left\{\gamma^\mu, \gamma^5\right\} = 0 \tag{3.1.10}$$

$$(\gamma^5)^2 = \mathbb{1}_4 \tag{3.1.11}$$

Plane wavesolutions

$$\psi(x) = \begin{cases} u(p)e^{-ipx} & \text{positive frequency} \\ v(p)e^{ipx} & \text{negative frequency} \end{cases} \tag{3.1.12}$$

$$u_s(p) = \sqrt{E_p + m}\begin{pmatrix} \chi_s \\ \frac{\boldsymbol{u}\cdot\boldsymbol{p}}{E_p+m}\chi_s \end{pmatrix} e^{-ipx} v_s(p) = \sqrt{E_p + m}\begin{pmatrix} \frac{\boldsymbol{u}\cdot\boldsymbol{p}}{E_p+m}\tilde{\chi}_s \\ \tilde{\chi}_s \end{pmatrix} e^{ipx} \tag{3.1.13}$$

with

$$\chi_{\frac{1}{2}} = \begin{pmatrix} 1 \\ 0 \end{pmatrix} \quad x_{-\frac{1}{2}} = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$$
$$s = \pm\frac{1}{2} \quad \tilde{\chi}_s = \chi_{-s}$$

Orthogonality of spinor

$$\bar{u}_s(p)u_{s'}(p) = -\bar{v}_s(p)v_{s'}(p) = 2m\delta_{ss'} \tag{3.1.14}$$
$$\bar{u}_s(p)v_{s'}(p) = 0 \tag{3.1.15}$$

Spin sums

$$\sum_s u_s(p)\bar{u}_s(p) = \not{p} + m \tag{3.1.16}$$
$$\sum_s v_s(p)\bar{v}_s(p) = \not{p} - m \tag{3.1.17}$$

3.2 Dirac Lagrangian and quantization

$$\mathcal{L} = \bar{\psi}(x)(i\not{\partial} - m)\psi(x) \tag{3.2.1}$$

Quantization

$$\left\{\psi_a(\boldsymbol{x}), \psi_b^\dagger(\boldsymbol{x}')\right\} = \delta_{ab}\delta^{(3)}(\boldsymbol{x} - \boldsymbol{x}') \tag{3.2.2}$$
$$\{\psi_a(\boldsymbol{x}), \psi_b(\boldsymbol{x}')\} = \left\{\psi_a^\dagger(\boldsymbol{x}), \psi_b^\dagger(\boldsymbol{x}')\right\} = 0 \tag{3.2.3}$$

Field operators

$$\psi(\boldsymbol{x}) = \int \frac{\mathrm{d}^3 p}{(2\pi)^3\sqrt{2E_p}} \sum_s (a_{\boldsymbol{p}}^s u_s(\boldsymbol{p})e^{i\boldsymbol{p}\cdot\boldsymbol{x}} + b_{\boldsymbol{p}}^{s\,\dagger} v_s(\boldsymbol{p})e^{-i\boldsymbol{p}\cdot\boldsymbol{x}}) \tag{3.2.4}$$

thus the anticommutations of ladder operators:

$$\left\{a_{\boldsymbol{p}}^s, a_{\boldsymbol{p}'}^{s'\,\dagger}\right\} = \left\{b_{\boldsymbol{p}}^s, b_{\boldsymbol{p}'}^{s'\,\dagger}\right\} = (2\pi)^3\delta_{ss'}\delta^{(3)}(\boldsymbol{p} - \boldsymbol{p}')$$
$$\{a, a\} = \left\{a^\dagger, a^\dagger\right\} = \cdots = 0$$

Hamiltonian in terms of Fourier modes (with normal ordering)

$$H = \int \frac{\mathrm{d}^3 p}{(2\pi)^3} \sum_s E_p (a_{\boldsymbol{p}}^{s\,\dagger} a_{\boldsymbol{p}}^s - b_{\boldsymbol{p}}^{s\,\dagger} b_{\boldsymbol{p}}^s) \tag{3.2.5}$$

3.3 Particles and antiparticles

$$Q = e \int \mathrm{d}^3 x \psi^\dagger(x)\psi(x) \tag{3.3.1}$$

$$: Q : = e \int \frac{\mathrm{d}^3 p}{(2\pi)^3} \sum_s (a_{\boldsymbol{p}}^{s\,\dagger} a_{\boldsymbol{p}}^s - b_{\boldsymbol{p}}^{s\,\dagger} b_{\boldsymbol{p}}^s) \tag{3.3.2}$$

3.4 Dirac propagator and anticommutators

$$\begin{aligned} S_{ab}(x-y) &= \{\psi_a(x), \bar{\psi}_b(y)\} \\ &= (i\not{\partial} + m)\left[D(x-y) - D(y-x)\right] \end{aligned} \tag{3.4.1}$$

Time ordering of Dirac fields

$$T(\phi_a(x)\bar{\psi}_b(y)) = \Theta(x^0 - y^0)\psi_a(x)\bar{\psi}_b(y) - \Theta(y^0 - x^0)\bar{\psi}_b(y)\psi_a(x) \tag{3.4.2}$$

Feynman propogator for the Dirac field

$$S_F(x-y) = \langle 0|T\psi(x)\bar{\psi}(y)|0\rangle = \int \frac{\mathrm{d}^4 p}{(2\pi)^4} \frac{i(\not{p} + m)}{p^2 - m^2 + i\epsilon} e^{-ip\cdot(x-y)} \tag{3.4.3}$$

3.5 Discrete symmetries of the Dirac Field

	orientation perserving	orientation not perserving
(ortho)chronous	$\mathcal{L}_+^\uparrow$	$\mathcal{L}_-^\uparrow = \mathcal{P}\mathcal{L}_+^\uparrow$
non-orthochronous	$\mathcal{L}_-^\downarrow = \mathcal{T}\mathcal{L}_+^\uparrow$	$\mathcal{L}_+^\downarrow = \mathcal{P}\mathcal{T}\mathcal{L}_+^\uparrow$

4 Interacting QFT

4.1 Introduction and examples

Theories discussed so far are Klein-Gordon theory with spin 0

$$\mathcal{L}_{KG} = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2$$

and Dirac theory spin $\frac{1}{2}$

$$\mathcal{L}_D = \bar{\psi}(i\partial\!\!\!/ - m)\psi$$

There is also $\mathcal{L}_{EM} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$ with $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ for a massless vector filed. Its quantiasation gives photon.

One thing they have in common is quadratic in the fields. As result:

- linear field equations
- exact quantisation
- multi-particle states without scattering or interaction
- linear fourier decompositions , no mementum changes

To have an interacting theory with scattering, need higher powers in the field in the Lagrangians. A few examples are following

Scalar ϕ^4 theory

$$\mathcal{L} = \mathcal{L}_{KG} + \frac{\lambda}{4!}\phi^4$$

need positive sign $\lambda > 0$ for a stable theory, otherwise classical energy can be arbitarily negative.

Equation of motions

$$(\partial^2 + m^2)\phi = -\frac{\lambda}{3!}\phi^3$$

is nonlinear, cannot be solved by Fourier decomposition.

Yukawa-theory

$$\mathcal{L} = \mathcal{L}_{KG} + \mathcal{L}_D - g\bar{\psi}\psi\phi$$

It is originally developed as a theory for nuclear forces with ψ nucleon, ϕ pion. In the Standard Model it is similar to interactions in Higgs mechanism.

Quantum Electrodynamics (QED)

$$\mathcal{L} = \mathcal{L}_{EM} + \mathcal{L}_D - eA_\mu \bar{\psi}\gamma^\mu\psi$$

descreibes electrons, their antiparticles positrons and photons.

Yang-Mills theory generalises $\mathcal{L}_{EM}$ with terms like A^4 or $A^2\partial A$

Scalar QED descreibes pions and photons

$$\begin{aligned}\mathcal{L} &= \mathcal{L}_{EM} + D_\mu\phi D^\mu\phi^* - m^2|\phi|^2 \\ &= \mathcal{L}_{EM} + \partial_\mu\phi\partial^\mu\phi^* - m^2\phi\phi^* + ieA_\mu(\phi\partial^\mu\phi^* - \phi^*\partial^\mu\phi) + e^2A_\mu A^\mu\phi\phi^*\end{aligned}$$

Remarks

1. Interaction terms in $H_{\text{int}} = \int \mathrm{d}^3x\mathcal{H}_{\text{int}} = -\int \mathrm{d}^3x\mathcal{L}_{\text{int}}$ always involves products of fields at the same point $\boldsymbol{x}$. It ensures causality, no "instant at a distance".

2. There are no derivative interactions. These may complicate quantisation as

$$\pi(\boldsymbol{x}) = \frac{\partial\mathcal{L}}{\partial(\partial_0\phi(\boldsymbol{x}))}$$

3. Why are we taking the examples above? There must be zillions of theories (Lagrangians)?
We have the criterion of **renormalizability**. Note the mass dimensions of fields;

$$[S] = 1 \text{ so } [\mathcal{L}] = [M]^4 \;\Rightarrow [\phi] = [M],\, [\psi] = [M]^{\frac{3}{2}},\, [A_\mu] = [M]$$

So in all the interaction terms indicated above, the coupling constant λ, e, g are all **dimensionless**!
Can add $-\frac{\mu}{3!}\phi^3$ to the ϕ^4 theory. This leads to $[\mu] = [M]$ and all these generate renormalisable interactions.
All higher interaction terms require coupling constants of **negative** mass dimension, e.g. $G\bar{\psi}\psi\bar{\psi}\psi$ and then $[G] = [M]^{-2}$. These are nonrenormalisable and create trouble when performing higher-order calculation in perturbation theory. (with energy cutoff; corrections $G\Lambda^2$, $\Lambda \to \infty$)

4. we haven't quantised the photon yet. The reason is that its is a vector field, i.e. 4 degrees of freedom, but photon has just 2 physical polarisaion states. It is linked to gauge symmetry and complicates quantisation somewhat.

4.2 The interaction picture

Consider the ϕ^4 theory,

$$\mathcal{L}_{int} = -\frac{\lambda}{4!}\phi(x)^4 \tag{4.2.1}$$

Hamiltonian $H = H_0 + H_{int}$ with

$$H_0 = \int \mathrm{d}^3x \left\{\frac{1}{2}\pi^2(x) + \frac{1}{2}(\boldsymbol{\nabla}\phi)^2 + \frac{1}{2}m^2\phi^2\right\} \tag{4.2.2}$$

$$H_{int} = -\int \mathrm{d}^3x\mathcal{L}_{int} = \frac{\lambda}{4!}\int \mathrm{d}^3x\phi^4 \tag{4.2.3}$$

Interaction picture means that operators evolve in time using H_0 (only), in particular

$$\phi_I(t,\boldsymbol{x}) = e^{iH_0t}\phi(\boldsymbol{x})e^{-iH_0t} \tag{4.2.4}$$

Time-dependence of the free field, obeys classical equation of motion $\left(\partial^2 + m^2\right)\phi_I(t,\boldsymbol{x}) = 0$. Solution in terms if fourier modes as before:

$$\phi_I = \int \frac{\mathrm{d}^3 p}{(2\pi)^3\sqrt{2E_p}}(a^I_{\boldsymbol{p}}e^{-ipx} + a^{I\,\dagger}_{\boldsymbol{p}}e^{+ipx}) \tag{4.2.5}$$

as in the free theory with standard commutation relations $[a^I_{\boldsymbol{p}}, a^{I\,\dagger}_{\boldsymbol{p}}] = (2\pi)^3\delta^{(3)}(\boldsymbol{p}-\boldsymbol{p}')$. The state satisfing $a^I_p\,|0\rangle = 0$ is the vacuum of the free, noninteratcting theory.

Relation between interaction and Schrödinger picure states:

$$|\phi_I(t)\rangle = e^{iH_0t}|\psi_S(t)\rangle \tag{4.2.6}$$

Schrödinger equation becomes:

$$\begin{aligned}
i\frac{\partial}{\partial t}\,|\psi_S\rangle &= (H_0 + H_{\text{int}})\,|\psi_S\rangle \\
\text{LHS} &= i\frac{\partial}{\partial t}\left(e^{-iH_0t}\,|\phi_I\rangle\right) = H_0e^{-iH_0t}\,|\phi_I\rangle + e^{-iH_0t}i\frac{\partial}{\partial t}\,|\phi_I\rangle \\
\text{RHS} &= (H_0 + H_{\text{int}})\,e^{-iH_0t}\,|\phi_I\rangle \\
\Rightarrow i\frac{\partial}{\partial t}\,|\phi_I\rangle &= e^{iH_0t}H_{int}e^{-iH_0t} = H_I(t)\,|\phi_I\rangle
\end{aligned} \tag{4.2.7}$$

with H_I interaction Hamiltonian in the interaction picture. Clearly

$$H_I = \frac{\lambda}{4!}\int \mathrm{d}^3x\phi_I^4(x)$$

What is the solution of 4.2.7 for the time evolution of $|\phi_I(t)\rangle$? Define time-evolution operator in the interaction picture.

$$|\phi_I(t)\rangle = U(t,t_0)\,|\phi_I(t_0)\rangle \tag{4.2.8}$$

$$\text{where } U(t,t_0) = e^{iH_0(t-t_0)}e^{-iH(t-t_0)} \tag{4.2.9}$$

With 4.2.7 and 4.2.8:

$$i\frac{\partial}{\partial t}U(t,t_0) = H_I(t)U(t,t_0) \tag{4.2.10}$$

To solve with boundary conditions: $U(t_0,t_0) = \mathbb{1}$. The formal solution:

$$U(t,t_0) = 1 - i\int_{t_0}^{t}\mathrm{d}t'H_I(t')U(t',t_0)$$

Substitute back in and we get:

$$U(t,t_0) = 1 - i\int_{t_0}^{t}\mathrm{d}t'H_I(t') + (-i)^2\int_{t_0}^{t}\mathrm{d}t'\int_{t_0}^{t'}\mathrm{d}t''H_I(t')H_I(t'') + \ldots$$

H_I inside the integral is automatically time-ordered. Ranges of integration is not.

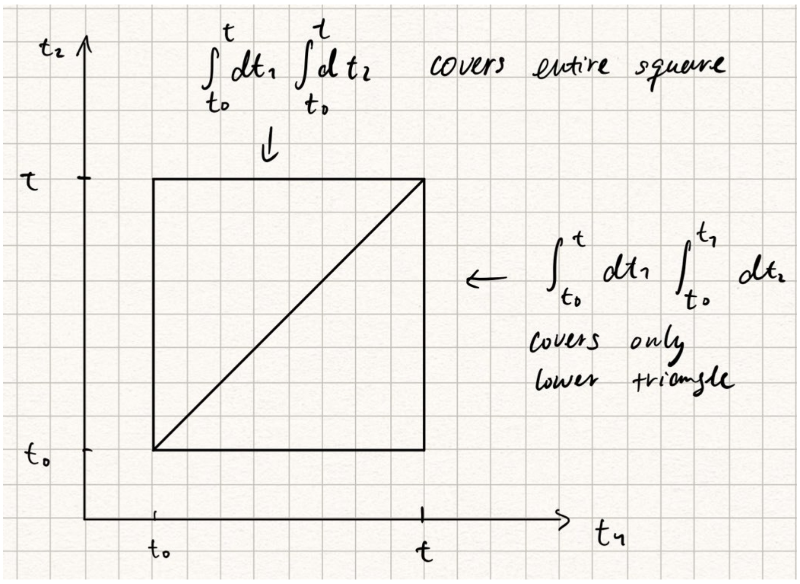


Figure 4.1: Time ordering

Upper triangle has the wrong time order. We are going to "repair" it by hand.

$$
\begin{aligned}
U(t,t_0) &= 1 - i\int_{t_0}^{t} \mathrm{d}t' H_I(t') + \frac{(-i)^2}{2}\int_{t_0}^{t}\mathrm{d}t'\int_{t_0}^{t}\mathrm{d}t'' T(H_I(t')H_I(t'')) + \dots \\
&= \sum_{n=0}^{\infty}\frac{(-i)^n}{n!}\int_{t_0}^{t}\mathrm{d}t_1 \cdots \int_{t_0}^{t}\mathrm{d}t_n T(H_I(t_1)\dots H_I(t_n)) \\
&= T\exp\left\{-i\int_{t_0}^{t}\mathrm{d}t' H_I(t')\right\} \qquad (4.2.11)
\end{aligned}
$$

It is interesing for scattering to transition into asymptotic state for $t \to \infty$

$$
\begin{aligned}
S \;&=\; \lim_{t\to\infty} U(t,-t) = T\exp\left\{-i\int_{-\infty}^{\infty}\mathrm{d}t H_I(t)\right\} \qquad (4.2.12) \\
&\overset{\phi^4}{=} T\exp\left\{-i\int \mathrm{d}^4x \frac{\lambda}{4!}\phi_I^4(x)\right\}
\end{aligned}
$$

Both U and S are formally unitary

Composition law for time evolution operator:

$$
U(t_2,t_0) = U(t_2,t_1)U(t,t_0) = U(t_2,t_1)U(t_0,t_1)^\dagger \qquad (4.2.13)
$$

4.2.1 Scattering amplitudes and the S-matrix

Take $|i\rangle$ the initial (multi-particle) state and $|f\rangle$ the final (multi-particle) state. Time evolution of $|i\rangle$ then is

$$
\lim(t\to\infty)U(t,-\infty)\,|i\rangle = S\;|i\rangle
$$

Probability that $|i\rangle$ evolves into $|f\rangle$ is proporional to the squared "**S-matrix element**"

$$
|\langle f,\ t\to\infty|i,t\to-\infty\rangle|^2 = |\langle f|S\,|i\rangle|^2 = \left|S_{fi}\right|^2 \qquad (4.2.14)
$$

The nontrivial part of the S-matrix is the T-matrix:

$$S_{fi} := \delta_{fi} + iT_{fi} \tag{4.2.15}$$

Use momentum conservation (from translation invariance) to define matrix element

$$S_{fi} = \delta_{fi} + i(2\pi)^4\delta^{(4)}(p_f - p_i)M_{fi} \tag{4.2.16}$$

M_{fi} measures "genuine scattering" from $|i\rangle$ to $|f\rangle$.
How are we going to calculate correlation functions in the interacting theory:

$$\langle\Omega|\, T\phi(x)\phi(y)|\Omega\rangle \tag{4.2.17}$$

or more generally $\langle\Omega|T\phi(x_1)\phi(x_2)\dots|\Omega\rangle$, where $|\Omega\rangle$ is the vaccum/ground state of the interacting theory and $\phi(x)$ the Heisenberg operators?

Ignore $|\Omega\rangle \neq |0\rangle$ for the moment other than saying: we want to study the time evolution from the vacuum at $t \to -\infty$ to $t \to +\infty$. So rewriting in terms $\phi_I(x)$, assuming $x^0 > y^0$ for now:

$$\langle 0|U(\infty, x^0)\phi_I(x^0)U(x^0, y^0)\phi_I(y^0)U(y^0, -\infty)|0\rangle = \langle 0|T(\phi_I(x)\phi_I(y)S)|0\rangle \tag{4.2.18}$$

still holds if $x^0 < y^0$ because of T.

Now $|\Omega\rangle \neq |0\rangle$: this can be taken care of by dividing out the time evolution of the (free) vacuum $\langle 0|S|0\rangle$, so

$$\begin{aligned}\langle\Omega|T(\phi(x)\phi(y))|\Omega\rangle&\\ &= \frac{\langle 0|T(\phi_I(x)\phi_I(y)S)|0\rangle}{\langle 0|S|0\rangle}\\ &\overset{\phi^4}{=} \frac{\langle 0|T\phi_I(x)\phi_I(y)\exp\left\{-i\int \mathrm{d}^4x' \frac{\lambda}{4!}\phi^4(x')\right\}|0\rangle}{\langle 0|T\exp\left\{-i\int \mathrm{d}^4x' \frac{\lambda}{4!}\phi^4(x')\right\}|0\rangle}\end{aligned} \tag{4.2.19}$$

Proof can be found in Peskin. It will also be illustrated parctically later ("vacuum bubbles").

Perturbation theory is viable when λ (or some other coupling) is "small" and then expands $U(t, t_0)$ or S in powers of λ.

4.3 Wick's theorem

From now on drop the subscript for interaction pictire fields $\phi_I(x) \to \phi(x)$.

Want to calculate stuff like $\langle 0|T\phi(x_1)\dots\phi(x_n)S|0\rangle$ in perturbation theory; so e.g. at order λ^n. So

$$\frac{1}{n!}\left(-i\frac{\lambda}{4!}\right)^n \int \mathrm{d}^4y_1\dots\mathrm{d}^4y_n\,\langle 0|T\phi(x_1)\dots\phi(x_n)\phi^4(y_1)\dots\phi^4(y_n)|0\rangle \tag{4.3.1}$$

is tough!

We know $\langle 0|T\phi(x_1)\phi(x_2)|0\rangle$ is the Feynman propagator!

Recall **normal ordering** with $\phi(x) = \phi^+(x) + \phi^-(x)$

$$:\phi^+\phi^- :=: \phi^-\phi^+ := \phi^-\phi^+ \tag{4.3.2}$$

where

$$\phi^+ = \int \frac{\mathrm{d}^3 p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\boldsymbol{p}}}} a_{\boldsymbol{p}} e^{-ip\cdot x}$$
$$\phi^- = \int \frac{\mathrm{d}^3 p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\boldsymbol{p}}}} a^\dagger_{\boldsymbol{p}} e^{+ip\cdot x}$$

Wick's therem expresses time-ordered products in terms of normal-ordered ones. Then it is easy to take vacuum expectation values, as $\langle 0| : \phi(x_1)\ldots\phi(x_n) : |0\rangle = 0$

Take two fields and $x^0 > y^0$:

$$\begin{aligned} T\phi(x)\phi(y) = \phi(x)\phi(y) &= \left(\phi^+(x) + \phi^-(x)\right)\left(\phi^+(y) + \phi^-(y)\right) \\ &= \phi^+(x)\phi^+(y) + \phi^-(x)\phi^-(y) + \phi^-(x)\phi^+(y) + \phi^-(x)\phi^+(y) + [\phi^+(x), \phi^-(y)] \\ &=: \phi(x)\phi(y) : + [\phi^+(x), \phi^-(y)] \end{aligned}$$

Particularly for $y^0 > x^0$:

$$T\phi(x)\phi(y) =: \phi(x)\phi(y) : + [\phi^+(y), \phi^-(x)]$$

Thus altogether:

$$T\phi(x)\phi(y) =: \phi(x)\phi(y) : + D_f(x-y) \tag{4.3.3}$$

as $\Theta(x^0 - y^0)[\phi^+(x), \phi^-(y)] + \Theta(y^0 - x^0)[\phi^+(y), \phi^-(x)] = D_F(x-y)$.

Worth noting that $D_F(x-y)$ is still a c-number, not operator (yet). Thus it can be pulled out of any matrix element or expectation value.

We now define "contraction":

$$\overbrace{\phi(x_1)\phi(x_2)} = D_F(x_1 - x_2) \tag{4.3.4}$$

Thus we can remove the fields from the product leaving only the propagators:

$$T\phi(x)\phi(y) =: \phi(x)\phi(y) : + \overbrace{\phi(x)\phi(y)} \tag{4.3.5}$$

General form of **Wick's theorem** for arbitary number of fields

$$T\phi(x_1)\ldots\phi(x_n) =: \phi(x_1)\ldots\phi(x_n) : + : (\text{sum over all possible contractions}) : \tag{4.3.6}$$

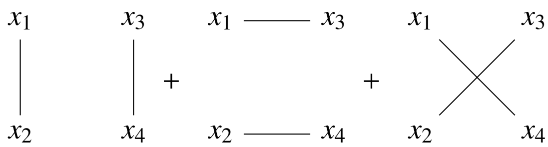
Example with four fields:

$$\begin{aligned} T(\phi_1\phi_2\phi_3\phi_4) =&: \phi_1\phi_2\phi_3\phi_4 : \\ &+ \overbrace{\phi_1\phi_2} : \phi_3\phi_4 : + \overbrace{\phi_1\phi_3} : \phi_2\phi_4 : + \overbrace{\phi_1\phi_4} : \phi_2\phi_3 : + \overbrace{\phi_2\phi_3} : \phi_1\phi_4 : + \overbrace{\phi_2\phi_4} : \phi_1\phi_3 : + \overbrace{\phi_3\phi_4} : \phi_1\phi_2 : \\ &+ \overbrace{\phi_1\phi_2}\overbrace{\phi_3\phi_4} + \overbrace{\phi_1\phi_3}\overbrace{\phi_2\phi_4} + \overbrace{\phi_1\phi_4}\overbrace{\phi_2\phi_3} \end{aligned}$$

Thus

$$\langle 0| T\left(\phi_1\phi_2\phi_3\phi_4\right)|0\rangle = D_F(x_1-x_2)D_F(x_3-x_4) + D_F(x_1-x_3)D_F(x_2-x_4) + D_F(x_1-x_4)D_F(x_2-x_3)$$

which can be visually represented as



Proof of the general theorem by *induction* in the number of fields (see exercise). The idea is to suppose it is true for $\phi_2 \dots \phi_m$, $x_1^0 > x_{k>1}^0$. Then

$$\begin{aligned} T\phi_1\phi_2 \dots \phi_m &= (\phi_1^+ + \phi_1^-)T\phi_2 \dots \phi_m \\ &= (\phi_1^+ + \phi_1^-)[: \phi_2 \dots \phi_m : + : \text{contractions} :] \end{aligned}$$

ϕ_1^- can stay as it is part of (: $\phi_1\phi_2 \dots \phi_m$:). But ϕ_1^+ needs to be comuted past all ϕ_1^- operators, giving rise to additional contractions $\phi_1\phi_2$.

Consequences

- $n = 2k + 1$, $k \in \mathbb{N}$
$$\langle 0|T\phi_1 \dots \phi_m|0\rangle = 0$$
- $n = 2k$, $k \in \mathbb{N}$
$$\langle 0|T\phi_1 \dots \phi_m|0\rangle = \sum_{\text{pairing of fields}} D_F(x_{i_1} - x_{i_2}) \dots D_F(x_{i_{m-1}} - x_{i_m})$$

4.3.1 Wick's theorem and the S-Matrix

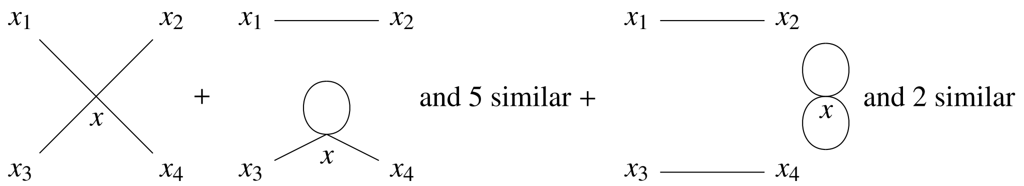
Apply Wick's theorem to correlation functions $\langle 0|T\{\phi_1 \dots \phi_m\}S|0\rangle$ n-th term in the perturbative expansion of S with $\phi(x_1) := \phi_1$.

$$\frac{1}{n!}\left(\frac{-i\lambda}{4!}\right)^n \int \mathrm{d}^4y_1 \dots \mathrm{d}^4y_n \langle 0|T\{\phi_1 \dots \phi_m\phi^4(y_1) \dots \phi^4(y_n)\}|0\rangle$$

Example with $m = 4$, $n = 1$

$$\begin{aligned} &-\frac{i\lambda}{4!}\int \mathrm{d}^4x \langle 0|T\phi_1\phi_2\phi_3\phi_4\phi^4(x)|0\rangle \\ &= -\frac{i\lambda}{4!}\int \mathrm{d}^4x \langle 0|\phi_1\phi_2\phi_3\phi_4\phi(x)\phi(x)\phi(x)\phi(x)|0\rangle + 23 \text{ permutations} \\ &\quad -\frac{i\lambda}{4!}\int \mathrm{d}^4x \langle 0|\phi_1\phi_2\phi_3\phi_4\phi(x)\phi(x)\phi(x)\phi(x)|0\rangle + 11 \text{ permutations} + 5 \text{ similar} \\ &\quad -\frac{i\lambda}{4!}\int \mathrm{d}^4x \langle 0|\phi_1\phi_2\phi_3\phi_4\phi(x)\phi(x)\phi(x)\phi(x)|0\rangle + 2 \text{ permutations} + 2 \text{ similar} \\ &= -i\lambda\int \mathrm{d}^4x D_F(x_1 - x)D_F(x_2 - x)D_F(x_3 - x)D_F(x_4 - x) \\ &\quad -\frac{i\lambda}{2}D_F(x_1 - x_2)\int \mathrm{d}^4x D_F(x_3 - x)D_F(x_4 - x)D_F(x - x) + 5 \text{ similar} \\ &\quad -\frac{i\lambda}{8}D_F(x_1 - x_2)D_F(x_3 - x_4)\int \mathrm{d}^4x D_F(x - x) + 2 \text{ similar} \end{aligned}$$

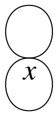
Permutation means permutation of $\phi(x)$ and similar means exchanging ϕ_i, $i \in 1, 2, 3, 4$ without changing the diagram. Represented in Feynman diagrams:



In fact $D_F(x-x) = D_F(0)$ diverges!

Example with $m=0$, $n=1$ vacuum diagram

$$
\begin{aligned}
&-\frac{i\lambda}{4!}\int \mathrm{d}^4x\,\langle 0|T\phi^4(x)|0\rangle \\
&= -\frac{i\lambda}{8}[D_F(0)]^2\int \mathrm{d}^4x \\
&=
\end{aligned}
$$



Example: 2nd order S-matrix term

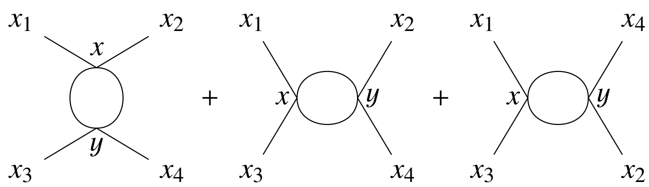
$$
\frac{1}{2!}\left(\frac{-i\lambda}{4!}\right)^4\int \mathrm{d}^4x\mathrm{d}^4y\,\langle 0|T\phi_1\phi_2\phi_3\phi_4\phi^4(x)\phi^4(y)|0\rangle
$$

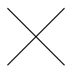
It has many contractions and some of the fully connected ones are of the type there are

(4×3)[choose $\phi(x)$] $\times$ (4×3)[choose $\phi(y)$] $\times$ 2[x-y-cont.] $\times$ 2(x-y-symm.) + 2 similar, exchanging external points

$$
= \frac{(-i\lambda)^2}{2}\int \mathrm{d}^4x\mathrm{d}^4y D_F(x_1-x)D_F(x_2-x)D_F(x_3-y)D_F(x_4-y)[D_F(x-y)]^2 + 2\text{ similar}
$$

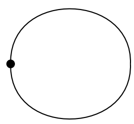
=

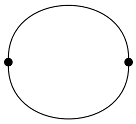
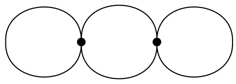


Symmetry factors A lot of the contractions eliminate the factors $\frac{1}{n!}\left(\frac{1}{4!}\right)^4$ in the denominators; the $\frac{1}{4!}$ was chosen as to yield  $\sim -i\lambda$

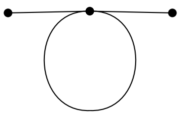
See examples above. Sometimes, factors are not completely cancelled and thus procedure gets "reversed". Divide diagrams by *symmetry factor* $\hat{=}$ the "missing factors".

Where does it come from?

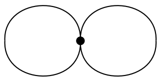
- factor 2 from the line that starts and ends at the same point. 

- two (or more) lines linking the same 2 points. 
- 2 vertices can be equivalent. 

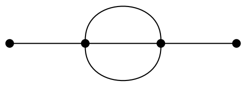
When in doubt, can always go back to Wick's theorem and count the contractions explicitely.
Examples:



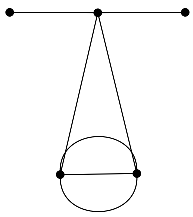
$S = 2$



$S = 8$



$S = 3! = 6$



$S = 3! = 6$

Summary of Feynman rules

$$\langle 0|T\phi_1 \dots \phi_m \exp\left(-\frac{i\lambda}{4!}\int d^4x \phi^4(x)\right)|0\rangle$$

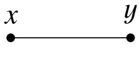
= sum of all diagrams with m external points;

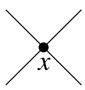
usually organised by number of internal points(i.e. power of λ).

Each diagram built cut of

- propagators
- vertices (n)
- external points (m)

Feynman rules in position space Analytic expression obtained by combining

- for each propagator 

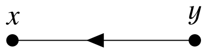
$= D_F(x - y)$
- for each vertex 

$= -i\lambda \int d^4x$

- for each external point $\overset{x}{\bullet}\!\!\rule{2cm}{0.4pt} = 1$
- divide diagram by its symmetry factor S

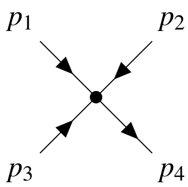
Since the propagator $D_F(x-y) = \int \frac{\mathrm{d}^4p}{(2\pi)^4} \frac{i}{p^2-m^2+i\epsilon} e^{-ip(x-y)}$. It is actually simpler to express these in momentum space instead.

The way to do it is to sssign a momentum p to each propagator. (direction arbitary)



- assign e^{ipy} to y-vertex (arrow out)
- assign e^{-ipx} to x-vertex (arrow in)
- $\frac{i}{p^2-m^2+i\epsilon}$ to the line and the integration $\int \frac{\mathrm{d}^4p}{(2\pi)^4}$

At vertex x:

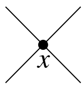


$$= -i\lambda \int \mathrm{d}^4x e^{-i(p_1+p_2+p_3)x+ip_4x}$$

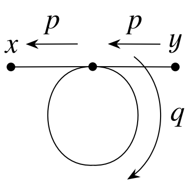
$$= -i\lambda(2\pi)^4\delta^{(4)}(p_1+p_2+p_3-p_4)$$

This imposes momentum conservation at vetex. $\delta^{(4)}$-functions make some of the momentum integrals trivial, always with $(2\pi)^4$ cancelled appropriately.

Momentum space Feynman rules

- propagator $\xrightarrow{p} = \frac{i}{p^2-m^2+i\epsilon}$
- vertex (position integrated out)  $= -i\lambda$
- external points $\begin{cases} e^{-ipx} & \text{incoming} \\ e^{+ipx} & \text{outgoing} \end{cases}$
- impose momentum conservation at each vertex
- integrate over each undetermined momentum $\int \frac{\mathrm{d}^4p}{(2\pi)^4}$
- divide by symmetry factor

e.g.:

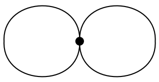
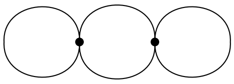


$$= (-i\lambda) \int \frac{\mathrm{d}^4p}{(2\pi)^4} \frac{\mathrm{d}^41}{(2\pi)^4} \left(\frac{i}{p^2-m^2+i\epsilon}\right)^2 \frac{i}{q^2-m^2+i\epsilon} e^{-ip(x-y)}$$

Vacuum diagrams

Disconnected pieces in Feynman diagrams are pretty bad. Not only $D_F(0) = \int \frac{\mathrm{d}^4 p}{(2\pi)^4} \frac{i}{p^2 - m^2 + i\epsilon}$ is divergent (that will be taken care of later), it also contains an integral $\int \mathrm{d}^4 x$ const. thus divergent once more!

Typical diagram contributing to 2-point function. one piece connected to x and y, plus disconnected pieces.

Call disconnected pieces $V_i \in \Big\{$  ,  , $\ldots \Big\}$. Points are connected internally, but not to external points.

V_i can occur n_i-times, then

$$[\text{diagram}] = [\text{connected pieces}] \times \prod_i \frac{1}{n!} (V_i)^{n_i}$$

The factorial is the symmetry factor of n_i disconnected copies of V_i.

Then

$$\begin{aligned}
&\langle 0 | T\phi_1 \ldots \phi_n S | 0 \rangle \\
&= \sum_{\text{connected}} \sum_{\text{all}\{n_i\}} [\text{connected}] \times \prod_i \frac{1}{n_i!} (V_i)^{n_i} \\
&= \Big(\sum [\text{connected}] \Big) \times \sum_{\text{all}\{n_i\}} \left(\prod (i) \frac{1}{n_i!} (V_i)^{n_i} \right) \\
&= \Big(\sum [\text{connected}] \Big) \times \prod_i \left(\sum_{n_i} \frac{1}{n_i!} (V_i)^{n_i} \right) \\
&= \Big(\sum [\text{connected}] \Big) \times \exp\left(\sum_i V_i \right)
\end{aligned}$$

Thus

$$\begin{aligned}
\text{sum of ALL diagrams} &= (\text{sum of all CONNECTED diagrams}) && (4.3.7) \\
&\quad \times \exp(\text{sum of all DISCONNECTED diagrams}) && (4.3.8)
\end{aligned}$$

Obvious from the above:

$$\langle 0 | S | 0 \rangle = \langle 0 | T \{ \exp\left(-\frac{i\lambda}{4!} \int \mathrm{d}^4 x \phi^4(x) \right) \} | 0 \rangle = \exp(\text{sum of all vacuum bubbles})$$

Conclusion from the (unproven) formula for n-point correlation functions in the true, interacting vacuum:

$$\begin{aligned}
\langle \Omega | T\phi_1 \ldots \phi_m | \Omega \rangle &= \frac{\langle 0 | T\phi_1 \ldots \phi_m S | 0 \rangle}{\langle 0 | S | 0 \rangle} && (4.3.9) \\
&= \sum (\text{connected diagrams with m external points}) && (4.3.10)
\end{aligned}$$

Here: "connected" means connected to any external point. External points do not have to linked to each other.

4.4 S-matrix elements and Feynman diagrams

What is the correlation function in interacting vacuum $\langle\Omega|T\phi_1\dots\phi_m|\Omega\rangle$ good for? For scattering, shouldn't we rather look at $\langle p_1\dots p_m|S|p_Ap_B\rangle$ with the perturbative expansion of S as before?

Decompose the S-matrix

$$S_{fi}=\delta_{fi}+iT_{fi}=\delta_{fi}+i(2\pi)^4\delta^{(4)}(p_f-p_i)M_{fi} \tag{4.4.1}$$

M_{fi} is the invariant matrix element, used to calculate cross section etc..

Zeroth term in the expansion of S

$$\begin{aligned}\langle p_1p_2|p_Ap_B\rangle &= \sqrt{2E_12E_22E_A2E_B}\,\langle 0|a_1a_2a_A^+a_B^+|0\rangle\\ &= 2E_A2E_B(2\pi)^6\left\{\delta^{(3)}(\boldsymbol{p}_A-\boldsymbol{p}_1)\delta^{(3)}(\boldsymbol{p}_B-\boldsymbol{p}_2)+\delta^{(3)}(\boldsymbol{p}_A-\boldsymbol{p}_2)\delta^{(3)}(\boldsymbol{p}_B-\boldsymbol{p}_1)\right\}\end{aligned}$$

This actually is "no scattering", part of the $\mathbb{1}$ in the S-matrix.

First term is

$$\begin{aligned}&\langle p_1p_2|T\left(-\frac{i\lambda}{4!}\int \mathrm{d}^4x\phi^4(x)\right)|p_Ap_B\rangle\\ &\overset{\text{wick}}{=}\langle p_1p_2|:\left(-\frac{i\lambda}{4!}\int \mathrm{d}^4x\phi^4(x)+\text{contractions}\right):|p_Ap_B\rangle\end{aligned}$$

However now the expectation value of a normal-ordered expression doesn't vanish!

$$\begin{aligned}\phi^+(x)\,|\boldsymbol{p}\rangle &= \int\frac{\mathrm{d}^3k}{(2\pi)^3\sqrt{2E_k}}a_{\boldsymbol{k}}e^{-ikx}\sqrt{2E_p}a^+_{\boldsymbol{p}}\,|0\rangle\\ &= \int\frac{\mathrm{d}^3k}{(2\pi)^3\sqrt{2E_k}}e^{-ikx}\sqrt{2E_p}\delta^{(3)}(\boldsymbol{k}-\boldsymbol{p})\,|0\rangle\\ &= e^{-ipx}\,|0\rangle\end{aligned}$$

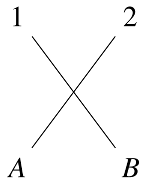
So in general, need two field operators to annihilate the in-state and m fields operators to create the out-states.

New type of Feynman diagram to deal with external states. Define contractions of field operators with external states according to

$$\begin{aligned}\overbracket{\phi(x)\,|\boldsymbol{p}}\rangle &= e^{-ipx}\,|0\rangle\\ \langle\overbracket{\boldsymbol{p}|\,\phi}(x) &= e^{+ipx}\,|0\rangle\end{aligned}$$

How does this work for $p_Ap_B\to p_1p_2$ in ϕ^4 at $O(\lambda)$? The above contains 3 types of terms: $\phi\phi\phi\phi$, $\overbracket{\phi\phi}\phi\phi$ and $\overbracket{\phi\phi}\,\overbracket{\phi\phi}$.

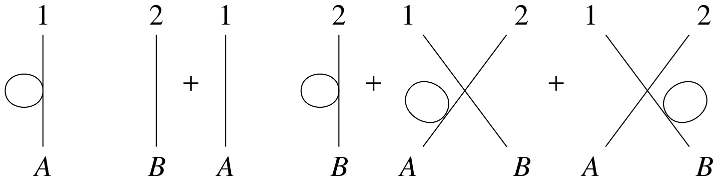
1. $\phi\phi\phi\phi$ allows full contractions with all external states. There is 4! possiblities



$$=4!\frac{-i\lambda}{4!}\int \mathrm{d}^4xe^{-i(p_A+p_B-p_1-p_2)x}=-i\lambda\underbrace{(2\pi)^4\delta^{(4)}(p_A+p_B-p_1-p_2)}_{\text{Prefactor in definition of }i\mathcal{M}}$$

$i\mathcal{M}$ receives a contributuon $-i\lambda$!

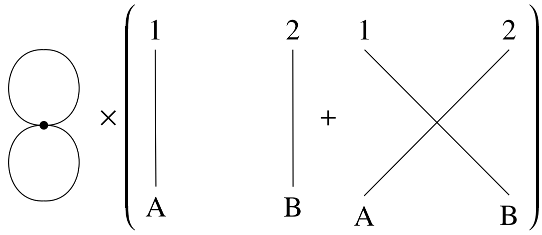
2. $\phi\phi\phi\phi$ leaves 2 operators to connect to external particles. Momentum conservation at each vertex. Still trivial!



Only fully connected Feynman diagrams contribute to $iT/i\mathcal{M}$!

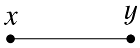
3.

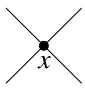
$$-\frac{i\lambda}{4!}\int \mathrm{d}^4x\,\langle p_1p_2|\phi\phi\phi\phi|p_Ap_B\rangle$$

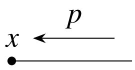
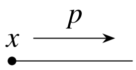
$=$ 

4.4.1 Feynman rules (with external lines)

Position space calculate iT by summing overall fully connected diagrams with

- propagator  $= D_F(x-y)$

- vertex  $= -i\lambda\int \mathrm{d}^4x$

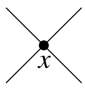
- external lines "in"  $= e^{-ip\cdot x}$;  $= e^{ip\cdot x}$

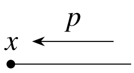
- divide diagram by its symmetry factor $\frac{1}{S}$

Momentum space We have seen it before. Now (with external lines) all positions are integrated over. Result is a function of external momenta only. Integrating out all momentum-conserving δ-distribution yields <u>overall</u> momentum conservation: $(2\pi)^4\delta^{(4)}(P_f-P_i)$

Momentum space Feynman rules for calculating iM:

- internal propagator  $= \frac{i}{p^2-M^2+i\epsilon}$

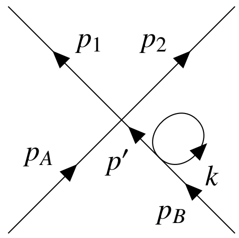
- vertex  $= -i\lambda$

- external lines ("in" or "out")  $= 1$

- impose 4-momentum conservation at each vertex

- integrate over all undetermined momenta $\int \frac{d^4p}{(2\pi)^4}$
- divide diagram by its symmetry factor $\frac{1}{S}$

There is still trouble in there. Consider the next-to-leading contribution to the scattering amplitude

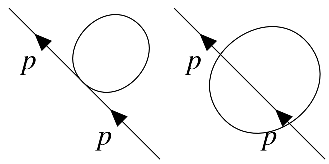


$$=\frac{1}{2}\int\frac{d^4p'}{(2\pi)^4}\frac{i}{p'^2-m^2}\int\frac{d^4k}{(2\pi)^4}\frac{i}{k^2-m^2}$$
$$\times(-i\lambda)(2\pi)^4\delta^{(4)}(p_A+p'-p_1-p_2)\cdot(-i\lambda)(2\pi)^4\delta^{(4)}(p_B-p')$$

This contains the internal propagator $\frac{i}{P_B^2-m^2+i\epsilon}$, but all the external particle are on their mass-shell, i.e.

$$P_A^2=P_B^2=P_1^2=P_2^2=m^2 \;\Rightarrow\; \frac{i}{P_B^2-m^2}=\frac{i}{0}$$

In Addition to having fully connected diagrams, also need to confine ourselves to amputated diagrams: disregard all these diagrams with loops attached to external legs.



These diagrams represent the transition from the free to the interacting asymptotic states.

Lehmann-Symanzik-Zimmermann (LSZ) reduction formula

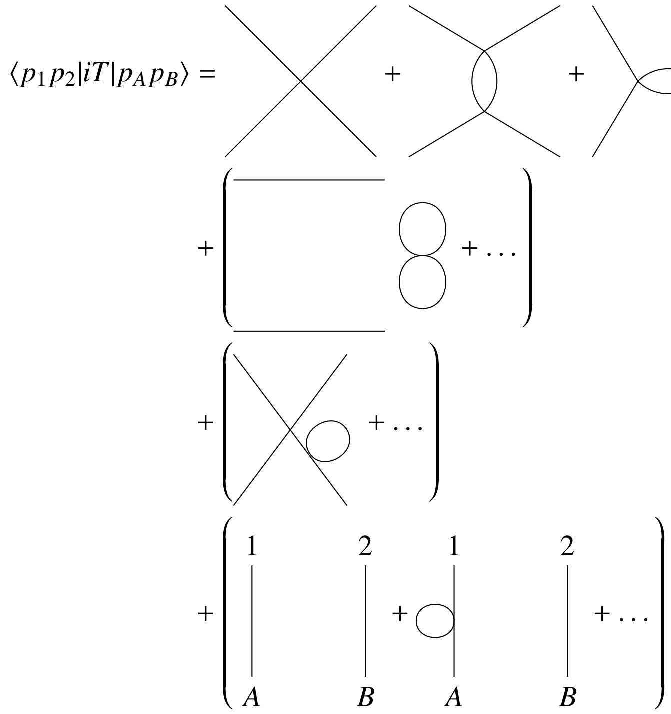
Proof on relation between correlation functions and S-matrix elements will be provided later.

$$\prod_{i=1}^{n}\int d^4x_i e^{ip_i\cdot x_i}\prod_{j=1}^{m}\int d^4y_j e^{-ik_j\cdot y_i}\langle\Omega|T\phi(x_1)\dots\phi(x_n)\phi(y_1)\dots\phi(y_m)|\Omega\rangle$$
$$\overset{\text{LSZ}}{=}(\text{disconnected stuff})+\underbrace{\prod_{i=1}^{n}\frac{\sqrt{z}i}{p_i^2-m^2+i\epsilon}\prod_{j=1}^{m}\frac{\sqrt{z}i}{k_j^2-m^2+i\epsilon}}_{\text{remove poles from external legs}}\langle p_1\dots p_n|S|k_i\dots k_m\rangle \tag{4.4.2}$$

z is the wave-function renormalization factor.

Then amend feynman rules above

consider only fully connected, amputated diagrams

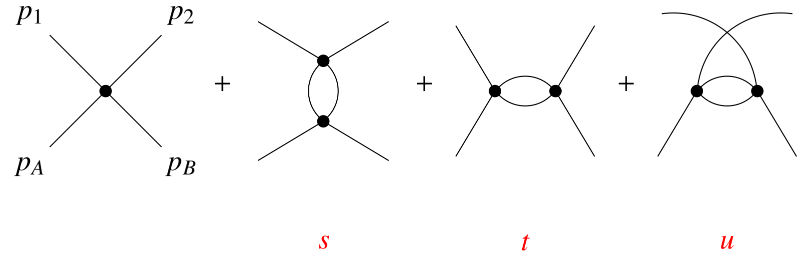


yields $|0\rangle \to |\Omega\rangle$

yields $|p\rangle_{\text{free}} \to |p\rangle_{\text{int}}$

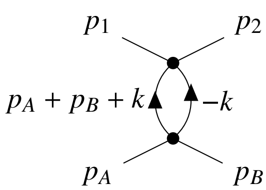
yields $\mathbb{1}$ in S-matrix

All allowed scattering diagrams $2 \to 2$ in ϕ^4 up to $O(\lambda^2)$:



Define the Lorentz-invariant quantities, *Mandelstam variables*:

$$s = (p_A + p_B)^2\,, \quad t = (p_A - p_1)^2\,, \quad u = (p_A - p_2)^2 \tag{4.4.3}$$



$$= \frac{1}{2}(-i\lambda)^2 \int \frac{\mathrm{d}^4 k}{(2\pi)^4} \frac{i}{k^2 - m^2 + i\epsilon} \frac{i}{(p_A + p_B + k)^2 - m^2 + i\epsilon} =: \frac{1}{2}(-i\lambda)^2 iJ(s)$$

Then the complete invariant amplitude is

$$M = -\lambda - \frac{\lambda^2}{2}\left(J(s) + J(t) + J(u)\right) \tag{4.4.4}$$

4.5 Scattering cross section

This section is based on Itzykson & Zuber, Chapter 5.1.
The aim is to relate (differential) cross section to reduced/invariant matrix element M_{fi}. First we describe the initial states not as momentum eigenstates $|p_A p_B\rangle$, but as wave packets.

$$|i\rangle = \int \frac{\mathrm{d}^3 k_A}{(2\pi)^3 2k_A^0} \frac{\mathrm{d}^3 k_B}{(2\pi)^3 2k_B^0} f(k_A) g(k_B) \, |k_A k_B\rangle$$

with $f(k_A)$, $g(k_B)$ strongly peaked at $k_A \approx p_A$, $k_B \approx p_B$.

We can write the transition amplitude to the final state $|f\rangle \propto |p_1 p_2\rangle$ (note: normalisation not the same)

$$\begin{aligned} A_{fi} &= \int \frac{\mathrm{d}^3 k_A}{(2\pi)^3 2k_A^0} \frac{\mathrm{d}^3 k_B}{(2\pi)^3 2k_B^0} f(k_A) g(k_B) \, \langle f | iT | k_A k_B \rangle \\ &= \int \frac{\mathrm{d}^3 k_A}{(2\pi)^3 2k_A^0} \frac{\mathrm{d}^3 k_B}{(2\pi)^3 2k_B^0} f(k_A) g(k_B) (2\pi)^4 \delta^{(4)}(\underbrace{p_f}_{=p_1+p_2} - k_A - k_B) iM(f, k_A, k_B) \end{aligned}$$

Thus the <u>transition probablity</u>:

$$\begin{aligned} \omega_{fi} = (2\pi)^8 &\int \frac{\mathrm{d}^3 k_A}{(2\pi)^3 2k_A^0} \frac{\mathrm{d}^3 k_B}{(2\pi)^3 2k_B^0} \frac{\mathrm{d}^3 q_A}{(2\pi)^3 2q_A^0} \frac{\mathrm{d}^3 q_B}{(2\pi)^3 2q_B^0} f(k_A) g(k_B) f(q_A)^* g(q_B)^* \\ \times \quad & \underbrace{\delta^{(4)}(p_f - k_A - k_B) \delta^{(4)}(p_f - q_A - q_B)}_{=\delta^{(4)}(q_A + q_B - k_A - k_B)\delta^{(4)}(p_f - p_A - p_B)} \underbrace{M(f, k_A, k_B) M^*(f, q_A, q_B)}_{\approx |M(f, p_A, p_B)|^2} \end{aligned}$$

Using the fourier representation of delta function $\delta^{(4)}(q_A + q_B - k_A - k_B) = (2\pi)^{-4} \int \mathrm{d}^4 x e^{i(k_A + k_B - q_A - q_B)\cdot x}$

$$\begin{aligned} &= \int \mathrm{d}^4 x \underbrace{\int \frac{\mathrm{d}^3 k_A}{(2\pi)^3 2k_A^0} \frac{\mathrm{d}^3 q_A}{(2\pi)^3 2q_A^0} e^{i(k_A - q_A)\cdot x} f(k_A) f^*(q_A)}_{:=|\tilde{f}(x)|^2} \\ &\times \quad \underbrace{\int \frac{\mathrm{d}^3 k_B}{(2\pi)^3 2k_B^0} \frac{\mathrm{d}^3 q_B}{(2\pi)^3 2q_B^0} e^{i(k_B - q_B)\cdot x} g(k_B) g^*(q_B)}_{:=|\tilde{g}(x)|^2} (2\pi)^4 \delta^{(4)}(p_f - p_A - p_B) \cdot |M(f, p_A, p_B)|^2 \end{aligned}$$

Using Fourier transformation $\tilde{g}(x) := \int \frac{\mathrm{d}^3 q}{(2\pi)^3 2q^0} e^{iq\cdot x} g(q)$

$$= \int \mathrm{d}^4 x |\tilde{f}(x)|^2 |\tilde{g}(x)|^2 (2\pi)^4 \delta^{(4)}(p_f - p_A - p_B) \cdot |M(f, p_A, p_B)|^2$$

note that $M(f, p_A, p_B)$ and $M(p_1, p_2, p_A, p_B)$ have different normalisation.

We now consider transition probabilty per unit volume per unit time:

$$\frac{\mathrm{d}\omega_{fi}}{\mathrm{d}V \mathrm{d}t} = (\text{incident flux}) \cdot (\text{target density}) \cdot \mathrm{d}\sigma$$

with $\mathrm{d}\sigma$ the infinitismal cross section for scattering into final state $\langle f|$.

Product (incident flux) · (target density) denotes overlap of wave function. Necessary condition!

Covariant renormalization of states $\langle \boldsymbol{p}|\boldsymbol{q}\rangle \sim 2p^0\delta^3(\boldsymbol{p}-\boldsymbol{q})$ means the number of particles per unit volume is $2p_A^0|\tilde{f}(x)|^2$ and $2p_B^0|\tilde{g}(x)|^2$, respectively.

Assume

$$\bullet \xrightarrow{p_A} \bullet\ \boldsymbol{p}_B = 0$$

in target rest frame. Then $2p_B^0 = 2m_B$ and target density $= 2m_B|\tilde{g}(x)|^2$

Incident flux $= |\boldsymbol{v}_A| \cdot 2p_A^0|\tilde{f}(x)|^2 = 2|\boldsymbol{p}_A||\tilde{f}(x)|^2$ since $|\boldsymbol{v}_A| = |\boldsymbol{p}_A|/p_A^0$. Then

$$\mathrm{d}\sigma = (2\pi)^4\delta^{(4)}(p_f - p_A - p_B)\frac{1}{4m_B|\boldsymbol{p_A}|}|M(f, p_A, p_B)|^2$$

for $A + B \to 1 + 2$ processes

$$= \int_\Delta \frac{\mathrm{d}^3p_1}{(2\pi)^3 2p_1^0}\frac{\mathrm{d}^3p_2}{(2\pi)^3 2p_2^0}(2\pi)^4\delta^{(4)}(p_1 + p_2 - p_A - p_B)\frac{1}{4m_B|\boldsymbol{p_A}|}|M(p_1, p_2, p_A, p_B)|^2$$

with Δ energy-momentum resolution of 4-momentum of final state $|f\rangle$.

Covariant form of

$$m_B \cdot |\boldsymbol{p}_A| = m_B\sqrt{(p_A^0)^2 - m_A^2} = \sqrt{(p_A \cdot p_B)^2 - m_A^2 m_B^2} =: F \tag{4.5.1}$$

This is scattering into arbitary final state subject to 4-momentum conservation: $p_A + p_B = p_1 + p_2$.

Consider now <u>differential</u> cross section for scattering into a particular infinitismal solid angle $\mathrm{d}\Omega$, hence specific momentum $\mathrm{d}p_1$, $\mathrm{d}p_2$ variations:

$$\begin{aligned}
\mathrm{d}\sigma &= \frac{1}{4F}\prod_f \frac{\mathrm{d}^3p_f}{(2\pi)^3 2p_f^0}(2\pi)^4\delta^{(4)}(p_A + p_B - \sum_f p_f)|M|^2\\
&\overset{f=1,2}{=} \frac{1}{4F}\frac{\mathrm{d}^3p_1}{(2\pi)^3 2p_1^0}\frac{\mathrm{d}^3p_2}{(2\pi)^3 2p_2^0}(2\pi)^4\delta^{(4)}(p_i - p_f)|M|^2\\
&= \frac{1}{64\pi^2 F}\frac{\mathrm{d}^3p_1}{E_1}\frac{\mathrm{d}^3p_2}{E_2}\delta^{(4)}(p_1 + p_2 - p_i)|M|^2
\end{aligned}$$

$$\boxed{\begin{aligned}
&\textstyle\int \frac{\mathrm{d}^3p_1}{E_1}\frac{\mathrm{d}^3p_2}{E_2}\delta^{(4)}(p_1 + p_2 - p_i)\\
&\overset{\text{CMS}}{=} \textstyle\int \mathrm{d}|\boldsymbol{p}_1|\mathrm{d}\Omega_1\frac{|\boldsymbol{p}_1|^2}{E_1E_2}\delta(E_1 + E_2 - E_i)\\
&= \textstyle\int \mathrm{d}(E_1 + E_2)\frac{\mathrm{d}|\boldsymbol{p}_1|}{\mathrm{d}(E_1+E_2)}\mathrm{d}\Omega_1\frac{|\boldsymbol{p}_1|^2}{E_1E_2}\delta(E_1 + E_2 - E_i)\\
&= \textstyle\frac{|\boldsymbol{p}_1|^2}{E_1E_2}\left(\frac{|\boldsymbol{p}_1|}{E_1} + \frac{|\boldsymbol{p}_1|}{E_2}\right)^{-1}\mathrm{d}\Omega_1\\
&= \textstyle\frac{|\boldsymbol{p}_1|\mathrm{d}\Omega_1}{E_1+E_2} = \frac{|\boldsymbol{p}_1|\mathrm{d}\Omega_1}{E_i}
\end{aligned}}$$

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\Omega} = \frac{1}{64\pi^2}\frac{|\boldsymbol{p}_1|}{F \cdot E_i}|M|^2 \tag{4.5.2}$$

Rewrite all kinematical factors in terms of $s = (p_A + p_B)^2 = (p_1 + p_2)^2$. Define the function

$$\lambda(x, y, z) := x^2 + y^2 + z^2 - 2(xy + xz + yz) \tag{4.5.3}$$

then

$$F = \sqrt{(p_A\cdot p_B)^2 - m_A^2 m_B^2} = \frac{1}{2}\lambda^{\frac{1}{2}}(s, m_A^2, m_B^2) = \sqrt{s}|\boldsymbol{p}_i|$$

$$\boxed{\begin{aligned} \lambda(s, m_A^2, m_B^2) &= s^2 - 2s(m_A^2 + m_B^2) - (m_A^2 - m_B^2)^2 = (s - (m_A + m_B)^2)(s - (m_A - m_B)^2) \\ &= (2p_A\cdot p_B - 2m_A\cdot m_B)\cdot(2p_A\cdot p_B + 2m_A\cdot m_B) = 4\left[(p_A p_B)^2 - m_A^2 m_B^2\right] \\ &\qquad \boxed{\begin{array}{ll} p_A = (c\sqrt{s}, \boldsymbol{p}_i),\ c\in[0,1] & \to m_A^2 = c^2 s - |\boldsymbol{p}_i|^2 \\ p_B = ((1-c)\sqrt{s}, -\boldsymbol{p}_i) & \to m_B^2 = (1-c)^2 s - |\boldsymbol{p}_i|^2 \end{array}} \\ &= 4\left[\left((c(1-c)s + p_i^2\right)^2 + (c^2 s - p_i^2)((1-c)^2 s - p_i^2)\right] = 4s|\boldsymbol{p}_i|^2 \end{aligned}}$$

$$|\boldsymbol{p}_f| = \sqrt{E_{1,2}^2 - m_{1,2}^2} = \frac{1}{2\sqrt{s}}\lambda^{\frac{1}{2}}(s, m_1^2, m_2^2)$$

$$E_i = \sqrt{s}$$

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\Omega_{CMS}} = \frac{1}{64\pi^2 s}\frac{|\boldsymbol{p}_f|}{|\boldsymbol{p}_i|}|M|^2 = \frac{1}{64\pi^2 s}\sqrt{\frac{\lambda(s, m_1^2, m_2^2)}{\lambda(s, m_A^2, m_A^2)}}|M|^2 \tag{4.5.4}$$

Decay rate instead of cross section means no "incident flux" to divide by, only "target density"

$$\mathrm{d}\Gamma = \frac{1}{2m_A}\prod_f \frac{\mathrm{d}^3 p_f}{(2\pi)^3 2p_f^0}(2\pi)^4\delta^{(4)}(p_A - \sum_f p_f)|M|^2 \tag{4.5.5}$$

Particles with spin (unpolarized): sum over outgoing or average over initial spins

$$|M|^2 \to \frac{1}{(2s_A+1)(2s_B+1)}\sum_{s_i,s_f}|M_{fi}|^2 \tag{4.5.6}$$

Symmetry factor $|M|^2 \to \frac{1}{s}|M|^2$ with $s = \prod_i k_i!$ if there are k_i identical particles of species i in the final states.

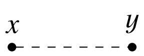
If 1 and 2 are identical, then facotr $\frac{1}{s} = \frac{1}{2}$ on the right hand side.

4.6 Feynman rules for fermions

Consider the simplest interacting theory with fermions, Yukawa-theory. We will treat QED later.

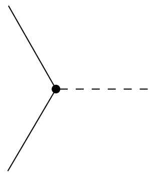
$$\mathcal{L} = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{M^2}{2}\phi^2 + \bar{\psi}(i\partial\!\!\!/ - m)\psi - g\bar{\psi}\psi\phi \tag{4.6.1}$$

Feynman rules will involve:

- scalar 

$= D_F(x-y) = \int \frac{\mathrm{d}^4}{(2\pi)^2}\frac{i}{p^2 - M^2 + i\epsilon}e^{-ip(x-y)}$
- fermions 

$= S_F(x-y)_{\alpha\beta} = \int \frac{\mathrm{d}^4}{(2\pi)^4}\frac{i(p\!\!\!/ + m)}{p^2 - m^2 + i\epsilon}e^{-ip(x-y)}$

- vertices  $= -ig\int d^4x$

What previous steps need reconsideration due to the anticommutating fermion operators? Interaction Hamiltonina $\sim \bar{\psi}\psi\phi$ and in general compose of even number of fermion fields (spin conservation and fermion number conservation). Thus there is no problem with time-ordered exponential in definition of S-matrix. (Time ordering always takes two or even number of fields.)

Remember the relation

$$T(\psi_\alpha(x)\bar{\psi}_\beta(x)) = -\bar{\psi}_\beta(x)\psi_\alpha(x) \text{ when } y^0 > x^0 \tag{4.6.2}$$

Similarly in normal product:

$$:\psi^+\psi^- = -\psi^-\psi^+: \tag{4.6.3}$$

Then Wick's theorem is formally the same as before

$$T(\psi_\alpha(x)\bar{\psi}_\beta(x)) =: \psi_\alpha(x)\bar{\psi}_\beta(x) : + \psi_\alpha(x)\bar{\psi}_\beta(x)$$

note by definition $\psi\psi = \bar{\psi}\bar{\psi} = 0$

Thus contractions inside normal-ordered products would be

$$:\psi_1\psi_2\bar{\psi}_3\bar{\psi}_4: = -\psi_1\bar{\psi}_3 : \psi_2\bar{\psi}_4 : = -S_F(x_1 - x_3) : \psi_2\bar{\psi}_4 :$$

because of the additional operator exchange.

We will want to consider fermion-(anti-)fermion scattering. Leading contribution at $O(g^2)$:

$$\frac{1}{2!}(-ig)^2 \int d^4x d^4y \langle p', k'|T\bar{\phi}(x)\phi(x)\phi(x)\bar{\phi}(y)\phi(y)|p, k\rangle$$

Contractions with initial-/final-state fermions?

$$\begin{aligned}\phi^+(x)\,|p, s\rangle &= \int \frac{d^3k}{(2\pi)^3\sqrt{2E_k}} \sum_r a_k^r u_r(k) e^{-ik\cdot x} \sqrt{2E_p} a_p^{s\dagger}\,|0\rangle \\ &= e^{-ip\cdot x} u_s(p)\,|0\rangle\end{aligned}$$

So define

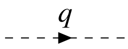
$$\begin{aligned}\psi(x)\,|p, s\rangle &= e^{-ip\cdot x} u_s(p) \\ \langle p, s|\,\bar{\psi}(x) &= e^{ip\cdot x} \bar{u}_s(p)\end{aligned} \tag{4.6.4}$$

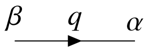
note, though, for antifermion states $|p', s'\rangle$:

$$\begin{aligned}\bar{\psi}(x)\,|p, s\rangle &= e^{-ip'\cdot x} \bar{v}_{s'}(p') \\ \langle p', s'|\,\psi(x) &= e^{ip'\cdot x} v_{s'}(p')\end{aligned} \tag{4.6.5}$$

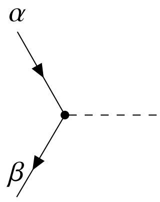
In short $\psi|\,\rangle$ contractes with a fermion, $\langle|\psi$ with an antifermion; vice verse for $\bar{\psi}$.

Momentum space feynman rule for iM

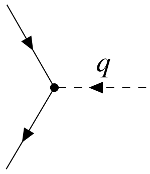
- internal propagators 

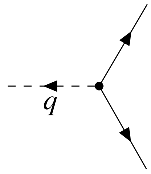
$= \frac{i}{q^2-M^2+i\epsilon}$; 

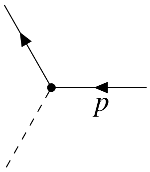
$= \frac{i}{q^2-M^2+i\epsilon} = \frac{i(\not{p}+m)_{\alpha\beta}}{p^2-m^2+i\epsilon}$

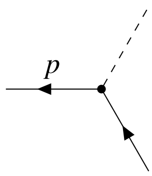
- vertex 

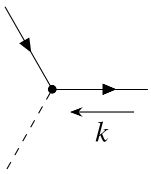
$= -ig\int \mathrm{d}^4x = ig\delta_{\beta\alpha}$

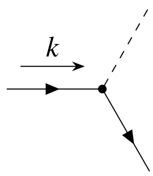
- external lines:
  - "in" and "out" scalar 

$=$ 

$= 1$
  - "in"-fermion 

$= u_s(p)$ and "out" fermion 

$= \bar{u}_s(p)$
  - "in"-antifermion 

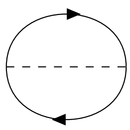
$= \bar{v}_s(k)$ and "out" antifermion 

$= v_s(k)$

- impose energy-momentum conservation at each vertex
- integrate over undetermined (loop) momenta
- include an overall sign for the diagram

note

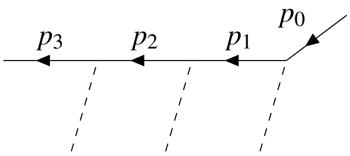
- Arrowas on the fermion lines by convention denote fermion (or charge) flow. They must flow consistently through the diagram. ($\equiv$ fermion number conservation) (Only potential confusion: external antifermion lines)

- No symmetry factors (except vacuum bubbles  $\frac{1}{s} = \frac{1}{2}$). $\bar{\psi}\psi\phi$ allows for unambiguous contractions.

- Dirac indices are summed over at each vertex

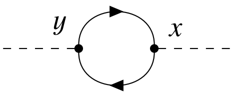
$$\mathcal{L}_{\text{int}} \approx \bar{\psi}_\alpha(x)\psi_\alpha(x)\phi(x)$$

$(\not{p}+m)$ terms in propagator are matrix-multiplied contracted with external spinors, e.g.



$$\sim \bar{u}_\alpha(p_3)\frac{i(\not{p}+m)_{\alpha\beta}}{p_2^2-m^2+i\epsilon}\frac{i(\not{p}+m)_{\beta\gamma}}{p_1^2-m^2+i\epsilon}u_\gamma(p_0)$$

- closed fermion loop



$$\sim \bar{\psi}_\alpha(x)\psi_\alpha(x)\bar{\psi}_\beta(y)\psi_\beta(y) = -\psi_\alpha(x)\bar{\psi}_\beta(y)\psi_\alpha(x)\bar{\psi}_\beta(y)$$

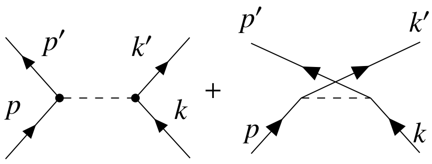
$$= -S_F(y-x)_{\beta\alpha}S_F(x-y)_{\alpha\beta} = -\mathrm{Tr}\left(S_F(y-x)S_F(x-y)\right)$$

It always (also with more propagators/couplings) involves an overal (-1) and a trace $\mathrm{Tr}(\dots)$.

Examples

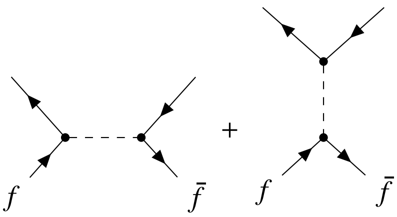
- fermion-fermion scattering to lowest order $O(g^2)$

$iM =$



$$= (-ig)^2\left\{\bar{u}(p')u(p)\frac{i}{\underbrace{(p'-p)^2}_{t}-M^2+i\epsilon}\bar{u}(k')u(k)-\bar{u}(p')u(k)\frac{i}{\underbrace{(p'-k)^2}_{u}-M^2+i\epsilon}\bar{u}(k')u(p)\right\}$$

- fermion-antifermion scattering



These are tree diagrams. Thus there is no undetermined momenta to integrate.

5 Quantum Electrodynamics (QED)

5.1 Classical Electrodynamics and Maxwell's equations

We have the gauge potential $A^\mu = (A^0, \boldsymbol{A}) = (\phi, \boldsymbol{A})$ & $A_\mu = (A^0, -\boldsymbol{A}) = (\phi, -\boldsymbol{A})$ and the field strength tensor $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$.

Then

- electric field $E_i = F_{0i} = \partial_0 A_i - \partial_i A_0 \to \boldsymbol{E} = -\dot{\boldsymbol{A}} - \boldsymbol{\nabla}\phi$
- magnetic field $B^i = -\frac{1}{2}\epsilon^{ijk}F_{jk} \to \boldsymbol{B} = \boldsymbol{\nabla} \times \boldsymbol{A}$

Lagrangian density $\mathcal{L}_{EM} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} = -\frac{1}{2}(\boldsymbol{E}\cdot\boldsymbol{E} - \boldsymbol{B}\cdot\boldsymbol{B})$. The field equation $\partial_\mu \left(\frac{\partial \mathcal{L}}{\partial(\partial_\mu A_\nu)}\right) - \frac{\partial \mathcal{L}}{\partial A_\nu} = 0$ leads to

$$\partial_\mu F^{\mu\nu} = 0 \tag{5.1.1}$$

it is half of Maxwell's equations (in vacuum).

The other half are Bianchi identities following from the definition of $F_{\mu\nu}$:

$$\begin{aligned} \partial_\lambda F_{\mu\nu} + \partial_\mu F_{\nu\lambda} + \partial_\nu F_{\lambda\mu} = 0 \Leftrightarrow \epsilon^{\sigma\lambda\mu\nu}\partial_\lambda F_{\mu\nu} = 0 \\ \text{or } \partial_\lambda \tilde{F}^{\sigma\lambda} = 0,\ \tilde{F}^{\sigma\lambda} = \frac{1}{2}\epsilon^{\sigma\lambda\mu\nu}F_{\mu\nu} \end{aligned}$$

In terms of $\boldsymbol{E}$ and $\boldsymbol{B}$:

$$\begin{aligned} &\boldsymbol{\nabla}\cdot\boldsymbol{E} = 0,\ \dot{\boldsymbol{E}} = \boldsymbol{\nabla}\times\boldsymbol{B} \quad \text{dynamical equations} \\ &\boldsymbol{\nabla}\cdot\boldsymbol{B} = 0,\ \dot{\boldsymbol{B}} = -\boldsymbol{\nabla}\times\boldsymbol{E} \quad \text{Bianchi identities} \end{aligned}$$

Remarks

- Lagrangian density does not depend on $\dot{A}_0$, since A_0 is not really dynamical.

  $$\boldsymbol{\nabla}\cdot\boldsymbol{E} = 0 \to \boldsymbol{\nabla}^2 A_0 + \boldsymbol{\nabla}\cdot\dot{\boldsymbol{A}} = 0$$

  Solve this Poisson equation for $A_0(\boldsymbol{x}, t) = \frac{1}{4\pi}\int d^3y \frac{\boldsymbol{\nabla}\cdot\dot{\boldsymbol{A}}(\boldsymbol{y},t)}{|\boldsymbol{y}-\boldsymbol{x}|}$. Thus A_0 is given in terms of the other components of A.

- gauge invariance: field strength tensor invariant under the transformation $A_\mu \longmapsto A_\mu - \partial_\mu \mathrm{X}$ due to commuting derivatives. This leads to gauge invariance of Maxwell equations.
  Choose X to satisfy $\partial_\mu\partial^\mu \mathrm{X} = \partial^2 \mathrm{X} = \partial_\mu A^\mu$ allows us to demand the condition (Lorenz condition)

  $$\partial_\mu A^\mu = 0 \tag{5.1.2}$$

  such that A_μ belongs to the "Lorenz gauge" and reduces the degrees of freedom from 4 to 3.

- Further freedom is eliminated by adding any X with $\partial^2 X = 0$, e.g. $\frac{\partial}{\partial t} X = A_0$. Then we get the Coulomb or radiation gauge

$$A_0 = 0, \ \nabla \cdot \boldsymbol{A} = 0 \tag{5.1.3}$$

Note: vice versa imposing $\nabla \cdot \boldsymbol{A} = 0$ first, yields $A_0 = 0$ (using Lorenz condition?).
In Coulomb gauge:

$$\begin{aligned}
&\boldsymbol{E} = -\dot{\boldsymbol{A}}. \ \boldsymbol{B} = \nabla \times \boldsymbol{A}, \ \nabla \times \boldsymbol{A} = 0 \\
&-\ddot{\boldsymbol{A}} = \dot{\boldsymbol{E}} \overset{\text{Maxwell}}{=} \nabla \times \boldsymbol{B} = \nabla \times (\nabla \times \boldsymbol{A}) = \nabla(\underbrace{\nabla \cdot \boldsymbol{A}}_{=0}) - \nabla^2 \boldsymbol{A} \\
&\Rightarrow \partial^2 \boldsymbol{A} = 0
\end{aligned}$$

This wave equation is massless KG equation for each spatial component.

Then the solutions are obvious: $\boldsymbol{A} = \boldsymbol{\epsilon} e^{-ik\cdot x}$ with $k^2 = 0$ and $\boldsymbol{\epsilon} \cdot \boldsymbol{k} = 0$. The polarization vector $\boldsymbol{\epsilon}$ is transverse to $\boldsymbol{k}$.

Can write the lagrangian in Coulmb gauge

$$\mathcal{L}_{\text{EM}} = \frac{1}{2} \dot{\boldsymbol{A}} \dot{\boldsymbol{A}} - \frac{1}{2} \boldsymbol{B} \cdot \boldsymbol{B}$$

Then the conjugate momentum to $\boldsymbol{A}$ is $\boldsymbol{\Pi} = \frac{\partial \mathcal{L}}{\partial \dot{A}} = \dot{\boldsymbol{A}} = -\boldsymbol{E}$. It has only 3 components, there is no conjugate momentum to A_0!. Because of Coulomb gauge $\boldsymbol{\Pi}$ is subject to the constraint $\nabla \cdot \boldsymbol{\Pi} = 0$

Hamiltonian

$$H_{\text{EM}} = \int \mathrm{d}^3 x \left(\frac{1}{2} \boldsymbol{\Pi} \cdot \boldsymbol{\Pi} + \frac{1}{2} \boldsymbol{B} \cdot \boldsymbol{B} \right)$$

5.2 Quantizing the Maxwell field

We would like to impose canonical commutation relations, à la

$$\begin{aligned}
&[A_i(\boldsymbol{x}), A_j(\boldsymbol{y})] = [\Pi_i(\boldsymbol{x}), \Pi_j(\boldsymbol{y})] = 0 \\
&[A_i(\boldsymbol{x}), \Pi_j(\boldsymbol{y})] = i\delta_{ij} \delta^{(3)}(\boldsymbol{x} - \boldsymbol{y})
\end{aligned}$$

However this cannot be true. Take either derivative of the last equation and it needs to vanish due to $\nabla \cdot \boldsymbol{A} = \nabla \cdot \boldsymbol{\Pi} = 0$. But

$$[\partial^i A_i(\boldsymbol{x}), \Pi_k(\boldsymbol{y})] = i\delta_{ij} \partial^i \delta^{(3)}(\boldsymbol{x} - \boldsymbol{y})$$

here the derivative is takev with respect to $\boldsymbol{x}$, i.e. $\partial^i = \frac{\partial}{\partial x_i}$.

Replace δ_{ij} by Δ_{ij}

$$\begin{aligned}
[\partial^i A_i(\boldsymbol{x}), \Pi_j(\boldsymbol{y})] &= i\Delta_{ij} \partial^i \frac{1}{(2\pi)^3} \int \mathrm{d}^3 k e^{i\boldsymbol{k}\cdot(\boldsymbol{x}-\boldsymbol{y})} \\
&= -\frac{1}{(2\pi)^3} \int \mathrm{d}^3 k (k^i \Delta_{ij}) e^{i\boldsymbol{k}\cdot(\boldsymbol{x}-\boldsymbol{y})} \overset{!}{=} 0
\end{aligned}$$

it works for $\Delta_{ij} = \delta_{ij} - \frac{k_i k_j}{\boldsymbol{k}^2}$ in momentum space or $\Delta_{ij} = \delta_{ij} - \nabla^{-2} \partial_i \partial_j$ in position space.

$$[A_i(\boldsymbol{x}), \Pi_j(\boldsymbol{y})] = i\left(\delta_{ij} - \boldsymbol{\nabla}^{-2}\partial_i\partial_j\right)\delta^{(3)}(\boldsymbol{x}-\boldsymbol{y}) \tag{5.2.1}$$

As before we have the mode expansion

$$\boldsymbol{A}(\boldsymbol{x}) = \int \frac{\mathrm{d}^3k}{(2\pi)^3\sqrt{2|\boldsymbol{k}|}}\left(\boldsymbol{a_k}e^{i\boldsymbol{k}\cdot\boldsymbol{x}} + \boldsymbol{a}_{\boldsymbol{k}}^{\dagger}e^{-i\boldsymbol{k}\cdot\boldsymbol{x}}\right)$$
$$\boldsymbol{\Pi}(\boldsymbol{x}) = \int \frac{\mathrm{d}^3k}{(2\pi)^3}(-i)\sqrt{\frac{|\boldsymbol{k}|}{2}}\left(\boldsymbol{a_k}e^{i\boldsymbol{k}\cdot\boldsymbol{x}} - \boldsymbol{a}_{\boldsymbol{k}}^{\dagger}e^{-i\boldsymbol{k}\cdot\boldsymbol{x}}\right)$$

with $\boldsymbol{k}\cdot\boldsymbol{a_k} = \boldsymbol{k}\cdot\boldsymbol{a}_{\boldsymbol{k}}^{\dagger} = 0$.

Introduce 2 orthogonal polarization vectors $\boldsymbol{\epsilon}^{(1)}(\boldsymbol{k})$ and $\boldsymbol{\epsilon}^{(2)}(\boldsymbol{k})$ for each $\boldsymbol{k}$.

$$\boldsymbol{a_k} = a_{\boldsymbol{k}}^{(1)}\boldsymbol{\epsilon}^{(1)} + a_{\boldsymbol{k}}^{(2)}\boldsymbol{\epsilon}^{(2)} = \sum_{\lambda=1}^{2} a_{\boldsymbol{k}}^{(\lambda)}\boldsymbol{\epsilon}^{(\lambda)}(\boldsymbol{k})$$
$$\text{with } \boldsymbol{k}\cdot\boldsymbol{\epsilon}^{(1)}(\boldsymbol{k}) = \boldsymbol{k}\cdot\boldsymbol{\epsilon}^{(2)}(\boldsymbol{k}) = 0,\ \boldsymbol{\epsilon}^{(\lambda)}\cdot\boldsymbol{\epsilon}^{(\lambda;)} = \delta_{\lambda\lambda'}$$

Creation and annihilation operator have the standard commutation relations

$$[a_{\boldsymbol{k}}^{(\lambda)}, a_{\boldsymbol{k}'}^{(\lambda')\dagger}] = (2\pi)^3\delta_{\lambda\lambda'}\delta^{(3)}(\boldsymbol{k}-\boldsymbol{k}') \tag{5.2.2}$$

and all other commutators vanish. Geometrically, still possible to write including the unphysical longitudinal components:

$$[\boldsymbol{a_k}, \boldsymbol{a_l}] = [\boldsymbol{a}_{\boldsymbol{k}}^{\dagger}, \boldsymbol{a}_{\boldsymbol{l}}^{\dagger}] = 0$$
$$[a_{\boldsymbol{k}}^{i}, a_{\boldsymbol{l}}^{j\dagger}] = (2\pi)^3\left(\delta^{ij} - \frac{k^ik^j}{\boldsymbol{k}^2}\right)\delta^{(3)}(\boldsymbol{k}-\boldsymbol{l})$$

$a_{\boldsymbol{k}}^{(\lambda)}$ and $a_{\boldsymbol{k}}^{(\lambda)\dagger}$ create and destroy photons of momentum $\boldsymbol{k}$, energy $|\boldsymbol{k}|$ and (electric) polarization along $\boldsymbol{\epsilon}^{(\lambda)}(\boldsymbol{k})$.

Next steps are analogout to KG theory.

Hamiltonian

$$H = \frac{1}{2}\int \mathrm{d}^3x\left(\boldsymbol{E}^2 + \boldsymbol{B}^2\right) = \frac{1}{2}\int \mathrm{d}^3x\left(\dot{\boldsymbol{A}}^2 + (\boldsymbol{\nabla}\times\boldsymbol{A})\cdot(\boldsymbol{\nabla}\times\boldsymbol{A})\right)$$

using identity $\boldsymbol{A}\cdot(\boldsymbol{B}\times\boldsymbol{C}) = \boldsymbol{B}\cdot(\boldsymbol{C}\times\boldsymbol{A})$

$$= \frac{1}{2}\int \mathrm{d}^3x\left(\dot{\boldsymbol{A}}^2 + \boldsymbol{A}\cdot\boldsymbol{\nabla}\times(\boldsymbol{\nabla}\times\boldsymbol{A})\right)$$

using the identity $\boldsymbol{\nabla}\times(\boldsymbol{\nabla}\times\boldsymbol{A}) = \boldsymbol{\nabla}(\boldsymbol{\nabla}\cdot\boldsymbol{A} - \boldsymbol{\nabla}^2\boldsymbol{A})$

$$= \frac{1}{2}\int \mathrm{d}^3x\left(\dot{\boldsymbol{A}}^2 - \boldsymbol{A}\cdot\boldsymbol{\nabla}^2\boldsymbol{A} + \boldsymbol{A}\cdot\boldsymbol{\nabla}(\boldsymbol{\nabla}\cdot\boldsymbol{A})\right)$$

using coulomb gauge condition

$$= \frac{1}{2}\int \mathrm{d}^3x\left(\dot{\boldsymbol{A}}^2 - \boldsymbol{A}\cdot\nabla^2\boldsymbol{A}\right)$$

the first term vanishes and use normal ordering

$$= \int \frac{\mathrm{d}^3k}{(2\pi)^3}|\boldsymbol{k}|\,\boldsymbol{a}_{\boldsymbol{k}}^{\dagger}\cdot\boldsymbol{a_k} = \sum_{\lambda=1}^{2}\int \frac{\mathrm{d}^3k}{(2\pi)^3}|\boldsymbol{k}|\,a_{\boldsymbol{k}}^{(\lambda\dagger)}a_{\boldsymbol{k}}^{\lambda}$$

Heisenberg field

$$\boldsymbol{A}(\boldsymbol{x},t) = \int \frac{\mathrm{d}^3k}{(2\pi)^3}\frac{1}{\sqrt{2|\boldsymbol{k}|}}\left(\boldsymbol{a}_{\boldsymbol{k}}e^{-ik\cdot x} + \boldsymbol{a}_{\boldsymbol{k}}^\dagger e^{ik\cdot x}\right)$$

Photon propagator

$$\langle 0|TA_i(x)A_j(y)|0\rangle =: D_{ij}^{\mathrm{tr}}(x-y) = \int \frac{\mathrm{d}^4k}{(2\pi)^4}\frac{i}{k^2+i\epsilon}\left(\delta_{ij} - \frac{k_ik_j}{|\boldsymbol{k}|^2}\right)e^{-ik\cdot(x-y)} \tag{5.2.3}$$

tr stands for transverse: photon polarization perpendicular to its momentum. This is NOT the final version of the photon propagator!

5.3 Inclusion of matter - QED

$$\mathcal{L}_{\mathrm{QED}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \bar{\psi}(i\not{D} - m)\psi \tag{5.3.1}$$

where $D_\mu = \partial_\mu + ieA_\mu$ is the (gauge) covariante derivative

$$= \mathcal{L}_{\mathrm{EM}} + \mathcal{L}_D - e\underbrace{\bar{\psi}\gamma^\mu\psi}_{j^\mu}A_\mu \tag{5.3.2}$$

Field equations would be

$$\partial_\mu F^{\mu\nu} = ej^\nu \qquad (i\not{D} - m)\psi = 0$$

where ej^ν is the electromagnetic 4-current.

Gauge invariance under the transformation

$$\begin{cases} \psi(x) \longmapsto \psi'(x) = e^{ie\chi(x)}\psi \\ A_\mu(x) \longmapsto A'_\mu(x) = A_\mu(x) - \partial_\mu\chi(x) \end{cases}$$

To check the consistence: cavariant derivative transforms like $D_\mu \longmapsto D'_\mu\psi'(x) = e^{ie\chi(x)}D_\mu\psi(x)$. Since the adjoint spinor transforms like $\bar{\psi}(x) \longmapsto \bar{\psi}'(x) = \bar{\psi}(x)e^{-ie\chi(x)}$, the Lagrangian and field equations are gauge invariant.

Again we choose Coulomb gauge $\boldsymbol{\nabla}\cdot\boldsymbol{A} = 0$, then equation for A^0:

$$\begin{aligned} \partial_i F^{i0} &= ej^0 \\ \Rightarrow -\boldsymbol{\nabla}^2 A^0 &= ej^0 = e\bar{\psi}\gamma^0\psi \\ &= e\bar{\psi}\gamma^0\psi = e\psi^\dagger\psi \\ &= e\rho(x) \\ A^0(\boldsymbol{x},t) &= e\int \mathrm{d}^3y\frac{\rho(\boldsymbol{y},t)}{4\pi\,|\boldsymbol{x}-\boldsymbol{y}|} \end{aligned} \tag{5.3.3}$$

We want to derive the interaction Hamiltonian. Note

$$\int \mathrm{d}^3x\frac{1}{2}\boldsymbol{E}^2 = \int \mathrm{d}^3x\frac{1}{2}\left(\dot{\boldsymbol{A}} + \boldsymbol{\nabla}A^0\right)^2$$

cross terms vanish after integration by parts due to $\nabla \cdot \dot{\boldsymbol{A}} = 0$

$$= \int \mathrm{d}^3 x \frac{1}{2}\left(\dot{\boldsymbol{A}}^2 + (\nabla A^0)^2\right)$$

$$= \int \mathrm{d}^3 x \frac{1}{2}\left(\dot{\boldsymbol{A}}^0 - A^0 \nabla^2 A^0\right)$$

$-e j^0 = -e\rho$

$$= \int \mathrm{d}^3 x \frac{1}{2}\dot{\boldsymbol{A}}^2 + \underbrace{\frac{e^2}{2}\int \mathrm{d}^3 x \mathrm{d}^3 y \frac{\rho(\boldsymbol{x})\rho(\boldsymbol{y})}{4\pi|\boldsymbol{x}-\boldsymbol{y}|}}_{=\frac{e^2}{2} j^0 A_0}$$

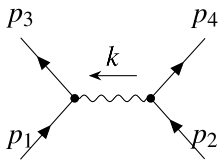
Combined Hamiltonian

$$H = \int \mathrm{d}^3 x \left\{\frac{1}{2}\boldsymbol{\Pi}\cdot\boldsymbol{\Pi} + \frac{1}{2}\boldsymbol{B}\cdot\boldsymbol{B} + i\bar{\psi}\boldsymbol{\gamma}\cdot\nabla\psi + m\bar{\psi}\psi\right\} \quad \text{free photon and fermion}$$
$$+ \frac{e^2}{2}\int \mathrm{d}^3 x \mathrm{d}^3 y \frac{\rho(\boldsymbol{x})\rho(\boldsymbol{y})}{4\pi|\boldsymbol{x}-\boldsymbol{y}|} - e\int \mathrm{d}^3 x \boldsymbol{j}\cdot\boldsymbol{A} \quad \text{interactions}$$

where $\rho = \psi^\dagger\psi = \bar{\psi}\gamma^0\psi$, $\boldsymbol{j} = \bar{\psi}\boldsymbol{\gamma}\psi$ for 2 types of interactions.

5.4 Lorentz-invariant propagator

Consider $e^- e^-$ scattering at $O(e^2)$



$$k = p_3 - p_1 = p_2 - p_4$$

We expect this to involve

- spinors for external fermions
- $-ie\gamma^\mu$
- Photon propagator $D_{\mu\nu}(x-y)$

What we have found in Coulomb gauge is actually

- vertices $ie\gamma^i$, transverse propagator $D^{\text{tr}}_{\mu\nu}(x-y)$
- vertices $\pm ie\gamma^0$, instantaneous Coumlomb interaction $\frac{1}{4\pi|\boldsymbol{x}-\boldsymbol{y}|}\delta(x^0-y^0)$

Effectively combine these propagators terms into $D^{\text{Coul}}_{\mu\nu}(x-y)$, where the $D^{\text{Coul}}_{00}(x-y) = \frac{1}{4\pi|\boldsymbol{x}-\boldsymbol{y}|}\delta(x^0-y^0)$. This component in momentum space is simply

$$\int \frac{\mathrm{d}^3 k}{(2\pi)^3}\frac{e^{i\boldsymbol{k}\cdot\boldsymbol{r}}}{|\boldsymbol{k}|^2} = \frac{1}{4\pi|\boldsymbol{r}|}$$

Therefore Coumlomb propagator in momentum space:

$$D_{\mu\nu}^{\text{Coul}}(k) = \begin{cases} \frac{i}{k^2+i\epsilon}\left(\delta_{ij} - \frac{k_i k_j}{|\boldsymbol{k}|^2}\right) & \mu = i, \nu = j \\ \frac{i}{|\boldsymbol{k}|^2} & \mu = \nu = 0 \\ 0 & \text{otherwise} \end{cases}$$

Consider contraction to scattering amplitude from vertex at x:

$$\sim e\bar{u}(p_3)\gamma^\mu u(p_1) e^{i(p_3-p_1)x}$$

current conservation $\partial_\mu j^\mu = 0$ written in momentum space

$$\underbrace{(p_3 - p_1)_\mu}_{k_\mu} \bar{u}(p_3)\gamma^\mu u(p_1) = 0$$

so in the complete diagram $D_{\mu\nu}^{\text{Coul}}$ occurs in a form

$$\begin{aligned} & a^\mu D_{\mu\nu}^{\text{Coul}}(k) b^\nu \\ =& a^0 \frac{i}{|\boldsymbol{k}|^2} b^0 + a^i \left[\frac{i}{k^2 + i\epsilon}\left(\delta_{ij} - \frac{k_i k_j}{|\boldsymbol{k}|^2}\right)\right] b^j \end{aligned}$$

where $k^\mu a^\mu = 0, k_\mu a^\mu = 0$

$$\begin{aligned} =& i\left[\frac{\boldsymbol{a}\cdot\boldsymbol{b}}{k^2} \underbrace{- \frac{k_0^2 a_0 b_0}{k^2|\boldsymbol{k}|^2} + \frac{a_0 b_0}{|\boldsymbol{k}|^2}}_{=\frac{-k_0^2 a_0 b_0 + a_0 b_0 (k_0^2 - |\boldsymbol{k}|^2)}{k^2|\boldsymbol{k}|}^2}\right] \\ =& \frac{i}{k^2}(\boldsymbol{a}\cdot\boldsymbol{b} - a_0 b_0) = -\frac{i}{k} a_\mu b^\mu \end{aligned}$$

Conclusion in this diagram (and in fact, in general), we may replace the $D_{\mu\nu}^{\text{Coul}}(k)$ by the manifestly Lorentz covariant propagator

$$D_{\mu\nu}(k) = -\frac{i g_{\mu\nu}}{k + i\epsilon} \tag{5.4.1}$$

This can be generalised to

$$D_{\mu\nu}(k) = -\frac{i}{k^2 + i\epsilon}\left(g_{\mu\nu} - (1-\alpha)\frac{k_\mu k_\nu}{k^2}\right) \tag{5.4.2}$$

as, by current consevation, additional term doesn't contribute.

Feynman gauge $\alpha = 1$; Landau gauge $\alpha = 0$.

Remark one can also try to quantise photons in a manifestly covariant way, imposing Lorentz gauge $\partial_\mu A^\mu = 0$

$$[A_\mu(\boldsymbol{x}), \Pi_\nu(\boldsymbol{y})] = i g_{\mu\nu} \delta^{(3)}(\boldsymbol{x} - \boldsymbol{y})$$

This is trouble since $\Pi^0 = \frac{\partial \mathcal{L}}{\partial \dot{A}_0} = 0$. This cannot hold!

We thus change the Lagrangian

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

with "gauge fixing term". The equation of motion from this is

$$\partial^2 A^\mu - \left(1 - \frac{1}{\alpha}\right)\partial^\mu(\partial_\lambda A^\lambda) = 0$$

e.g. $\alpha = 1$ is the Feynman gauge.

With this Lagrangian we can the 0th component of conjugate momentum

$$\Pi^0 = -\frac{1}{\alpha}\partial_\mu A^\mu$$

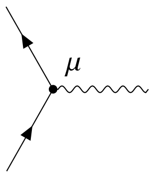
but this seems as bad as before!

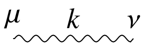
We cannot impose Coulomb gauge condition $\partial_\mu A^\mu = 0$ as an operator identity. Instead demand a weaker condition $\langle \text{out}|\partial_\mu A^\mu|\text{in}\rangle = 0$ for all physical states.

This in turn tells us which states are actually physical. The 4 polarisation states consist of physical, timelike(scalar) and longitudinal states. The negative-norm states cancel each ther out (Gopta-Bleuler formalism).

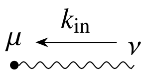
Feynman rules for QED diagrams constucted from electron (positron)  and photon ; rules for fermions are valid as before.

In addition

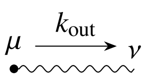
- vertex 

$= -ie\gamma^\mu$;
- photon propagator 

$= -\frac{ig_{\mu\nu}}{k^2+i\epsilon}$
- external photons



$= \epsilon_\mu$

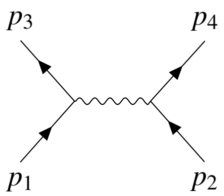


$= \epsilon_\nu^*$

ϵ_μ polarisation vector of in/out photon and ϵ_μ^* for out photon required for complex (circular) polarisation.

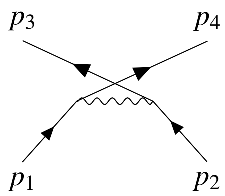
5.5 QED process at tree level

Example $e^-e^- \to e^-e^-$



$$\sim \frac{1}{(p_1 - p_3)^2} = \frac{1}{t}$$

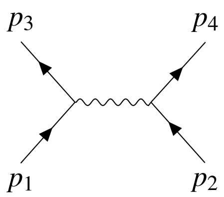
t-channel



$$\sim \frac{1}{(p_1 - p_4)^2} = \frac{1}{u}$$

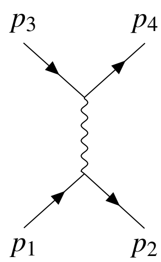
u-channel

Example $e^-e^+ \to e^-e^+$



$$\sim \frac{1}{(p_1 - p_3)^2} = \frac{1}{t}$$

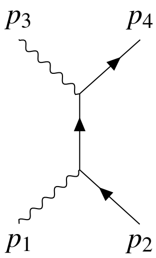
t-channel



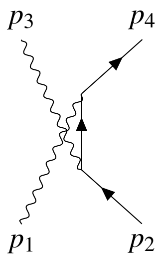
$$\sim \frac{1}{(p_1 + p_2)^2} = \frac{1}{s}$$

s-channel

Compton scattering $\gamma e^- \to \gamma e^-$

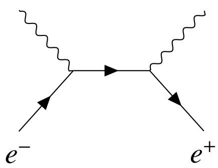


s-channel



u-channel

Example $e^+e^- \to \gamma\gamma$



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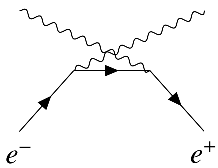
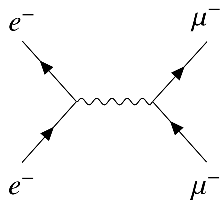


u-channel

These are important for lifetime of positronium.

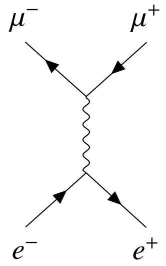
All these amplitudes are $O(e^2)$, $\alpha = \frac{e^2}{4\pi} = \frac{1}{137.036}$ the fine structure constant.
Muons $\mu^\pm$, like electrons, just ca. 200 times heavier.

$e^-\mu^- \to e^-\mu^- \Rightarrow$



Just one diagram

$e^+e^- \to \mu^+\mu^- \Rightarrow$



Just one diagram

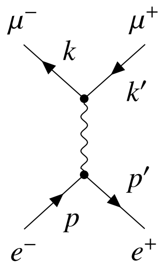
$\mu^\pm$ decay into $e^\pm$ and neutrinos in weak interactions.

For tree level diagrams the photon propagator does not need to have the $i\epsilon$ in the denominator, since we will never be able to see a singularity/pole.

5.5.1 Some hints and tricks for cross section calculations

See application in exercises!

Example $e^+e^- \to \mu^+\mu^-$



$$\begin{aligned} iM &= \bar{v}_e^{s'}(-ie\gamma^\mu)u_e^s(p) \left.\frac{-ig_{\mu\nu}}{s}\right|_{s=q^2} \bar{u}_\mu^r(k)(-ie\gamma^\nu)v_\mu^{r'}(k') \\ &= \frac{ie^2}{s}\left(\bar{v}_e(p')\gamma^\mu u_e(p)\right)\left(\bar{u}_\mu(k)\gamma_\mu v_\mu(k')\right) \end{aligned}$$

See section 4.5, $|M|^2$ is needed for cross section. M^* involves things like

$$\begin{aligned} (\bar{v}\gamma^\mu u)^* &= (\bar{v}\gamma^\mu u)^\dagger = u^\dagger\gamma^{\mu\dagger}\gamma_0^\dagger v \\ &= u^\dagger\gamma_0\gamma^\mu\gamma_0\gamma_0 v = \bar{u}\gamma^\mu v \end{aligned}$$

So

$$|M|^2 = \frac{e^4}{s^2}\left[\bar{v}(p')\gamma^\mu u(p)\bar{u}(p)\gamma^\nu v(p')\right]_{e^\pm} \cdot \left[\bar{u}(k)\gamma_\mu v(p)\bar{v}(k')\gamma_\nu u(k)\right]_{\mu^\pm}$$

Unpolarized scattering= $\frac{1}{4}\sum_{r,s,r',s'}|M|^2$.

Now $\bar{v}\gamma^\mu u$, $\bar{u}\gamma^\nu v$ etc. are scalars in Dirac/spinor space:

$$
\begin{aligned}
& \sum_{s,s'} \bar{v}_{s'} p' \gamma^\mu u_s(p) \bar{u}_s(p) \gamma^\nu v_{s'}(p') \\
\text{(taking trace of scalar)} \; &= \sum_{s,s'} \mathrm{Tr}\left(\bar{v}_{s'} p' \gamma^\mu u_s(p) \bar{u}_s(p) \gamma^\nu v_{s'}(p')\right) \\
&= \sum_{s,s'} \mathrm{Tr}\left(v_{s'}(p') \bar{v}_{s'}(p') \gamma^\mu u_s(p) \bar{u}_s(p) \gamma^\nu\right)
\end{aligned}
$$

using spin sums

$$
= \mathrm{Tr}\left((\not{p}' - m)\gamma^\mu (\not{p} + m)\gamma^\nu\right)
$$

Trace technology

- remember $\mathrm{Tr}\gamma_\mu = 0$
- $\mathrm{Tr}(\gamma^\mu \gamma^\nu) = 4g_{\mu\nu}$
- $\mathrm{Tr}(\text{odd number of } \gamma) = 0$
- $\mathrm{Tr}(\gamma_\mu \gamma_\nu \gamma_\alpha \gamma_\beta) = 4\left(g_{\mu\nu} g_{\alpha\beta} + g_{\mu\beta} g_{\nu\alpha} - g_{\mu\alpha} g_{\nu\beta}\right)$
- more rules involving γ_5 (weak interactions!)

So

$$
\mathrm{Tr}\left((\not{p}' - m)\gamma^\mu (\not{p} + m)\gamma^\nu\right) = 4\left(p'^\mu p^\nu + p'^\nu p^\mu - g^{\mu\nu}(p \cdot p' + m^2)\right)
$$

Mandelstam variables with 4 equal masses, center-of-mass system (CMS):

$$
p = (E, \boldsymbol{p}), \quad p' = (E, -\boldsymbol{p}), \quad k = (E, \boldsymbol{k}), \quad \theta = \measuredangle(\boldsymbol{p}, \boldsymbol{k})
$$

$$
s \overset{\text{CMS}}{=} (p + p')^2 = 4E^2 \tag{5.5.1}
$$

$$
t = (p - k)^2 = -(\boldsymbol{p} - \boldsymbol{k})^2 = -2|\boldsymbol{p}|^2(1 - \cos\theta) \tag{5.5.2}
$$

$$
u = (p' - k)^2 = -2|\boldsymbol{p}|^2(1 + \cos\theta) \tag{5.5.3}
$$

$$
|\boldsymbol{p}|^2 = E^2 - m^2 = \frac{s}{4} - m^2 \tag{5.5.4}
$$

Only 2 Mandelstam variables are independent.

$$
\begin{aligned}
s + t + u &= (p_1 + p_2)^2 + (p_1 - p_3)^2 + (p_1 - p_4)^2 \\
&= p_1^2 + p_2^2 + p_3^2 + p_4^2 + 2p_1 \underbrace{(p_1 + p_2 - p_3 - p_4)}_{=0} \\
&= \sum_i m_i^2 = \text{const}
\end{aligned}
$$

photon polarisation sums Analogy to fermion spin sums before, Feynman rules for external photons involve $\epsilon^(*)_\mu$; e.g. Compton amplitude of the form

$$M \sim \epsilon^*_\mu(p_3)\epsilon_\nu(p_1)T^{\mu\nu}$$

Thus

$$\sum_{\text{spin, pol.}} |M|^2 = \sum_{\text{spin, pol.}} \epsilon^*_\mu(p_3)\epsilon_\alpha(p_3)\epsilon^*_\beta(p_1)\epsilon_\nu(p_1)T^{\mu\nu}T^{\alpha\beta^*}$$

How can we simplify $\sum_{pol} \epsilon^*_\mu(k)\epsilon_\nu(k)$? Again we have only 2 physical polarisation states, but want to do it in a covariant form.

Assume a simpler process (than Compton) with a single external photon, $\epsilon^*_\mu(k)M^\mu$. Choose

$$k^\mu = (k,0,0,k), \quad \epsilon^\mu_{(1)} = (0,1,0,0), \quad \epsilon^\mu_{(2)} = (0,0,1,0)$$

so $\sum_{\text{pol}} |\epsilon^*_\mu(k)M^\mu|^2 = |M_1|^2 + |M_2|^2$

Remember that photon coupled source j^μ, current sonvervation $\partial_\mu j^\mu = 0$. We will see (next term) this holds in general as Ward identity

$$k_\mu M^\mu = 0 \tag{5.5.5}$$

In exercises, show $p_{3\mu}T^{\mu\nu} = 0 = p_{1\nu}T^{\mu\nu}$ for Compton

Here $kM^0 - kM^3 = 0 \Rightarrow M^0 = M^3$ and we can rewrite

$$\sum_{pol} \epsilon^*_\mu\epsilon_\nu M^\mu M^{*\nu} = |M_1|^2 + |M_2|^2 + \underbrace{|M_3|^2 - |M_0|^2}_{=0} = -g_{\mu\nu}M^\mu M^{*\nu}$$

so effectively

$$\sum_{pol} \epsilon^*_\mu(k)\epsilon_\nu(k) = -g_{\mu\nu} \tag{5.5.6}$$

side remark

- KG propagator $\frac{i}{p^2-M^2+i\epsilon}$
- Dirac propagator $\frac{i(\not{p}+m)}{p^2-m^2+i\epsilon} = \frac{i\sum_s u_s(p)\bar{u}_s(p)}{p^2-m^2+i\epsilon}$
- Photon propagator $\frac{-ig_{\mu\nu}}{p^2+i\epsilon} = \frac{i\sum_{pol}\epsilon^*_\mu(p)\epsilon_\nu(p)}{p^2+i\epsilon}$