

# Quantum Field Theory II

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\*<https://github.com/M-a-s-o/notes>

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## Lecture 1

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**Topics.** Generalize the formalism for complex scalar fields. Introduce fermions and their functional quantization. Discrete symmetries and PCT theorem. Gauge theories and their functional quantization. Anomalies. Possibly instantons, applications of RG flows, etc.

## 1 Complex scalar boson fields

[r] sources?

### 1.1 Generalization

Consider a complex scalar field  $\varphi$ . It has two degrees of freedom: either the fields  $\varphi$  and  $\varphi^*$  or the real and imaginary parts of the field  $\varphi$ . Both are useful for computation

$$\varphi = \frac{1}{\sqrt{2}}(\varphi_1 + i\varphi_2), \quad \varphi^* = \frac{1}{\sqrt{2}}(\varphi_1 - i\varphi_2)$$

The action of the  $\lambda\varphi^4$  theory is

$$\begin{aligned} S &= \int d^4x [\partial_\mu \varphi \partial^\mu \varphi^* - m^2 \varphi \varphi^* - \lambda(\varphi \varphi^*)^2] \\ &= \int d^4x \left[ \frac{1}{2} \partial_\mu \varphi_1 \partial^\mu \varphi_1 - \frac{1}{2} m^2 \varphi_1^2 + \frac{1}{2} \partial_\mu \varphi_2 \partial^\mu \varphi_2 - \frac{1}{2} m^2 \varphi_2^2 - \frac{\lambda}{4} (\varphi_1^2 + \varphi_2^2)^2 \right] \end{aligned}$$

Each degree of freedom has an equation of motion. Varying the action with respect to  $\varphi$  gives the equations of motion for  $\varphi^*$  and vice versa. For  $\lambda = 0$ , the fields components satisfy the Klein–Gordon equation.

Since the interaction term is real, it corresponds to a quartic vertex  $\varphi\varphi^*\varphi\varphi^*$ . For  $\lambda = 0$ , one can rewrite the action in terms of real  $\varphi_1$  and imaginary  $\varphi_2$  parts and see that the two degrees of freedom decouple since there is no mixed term  $\varphi_1^2\varphi_2^2$ .

With two degrees of freedom, the Euclidean generating functional depends on two source terms

$$\begin{aligned} W[J, J^*] &= \int [\mathcal{D}\varphi \mathcal{D}\varphi^*] \exp \left[ - \int d^4x [\mathcal{L}(\varphi, \varphi^*) - J\varphi - J^*\varphi^*] \right] \\ W[J_1, J_2] &= \int [\mathcal{D}\varphi_1 \mathcal{D}\varphi_2] \exp \left[ - \int d^4x [\mathcal{L}(\varphi_1, \varphi_2) - J_1\varphi_1 - J_2\varphi_2] \right] \end{aligned}$$

The Euclidean Green's functions are

$$\begin{aligned} \langle 0 | \mathcal{T} \{ \varphi(x_1) \cdots \varphi(x_l) \varphi^*(x_{l+1}) \cdots \varphi^*(x_n) \} | 0 \rangle &= \\ &= \frac{\delta^n W[J, J^*]}{\delta J(x_1) \cdots \delta J(x_l) \delta J^*(x_{l+1}) \cdots \delta J^*(x_n)} \Big|_{J=J^*=0} \end{aligned}$$

Similarly

$$\begin{aligned} \langle 0 | \mathcal{T} \{ \varphi_1(x_1) \cdots \varphi_1(x_l) \varphi_2(x_{l+1}) \cdots \varphi_2(x_n) \} | 0 \rangle &= \\ &= \frac{\delta^n W[J_1, J_2]}{\delta J_1(x_1) \cdots \delta J_1(x_l) \delta J_2(x_{l+1}) \cdots \delta J_2(x_n)} \Big|_{J_1=J_2=0} \end{aligned}$$

**Free theory.** For a free real scalar theory, i.e. setting  $\lambda = 0$ , one can compute the free propagator exactly. In this case the situation is slightly different? [r]. In the complex fields formulation, one has to compute a complex Gaussian integral of the form

$$\int \left[ \prod_{j=1}^n dz_j dz_j^* \right] e^{-z^* A z} = \frac{\pi^n}{\det A}$$

where  $z \equiv (z_1, \dots, z_n)$  and  $A$  is a square matrix of dimension  $n$  (with hermitian part positive-definite). The Euclidean generating functional is then

$$\begin{aligned} W[J_1, J_2] &= \int [\mathcal{D}\varphi_1 \mathcal{D}\varphi_2] e^{-S_0[\varphi_1] - S_0[\varphi_2]} \exp \left[ \int d^4x (J_1 \varphi_1 + J_2 \varphi_2) \right] \\ &= \int [\mathcal{D}\varphi_1] e^{-S[\varphi_1] + \int d^4x J_1 \varphi_1} \int [\mathcal{D}\varphi_2] e^{-S[\varphi_2] + \int d^4x J_2 \varphi_2} \\ &= W_1[J_1] W_2[J_2] \end{aligned}$$

The theory factorizes.

The Lagrangian above has a  $U(1)$  symmetry of the fields

$$\varphi' = e^{i\alpha} \varphi, \quad \varphi'^* = e^{-i\alpha} \varphi^*, \quad \alpha \in \mathbb{R}$$

Equivalently, the Lagrangian is invariant under  $SO(2)$  of the components  $\varphi_1$  and  $\varphi_2$ .

**Exercise.** Write the most general  $SO(2)$  transformation.

**Exercise.** See Peskin, Problem 12.3, p. 428. The interaction in the action  $S[\varphi_1, \varphi_2]$  can be split into self-interaction and coupling between the fields. This form can be generalized

$$\mathcal{L} = \frac{1}{2} \partial_\mu \varphi_1 \partial^\mu \varphi_1 + \frac{1}{2} \partial_\mu \varphi_2 \partial^\mu \varphi_2 - \frac{\lambda}{4!} (\varphi_1^4 + \varphi_2^4) - \frac{2\rho}{4!} \varphi_1^2 \varphi_2^2$$

Do the following:

- Write the Euclidean Feynman rules.
- Compute all one-loop corrections to the  $\lambda$ -vertex and the  $\rho$ -vertex in dimensional regularization  $D = 4 - 2\varepsilon$ , use massless integrals (discussed in the following).
- Impose the normalization conditions at  $s = t = u = \Lambda^2$  such that

$$\Gamma^{(4)}(s = t = u = \Lambda^2) = -\lambda, \quad \Gamma_{\text{mixed}}^{(4)}(s = t = u = \Lambda^2) = -2\rho$$

[r] minus?

Notice that this implies that there are no one-loop finite corrections to the couplings.

- Compute the beta functions  $\beta_\lambda$ ,  $\beta_\rho$  and

$$\beta_{\frac{\lambda}{\rho}} = d_t \frac{\lambda}{\rho} = \frac{1}{\rho} d_t \lambda - \frac{1}{\rho^2} \lambda d_t \rho = \frac{1}{\rho} \beta_\lambda - \frac{\lambda}{\rho^2} \beta_\rho$$

- Find the fixed points and describe the RG flow.
- Explain what happens for

$$\frac{1}{3} < \frac{\lambda}{\rho} < 1$$

and for  $\lambda = \rho$ .

## 1.2 Massless integrals in dimensional regularization

**Proposition.** In dimensional regularization, it holds

$$I_\alpha = \int \frac{d^D p}{(2\pi)^D} \frac{1}{(p^2)^\alpha} = 0, \quad \alpha \in \mathbb{C}$$

Notice that it is null also for  $\alpha = 0$ . This integral describes tadpoles for  $\alpha = 1$ : the two-point Green's function receives no correction at one-loop if the fields are massless.

The intuitive argument of why this is true can be seen for  $D \neq 2\alpha$  where the integral is dimensionful: the result has to be dimensionful as well and must be written in terms of dimensionful parameters, however the integral does not depend on any parameter and so the result must be zero.

*Proof.* Consider the theorem below and the Euclidean formalism. One may rewrite the integrand

$$\frac{1}{(p^2)^\alpha} = \frac{1}{(p^2)^\alpha} \frac{p^2 + m^2}{p^2 + m^2} = \frac{1}{(p^2)^\alpha} \frac{m^2}{p^2 + m^2} + \frac{1}{(p^2)^{\alpha-1}} \frac{1}{p^2 + m^2}$$

The region of convergence for the integral

$$\int d^D p \frac{1}{(p^2)^\alpha} \frac{m^2}{p^2 + m^2}$$

is given by the following limits. At infinity, one has

$$|p| \rightarrow \infty, \quad I \sim \frac{1}{p^{2\alpha+2-D}}$$

There is no singularity for  $2\alpha + 2 - \text{Re } D > 0$ . At the origin, one has

$$|p| \rightarrow 0, \quad I \sim p^{D-2\alpha}$$

There is no singularity for  $\text{Re } D - 2\alpha > 0$ . The integral of the first addendum is well-defined for

$$2\alpha < \text{Re } D < 2\alpha + 2$$

Similarly, the region of convergence of the integral of the second addendum is

$$2\alpha - 2 < \text{Re } D < 2\alpha$$

The two regions do not overlap and, by the theorem below, one has

$$I_\alpha = \int \frac{d^D p}{(2\pi)^D} \frac{m^2}{(p^2)^\alpha (p^2 + m^2)} + \int \frac{d^D p}{(2\pi)^D} \frac{1}{(p^2)^{\alpha-1} (p^2 + m^2)} \equiv I_\alpha^{(1)} + I_\alpha^{(2)}$$

Looking at tables of integrals [r], one has a general formula

$$\int \frac{d^D p}{(2\pi)^D} \frac{1}{(p^2)^a (p^2 + m^2)^b} = (m^2)^{\frac{D}{2}-a-b} \frac{\Gamma(D/2-a)\Gamma(a+b-D/2)}{(4\pi)^{\frac{D}{2}} \Gamma(b)\Gamma(D/2)}$$

For the first integral, one has  $a = \alpha$ ,  $b = 1$  and for the second  $a = \alpha - 1$  and  $b = 1$ . Therefore

$$\begin{aligned} I_\alpha &= m^{D-2\alpha} \frac{\Gamma(D/2-\alpha)\Gamma(1+\alpha-D/2)}{(4\pi)^{\frac{D}{2}} \Gamma(D/2)} + m^{D-2\alpha} \frac{\Gamma(D/2+1-\alpha)\Gamma(\alpha-D/2)}{(4\pi)^{\frac{D}{2}} \Gamma(D/2)} \\ &= \frac{m^{D-2\alpha}}{(4\pi)^{\frac{D}{2}} \Gamma(D/2)} \Gamma(\alpha-D/2)\Gamma(D/2-\alpha) \left[ \alpha - \frac{D}{2} + \frac{D}{2} - \alpha \right] = 0 \end{aligned}$$

where one applies  $\Gamma(z+1) = z\Gamma(z)$  to the Gamma function with three addenda in its argument.  $\square$

**Theorem.** If an integrand can be written as a sum of terms with non-overlapping regions of convergence (for their integrals), then the integral is a sum of their integrals.

One a-posteriori argument to see why the integral is null follows. One may rescale  $p = sp'$  to have

$$I_\alpha = \int \frac{d^D p}{(2\pi)^D} \frac{1}{(p^2)^\alpha} = s^{D-2\alpha} \int \frac{d^D p'}{(2\pi)^D} \frac{1}{(p'^2)^\alpha} \implies I_\alpha = s^{D-2\alpha} I_\alpha$$

For  $D \neq 2\alpha$  then  $I_\alpha = 0$ . For the case  $D = 2\alpha$ , one can argue that there is a fine-tuning between the ultraviolet and the infrared divergences. In this case, there divergences are both present and give zero when summed. Consider the particular example of  $D = 2 - \varepsilon$  and  $\alpha = 1$ . Then, the integral becomes

$$\begin{aligned} I_1 &= \int \frac{d^D p}{(2\pi)^D} \frac{1}{p^2} = - \int \frac{d^D p}{(2\pi)^D} \int_0^\infty d\lambda \frac{1}{(p^2 - \lambda)^2} = - \int_0^\infty d\lambda \int \frac{d^{2-\varepsilon} p}{(2\pi)^{2-\varepsilon}} \frac{1}{(p^2 - \lambda)^2} \\ &= - \int_0^\infty d\lambda \frac{\Gamma(1+\varepsilon/2)}{(4\pi)^{1-\varepsilon/2} \Gamma(2)} (-\lambda)^{-\frac{\varepsilon}{2}-1} = (-1)^{\frac{\varepsilon}{2}} \frac{\Gamma(1+\varepsilon/2)}{(4\pi)^{1-\varepsilon/2}} \int_0^\infty \frac{d\lambda}{\lambda^{1+\frac{\varepsilon}{2}}} \end{aligned}$$

Ignoring the coefficients, the integral is

$$\begin{aligned} I_1 &\propto \int_{a_{\text{IR}}}^1 \frac{d\lambda}{\lambda^{1+\frac{\varepsilon}{2}}} + \int_1^{a_{\text{UV}}} \frac{d\lambda}{\lambda^{1+\frac{\varepsilon}{2}}} = -\frac{2}{\varepsilon} \lambda^{-\frac{\varepsilon}{2}} \Big|_{a_{\text{IR}}}^1 - \frac{2}{\varepsilon} \lambda^{-\frac{\varepsilon}{2}} \Big|_1^{a_{\text{UV}}} = -\frac{2}{\varepsilon} [-a_{\text{IR}}^{-\frac{\varepsilon}{2}} + a_{\text{UV}}^{-\frac{\varepsilon}{2}}] \\ &= -\frac{2}{\varepsilon} \left[ -\frac{\varepsilon}{2} (-\ln a_{\text{IR}} + \ln a_{\text{UV}}) + o(\varepsilon) \right] = \ln a_{\text{UV}} - \ln a_{\text{IR}} + o(\varepsilon^0) = 0 \end{aligned}$$

At the second line, one has integrated up to a cutoff [r]. From this one sees that  $a_{\text{IR}} = a_{\text{UV}}$  and  $I_\alpha = 0$  is the product of a cancellation between the ultraviolet and infrared divergences.

When dealing with massless theories in two dimensions  $D = 2$ , the cancellation of the tadpole is due to a balance between infrared and ultraviolet divergences. When interested in either divergence, one has to remove the other divergence in order to renormalize the one of interest. For example, one replaces

$$I_1 = \int \frac{d^{2-\varepsilon}p}{(2\pi)^{2-\varepsilon}} \frac{1}{p^2} \rightarrow \int \frac{d^{2-\varepsilon}p}{(2\pi)^{2-\varepsilon}} \frac{1}{p^2 + \mu^2} = \frac{\Gamma(\varepsilon/2)}{(4\pi)^{1-\frac{\varepsilon}{2}}} \mu^{-\varepsilon} \sim \frac{2}{\varepsilon} e^{-\varepsilon \ln \mu + \dots} \sim \frac{2}{\varepsilon}, \quad \varepsilon \rightarrow 0$$

where  $\mu^2$  is an infrared regulator. This is the ultraviolet divergence of the tadpole in two dimensions.

Therefore, in dimensional regularization, massless tadpoles can be ignored when  $D \neq 2$ . In  $D = 2$  the tadpole is dimensionless and gives a contribution when removing one divergence.

## Part I

# Spin-half fermion fields

## 2 Introduction

**Review – classical fields.** See Srednicki. A fermion field is a field describing Dirac spinors. A Dirac spinor in the Weyl basis is comprised of two fixed-chirality Weyl spinors

$$\psi = \begin{bmatrix} \chi_L \\ \chi_R \end{bmatrix}$$

where the left-chiral spinor belongs to the  $(\frac{1}{2}, 0)$  representation of the Lorentz group  $SO(1, 3)$ , while the right-chiral spinor belongs to  $(0, \frac{1}{2})$ .

The equation of motion of the Dirac field is the Dirac equation

$$(i \not{\partial} - m)\psi(x) = 0$$

where the Dirac matrices are four square matrices of dimension four defining the Dirac algebra  $Cl_{1,3}(\mathbb{C})$

$$\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu}$$

The equation of motion can be obtained from the action principle with the Lagrangian

$$\mathcal{L} = i\bar{\psi} \not{\partial} \psi - m\bar{\psi}\psi = \bar{\psi}(i \not{\partial} - m)\psi, \quad \bar{\psi} = \psi^\dagger \gamma^0$$

If the field  $\psi$  satisfies the Dirac equation of motion, then each of its four components (each two of the Weyl spinors) satisfies the Klein–Gordon equation

$$(\square + m^2)\psi_j(x) = 0$$

In momentum space, the above is an algebraic equation that gives the dispersion relation

$$(-p^2 + m^2)\psi(p) = 0 \implies p^2 = m^2 \iff p^0 = \sqrt{|\mathbf{p}|^2 + m^2} \equiv \omega$$

The most general solution of the Klein–Gordon equation has the form<sup>1</sup>

$$\psi(x) = \int \frac{d^3p}{(2\pi)^3 2\omega} [u(\mathbf{p})e^{-ipx} + v(\mathbf{p})e^{ipx}]_{p^0=\omega}$$

Imposing the Dirac equation one finds

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

One solves these equation by going to the rest frame  $\mathbf{p} = 0$  for which

$$\not{p} = \gamma^0 p_0 = \gamma^0 m \implies (\gamma^0 - 1)u(0) = 0, \quad (\gamma^0 + 1)v(0) = 0$$

<sup>1</sup>Understood as a group of four components that solve the equation.

Each equation has two independent solutions  $u_{\pm}$  and  $v_{\pm}$ . To get the solution for an arbitrary momentum one has to perform a boost. Therefore, the general solution to the Dirac equation is

$$\psi(x) = \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [b_s(\mathbf{p}) u_s(\mathbf{p}) e^{-ipx} + d_s^\dagger v_s(\mathbf{p}) e^{ipx}]_{p^0=\omega}$$

where  $b$  and  $d$  are numbers.

The two-spinors  $u_s$  and  $v_s$  satisfy several identities (see Peskin, p. 48). The normalization is chosen to be

$$\bar{u}_r(\mathbf{p}) u_s(\mathbf{p}) = 2m \delta_{rs}, \quad \bar{v}_r(\mathbf{p}) v_s(\mathbf{p}) = -2m \delta_{rs}$$

Equivalent to

$$u_r^\dagger(\mathbf{p}) u_s(\mathbf{p}) = 2p^0 \delta_{rs}, \quad v_r^\dagger(\mathbf{p}) v_s(\mathbf{p}) = 2p^0 \delta_{rs}$$

The spinors are orthogonal

$$\bar{u}_r(\mathbf{p}) v_s(\mathbf{p}) = \bar{v}_r(\mathbf{p}) u_s(\mathbf{p}) = 0, \quad u_r^\dagger(\mathbf{p}) v_s(-\mathbf{p}) = v_r^\dagger(-\mathbf{p}) u_s(\mathbf{p}) = 0$$

The last two relations are not zero for both momenta being  $+\mathbf{p}$ . With these, one can obtain

$$b_s(\mathbf{p}) = \int d^3x e^{ipx} u_s^\dagger(\mathbf{p}) \psi(x), \quad d_s(\mathbf{p}) = \int d^3x e^{ipx} \psi^\dagger(x) v_s(\mathbf{p})$$

To see this, it is sufficient to Fourier transform  $u_s^\dagger(\mathbf{p}) \psi(x)$ , substitute  $\psi(x)$  and apply the above rules. Notice that

$$\int d^3x e^{ix(p-q)} = e^{ix^0(p^0-q^0)} (2\pi)^3 \delta^3(\mathbf{p}-\mathbf{q})$$

and that

$$|\mathbf{p}| = |\mathbf{q}| \implies p^0 = q^0$$

## Lecture 2

**Canonical quantization.** See Srednicki, §39. The classical fields are promoted to operator fields by imposing a suitable set of equal-time anti-commutation rules (ACR)

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$$\{\psi_\alpha(x), \psi_\beta(y)\} = 0, \quad \{\psi_\alpha(x), \psi_\beta^\dagger(y)\} = \delta_{\alpha\beta} \delta^{(3)}(\mathbf{x}-\mathbf{y})$$

From these, one obtains the rules for the Fourier coefficients which are promoted to operators

$$\{b_r(\mathbf{p}), b_s(\mathbf{q})\} = \{d_r(\mathbf{p}), d_s(\mathbf{q})\} = 2\omega (2\pi)^3 \delta_{rs} \delta^{(3)}(\mathbf{p}-\mathbf{q})$$

One defines the vacuum state as the state destroyed by every annihilation operator

$$b_s(\mathbf{p}) |0\rangle = d_s(\mathbf{p}) |0\rangle = 0, \quad \forall b, d$$

The excited states are obtained by acting on the vacuum with the creation operators  $b_s^\dagger$  and  $d_s^\dagger$ . One interprets  $b_s^\dagger(\mathbf{p}) |0\rangle$  as a single-particle state with momentum  $\mathbf{p}$ , energy  $\omega = \sqrt{|\mathbf{p}|^2 + m^2}$  and projection  $s$  (in units) of the spin along the  $z$ -direction,  $S_z = \frac{1}{2}s$ . The difference between the two creation operators is not yet clear.

The use of anti-commutation rules implies the Pauli exclusion principle: there cannot be two fermions with the same quantum numbers. This is reflected in the spectrum of the number operator and the relation

$$\{b_s^\dagger(\mathbf{p}), b_s^\dagger(\mathbf{p})\} = 0 \implies [b_s^\dagger(\mathbf{p})]^2 = 0$$

The Hamiltonian density is obtained by Legendre transforming the Lagrangian density. By substituting the fields  $\psi$  and  $\bar{\psi}$ , one obtains the energy (Hamiltonian) of the system

$$H = \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} \omega [N_s^b(\mathbf{p}) + N_s^d(\mathbf{p})]$$

where  $N_s^a$  is the number operator of spin  $s$  for the ladder operator  $a$

$$N_s^a(\mathbf{p}) \equiv a_s^\dagger(\mathbf{p})a_s(\mathbf{p})$$

The anti-commutation relations imply a finite spectrum

$$[N_s^a(\mathbf{p})]^2 = N_s^a(\mathbf{p}) \implies n = 0, 1$$

The Dirac Lagrangian exhibits a global U(1) symmetry

$$\psi' = e^{i\alpha}\psi, \quad \bar{\psi}' = e^{-i\alpha}\bar{\psi}, \quad \alpha \in \mathbb{R}$$

Noether's theorem implies a conserved current

$$\partial_\mu J^\mu = 0, \quad J^\mu = \bar{\psi}\gamma^\mu\psi$$

The conserved charge is

$$Q = \int d^3x J^0(x) = \int d^3x \bar{\psi}\gamma^0\psi = \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [N_s^b(\mathbf{p}) - N_s^d(\mathbf{p})]$$

The global symmetry gives a minus sign for the operator  $d$ : the creation operators have different meanings. The operator  $d$  treats particles with opposite U(1) charge of the operator  $b$ . The  $b$ -particles, with charge  $Q = 1$ , are simply called particles; while the  $d$ -particles, with  $Q = -1$ , are called anti-particles.

**Free fermion propagator.** See Srednicki, §42. The free fermion propagator is the inverse of the kinetic term in the Lagrangian

$$S_{\alpha\beta}(x-y) \equiv -i \langle 0 | \mathcal{T} \{ \psi_\alpha(x) \bar{\psi}_\beta(y) \} | 0 \rangle = \int \frac{d^4p}{(2\pi)^4} \frac{(\not{p} + m)_{\alpha\beta}}{p^2 - m^2 + i\varepsilon} e^{ip(x-y)}$$

**Exercise.** Check that the propagator is the Green's function of the Dirac operator

$$(i\not{\partial} - m)_{\alpha\beta} S_{\beta\gamma}(x) = \delta_{\alpha\gamma} \delta^{(4)}(x)$$

### 3 LSZ reduction formula

See Srednicki, §41. The physical observables are the cross-sections which are expressed in terms of scattering amplitudes  $\langle f|i \rangle$ . The initial and final states are defined asymptotically in the distant past and future,  $t \rightarrow \mp\infty$ . One utilizes the adiabatic hypothesis: the region where the interaction is non-trivial is finite and, beyond, the theory is essentially free. The hypothesis has to be supported by the ansatz that the free regions are reached in the distant past and future.

One-particle states can be constructed by applying one a creation operator

$$a_s^\dagger(\mathbf{p}) |0\rangle = |\mathbf{p}, s\rangle$$

Due to the Heisenberg uncertainty principle, a more physical particle is the one comprised of a wave-packet of momenta. The physical creation operator is a Gaussian distribution centered around a central momentum  $\mathbf{q}$

$$a_s^\dagger = \int d^3p f(\mathbf{p}) a_s^\dagger(\mathbf{p}), \quad f(\mathbf{p}) \propto e^{-\frac{(\mathbf{p}-\mathbf{q})^2}{4\sigma^2}}$$

The operation of applying the creation operator to the vacuum to obtain a particle is valid and meaningful in the asymptotic limit, where the states and theory are free. This holds since the vacuum defined by the annihilation operators is the free vacuum. By introducing an interaction, the vacuum is interacting. One assumes to be able to define a vacuum in the interacting theory and that the particles can be obtained in the same way as the free theory. For instance, a one-particle initial and final state are taken to be

$$|i\rangle = \lim_{t \rightarrow -\infty} b_r^\dagger(t) |0\rangle, \quad |f\rangle = \lim_{t \rightarrow \infty} b_s^\dagger(t) |0\rangle$$

In the interacting theory, the ladder operators depend on time.



**Time dependence.** To obtain a useful expression for the scattering amplitude, one needs to study the time dependence of the ladder operators. As with the scalar fields, consider

$$\begin{aligned}
 b_s(+\infty) - b_s(-\infty) &= \int_{-\infty}^{\infty} dt \partial_0 b_s(t) = \int d^3p f(\mathbf{p}) \int_{-\infty}^{\infty} dt \partial_0 b_s(\mathbf{p}) \\
 &= \int d^3p f(\mathbf{p}) \int_{-\infty}^{\infty} dt \int d^3x \partial_0 [e^{ipx} u_s^\dagger(\mathbf{p}) \psi(x)] \\
 &= \int d^3p f(\mathbf{p}) \int d^4x \partial_0 [e^{ipx} \bar{u}_s(\mathbf{p}) \gamma^0 \psi(x)] \\
 &= \int d^3p f(\mathbf{p}) \int d^4x e^{ipx} \bar{u}_s(\mathbf{p}) [ip_0 \gamma^0 + \gamma^0 \partial_0] \psi(x) \\
 &= \int d^3p f(\mathbf{p}) \int d^4x e^{ipx} \bar{u}_s(\mathbf{p}) [-ip_j \gamma^j + im + \gamma^0 \partial_0] \psi(x) \\
 &= \int d^3p f(\mathbf{p}) \int d^4x \bar{u}_s(\mathbf{p}) [-\gamma^j \partial_j e^{ipx} + e^{ipx} (im + \gamma^0 \partial_0)] \psi(x) \\
 &= \int d^3p f(\mathbf{p}) \int d^4x e^{ipx} \bar{u}_s(\mathbf{p}) [\gamma^j \partial_j + im + \gamma^0 \partial_0] \psi(x) \\
 &= -i \int d^3p f(\mathbf{p}) \int d^4x e^{ipx} \bar{u}_s(\mathbf{p}) [i \not{\partial} - m] \psi(x)
 \end{aligned}$$

At the first line, one inserts the expression of the wave-packet operator  $b_1$ . At the second line, one inserts the expression of the annihilation operator  $b_s(\mathbf{p})$ . At the fifth line, one has remembered that

$$\bar{u}(\mathbf{p})(\not{p} - m) = 0 \implies \bar{u}(\mathbf{p})(p_0 \gamma^0 + p_j \gamma^j - m) = 0 \implies \bar{u}(\mathbf{p}) p_0 \gamma^0 = \bar{u}(\mathbf{p}) (-p_j \gamma^j + m)$$

At the penultimate line, one has integrated by parts the first addendum.

The integrand is zero if the theory is free, but the theory is interacting and so

$$(i \not{\partial} - m) \psi(x) = -\delta_{\bar{\psi}(x)} \mathcal{L}_{\text{int}} \neq 0$$

and the ladder operators are functions of time. For the creation operators, one has

$$b_s^\dagger(+\infty) - b_s^\dagger(-\infty) = -i \int d^3p f(\mathbf{p}) \int d^4x \bar{\psi}(x) (i \overleftarrow{\not{\partial}} + m) u_s(\mathbf{p}) e^{-ipx}$$

[r]

**Two-by-two scattering.** Consider a two-by-two scattering  $p_1 p_2 \rightarrow p_1' p_2'$ . The scattering amplitude is

$$\langle f|i \rangle = \langle 0| b_{2'}(+\infty) b_{1'}(+\infty) b_1^\dagger(-\infty) b_2^\dagger(-\infty) |0 \rangle$$

where the index of the spin projection is dropped and 1, 2 or 1', 2' indicates the particle. Thanks to the ordering of the operators, the above is equal to

$$\langle f|i \rangle = \langle 0| \mathcal{T} \{ b_{2'}(+\infty) b_{1'}(+\infty) b_1^\dagger(-\infty) b_2^\dagger(-\infty) \} |0 \rangle$$

In the limit  $\sigma \rightarrow 0$ , the wave-packet is a Dirac delta giving

$$b_s^\dagger(+\infty) - b_s^\dagger(-\infty) = -i \int d^4x e^{-ipx} \bar{\psi}(x) (i \overleftarrow{\not{\partial}} + m) u_s(\mathbf{p})$$

One would like to replace  $b_j^\dagger(-\infty)$  with  $b_j^\dagger(+\infty)$  and the integral above. Due to the time-ordered product, the operator  $b_j^\dagger(+\infty)$  is moved to the left and gives zero when acting on the

bra  $[r]$ . Similarly happens for the annihilation operators. Therefore the scattering amplitude is

$$\begin{aligned} \langle f|i \rangle &= (-1)^2 i^4 \langle 0| \mathcal{T} \left\{ \int d^4x_1 d^4x_2 d^4x_{1'} d^4x_{2'} \right. \\ &\quad \times e^{ip_{2'}x_{2'}} \bar{u}_{s_{2'}}(\mathbf{p}_{2'}) (i\bar{\not{\partial}} - m) \psi(x_{2'}) e^{ip_{1'}x_{1'}} \bar{u}_{s_{1'}}(\mathbf{p}_{1'}) (i\bar{\not{\partial}} - m) \psi(x_{1'}) \\ &\quad \times \bar{\psi}(x_1) (i\overleftarrow{\not{\partial}} + m) u_{s_1}(\mathbf{p}_1) e^{-ip_1x_1} \bar{\psi}(x_2) (i\overleftarrow{\not{\partial}} + m) u_{s_2}(\mathbf{p}_2) e^{-ip_2x_2} \left. \right\} |0 \rangle \\ &= \int d^4x_1 d^4x_2 d^4x_{1'} d^4x_{2'} e^{ip_{2'}x_{2'}} \bar{u}_{s_{2'}}(\mathbf{p}_{2'}) (i\bar{\not{\partial}} - m)_{2'} e^{ip_{1'}x_{1'}} \bar{u}_{s_{1'}}(\mathbf{p}_{1'}) (i\bar{\not{\partial}} - m)_{1'} \\ &\quad \times \langle 0| \mathcal{T} \{ \psi(x_{2'}) \psi(x_{1'}) \bar{\psi}(x_1) \bar{\psi}(x_2) \} |0 \rangle (i\overleftarrow{\not{\partial}} + m)_1 u_{s_1} e^{ip_1x_1} (i\overleftarrow{\not{\partial}} + m)_2 u_{s_2} e^{ip_2x_2} \end{aligned}$$

[r] This is the LSZ reduction formula for fermions. In general, the computation of a scattering amplitude can be expressed as a computation of correlation functions

$$G^{(2n)}(x_1, \dots, x_n, x'_1, \dots, x'_n) = \langle 0| \mathcal{T} \{ \psi(x'_1) \cdots \psi(x'_n) \bar{\psi}(x_1) \cdots \bar{\psi}(x_n) \} |0 \rangle$$

[r] One needs to develop a functional approach to compute the above Green's functions.

## 4 Functional quantization

See Srednicki, §43. Consider an interacting fermionic theory

$$\mathcal{L}(\psi, \bar{\psi}) = \bar{\psi}(i\bar{\not{\partial}} - m)\psi + \mathcal{L}_{\text{int}}(\psi, \bar{\psi})$$

One needs to generalize the generating functional for fermions. One introduces the spinorial source terms  $\eta$  and  $\bar{\eta}$  to formally write

$$W[\eta, \bar{\eta}] = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ i \int d^4x [\mathcal{L}(x) + \bar{\psi}\eta + \bar{\eta}\psi] \right]$$

where the spinorial product is defined using the van der Waerden notation<sup>2</sup>. The integrand must be a real quantity. In fact

$$(\bar{\psi}\eta)^\dagger = (\psi^\dagger \gamma^0 \eta)^\dagger = \eta^\dagger \gamma^0 \psi = \bar{\eta}\psi$$

which implies that  $\bar{\psi}\eta + \bar{\eta}\psi$  is real. Since spinors anti-commute, then

$$\bar{\psi}\eta = -\eta\bar{\psi}$$

Therefore, when differentiating one has to be careful about the signs

$$\delta_{\eta(x)} \int d^4y \bar{\psi}(y) \eta(y) = -\bar{\psi}(x)$$

The Green's function is

$$\begin{aligned} G^{(2n)}(x_1, \dots, x_n, y_1, \dots, y_n) &= \langle 0| \mathcal{T} \{ \psi_{\alpha_1}(x_1) \cdots \psi_{\alpha_n}(x_n) \bar{\psi}_{\beta_1}(y_1) \cdots \bar{\psi}_{\beta_n}(y_n) \} |0 \rangle \\ &= (-1)^n i^{2n} \frac{\delta^{2n} W[\eta, \bar{\eta}]}{\delta \bar{\eta}_{\alpha_1}(x_1) \cdots \delta \bar{\eta}_{\alpha_n}(x_n) \delta \eta_{\beta_1}(y_1) \cdots \delta \eta_{\beta_n}(y_n)} \Big|_{\eta=\bar{\eta}=0} \\ &= \frac{\delta^{2n} W[\eta, \bar{\eta}]}{\delta \bar{\eta}_{\alpha_1}(x_1) \cdots \delta \bar{\eta}_{\alpha_n}(x_n) \delta \eta_{\beta_1}(y_1) \cdots \delta \eta_{\beta_n}(y_n)} \Big|_{\eta=\bar{\eta}=0} \end{aligned}$$

Each derivative with respect to  $\bar{\eta}$  brings a  $-i$  while each derivative with respect to  $\eta$  brings an  $i$ .

The Euclidean functional integral is

$$W_E[\eta, \bar{\eta}] = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ - \int d^4x [\mathcal{L}(x) - \bar{\psi}\eta - \bar{\eta}\psi] \right]$$

<sup>2</sup>See also eigenchris, Spinors for Beginners 9, <https://youtu.be/4NJBvkjpc3E?t=2340>, at minute 39:00 as well as A. Steane, An Introduction to Spinors, <https://arxiv.org/abs/1312.3824>.

and the Euclidean Green's function is [r]

$$G_E^{(2n)}(x_1, \dots, x_n, y_1, \dots, y_n) = (-1)^n \frac{\delta^{2n} W_E[\eta, \bar{\eta}]}{\delta \bar{\eta}_{\alpha_1}(x_1) \cdots \delta \bar{\eta}_{\alpha_n}(x_n) \delta \eta_{\beta_1}(y_1) \cdots \delta \eta_{\beta_n}(y_n)} \Big|_{\eta=\bar{\eta}=0}$$

The functional integral involves spinorial variables. This kind of integration has to be properly defined.

**Exercise.** Let  $n = 1$ , then

$$\begin{aligned} G^{(2)}(x, y) &= (-1) \frac{\delta^2 W_E[\eta, \bar{\eta}]}{\delta \bar{\eta}_\alpha(x) \delta \eta_\beta(y)} \Big|_{\eta=\bar{\eta}=0} \\ &= (-1) \frac{\delta}{\delta \bar{\eta}_\alpha(x)} \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ - \int d^4 x' [\mathcal{L} - \bar{\psi} \eta - \bar{\eta} \psi] \right] [-\bar{\psi}_\beta(y)] \Big|_{\eta=\bar{\eta}=0} \\ &= \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ - \int d^4 x' [\mathcal{L} - \bar{\psi} \eta - \bar{\eta} \psi] \right] \psi_\alpha(x) \bar{\psi}_\beta(y) \Big|_{\eta=\bar{\eta}=0} \\ &= \langle 0 | \mathcal{T} \{ \psi_\alpha(x) \bar{\psi}_\beta(y) \} | 0 \rangle \end{aligned}$$

## Lecture 3

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### 4.1 Grassmann algebra

See Cheng, §1.3. See also DeWitt, Supermanifolds. The fields  $\psi$  and  $\bar{\psi}$  are classical spinor fields with four components in Dirac's notation

$$\psi_\alpha, \quad \alpha = 1, 2, 3, 4$$

Each component is a Grassmann-odd number<sup>3</sup>: it anti-commutes with itself and other Grassmann-odd numbers. For example

$$\psi_1 \psi_2 = -\psi_2 \psi_1$$

Grassmann numbers are needed when taking the classical limit of the quantized spinor fields which are anti-commuting. When using a functional definition of quantum field theory, the quantities appearing inside the path integral are all classical, not operators, so they have to be Grassmann numbers.

The path integral is the continuum limit of an ordinary integral on a lattice. A field is evaluated only at the sites of the lattice

$$\psi_\alpha(x^i) \equiv \psi_\alpha^i$$

The path integral on the lattice is an ordinary integral but on Grassmann-odd variables<sup>4</sup>

$$\int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \leftrightarrow \int \prod_{\alpha, i} d\psi_\alpha^i \prod_{\alpha, i} d\bar{\psi}_\alpha^i$$

One needs to define analysis for Grassmann numbers.

**Algebra.** An  $n$ -dimensional Grassmann algebra is generated by  $n$  Grassmann variables  $\theta_i$  that anti-commute

$$\{\theta_i, \theta_j\} = 0 \implies \theta_i^2 = 0$$

A generic element of the algebra can be expanded in a finite Taylor series

$$f(\theta_1, \dots, \theta_n) = f_0 + f_{1,i} \theta^i + f_{2,ij} \theta^i \theta^j + \cdots + f_{n,i_1 \dots i_n} \theta^{i_1} \cdots \theta^{i_n}$$

<sup>3</sup>Odd or even refers to the number of Grassmann variables  $\theta_i$  in the expansion of a Grassmann number  $z$  in terms of such Grassmann variables  $\theta_i$ . A Grassmann number may not have definite Grassmann parity, but can be separated into odd and even parts. Odd Grassmann numbers anti-commute between themselves and even Grassmann numbers commute with every Grassmann number.

<sup>4</sup>Grassmann numbers are individual elements of the exterior algebra generated by a set of  $n$  Grassmann variables.

This is because if there is any repeated variable, then it is zero  $\theta_j^2 = 0$ . The coefficients of the expansion are complex numbers that are completely anti-symmetric in the indices  $i_k$ .

Consider the expansion of an element of the algebra by writing explicitly the dependence on one particular<sup>5</sup> Grassmann variable  $\theta$ . In such variable, the expansion is at most linear

$$f(\theta) = f_0 + f_1\theta$$

The coefficients  $f_0$  and  $f_1$  are independent of  $\theta$  and depend on the other  $n - 1$  Grassmann variables.

In general, one has to consider a function also of space-time

$$f(x, \theta) = f_0(x) + f_1(x)\theta$$

The coefficient of the  $\theta$ -expansion are space-time fields. If  $f(x, \theta)$  is a Grassmann-even field, then the Grassmann parity must be the same on either side of the equation, so  $f_0(x)$  is a Grassmann-even field and  $f_1(x)$  is a Grassmann-odd field.

The addition of Grassmann variables to space-time gives superspace<sup>6</sup>. A field on superspace is a superfield and contains different ordinary fields. In supersymmetry, a superfield is a representation of the supersymmetry algebra.

**Differentiation.** The left and right derivatives are defined as

$$d_\theta \theta = \theta \overleftarrow{d}_\theta = 1$$

Therefore, the left derivative of the above element of the algebra is

$$d_\theta f(\theta) = \begin{cases} +f_1, & f \text{ odd} \\ -f_1, & f \text{ even} \end{cases}$$

while the right derivative is simpler

$$f(\theta) \overleftarrow{d}_\theta = f_1$$

**Integration.** The integration over Grassmann-odd variables is called Berezin integration<sup>7</sup> which differs from Lebesgue's.

Typically, the integral is the inverse operator of the derivative. Since the partial derivative with respect to a Grassmann variable is a Grassmann-odd operator then the second derivative is zero<sup>8</sup>. The derivative is not an invertible operator. The integral must be defined in another way

$$\int d\theta f(\theta)$$

One imposes two properties: linearity

$$\int d\theta (f + \alpha g) = \int d\theta f + \alpha \int d\theta g, \quad \alpha \in \mathbb{C}$$

and translational invariance

$$\int d\theta f(\theta) \equiv \int d(\theta + \eta) f(\theta + \eta)$$

where  $\eta$  is Grassmann-odd number independent of the particular variable  $\theta$ .

<sup>5</sup>This means that the dependence on the other  $n - 1$  variables is hidden inside the expansion coefficients which are no longer complex numbers, but Grassmann numbers.

<sup>6</sup>In particular Minkowski superspace: Minkowski space is extended with anti-commuting fermionic degrees of freedom, taken to be anti-commuting Weyl spinors from the Clifford algebra associated to the Lorentz group.

<sup>7</sup>See Berezin, F. A. (1966). The Method of Second Quantization. Pure and Applied Physics. Vol. 24. New York. ISSN 0079-8193. <https://www.sciencedirect.com/bookseries/pure-and-applied-physics/vol/24>.

<sup>8</sup>This is a consequence of  $\partial_z z = 1$ .

The second property can be expanded to have

$$\int d\theta f(\theta) = \int d(\theta + \eta) f(\theta + \eta) = \int d\theta f(\theta + \eta) \implies \int d\theta (f_0 - \theta f_1) = \int d\theta [f_0 - (\theta + \eta)f_1]$$

Applying linearity, one finds

$$\int d\theta \eta f_1 = 0 \implies \eta f_1 \int d\theta = 0 \implies \boxed{\int d\theta = 0}$$

since the above has to hold for every  $f_1$  and  $\eta$  (which do not depend on  $\theta$ ). Using this result, one obtains

$$\int d\theta f(\theta) = \int d\theta f_0 + f_1\theta = \int d\theta f_1\theta = \pm f_1 \int d\theta \theta \equiv \pm f_1 \implies \boxed{\int d\theta \theta = 1}$$

where  $+$  is for  $f$  odd and  $-$  for  $f$  even. One may notice that Berezin integration is equivalent to differentiation

$$\int d\theta f(\theta) = \pm f_1 = d_\theta f$$

Therefore, the general definition of integration in a Grassmann-odd variable is

$$\boxed{\int d\theta f(\theta) \equiv d_\theta f|_{\theta=0}}$$

[r] why evaluated at  $\theta = 0$ ? for defining the Berezin integral also for ordinary functions?

**Change of variables.** When performing a change of variables, one has to account for the Jacobian. For the Berezin integral, the inverse Jacobian is produced instead.

For ordinary (Grassmann-even, Riemann or Lebesgue) integration, a change of variable produces

$$\int dy f(y), \quad y = g(x) \implies \int dy f(y) = \int dx |\partial_x y| f(g(x))$$

For Grassmann-odd integration, a change of variables is

$$\theta' = g(\theta) = a + b\theta, \quad d_\theta \theta' = b$$

Notice that  $a$  is odd and  $b$  is even. Knowing that

$$\int d\theta' f(\theta') = \int d\theta' (f_0 + f_1\theta') = \pm f_1$$

one may consider the integration without Jacobian

$$\int d\theta f(\theta'(\theta)) = \int d\theta f(a + b\theta) = \int d\theta [f_0 + f_1(a + b\theta)] = \pm f_1 b = \pm f_1 d_\theta \theta'$$

Therefore, when performing the change of variables, one has to use the inverse Jacobian

$$\int d\theta' f(\theta') = \int d\theta |d_\theta \theta'|^{-1} f(\theta'(\theta))$$

**Generalization to more variables.** The differentiation and integration results above can be generalized to more Grassmann variables.

**Differentiation.** The left and right derivatives are

$$d_{\theta_i} \theta_j = \theta_j \overleftarrow{d}_{\theta_i} = \delta_{ij}$$

The derivative are Grassmann-odd operators

$$\{d_{\theta_i}, d_{\theta_j}\} = 0, \quad \{d_{\theta_i}, \theta_j\} = 0$$

where, in the second equality, the operator is understood as not acting on  $\theta_j$ : when computing the derivative of a product, a minus sign appears every time the derivative goes through an odd variable. As such, the derivative of a product is

$$d_{\theta_i}(\theta_1 \cdots \theta_n) = \delta_{i1} \theta_2 \cdots \theta_n - \theta_1 \delta_{i2} \theta_3 \cdots \theta_n + \cdots + (-1)^{n-1} \theta_1 \cdots \theta_{n-1} \delta_{in} = \sum_{j=1}^n (-1)^{j-1} \delta_{ij} \prod_{k \neq j}^n \theta_k$$

while the right derivative is similar

$$(\theta_1 \cdots \theta_n) \overleftarrow{d}_{\theta_i} = \theta_1 \cdots \theta_{n-1} \delta_{in} - \theta_1 \cdots \delta_{i,n-1} \theta_n + \cdots + (-1)^{n-1} \delta_{i1} \theta_2 \cdots \theta_n = \sum_{j=1}^n (-1)^{n-j} \delta_{ij} \prod_{k \neq j}^n \theta_k$$

**Integration.** For integration, the measures anti-commute

$$\{d\theta_i, d\theta_j\} = 0$$

The fundamental integral properties are

$$\int d\theta_i 1 = 0, \quad \int d\theta_i \theta_j = \delta_{ij}$$

The integral in one variable of an element of the algebra is

$$\int d\theta_i f(\theta_1, \dots, \theta_n) = d_{\theta_i} f(\theta_1, \dots, \theta_n)|_{\theta_i=0}$$

The integral in multiple variables is

$$\int d\theta_n \cdots d\theta_1 f(\theta_1, \dots, \theta_n) = d_{\theta_n} \cdots d_{\theta_1} f|_{\theta_1=\dots=\theta_n=0}$$

**Change of variables.** For a change of variables  $\theta'_i = b_{ij} \theta_j$ , one has

$$\int d\theta'_n \cdots d\theta'_1 f(\theta'_1, \dots, \theta'_n) = \int d\theta_n \cdots d\theta_1 (\det b)^{-1} f(\theta'_1(\theta), \dots, \theta'_n(\theta))$$

**Gaussian integrals.** The Gaussian integral is

$$G(A) = \int d\theta_n \cdots d\theta_1 e^{\frac{1}{2} \theta_i A_{ij} \theta_j}$$

where  $A$  is an anti-symmetric square matrix of size  $n$ ,  $A_{ij} = -A_{ji}$ .

Let  $n = 2$ . The exponent is

$$\theta_i A_{ij} \theta_j = \begin{bmatrix} \theta_1 & \theta_2 \end{bmatrix} \begin{bmatrix} 0 & A_{12} \\ -A_{12} & 0 \end{bmatrix} \begin{bmatrix} \theta_1 \\ \theta_2 \end{bmatrix}$$

The integral is

$$\begin{aligned} G(A) &= \int d\theta_2 d\theta_1 \exp \left[ \frac{1}{2} (\theta_1 A_{12} \theta_2 - \theta_2 A_{12} \theta_1) \right] = \int d\theta_2 d\theta_1 e^{\theta A_{12} \theta} \\ &= \int d\theta_2 d\theta_1 (1 + \theta_1 A_{12} \theta_2) = \int d\theta_2 d\theta_1 \theta_1 A_{12} \theta_2 = A_{12} \int d\theta_2 \theta_2 = A_{12} = \sqrt{\det A} \end{aligned}$$

This is true for any  $n$ . For Berenzin integrals, the determinant is in the numerator of the result and not the denominator.

**Table of Gaussian integrals.** For real Grassmann-even variables

$$\int dx_1 \cdots dx_n e^{-\frac{1}{2}x^\top A x} = \frac{(2\pi)^{\frac{n}{2}}}{\sqrt{\det A}}$$

For complex Grassmann-even variables

$$\int \left[ \prod_{j=1}^n dz_j d\bar{z}_j \right] e^{-z^\dagger A z} = \frac{\pi^{\frac{n}{2}}}{\det A}$$

[r] For real Grassmann-odd variables

$$\int d\theta_1 \cdots d\theta_n e^{\frac{1}{2}\theta A \theta} = \sqrt{\det A}$$

For complex Grassmann-odd variables

$$\int \left[ \prod_{j=1}^n d\theta_j d\bar{\theta}_j \right] e^{\bar{\theta}^\top A \theta} = \det A$$

**Functional integrals.** The generalization to functional integration can be defined on the lattice. The continuum theory is retrieved in the continuum limit of the lattice.

**Example.** Consider the free generating functional of a scalar theory

$$W_0[0] = \int [\mathcal{D}\varphi] \exp \left[ - \int d^4x \frac{1}{2} \varphi (\square + m^2) \varphi \right] \propto [\det(\square + m^2)]^{-\frac{1}{2}}$$

where the determinant of an operator is the product of its eigenvalues. For a spinor theory, one has

$$W_0[0] = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ - \int d^4x \frac{1}{2} \bar{\psi} (i\cancel{\partial} - m) \psi \right] \propto \det(i\cancel{\partial} - m)$$

## 4.2 Free field theory

For a free theory, the functional integral can be computed exactly. For a weak interacting theory, one may apply perturbation theory.

Through Grassmann numbers, one may give meaning to the path integral for fermions

$$W[\eta, \bar{\eta}] = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ i \int d^4x [\mathcal{L}(x) + \bar{\psi}\eta + \bar{\eta}\psi] \right]$$

Consider the free Lagrangian

$$\mathcal{L}_0(\psi, \bar{\psi}) = \bar{\psi} (i\cancel{\partial} - m) \psi$$

The generating functional is

$$W_0[\eta, \bar{\eta}] = N \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ i \int d^4x [\bar{\psi} (i\cancel{\partial} - m) \psi + \bar{\psi}\eta + \bar{\eta}\psi] \right]$$

The propagator is the inverse of the kinetic term. Using the fact that the propagator is Green's function of the Dirac operator, one gets

$$\int d^4x (i\cancel{\partial} - m)_{\alpha\beta} S_{\beta\gamma}(x) = \delta_{\alpha\gamma} \implies S = (i\cancel{\partial} - m)^{-1}$$

One may complete the square in the exponent of the generating functional

$$\begin{aligned} -iE &= \int d^4x [\bar{\psi} (i\cancel{\partial} - m) \psi + \bar{\psi}\eta + \bar{\eta}\psi] \\ &= \int d^4x \left\{ \left[ \bar{\psi} + \int d^4y \bar{\eta}(y) (i\cancel{\partial} - m)_y^{-1} \right] (i\cancel{\partial} - m)_x \left[ \psi + \int d^4y' (i\cancel{\partial} - m)_{y'}^{-1} \eta(y') \right] \right. \\ &\quad \left. - \int d^4y d^4y' \bar{\eta}(y) (i\cancel{\partial} - m)_y^{-1} (i\cancel{\partial} - m)_x (i\cancel{\partial} - m)_{y'}^{-1} \eta(y') \right\} \\ &= \int d^4x \bar{\chi} (i\cancel{\partial} - m)_x \chi - \int d^4x d^4y' \bar{\eta}(x) (i\cancel{\partial} - m)_{y'}^{-1} \eta(y') \end{aligned}$$

[r] where

$$\chi \equiv \psi + \int d^4 y' (\mathbf{i} \not{\partial} - m)_{y'}^{-1} \eta(y'), \quad (\mathbf{i} \not{\partial} - m)_y^{-1} (\mathbf{i} \not{\partial} - m)_x = \delta^{(4)}(x - y)$$

The generating functional becomes

$$\begin{aligned} W_0[\eta, \bar{\eta}] &= N \int [\mathcal{D}\chi \mathcal{D}\bar{\chi}] \exp \left[ \mathbf{i} \int d^4 x \bar{\chi} (\mathbf{i} \not{\partial} - m)_x \chi - \mathbf{i} \int d^4 x d^4 y' \bar{\eta}(x) (\mathbf{i} \not{\partial} - m)_{y'}^{-1} \eta(y') \right] \\ &= \exp \left[ - \mathbf{i} \int d^4 x d^4 y' \bar{\eta}(x) (\mathbf{i} \not{\partial} - m)_{y'}^{-1} \eta(y') \right] W_0[0, 0] \\ &= \exp \left[ - \mathbf{i} \int d^4 x d^4 y' \bar{\eta}_\alpha(x) S_{\alpha\beta}(x - y') \eta_\beta(y') \right] \end{aligned}$$

At the first line, the path integral of the first addendum of the exponent is a Gaussian integral and can be absorbed into the normalization constant  $N$  to give a normalization of  $W_0[\eta = \bar{\eta} = 0] = 1$  [r].

**Remark.** The two-point Green's function for the free theory is

$$\begin{aligned} G^{(2)}(x_1, x_2) &= \frac{\delta^2 W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_\gamma(x_1) \delta \eta_\delta(x_2)} \Big|_{\eta=\bar{\eta}=0} = \frac{\delta}{\delta \bar{\eta}_\gamma(x_1)} \left[ \mathbf{i} \int d^4 x \bar{\eta}_\alpha(x) S_{\alpha\delta}(x - x_2) W_0 \right] \Big|_{\eta=\bar{\eta}=0} \\ &= \mathbf{i} S_{\gamma\delta}(x_1 - x_2) \equiv \langle 0 | \mathcal{T} \{ \psi_\gamma(x_1) \bar{\psi}_\delta(x_2) \} | 0 \rangle \end{aligned}$$

### 4.3 Interacting field theory

The Lagrangian is

$$\mathcal{L}(\psi, \bar{\psi}) = \bar{\psi} (\mathbf{i} \not{\partial} - m) \psi + \mathcal{L}_{\text{int}}(\psi, \bar{\psi})$$

Since the Lagrangian is a scalar, it has to be a function only of bilinears of spinors. The generating functional is

$$W[\eta, \bar{\eta}] = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ \mathbf{i} \int d^4 x [\mathcal{L}_0(x) + \mathcal{L}_{\text{int}}(\psi, \bar{\psi}) + \bar{\psi} \eta + \bar{\eta} \psi] \right]$$

Noting that

$$\delta_{\eta(y)} \int d^4 x \bar{\psi}(x) \eta(x) = -\bar{\psi}(y), \quad \delta_{\bar{\eta}(x)} \int d^4 x \bar{\eta}(x) \psi(x) = \psi(x)$$

one may apply the useful property of functional integrals<sup>9</sup> and rewrite the fields in terms of derivatives

$$\mathcal{L}_{\text{int}}(\psi, \bar{\psi}) \exp \left[ \mathbf{i} \int d^4 x (\bar{\psi} \eta + \bar{\eta} \psi) \right] = \mathcal{L}_{\text{int}}(-\mathbf{i} \delta_{\bar{\eta}(x)}, \mathbf{i} \delta_{\eta(x)}) \exp \left[ \mathbf{i} \int d^4 x (\bar{\psi} \eta + \bar{\eta} \psi) \right]$$

Therefore the generating functional is

$$\begin{aligned} W[\eta, \bar{\eta}] &= \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ \mathbf{i} \int d^4 x \mathcal{L}_{\text{int}}(-\mathbf{i} \delta_{\bar{\eta}}, \mathbf{i} \delta_{\eta}) \right] \exp \left[ \mathbf{i} \int d^4 x [\mathcal{L}_0(x) + \bar{\psi} \eta + \bar{\eta} \psi] \right] \\ &= \exp \left[ \mathbf{i} \int d^4 x \mathcal{L}_{\text{int}}(-\mathbf{i} \delta_{\bar{\eta}}, \mathbf{i} \delta_{\eta}) \right] W_0[\eta, \bar{\eta}] \end{aligned}$$

If the coupling constant is weak, then one may expand the exponential of the interaction in a perturbative series in powers of the coupling constant.

<sup>9</sup>See QFT I.



**Wick's theorem.** To see the effect of Wick's theorem, one may compute the four-point Green's function in the free theory

$$\begin{aligned}
 G_{\alpha\beta\gamma\delta}^{(4)}(x_1, x_2, x_3, x_4) &= \langle 0 | \mathcal{T} \{ \psi_\alpha(x_1) \psi_\beta(x_2) \bar{\psi}_\gamma(x_3) \bar{\psi}_\delta(x_4) \} | 0 \rangle \\
 &= \frac{\delta^4 W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_\alpha(x_1) \delta \bar{\eta}_\beta(x_2) \delta \eta_\gamma(x_3) \delta \eta_\delta(x_4)} \Big|_{\eta=\bar{\eta}=0} \\
 &= \frac{\delta^3}{\delta \bar{\eta}_\alpha(x_1) \delta \bar{\eta}_\beta(x_2) \delta \eta_\gamma(x_3)} \left\{ \exp \left[ -i \int d^4x d^4y \bar{\eta}_\varepsilon(x) S_{\varepsilon\zeta}(x-y) \eta_\zeta(y) \right] \right. \\
 &\quad \left. \times i \int d^4x \bar{\eta}_\varepsilon S_{\varepsilon\delta}(x-x_4) \right\} \Big|_{\eta=\bar{\eta}=0} \\
 &= \frac{\delta^2}{\delta \bar{\eta}_\alpha(x_1) \delta \bar{\eta}_\beta(x_2)} \left[ W_0[\eta, \bar{\eta}] i \int d^4y \bar{\eta}_\varepsilon S_{\varepsilon\gamma}(y-x_3) \right. \\
 &\quad \left. \times i \int d^4x \bar{\eta}_\varepsilon S_{\varepsilon\delta}(x-x_4) \right] \Big|_{\eta=\bar{\eta}=0} \\
 &= \frac{\delta}{\delta \bar{\eta}_\alpha(x_1)} \left[ \frac{\delta W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_\beta(x_2)} i \int d^4y \bar{\eta}_\varepsilon S_{\varepsilon\gamma}(y-x_3) i \int d^4x \bar{\eta}_\varepsilon S_{\varepsilon\delta}(x-x_4) \right. \\
 &\quad + W_0[\eta, \bar{\eta}] i S_{\beta\gamma}(x_2-x_3) i \int d^4x \bar{\eta}_\varepsilon S_{\varepsilon\delta}(x-x_4) \\
 &\quad \left. - W_0[\eta, \bar{\eta}] i \int d^4y \bar{\eta}_\varepsilon S_{\varepsilon\gamma}(y-x_3) i S_{\beta\delta}(x_2-x_4) \right] \Big|_{\eta=\bar{\eta}=0} \\
 &= \frac{\delta}{\delta \bar{\eta}_\alpha(x_1)} \left[ \frac{\delta W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_\beta(x_2)} i \int d^4y \bar{\eta}_\varepsilon S_{\varepsilon\gamma}(y-x_3) i \int d^4x \bar{\eta}_\varepsilon S_{\varepsilon\delta}(x-x_4) \right] \Big|_{\eta=\bar{\eta}=0} \\
 &\quad + \left[ \frac{\delta W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_\alpha(x_1)} i S_{\beta\gamma}(x_2-x_3) i \int d^4x \bar{\eta}_\varepsilon S_{\varepsilon\delta}(x-x_4) \right. \\
 &\quad \left. - \frac{\delta W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_\alpha(x_1)} i \int d^4y \bar{\eta}_\varepsilon S_{\varepsilon\gamma}(y-x_3) i S_{\beta\delta}(x_2-x_4) \right] \Big|_{\eta=\bar{\eta}=0} \\
 &\quad + i^2 \left[ W_0[\eta, \bar{\eta}] S_{\beta\gamma}(x_2-x_3) S_{\alpha\delta}(x_1-x_4) \right. \\
 &\quad \left. - W_0[\eta, \bar{\eta}] S_{\alpha\gamma}(x_1-x_3) S_{\beta\delta}(x_2-x_4) \right] \Big|_{\eta=\bar{\eta}=0} \\
 &= S_{\alpha\gamma}(x_1-x_3) S_{\beta\delta}(x_2-x_4) - S_{\beta\gamma}(x_2-x_3) S_{\alpha\delta}(x_1-x_4)
 \end{aligned}$$

In terms of diagrams, one has

$$G_{\alpha\beta\gamma\delta}^{(4)}(x_1, x_2, x_3, x_4) = \begin{array}{c} \begin{array}{cc} 1, \alpha & 2, \beta \\ \bullet & \bullet \\ \swarrow & \searrow \\ \bullet & \bullet \\ 4, \delta & 3, \gamma \end{array} \end{array} - \begin{array}{c} \begin{array}{cc} 1, \alpha & 2, \beta \\ \bullet & \bullet \\ \uparrow & \uparrow \\ \bullet & \bullet \\ 4, \delta & 3, \gamma \end{array} \end{array}$$

where every directed line is a propagator. In the Feynman diagrams, scalars are represented as undirected dashed lines, while fermions are represented as directed solid lines.

## Lecture 4

The different sign can be interpreted using Wick's theorem: the only non-vanishing contributions are the ones where the fields are contracted completely. The minus sign appears in the second diagram since one has to contract the fields when they are not next to each other

$$\langle 0 | \mathcal{T} \{ \psi_1 \psi_2 \bar{\psi}_3 \bar{\psi}_4 \} | 0 \rangle = : \overbrace{\psi_1 \psi_2 \bar{\psi}_3 \bar{\psi}_4} : + : \overbrace{\psi_1 \psi_2 \bar{\psi}_3 \bar{\psi}_4} :$$

This gives a general rule: when performing contractions of spinor fields one has to be careful about the signs.

### 4.3.1 Yukawa theory

See Srednicki, §§45, 51. The Yukawa theory treats the coupling between a massive real scalar field  $\varphi$  and a massive spinor field  $\psi$ . This theory is the analogue of the  $\lambda\varphi^4$  theory for scalar fields.

**Lagrangian.** Let  $M$  be the mass of the boson field and  $m$  the mass of the fermion field. The Yukawa interaction Lagrangian is

$$\mathcal{L}_Y = g\varphi\bar{\psi}\psi$$

[r] One would like to write the simplest Lagrangian. To this end, one may impose a  $U(1)$  symmetry on the spinor field along with Lorentz symmetry and the three discrete symmetries of parity, time reversal and charge conjugation. The scalar terms allowed are all the powers of field with non-negative coupling constant:  $\varphi$ ,  $\varphi^2$ ,  $\varphi^3$  and  $\varphi^4$ . While no further spinor terms are allowed other than the kinetic term, the three scalar interaction terms above are too many (notice that  $\varphi^2$  is a mass term). One may modify the Yukawa interaction

$$\mathcal{L}_Y = ig\varphi\bar{\psi}\gamma_5\psi$$

The imaginary unit makes the interaction Lagrangian real. In fact, consider

$$(\bar{\psi}\gamma_5\psi)^\dagger = (\psi^\dagger\gamma_0\gamma_5\psi)^\dagger = \psi^\dagger\gamma_5^\dagger\gamma_0^\dagger\psi = \psi^\dagger\gamma_5\gamma_0\psi = -\bar{\psi}\gamma_5\psi$$

This term is anti-hermitian and therefore is purely imaginary. Since the scalar field is real, then an imaginary unit is needed to keep the Lagrangian real.

One notices that under a parity transformation, the interaction term changes sign

$$P(\bar{\psi}\gamma_5\psi) = -\bar{\psi}\gamma_5\psi$$

while the kinetic fermionic term is invariant. For the Lagrangian to conserve parity and in particular  $\varphi \rightarrow -\varphi$ , the scalar field  $\varphi$  must be a pseudo-scalar field and the linear and cubic scalar interactions along with their ultraviolet divergences cannot appear:

$$\mathcal{L} = \frac{1}{2}\partial_\mu\varphi\partial^\mu\varphi - \frac{1}{2}M^2\varphi^2 + \bar{\psi}(i\partial\!\!\!/ - m)\psi - \frac{\lambda}{4!}\varphi^4 + ig\varphi\bar{\psi}\gamma_5\psi$$

**Perturbation theory and Feynman rules.** For a weak interacting theory,  $\lambda, g \ll 1$ , one may apply perturbation theory. Trading

$$\bar{\psi} \rightarrow i\delta_\eta, \quad \psi \rightarrow -i\delta_{\bar{\eta}}, \quad \varphi \rightarrow -i\delta_J$$

one can write the generating functional

$$W[J, \eta, \bar{\eta}] = \exp \left\{ i \int d^4x \left[ -\frac{\lambda}{4!}(-i\delta_{J(x)})^4 + ig(-i\delta_{J(x)})(i\delta_{\eta(x)})\gamma_5(-i\delta_{\bar{\eta}(x)}) \right] \right\} W_0[J, \eta, \bar{\eta}]$$

where the free generating functional is

$$\begin{aligned} W_0[J, \eta, \bar{\eta}] &= \int [\mathcal{D}\varphi] \exp \left[ i \int d^4x \left( \frac{1}{2}\partial_\mu\varphi\partial^\mu\varphi - \frac{1}{2}M^2\varphi^2 + J\varphi \right) \right] \\ &\quad \times \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ i \int d^4x \bar{\psi}(i\partial\!\!\!/ - m)\psi + \bar{\psi}\eta + \bar{\eta}\psi \right] \\ &= \exp \left[ i \int d^4x d^4y J(x)\Delta(x-y)J(y) \right] \\ &\quad \times \exp \left[ -i \int d^4x d^4y \bar{\eta}_\alpha(x)S_{\alpha\beta}(x-y)\eta_\beta(y) \right] \end{aligned}$$

and the propagators are

$$\Delta(x-y) = \int \frac{d^4k}{(2\pi)^4} \frac{e^{-ik(x-y)}}{k^2 - M^2 + i\epsilon}, \quad S_{\alpha\beta}(x-y) = \int \frac{d^4k}{(2\pi)^4} \frac{(\not{k} + m)_{\alpha\beta}}{k^2 - m^2 + i\epsilon} e^{-ik(x-y)}$$

By computing the generating functional at one-loop, one finds

$$W[J, \eta, \bar{\eta}] = W_0[J, \eta, \bar{\eta}] + i \int d^4x \left[ -\frac{\lambda}{4!} (-i \delta_{J(x)})^4 + ig(-i \delta_{J(x)})(i \delta_{\eta(x)}) \gamma_5 (-i \delta_{\bar{\eta}(x)}) \right] W_0 + o(\lambda, g)$$

Proceeding as the  $\lambda\varphi^4$  theory, the expression above can be expressed in terms of Feynman diagrams according to the following rules:

- Scalar propagator

$$\Delta(x - y) \equiv x \bullet \text{-----} \bullet y$$

- Fermion propagator

$$S_{\alpha\beta}(x - y) \equiv \psi_\alpha(x) \bullet \text{-----} \bullet \bar{\psi}_\beta(y)$$

Notice the addition of an arrow to denote direction and distinguish between the field  $\psi$  and its Dirac adjoint  $\bar{\psi}$ .

- Internal point

$$ig \int d^4x (\gamma_5)_{\alpha\beta} \equiv \text{-----} \bullet \begin{array}{l} \nearrow \alpha \\ \searrow \beta \end{array}$$

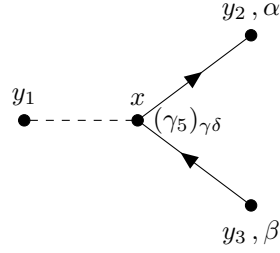
- External scalar point

$$i \int d^4y_j J(y_j) \rightarrow \bullet \text{-----}$$

- External fermionic point

$$i \int d^4y_j \eta_\alpha(y_j) \rightarrow \bullet \text{-----} \bullet, \quad i \int d^4y_j \bar{\eta}_\alpha(y_j) \rightarrow \bullet \text{-----} \bullet$$

**Example.** The tree-level diagram for the Yukawa interaction is

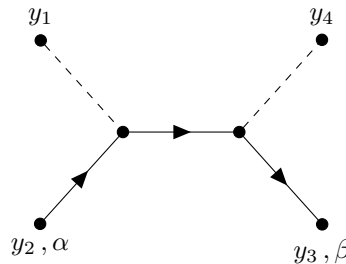


which corresponds to

$$i^3 \int d^4y_1 d^4y_2 d^4y_3 d^4x \bar{\eta}_\alpha(y_2) S_{\alpha\gamma}(y_2 - x) (ig)(\gamma_5)_{\gamma\delta} S_{\delta\beta}(x - y_3) \eta_\beta(y_3) \Delta(x - y_1) J(y_1)$$

[r] Since order matters for spinors, to write the integral without indices one has to go from right to left and follow against the direction of the arrows.

**Example.** With two Yukawa vertices, one can draw



[r] When writing amputated Green's function, the external lines are cut and only the internal part of the diagram is computed.

**Remark.** One neglects vacuum diagrams by choosing a suitable normalization

$$W[J = \eta = \bar{\eta} = 0] = 1$$

**Remark.** One is interested in connected Green's functions because the non-connected diagrams can be written as products of connected ones. The generating functional for connected Green's functions is

$$Z[J, \eta, \bar{\eta}] = \ln W[J, \eta, \bar{\eta}]$$

**Remark.** In the interest of renormalization, one computes amputated Green's functions  $\Gamma^{(n)}$  which are the quantum vertices in the effective action and generated by it.

**Remark.** Computations are done in momentum space.

**Feynman rules in momentum space.** The rules are

- Scalar propagator

$$\frac{1}{k^2 - M^2 + i\varepsilon} \equiv \bullet \text{-----} \bullet$$

- Fermion propagator

$$\frac{(\not{k} + m)_{\alpha\beta}}{k^2 - m^2 + i\varepsilon} \equiv \begin{array}{c} \xleftarrow{k} \\ \bullet \quad \bullet \\ \alpha \quad \beta \end{array}$$

Notice that the momentum arrow follows the line's arrow for particles, but it is in the opposite direction for anti-particles.

- Yukawa vertex

$$ig(\gamma_5)_{\alpha\beta} \equiv \text{-----} \bullet \begin{array}{l} \nearrow \beta \\ \searrow \alpha \end{array}$$

[r]

**Symmetries.** The Lagrangian is invariant under a global U(1) transformation of the fermionic fields

$$\psi' = e^{iq}\psi, \quad \bar{\psi} = e^{-iq}\bar{\psi}, \quad \varphi' = \varphi$$

**Renormalizability.** To study renormalizability one has to look at the mass dimension of the coupling constants. One already knows that  $\lambda$  is dimensionless and  $\dim \varphi = 1$ . From the fermionic mass term, one obtains

$$\dim(\bar{\psi}m\psi) = 4 \implies \dim \psi = \frac{3}{2}$$

The Yukawa coupling is

$$\dim(g\varphi\bar{\psi}\gamma_5\psi) = 4 \implies \dim g = 0$$

Therefore the theory is renormalizable.

By applying the BPHZ renormalization, one finds that the fields are renormalized as

$$\varphi_0 = Z_\varphi^{\frac{1}{2}}\varphi, \quad \psi_0 = Z_\psi^{\frac{1}{2}}\psi, \quad \bar{\psi}_0 = Z_\psi^{\frac{1}{2}}\bar{\psi}$$

The parameters are renormalized as

$$\lambda_0 = Z_\lambda Z_\varphi^{-2}\lambda, \quad g_0 = Z_g Z_\varphi^{-\frac{1}{2}} Z_\psi^{-1}g, \quad M_0^2 = Z_M Z_\varphi^{-1}M^2, \quad m_0^2 = Z_m Z_\psi^{-1}m^2$$

The bare Lagrangian can be split into a renormalized Lagrangian and the counter terms

$$\begin{aligned}\mathcal{L}_0 = & \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi - \frac{1}{2} M^2 \varphi^2 - \frac{\lambda}{4!} \varphi^4 + \bar{\psi} (i \not{\partial} - m) \psi + i g \varphi \bar{\psi} \gamma_5 \psi \\ & + (Z_\varphi - 1) \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi - \frac{1}{2} (Z_M - 1) M^2 \varphi^2 + (Z_\psi - 1) \bar{\psi} i \not{\partial} \psi - (Z_m - 1) m \bar{\psi} \psi \\ & - (Z_\lambda - 1) \frac{\lambda}{4!} \varphi^4 - (Z_g - 1) i g \varphi \bar{\psi} \gamma_5 \psi\end{aligned}$$

The counter terms (second and third line) produce divergent contributions that cancel the divergent contributions given by the renormalized Lagrangian. Therefore, one has to add the counter term vertices to the one already present. The first two addenda of the second line are a propagator



The last two addenda on the second line are another propagator



The third line corresponds to two vertices



The counter terms are fixed in order to cancel the divergences up to a finite part (which is scheme dependent).

## 5 One-loop contributions

One applies power counting to find the divergent contributions. From the  $\lambda\varphi^4$  theory, one already knows that the two-point and four-point scalar Green's function,  $\Gamma_\varphi^{(2)}$  and  $\Gamma_\varphi^{(4)}$ , are divergent. One expects a divergence in the two-point fermionic Green's function  $\Gamma_\psi^{(2)}$  and the three-point Yukawa Green's function  $\Gamma_Y^{(3)}$ .

**Exercise.** Find the superficial degree of divergence for a generic diagram at  $L$  loops with  $n_\varphi$  external scalar lines,  $2n_\psi$  external fermion lines,  $I_\varphi$  internal scalar propagators,  $I_\psi$  internal fermion propagators,  $V_Y$  Yukawa vertices,  $V_\lambda$  scalar  $\lambda$  vertices.

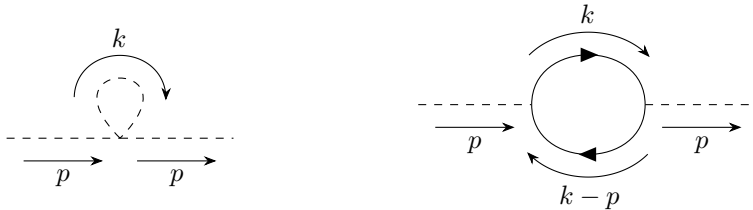
From this one finds that only a particular topology of diagrams is divergent, which are the ones above.

### 5.1 Two-point scalar Green's function

Recall that the two-point scalar Green's function up to one-loop is

$$\Gamma_\varphi^{(2)} = p^2 + m^2 - \Sigma_\varphi, \quad \Gamma_\varphi^{(2)}|_{1L} = -\Sigma_\varphi$$

The one-loop contributions are given by the diagrams



**First diagram.** For the first diagram, one has found

$$\Sigma_\varphi(\text{I}) = -\frac{\lambda M^2}{32\pi^2} \left[ \frac{1}{\varepsilon} + 1 - \gamma - \ln \frac{M^2}{4\pi k^2} + o(\varepsilon^0) \right]$$

[r] where the overall minus is due to working in Minkowski (see Srednicki, eq. 51.21, but remember that their conventions are a spacelike metric,  $D = 4 - \varepsilon$  and  $4\pi k^2 = \mu^2$ ).

## Lecture 5

**Second diagram.** From two Yukawa vertices one has to produce two fermionic propagators through contractions. In one of them, the spinor fields have to be swapped and one gets a negative sign. This is a general rule: alla fermionic loops bring a  $-1$ . Therefore, the one-loop contribution to the two-point scalar Green's function is

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$$\begin{aligned} i\Sigma_\varphi(\text{II}) &= \frac{i^2}{2!} (ig)^2 (-1) \cdot 2 \int \frac{d^4 k}{(2\pi)^4} \frac{(\not{k} + m)_{\delta\alpha}}{k^2 - m^2 + i\varepsilon} (\gamma_5)_{\alpha\beta} \frac{(\not{k} - \not{p})_{\beta\gamma} + m}{(k - p)^2 - m^2 + i\varepsilon} (\gamma_5)_{\gamma\delta} \\ &= -g^2 \int \frac{d^4 k}{(2\pi)^4} \frac{\text{Tr}[(\not{k} + m)\gamma_5(\not{k} - \not{p} + m)\gamma_5]}{(k^2 - m^2 + i\varepsilon)[(k - p)^2 - m^2 + i\varepsilon]} \\ &= -g^2 \int \frac{d^4 k}{(2\pi)^4} \frac{\text{Tr}[(\not{k} + m)(\not{p} - \not{k} + m)]}{(k^2 - m^2 + i\varepsilon)[(k - p)^2 - m^2 + i\varepsilon]} \\ &= -g^2 \int \frac{d^4 k}{(2\pi)^4} \frac{\text{Tr}[\gamma^\mu \gamma^\nu k_\mu (p - k)_\nu + m^2]}{(k^2 - m^2 + i\varepsilon)[(k - p)^2 - m^2 + i\varepsilon]} \\ &= -4g^2 \int \frac{d^4 k}{(2\pi)^4} \frac{k(p - k) + m^2}{(k^2 - m^2 + i\varepsilon)[(k - p)^2 - m^2 + i\varepsilon]} \end{aligned}$$

[r] why 2 in 1st line?

At the penultimate line, one applies

$$\text{Tr}[\gamma^\mu \gamma^\nu] = 4\eta^{\mu\nu}, \quad \text{Tr } I = 4$$

Since the Dirac matrices disappeared, one can operate a Wick rotation to have

$$k_0 = ik_0^E, \quad k^2 = -k_E^2, \quad d^4 k = i d^4 k_E$$

Therefore

$$k(k - p) = k_0(k - p)^0 - \mathbf{k} \cdot (\mathbf{k} - \mathbf{p}) \rightarrow -k(k - p)$$

One could begin from Euclidean space instead and rotate the Dirac matrices and the Clifford algebra.

The one-loop contribution is

$$\begin{aligned} \Sigma_\varphi(\text{II}) &= -4g^2 \int \frac{d^4 k}{(2\pi)^4} \frac{-k(p - k) + m^2}{(k^2 + m^2)[(p - k)^2 + m^2]} \\ &= 4g^2 \int_0^1 dx \int \frac{d^4 q}{(2\pi)^4} \frac{(q + xp)(p - q - xp) - m^2}{(q^2 + D)^2} \\ &= 4g^2 \int_0^1 dx \int \frac{d^4 q}{(2\pi)^4} \frac{-q^2 + (1 - 2x)qp + x(1 - x)p^2 - m^2}{(q^2 + D)^2} \\ &= 4g^2 \int_0^1 dx \int \frac{d^4 q}{(2\pi)^4} \frac{-q^2 + x(1 - x)p^2 - m^2 + D - D}{(q^2 + D)^2} \\ &= 4g^2 \int_0^1 dx \int \frac{d^4 q}{(2\pi)^4} \frac{x(1 - x)p^2 - m^2 + D}{(q^2 + D)^2} - \frac{1}{q^2 + D} \\ &= 4g^2 \int_0^1 dx \int \frac{d^4 q}{(2\pi)^4} \frac{2x(1 - x)p^2}{(q^2 + D)^2} - \frac{1}{q^2 + D} \\ &\equiv \Sigma_2(\text{II}) + \Sigma_1(\text{II}) \end{aligned}$$

At the second line, one has applied Feynman combining

$$x[(p-k)^2 + m^2] + (1-x)(k^2 + m^2) = k^2 + xp^2 - 2xpk + m^2 = (k-xp)^2 + m^2 + x(1-x)p^2$$

and one lets

$$q \equiv k - xp, \quad D \equiv m^2 + x(1-x)p^2$$

At the third line, one notices that linear terms in  $q$  vanish due to parity.

The two integral are divergent since their dimensions are 2 and 0. Using dimensional regularization, the first integral is

$$\begin{aligned} \Sigma_1(\text{II}) &= -4g^2 k^{2\varepsilon} \int_0^1 dx \int \frac{d^n q}{(2\pi)^n} \frac{1}{q^2 + D} = -4g^2 \int_0^1 dx \frac{\Gamma(-1+\varepsilon)}{(4\pi)^{2-\varepsilon}\Gamma(1)} \frac{1}{D^{-1+\varepsilon}} \\ &= -\frac{4g^2 k^{2\varepsilon}}{(4\pi)^{2-\varepsilon}} \Gamma(-1+\varepsilon) \int_0^1 \frac{dx}{[x(1-x)p^2 + m^2]^{-1+\varepsilon}} \\ &= -\frac{4g^2}{(4\pi)^2} (4\pi)^\varepsilon \frac{\Gamma(1+\varepsilon)}{(-1+\varepsilon)} \frac{k^{2\varepsilon}}{\varepsilon} \int_0^1 \frac{dx}{[x(1-x)p^2 + m^2]^{-1+\varepsilon}} \\ &= \frac{4g^2}{(4\pi)^2} (4\pi k^2)^\varepsilon \frac{\Gamma(1+\varepsilon)}{\varepsilon(1-\varepsilon)} \int_0^1 dx \frac{x(1-x)p^2 + m^2}{[x(1-x)p^2 + m^2]^\varepsilon} \\ &= \frac{4g^2}{(4\pi)^2} \frac{\Gamma(1+\varepsilon)}{\varepsilon(1-\varepsilon)} \int_0^1 dx [x(1-x)p^2 + m^2] \exp \left[ -\varepsilon \ln \frac{x(1-x)p^2 + m^2}{4\pi k^2} \right] \\ &= \frac{g^2}{4\pi^2} \frac{1}{\varepsilon} \int_0^1 dx [x(1-x)p^2 + m^2] + o(\varepsilon^{-1}) \\ &= \frac{g^2}{4\pi^2} \frac{1}{\varepsilon} \left[ \frac{p^2}{6} + m^2 \right] + o(\varepsilon^{-1}) \end{aligned}$$

At the first line, one has applied

$$\int \frac{d^n q}{(2\pi)^n} \frac{(q^2)^a}{(q^2 + D)^b} = \frac{1}{D^{b-a-\frac{n}{2}}} \frac{\Gamma(b-a-\frac{n}{2})\Gamma(a+\frac{n}{2})}{(4\pi)^{\frac{n}{2}}\Gamma(b)\Gamma(\frac{n}{2})}$$

with  $a = 0$ ,  $b = 1$  and  $\frac{n}{2} = 2 - \varepsilon$ . At the third line one has applied  $\Gamma(z+1) = z\Gamma(z)$ .

The second integral is

$$\begin{aligned} \Sigma_2(\text{II}) &= 8g^2 \int_0^1 dx \int \frac{d^4 q}{(2\pi)^4} \frac{x(1-x)p^2}{(q^2 + D)^2} \\ &= 8g^2 \frac{k^{2\varepsilon}}{(4\pi)^{2-\varepsilon}} \Gamma(\varepsilon) \int_0^1 dx \frac{x(1-x)p^2}{[x(1-x)p^2 + m^2]^\varepsilon} \\ &= \frac{2g^2}{4\pi^2} \frac{1}{\varepsilon} \int_0^1 dx x(1-x)p^2 + o(\varepsilon^{-1}) \\ &= \frac{2g^2}{4\pi^2} \frac{1}{\varepsilon} \frac{p^2}{6} + o(\varepsilon^{-1}) = \frac{g^2}{4\pi^2} \frac{1}{\varepsilon} \frac{p^2}{3} + o(\varepsilon^{-1}) \end{aligned}$$

At the second line, one has applied the integral formula above with  $a = 1$ ,  $b = 2$  and  $\frac{n}{2} = 2 - \varepsilon$ .

The total one-loop contribution to the two-point scalar Green's function coming from the second diagram is

$$\begin{aligned} \Sigma_\varphi(\text{II}) &= \Sigma_1(\text{II}) + \Sigma_2(\text{II}) = \frac{g^2}{4\pi^2} \frac{1}{\varepsilon} \left[ \frac{p^2}{6} + m^2 \right] + \frac{g^2}{4\pi^2} \frac{1}{\varepsilon} \frac{p^2}{3} + \text{finite} \\ &= \frac{g^2}{4\pi^2} \frac{1}{\varepsilon} \left[ \frac{p^2}{2} + m^2 \right] \end{aligned}$$

Notice that, at one-loop, there is a dependence on the momentum as opposed to the  $\lambda\varphi^4$  theory. The first addendum of the above contribution is cancelled by a counter term of the kinetic term, while the second addendum is added to the first diagram's contribution and is cancelled by a counter term of the mass term.

**Total contribution.** The total contribution at one-loop is

$$\Sigma_\varphi|_{1L} = \Sigma_\varphi(I) + \Sigma_\varphi(II) = \frac{g^2}{8\pi^2} \frac{1}{\varepsilon} p^2 + \frac{1}{\varepsilon} \left[ \frac{g^2}{4\pi^2} m^2 - \frac{\lambda M^2}{32\pi^2} + \text{finite} \right]$$

This total contributing divergence can be cancelled by a counter term of the kinetic term and mass term that give contributions

$$(Z_\varphi - 1) \partial_\mu \varphi \partial^\mu \varphi - (Z_M - 1) M^2 \varphi^2$$

This counter term is computed in configuration space and not momentum space, like the one-loop contribution above. Integrating by parts, one has

$$(Z_\varphi - 1) \partial_\mu \varphi \partial^\mu \varphi = -(Z_\varphi - 1) \varphi \square \varphi$$

Therefore, the counter term contribution in momentum space is

$$(Z_\varphi - 1) p^2 = -(Z_\varphi - 1) p_E^2$$

[r] Wick rotation of the field, do computations? Since one is expanding  $e^{-S}$  then the counter term gives

$$(Z_\varphi - 1) \varphi p_E^2 \varphi$$

Therefore

$$\frac{g^2}{8\pi^2} \frac{1}{\varepsilon} p_E^2 + (Z_\varphi - 1) p_E^2 = \text{finite} \implies \boxed{Z_\varphi|_{1L} = 1 - \frac{g^2}{8\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

The mass term for the scalar field is

$$\left[ \frac{g^2}{4\pi^2} m^2 + \frac{\lambda M^2}{32\pi^2} + \text{finite} \right] \frac{1}{\varepsilon} - (Z_M - 1) M^2 \equiv \text{finite}$$

from which one has

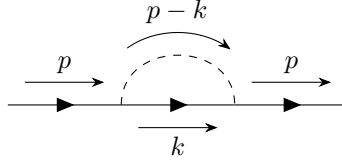
$$\boxed{Z_M|_{1L} = 1 + \left[ \frac{\lambda}{32\pi^2} + \frac{g^2}{4\pi^2} \frac{m^2}{M^2} \right] \frac{1}{\varepsilon} + \text{finite}}$$

## Lecture 6

### 5.2 Two-point spinor Green's function

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Consider the diagram



One has to contract to Yukawa vertices

$$ig\varphi\bar{\psi}_\gamma(\gamma_5)_{\gamma\delta}\psi_\delta, \quad ig\varphi\bar{\psi}_\alpha(\gamma_5)_{\alpha\beta}\psi_\beta$$

to produce a fermion propagator and a scalar propagator (recall that the diagram has to be read in the opposite direction when writing its associated integral). The one-loop contribution to the two-point spinor Green's function is

$$\begin{aligned} [i\Sigma_\psi^{(2)}]_{\alpha\delta} &= \frac{i^2}{2!} (ig)^2 2 \int \frac{d^4 k}{(2\pi)^4} \frac{(\gamma_5)_{\alpha\beta} (\not{k} + m)_{\beta\gamma} (\gamma_5)_{\gamma\delta}}{[(p-k)^2 - M^2 + i\varepsilon](k^2 - m^2 + i\varepsilon)} \\ &= g^2 \int \frac{d^4 k}{(2\pi)^4} \frac{[\gamma_5(\not{k} + m)\gamma_5]_{\alpha\delta}}{[(p-k)^2 - M^2 + i\varepsilon](k^2 - m^2 + i\varepsilon)} \end{aligned}$$



One applies dimensional regularization to have

$$\begin{aligned}
 [i\Sigma_\psi^{(2)}]_{\alpha\delta} &= g^2 k^{2\varepsilon} \int \frac{d^4 k}{(2\pi)^4} \frac{(-\not{k} + m)_{\alpha\delta}}{[(p-k)^2 - M^2](k^2 - m^2)} \\
 &= g^2 k^{2\varepsilon} \int_0^1 dx \int \frac{d^n q}{(2\pi)^n} \frac{(-\not{q} - x\not{p} + m)_{\alpha\delta}}{(q^2 + D)^2} \\
 &= g^2 k^{2\varepsilon} \int_0^1 dx (m - x\not{p})_{\alpha\delta} \int \frac{d^n q}{(2\pi)^n} \frac{1}{(q^2 + D)^2} \\
 &= g^2 k^{2\varepsilon} \frac{1}{(4\pi)^{2-\varepsilon}} \Gamma(\varepsilon) \int_0^1 dx \frac{(m - x\not{p})_{\alpha\delta}}{[x(1-x)p^2 - m^2(1-x) - M^2x]^\varepsilon} \\
 &= \frac{g^2}{16\pi^2} \frac{\Gamma(1+\varepsilon)}{\varepsilon} \int_0^1 dx \frac{(m - x\not{p})_{\alpha\delta}}{[\frac{x(1-x)p^2 - m^2(1-x) - M^2x}{4\pi k^2}]^\varepsilon} \\
 &= \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} \int_0^1 dx (m - x\not{p})_{\alpha\delta} + o(\varepsilon^{-1}) \\
 &= \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} \left[ m - \frac{1}{2}\not{p} \right]_{\alpha\delta}
 \end{aligned}$$

[r] check. At the second line, one has applied Feynman combining

$$\text{Den} = k^2 - 2xkp + xp^2 - m^2(1-x) - M^2x = (k - xp)^2 + x(1-x)p^2 - m^2(1-x) - M^2x$$

and one lets

$$q \equiv k - xp, \quad D \equiv x(1-x)p^2 - m^2(1-x) - M^2x$$

**Renormalization.** The mass contribution is [r]

$$[\Sigma_\psi^{(2)}]_m = -i \frac{g^2}{16\pi^2} \frac{m}{\varepsilon}$$

From the counter term Lagrangian

$$\exp[iS_{\text{ct}}] \rightarrow \exp \left[ i \int d^4x [-(Z_m - 1)m\bar{\psi}\psi] \right]$$

one has

$$-i \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} - i(Z_m - 1) \equiv \text{finite} \implies \boxed{Z_m = 1 - \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

The momentum contribution is

$$[\Sigma_\psi^{(2)}]_p = i \frac{g^2}{32\pi^2} \frac{1}{\varepsilon} \not{p}$$

Therefore

$$e^{iS} \rightarrow \exp \left[ i \int d^4x \bar{\psi} i \not{\partial} \psi (Z_\psi - 1) \right] \rightarrow \exp \left[ -i \int \frac{d^4p}{(2\pi)^4} \bar{\psi}(-p) \not{p} \psi(p) (Z_\psi - 1) \right]$$

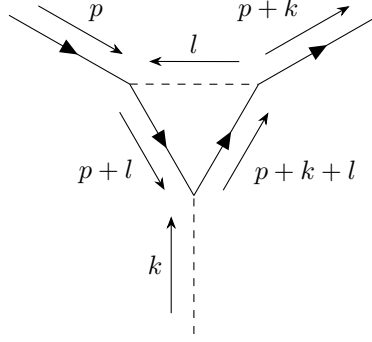
[r]  $\not{p} = -i \not{\partial}$ ?? check Fourier convention

from which

$$i \frac{g^2}{32\pi^2} \frac{1}{\varepsilon} \not{p} - i(Z_\psi - 1)\not{p} = \text{finite} \implies \boxed{Z_\psi = 1 + \frac{g^2}{32\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

### 5.3 Three-point Yukawa Green's function

One would like to compute the one-loop correction to the Yukawa vertex  $\Gamma_Y^{(3)}$ . Consider the following diagram



Let  $p' = p + k$ , the one-loop contribution to the three-point Green's function is

$$\begin{aligned}
 \Gamma_Y^{(3)}|_{1L} &= \frac{i^3}{3!} (ig)^3 (3 \cdot 2) \int \frac{d^4 l}{(2\pi)^4} \frac{(\gamma_5)_{\alpha\beta} (\not{p}' + \not{l} + m)_{\beta\gamma} (\gamma_5)_{\gamma\delta} (\not{p} + \not{l} + m)_{\delta\eta} (\gamma_5)_{\eta\rho}}{(l^2 - M^2 + i\varepsilon)[(p+l)^2 - m^2 + i\varepsilon][(p'+l)^2 - m^2 + i\varepsilon]} \\
 &= -g^3 \int \frac{d^4 l}{(2\pi)^4} \frac{[(\not{p}' + \not{l} + m)(-\not{p} - \not{l} + m)\gamma_5]_{\alpha\rho}}{(l^2 - M^2 + i\varepsilon)[(p+l)^2 - m^2 + i\varepsilon][(p'+l)^2 - m^2 + i\varepsilon]} \\
 &= -g^3 k^{3\varepsilon} \int \frac{d^n l}{(2\pi)^n} \frac{[(\not{p}' + \not{l} + m)(-\not{p} - \not{l} + m)\gamma_5]_{\alpha\rho}}{(l^2 - M^2 + i\varepsilon)[(p+l)^2 - m^2 + i\varepsilon][(p'+l)^2 - m^2]} \\
 &= -g^3 k^{3\varepsilon} \int_0^1 dx_1 dx_2 \int \frac{d^n q}{(2\pi)^n} \frac{(-q^2 \gamma_5 + \tilde{N} + \text{linear in } q)_{\alpha\rho}}{(q^2 + D)^3} \\
 &= g^3 k^{3\varepsilon} \int_0^1 dx_1 dx_2 \int \frac{d^n q}{(2\pi)^n} \frac{q^2}{(q^2 + D)^3} (\gamma_5)_{\alpha\rho} + \text{finite} \\
 &= g^3 k^{3\varepsilon} \frac{1}{(4\pi)^{2-\varepsilon}} \frac{\Gamma(\varepsilon)\Gamma(3-\varepsilon)}{\Gamma(3)\Gamma(2-\varepsilon)} \int_0^1 dx_1 dx_2 D^{-\varepsilon} (\gamma_5)_{\alpha\rho} + \text{finite} \\
 &= \frac{g^3 k^{3\varepsilon}}{16\pi^2} (4\pi k^2)^\varepsilon \frac{1}{\varepsilon} \frac{\Gamma(1-\varepsilon)\Gamma(3-\varepsilon)}{\Gamma(3)\Gamma(2-\varepsilon)} \int_0^1 dx_1 dx_2 D^{-\varepsilon} (\gamma_5)_{\alpha\rho} + \text{finite} \\
 &= \frac{g^3}{16\pi^2} \frac{1}{\varepsilon} (\gamma_5)_{\alpha\rho} + o(\varepsilon^{-1})
 \end{aligned}$$

[r] At the fourth line, one has applied Feynman combining. The denominator is

$$\text{Den} = (1 - x_1 x_2)(l^2 - M^2) + x_1[(l + p)^2 - m^2] + x_2[(l + p')^2 - m^2] \equiv q^2 + D$$

where one has

$$q \equiv l + x_1 p + x_2 p'$$

$$D \equiv -(1 - x_1 - x_2)M^2 - (x_1 + x_2)m^2 + x_1(1 - x_1)p^2 + x_2(1 - x_2)p'^2 - 2x_1 x_2 p p'$$

The numerator is

$$\begin{aligned}
 \text{Num} &= (-\not{q} + x_1 \not{p} + x_2 \not{p}' - \not{p} + m)(\not{q} - x_1 \not{p} - x_2 \not{p}' + \not{p}' + m)\gamma_5 \\
 &= -\not{q}\gamma_5 + [x_2 \not{p}' - (1 - x_1)\not{p} + m][(1 - x_2)\not{p}' - x_1 \not{p} + m]\gamma_5 + (\text{linear in } q) \\
 &= -q^2 \gamma_5 + \tilde{N} + (\text{linear in } q)
 \end{aligned}$$

where one has

$$\not{q}\not{q} = \gamma^\mu q_\mu \gamma^\nu q_\nu = \frac{1}{2} q_\mu q_\nu \{\gamma^\mu, \gamma^\nu\} = \frac{1}{2} q_\mu q_\nu 2\eta^{\mu\nu} = q^2$$

**Renormalization.** From the exponential, one gets the counter term

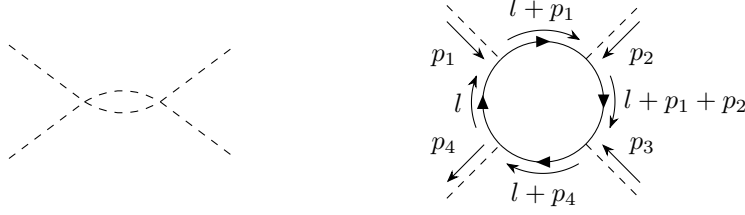
$$e^{iS} \rightarrow \exp \left[ i \int d^4 x i g (Z_g - 1) \varphi \bar{\psi} \gamma_5 \psi \right] \rightarrow -g (Z_g - 1) \varphi \bar{\psi} \gamma_5 \psi$$

Therefore

$$\frac{g^3}{16\pi^2} \frac{1}{\varepsilon} - g(Z_g - 1) \equiv \text{finite} \implies \boxed{Z_g = 1 + \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

## 5.4 Four-point Green's function

The vertex function for the  $\lambda$ -vertex gets two contributions:



In principle, the Yukawa theory contains the coupling  $g$ , but the theory is renormalizable only if one includes the  $\lambda\varphi^4$  term. [r] In fact, the second diagram produces a divergent term proportional to  $\lambda^4$ .

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**Second diagram.** The integral associated with the diagram is

$$\begin{aligned} i\Gamma_{\varphi}^{(4)}(\text{II}) &= \frac{i^4}{4!} (ig)^4 4! (-1) \\ &\times \int \frac{d^4 l}{(2\pi)^4} \frac{\text{Tr}[(\not{l} + m)\gamma_5(\not{l} + \not{p}_4 + m)\gamma_5(\not{l} + \not{p}_1 + \not{p}_2 + m)\gamma_5(\not{l} + \not{p}_1 + m)\gamma_5]}{(l^2 - m^2)[(l + p_4)^2 - m^2][(l + p_1 + p_2)^2 - m^2][(l + p_1)^2 - m^2]} \\ &+ 5 \text{ permutations of } (k_2, k_3, k_4) \end{aligned}$$

where  $i\epsilon$  is understood in all propagators. Since one is interested only in the divergent part, one selects only the  $l^4$  term in the numerator. One may set all external momenta to zero  $p_1 = p_2 = p_3 = p_4 = 0$  since the integral has superficial degree of divergence equal to zero and so it does not depend on momentum. The trace is

$$\text{Tr}[\not{l}\not{l}\not{l}\not{l}] = l_{\mu}l_{\nu}l_{\rho}l_{\sigma} \text{Tr}[\gamma^{\mu}\gamma^{\nu}\gamma^{\rho}\gamma^{\sigma}] = 4[l^4 - l^4 + l^4] = 4l^4$$

Therefore

$$i\Gamma_{\varphi}^{(4)}(\text{II}) = -4g^4 \cdot 6 \int \frac{d^4 l}{(2\pi)^4} \frac{l^4}{(l^2 - m^2)^4} + \text{finite} = -\frac{3g^4}{2\pi^2} \frac{1}{\epsilon} + \text{finite}$$

The coefficient 6 comes from the inequivalent permutations of the momenta of the diagram.

**Total contribution.** Recalling that the first diagram gives

$$i\Gamma_{\varphi}^{(4)}(\text{I}) = \frac{3\lambda^2}{32\pi^2} \frac{1}{\epsilon} + \text{finite}$$

[r] then one finds

$$Z_{\lambda} = 1 + \left[ \frac{3\lambda}{32\pi^2} - \frac{3}{2\pi^2} \frac{g^4}{\lambda} \right] \frac{1}{\epsilon} + \text{finite}$$

## 6 Renormalization group

See Srednicki, §52.

**Renormalization functions.** The renormalization functions at one-loop are

$$\begin{aligned} Z_{\varphi} &= 1 - \frac{g^2}{8\pi^2} \frac{1}{\epsilon} + \text{finite} \\ Z_{\psi} &= 1 - \frac{g^2}{32\pi^2} \frac{1}{\epsilon} + \text{finite} \\ Z_g &= 1 + \frac{g^2}{16\pi^2} \frac{1}{\epsilon} + \text{finite} \\ Z_{\lambda} &= 1 + \left[ \frac{3\lambda}{32\pi^2} - \frac{3}{2\pi^2} \frac{g^4}{\lambda} \right] \frac{1}{\epsilon} + \text{finite} \end{aligned}$$

[r] wrong sign  $Z_\psi$ , affects following argument  
 The bare coupling constant are

$$g_0 = Z_g Z_\varphi^{-\frac{1}{2}} Z_\psi^{-1} g k^\varepsilon, \quad \lambda_0 = Z_\lambda Z_\varphi^{-2} \lambda k^{2\varepsilon}$$

Using the minimal subtraction scheme, one writes the renormalization functions as

$$Z = 1 + \sum_n \frac{a_n(g, \lambda)}{\varepsilon^n}$$

Therefore, the coupling constants are

$$g_0 = g k^\varepsilon \left[ 1 + \sum_{n=1}^{\infty} \frac{\tilde{G}_n(g, \lambda)}{\varepsilon^n} \right], \quad \lambda_0 = \lambda k^{2\varepsilon} \left[ 1 + \sum_{n=1}^{\infty} \frac{\tilde{L}_n(g, \lambda)}{\varepsilon^n} \right]$$

Comparing these series with the above expressions for the coupling constants, one sees that the brackets are the products of the renormalization functions

$$\begin{aligned} 1 + \sum_{n=1}^{\infty} \frac{\tilde{G}_n(g, \lambda)}{\varepsilon^n} &= Z_g Z_\varphi^{-\frac{1}{2}} Z_\psi^{-1} \\ &= \left[ 1 + \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} + \dots \right] \left[ 1 + \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} + \dots \right] \left[ 1 + \frac{g^2}{32\pi^2} \frac{1}{\varepsilon} + \dots \right] \\ &= 1 + \frac{5g^2}{32\pi^2} \frac{1}{\varepsilon} + \dots \end{aligned}$$

where the ellipses contain higher-order corrections in  $g$  and  $\lambda$  (and thus higher powers of  $\varepsilon^{-1}$ ).  
 Therefore

$$\tilde{G}_1(g, \lambda)|_{1L} = \frac{5g^2}{32\pi^2}$$

Similarly

$$\tilde{L}_1(g, \lambda)|_{1L} = \frac{3\lambda}{32\pi^2} + \frac{g^2}{4\pi^2} - \frac{3}{2\pi^2} \frac{g^4}{\lambda}$$

One notices that  $\tilde{G}_1$  depends only on one coupling constant.

**Beta functions.** Recall that the beta functions are

$$\beta_g = k \, d_k g, \quad \beta_\lambda = k \, d_k \lambda$$

The bare Yukawa coupling constant can be rewritten as

$$\begin{aligned} g_0 &= g k^\varepsilon \left[ 1 + \sum_{n=1}^{\infty} \frac{\tilde{G}_n(g, \lambda)}{\varepsilon^n} \right] = g e^{\varepsilon \ln k} \exp \left[ \sum_{n=1}^{\infty} \frac{G_n(g, \lambda)}{\varepsilon^n} \right] \\ \lambda_0 &= \lambda k^{2\varepsilon} \left[ 1 + \sum_{n=1}^{\infty} \frac{\tilde{L}_n(g, \lambda)}{\varepsilon^n} \right] = \lambda e^{2\varepsilon \ln k} \exp \left[ \sum_{n=1}^{\infty} \frac{L_n(g, \lambda)}{\varepsilon^n} \right] \end{aligned}$$

Taking the logarithm, one has

$$\begin{aligned} \ln g_0 &= \ln g + \varepsilon \ln k + \sum_{n=1}^{\infty} \frac{G_n(g, \lambda)}{\varepsilon^n} \\ \ln \lambda_0 &= \ln \lambda + 2\varepsilon \ln k + \sum_{n=1}^{\infty} \frac{L_n(g, \lambda)}{\varepsilon^n} \end{aligned}$$

The right-hand side depends on the mass scale  $k$ , but not the left-hand side. Applying  $k \, d_k$ , one has

$$0 = k \, d_k \ln g + \varepsilon + \sum_{n=1}^{\infty} \frac{1}{\varepsilon^n} [\partial_g G_n k \, \partial_k g + \partial_\lambda G_n k \, \partial_k \lambda]$$

Multiplying by the coupling constant  $g$ , one has

$$0 = g\varepsilon + k \, d_k g + \sum_{n=1}^{\infty} \frac{1}{\varepsilon^n} [g \partial_g G_n k \partial_k g + g \partial_\lambda G_n k \partial_k \lambda]$$

This expression contains the beta function. Since the beta functions are finite in the limit  $\varepsilon \rightarrow 0$ , one may write

$$k \, d_k g = \beta_g(g, \lambda) - \varepsilon g, \quad k \, d_k \lambda = \beta_\lambda(g, \lambda) - 2\varepsilon \lambda$$

[r] Substituting this inside the expression above, one obtains

$$\begin{aligned} 0 &= \beta_g + \sum_{n=1}^{\infty} \frac{1}{\varepsilon^n} [g \partial_g G_n (-g\varepsilon + \beta_g) + g \partial_\lambda G_n (-2\varepsilon \lambda + \beta_\lambda)] \\ &= \beta_g - g^2 \partial_g G_1 - 2g\lambda \partial_\lambda G_1 + 0 \cdot \text{poles in } \frac{1}{\varepsilon} \end{aligned}$$

From this one has

$$\boxed{\beta_g(g, \lambda) = g[g \partial_g G_1 + 2\lambda \partial_\lambda G_1]}$$

Similarly, one has

$$\boxed{\beta_\lambda = \lambda[g \partial_g L_1 + 2\lambda \partial_\lambda L_1]}$$

Since  $\tilde{G}_1 = G_1$  and  $\tilde{L}_1 = L_1$ , then the one-loop beta functions are

$$\beta_g|_{1L} = \frac{5}{16\pi^2} g^3, \quad \beta_\lambda|_{1L} = \frac{1}{16\pi^2} (3\lambda^2 + 8g^2\lambda - 48g^4)$$

## 6.1 Fixed points and RG flow

One introduces a new coupling  $\rho$  such that  $\lambda = \rho g^2$  and studies the flow of  $g$  and  $\rho$ . The beta function is

$$\beta_\rho = k \, d_k \rho = k \, d_k \frac{\lambda}{g^2} = \frac{1}{g^2} \beta_\lambda - \frac{2\lambda}{g^3} \beta_g$$

At one-loop, one finds

$$\beta_\rho|_{1L} = \frac{g^2}{16\pi^2} (3\rho^2 - 2\rho - 48)$$

This expression is a product of polynomials of the coupling constants as opposed to the beta function of  $\lambda$  that is a sum of polynomials.

The only fixed point of the flow of  $g$  is

$$\beta_g = 0 \implies g = 0$$

In order to study the flow, one considers small perturbations away from the fixed point,  $|g| \ll 1$ . The fixed points for  $\rho$  are

$$\beta_\rho = 0 \implies \rho_{1,2} = \frac{1 \pm \sqrt{145}}{3}$$

The first is positive and the second is negative. The small perturbations around the solutions  $\rho_j$  can be studied by Taylor expanding

$$k \, d_k \rho = \beta_\rho = \beta_\rho(\rho_1) + \beta'_\rho(\rho_1)(\rho - \rho_1) + o(\Delta\rho) = \beta'_\rho(\rho_1)(\rho - \rho_1) + o(\Delta\rho)$$

From the sign of the derivative, one obtains the flow. If the derivative is positive, then for  $\rho > \rho_j$  the derivative of the coupling with respect to the energy is positive, so the energy is increasing, while for  $\rho < \rho_j$ , the derivative of the coupling is negative so the energy decreases; vice versa, if the derivative is negative, then for  $\rho > \rho_j$ , the energy decreases, while for  $\rho < \rho_j$  it increases. The derivative is

$$\beta'_\rho = \frac{g^2}{16\pi^2} (6\rho - 2)$$

from which

$$\beta'_\rho(\rho_1) = \frac{g^2}{8\pi^2} \sqrt{145} > 0, \quad \beta'_\rho(\rho_2) = -\frac{g^2}{8\pi^2} \sqrt{145} < 0$$

[r] The flow is away from  $\rho_1$  and towards  $\rho_2$ . The first is an infrared-stable fixed point while the second is an ultraviolet-stable fixed point.

**Trajectories.** One may study the trajectory of the theory in the  $(\rho, g)$  plane. To this end, one needs to find  $g(\rho)$ . One has to solve

$$d_\rho g = d_{\ln k} g d_\rho \ln k = \frac{\beta_g}{\beta_\rho} = \frac{5}{16\pi^2} g^3 \left[ \frac{g^2}{16\pi^2} 3(\rho - \rho_1)(\rho - \rho_2) \right]^{-1} = \frac{5g}{3(\rho - \rho_1)(\rho - \rho_2)}$$

Therefore

$$\frac{1}{g} d_\rho g = d_\rho \ln g = \frac{5}{3(\rho - \rho_1)(\rho - \rho_2)}$$

after integrating one gets

$$\ln \frac{g(\rho)}{g(0)} = \frac{5}{3} \int_0^\rho \frac{dx}{(x - \rho_1)(x - \rho_2)} = \frac{5}{3(\rho_1 - \rho_2)} \ln \left| \frac{\rho - \rho_1}{\rho - \rho_2} \frac{\rho_2}{\rho_1} \right|$$

from which

$$g(\rho) = g(0) \left| \frac{\rho - \rho_1}{\rho - \rho_2} \right|^{\frac{5}{3(\rho_1 - \rho_2)}} = g_0 \left| \frac{\rho - \rho_1}{\rho - \rho_2} \right|^{\frac{5}{3(\rho_1 - \rho_2)}}$$

Due to the absolute value, the constant-sign domains are

$$\rho < \rho_2, \quad \rho_2 < \rho < \rho_1, \quad \rho > \rho_1$$

[r] diagr. Let

$$\nu = \frac{5}{3(\rho_1 - \rho_2)} > 0$$

then the derivatives are

$$g'(\rho) = g_0 \nu \left| \frac{\rho - \rho_1}{\rho - \rho_2} \right|^{\nu-1} \frac{1}{(\rho - \rho_2)^2} \begin{cases} \rho_1 - \rho_2 > 0, & \rho < \rho_2 \vee \rho > \rho_1 \\ \rho_2 - \rho_1 < 0, & \rho_2 < \rho < \rho_1 \end{cases}$$

[r] With two coupling constants, it is interesting to study how the coupling constants affect each other.

**Exercise.** Compute the anomalous dimensions. Recall that

$$\gamma_\varphi = \frac{1}{2} k d_k \ln Z_\varphi, \quad \gamma_\psi = \frac{1}{2} k d_k \ln Z_\psi$$

**Solution.** Recalling that up to one-loop one has

$$Z_\varphi = 1 - \frac{g^2}{8\pi^2} \frac{1}{\varepsilon} \implies \ln Z_\varphi = \sum_{n=1}^{\infty} \frac{F_n(g, \lambda)}{\varepsilon^n} \implies F_1(g, \lambda) = -\frac{g^2}{8\pi^2}$$

for the scalar field follows

$$\begin{aligned} \gamma_\varphi &= \frac{1}{2} k d_k \ln Z_\varphi = \frac{1}{2} [(d_g \ln Z_\varphi)(d_{\ln k} g) + (d_\lambda \ln Z_\varphi)(d_{\ln k} \lambda)] \\ &= \frac{1}{2} [(d_g \ln Z_\varphi)(\beta_g - \varepsilon g) + (d_\lambda \ln Z_\varphi)(\beta_\lambda - 2\varepsilon \lambda)] \\ &= \frac{1}{2} (d_g \ln Z_\varphi)(\beta_g - \varepsilon g) + 0 = \frac{1}{2} d_g \left[ -\frac{g^2}{8\pi^2} \frac{1}{\varepsilon} + \dots \right] (\beta_g - \varepsilon g) \\ &= \frac{g^2}{8\pi^2} + \dots \end{aligned}$$

where the ellipses are higher powers of  $\varepsilon^{-1}$ .

For the spinor field, one has

$$Z_\psi = 1 - \frac{g^2}{32\pi^2} \frac{1}{\varepsilon} \implies P_1(g, \lambda) = -\frac{g^2}{32\pi^2}$$

The procedure is exactly the same

$$\begin{aligned}
 \gamma_\psi &= \frac{1}{2} k \, d_k \ln Z_\psi = \frac{1}{2} [(d_g \ln Z_\psi)(d_{\ln k} g) + (d_\lambda \ln Z_\psi)(d_{\ln k} \lambda)] \\
 &= \frac{1}{2} [(d_g \ln Z_\psi)(\beta_g - \varepsilon g) + (d_\lambda \ln Z_\psi)(\beta_\lambda - 2\varepsilon \lambda)] \\
 &= \frac{1}{2} (d_g \ln Z_\psi)(\beta_g - \varepsilon g) + 0 = \frac{1}{2} d_g \left[ -\frac{g^2}{32\pi^2} \frac{1}{\varepsilon} + \dots \right] (\beta_g - \varepsilon g) \\
 &= \frac{g^2}{32\pi^2} + \dots
 \end{aligned}$$

## Lecture 8

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## 7 Functional determinants

See Srednicki, §53. Functional determinants are useful to motivate the use of an infinite sum of Feynman diagrams.

**Example.** Consider a complex scalar field  $\chi$  coupled to a background real scalar field  $\varphi$

$$\mathcal{L} = \partial_\mu \bar{\chi} \partial^\mu \chi - m^2 \bar{\chi} \chi + g \varphi \bar{\chi} \chi$$

The generating functional with no source is

$$e^{i\Gamma[\varphi]} = \int [\mathcal{D}\chi \mathcal{D}\bar{\chi}] e^{i \int d^4x \mathcal{L}}$$

where  $\Gamma[\varphi]$  is the effective action for the field  $\varphi$ . The action is quadratic in the complex field and one may rewrite

$$\begin{aligned}
 \int d^4x \mathcal{L} &= \int d^4x [-\bar{\chi}(\Box + m^2)\chi + g\varphi \bar{\chi} \chi] \\
 &= - \int d^4x d^4y \bar{\chi}(x) [(\Box + m^2) - g\varphi(x)] \delta^{(4)}(x - y) \chi(y) \\
 &= - \int d^4x d^4y \bar{\chi}(x) M(x, y) \chi(y)
 \end{aligned}$$

where the differential operator is

$$M(x, y) \equiv [(\Box + m^2) - g\varphi(x)] \delta^{(4)}(x - y)$$

The generating functional is

$$e^{i\Gamma[\varphi]} = \int [\mathcal{D}\chi \mathcal{D}\bar{\chi}] \exp \left[ -i \int d^4x \int d^4y \bar{\chi}(x) M(x, y) \chi(y) \right]$$

One recognizes the functional generalization of the Gaussian integral

$$\int d^n z d^n \bar{z} e^{-i \bar{z}_i M_{ij} z_j} \propto \frac{1}{\det M}$$

One may rewrite the differential operator as

$$M(x, y) = \int d^4z [(\Box + m^2) \delta^{(4)}(x - z)] [\delta^{(4)}(z - y) - g \Delta(z - y) \varphi(y)] = \int d^4z M_0(x, z) \tilde{M}(z, y)$$

where

$$M_0(x, z) \equiv (\Box + m^2) \delta^{(4)}(x - z), \quad \tilde{M}(z, y) \equiv \delta^{(4)}(z - y) - g \Delta(z - y) \varphi(y)$$

and

$$\Delta(x) = \frac{1}{\Box + m^2}$$

One may notice that the operator  $M_0$  is the operator  $M$  with no background,  $\varphi = 0$ , and the structure of the operator  $\tilde{M}$  is  $\tilde{M} = I - G$ .

In general, the determinant of a linear operator is the product of its eigenvalues. The determinant also satisfies the typical matrix relations [r]

$$\det M = \det (M_0 \tilde{M}) = (\det M_0)(\det \tilde{M})$$

Therefore, the generating functional is

$$e^{i\Gamma[\varphi]} \propto \frac{1}{(\det M_0)(\det \tilde{M})} \propto \frac{1}{\det \tilde{M}}$$

Since the operator  $M_0$  does not depend on the background field, then its determinant  $\det M_0$  is a constant that can be absorbed into the normalization constant. Knowing that

$$\det \tilde{M} = e^{\text{Tr} \ln \tilde{M}}$$

the generating functional is

$$e^{i\Gamma[\varphi]} \propto e^{-\text{Tr} \ln \tilde{M}} = e^{-\text{Tr} \ln (I - G)} = \exp \left[ \text{Tr} \sum_{n=1}^{\infty} \frac{G^n}{n} \right]$$

where one Taylor-expands the logarithm assuming that the theory is weakly interacting,  $g \ll 1$ . Recalling that  $\text{Tr}(A + B) = \text{Tr} A + \text{Tr} B$ , then one may write

$$\text{Tr} G^n = g^n \int \left[ \prod_{j=1}^n d^4 x_j \right] \Delta(x_1 - x_2) \varphi(x_2) \Delta(x_2 - x_3) \varphi(x_3) \cdots \Delta(x_n - x_1) \varphi(x_1)$$

This trace is represented as a loop with  $n$  external  $\varphi$  lines and  $n$  internal  $\chi$  propagators. Therefore the effective action is

$$i\Gamma[\varphi] \propto \sum_{n=1}^{\infty} \frac{\text{Tr} G^n}{n}$$

This is the typical expansion in powers of the coupling constant  $g$  of Feynman diagrams.

**Example.** Consider a Dirac spinor field  $\psi$  with a Yukawa interaction and a background real scalar field  $\varphi$

$$\mathcal{L} = \bar{\psi}(i\partial - m)\psi + g\varphi\bar{\psi}\psi$$

The generating functional is

$$e^{i\Gamma[\varphi]} = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] e^{i \int d^4 x \mathcal{L}}$$

The action can be rewritten as

$$\begin{aligned} \int d^4 x \mathcal{L} &= \int d^4 x [\bar{\psi}(i\partial - m)\psi + g\varphi\bar{\psi}\psi] \\ &= - \int d^4 x d^4 y \bar{\psi}_\alpha(x) [(-i\partial + m)_{\alpha\beta} - g\varphi(x)\delta_{\alpha\beta}] \delta^{(4)}(x - y) \psi_\beta(y) \\ &= - \int d^4 x d^4 y \bar{\psi}_\alpha(x) M_{\alpha\beta}(x, y) \psi_\beta(y) \end{aligned}$$

where one has

$$M_{\alpha\beta}(x, y) \equiv (-i\partial + m)_{\alpha\beta} - g\varphi(x)\delta_{\alpha\beta}$$

The generating functional becomes

$$e^{i\Gamma[\varphi]} = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ -i \int d^4 x d^4 y \bar{\psi} M \psi \right] \propto \det M$$

recalling that a Grassmann-odd Gaussian integral is proportional to the determinant and not its inverse.



The differential operator may be rewritten as

$$\begin{aligned} M_{\alpha\beta}(x, y) &= \int d^4z [(-i\not{\partial} + m)_{\alpha\gamma}\delta^{(4)}(x - y)][\delta_{\gamma\beta}\delta^{(4)}(y - z) - g\varphi(z)S_{\gamma\beta}(y - z)] \\ &= \int d^4z M_0(x, z)_{\alpha\gamma}\tilde{M}(z, y)_{\gamma\beta} \end{aligned}$$

where one defines

$$\begin{aligned} M_0(x, z)_{\alpha\gamma} &\equiv (-i\not{\partial} + m)_{\alpha\gamma}\delta^{(4)}(x - y) \\ \tilde{M}(z, y)_{\gamma\beta} &\equiv \delta_{\gamma\beta}\delta^{(4)}(y - z) - g\varphi(z)S_{\gamma\beta}(y - z) = I - \tilde{G} \end{aligned}$$

Like the previous example, the generating functional is

$$e^{i\Gamma[\varphi]} \propto \det M_0 \det \tilde{M} \propto \det \tilde{M} = e^{\text{Tr} \ln \tilde{M}} = e^{\text{Tr} \ln (I - \tilde{G})} = \exp \left[ - \sum_{n=1}^{\infty} \frac{\text{Tr} \tilde{G}^n}{n} \right]$$

where one has

$$\text{Tr} \tilde{G}^n = g^n \int \left[ \prod_{j=1}^n d^n x_j \right] S_{\alpha_1\alpha_2}(x_1 - x_2)\varphi(x_2)S_{\alpha_2\alpha_3}(x_2 - x_3)\varphi(x_3) \cdots S_{\alpha_n\alpha_1}(x_n - x_1)\varphi(x_1)$$

The corresponding diagram is again a loop with  $n$  external scalar lines and  $n$  internal spinor propagators. Notice that the negative sign associated with fermion loops is the one present in the exponential.

The effective action is

$$i\Gamma[\varphi] \propto - \sum_{n=1}^{\infty} \frac{\text{Tr} \tilde{G}^n}{n}$$

**Conclusion.** The scalar effective action  $\Gamma[\varphi]$  is given by an infinite sum of one-loop diagrams with external scalar lines.

## 8 Discrete symmetries

See Srednicki, §§2, 23, 40. There are many theories forbidden by parity, time reversal and charge conjugation.

**Lorentz group.** The Lorentz group is

$$O(1, 3) = \{ \Lambda \in \text{GL}(4, \mathbb{R}) \mid \Lambda^\top \eta \Lambda = \eta \} = \mathcal{L}$$

where the Minkowski metric  $\eta$  is timelike. The above condition is

$$\Lambda^\mu{}_\nu \eta_{\mu\rho} \Lambda^\rho{}_\sigma = \eta_{\nu\sigma}$$

Relativistic field theories have to be invariant under Lorentz changes of coordinates

$$x'^\mu = \Lambda^\mu{}_\nu x^\nu$$

From the definition of the group, it follows

$$\det \Lambda = \pm 1$$

and

$$\Lambda^\mu{}_0 \eta_{\mu\nu} \Lambda^\nu{}_0 = 1 \implies (\Lambda^0{}_0)^2 \geq 1$$

In this way, the Lorentz group can be split into four disjoint subsets, but only one is a subgroup<sup>10</sup>

$$\begin{aligned}\mathcal{L}_+^\uparrow &= \{\Lambda \in O(1,3) \mid \det \Lambda = 1, \Lambda^0_0 \geq 1\} \\ \mathcal{L}_+^\downarrow &= \{\Lambda \in O(1,3) \mid \det \Lambda = 1, \Lambda^0_0 \leq -1\} \\ \mathcal{L}_-^\uparrow &= \{\Lambda \in O(1,3) \mid \det \Lambda = -1, \Lambda^0_0 \geq 1\} \\ \mathcal{L}_-^\downarrow &= \{\Lambda \in O(1,3) \mid \det \Lambda = -1, \Lambda^0_0 \leq -1\}\end{aligned}$$

The first one is a group and it is generated by proper rotations and Lorentz boosts. Since it is a subgroup it contains the identity. Furthermore, it is continuously connected to it: transformations can be written infinitesimally close to the identity.

The other subsets contain discrete transformations. One such transformation is parity

$$\Lambda_P = \text{diag}(1, -1, -1, -1) \in \mathcal{L}_-^\uparrow$$

It mirrors all spatial coordinates. Another is time reversal

$$\Lambda_T = \text{diag}(-1, 1, 1, 1) \in \mathcal{L}_-^\downarrow$$

It mirrors time. The combination of these two gives

$$\Lambda_{PT} = \text{diag}(-1, -1, -1, -1) \in \mathcal{L}_+^\downarrow$$

The set of the transformations

$$I = \{1, \Lambda_P, \Lambda_T, \Lambda_{PT}\}$$

is a subgroup of the Lorentz group.

Any transformation in the three non-group subsets of the Lorentz group can be obtained by combining a proper Lorentz transformation with a transformation in group  $I$ .

One would like to study how the fields (which are finite-dimensional, non-unitary irreducible representations of the proper Lorentz group) transform under discrete symmetries.

## 8.1 Representation on scalar fields

Scalar fields are the representation with spin  $s = 0$ . The transformation of a scalar field under a proper<sup>11</sup> Lorentz transformation is

$$U^{-1}(\Lambda)\varphi(x)U(\Lambda) = \varphi(\Lambda^{-1}x), \quad U(\Lambda) = e^{\frac{i}{2}\varepsilon^{\mu\nu}L_{\mu\nu}}$$

where  $L_{\mu\nu}$  are the generators Lorentz group.

Similarly, a discrete transformation has a representation in terms of the field.

**Parity.** For parity, one has

$$U^{-1}(\Lambda_P)\varphi(x)U(\Lambda_P) = \eta_P\varphi(\Lambda_P x)$$

where  $\eta_P$  is a constant. Applying twice a parity transformation, one gets unity

$$U^{-2}(\Lambda_P)\varphi(x)U^2(\Lambda_P) = \eta_P U^{-1}(\Lambda_P)\varphi(\Lambda_P x)U(\Lambda_P x) = \eta_P^2 \varphi(\Lambda_P^2 x) = \eta_P^2 \varphi(x) \equiv \varphi(x)$$

Therefore, one obtains

$$\eta_P^2 = 1 \implies \eta_P = \pm 1$$

For  $\eta_P = 1$ , the field is parity-even and is a scalar field<sup>12</sup>, while  $\eta_P = -1$  the field is parity-odd and is a pseudo-scalar field.

<sup>10</sup>Called restricted proper Lorentz group, cf. Mathematical Methods for Physics.

<sup>11</sup>For this section, the word “proper” includes the adjective restricted when referring to Lorentz transformations.

<sup>12</sup>Notice that pseudo-tensors transform like tensors under proper rotations, but additionally change sign under improper rotations. For example, under improper rotations, scalars do not change sign, while pseudo-scalars do; vectors change sign, while pseudo-vectors do not.

**Time reversal.** For time reversal, one has

$$U^{-1}(\Lambda_T)\varphi(x)U(\Lambda_T) = \eta_T\varphi(\Lambda_T x)$$

where  $\eta_T$  is a constant. Like before, applying twice the transformation, one gets unity. Thus

$$\eta_T = \pm 1$$

The field can be time reversal-even or time reversal-odd.

**Physical constraints.** To decide whether the scalar fields are even or odd under a discrete symmetry, one utilizes the physical input: the Lagrangian of a scalar field is parity- and time reversal-even.

### 8.1.1 Unitarity

The operator associated to parity is unitary

$$U^\dagger(\Lambda_P)U(\Lambda_P) = U(\Lambda_P)U^\dagger(\Lambda_P) = 1$$

while the operator associated to time reversal is anti-unitary

$$\langle Ux|Uy \rangle = \langle x|U^\dagger U|y \rangle \equiv \langle x|y \rangle^*$$

or equivalently

$$U^{-1}(\Lambda_T)iU(\Lambda_T) = -i$$

**Motivation.** Consider a proper Lorentz transformation,  $\Lambda \in \mathcal{L}_+^\uparrow$ , and a four-momentum  $p^\mu$ . Under such transformation, one has

$$U^{-1}(\Lambda)p^\mu U(\Lambda) = \Lambda^\mu{}_\nu p^\nu$$

Under a parity transformation and a time reversal transformation, one obtains

$$U^{-1}(\Lambda_P)p^\mu U(\Lambda_P) = (\Lambda_P)^\mu{}_\nu p^\nu \implies p^\mu = (p^0, p^i) \rightarrow (p^0, -p^i)$$

$$U^{-1}(\Lambda_T)p^\mu U(\Lambda_T) = (\Lambda_T)^\mu{}_\nu p^\nu \implies p^\mu = (p^0, p^i) \rightarrow (-p^0, p^i)$$

Since  $p^0$  is the Hamiltonian density, then time reversal can be a symmetry if and only if  $H = 0$ . One then must have

$$U^{-1}(\Lambda_T)p^\mu U(\Lambda_T) = -(\Lambda_T)^\mu{}_\nu p^\nu$$

and as such one must consider an anti-unitary operator.

**Proof of anti-unitarity.** The existence of anti-unitarity is proven as follows. Consider the space-time translation operator

$$T(a) = e^{-ia_\mu p^\mu} = 1 - ia_\mu p^\mu + o(a_\mu), \quad a_\mu \ll 1$$

A translation of the field is

$$T^{-1}(a)\varphi(x)T(a) = \varphi(x - a)$$

Applying a proper Lorentz transformation gives

$$\begin{aligned} U^{-1}(\Lambda)\varphi(x - a)U(\Lambda) &= U^{-1}(\Lambda)T^{-1}(a)\varphi(x)T(a)U(\Lambda) \\ &= (U^{-1}T^{-1}U)(U^{-1}\varphi U)(U^{-1}TU) \\ \varphi(\Lambda^{-1}(x - a)) &= (U^{-1}T^{-1}U)\varphi(\Lambda^{-1}x)(U^{-1}TU) \\ \varphi(\Lambda^{-1}x - \Lambda^{-1}a) &= (U^{-1}TU)^{-1}\varphi(\Lambda^{-1}x)(U^{-1}TU) \\ \varphi(y - \Lambda^{-1}a) &= (U^{-1}TU)^{-1}\varphi(y)(U^{-1}TU) \end{aligned}$$

Therefore, the operator that translates by  $\Lambda^{-1}a$  is

$$T(\Lambda^{-1}a) = U^{-1}(\Lambda)T(a)U(\Lambda)$$

which can be written as an infinitesimal transformation

$$U^{-1}(\Lambda)(I - ia_\mu p^\mu)U(\Lambda) = I - i(\Lambda^{-1}a)_\mu p^\mu = I - i(\Lambda^{-1})_\mu^\nu a_\nu p^\mu$$

Therefore, for a generic Lorentz transformation, one finds

$$U^{-1}(\Lambda)(ia_\mu p^\mu)U(\Lambda) = i(\Lambda^{-1})_\mu^\nu a_\nu p^\mu$$

Consider a proper Lorentz transformation and the fact that

$$(\Lambda^{-1})_\mu^\nu = (\Lambda^\top)_\mu^\nu = \Lambda^\nu_\mu$$

then one has

$$U^{-1}(\Lambda)(ia_\mu p^\mu)U(\Lambda) = ia_\nu \Lambda^\nu_\mu p^\mu$$

Knowing that

$$\Lambda^\nu_\mu p^\mu = U^{-1}(\Lambda)p^\nu U(\Lambda)$$

one finds

$$U^{-1}(\Lambda)iU(\Lambda) = i$$

Consider a time reversal transformation

$$U^{-1}(\Lambda_T)(ia_\mu p^\mu)U(\Lambda_T) = ia_\nu (\Lambda_T)^\nu_\mu p^\mu$$

Assuming that

$$U^{-1}(\Lambda_T)iU(\Lambda_T) = i\eta, \quad \eta = \pm 1$$

one has

$$ia_\nu \eta U^{-1}(\Lambda_T)p^\nu U(\Lambda_T) = ia_\nu (\Lambda_T)^\nu_\mu p^\mu$$

One requires

$$U^{-1}(\Lambda_T)p^\nu U(\Lambda_T) = -(\Lambda_T)^\nu_\mu p^\mu$$

and there one obtains

$$\eta = -1$$

from which  $U(\Lambda_T)$  is anti-unitary.

## Lecture 9

### 8.1.2 Charge conjugation

Charge conjugating is a  $\mathbb{Z}_2$  symmetry:

$$Z^{-1}\varphi(x)Z = \pm\varphi(x), \quad Z^2 = 1$$

For real scalar fields there is no charge and the discussion ends here. For complex scalar fields, a Lagrangian may be

$$\begin{aligned} \mathcal{L} &= \partial_\mu \varphi^\dagger \partial^\mu \varphi - m^2 \varphi^\dagger \varphi - \frac{\lambda}{4!} (\varphi^\dagger \varphi)^2 \\ &= \frac{1}{2} \partial_\mu \varphi_1 \partial^\mu \varphi_1 - \frac{1}{2} m^2 \varphi_1^2 + \frac{1}{2} \partial_\mu \varphi_2 \partial^\mu \varphi_2 - \frac{1}{2} m^2 \varphi_2^2 - \frac{\lambda}{4!} (\varphi_1 + \varphi_2)^4 \end{aligned}$$

where one has

$$\varphi = \frac{1}{\sqrt{2}}(\varphi_1 + i\varphi_2), \quad \varphi^\dagger = \frac{1}{\sqrt{2}}(\varphi_1 - i\varphi_2)$$

The Lagrangian written with complex fields has a global  $U(1)$  symmetry

$$\varphi' = e^{iq}\varphi, \quad \varphi'^\dagger = e^{-iq}\varphi^\dagger$$

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The field has charge  $q$  while the conjugate field has charge  $-q$ . They are associated to particles and anti-particles. The Lagrangian written in terms of real fields has a global  $\text{SO}(2)$  symmetry.

Along the previous symmetry, the complex-field Lagrangian exhibits also a  $\mathbb{Z}_2$  symmetry,  $\varphi \leftrightarrow \varphi^\dagger$ . The generator of this symmetry is given by the operator  $C$ :

$$C^{-1}\varphi(x)C = \varphi^\dagger(x), \quad C\varphi^\dagger(x)C^{-1} = \varphi(x) \implies C^{-1}\mathcal{L}C = \mathcal{L}$$

For the real-field Lagrangian, one has

$$C^{-1}\varphi_1(x)C = \varphi_1(x), \quad C^{-1}\varphi_2(x)C = -\varphi_2(x)$$

In matrix form, one has

$$C = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix}, \quad \det C = -1$$

Since this operator has negative determinant, one has to enlarge the symmetry group to  $\text{O}(2)$ .

The operator  $C$  is responsible for charge conjugation since it exchanges particles of charge  $q$  with anti-particles of charge  $-q$ .

## 8.2 Representation on spinor fields

Consider spinor fields, that is spin-half fermionic fields,  $s = \frac{1}{2}$ . Recall that, in canonical quantization, the most general solution to the free Dirac equation is

$$\psi_\alpha(x) = \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [b_s(\mathbf{p})u_{s\alpha}(\mathbf{p})e^{-ipx} + d_s^\dagger(\mathbf{p})v_{s\alpha}(\mathbf{p})e^{ipx}], \quad p_0 \equiv \omega = \sqrt{|\mathbf{p}|^2 + m^2}$$

In the Weyl basis, the Dirac matrices are

$$\gamma^0 = \begin{bmatrix} 0 & I_2 \\ I_2 & 0 \end{bmatrix}, \quad \gamma^k = \begin{bmatrix} 0 & \sigma^k \\ -\sigma^k & 0 \end{bmatrix}, \quad \gamma^5 = \begin{bmatrix} -I_2 & 0 \\ 0 & I_2 \end{bmatrix}$$

where  $\sigma^k$  are the Pauli matrices. One introduces the charge conjugation matrix<sup>13</sup>

$$C = i\gamma^0\gamma^2 = \begin{bmatrix} -i\sigma^2 & 0 \\ 0 & i\sigma^2 \end{bmatrix} = \begin{bmatrix} 0 & -1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & -1 & 0 \end{bmatrix}$$

**Identities.** For the plane-wave spinor solutions (see Peskin, eqs. 3.59, 3.62)

$$u_s(\mathbf{p}) = \begin{bmatrix} \sqrt{\mathbf{p} \cdot \boldsymbol{\sigma}} \xi_s \\ \sqrt{\mathbf{p} \cdot \boldsymbol{\sigma}} \xi_s \end{bmatrix}, \quad v_s(\mathbf{p}) = \begin{bmatrix} \sqrt{\mathbf{p} \cdot \boldsymbol{\sigma}} \eta_s \\ -\sqrt{\mathbf{p} \cdot \boldsymbol{\sigma}} \eta_s \end{bmatrix}$$

it holds (see Srednicki, §38, eqs. 38.32, 38.40)

$$u_s(-\mathbf{p}) = \gamma^0 u_s(\mathbf{p}), \quad v_s(-\mathbf{p}) = -\gamma^0 v_s(\mathbf{p})$$

also

$$u_{-s}^*(-\mathbf{p}) = -sC\gamma^5 u_s(\mathbf{p}), \quad v_{-s}^*(-\mathbf{p}) = -sC\gamma^5 v_s(\mathbf{p})$$

**Proper Lorentz transformation.** Under a proper Lorentz transformation,  $\Lambda \in \mathcal{L}_+^\uparrow$ , a spinor transforms as

$$U^{-1}(\Lambda)\psi(x)U(\Lambda) = D(\Lambda)\psi(\Lambda^{-1}x)$$

with

$$D(\Lambda) = e^{\frac{i}{2}\omega_{\mu\nu}S^{\mu\nu}}, \quad S^{\mu\nu} = \frac{i}{4}[\gamma^\mu, \gamma^\nu]$$

where  $S^{\mu\nu}$  is the spin representation of the Lorentz generators.

<sup>13</sup>Typically the convention is to define  $C = i\gamma^2\gamma_0$ . In this convention the following identities change too.

### 8.2.1 Parity

One makes the ansatz that

$$U^{-1}(\Lambda_P)\psi(x)U(\Lambda_P) = D(P)\psi(\Lambda_P x)$$

where  $D(P)$  is an operator to be determined in the following. By applying the relation above twice, one obtains  $-\psi(x)$ . One does not require unity because spinors are not observables, only bilinear functions made from them are

$$U^{-2}(\Lambda_P)\psi(x)U^2(\Lambda_P) = D^2(P)\psi(\Lambda_P x) \equiv \pm\psi(x) \implies D^2(P) = \pm 1$$

Remembering that, under parity, the momentum  $\mathbf{p}$  changes sign while the angular momentum  $\mathbf{J}$  does not, then one expects that

$$U^{-1}(\Lambda_P)b_s^\dagger(\mathbf{p})U(\Lambda_P) = \eta b_s^\dagger(-\mathbf{p}), \quad U^{-1}(\Lambda_P)d_s^\dagger(\mathbf{p})U(\Lambda_P) = \eta d_s^\dagger(-\mathbf{p})$$

with  $\eta^2 = \pm 1$ . One introduces the same constant  $\eta$  for  $b$  and  $d$  to be able to relate the operators when working with Majorana fermions.

[r] Therefore, substituting the wave expansion — and recalling that the dispersion relation holds —, one has

$$\begin{aligned} U^{-1}(\Lambda_P)\psi(x)U(\Lambda_P) &= \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [\eta^* b_s(-\mathbf{p}) u_s(\mathbf{p}) e^{ipx} + \eta d_s^\dagger(-\mathbf{p}) v_s(\mathbf{p}) e^{ipx}]_{p^0=\omega} \\ &= \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [\eta^* b_s(\mathbf{p}) u_s(-\mathbf{p}) e^{-ip\Lambda_P x} + \eta d_s^\dagger(\mathbf{p}) v_s(-\mathbf{p}) e^{ip\Lambda_P x}] \\ &= \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [\eta^* b_s(\mathbf{p}) \gamma^0 u_s(\mathbf{p}) e^{-ip\Lambda_P x} - \eta d_s^\dagger(\mathbf{p}) \gamma^0 v_s(\mathbf{p}) e^{ip\Lambda_P x}] \\ &= i\gamma^0 \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [b_s(\mathbf{p}) u_s(\mathbf{p}) e^{-ip\Lambda_P x} + \eta d_s^\dagger(\mathbf{p}) v_s(\mathbf{p}) e^{ip\Lambda_P x}] \\ &= i\gamma^0 \psi(\Lambda_P x) \end{aligned}$$

At the second line, one has sent  $\mathbf{p} \rightarrow -\mathbf{p}$  in the integral. The spatial components of the momentum in the exponential gain an extra minus sign which can be put into the position  $\mathbf{x}$  to have

$$e^{-ipx} \rightarrow e^{-ip\Lambda_P x}$$

Notice that there is no overall minus sign since the integration limits are mapped into one another and the sign from the measure is absorbed to swap them. At the third line, one has applied the identities above. At the fourth line, one has required that

$$\eta^* = -\eta \implies \eta \in i\mathbb{R}$$

so that one obtains a spinor field again. One chooses  $\eta = -i$ .

From this, one sees that

$$\boxed{D(P) = i\gamma^0} \implies D^2(P) = -I$$

**Spinor transformation.** Consider a spinor<sup>14</sup>

$$\psi = \begin{bmatrix} \psi_L \\ \psi_R \end{bmatrix} = \begin{bmatrix} \chi_\alpha \\ \xi^{\dot{\alpha}} \end{bmatrix}$$

The spinor transforms as

$$D(P)\psi = i\gamma^0\psi = i \begin{bmatrix} 0 & I_2 \\ I_2 & 0 \end{bmatrix} \begin{bmatrix} \psi_L \\ \psi_R \end{bmatrix} = i \begin{bmatrix} \psi_R \\ \psi_L \end{bmatrix}$$

The chiral components are swapped.

<sup>14</sup>The second equality is in the van der Waerden notation.

### 8.2.2 Time reversal

One assumes that under a time reversal transformation, the spinor field changes as

$$U^{-1}(\Lambda_T)\psi(x)U(\Lambda_T) = D(T)\psi(\Lambda_T x)$$

As before, one must have  $D^2(T) = \pm 1$ .

Noting that, under time reversal, the momentum  $\mathbf{p}$  and the angular momentum  $\mathbf{J}$  both change sign, one has

$$U^{-1}(\Lambda_T)b_s^\dagger(\mathbf{p})U(\Lambda_T) = \zeta_s b_{-s}^\dagger(-\mathbf{p}), \quad U^{-1}(\Lambda_T)d_s^\dagger(\mathbf{p})U(\Lambda_T) = \zeta_s d_{-s}^\dagger(-\mathbf{p})$$

with  $\zeta_s^2 = \pm 1$ . When one applies the time reversal transformation to the spinor field, one produces the term

$$U^{-1}(\Lambda_T)\psi(x)(\mathbf{p})U(\Lambda_T) \propto U^{-1}(\Lambda_T)b_s(\mathbf{p})u_s(\mathbf{p})e^{-ipx}U(\Lambda_T) = U^{-1}(\Lambda_T)b_s(\mathbf{p})U(\Lambda_T)u_s^*(\mathbf{p})e^{ipx}$$

where one inserts  $UU^{-1}$  between  $bu$  and recalls that  $U(\Lambda_T)$  is anti-unitary. Therefore

$$\begin{aligned} U^{-1}(\Lambda_T)\psi(x)U(\Lambda_T) &= \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [\zeta_s^* b_{-s}(-\mathbf{p})u_s^*(\mathbf{p})e^{ipx} + \zeta_s d_{-s}^\dagger(-\mathbf{p})v_s^*(\mathbf{p})e^{-ipx}]_{p^0=\omega} \\ &= \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [\zeta_{-s}^* b_s(\mathbf{p})u_{-s}^*(-\mathbf{p})e^{-ip\Lambda_T x} + \zeta_{-s} d_s^\dagger(\mathbf{p})v_{-s}^*(-\mathbf{p})e^{ip\Lambda_T x}] \\ &= \sum_{s=\pm} \int \frac{d^3p}{(2\pi)^3 2\omega} [\zeta_{-s}^* b_s(\mathbf{p})(-sC\gamma^5 u_s(\mathbf{p}))e^{-ip\Lambda_T x} \\ &\quad + \zeta_{-s} d_s^\dagger(\mathbf{p})(-sC\gamma^5 v_s(\mathbf{p}))e^{ip\Lambda_T x}] \\ &= C\gamma^5 \psi(\Lambda_T x) \end{aligned}$$

At the second line, one has interchanged  $\mathbf{p} \rightarrow -\mathbf{p}$  and  $s \rightarrow -s$ . Similar to before, the spatial components of the momentum in the exponential gain a minus sign which is put in front along with the minus sign of the time-like component coming from  $\Lambda_T x$ . At the third line, one has applied the identities above. At the fourth line, to obtain a spinor field, one has to choose

$$\zeta_s = s \implies \zeta_{-s} = -s \implies (-s)^2 = 1$$

Therefore

$$\boxed{D(T) = C\gamma^5} \implies D^2(T) = -I$$

**Spinor transformation.** Noting that

$$C = \begin{bmatrix} 0 & -1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & -1 & 0 \end{bmatrix} = \begin{bmatrix} -\varepsilon^{\alpha\beta} & 0 \\ 0 & -\varepsilon_{\dot{\alpha}\dot{\beta}} \end{bmatrix}, \quad \gamma^5 = \begin{bmatrix} -I_2 & 0 \\ 0 & I_2 \end{bmatrix} = \begin{bmatrix} -\delta_\beta^\gamma & 0 \\ 0 & \delta_\dot{\beta}^{\dot{\gamma}} \end{bmatrix}$$

one finds

$$C\gamma^5 = \begin{bmatrix} 0 & 1 & 0 & 0 \\ -1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & -1 & 0 \end{bmatrix} = \begin{bmatrix} \varepsilon^{\alpha\gamma} & 0 \\ 0 & -\varepsilon_{\dot{\alpha}\dot{\gamma}} \end{bmatrix}$$

from which

$$C\gamma^5 \begin{bmatrix} \psi_L \\ \psi_R \end{bmatrix} = \begin{bmatrix} \varepsilon & 0 \\ 0 & -\varepsilon \end{bmatrix} \begin{bmatrix} \psi_L \\ \psi_R \end{bmatrix} = \begin{bmatrix} \varepsilon \psi_L \\ -\varepsilon \psi_R \end{bmatrix}$$

Time reversal does not exchange chiral components, but changes each component into its dual.  
[r]

### 8.2.3 Charge conjugation

Let the charge conjugation matrix be  $C$  and the charge conjugation generator on the field  $\psi$  be  $\mathcal{C}$ . Then a charge conjugation transformation is

$$\mathcal{C}^{-1}\psi(x)\mathcal{C} = C\bar{\psi}^\top(x) = C(\psi^\dagger\gamma^0)^\dagger = C\gamma^0\psi^*$$

Similarly

$$\mathcal{C}^{-1}\bar{\psi}(x)\mathcal{C} = \psi^\top(x)C$$

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The charge conjugation matrix satisfies

$$C^\top = C^\dagger = C^{-1} = -C, \quad \{\gamma^0, C\} = 0, \quad C^{-1}\gamma^\mu C = -(\gamma^\mu)^\top, \quad C^{-1}\gamma^5 C = \gamma^5$$

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### 8.2.4 Majorana spinors

**Definition.** A Majorana spinor is a bi-spinor satisfying

$$\psi \equiv \psi^c = C\bar{\psi}^\top = C\gamma^0\psi^*$$

Recalling that

$$\gamma^0 = \begin{bmatrix} 0 & I_2 \\ I_2 & 0 \end{bmatrix} \implies C\gamma^0 = \begin{bmatrix} 0 & -i\sigma^2 \\ i\sigma^2 & 0 \end{bmatrix}$$

The Majorana condition is equivalent to

$$C\gamma^0\psi^* = \psi \iff \begin{bmatrix} 0 & -i\sigma^2 \\ i\sigma^2 & 0 \end{bmatrix} \begin{bmatrix} \psi_L^* \\ \psi_R^* \end{bmatrix} = \begin{bmatrix} \psi_L \\ \psi_R \end{bmatrix} \implies \begin{cases} -i\sigma^2\psi_L^* = \psi_L \\ i\sigma^2\psi_R^* = \psi_R \end{cases}$$

The degrees of freedom are half the ones for a Dirac spinor.

**Summary.** A brief summary of parity, time reversal and charge conjugation follows

$$\begin{aligned} U^{-1}(\Lambda_P)\psi(x)U(\Lambda_P) &= i\gamma^0\psi(\Lambda_P x) \\ U^{-1}(\Lambda_T)\psi(x)U(\Lambda_T) &= C\gamma^5\psi(\Lambda_T x) \\ \mathcal{C}^{-1}\psi(x)\mathcal{C} &= C\gamma^0\psi^*(x) \end{aligned}$$

One may consider the transformations for the Dirac adjoint field  $\bar{\psi}$ . Parity gives

$$\begin{aligned} U^{-1}(\Lambda_P)\bar{\psi}(x)U(\Lambda_P) &= U^{-1}(\Lambda_P)\psi^\dagger(x)\gamma^0U(\Lambda_P) = U^{-1}(\Lambda_P)\psi^\dagger(x)\gamma^0U(\Lambda_P) \\ &= -i\psi^\dagger(\Lambda_P x)\gamma^0\gamma^0 = -i\bar{\psi}(\Lambda_P x)\gamma^0 \end{aligned}$$

For time reversal, one finds

$$U^{-1}(\Lambda_T)\bar{\psi}(x)U(\Lambda_T) = \bar{\psi}(\Lambda_T x)\gamma^5 C^{-1}$$

Therefore, the summary for the Dirac adjoint is

$$\begin{aligned} U^{-1}(\Lambda_P)\bar{\psi}(x)U(\Lambda_P) &= -i\bar{\psi}(\Lambda_P x)\gamma^0 \\ U^{-1}(\Lambda_T)\bar{\psi}(x)U(\Lambda_T) &= \bar{\psi}(\Lambda_T x)\gamma^5 C^{-1} \\ \mathcal{C}^{-1}\bar{\psi}(x)\mathcal{C} &= \psi^\top(x)C \end{aligned}$$

## 8.3 Spinor bilinears

A bilinear in the spinor field is a real expression of the type  $\bar{\psi}A\psi$  where  $A$  is a product of Dirac matrices such that

$$(\bar{\psi}A\psi)^\dagger = \bar{\psi}A\psi \implies \bar{A} = \gamma^0 A^\dagger \gamma^0 = A$$

where the bar denotes the Dirac adjoint of a matrix.



**Parity.** A parity transformation of a bilinear is

$$\begin{aligned} U^{-1}(\Lambda_P)(\bar{\psi}A\psi)U(\Lambda_P) &= [U^{-1}(\Lambda_P)\bar{\psi}U(\Lambda_P)][U^{-1}(\Lambda_P)AU(\Lambda_P)][U^{-1}(\Lambda_P)\psi U(\Lambda_P)] \\ &= [-i\bar{\psi}\gamma^0 A i\gamma^0 \psi](\Lambda_P x) = \bar{\psi}\gamma^0 A \gamma^0 \psi = \bar{\psi}\tilde{A}\psi \end{aligned}$$

To go further, one needs to specify the bilinear and in particular  $A$ .

Let  $A = I$ , one has a scalar

$$U^{-1}(\Lambda_P)\bar{\psi}\psi U(\Lambda_P) = \bar{\psi}\psi$$

Let  $A = i\gamma^5$ , one has a pseudo-scalar

$$U^{-1}(\Lambda_P)i\bar{\psi}\gamma^5\psi U(\Lambda_P) = -i\bar{\psi}\gamma^5\psi$$

Let  $A = \gamma^\mu$ , one has a vector (also called polar vector)

$$U^{-1}(\Lambda_P)\bar{\psi}\gamma^\mu\psi U(\Lambda_P) = (\Lambda_P)^\mu_\nu(\bar{\psi}\gamma^\nu\psi)$$

where the parity transformation  $\Lambda_P$  encodes the following relations

$$\bar{\psi}\gamma^0\gamma^\mu\gamma^0\psi = \begin{cases} \bar{\psi}\gamma^0\psi, & \mu = 0 \\ -\bar{\psi}\gamma^i\psi, & \mu = i \end{cases}$$

Let  $A = \gamma^\mu\gamma^5$ , one has a pseudo-vector (also called axial vector)

$$U^{-1}(\Lambda_P)\bar{\psi}\gamma^\mu\gamma^5\psi U(\Lambda_P) = -(\Lambda_P)^\mu_\nu(\bar{\psi}\gamma^\nu\gamma^5\psi)$$

**Time reversal.** The transformation of a bilinear is

$$\begin{aligned} U^{-1}(\Lambda_T)(\bar{\psi}A\psi)U(\Lambda_T) &= [U^{-1}(\Lambda_T)\bar{\psi}U(\Lambda_T)][U^{-1}(\Lambda_T)AU(\Lambda_T)][U^{-1}(\Lambda_T)\psi U(\Lambda_T)] \\ &= \bar{\psi}\gamma^5 C^{-1} A^* C \gamma^5 \psi \end{aligned}$$

The second bracket evaluates to  $A^*$  since the operator  $U(\Lambda_T)$  is anti-unitary. For the following it is useful to remember that

$$\gamma_{0,5} = \gamma_{0,5}^\dagger, \quad \gamma_i = -\gamma_i^\dagger$$

which can be shortened to

$$\gamma_5 = \gamma_5^\dagger, \quad (\gamma^\mu)^\dagger = \gamma^0\gamma^\mu\gamma^0$$

Let  $A = I$ , the bilinear is T-even

$$U^{-1}(\Lambda_T)(\bar{\psi}\psi)U(\Lambda_T) = \bar{\psi}\gamma^5 C^{-1} C \gamma^5 \psi = \bar{\psi}\psi$$

Let  $A = i\gamma^5$ , the bilinear is T-odd

$$U^{-1}(\Lambda_T)(\bar{\psi}i\gamma^5\psi)U(\Lambda_T) = \bar{\psi}\gamma^5 C^{-1}(-i\gamma^5)C\gamma^5\psi = -i\bar{\psi}(\gamma^5)^3\psi = -i\bar{\psi}\gamma^5\psi$$

where one applies the properties of the charge conjugation matrix. Let  $A = \gamma^\mu$ , the bilinear is T-odd

$$U^{-1}(\Lambda_T)(\bar{\psi}\gamma^\mu\psi)U(\Lambda_T) = \bar{\psi}\gamma^5 C^{-1}(\gamma^\mu)^* C \gamma^5 \psi = -\bar{\psi}\gamma^5(\gamma^\mu)^\dagger\gamma^5\psi = \bar{\psi}(\gamma^\mu)^\dagger\psi = -(\Lambda_T)^\mu_\nu\bar{\psi}\gamma^\nu\psi$$

The sign difference is encoded into the time reversal transformation  $\Lambda_T$ .

**Exercise.** Check that  $\bar{\psi}\gamma^\mu\gamma^5\psi$  is also T-odd.

**Charge conjugation.** The transformation of a bilinear is

$$\begin{aligned}\mathcal{C}^{-1}(\bar{\psi}A\psi)\mathcal{C} &= [\mathcal{C}^{-1}\bar{\psi}\mathcal{C}][\mathcal{C}^{-1}A\mathcal{C}][\mathcal{C}^{-1}\psi\mathcal{C}] = \psi^\top CAC\bar{\psi}^\top \\ &= (\psi^\top CAC\bar{\psi}^\top)^\top = -\bar{\psi}(CAC)^\top\psi = \bar{\psi}C^{-1}A^\top C\psi\end{aligned}$$

at the second line one has noticed that the transpose of a bilinear gains a minus sign due to interchanging the spinor fields<sup>15</sup>.

Let  $A = I$ , the bilinear is C-even

$$\mathcal{C}^{-1}\bar{\psi}\psi\mathcal{C} = \bar{\psi}\psi$$

Let  $A = i\gamma^5$ , the bilinear is C-even

$$\mathcal{C}^{-1}(\bar{\psi}i\gamma^5\psi)\mathcal{C} = \bar{\psi}i\gamma^5\psi$$

Let  $A = \gamma^\mu$ , the bilinear is C-odd

$$\mathcal{C}^{-1}(\bar{\psi}\gamma^\mu\psi)\mathcal{C} = -\bar{\psi}\gamma^\mu\psi$$

Let  $A = \gamma^\mu\gamma^5$ , the bilinear is C-even

$$\mathcal{C}^{-1}(\bar{\psi}\gamma^\mu\gamma^5\psi)\mathcal{C} = \bar{\psi}\gamma^\mu\gamma^5\psi$$

**Remark.** If the spinor field is Majorana, then

$$\psi^c = \psi \iff \mathcal{C}^{-1}\psi\mathcal{C} = \psi, \quad \mathcal{C}^{-1}\bar{\psi}\mathcal{C} = \bar{\psi}$$

Therefore, one has [r]

$$\mathcal{C}^{-1}\bar{\psi}A\psi\mathcal{C} = \bar{\psi}A\psi$$

This means that the odd bilinears have to be zero. In particular

$$\bar{\psi}\gamma^\mu\psi = 0$$

## 8.4 Representation on vector fields

[r] Source?. Each component of a vector field  $V^\mu(x)$  behaves as a scalar field under charge conjugation.

**Parity.** Under parity, there are vectors and pseudo-vectors

$$U^{-1}(\Lambda_P)\mathbf{V}(x)U(\Lambda_P) = -\mathbf{V}(\Lambda_P x), \quad U^{-1}(\Lambda_P)\mathbf{A}(x)U(\Lambda_P) = \mathbf{A}(\Lambda_P x)$$

[r]

**Time reversal.** For time reversal one has T-even and T-odd vectors, similar to vector bilinears.

**Theorem** (Furry's in QED). In QED, the interaction term is

$$\mathcal{L}_{\text{int}} = J^\mu A_\mu, \quad J^\mu \propto \bar{\psi}\gamma^\mu\psi$$

A Lagrangian invariant under charge conjugation implies that the four-potential  $A^\mu$  has to be odd

$$\mathcal{C}^{-1}A^\mu\mathcal{C} = -A^\mu$$

since the current  $J^\mu$  is odd. This implies that the correlation functions of an odd number of gauge fields  $A^\mu$  are zero.

**Example.** For example, the triangle diagram has two configurations depending on the orientation of the internal fermionic propagators. These two configurations sum up to zero.



<sup>15</sup>See <https://physics.stackexchange.com/q/458451>.

## 8.5 CPT theorem

See Weinberg, vol. 1, §5.8, Srednicki, §40.

**Theorem.** The action of a local and hermitian<sup>16</sup> Lagrangian is symmetric under a simultaneous transformation of charge conjugation, parity and time reversal

$$(CPT)^{-1}\mathcal{L}(x)(CPT) = \mathcal{L}(-x) \implies S = \int d^4x \mathcal{L}(-x) = \int d^4x \mathcal{L}(x)$$

*Proof.* For scalar fields, one has

$$P^{-1}\varphi P = \pm\varphi = \eta\varphi, \quad T^{-1}\varphi T = \pm\varphi = \zeta^*\varphi, \quad C^{-1}\varphi C = \pm\varphi^\dagger = \xi\varphi^\dagger$$

Therefore

$$(CPT)^{-1}\varphi(CPT) = \eta\zeta^*\xi\varphi^\dagger(-x) = \varphi^\dagger(-x)$$

where one chooses the inversion phases such that

$$\eta\zeta^*\xi = 1$$

For spinor fields, one has

$$\begin{aligned} (CPT)^{-1}\bar{\psi}\psi(CPT) &= \bar{\psi}\psi \\ (CPT)^{-1}\bar{\psi}i\gamma^5\psi(CPT) &= \bar{\psi}i\gamma^5\psi \\ (CPT)^{-1}\bar{\psi}\gamma^\mu\psi(CPT) &= -\bar{\psi}\gamma^\mu\psi \\ (CPT)^{-1}\bar{\psi}\gamma^\mu\gamma^5\psi(CPT) &= -\bar{\psi}\gamma^\mu\gamma^5\psi \end{aligned}$$

where the first two are CPT-even and the other two are CPT-odd. This is true for any bilinear. In general, bilinears with an even number of vector indices are even under CPT; while for an odd number, they are for CPT-odd. This is true also for four-derivatives. In fact

$$(CPT)^{-1}\partial_\mu\varphi(x)(CPT) = -\partial_\mu\varphi^\dagger(-x)$$

On to the actual proof. Since a Lagrangian is a Lorentz scalar, it is CPT-invariant [r]. Since the Lagrangian is hermitian, it contains the same number of fields and their conjugates, therefore the Lagrangian is CPT-even. This means that the Lagrangian is CPT-invariant even if  $\varphi \leftrightarrow \varphi^\dagger$ .  $\square$

## Lecture 11

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## Part II

# Gauge fields

**Review of classical fields – abelian gauge theory.** See Cheng, §8.1. Consider QED as an abelian gauge theory. The free Lagrangian is

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

The Lagrangian is invariant under a global U(1) transformation

$$\psi' = e^{i\alpha}\psi, \quad \bar{\psi}' = \bar{\psi}e^{-i\alpha}$$

The gauge principle involves promoting the global parameter  $\alpha$  to a function of space-time  $\alpha(x)$ . By requiring a local U(1) invariance, one introduces a gauge covariant derivative<sup>17</sup> and a gauge field

$$D_\mu = \partial_\mu + iqA_\mu, \quad A'_\mu = A_\mu - \frac{1}{q}\partial_\mu\alpha$$

<sup>16</sup>Hermitian so that it guarantees unitarity.

<sup>17</sup>Note that Cheng uses  $q = e$  where, for an electron,  $e < 0$ .

where  $q$  is the charge (with sign) of the bispinor field. The gauge-invariant Lagrangian has an interaction term

$$\mathcal{L} = \bar{\psi}(i\not{D} - m)\psi = \bar{\psi}(i\not{\partial} - m)\psi - qA_\mu\bar{\psi}\gamma^\mu\psi$$

If there is no kinetic term for the gauge field, then the field is a background field. To make it dynamical, one has to insert such kinetic term

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu}, \quad F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$$

where the field strength tensor is defined by

$$i q F_{\mu\nu} = [D_\mu, D_\nu]$$

and it is invariant under a gauge transformation.

One observes that gauge invariance forbids any mass term for the gauge field. The gauge Lagrangian term is a free Lagrangian — meaning it does not contain any interaction vertex — since the gauge field  $A_\mu$  is neutral (because it is real) and cannot couple to itself.

**Review of classical fields – non-abelian gauge theory.** An example of non-abelian gauge theory is QCD. Consider the gauge group  $G = \text{SU}(n)$  (so the theory is a Yang–Mills theory). The gauge group has  $n^2 - 1$  generators that satisfy the (fundamental representation of the)  $\mathfrak{su}(n)$  Lie algebra

$$[T^a, T^b] = i c^{abc} T^c$$

where  $c^{abc}$  are the structure constants,  $abc$  are indices labelling the dimensions of the algebra and the generators are traceless Hermitian complex  $n \times n$  matrices<sup>18</sup>. Consider a set of Dirac spinors in some representation of the algebra  $\mathfrak{su}(n)$ . The Lagrangian is

$$\mathcal{L} = \bar{\psi}_j(i\not{\partial} - m)\psi_j$$

It has a global  $\text{SU}(n)$  symmetry

$$\psi'_j = U_{jk}(\theta)\psi_k, \quad \bar{\psi}'_j = \bar{\psi}_k U_{kj}^\dagger(\theta)$$

where the transformation matrix is

$$U_{jk}(\theta) = [e^{-i\theta^a T^a}]_{jk}, \quad UU^\dagger = U^\dagger U = I$$

By applying the gauge principle, one obtains  $\theta = \theta(x)$ . The requirement of local  $\text{SU}(n)$  symmetry implies

$$D_\mu = \partial_\mu - i g T^a A_\mu^a, \quad (A_\mu^a)' = A_\mu^a + c^{abc} \theta^b A_\mu^c - \frac{1}{g} \partial_\mu \theta^a$$

Let  $A_\mu = T^a A_\mu^a$ , then the infinitesimal transformation is

$$A'_\mu = A_\mu - i[\theta, A_\mu] - \frac{1}{g} \partial_\mu \theta$$

To make the gauge fields dynamical, one introduces the field strengths

$$[D_\mu, D_\nu] = -i g T^a F_{\mu\nu}^a, \quad F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + g c^{abc} A_\mu^b A_\nu^c$$

The field strength is not gauge-invariant, but transforms infinitesimally as

$$F_{\mu\nu}^{'a} = F_{\mu\nu}^a + c^{abc} \theta^b F_{\mu\nu}^c$$

For a finite transformation, it is a covariant tensor

$$F_{\mu\nu}' = U^\dagger(\theta) F_{\mu\nu} U(\theta)$$

<sup>18</sup>In the physicists' convention. In the mathematicians' convention they are traceless anti-Hermitian matrices. The two conventions differ by a factor of  $i$ .

The gauge kinetic term is then

$$\begin{aligned}\mathcal{L}_{\text{gauge}} &= -\frac{1}{4} \text{Tr}(F_{\mu\nu} F^{\mu\nu}) = -\frac{1}{4} F_{\mu\nu}^a F^{b\mu\nu} \text{Tr}(T^a T^b) = -\frac{1}{4} F_{\mu\nu}^a F^{b\mu\nu} \delta^{ab} \\ &= -\frac{1}{2} \partial_\mu A_\nu^a (\partial^\mu A^{a\nu} - \partial^\nu A^{a\mu}) - g c^{abc} (\partial_\mu A_\nu^a) A^{b\mu} A^{c\nu} - \frac{g^2}{4} c^{abc} c^{adf} A_\mu^b A_\nu^c A^{d\mu} A^{f\nu}\end{aligned}$$

The convention of the trace may be different. When inserting the explicit expression for the field strength one finds the following general structure

$$\text{kinetic} + g \partial A A^2 + g^2 A^4$$

There is derivative three-point vertex and a four-point vertex. The gauge Lagrangian term is not free Lagrangian since it contains two vertices.

One observes that the covariant derivative introduces an interaction term

$$g A_\mu^a \bar{\psi}_j T_{jk}^a \gamma^\mu \psi_k = g A_\mu^a J^{a\mu}$$

where  $jk$  label the bispinor fields of the theory and  $a$  is an  $\mathfrak{su}(n)$  index. [r]

While in abelian gauge theories, one may have different matter sectors coupled to the same gauge field with different coupling constants, in this case the coupling constant  $g$  is fixed when choosing the gauge group and it has to be the same for all matter fields since it appears in the field strength.

For each simple gauge group (e.g.  $\text{SU}(n)$ ), there is a corresponding single coupling constant.

**Canonical quantization.** Consider QED. Out of the four components of a gauge field, only two are physical degrees of freedom, typically  $A^1$  and  $A^2$ . This is due to gauge invariance which can be used to eliminate non-physical degrees of freedom.

One may use gauge invariance to fix the Lorenz gauge. If there are no sources in the free theory, one may perform a further gauge transformation while being in the Lorenz gauge to set  $A_3 = 0$ . In fact, consider a generic gauge

$$\partial_\mu A^\mu \neq 0$$

One may transform the gauge field to go to the Lorenz gauge

$$A'_\mu(x) = A_\mu(x) - \partial_\mu \Lambda(x), \quad 0 = \partial_\mu A'^\mu = \partial_\mu A^\mu - \square \Lambda \implies \square \Lambda = \partial_\mu A^\mu$$

From gauge field  $A'$ , one can always choose another gauge field within the Lorenz gauge

$$A''_\mu(x) = A'_\mu(x) - \partial_\mu \lambda(x), \quad 0 = \partial^\mu A''_\mu = \partial_\mu A'^\mu - \square \lambda \implies \square \lambda = 0$$

Since  $\lambda$  is an arbitrary superposition of plane waves, one may choose

$$A''_3 = A'_3 - \partial_3 \lambda = 0 \implies \partial_3 \lambda = A'_3$$

There are two ways to perform canonical quantization.

1. One eliminates the non-physical degrees of freedom and quantizes only the physical ones. This method is not Lorentz-invariant because one fixes a component,  $A_3 = 0$ .
2. The other possibility follows the approach by Gupta-Bleuler<sup>19</sup> in quantizing all four degrees of freedom and removing the non-physical ones at the end. This approach produces negative-norm states that can be eliminated from the spectrum of physical states by defining the physical states as

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

where the superscript  $(+)$  denotes the positive frequencies of the gauge field, those corresponding to the annihilation operators,  $a_r(\mathbf{k})$ .

<sup>19</sup>See Theoretical Physics II for a covariant quantization of the electromagnetic field.

## 9 Functional quantization

See Cheng, §9. The method of functional quantization is an alternative to canonical quantization. The LSZ reduction formula can also be formulated for vector fields. The scattering amplitudes are directly related to the computation of correlation functions.

For a pure gauge theory (i.e. no matter fields), one would like to construct a generating functional that provides all the possible correlation functions

$$W[J] = \int [\mathcal{D}A_\mu] \exp \left[ i \int d^4x (\mathcal{L}_g + J_\mu^a A^{a\mu}) \right]$$

The Green's functions are

$$\langle 0 | \mathcal{T} \{ A_{\mu_1}^{a_1}(x) \cdots A_{\mu_n}^{a_n}(x_n) \} | 0 \rangle = (-i)^n \frac{\delta^n W[J]}{\delta J_{a_1 \mu_1}(x_1) \cdots \delta J_{a_n \mu_n}(x_n)} \Big|_{J=0}$$

### 9.1 Free theory

The free Lagrangian (which does not contain the self-interactions) is

$$\mathcal{L}_0 = -\frac{1}{2} \partial_\mu A_\nu^a (\partial^\mu A^{a\nu} - \partial^\nu A^{a\mu})$$

The associated generating functional is

$$W_0[J] = \int [\mathcal{D}A_\mu] \exp \left[ i \int d^4x (\mathcal{L}_0 + J_\mu^a A_\mu^a) \right]$$

[r] The free action is

$$S_0 = -\frac{1}{2} \int d^4x \partial_\mu A_\nu^a (\partial^\mu A^{a\nu} - \partial^\nu A^{a\mu}) = \frac{1}{2} \int d^4x A_\mu^a (\eta^{\mu\nu} \square - \partial^\mu \partial^\nu) A_\nu^a = \frac{1}{2} \int d^4x A K A$$

The differential operator is the kinetic operator

$$K^{\mu\nu} \equiv \eta^{\mu\nu} \square - \partial^\mu \partial^\nu$$

This operator is not invertible. If it could be inverted, one would complete the square in the exponent of the generating functional

$$\int d^4x (A K A + 2 J A) = \int d^4x [(A + J K^{-1}) K (A + K^{-1} J) - J K^{-1} J] = \int d^4x (\tilde{A} K \tilde{A} - J K^{-1} J)$$

therefore the generating functional would be

$$\begin{aligned} W_0[J] &= \exp \left[ -\frac{i}{2} \int d^4x J K^{-1} J \right] \int [\mathcal{D}\tilde{A}_\mu] \exp \left[ \frac{i}{2} \int d^4x \tilde{A} K \tilde{A} \right] \\ &\propto \frac{1}{\sqrt{\det K}} \exp \left[ -\frac{i}{2} \int d^4x J K^{-1} J \right] \end{aligned}$$

However, the inverse of the kinetic operator does not exist. Consider

$$K_{\mu\nu} K^{\nu\rho} = (\eta_{\mu\nu} \square - \partial_\mu \partial_\nu) (\delta_\rho^\nu \square - \partial^\nu \partial_\rho) = \eta_{\mu\rho} \square^2 - \square \partial_\mu \partial_\rho = \square (\eta_{\mu\rho} \square - \partial_\mu \partial_\rho) = \square K_{\mu\rho}$$

From this one sees that the operator  $K$  behaves like a projection operator and as such is not invertible (because it is not injective).

One looks for the Green's function  $G$  of the kinetic operator

$$K_{\mu\nu} G^{\nu\rho} = \delta_\mu^\rho \iff (\eta_{\mu\nu} \square - \partial_\mu \partial_\nu) G^{\nu\rho}(x-y) = \delta_\mu^\rho \delta^{(4)}(x-y)$$

By Fourier transformation, one has

$$G^{\nu\rho}(x-y) = \int \frac{d^4k}{(2\pi)^4} e^{ik(x-y)} \tilde{G}^{\nu\rho}(k)$$

from which the equation above becomes

$$(-k^2 \eta_{\mu\nu} + k_\mu k_\nu) \tilde{G}^{\nu\rho}(k) = \delta_\mu^\rho$$

The most general type-(2,0) tensor field, function of the momentum  $k$  has the form

$$\tilde{G}^{\nu\rho}(k) = A(k^2) \eta^{\mu\nu} + B(k^2) k^\nu k^\rho$$

Inserting this expression above, one finds no possible functions  $A$  and  $B$ .

**Remark.** Since the kinetic term is not invertible, then at least one eigenvalue is zero. The operator has a zero mode. Its determinant is zero and the generating functional is ill-defined. This is because the functional integral goes over all configurations of the gauge field. Recalling that the field is gauge invariant, the exponential in the generating functional is the same for each field connected by a gauge transformation. One is summing the same phase.

**Remark.** If one were to fix the gauge, one could modify the kinetic term to obtain an invertible operator.

## Lecture 12

### 9.1.1 Faddeev–Popov prescription

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All the field configurations linked by a gauge transformation constitute a gauge orbit. An orbit has infinite volume which makes the generating functional ill-defined. The problem can be solved by employing the Faddeev–Popov prescription. One rewrites the generating functional as

$$W[J] \sim e^{iS} \times (\text{volume gauge orbit})$$

and removes the infinite volume of the gauge orbits. One declares this as the well-defined prescription.

**Example.** An illustrative example with ordinary integrals is the following. Consider a Lebesgue bidimensional integral

$$W = \int dx dy e^{iS(x,y)}$$

Going to polar coordinates  $(r, \theta)$

$$x = r \cos \theta, \quad y = r \sin \theta, \quad dx dy = r dr d\theta$$

one has

$$W = \int r dr d\theta e^{iS(r,\theta)}$$

To each fixed radius corresponds a circle. The prescription is read as integrating along a circle and then integrating over all circles. Let  $S$  be a function invariant under rotation  $\theta \rightarrow \theta + \varphi$ . This invariance is the analogue of gauge invariance in QFT. This implies that the function depends only on the radius

$$S(r, \theta) = S(r)$$

The integral is

$$W = \int r dr d\theta e^{iS(r)} = 2\pi \int r dr e^{iS(r)}$$

The factor  $2\pi$  is the volume of the circle, which corresponds to the volume of the orbit. In this case, the volume of the orbit is finite, but in QFT it is infinite.

One may compute the integral in an alternative way. One inserts unity into the initial integral

$$1 = \int_0^{2\pi} d\varphi \delta(\theta - \varphi)$$

therefore

$$W = \int_0^{2\pi} d\varphi \int r dr d\theta \delta(\theta - \varphi) e^{iS(r)} \equiv \int_0^{2\pi} d\varphi W_\varphi$$

This corresponds to fixing an angle  $\theta = \varphi$ , integrating over the radius and then sum over all angles. The integral  $W_\varphi$  is invariant under a rotation by an angle  $\varphi$  (because the integrand is invariant). Therefore  $W_\varphi$  does not depend on such angle  $\varphi$ . Thus, the whole integral may be written as

$$W = \int_0^{2\pi} d\varphi W_\varphi = W_\varphi \int_0^{2\pi} d\varphi = 2\pi W_\varphi$$

The  $\varphi$ -integral is the integral on the orbit and the original integral  $W$  can be normalized by the volume

$$\frac{W}{\text{volume orbit}} = W_\varphi$$

and one considers this renormalized function.

The insertion of the delta function means integrating first over the radius and then over the angle. One can generalize the approach by integrating first over an arbitrary line (i.e. not straight) that intersects each orbit only once and then integrate over the angle. This line is given by a gauge function  $g(r, \theta)$  and is such that, for any  $r$ ,

$$g(r, \theta + \varphi) = 0$$

has a unique solution. One generalizes unity as

$$1 = \Delta_g(\mathbf{r}) \Delta_g^{-1}(\mathbf{r}), \quad \mathbf{r}_\varphi = (r, \theta + \varphi)$$

where one has

$$\Delta_g^{-1}(\mathbf{r}) = \int_0^{2\pi} d\varphi \delta(g(\mathbf{r}_\varphi)) = \int dg \partial_g \varphi \delta(g) = \partial_g \varphi|_{g=0} \implies \Delta_g(\mathbf{r}) = \partial_\varphi g|_{g=0}$$

A property of this object is gauge invariance under rotations

$$\Delta_g^{-1}(\mathbf{r}_{\varphi'}) = \int_0^{2\pi} d\varphi \delta(g(\mathbf{r}_{\varphi+\varphi'})) = \int_0^{2\pi} d(\varphi + \varphi') \delta(g(\mathbf{r}_{\varphi'+\varphi})) = \int_0^{2\pi} d\varphi'' \delta(g(\mathbf{r}_{\varphi''})) = \Delta_g^{-1}(\mathbf{r})$$

where one sets  $\varphi'' = \varphi + \varphi'$ .

Inserting this tool inside the integral  $W$ , one finds

$$\begin{aligned} W &= \int r dr d\theta e^{iS(r)} = \int r dr d\theta e^{iS(r)} \Delta_g(\mathbf{r}) \Delta_g^{-1}(\mathbf{r}) = \\ &= \int_0^{2\pi} d\varphi \int r dr d\theta \Delta_g(\mathbf{r}) \delta(g(\mathbf{r}_\varphi)) e^{iS(r)} \\ &= \int_0^{2\pi} d\varphi \int r dr d(\theta + \varphi) \Delta_g(\mathbf{r}_\varphi) \delta(g(\mathbf{r}_\varphi)) e^{iS(\mathbf{r}_\varphi)} \\ &= \int_0^{2\pi} d\varphi \int r dr d\theta' \Delta_g(\mathbf{r}') \delta(g(\mathbf{r}')) e^{iS(\mathbf{r}')} \\ &\equiv \int_0^{2\pi} d\varphi W_\varphi \end{aligned}$$

At the third line, one has applied

$$\Delta_g(\mathbf{r}) = \Delta(\mathbf{r}_\varphi), \quad S(r) = S(\mathbf{r}) = S(\mathbf{r}_\varphi), \quad r dr d\theta = r dr d(\theta + \varphi)$$

At the fourth line one renames

$$\mathbf{r}_\varphi = (r, \theta + \varphi) \rightarrow \mathbf{r}' = (r, \theta')$$

The quantity  $W_\varphi$  is independent of the angle  $\varphi$  (due to the integral over  $\theta'$ ), therefore one can normalize the integral and consider this normalized quantity

$$W \rightarrow \frac{W}{\int_0^{2\pi} d\varphi} = W_\varphi = \int r dr d\theta \Delta_g(\mathbf{r}) \delta(g(\mathbf{r})) e^{iS(r)}$$

The reason why one is allowed to normalize by the orbit volume is because the path integral of the generating functional counts infinitely many times the same phase. The normalization removes this over-counting.



**Functional quantization.** The discussion proceeds by analogy to the previous example. Consider gauge invariance under finite gauge transformations

$$A_\mu^\theta = U(\theta)A_\mu U^{-1}(\theta) - \frac{i}{g} [\partial_\mu U(\theta)]U^{-1}(\theta), \quad U(\theta) = e^{-i\theta^a T^a}$$

The generating functional with no source is

$$W = \int [\mathcal{D}A_\mu] e^{iS_g(A_\mu)}, \quad S_g(A_\mu) = S_g(A_\mu^\theta), \quad \forall \theta$$

where the index  $g$  stands for “gauge”. Notice that the action is gauge-invariant.

One chooses a set of  $\dim G$  gauge functions and imposes the gauge-fixing condition

$$f_a(A_\mu) = 0$$

These functions are such that, for each  $A_\mu$ ,

$$f_a(A_\mu^\theta) = 0$$

has a unique solution  $\theta(x)$ . The gauge functions intersect each orbit only once.

The operator providing unity is

$$\Delta_f^{-1}[A_\mu] = \int [d\theta(x)] \prod_a \delta(f_a(A_\mu^\theta(x))), \quad [d\theta(x)] = \prod_a d\theta_a(x)$$

[r] One requires the measure  $[d\theta(x)]$  to respect the group’s composition law

$$U(\theta)U(\theta') = U(\theta\theta') \implies [d\theta] = [d(\theta\theta')]$$

[r] Therefore

$$\Delta_f^{-1}[A_\mu] = \int \prod_a [df_a] |\det \partial_{f_a} \theta_b| \prod_a \delta(f_a) \implies \Delta_f^{-1} = \det M_{ab}^{-1}$$

from which

$$\Delta_f = |\det \partial_{\theta_b} f_a(A_\mu^\theta)|_{f_a=0} = \det M_{ab}$$

An important property is gauge invariance

$$\Delta_f[A_\mu] = \Delta_f[A_\mu^\theta], \quad \forall \theta$$

This can be checked as follows

$$\begin{aligned} \Delta_f^{-1}[A_\mu^\theta] &= \int [d\theta'(x)] \prod_a \delta(f_a(A_\mu^{\theta\theta'})) = \int [d\theta\theta'(x)] \prod_a \delta(f_a(A_\mu^{\theta\theta'})) \\ &= \int [d\theta''(x)] \prod_a \delta(f_a(A_\mu^{\theta''})) = \Delta_f^{-1}[A_\mu] \end{aligned}$$

where  $d\theta' = d(\theta\theta')$  and  $\theta'' = \theta\theta'$ . Therefore, the generating functional is

$$\begin{aligned} W &= \int [\mathcal{D}A_\mu] e^{iS[A_\mu]} \Delta_f[A_\mu] \int [d\theta(x)] \prod_a \delta(f_a(A_\mu^\theta)) \\ &= \int [d\theta(x)] \int [\mathcal{D}A_\mu] e^{iS[A_\mu]} \Delta_f[A_\mu] \prod_a \delta(f_a(A_\mu^\theta(x))) \\ &= \int [d\theta(x)] \int [\mathcal{D}A_\mu^\theta] e^{iS[A_\mu^\theta]} \Delta_f[A_\mu^\theta] \prod_a \delta(f_a(A_\mu^\theta(x))) \\ &= \int [d\theta(x)] \int [\mathcal{D}A'_\mu] e^{iS[A'_\mu]} \Delta_f[A'_\mu] \prod_a \delta(f_a(A'_\mu(x))) \\ &= \infty \times (\text{smth}) \end{aligned}$$

at the third line, one notices that the action is gauge invariant [r]. At the fourth line, one has renamed the field  $A_\mu^\theta = A'_\mu$ .

One notices that the second integral does not depend on the parameter  $\theta$  and the first integral is the infinite volume of the gauge orbits. One normalizes the generating functional by the orbits' volume and considers the result as the well-defined generating functional for gauge theories

$$W_f = \int [\mathcal{D}A_\mu] \Delta_f[A_\mu] \prod_a \delta(f_a(A(x))) e^{iS[A_\mu]} = \int [\mathcal{D}A_\mu] (\det M_f) \prod_a \delta(f_a(A(x))) e^{iS[A_\mu]}$$

where

$$(M_f)_{ab} = \left. \frac{\delta f_a(A_\mu^\theta)}{\delta \theta_b} \right|_{f_a=0}$$

[r] where absolute value?

This is the Faddeev–Popov prescription. One is not integrating over every field configuration due to the presence of the delta function: one is summing only one single representative of each orbit.

**Determinant.** One may rewrite the coefficients of the exponential as phases. Recall

$$\int \left[ \prod_{j=n}^1 d\theta_j d\bar{\theta}_j \right] e^{i\bar{\theta}_a M_{ab} \theta_b} \propto \det M$$

One introduces a set of Grassmann-odd fields

$$c_a(x), \quad \bar{c}_a(x) = c_a^\dagger(x), \quad a = 1, \dots, \dim G$$

so the determinant is

$$\det M_f \propto \int [dc d\bar{c}] \exp \left[ i \int d^4x d^4y \bar{c}_a(x) [M_f(x, y)]_{ab} c_b(y) \right]$$

These new fields are non-physical degrees of freedom called Faddeev–Popov ghosts.

**Delta function.** One generalizes the gauge-fixing condition from  $f_a(A_\mu) = 0$  to

$$f_a(A_\mu(x)) = B_a(x)$$

where  $B_a(x)$  is a smooth function independent of the gauge field  $A_\mu$ . The delta function is then

$$\delta(f_a(A_\mu) - B_a)$$

Considering that the original generating functional  $W$  is independent of  $B_a$  due to the dependence being inside  $\Delta \Delta^{-1}$ , one may change the normalization by a constant

$$\text{const} \equiv \int [\mathcal{D}B_a] \exp \left[ -\frac{i}{2\xi} \int d^4x B_a(x) B_a(x) \right]$$

Therefore

$$\begin{aligned} W[J] &= \int [\mathcal{D}B_a] [\mathcal{D}A_\mu] (\det M_f) \prod_a \delta(f_a(A_\mu) - B_a) \exp \left[ -\frac{i}{2\xi} \int d^4x B_a(x) B_a(x) \right] e^{iS_g[A_\mu]} \\ &= \int [\mathcal{D}A_\mu] (\det M_f) \exp \left[ -\frac{i}{2\xi} \int d^4x [f_a(A_\mu)]^2 \right] e^{iS_g[A_\mu]} \\ &= \int [\mathcal{D}A_\mu \mathcal{D}c_a \mathcal{D}\bar{c}_a] e^{iS_g[A_\mu]} \\ &\quad \times \exp \left[ -\frac{i}{2\xi} \int d^4x [f_a(A_\mu)]^2 + i \int d^4x d^4y \bar{c}_a(x) (M_f)_{ab}(x, y) c_b(y) \right] \end{aligned}$$

At the second line, one integrates using the delta function. From the third line, one sees that ghosts are propagating fields, but just quantum [r].

The total action is therefore

$$S = S_{\text{gauge}} + S_{\text{gauge-fixing}} + S_{\text{ghost}}$$

where one has

$$\begin{aligned} S_g(A^\mu) &= -\frac{1}{4} \int d^4x \operatorname{Tr}(F_{\mu\nu} F^{\mu\nu}) \\ S_{\text{gf}}(A^\mu) &= -\frac{1}{2\xi} \int d^4x [f_a(A^\mu)]^2 \\ S_{\text{gh}}(A^\mu, c, \bar{c}) &= \int d^4x d^4y \bar{c}_a(x) (M_f)_{ab}(x, y) c_b(y) \end{aligned}$$

with  $\xi$  an arbitrary gauge parameter and

$$f_a(A^\mu) = B_a(x), \quad (M_f)_{ab}(x, y) = \left. \frac{\delta f_a(A_\mu^\theta(x))}{\delta \theta_b(y)} \right|_{f^a=0}$$

**Remark.** The total action is not gauge-invariant due to the presence of the gauge-fixing term.

## Lecture 13

**Example** (Axial gauge). Consider the axial gauge, also called Arnowitt–Fickler gauge<sup>20</sup>,

$$f_a(A^\mu) = A_a^3$$

An infinitesimal gauge transformation by  $\theta_a(x)$  is

$$f^a(A_\mu^\theta) = A_3^a + c^{abc} \theta^b A_3^c - \frac{1}{g} \partial_3 \theta^a \implies f_a(A_\mu^\theta) = -\frac{1}{g} \partial_3 \theta_a$$

[r] where one applies the gauge-fixing condition  $f^a = 0$ . The matrix  $M$  is

$$M_{ab}(x, y) = \left. \frac{\delta f_a(A_\mu(x))}{\delta \theta_b(y)} \right|_{f_a=0} = -\frac{1}{g} \partial_3 [\delta^{(4)}(x - y) \delta_{ab}]$$

Notice how it does not depend on the gauge field  $A_\mu^a$ . This implies that the ghost Lagrangian depends only on the ghosts: the gauge field decouples from the ghosts. The functional integral of the ghosts in the generating functional can be performed and absorbed into the normalization constant, so that the ghosts do not appear anymore.

From this follows that the ghost action describes non-propagating degrees of freedom<sup>21</sup>

$$\begin{aligned} S_{\text{gh}} &= \int d^4x d^4y \bar{c}_a(x) (M_f)_{ab}(x, y) c_b(y) = -\frac{1}{g} \int d^4x d^4y \bar{c}_a \partial_3 [\delta^{(4)}(x - y) \delta_{ab}] c_b(y) \\ &= \frac{1}{g} \int d^4x \partial_3 \bar{c}_a(x) c_a(x) \end{aligned}$$

where one applies the properties of the Dirac delta function under differentiation noting that  $\partial_3$  is done with respect to  $x$ . Since the ghost fields are non-propagating, they can be ignored. In the axial gauge there are no ghosts. This result may be generalized: a gauge-fixing condition that is linear in the gauge field implies that  $M_{ab}$  is independent of the field and the ghosts decouple from the it.

<sup>20</sup>Generally it is written as  $n^\mu A_\mu$  with  $n^\mu$  a unit four-vector specifying the direction.

<sup>21</sup>Notice how there is only one derivative. The equations of motion imply no wave equation:  $\partial_3 \bar{c} = \partial_3 c = 0$ . Any function independent of  $x^3$  is a solution, and there is no sense that the past determines future.

**Example** (Abelian theory). Consider an abelian theory. The gauge transformation is

$$A_\mu^\theta = A_\mu - \frac{1}{g} \partial_\mu \theta$$

Taking a gauge-fixing condition linear in the gauge field permits to decouple the ghosts. This is why QED can be a theory without ghosts. An example of non-linear condition is

$$A_\mu A^\mu = 0$$

**Example** (Abelian Lorenz gauge). Consider the Lorenz gauge

$$f(A_\mu) = \partial_\mu A^\mu = 0$$

The gauge transformation is

$$f_a(A_\mu^\theta) = \partial_\mu A_\theta^\mu = \partial_\mu A^\mu - \frac{1}{g} \square \theta = -\frac{1}{g} \square \theta$$

The matrix is

$$M(x, y) = -\frac{1}{g} \partial_{\theta(x)} [\square \theta(y)] = -\frac{1}{g} \square [\delta^{(4)}(x - y)]$$

The ghost action is

$$S_{\text{gh}} = -\frac{1}{g} \int d^4x \bar{c}(x) \square c(x)$$

This is a free action, describing propagating but non-interacting fields.

**Example** (Non-abelian Lorenz gauge). Consider the Lorenz gauge like before

$$f_a(A_\mu) = \partial_\mu A_a^\mu = 0$$

The gauge function for the transformed field is

$$\begin{aligned} f^a(A_\mu^\theta) &= \partial_\mu \left[ A^{a\mu} + c^{abc} \theta^b A^{c\mu} - \frac{1}{g} \partial^\mu \theta^a \right] = -\frac{1}{g} \partial_\mu (\partial^\mu \theta^a - g c^{abc} \theta^b A^{c\mu}) \\ &= -\frac{1}{g} \partial_\mu (\partial^\mu \theta^a + g c^{abc} A^{b\mu} \theta^c) = -\frac{1}{g} \partial_\mu (D^\mu \theta)^a \end{aligned}$$

where the covariant derivative  $D_\mu$  acts on a field  $\theta_a(x)$  in the adjoint representation<sup>22</sup>. In fact, recalling that

$$D_\mu = \partial_\mu - ig A_\mu^a T^a$$

the covariant derivative acts on the adjoint representation of a field  $\varphi$  as

$$[D_\mu, \varphi] = [\partial_\mu - ig A_\mu^a T^a, \varphi^b T^b] = (\partial_\mu \varphi^a + g c^{abc} A_\mu^b \varphi^c) T^a$$

One obtains

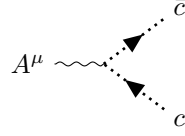
$$M_{ab}(x, y) = -\frac{1}{g} \partial_{\theta^b(y)} [\partial_\mu (D^\mu \theta(x))_a] |_{f_a=0} = -\frac{1}{g} \partial^\mu (D_\mu)_{ab} \delta^{(4)}(x - y)$$

The ghost action is

$$\begin{aligned} S_{\text{gh}} &= -\frac{1}{g} \int d^4x d^4y \bar{c}_a(x) [\partial^\mu D_\mu^{ab} \delta^{(4)}(x - y)] c_b(y) = \frac{1}{g} \int d^4x \partial^\mu \bar{c}_a(x) D_\mu^{ab} c_b(x) \\ &= \frac{1}{g} \int d^4x [\partial^\mu \bar{c}_a(x) \partial_\mu c_a(x) + g c^{acb} \partial_\mu \bar{c}_a(x) A^{c\mu}(x) c^b(x)] \end{aligned}$$

where one integrates by parts. The second addendum is present only for non-abelian theories: it is a three-point vertex between the gauge field and two ghosts:

<sup>22</sup>For the meaning of a field to be in a representation, see <https://physics.stackexchange.com/q/412232>.



The complete action is then

$$S = S_g + S_{gf} + S_{gh} = S_0 + S_I$$

where the free and interacting parts are

$$S_0 = \int d^4x \left[ -\frac{1}{2} \partial_\mu A_\nu^a (\partial^\mu A^{a\nu} - \partial^\nu A^{a\mu}) - \frac{1}{2\xi} (\partial_\mu A^{a\mu})^2 - \bar{c}_a \square c_a \right]$$

$$S_I = \int d^4x \left[ -c^{abc} \partial_\mu A_\nu^a A^{b\mu} A^{c\nu} - \frac{g^2}{4} c^{abc} c^{ade} A_\mu^b A_\nu^c A^{d\mu} A^{e\nu} - g c^{abc} (\partial_\mu \bar{c}^a) c^b A^{c\mu} \right]$$

where the fields  $c$  and  $\bar{c}$  have been redefined to

$$c_a \rightarrow \sqrt{g} c_a$$

The interacting part includes a three-point and a four-point self-coupling vertices and a three-point ghost-field vertex.

## 9.2 Perturbation theory

[r] The kinetic term found using ghosts is invertible. The generating functional is

$$W[J, \eta, \bar{\eta}] = \int [\mathcal{D}A_\mu \mathcal{D}c \mathcal{D}\bar{c}] \exp \left[ iS + i \int d^4x (J_\mu^a A_\mu^a + \bar{\eta}^a c^a + \bar{c}^a \eta^a) \right]$$

As with scalar and spinor fields, one may replace

$$A_\mu^a(x) \rightarrow -i \frac{\delta}{\delta J_\mu^a(x)}, \quad c^a \rightarrow -i \frac{\delta}{\delta \bar{\eta}^a}, \quad \bar{c}^a \rightarrow i \frac{\delta}{\delta \eta^a}$$

to have

$$W[J, \eta, \bar{\eta}] = \exp \left[ iS_I(-i\partial_J, -i\delta_{\bar{\eta}}, i\delta_\eta) \right] W_0[J, \eta, \bar{\eta}]$$

noting that the free generating functionals split

$$W_0[J, \eta, \bar{\eta}] = W_0[J] W_0[\eta, \bar{\eta}]$$

with

$$W_0[J] = \int [\mathcal{D}A_\mu] \exp \left[ i(S_g^{(0)} + S_{gf}) + i \int d^4x J_\mu^a A_\mu^a \right]$$

$$W_0[\eta, \bar{\eta}] = \int [\mathcal{D}c \mathcal{D}\bar{c}] \exp \left[ iS_{gh}^{(0)} + i \int d^4x (\bar{\eta} c + \bar{c} \eta) \right]$$

where the superscript (0) indicates the free part.

**Gauge field propagator.** Consider the free action of the gauge boson field

$$S_g^{(0)} + S_{gf} = \int d^4x \left[ -\frac{1}{2} \partial_\mu A_\nu^a (\partial^\mu A^{a\nu} - \partial^\nu A^{a\mu}) - \frac{1}{2\xi} (\partial_\mu A^{a\mu}) (\partial_\nu A^{a\nu}) \right]$$

$$= \int d^4x \frac{1}{2} A_\nu^a \left[ \eta^{\mu\nu} \square - \left( 1 - \frac{1}{\xi} \right) \partial^\mu \partial^\nu \right] A_\mu^a$$

$$= \int d^4x d^4y \frac{1}{2} A_\nu^a(x) \left[ \eta^{\mu\nu} \square - \left( 1 - \frac{1}{\xi} \right) \partial^\mu \partial^\nu \right] \delta^{(4)}(x-y) \delta_b^a A_\mu^b(y)$$

$$= \int d^4x d^4y \frac{1}{2} A_\nu^a(x) K_{ab}^{\mu\nu}(x, y) A_\mu^b(y)$$

At the second line one has integrated by parts.

The kinetic term is an invertible operator

$$K_{ab}^{\mu\nu}(x-y) = \left[ \eta^{\mu\nu} \square - \left( 1 - \frac{1}{\xi} \right) \partial^\mu \partial^\nu \right] \delta^{(4)}(x-y) \delta_{ab}$$

[r] In fact, its Green's function  $G_{\mu\nu}(x-y)$  that satisfies

$$K^{\mu\nu}(x-y) G_{\nu\rho}(y-z) = \delta_\rho^\mu \delta^{(4)}(x-z)$$

Applying the Fourier transform

$$G_{\nu\rho}(y-z) = \int \frac{d^4 k}{(2\pi)^4} e^{ik(y-z)} G_{\nu\rho}(k)$$

gives

$$\left[ -k^2 \eta^{\mu\nu} + \left( 1 - \frac{1}{\xi} \right) k^\mu k^\nu \right] G_{\nu\rho}(k) = \delta_\rho^\mu$$

which has solution

$$G_{\nu\rho}^{ab}(k) = - \left[ \eta_{\nu\rho} - (1 - \xi) \frac{k_\nu k_\rho}{k^2} \right] \frac{\delta^{ab}}{k^2 + i\varepsilon}$$

Completing the square of the generating functional  $W_0[J]$  gives

$$W_0[J] = \exp \left[ -\frac{i}{2} \int d^4 x d^4 y J_\mu^a(x) G_{ab}^{\mu\nu}(x, y) G_\nu^b(y) \right]$$

The two-point Green's function is

$$G_{ab}^{\mu\nu}(x-y) = \langle 0 | \mathcal{T} \{ A_\mu^a(x) A_\nu^b(y) \} | 0 \rangle = (-i)^2 \frac{\delta^2 W_0[J]}{\delta J_\mu^a(x) \delta J_\nu^b(y)} \Big|_{J=0} = i \Delta_{ab}^{\mu\nu}(x-y)$$

and in Fourier space it is

$$G_{\mu\nu}^{ab}(k) = \langle 0 | \mathcal{T} \{ A_\mu^a(k) A_\nu^b(-k) \} | 0 \rangle = -i \left[ \eta_{\mu\nu} - (1 - \xi) \frac{k_\mu k_\nu}{k^2} \right] \frac{\delta^{ab}}{k^2 + i\varepsilon} = i \Delta_{\mu\nu}^{ab}(k)$$

**Ghost propagator.** The free generating functional is

$$W_0[\eta, \bar{\eta}] = \int [\mathcal{D}c \mathcal{D}\bar{c}] \exp \left[ -i \int d^4 x [\bar{c}_a \square c_a - \bar{\eta}_a c_a - \bar{c}_a \eta_a] \right]$$

Knowing that the kinetic term in momentum space is

$$\delta_{ab} \square \implies -\delta_{ab} k^2$$

one obtains

$$G_{ab}(x-y) = - \int \frac{d^4 k}{(2\pi)^4} \frac{e^{ik(x-y)}}{k^2 + i\varepsilon} \delta_{ab}$$

[r] Completing the quare in the exponent gives

$$W_0[\eta, \bar{\eta}] = \exp \left[ -i \int d^4 x d^4 y \bar{\eta}_a G_{ab}(x-y) \eta_b(y) \right]$$

The two-point Green's function is

$$G^{ab}(x-y) = \langle 0 | \mathcal{T} \{ \bar{c}^a(x) c^b(y) \} | 0 \rangle = \frac{\delta^2 W_0[\eta, \bar{\eta}]}{\delta \bar{\eta}_a(x) \delta \eta_b(y)} \Big|_{\eta=\bar{\eta}=0} = i \Delta^{ab}(x-y)$$

and in momentum space, one has

$$G_{ab}(k) = -\frac{i\delta_{ab}}{k^2 + i\varepsilon} = i\Delta_{ab}(k)$$

which corresponds to the line



## 10 Becchi–Rouet–Stora–Tyutin quantization

[r] See Cheng, §9.3. Adding a gauge-fixing Lagrangian implies that the action is no longer gauge-invariant. One would like to find another symmetry (and the corresponding transformations) for the gauge-fixed action. This symmetry is the Becchi–Rouet–Stora–Tyutin (BRST) symmetry.

Consider a gauge field minimally coupled to some spinor fields also called matter fields. The most general Yang–Mills Lagrangian (gauge group  $SU(n)$ ) is

$$\mathcal{L} = -\frac{1}{4} \text{Tr}(F_{\mu\nu} F^{\mu\nu}) + \bar{\psi}(i \not{D} - m)\psi - \frac{1}{2\xi}(\partial_\mu A^{a\mu})^2 + \bar{c}_a \partial^\mu (D_\mu)_{ab} c_b$$

where

$$F_{\mu\nu}^a = \partial_{[\mu} A_{\nu]}^a + g c^{abc} A_\mu^b A_\nu^c, \quad D_\mu \psi = (\partial_\mu - ig A_\mu^a T^a) \psi, \quad (D_\mu)_{ab} c_b = \partial_\mu c_a + g c^{acb} A_\mu^c c^b$$

It is convenient to express the complex Grassmann-odd ghost fields as real Grassmann-odd fields

$$c_a = \frac{1}{\sqrt{2}}(\rho_a + i\sigma_a), \quad \bar{c}_a = \frac{1}{\sqrt{2}}(\rho_a - i\sigma_a)$$

This implies that the ghost Lagrangian is

$$\mathcal{L}_{\text{gh}} = -i \partial^\mu \rho^a (D_\mu)_{ab} \sigma^b$$

which can be obtained noting that

$$\rho^a \partial^\mu (D_\mu)_{ab} \rho^b = \sigma^a \partial^\mu (D_\mu)_{ab} \sigma^b = 0$$

**Proposition.** The action is invariant under the BRST transformation

$$\begin{aligned} \delta A_\mu^a &= \omega D_\mu^{ab} \sigma_b \\ \delta \rho^a &= -\frac{i}{\xi} \omega \partial^\mu A_\mu^a \\ \delta \sigma^a &= -\frac{g}{2} \omega c^{abc} \sigma^b \sigma^c \\ \delta \psi &= ig \omega (T^a \sigma^a) \psi \end{aligned}$$

where  $\omega$  is a constant Grassmann-odd number and  $\xi$  is the gauge parameter.

In general, infinitesimal transformations can be written as the commutator of a parameter  $\varepsilon$  and the generator  $Q$  of a symmetry

$$\delta \chi = [\varepsilon Q, \chi]$$

In this case, the parameter  $\varepsilon = \omega$  and the generator are Grassmann-odd quantities. This also implies  $Q^2 = 0$  and that applying twice the transformations gives zero.

*Proof.* The action is

$$S = (S_g + S_\psi) + S_{\text{gf}} + S_{\text{gh}}$$

The first two addenda are gauge-invariant. The transformations of the gauge field and the matter fields are the typical (infinitesimal) gauge transformations with parameter  $\theta^a = -g\omega\sigma^a$ . Thus one has

$$\delta_{\text{BRST}}(S_g + S_\psi) = 0$$

The gauge-fixing Lagrangian transforms as

$$\delta L_{\text{gf}} = -\frac{1}{2\xi} \delta(\partial_\mu A^{a\mu})^2 = -\frac{1}{\xi} \partial_\mu A^{a\mu} \partial_\nu (\delta A^{a\nu}) = -\frac{1}{\xi} \partial_\mu A^{a\mu} \partial_\nu [\omega (D^\nu)^{ab} \sigma^b]$$

The ghost Lagrangian transforms as

$$\begin{aligned} \delta \mathcal{L}_{\text{gh}} &= \delta[-i \partial^\mu \rho_a D_\mu^{ab} \sigma_b] = -i \partial^\mu (\delta \rho_a) D_\mu^{ab} \sigma_b - i (\partial^\mu \rho_a) \delta[D_\mu^{ab} \sigma_b] \\ &= -\frac{1}{\xi} \omega \partial^\mu (\partial^\nu A_\nu^a) D_\mu^{ab} \sigma_b - i (\partial^\mu \rho_a) \delta[D_\mu^{ab} \sigma_b] \end{aligned}$$

At the second line, one may integrate by parts the first addendum and see that it cancels the variation of the gauge-fixing Lagrangian.

The variation of the total Lagrangian (up to four-divergences) is

$$\delta\mathcal{L} = -i(\partial^\mu \rho_a) \delta[D_\mu^{ab} \sigma_b]$$

One needs to show that the last variation is zero

$$\begin{aligned} \delta[D_\mu^{ab} \sigma_b] &= \delta[\partial_\mu \sigma^a - g c^{abc} \sigma^b A_\mu^c] \\ &= \partial_\mu (\delta \sigma^a) - g c^{abc} (\delta \sigma^b) A_\mu^c - g c^{abc} \sigma^b (\delta A_\mu^c) \\ &= -\frac{g}{2} \omega c^{abc} \partial_\mu (\sigma^b \sigma^c) + g c^{abc} \frac{g}{2} \omega c^{bef} \sigma^e \sigma^f A_\mu^c - g c^{abc} \sigma^b \omega D_\mu^{cd} \sigma^d \\ &= -\frac{g}{2} \omega c^{abc} \partial_\mu (\sigma^b \sigma^c) + \frac{g^2}{2} \omega c^{abc} c^{bef} \sigma^e \sigma^f A_\mu^c + g \omega c^{abc} \sigma^b (\partial_\mu \sigma^c - g c^{cd} \sigma^d A_\mu^c) \\ &= \frac{g^2}{2} \omega c^{abc} c^{bef} \sigma^e \sigma^f A_\mu^c - g^2 \omega c^{abc} c^{cef} \sigma^b \sigma^e A_\mu^f \\ &= \frac{g^2}{2} \omega c^{abe} c^{bcf} \sigma^c \sigma^f A_\mu^e - g^2 \omega c^{acb} c^{bfe} \sigma^c \sigma^f A_\mu^e \\ &= \frac{g^2}{2} \omega (c^{abe} c^{bcf} + c^{abc} c^{bfe} - c^{abf} c^{bce}) \sigma^c \sigma^f A_\mu^e \\ &= 0 \end{aligned}$$

[r] At the fourth line, the last addendum has a positive sign since  $\omega$  and  $\sigma$  are both Grassmann-odd. At the fifth line, one has considered the two terms without  $A_\mu$  and has recalled that the structure constants  $c^{abc}$  are totally anti-symmetric. At the sixth line, one renames  $c \leftrightarrow e$  in the first addendum, and  $c \leftrightarrow b$  and  $f \leftrightarrow e$  in the second addendum. At the penultimate line, the third addendum one exchanges the two  $\sigma$ s and renames  $c \leftrightarrow f$ .

The result is zero due to the Jacobi identity. This means that the variation of the Lagrangian is zero under BRST transformations.  $\square$

## Lecture 14

**Remark.** The above variation can be written as

$$\delta_{\text{BRST}}(D_\mu \sigma)^a = 0$$

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### 10.1 Slavnov–Taylor identities

The Slavnov–Taylor identities are the generalization of the Ward–Takahashi identity (originally derived in QED) to non-abelian gauge theories.

**Review – Ward identity in QED.** Quantum electrodynamics is a renormalizable theory. Its ultraviolet divergences can be taken care of using the renormalization functions

$$\psi_0 = Z_2^{\frac{1}{2}} \psi, \quad A_0^\mu = Z_3^{\frac{1}{2}} A^\mu, \quad q_0 = k^\varepsilon Z_1 Z_2^{-1} Z_3^{-\frac{1}{3}} q, \quad m_0 = Z_m Z_2^{-1} m$$

Gauge invariance imposes the use of covariant derivatives

$$D_\mu = \partial_\mu + iq A_\mu, \quad g = |q|$$

where  $q$  is the charge (with sign) of the matter field. Assuming that gauge invariance is preserved after quantization implies that the covariant derivative of a matter field  $D_\mu \psi$  has to be protected against quantum corrections

$$D_{\text{R}}^\mu \equiv D_{\text{bare}}^\mu \iff (D_\mu \psi)_{\text{R}} \equiv Z_2^{-\frac{1}{2}} (D_\mu \psi)_{\text{bare}}$$

This implies that

$$(q A_\mu)_{\text{R}} = (q A_\mu)_{\text{bare}} \implies (q A_\mu)_{\text{bare}} = q_0 A_{0\mu} = k^\varepsilon \frac{Z_1}{Z_2} q A_\mu \equiv q A_\mu$$

from which one obtains the Ward–Takahashi identity in QED

$$\boxed{Z_1 = Z_2}$$



**Non-abelian theories.** One assumes that the BRST symmetry is preserved after quantization

$$\delta S = 0 \implies \delta W = 0$$

If the classical action  $S$  is invariant, then so is the effective action (equivalently the generating functional  $W$ ). Since the generating functional contains the classical action, then the phase  $e^{iS}$ , is invariant under BRST transformations. The above statement implies that the functional measure must be invariant.

Consider a generating functional containing the sources for the physical fields, the ghost fields and the composite fields that appear in the BRST transformations

$$W[J, \alpha, \beta, \chi, \bar{\chi}, \kappa, \nu, \lambda, \bar{\lambda}] = \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi} \mathcal{D}\rho \mathcal{D}\sigma] \exp \left[ i \int d^4x (\mathcal{L} + \Sigma) \right]$$

where the source term is

$$\Sigma = J_\mu^a A_\mu^a + \alpha_a \rho_a + \beta_a \sigma_a + \bar{\chi} \psi + \bar{\psi} \chi + \kappa_\mu^a (D^\mu \sigma)^a + \frac{1}{2} \nu_a c^{abc} \sigma_b \sigma_c + \bar{\lambda} T^a \sigma^a \psi + \bar{\psi} T^a \sigma^a \lambda$$

notice that the last four addenda involve composite fields and, within their sources, only  $\kappa$  is Grassmann-odd [r]. The assumption that the generating functional is invariant implies that

$$\delta W = \int d^4x \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi} \mathcal{D}\rho \mathcal{D}\sigma] \delta \Sigma \exp \left[ i \int d^4x (\mathcal{L} + \Sigma) \right] = 0 \implies \delta \Sigma = 0$$

If the transformation of the source term  $\Sigma$  is non-trivial, then the equation above gives useful relations between composite fields. The variation is

$$\begin{aligned} \delta \Sigma &= J_\mu^a \delta A_\mu^a + \alpha_a \delta \rho_a + \beta_a \delta \sigma_a + \bar{\chi} \delta \psi + \delta \bar{\psi} \chi + \kappa_\mu^a \delta (D^\mu \sigma)^a \\ &\quad + \frac{1}{2} \nu_a \delta (c^{abc} \sigma_b \sigma_c) + \bar{\lambda} T^a \delta (\sigma^a \psi) + \delta (\bar{\psi} T^a \sigma^a) \lambda \end{aligned}$$

From the proof of invariance the gauge-fixed action one has  $\delta (D_\mu \sigma)^a = 0$ . The variation of all the composite operators is zero

$$\delta (c^{abc} \sigma_b \sigma_c) = T^a \delta (\sigma^a \psi) = \delta (\bar{\psi} T^a \sigma^a) = 0$$

See Cheng, p. 276 for computations. Inserting the BRST transformations, one finds

$$\delta \Sigma = \omega \left[ J_\mu^a (D^\mu \sigma)^a + \frac{i}{\xi} \alpha_a \partial_\mu A_\mu^a + \frac{g}{2} \beta_a c^{abc} \sigma_b \sigma_c - i g \bar{\chi} (T^a \sigma^a) \psi - i g \bar{\psi} (T^a \sigma^a) \chi \right]$$

In the functional integral, the fields can be replaced with the derivative with respect to the corresponding source. In this way, the source term  $\Sigma$  is no longer function of the integrated fields and can be brought outside the functional integral. The variation of the generating functional is then

$$\delta W = \int d^4x \delta \Sigma W[J, \alpha, \beta, \chi, \bar{\chi}, \kappa, \nu, \lambda, \bar{\lambda}] = 0$$

which implies<sup>23</sup>

$$\int d^4x \omega \left[ J_\mu^a \frac{\delta}{\delta \kappa_\mu^a} + \frac{i}{\xi} \alpha_a \partial_\mu \frac{\delta}{\delta J_\mu^a} + g \beta_a \frac{\delta}{\delta \nu_a} - i g \bar{\chi} \frac{\delta}{\delta \bar{\lambda}} - i g \frac{\delta}{\delta \lambda} \chi \right] W[J, \alpha, \beta, \chi, \bar{\chi}, \kappa, \nu, \lambda, \bar{\lambda}] = 0$$

This is the Slavnov–Taylor identity.

From the identity one may obtain relations between different types of Green’s functions by applying derivatives with respect to the sources and setting them to zero afterwards. This procedure is equivalent to the implication that, since the generating functional  $W$  is invariant under BRST transformations, so are the Green’s function obtained from it

$$\delta W = 0 \implies \delta G^{(n)} = 0$$

<sup>23</sup>The following integral is the variation  $\delta W$  when the integral is multiplied by  $-i$  which cancels the imaginary unit coming from differentiating the exponential.

**Example.** Consider a two-point Green's function for the gauge field

$$(G_{\mu\nu}^{ab})^{(2)}(x) = \langle 0 | \mathcal{T} \{ A_\mu^a(x) A_\nu^b(0) \} | 0 \rangle$$

If the BRST symmetry holds after quantization, one has to find

$$\delta G^{(2)}(x) = 0$$

Explicitly calculating the variation gives

$$\begin{aligned} \delta(G_{\mu\nu}^{ab})^{(2)}(x) &= \delta \langle 0 | \mathcal{T} \{ A_\mu^a(x) A_\nu^b(0) \} | 0 \rangle \\ &= \langle 0 | \mathcal{T} \{ \delta A_\mu^a(x) A_\nu^b(0) \} | 0 \rangle + \langle 0 | \mathcal{T} \{ A_\mu^a(x) \delta A_\nu^b(0) \} | 0 \rangle \\ &= \langle 0 | \mathcal{T} \{ \omega(D_\mu \sigma)^a(x) A_\nu^b(0) \} | 0 \rangle + \langle 0 | \mathcal{T} \{ A_\mu^a(x) \omega(D_\nu \sigma)^b(0) \} | 0 \rangle \\ &= 0 \end{aligned}$$

This implies that the two addenda are related.

**Example.** Consider a four-point Green's function

$$G^{(4)} = \langle 0 | \mathcal{T} \{ \rho^a A_\mu^a \bar{\psi} \psi \} | 0 \rangle$$

The variation is

$$\begin{aligned} \delta G^{(4)} &= -\frac{i}{\xi} \omega \langle \mathcal{T} \{ \partial_\nu A_\mu^\nu A_\mu^a \bar{\psi} \psi \} \rangle - \omega \langle \mathcal{T} \{ \rho^a (D_\mu \sigma)^a \bar{\psi} \psi \} \rangle \\ &\quad + i g \omega \langle \mathcal{T} \{ \rho^a A_\mu^a \bar{\psi} (T^b \sigma^b) \psi \} \rangle + i g \omega \langle \mathcal{T} \{ \rho^a A_\mu^a \bar{\psi} (T^b \sigma^b) \psi \} \rangle \\ &= 0 \end{aligned}$$

This linear combination of different Green's function has to be zero.

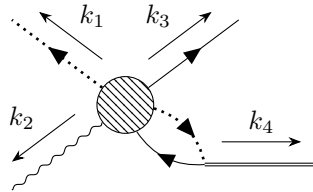
## 10.2 Unitarity

From the relation in the above example, one may derive physically important identities<sup>24</sup>. Thanks to the LSZ reduction formula, one may compute scattering amplitudes with an integral expression of kinetic operators applied to a Green's function (which can be a linear combination of different kinds of Green's functions). The LSZ formula gives a non-trivial expression for a scattering amplitude when the Green's function contains propagators that simplify the kinetic terms. This is necessary since external particles are on-shell: the kinetic terms are zero, while the propagators diverge. One then would like to understand which terms in a Green's function contribute.

Consider the above example. When setting the external particles on-shell, the terms in the second line do not contribute. The reasoning is the following. Consider the Green's function

$$\langle \mathcal{T} \{ \rho(x_1) A_\mu(x_2) \bar{\psi}(x_3) T^a(\sigma^a \psi)(x_4) \} \rangle$$

Notice that  $\sigma^a \psi$  is composite field. In momentum space, there are four momenta associated to the above Green's function



Since the momentum  $k_4$  splits into the sum of two momenta, one never gets a propagator like

$$\frac{1}{\not{k}_4 - m}$$

<sup>24</sup>Since they come from the Slavnov–Taylor identity they are referred to as (being part of the) Slavnov–Taylor identities.

which cancels the associated kinetic term of the form  $\not{k}_4 - m$  in the LSZ reduction formula. Since, on-shell it holds  $\not{k}_4 - m = 0$ , then this Green's function does not contribute on-shell to the scattering amplitude. In general, only the terms linear in each field survive.

Therefore, on-shell the four-point Green's function above gives

$$\frac{i}{\xi} \langle \mathcal{T} \{ \partial_\nu A_a^\nu A_\mu^a \bar{\psi} \psi \} \rangle + \langle \mathcal{T} \{ \rho_a (D_\mu \sigma)_a \bar{\psi} \psi \} \rangle = 0$$

Since the above discussion also applies for the composite operator  $\sigma A_\mu$ , then the covariant derivative can be traded for the ordinary derivative

$$\frac{i}{\xi} \langle \mathcal{T} \{ \partial_\nu A_a^\nu A_\mu^a \bar{\psi} \psi \} \rangle + \langle \mathcal{T} \{ \rho_a \partial_\mu \sigma_a \bar{\psi} \psi \} \rangle = 0$$

In momentum space, the derivatives become momenta. In the 't Hooft–Feynman gauge,  $\xi = 1$ , one may define

$$iT_{\nu\mu}^{ab} \equiv \langle 0 | \mathcal{T} \{ A_\nu^a A_\mu^b \bar{\psi} \psi \} | 0 \rangle, \quad iS^{ab} \equiv \langle 0 | \mathcal{T} \{ \rho^a \sigma^b \bar{\psi} \psi \} | 0 \rangle$$

[r] source?

With these definitions, the Ward identity in momentum space becomes

$$k_1^\nu T_{\nu\mu}^{ab} = iS^{ab}(k_2)_\mu$$

Multiplying by  $k_2^\mu$  gives

$$k_1^\nu T_{\nu\mu}^{ab} k_2^\mu = 0$$

because  $(k_2)^2 = 0$  on-shell.

## Lecture 15

From these two identities, one may show that unitarity is preserved. Unitarity is equivalent to the conservation of probability which implies that the  $S$ -matrix is unitary. The scattering matrix enables one to write

$$|f\rangle = S |i\rangle$$

The unitarity of the  $S$ -matrix can be written as

$$SS^\dagger = I \implies S_{ac} S_{bc}^* = \delta_{ab}$$

The scattering matrix can be written as

$$S_{ab} = \delta_{ab} + i(2\pi)^4 \delta^{(4)}(p_a - p_b) T_{ab}$$

where  $T_{ab}$  is the transition amplitude. Inserting this expression into the unitarity constraint, one obtains the optical theorem

$$\text{Im } T_{ab} = \frac{1}{2} T_{ac} T_{bc}^* (2\pi)^4 \delta^{(4)}(p_a - p_c)$$

The imaginary part of a transition amplitude is the product of the transition amplitudes associated to the transitions to all possible intermediate physical states.

**Optical theorem.** The optical theorem is satisfied through the two Slavnov–Taylor identities above. Consider a fermion-fermion  $\bar{\psi}\psi \rightarrow \bar{\psi}\psi$  scattering in a Yang-Mills theory. To compute the imaginary part of the transition amplitude, one considers the spinor, vector and ghost propagators. These last two are

$$\Delta_{\mu\nu}^{ab}(k) = -\frac{\delta^{ab} \eta_{\mu\nu}}{k^2 + i\varepsilon}, \quad \Delta^{ab}(k) = -\frac{\delta^{ab}}{k^2 + i\varepsilon}$$

One employs the Sokhotski–Plemelj theorem

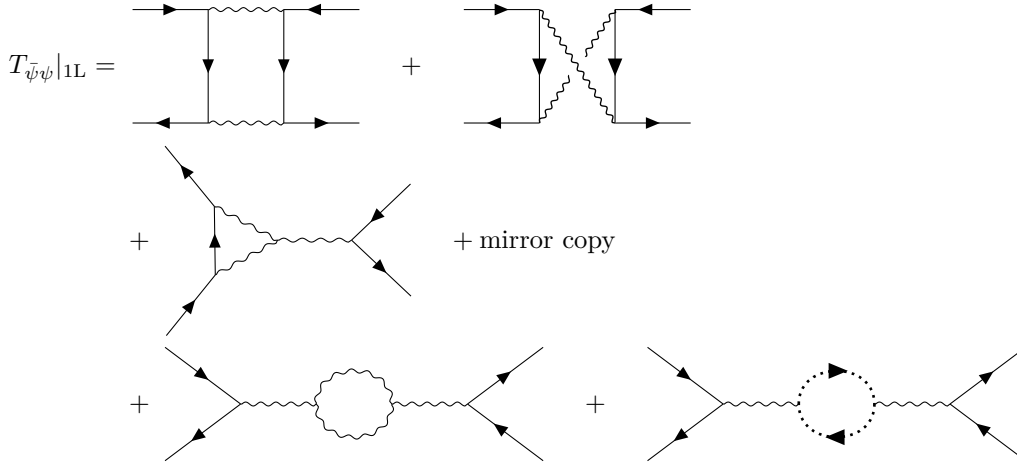
$$\lim_{\varepsilon \rightarrow 0^+} \frac{1}{k^2 + i\varepsilon} = P \frac{1}{k^2} - i\pi\delta(k^2)$$

where  $P$  is the Cauchy principal value. Close to the  $k^0$  singularity, one has

$$\text{Im } \Delta_{\mu\nu}^{ab} = \pi \eta_{\mu\nu} \delta^{ab} \delta(k^2) \theta(k^0), \quad \text{Im } \Delta^{ab} = \pi \delta^{ab} \delta(k^2) \theta(k^0)$$

To compute the imaginary part of a scattering amplitude, one replaces the propagators in the intermediate states by their imaginary parts and multiplies them by the on-shell scattering amplitudes. This is the Cutkowski rule. Concretely, the imaginary part of an amplitude is obtained by cutting the internal propagators of a diagram and then multiply together the on-shell amplitudes obtained from the cuts.

**Example.** The one-loop contributions to the fermion-fermion scattering are



To compute only the imaginary part, one has to cut the internal propagators. The imaginary part is the product of tree-level amplitudes (see Peskin for proof). In the first five diagrams, there are products of amplitudes of the type  $\bar{\psi}\psi AA$  which corresponds to the transition amplitude  $T$ , while the last diagram involves the amplitude  $\bar{\psi}\psi\rho\sigma$  which corresponds to the scattering amplitude  $S$ . Therefore, the left-hand side of the optical theorem is

$$\text{Im } T_{\bar{\psi}\psi} = \int d\phi_2 \left[ \frac{1}{2} T_{\mu\nu}^{ab} (T^*)_{\rho\sigma}^{ab} \eta^{\mu\rho} \eta^{\nu\sigma} - S^{ab} (S^*)^{ab} \right]$$

where the measure  $d\phi_2$  denotes the massless two-particle phase space. The minus comes from cutting a fermionic loop.

For the right-hand side of the optical theorem, the intermediate states cannot be ghosts, but only vector bosons. One builds the diagrams with these amplitudes



and sums over all the intermediate physical states. The two diagrams correspond both to the transition amplitude  $T$ . In the sum, one has to consider the polarization of the vector bosons. The general classical solution to the massless Klein–Gordon equation for vector bosons is

$$A^\mu(x) = \sum_{r=0}^3 \int \frac{d^3k}{(2\pi)^3} \frac{1}{\sqrt{2\omega_k}} [\varepsilon_r^\mu(\mathbf{k}) a_r(\mathbf{k}) e^{-ikx} + \varepsilon_r^\mu(\mathbf{k}) a_r^\dagger(\mathbf{k}) e^{ikx}]_{k^0=\omega}$$

where  $\varepsilon_r$  are the polarization four-vectors. The physical polarization are  $\varepsilon_1$  and  $\varepsilon_2$ . Therefore, the sum over the intermediate physical states involves only on the physical polarizations. In the

sum there is a product of the polarization vectors. In the end, one obtains a right-hand side equal to

$$\text{RHS} = \frac{1}{2} \int d\phi_2 T_{\mu\nu}^{ab}(T^*)_{\rho\sigma}^{ab} P^{\mu\rho}(k_1) P^{\nu\sigma}(k_2)$$

where one has

$$P^{\mu\rho}(k_j) = \sum_{r=1,2} \varepsilon_r^\mu(\mathbf{k}_j) \varepsilon_r^\rho(\mathbf{k}_j)$$

One may check the validity of the optical theorem for this particular example

$$\text{Im } T \stackrel{?}{=} \text{RHS}$$

A useful choice of polarization basis is

$$\varepsilon_0^\mu = (1, 0, 0, 0) = n^\mu, \quad \varepsilon_i^\mu = (0, \varepsilon_i)$$

where

$$\varepsilon_i \cdot \varepsilon_j = \delta_{ij}, \quad \varepsilon_{1,2} \cdot \mathbf{k} = 0$$

One sets the longitudinal polarization vector to be

$$\varepsilon_3^\mu = \frac{k^\mu - (nk)n^\mu}{[(nk)^2 - k^2]^{\frac{1}{2}}}$$

The first vector  $\varepsilon_0$  is the scalar polarization, while the transverse polarizations are  $\varepsilon_{1,2}$ . The polarization vectors satisfy

$$\sum_{r=0}^3 \zeta_r \varepsilon_r^\mu(\mathbf{k}) \varepsilon_r^\nu(\mathbf{k}) = -\eta^{\mu\nu}, \quad \zeta_r = \begin{cases} -1, & r = 0 \\ 1, & r = 1, 2, 3 \end{cases}$$

Therefore, one has

$$\begin{aligned} P^{\mu\rho}(k_1) &= \sum_{r=1,2} \varepsilon_r^\mu(\mathbf{k}_1) \varepsilon_r^\rho(\mathbf{k}_1) = \sum_{r=0}^3 \zeta_r \varepsilon_r^\mu(\mathbf{k}_1) \varepsilon_r^\rho(\mathbf{k}_1) + \varepsilon_0^\mu(\mathbf{k}_1) \varepsilon_0^\rho(\mathbf{k}_1) - \varepsilon_3^\mu(\mathbf{k}_1) \varepsilon_3^\rho(\mathbf{k}_1) \\ &= -\eta^{\mu\rho} + \varepsilon_0^\mu(\mathbf{k}_1) \varepsilon_0^\rho(\mathbf{k}_1) - \varepsilon_3^\mu(\mathbf{k}_1) \varepsilon_3^\rho(\mathbf{k}_1) = -\eta^{\mu\rho} + \frac{(nk)(k^\mu n^\rho + n^\mu k^\rho) - k^\mu k^\rho}{(kn)^2} \end{aligned}$$

The second addendum is the contribution of non-physical polarizations. Inserting the above expression into the right-hand side of the optical theorem, and applying the Ward identities, one finds the left-hand side,  $\text{Im } T$ .

**Remark.** The optical theorem holds thanks to the Ward identities which arise from the BRST symmetry being preserved in the quantum theory.

**Remark.** Non-physical degrees of freedom are exactly compensated by ghosts. The term  $SS$  present in the imaginary part  $\text{Im } T$  can be rewritten on the right-hand side to see that it balances the non-physical polarizations.

## 11 Quantum electrodynamics

See Ramond, §8.2 and consider the Euclidean formalism.

### 11.1 One-loop structure

Consider an abelian gauge theory based on  $U(1)$  with one massive Dirac spinor field  $\psi$  minimally coupled to the gauge field  $A^\mu$ . The covariant derivative is<sup>25</sup>

$$D_\mu = \partial_\mu + iqA_\mu$$

<sup>25</sup>Note that Ramond uses  $q = -e$  for electrons, with  $e > 0$ .

The Lagrangian is

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \bar{\psi}(i\not{D} - m)\psi - \frac{1}{2\xi}(\partial_\mu A^\mu)^2$$

The generating functional is

$$W[J, \eta, \bar{\eta}] = \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ i \int d^4x (\mathcal{L} + J_\mu A^\mu + \bar{\psi}\eta + \bar{\eta}\psi) \right]$$

**Wick rotation.** See Ramond, §5.2. When going to Euclidean space, one would like to keep the covariant derivative

$$x^0 \equiv -ix_E^0, \quad \partial_0 \equiv i\partial_{0E}, \quad A_0 \equiv iA_{0E}$$

The Dirac matrices are rotated as

$$\gamma_0 \equiv i\gamma_{0E} \implies \{\gamma_E^\mu, \gamma_E^\nu\} = -2\delta^{\mu\nu}$$

The Euclidean generating functional becomes

$$W_E[J, \eta, \bar{\eta}] = \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ - \int d^4x (\mathcal{L}_E + J_\mu A^\mu - i\bar{\eta}\psi - i\bar{\psi}\eta) \right]$$

where the Euclidean Lagrangian is<sup>26</sup>

$$\mathcal{L}_E = \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2\xi}(\partial_\mu A^\mu)^2 + \bar{\psi}(\not{D} + im)\psi - iqA_\mu(\bar{\psi}\gamma^\mu\psi)$$

**Euclidean Feynman rules.** The gauge field propagator in the 't Hooft–Feynman gauge,  $\xi = 1$ , is

$$\frac{\delta_{\mu\nu}}{p^2} = \begin{array}{c} \mu \quad \xrightarrow{p} \quad \nu \\ \bullet \text{---} \text{wavy line} \text{---} \bullet \end{array}$$

The spinor propagator is

$$-\frac{i}{\not{p} + m} = \begin{array}{c} \bar{\psi} \quad \longrightarrow \quad \psi \\ \bullet \text{---} \text{arrow} \text{---} \bullet \end{array}$$

The vertex is

$$iqk^\varepsilon\gamma^\mu = \begin{array}{c} \swarrow \quad \searrow \\ \psi \quad \bar{\psi} \\ \text{---} \text{wavy line} \end{array}$$

One has to use Furry's theorem also: all odd-fermion loops are zero.

**Power counting.** The superficial degree of divergence obtained through power counting is

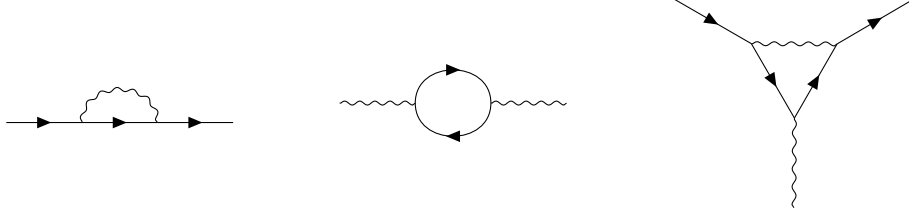
$$D = 4 - E_B - \frac{3}{2}E_F$$

where  $E_B$  is the number of external boson lines while  $E_F$  is the number of external fermion lines. Pure vector diagrams,  $E_F = 0$ , are divergent for  $E_B = 2$  and  $E_B = 4$ : the self-energy of the gauge field and the light-light diagram. For two external fermion lines,  $E_F = 2$ , the only divergent diagrams have  $E_B = 1$  and  $E_B = 0$ : the vertex correction and the self-energy of the fermion field.

<sup>26</sup>Notice that the partial derivative and the interaction term have different signs. The Euclidean Lagrangian is defined starting from the Euclidean action where one sends  $x_E \rightarrow x_E$  in the integral for the derivative. See notes last year, Lecture 14, p. 18.

### 11.1.1 Regularization

At one-loop, one has to evaluate the diagrams



To perform the computations, one has to

- compute the combinatorial factor,
- apply Feynman combining,
- move to spherical coordinates,
- evaluate momentum integrals of the form

$$\int \frac{d^d l}{(2\pi)^d} \frac{(l^2)^m}{(l^2 + A^2)^n}$$

See Ramond, appendix B.

- consider the simple poles of the regularization parameter  $\varepsilon$ .

**One-loops contributions.** The fermion self-energy is

$$\Sigma(p) = -\frac{i}{\varepsilon} \frac{q^2}{16\pi^2} (\not{p} + 4m) + \text{finite}$$

The gauge field self-energy is

$$\Pi_{\mu\nu}(p) = \frac{q^2}{12\pi^2} (p_\mu p_\nu - p^2 \delta_{\mu\nu}) \frac{1}{\varepsilon} + \text{finite}$$

The vertex function is

$$\Gamma_\mu(p, q) = i \frac{q^3}{16\pi^2} k^\varepsilon \gamma_\mu \frac{1}{\varepsilon} + \text{finite}$$

### 11.1.2 Renormalization

One may apply the BPHZ renormalization. Let the bare quantities be

$$\psi_0 = Z_2^{\frac{1}{2}} \psi, \quad A_0^\mu = Z_3^{\frac{1}{2}} A^\mu, \quad q_0 = Z_1 Z_2^{-1} Z_3^{-\frac{1}{2}} q k^\varepsilon, \quad m_0 = Z_m Z_2^{-1} m, \quad \xi_0^{-1} = Z_\xi Z_3^{-1} \xi^{-1}$$

The bare Lagrangian is

$$\begin{aligned} \mathcal{L}_0 &= \frac{1}{4} (\partial_\mu A_\nu - \partial_\nu A_\mu)^2 + \bar{\psi} \gamma^\mu \partial_\mu \psi + i m \bar{\psi} \psi - i q A_\mu \bar{\psi} \gamma^\mu \psi + \frac{1}{2\xi} (\partial_\mu A^\mu)^2 \\ &\quad + (Z_3 - 1) (\partial_\mu A_\nu - \partial_\nu A_\mu)^2 + (Z_2 - 1) \bar{\psi} \gamma^\mu \partial_\mu \psi + (Z_m - 1) i m \bar{\psi} \psi \\ &\quad - (Z_1 - 1) i q A_\mu \bar{\psi} \gamma^\mu \psi + (Z_\xi - 1) \frac{1}{2\xi} (\partial_\mu A^\mu)^2 \\ &= \mathcal{L}_R + \mathcal{L}_{\text{ct}} \end{aligned}$$

In the following, let  $K_j = Z_j - 1$ .

**Fermion field: mass term.** The divergence of the mass term from the fermion self-energy is cancelled by the mass counter term  $imK_m$  to have

$$-\frac{i}{\varepsilon} \frac{q^2}{4\pi^2} - imK_m = \text{finite} \implies k_m = -\frac{q^2}{4\pi^2} \frac{1}{\varepsilon} + \text{finite} \implies \boxed{Z_m = 1 - \frac{q^2}{4\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

Remember that the minus sign of the counter term comes from the minus sign of the action in the exponential of the generating functional.

**Fermion field: kinetic term.** From the self-energy one has

$$-\frac{i}{\varepsilon} \frac{q^2}{16\pi^2} \not{p} - iK_2 \not{p} = \text{finite} \implies k_2 = -\frac{q^2}{16\pi^2} \frac{1}{\varepsilon} + \text{finite} \implies \boxed{Z_2 = 1 - \frac{q^2}{16\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

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**Gauge field.** [r] From the counter term Lagrangian, one reads

$$\begin{aligned} \mathcal{L}_{\text{ct}} &\propto (Z_3 - 1) \frac{1}{4} (\partial_\mu A_\nu - \partial_\nu A_\mu)^2 + \frac{1}{2\xi} (Z_\xi - 1) (\partial_\mu A^\mu)^2 \\ &= \frac{1}{2} K_3 \partial_\mu A_\nu \partial^\mu A^\nu - \frac{1}{2} K_3 \partial_\mu A_\nu \partial^\nu A^\mu + \frac{1}{2\xi} K_\xi \partial_\mu A^\mu \partial_\nu A^\nu \\ &= -\frac{1}{2} K_3 A^\nu \square A_\nu - \frac{1}{2} \left[ \frac{K_\xi}{\xi} - K_3 \right] A^\nu \partial_\mu \partial_\nu A^\mu + 4\text{-div} \end{aligned}$$

At the third line, one has integrated all terms by parts. Remembering the minus sign coming from the action in the exponential of the generating functional

$$-S \propto \frac{1}{2} K_3 A^\nu \square A_\nu + \frac{1}{2} \left[ \frac{K_\xi}{\xi} - K_3 \right] A^\nu \partial_\mu \partial_\nu A^\mu$$

In momentum space, this becomes

$$\frac{1}{2} K_3 A_\nu (-p^2) A^\nu + \frac{1}{2} \left[ \frac{K_\xi}{\xi} - K_3 \right] A^\nu (-p_\mu p_\nu) A^\mu$$

The counter term contribution is then

$$-\frac{1}{2} 2 \left[ \delta_{\mu\nu} K_3 p^2 + \left( \frac{K_\xi}{\xi} - K_3 \right) p_\mu p_\nu \right]$$

where the 2 is a symmetry factor [r]. Therefore, the total contribution is

$$\frac{q^2}{12\pi^2} (p_\mu p_\nu - p^2 \delta_{\mu\nu}) \frac{1}{\varepsilon} - \left[ \delta_{\mu\nu} K_3 p^2 + \left( \frac{K_\xi}{\xi} - K_3 \right) p_\mu p_\nu \right] = \text{finite}$$

The terms proportional to  $p^2$  give

$$-p^2 \delta_{\mu\nu} \left[ \frac{q^2}{12\pi^2} \frac{1}{\varepsilon} + K_3 \right] \implies K_3 = -\frac{q^2}{12\pi^2} \frac{1}{\varepsilon} \implies \boxed{Z_3 = 1 - \frac{q^2}{12\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

while the ones proportional to  $p_\mu p_\nu$  give

$$\left[ \frac{q^2}{12\pi^2} \frac{1}{\varepsilon} - \left( \frac{K_\xi}{\xi} - K_3 \right) \right] p_\mu p_\nu = \text{finite} \implies \frac{q^2}{12\pi^2} \frac{1}{\varepsilon} + K_3 - K_\xi = \text{finite}, \quad \xi = 1$$

In the 't Hooft–Feynman gauge, the above relation implies no gauge parameter renormalization

$$K_\xi = \text{finite} = 0 \implies Z_\xi = 1$$

The divergences of the two terms of self-energy  $\Pi_{\mu\nu}$  are both cancelled by  $K_3$  in the 't Hooft–Feynman gauge.



**Vertex function.** The vertex counter term is

$$-S \propto K_1 i q k^\varepsilon A_\mu \bar{\psi} \gamma^\mu \psi$$

which leads to a counter term

$$K_1 i q k^\varepsilon \gamma^\mu$$

From this one finds

$$K_1 = -\frac{q^2}{16\pi^2} \frac{1}{\varepsilon} + \text{finite} \implies \boxed{Z_1 = 1 - \frac{q^2}{16\pi^2} \frac{1}{\varepsilon} + \text{finite}}$$

**Remark.** Notice that at one-loop, the Ward identity of QED holds:  $Z_1 = Z_2$ .

### 11.1.3 Beta function

The coupling constant is the elementary charge  $g = |q|$ . In the minimal subtraction scheme (i.e. no finite parts), one has

$$g_0 = Z_1 Z_2^{-1} Z_3^{-\frac{1}{2}} g k^\varepsilon = Z_3^{-\frac{1}{2}} g k^\varepsilon = \left[ 1 - \frac{g^2}{12\pi^2} \frac{1}{\varepsilon} \right]^{-\frac{1}{2}} g k^\varepsilon = \left[ 1 + \frac{g^2}{24\pi^2} \frac{1}{\varepsilon} + o(g^2) \right] g k^\varepsilon$$

From this, one notices the coefficient of the simple pole

$$a_1 = \frac{g^3}{24\pi^2}$$

Recall that for the  $\lambda\varphi^4$  theory it holds

$$\lambda_0 = k^{2\varepsilon} \left[ \lambda + \sum_{r=1}^{\infty} \frac{a_r(\lambda)}{\varepsilon^r} \right] \implies \beta(\lambda) = 2[\lambda d_\lambda - 1] a_1$$

where the factor of 2 comes from the 2 in the exponent of the energy scale. For QED, one has

$$\beta(g) = [g d_g - 1] a_1 = 2a_1 = \frac{g^3}{12\pi^2}$$

The beta function enters in the renormalization group flow equation

$$k d_k g = \beta(g)$$

Since the coupling constant is positive, so is the beta function and as such at higher energies the coupling constant gets stronger.

**Landau pole.** One may solve the above equation. One considers  $q^2(k)$ . The equation is

$$k d_k g^2 = 2g k d_k g = \frac{g^4}{6\pi^2} \implies \frac{1}{g^4} d_k g^2 = \frac{1}{6\pi^2 k}$$

Integrating from  $g(k = k_0) = g_0$  to  $g$  gives

$$\boxed{g^2(k) = \frac{g_0^2}{1 - \frac{g_0^2}{6\pi^2} \ln \frac{k}{k_0}}}$$

The running coupling constant of QED has a Landau pole at

$$k = k_0 e^{\frac{6\pi^2}{g_0^2}}$$

When getting close to the pole, the coupling constant gets bigger than 1 and perturbation theory cannot be applied, so this result should not be trusted in such region.

There is no asymptotic freedom in QED because the coupling constant grows stronger with the energy scale. This lack of asymptotic freedom implies a well-defined notion of asymptotic states (of the  $S$ -matrix): at large distances, which correspond to low energies, the coupling constant is weak [r].

## 11.2 Ward identities

See Ramond, §8.3. Since QED is an abelian theory, one may use the results for non-abelian theories and set the structure constants to zero,  $c^{abc} = 0$ . The generating functional is

$$W[J, \eta, \bar{\eta}] = \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi}] \exp \left[ - \int d^4x (\mathcal{L} + J_\mu A^\mu - i\bar{\eta}\psi - i\bar{\psi}\eta) \right] \\ \times \int [\mathcal{D}\rho \mathcal{D}\sigma] \exp \left[ - \int d^4x (\mathcal{L}_{\text{gh}} - \alpha\rho - \beta\sigma) \right]$$

Since ghosts decouple, one introduces the ghost fields with Lagrangian<sup>27</sup>

$$\mathcal{L}_{\text{gh}} = \rho \square \sigma$$

The total Lagrangian

$$\mathcal{L} = \mathcal{L}_g + \mathcal{L}_{\text{gf}} + \mathcal{L}_{\text{gh}} + \mathcal{L}_\psi$$

is invariant under BRST transformations<sup>28</sup>

$$\begin{aligned} \delta A_\mu &= \omega \partial_\mu \sigma \\ \delta \rho &= -\frac{\omega}{\xi} (\partial_\mu A^\mu) \\ \delta \sigma &= 0 \\ \delta \psi &= i g \omega \sigma \psi \\ \delta \bar{\psi} &= -i g \omega \sigma \bar{\psi} \end{aligned}$$

Recall that  $\mathcal{L}_g + \mathcal{L}_\psi$  is invariant thanks to gauge invariance. The remaining terms are also invariant

$$\delta(\mathcal{L}_{\text{gf}} + \mathcal{L}_{\text{gh}}) = 0$$

With these transformations, one may obtain the Ward identity imposing that

$$\delta W[J, \eta, \bar{\eta}, \alpha, \beta] = 0$$

where one has

$$W[J, \eta, \bar{\eta}, \alpha, \beta] = \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi} \mathcal{D}\rho \mathcal{D}\sigma] e^{-S} \exp \left[ - \int d^4x (J_\mu A^\mu - i\bar{\eta}\psi - i\bar{\psi}\eta - \alpha\rho - \beta\sigma) \right]$$

The variation of the generating functional gives

$$\begin{aligned} 0 &= -J_\mu \omega \partial_\mu \sigma + i\bar{\eta}(i g \omega \sigma \psi) + i(-i g \omega \sigma \bar{\psi})\eta - \alpha \frac{\omega}{\xi} (\partial_\mu A^\mu) \\ &= -J_\mu \omega \partial_\mu \sigma - g \omega \sigma \bar{\eta} \psi + g \omega \sigma \bar{\psi} \eta - \alpha \frac{\omega}{\xi} (\partial_\mu A^\mu) \end{aligned}$$

It is useful to rewrite this Ward identity in terms of the effective action  $\Gamma$ . In the  $\lambda\varphi^4$  theory, one has

$$\Gamma[\varphi_c] = -Z[J] + \int d^4x J \varphi_c, \quad J(x) = \delta_{\varphi(x)} \Gamma[\varphi]$$

Doing the same for QED, one has

$$\Gamma[A_\mu, \psi, \bar{\psi}, \rho, \sigma] = -Z[J, \eta, \bar{\eta}, \alpha, \beta] + \int d^4x [-J_\mu A^\mu + i\bar{\eta}\psi + i\bar{\psi}\eta + \alpha\rho + \beta\sigma]$$

where all the fields are classical and a subscript  $c$  is understood. One has

$$J_\mu = -\delta_{A^\mu} \Gamma, \quad i\eta = \delta_{\bar{\psi}} \Gamma, \quad i\bar{\eta} = -\delta_\psi \Gamma, \quad \alpha = -\delta_\rho \Gamma, \quad \beta = -\delta_\sigma \Gamma$$

[r] Imposing the invariance under the BRST transformations of the effective action  $\Gamma$  above gives the Ward identity

$$\delta_{A_\mu} \Gamma \partial_\mu \sigma + i g \delta_\psi \Gamma \sigma \psi - i g \sigma \bar{\psi} \delta_{\bar{\psi}} \Gamma - \frac{1}{\xi} (\partial_\mu A^\mu) \delta_\rho \Gamma = 0$$

where one recalls that  $\delta Z = 0$ .

<sup>27</sup>Where one has renamed  $i\rho \rightarrow \rho$ .

<sup>28</sup>See previous note. Notice a missing factor of  $i$  in the transformation of  $\rho$  and that  $(\partial_\mu A^\mu)_M \rightarrow -(\partial_\mu A_\mu)_E$  which preserves the minus sign when performing a Wick rotation.

**Effective action of QED.** Recall that

$$e^{-\Gamma} = \int [\mathcal{D}A_\mu \mathcal{D}\psi \mathcal{D}\bar{\psi} \mathcal{D}\sigma \mathcal{D}\rho] \exp \left[ - \int d^4x (\mathcal{L}_{\text{QED}} + \mathcal{L}_{\text{gh}}) \right]$$

The effective action is

$$\begin{aligned} \Gamma[A_\mu, \psi, \bar{\psi}, \rho, \sigma] = & \int d^4x d^4y \left[ \rho(x) \Delta^{-1}(x-y) \sigma(y) \right. \\ & + \frac{1}{2} A_\mu(x) \Delta_{\mu\nu}^{-1}(x-y) A_\nu(y) + \bar{\psi}(x) S^{-1}(x-y) \psi(y) \\ & \left. + g \int d^4z \bar{\psi}(x) A^\rho(y) \Gamma_\rho(x, y, z) \psi(z) + o(g) \right] \end{aligned}$$

Since the ghosts do not interact with the fields, then their inverse propagator  $\Delta^{-1} = \square$  gets no loop corrections. The other propagators and the vertex function are all corrected.

**Kinetic term of the gauge field.** One may use the Ward identity to find the general expression of the kinetic term of the gauge field  $\Delta_{\mu\nu}^{-1}$ . One may set the fermion fields to zero. After an integration by parts of the first term, the Ward identity becomes

$$0 = [\partial_\mu \delta_{A_\mu} \Gamma] \sigma + \frac{1}{\xi} (\partial_\mu A^\mu) \delta_\rho \Gamma = \partial_\mu (\Delta_{\mu\nu}^{-1} A_\nu) \sigma + \frac{1}{\xi} (\partial_\mu A^\mu) \Delta^{-1} \sigma$$

Recalling that  $\Delta^{-1} = \square$ , in momentum space one has

$$ik_\mu \Delta_{\mu\nu}^{-1} A_\nu \sigma + \frac{1}{\xi} ik_\mu A^\mu (-k^2) \sigma = 0 \implies \boxed{k^\mu \Delta_{\mu\nu}^{-1}(k) - \frac{1}{\xi} k_\nu k^2 = 0}$$

This constraint must hold at every loop order.

One may obtain the general form of the kinetic term. Starting from the ansatz

$$\Delta_{\mu\nu}^{-1}(k) = A(k^2) \delta_{\mu\nu} + B(k^2) k_\mu k_\nu$$

From the constraint above, one obtains

$$A(k^2) = [1 - B(k^2)] k^2, \quad \xi = 1$$

Therefore

$$\boxed{\Delta_{\mu\nu}^{-1}(k) = k^2 \delta_{\mu\nu} - (\delta_{\mu\nu} k^2 - k_\mu k_\nu) B(k^2)}$$

When computing quantum corrections to the kinetic term, one has to determine  $B(k^2)$  which is always multiplied by the projection operator

$$\Pi_{\mu\nu} = \delta_{\mu\nu} k^2 - k_\mu k_\nu$$

that satisfies

$$\Pi_{\mu\nu} \Pi_{\nu\rho} = \Pi_{\mu\rho} k^2$$

The projector is not invertible, but the first addendum in the kinetic term  $\Delta^{-1}$  comes from gauge-fixing and makes the kinetic term invertible.

**Renormalization functions.** Another application of the Ward identity is the following. Let  $A_\mu = 0$ . Then the Ward identity is

$$\delta_{A_\mu} \Gamma|_{A_\mu=0} \partial_\mu \sigma + iq \delta_\psi \Gamma \sigma \psi + iq \bar{\psi} \sigma \delta_{\bar{\psi}} \Gamma = 0$$

One gets a non-trivial contribution from the vertex function. Inserting the effective action gives

$$\begin{aligned} 0 = & g \int d^4x d^4z \bar{\psi}(x) \Gamma_\rho(x, w, z) \psi(z) \partial_\rho \sigma(w) - ig \int d^4x \bar{\psi}(x) S^{-1}(x-w) \sigma(w) \psi(w) \\ & + ig \int d^4y \bar{\psi}(w) \sigma(w) S^{-1}(w-y) \psi(y) \end{aligned}$$

In momentum space, one has

$$S^{-1}(p) - S^{-1}(q) = (q - p)^\rho \Gamma_\rho(p, q - p, q)$$

[r] check. Knowing that  $S^{-1}$  is renormalized by  $Z_2$  and  $\Gamma$  is renormalized by  $Z_1$ , this Ward identity gives

$$Z_1 = Z_2$$

## Lecture 17

## 12 Yang–Mills gauge theories

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See Ramond, §§8.5, 8.6. The Euclidean Lagrangian is

$$\begin{aligned} \mathcal{L} = & \frac{1}{4} \text{Tr}(F_{\mu\nu} F^{\mu\nu}) + \frac{1}{2\xi} \text{Tr}[(\partial_\mu A^\mu)^2] + i \partial_\mu \bar{c}^a \partial_\mu c^a - \frac{i}{2} g k^\varepsilon c^{abc} A_\mu^c \bar{c}^a \partial_\mu c^b \\ & - \frac{i}{2} g k^\varepsilon c^{abc} \bar{c}^a c^b \partial_\mu A_\mu^c + \bar{\psi}(\not{\partial} + im)\psi + i g k^\varepsilon A_\mu^a \bar{\psi}_i \gamma^\mu T_{ij}^a \psi_j \end{aligned}$$

**Ghost Lagrangian.** The ghost Lagrangian is

$$\mathcal{L}_{\text{gh}} \propto \partial_\mu \bar{c}^a D_\mu c^a, \quad D_\mu = \partial_\mu + i g T^a A_\mu^a$$

From the second addendum in the covariant derivative — recalling that it acts on the adjoint representation —, one obtains

$$\begin{aligned} \partial_\mu \bar{c}^a c^{abc} c^b A_\mu^c &= \frac{1}{2} \partial_\mu \bar{c}^a c^{abc} c^b A_\mu^c + \frac{1}{2} \partial_\mu \bar{c}^a c^{abc} c^b A_\mu^c \\ &= \frac{1}{2} c^{abc} (\partial_\mu \bar{c}^a c^b - \bar{c}^a \partial_\mu c^b) A_\mu^c - \frac{1}{2} c^{abc} \bar{c}^a c^b \partial_\mu A_\mu^c \\ &= \frac{1}{2} c^{abc} \bar{c}^a \partial_\mu c^b A_\mu^c - \frac{1}{2} c^{abc} \bar{c}^a c^b \partial_\mu A_\mu^c \end{aligned}$$

At the second line, one has integrated by parts the second addendum.

One defines

$$\partial_\mu \bar{c}^a c^b - \bar{c}^a \partial_\mu c^b = \bar{c}^a \partial_\mu c^b, \quad A \equiv -\frac{i}{2} g k^\varepsilon c^{abc} \bar{c}^a c^b \partial_\mu A_\mu^c$$

One notices that the remaining term  $A$  is proportional to the gauge term [r] and is zero when in the Lorenz gauge.

### 12.1 Feynman rules

**Propagators.** The fermion propagator is

$$-i \frac{\delta_{ij}}{\not{p} + m} = i \begin{array}{c} \xrightarrow{p} \\ \longrightarrow \end{array} \begin{array}{c} i \\ j \end{array}$$

The gauge field propagator is

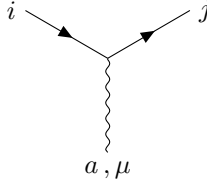
$$\frac{\delta_{ab}}{p^2} \left[ \delta_{\mu\nu} - (1 - \xi) \frac{p_\mu p_\nu}{p^2} \right] = \begin{array}{c} \mu \xrightarrow{p} \nu \\ a \text{-----} b \end{array}$$

where the  $a, b$  indices are the indices in the adjoint representation  $a, b = 1, \dots, n^2 - 1$ . The ghost propagator is

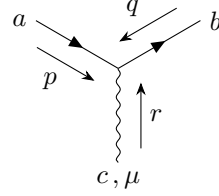
$$-i \frac{\delta_{ab}}{p^2} = \begin{array}{c} \xrightarrow{p} \\ a \text{-----} b \end{array}$$

See Ramond, appendix C.

**Vertices.** Consider the vertices from the expansion of the exponential  $e^{-S}$ . The fermion-boson vertex is

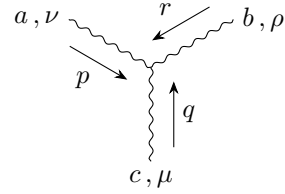
$$-igk^\varepsilon \gamma_\mu (T^a)_{ij} =$$


The ghost-boson vertex is

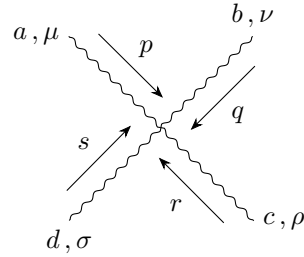


$$= \begin{cases} -gk^\varepsilon c^{abc} q_\mu, & \text{with } A \\ \frac{1}{2} gk^\varepsilon c^{abc} (p - q)_\mu, & \text{without } A \end{cases}$$

with all momenta entering. The adjoint indices have to be put in the same order every time: in this case clockwise. The three-point self-coupling of the gauge field is

$$-igk^\varepsilon c^{abc} [(r - q)_\mu \delta_{\nu\rho} + (q - p)_\rho \delta_{\mu\nu} + (p - r)_\nu \delta_{\rho\mu}] =$$


The four-point vertex is



$$= -g^2 k^{2\varepsilon} [c^{abe} c^{cde} (\delta_{\mu\rho} \delta_{\nu\sigma} - \delta_{\nu\rho} \delta_{\mu\sigma}) + c^{cbe} c^{ade} (\delta_{\mu\rho} \delta_{\nu\sigma} - \delta_{\nu\mu} \delta_{\rho\sigma}) + c^{dbe} c^{cae} (\delta_{\rho\sigma} \delta_{\nu\mu} - \delta_{\nu\rho} \delta_{\mu\sigma})]$$

## 12.2 Renormalizability

**Power counting.** Consider a diagram with

- $L$  loops,
- $V_1$  first-kind vertices,
- $V_2$  second-kind vertices,
- $V_3$  third-kind vertices,
- $V_4$  fourth-kind vertices,
- $I$  internal gauge propagators  $A^\mu$ ,
- $i$  internal fermion propagators  $\psi$ ,
- $I_g$  internal ghost propagators  $c$ ,
- $E$  external gauge lines  $A_\mu$ ,
- $e$  external fermion lines  $\psi$ ,

- $E_g$  external ghost lines  $c$ .

The number of loops is given by

$$L = I + i + I_g - (V_1 + V_2 + V_3 + V_4) + 1$$

There are various constraints

$$2V_1 = 2i + e \implies i = V_1 - \frac{1}{2}e$$

$$2V_2 = 2I_g + E_g$$

$$V_1 + V_2 + 3V_3 + 4V_4 = 2I + E$$

The superficial degree of divergence is

$$\begin{aligned} D &= 4L + V_2 + V_3 - 2I - i - 2I_g \\ &= 4(I + i + I_g - V_1 - V_2 - V_3 - V_4 + 1) - 2I - i - 2I_g + V_2 + V_3 \\ &= 4 + 2I + 3i + 2I_g - 4V_1 - 3V_2 - 3V_3 - 4V_4 \\ &= 4 - E - \frac{3}{2}e - E_g \end{aligned}$$

Notice that it does not depend on the internal lines, but only on the external lines: there is a finite number of topologies of Green's functions that diverge.

**Topologies.** There are five different topologies

1. Consider pure gauge,  $e = E_g = 0$ . Then  $d = 4 - E \geq 0$  for  $E = 1, 2, 3, 4$ . Tadpoles,  $E = 1$ , are not considered [r].
2. Consider pure fermions,  $E = E_g = 0$ . Then  $D = 4 - \frac{3}{2}e \geq 0$ . Since fermions must appear as bilinears, then  $e = 2$ .
3. Let  $E = e = 0$ . Then  $D = 4 - E_g \geq 0$  for  $E_g = 2, 4$  since ghosts are also Grassmann-odd fields and must come in pairs. [r] diagr, there is no four-point ghost vertex, so if the second diagram diverges, it is a problem; however it is not due to Ward identity.
4. Let  $e = 2$  and  $E_g = 0$ . Then  $D = 1 - E \geq 0$  for  $E = 1$ .
5. Let  $e = 0$  and  $E_g = 2$ . Then  $D = 2 - E \geq 0$  for  $E = 1$ .

The topologies resemble terms in the Lagrangian so they can be renormalized through counter terms. Yang-Mills theories are renormalizable. These theories describe three of the four fundamental interactions.

### 12.3 One-loop regularization

The renormalization is done after using dimensional regularization. This implies that massless tadpoles are zero. The one-loop diagrams are the following.

**Gauge propagator.** The diagrams contributing to the gauge propagator are



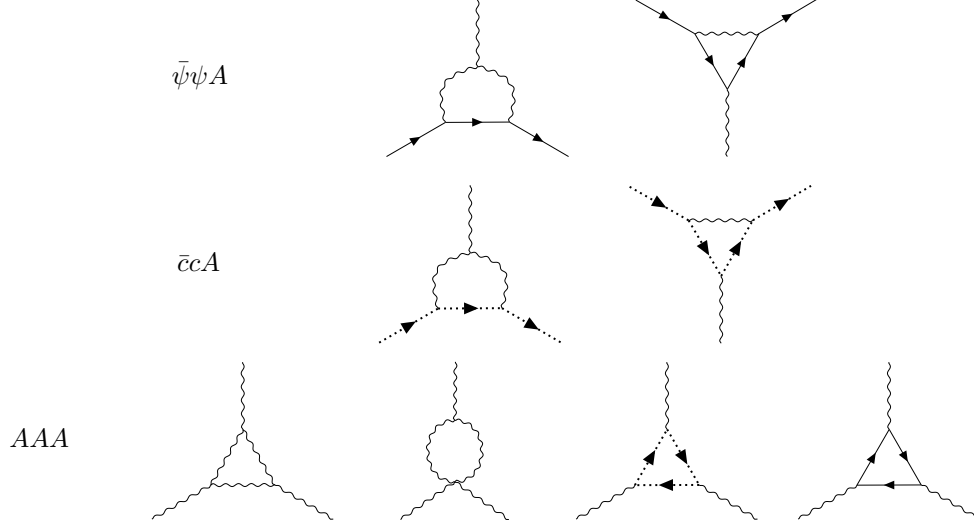
**Fermion propagator.** The diagram contributing to the fermion propagator is



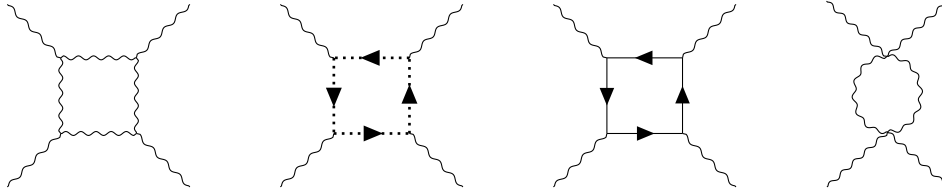
**Ghost propagator.** The diagram contributing to the ghost propagator is



**Three-point vertex.** There are eight diagrams



**Four-point vertex.** There are four diagrams



**BPHZ renormalization.** Starting from the bare Lagrangian, one has

$$\begin{aligned}
 \mathcal{L}_0 = & \frac{1}{4}(\partial_\mu A_{0\nu}^a - \partial_\nu A_{0\mu}^a)^2 - g'_0 c^{abc} \partial_\mu A_{0\nu}^a A_{0\mu}^b A_{0\nu}^c + \frac{1}{2} g_0''^2 c^{abc} c^{ade} A_{0\mu}^b A_{0\nu}^c A_{0\mu}^d A_{0\nu}^e \\
 & + \frac{1}{2\xi_0} (\partial_\mu A_{0\mu}^a)^2 + i \partial_\mu \bar{c}_0^a \partial_\mu c_0^a - \frac{i}{2} g_0''' c^{abc} A_{0\mu}^c \bar{c}_0^a \partial_\mu c_0^b - \frac{i}{2} g_0^{(4)} k^\varepsilon c^{abc} \bar{c}_0^a c_0^b \partial_\mu A_{0\mu}^c \\
 & + \bar{\psi}_0 \not{\partial} \psi_0 + i m_0 \bar{\psi}_0 \psi_0 + i g_0 A_{0\mu}^a \bar{\psi}_0 \gamma_\mu T^a \psi_0 \\
 = & \mathcal{L}_R + \mathcal{L}_{ct}
 \end{aligned}$$

The most general Lagrangian has bare different couplings for the interactions. The renormalized quantities are given by

$$\begin{aligned}
 \psi_0 &= Z_2^{\frac{1}{2}} \psi \\
 A_{0\mu}^a &= Z_3^{\frac{1}{2}} A_\mu^a \\
 c_0^a &= Z_6^{\frac{1}{2}} c^a \\
 m_0 &= Z_m Z_2^{-1} m \\
 \frac{1}{\xi_0} &= \frac{1}{\xi} Z_\xi Z_3^{-1} \\
 g_0 &= Z_1 Z_2^{-1} Z_3^{-\frac{1}{2}} g k^\varepsilon \\
 g'_0 &= Z_4 Z_3^{-\frac{3}{2}} g k^\varepsilon \\
 g_0'' &= Z_5^{\frac{1}{2}} Z_3^{-1} g k^\varepsilon
 \end{aligned}$$

$$g_0''' = Z_7 Z_6^{-1} Z_3^{-\frac{1}{2}} g k^\varepsilon$$

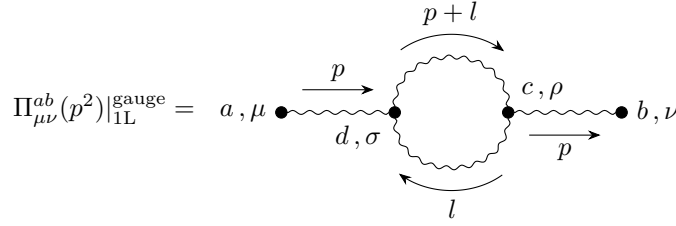
$$g_0^{(4)} = Z_8 Z_6^{-1} Z_3^{-\frac{1}{2}} g k^\varepsilon$$

The counter term Lagrangian has a similar form to the bare and renormalized Lagrangians.

### 12.3.1 Two-point gauge Green's function

The two-point gauge Green's function gets contributions from three diagrams, one for each kind of loop.

**Gauge loop.** Consider the two-point gauge Green's function with a gauge loop



One has to multiply two three-point vertices. The colour indices of the two are taken both clockwise<sup>29</sup>. Recall that the formula given previously is for entering momenta. Therefore

$$\begin{aligned} \Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{gauge}} &= \frac{1}{2!} (-igk^\varepsilon)^2 c^{acd} c^{bdc} [(-p-l-p)_\sigma \delta_{\mu\rho} + (p-l)_\rho \delta_{\mu\sigma} + (l-(-p-l))_\mu \delta_{\rho\sigma}] \\ &\quad \times [(p-l)_\rho \delta_{\nu\sigma} + (-p-l-p)_\sigma \delta_{\nu\rho} + (l+p+l)_\nu \delta_{\rho\sigma}] \\ &\quad \times \int \frac{d^d l}{(2\pi)^d} \frac{1}{l^2(p+l)^2} \\ &= -\frac{1}{2} g^2 k^{2\varepsilon} c^{acd} c^{bdc} \int \frac{d^d l}{(2\pi)^d} \frac{N_{\mu\nu}(l, p)}{l^2(p+l)^2} \end{aligned}$$

with

$$N_{\mu\nu}(l, p) = (10 - 8\varepsilon) l_\mu l_\nu + (5 - 4\varepsilon)(l_\mu p_\nu + l_\nu p_\mu) - 2(1 + \varepsilon) p_\mu p_\nu + [(p-l)^2 + (2p+l)^2] \delta_{\mu\nu}$$

[r] check. Notice that the combinatorial factor is inside the three-point vertex. Applying Feynman combining and using a table of integrals gives

$$\begin{aligned} \Pi_{\mu\nu}^{ab}(p)|_{1L}^{\text{gauge}} &= \int_0^1 dx \left\{ \delta_{\mu\nu} (9 - 3\varepsilon) \frac{\Gamma(-1 + \varepsilon)}{[p^2 x(1-x)]^{-1+\varepsilon}} \right. \\ &\quad \left. + \frac{\Gamma(\varepsilon)}{[p^2 x(1-x)]^\varepsilon} [\delta_{\mu\nu} p^2 (5 - 2x(1-x)) + p_\mu p_\nu (-2(1 + \varepsilon) - (10 - 8\varepsilon)x(1-x))] \right\} \end{aligned}$$

[r] Noting that

$$\Gamma(-1 + \varepsilon) = \frac{\Gamma(1 + \varepsilon)}{\varepsilon(-1 + \varepsilon)} \sim -\frac{1}{\varepsilon}, \quad \Gamma(\varepsilon) \sim \frac{1}{\varepsilon}$$

one may find

$$\Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{gauge}} = \frac{g^2}{32\pi^2} c^{acd} c^{bcd} \frac{1}{\varepsilon} \left[ \frac{19}{6} \delta_{\mu\nu} p^2 - \frac{11}{3} p_\mu p_\nu \right] + o(\varepsilon^{-1})$$

<sup>29</sup>When writing the structure constants. Notice how the pairs of momenta are taken counter-clockwise.



**Ghost loop.** Consider the two-point gauge Green's function with a ghost loop

$$\Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{ghost}} = \begin{array}{c} \text{Diagram: A two-point function with external wavy lines labeled } a, \mu \text{ and } b, \nu. \text{ The left vertex is connected to a ghost loop (dashed circle) with momentum } p \text{ entering and } p+l \text{ leaving. The right vertex is connected to the same loop with momentum } l \text{ entering and } p \text{ leaving. The loop has two vertices labeled } d \text{ and } c. \end{array}$$

This diagram is computed including the  $A$  term. The left vertex gives

$$-gk^\varepsilon c^{dca}(-p-l)_\mu$$

the right one gives

$$-gk^\varepsilon c^{cdb}(-l)_\nu$$

Therefore

$$\begin{aligned} \Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{ghost}} &= g^2 k^{2\varepsilon} c^{dca} c^{cdb} (-i)^2 (-1) \int \frac{d^d l}{(2\pi)^d} \frac{(p+l)_\mu l_\nu}{l^2 (p+l)^2} \\ &= \boxed{\frac{g^2}{32\pi^2} c^{acd} c^{bcd} \left[ \frac{1}{6} p^2 \delta_{\mu\nu} + \frac{1}{3} p_\mu p_\nu \right] \frac{1}{\varepsilon} + o(\varepsilon^{-1})} \end{aligned}$$

The  $-i$  comes from the ghost propagators, while  $-1$  comes from the Grassmann-odd loop integral.

The computation without the  $A$  term is slightly different. The two vertices contribute with

$$\frac{1}{2} gk^\varepsilon c^{dca} (2l+p)_\mu, \quad \frac{1}{2} gk^\varepsilon c^{cdb} (2l+p)_\nu$$

Thus

$$\begin{aligned} \Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{ghost}} &= \frac{g^2}{4} k^{2\varepsilon} c^{dca} c^{cdb} \int \frac{d^d l}{(2\pi)^d} \frac{(2l+p)_\mu (2l+p)_\nu}{l^2 (l+p)^2} \\ &= \frac{g^2}{32\pi^2} c^{acd} c^{bcd} \left[ \frac{1}{6} p^2 \delta_{\mu\nu} - \frac{2}{3} p_\mu p_\nu \right] \frac{1}{\varepsilon} + o(\varepsilon^{-1}) \end{aligned}$$

The only difference is in the coefficient of the longitudinal part: there is one term  $p_\mu p_\nu$  less. This part is taken care of by the gauge-fixing part. Therefore, the fact that it is different does not matter.

## Lecture 18

**Sum of contributions.** The sum of the two above contributions is

$$\Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{gauge}} + \Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{ghost}} = \frac{g^2}{32\pi^2} c^{acd} c^{bcd} \frac{10}{3} (p^2 \delta_{\mu\nu} - p_\mu p_\nu) \frac{1}{\varepsilon} + o(\varepsilon^{-1})$$

where one has included the  $A$  term. Notice that the parenthesis is the projection operator on transverse degrees of freedom.

**Fermion loop.** Consider the two-point gauge Green's function with a fermion loop

$$\Pi_{\mu\nu}^{ab}(p^2)|_{1L}^{\text{fermion}} = \begin{array}{c} \text{Diagram: A two-point function with external wavy lines labeled } a, \mu \text{ and } b, \nu. \text{ The left vertex is connected to a fermion loop (solid circle) with momentum } p \text{ entering and } p+l \text{ leaving. The right vertex is connected to the same loop with momentum } l \text{ entering and } p \text{ leaving. The loop has two vertices labeled } j \text{ and } i. \end{array}$$

The two vertices correspond to

$$-igk^\varepsilon \gamma_\mu (T^a)_{ji}, \quad -igk^\varepsilon \gamma_\nu (T^b)_{ij}$$

They are similar to the ones of QED up to a colour term  $T$ . Therefore

$$\begin{aligned} \Pi_{\mu\nu}^{ab}(p^2)|_{\text{1L}}^{\text{fermion}} &= (T_f^a)_{ji} (T_f^b)_{ij} \Pi_{\mu\nu}^{\text{QED}} \\ &= -\text{Tr}(T_f^a T_f^b) \frac{g^2}{16\pi^2} (\delta_{\mu\nu} p^2 - p_\mu p_\nu) \frac{4}{3} \frac{1}{\varepsilon} + o(\varepsilon^{-1}) \end{aligned}$$

where the generators  $T_f$  are in a generic representation  $f$  of the fermions.

**Total contribution.** The total contribution is

$$\begin{aligned} \Pi_{\mu\nu}^{ab}(p^2)|_{\text{1L}} &= \frac{g^2}{16\pi^2} (p^2 \delta_{\mu\nu} - p_\mu p_\nu) \left[ \frac{5}{3} c^{acd} c^{bcd} - \frac{4}{3} \text{Tr}(T_f^a T_f^b) \right] \frac{1}{\varepsilon} + o(\varepsilon^{-1}) \\ &= \frac{g^2}{16\pi^2} (p^2 \delta_{\mu\nu} - p_\mu p_\nu) \left[ \frac{5}{3} c_{\text{adj}} - \frac{4}{3} c_f \right] \delta^{ab} \frac{1}{\varepsilon} + o(\varepsilon^{-1}) \end{aligned}$$

At the second line, one has used the colour structure of the  $\text{SU}(n)$  group.

### 12.3.2 Color structure

See Peskin, §15.4.

**Proposition** (Dynkin index). [r] The Dynkin index is  $c_f$  defined through

$$\text{Tr}(T_f^a T_f^b) = c_f \delta^{ab}$$

**Proposition.** It holds

$$c^{acd} c^{bcd} = c_{\text{adj}} \delta^{ab}$$

*Proof.* In the adjoint representation, it holds

$$(T^a)_{bc} = -ic^{abc}$$

therefore one has

$$\text{Tr}(T^a T^b) = (T^a)_{cd} (T^b)_{dc} = -ic^{acd} (-ic^{bdc}) \implies c_{\text{adj}} \delta^{ab} = c^{acd} c^{bcd}$$

□

**Proposition** (Casimir operator). Knowing

$$(T^2)_{ij} = (T^a T^a)_{ij} = T_{ik}^a T_{kj}^a$$

the Casimir operator  $T^2$  is such that

$$[T^2, T^b] = 0, \quad \forall b$$

The Casimir index  $c_2$  is defined through

$$T_f^2 = c_2(f) I$$

**Proposition.** For the adjoint representation, the Casimir index is the Dynkin index

$$c_2(\text{adj}) = c_{\text{adj}}$$

**Proposition.** For a representation  $f$  it holds

$$c_2(f) \dim f = c_f \dim G$$

*Proof.* From the Dynkin index, one traces (also) over the group indices

$$\text{Tr}(T_f^a T_f^a) = c_f \dim G$$

Tracing over the definition of Casimir index, one gets

$$\text{Tr}(T_f^2) = \text{Tr}(T_f^a T_f^a) = c_2(f) \dim G$$

Thus the thesis follows.  $\square$

**Example.** Consider  $G = \text{SU}(n)$  and its fundamental representation  $f$ . The dimensions are

$$\dim G = n^2 - 1, \quad \dim f = n$$

From the identity above, one has

$$c_2(\text{fond})n = c_{\text{fond}}(n^2 - 1)$$

It is convention to consider  $c_{\text{fond}} = \frac{1}{2}$  and

$$\text{Tr}(T_{\text{fond}}^a T_{\text{fond}}^a) = \frac{1}{2} \delta^{ab}$$

Therefore

$$c_2(\text{fond}) = \frac{n^2 - 1}{2n}$$

**Proposition.** In the adjoint representation of  $\text{SU}(n)$ , it holds

$$c_2(\text{adj}) = c_{\text{adj}} = n$$

*Proof.* The adjoint representation can be built from the fundamental representation  $n$  and the anti-fundamental representation  $\bar{n}$  recalling that the product of such two representations can be decomposed as the direct sum of irreducible representations

$$n \otimes \bar{n} = \sum_i r_i = \text{trivial} \oplus \text{adj}$$

The object that transform under  $n \times \bar{n}$  carry an index for each  $A_{a\bar{a}}$ . Therefore, a generator can be written as

$$T_{n \otimes \bar{n}}^a = T_n^a \otimes I_{\bar{n}} + I_n \otimes T_{\bar{n}}^a$$

Therefore, the Casimir operator is

$$T_{n \times \bar{n}}^2 = T_{n \times \bar{n}}^a T_{n \times \bar{n}}^a = T_n^2 \otimes I + I \otimes T_{\bar{n}}^2 + 2T_n^a \otimes T_{\bar{n}}^a$$

Taking the trace gives

$$\text{Tr}(T_{n \times \bar{n}}^2) = [c_2(n) + c_2(\bar{n})] \dim n \dim \bar{n} = [c_2(n) + c_2(\bar{n})]n^2 = 2c_2(n)n^2$$

The trace can also be written differently

$$\text{Tr}(T_{n \times \bar{n}}^2) = \sum_i c_2(r_i) \dim r_i = c_2(1) \dim 1 + c_2(\text{adj}) \dim \text{Adj} = 0 + c_2(\text{adj})(n^2 - 1)$$

Comparing these two, one finds

$$2c_2(n)n^2 = c_2(\text{adj})(n^2 - 1) \implies \frac{n^2 - 1}{n}n^2 = c_2(\text{adj})(n^2 - 1) \implies c_2(\text{adj}) = n$$

Noting that  $c_2(\text{adj}) = c_{\text{adj}}$  due to a previous proposition, one obtains the thesis.  $\square$

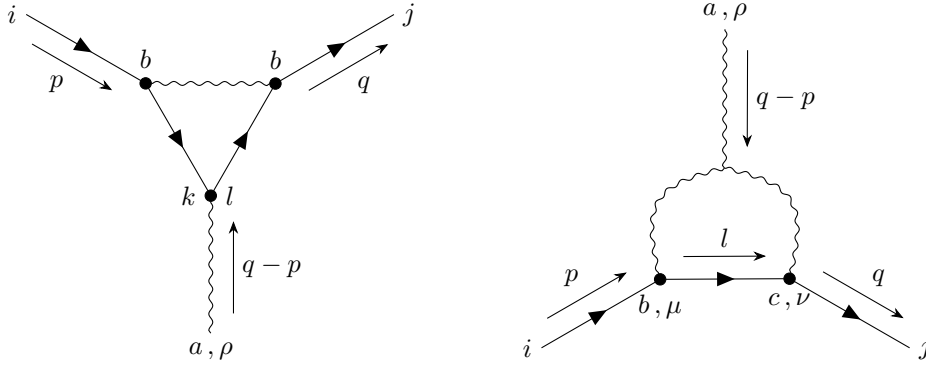
Consider the two-point fermion Green's function

The diagram shows a horizontal line with two vertices labeled 'a' and 'k'. A wavy line connects the two vertices. The incoming momentum is 'i' and the outgoing momentum is 'j'. The internal momentum is 'p'.

Apart from the colour factors, the computation is the same as in QED. Therefore

$$\begin{aligned}\Sigma_{ij}(p^2)|_{1\text{L}} &= (T_f^a)_{ik}(T_f^a)_{kj}\Sigma^{\text{QED}}(p^2) = -(T_f^a T_f^a)_{ij} \frac{i}{\varepsilon} \frac{g^2}{16\pi^2} (\not{p} + 4m) + o(\varepsilon^{-1}) \\ &= \left[ -\frac{ig^2}{16\pi^2} c_2(f) \delta_{ij} (\not{p} + 4m) \frac{1}{\varepsilon} + o(\varepsilon^{-1}) \right]\end{aligned}$$

**Spinor-vector vertex.** The spinor-vector vertex gets contributions from two diagrams



The first diagram is

$$\Gamma_{\rho ij}^a(p, q)|_{1\text{L}} = (T_f^b)_{ik}(T_f^a)_{kl}(T_f^b)_{lj}\Gamma_{\rho}^{\text{QED}} = (T_f^b T_f^a T_f^b)_{ij}\Gamma_{\rho}^{\text{QED}}$$

The colour structure is

$$\begin{aligned} T_f^b T_f^a T_f^b &= [T_f^b, T_f^a] T_f^b + T_f^a T_f^b T_f^b = i c^{bad} T_f^d T_f^b + T_f^a c_2(f) \\ &= \frac{i}{2} c^{bad} [T_f^d, T_f^b] + c_2(f) T_f^a = -\frac{1}{2} c^{bad} c^{dbg} T_f^g + c_2(f) T_f^a \\ &= -\frac{1}{2} c_{\text{adj}} T_f^a + c_2(f) T_f^a \end{aligned}$$

Therefore

$$\Gamma_{\rho ij}^a(p, q)|_{1L} = -igk^\varepsilon \gamma_\rho T_f^a \left[ c_2(f) - \frac{1}{2}c_{\text{adj}} \right] \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} + o(\varepsilon^{-1})$$

The second diagram uses the vertices

$$-igk^\varepsilon \gamma_\mu (T^b)_{ij}, \quad -igk^\varepsilon c^{abc}[(r-q)_\mu \delta_{\nu\rho} + (q-p)_\rho \delta_{\mu\nu} + (p-r)_\nu \delta_{\rho\mu}]$$

Therefore

$$\Gamma_{\rho ij}^a(p, q)|_{\text{IL}}^{\text{II}} = -g^3 k^{3\varepsilon} c^{abc} (T^b)_{ik} (T^c)_{kj} \int \frac{d^d l}{(2\pi)^d} \gamma^\mu \frac{1}{\not{p} + m} \gamma^\nu \frac{N_{\mu\nu\rho}(p, q)}{(p-l)^2 (q-l)^2}$$

where one has

$$N_{\mu\nu\rho}(p, q) = (q + p - 2l)_\rho \delta_{\mu\nu} + (l - 2p + q)_\nu \delta_{\mu\rho} + (l + p - 2q)_\mu \delta_{\nu\rho}$$

The colour structure is

$$c^{abc}(T_f^b)_{ik}(T_f^c)_{kj} = \frac{1}{2}c^{abc}[T_f^b, T_f^c] = \frac{i}{2}c^{abc}c^{bcd}T_f^d = \frac{i}{2}c_{\text{adj}}T_f^a$$

See Peskin also. One obtains

$$\Gamma_{\rho ij}^a(p, q)|_{1\text{L}}^{\text{II}} = -igk^\varepsilon \gamma_\rho T_f^a \frac{3}{2}c_{\text{adj}} \frac{g^2}{16\pi^2} \frac{1}{\varepsilon} + o(\varepsilon^{-1})$$

The sum of the two diagrams is

$$\boxed{\Gamma_{\rho ij}^a(p, q)|_{1\text{L}} = -igk^\varepsilon \gamma_\rho (T_f^a)_{ij} \frac{g^2}{16\pi^2} [c_{\text{adj}} + c_2(f)] \frac{1}{\varepsilon} + o(\varepsilon^{-1})}$$

**Other vertices.** The other vertices produce divergences that are already taken care of the counter terms of the above divergences. [r]

## 12.4 One-loop renormalization

The bare Lagrangian is split

$$\mathcal{L}_0 = \mathcal{L}_R + \mathcal{L}_{\text{ct}}$$

where the counter term Lagrangian is

$$\mathcal{L}_{\text{ct}} = K_3 \frac{1}{4} (\partial_\mu A_\nu^a - \partial_\nu A_\mu^a)^2 + K_2 \bar{\psi} \not{\partial} \psi + K_m i m \bar{\psi} \psi + K_\xi \frac{1}{2} (\partial_\mu A_\mu^a)^2 + K_1 i g k^\varepsilon A_\mu^a \bar{\psi} \gamma_\mu T_f^a \psi$$

Let  $K_j \equiv Z_j - 1$ .

**Fermion mass.** The mass gets a contribution from the fermion Green's function at one-loop

$$\Sigma_{ij}(p^2)|_{1\text{L}} \implies -i \frac{g^2}{16\pi^2} c_2(f) 4m \frac{1}{\varepsilon} = -i \frac{g^2}{4\pi^2} c_2(f) m \frac{1}{\varepsilon}$$

Recalling the minus sign from the exponential of the counter term action, one obtains

$$-im \left[ K_m + \frac{g^2}{4\pi^2} c_2(f) \frac{1}{\varepsilon} \right] = \text{finite} \implies \boxed{Z_m = 1 - \frac{g^2}{4\pi^2} c_2(f) \frac{1}{\varepsilon} + \text{finite}}$$

**Fermion field.** The kinetic term gets a contribution from the fermion Green's function

$$\Sigma_{ij}(p^2)|_{1\text{L}} \implies -i \frac{g^2}{4\pi^2} c_2(f) \frac{1}{\varepsilon} \not{p}$$

[r] Recalling that  $\partial_\mu \leftrightarrow ip_\mu$  when doing Fourier transformations, the counter term contribution is  $-iK_2 \not{p}$ , already including the minus sign from the exponential. Therefore

$$-i \not{p} \left[ K_2 + \frac{g^2}{16\pi^2} c_2(f) \frac{1}{\varepsilon} \right] = \text{finite} \implies \boxed{Z_2 = 1 - \frac{g^2}{16\pi^2} c_2(f) \frac{1}{\varepsilon} + \text{finite}}$$

**Gauge field.** The divergent contribution is

$$\Pi_{\mu\nu}^{ab}(p^2)|_{1\text{L}} = \frac{g^2}{16\pi^2} (p^2 \delta_{\mu\nu} - p_\mu p_\nu) \left[ \frac{5}{3} c_{\text{adj}} - \frac{4}{3} c_f \right] \delta^{ab} \frac{1}{\varepsilon} + o(\varepsilon^{-1})$$

From the counter term Lagrangian

$$\mathcal{L}_{\text{ct}} \propto K_3 \frac{1}{4} (\partial_\mu A_\nu^a - \partial_\nu A_\mu^a)^2$$

One integrates by parts and goes to momentum space

$$\mathcal{L}_{\text{ct}} \rightsquigarrow \frac{1}{2} K_3 A_\mu^a (p^2 \delta_{\mu\nu} - p_\mu p_\nu) A_\nu^a$$

The counter term renormalization gives

$$(p^2 \delta_{\mu\nu} - p_\mu p_\nu) \delta^{ab} \left[ -\frac{1}{2} K_3 \cdot 2 + \frac{g^2}{16\pi^2} \left( \frac{5}{3} c_{\text{adj}} - \frac{4}{3} c(f) \right) \frac{1}{\varepsilon} \right] = \text{finite}$$

from which one has

$$Z_3 = 1 + \frac{g^2}{16\pi^2} \left( \frac{5}{3} c_{\text{adj}} - \frac{4}{3} c(f) \right) \frac{1}{\varepsilon} + \text{finite}$$

The gauge parameter  $\xi$  does not need to be renormalized when considering the  $A$  term. One obtains

$$\Pi_{\mu\nu}^{ab}(p^2)|_{1\text{L}} = \frac{g^2}{16\pi^2} \left[ \frac{5}{3} c_{\text{adj}} + ? c(f) \right] \frac{1}{\varepsilon} (p^2 \delta_{\mu\nu} - 2p_\mu p_\nu) \delta^{ab}$$

The extra factor of 2 in the projection operator (the parenthesis) means renormalizing the gauge parameter. In fact, the counter term gives

$$K_\xi \frac{1}{2} (\partial_\mu A_\mu^a)^2 \rightarrow -\frac{1}{2} K_\xi A_\nu^a \partial_\mu \partial_\nu A_\mu^a \rightarrow \frac{1}{2} K_\xi A_\nu^a p_\mu p_\nu A_\mu^a$$

One may separate the extra term of the projection operator from the actual projection operator and renormalize the gauge parameter

$$Z_\xi = 1 - \frac{g^2}{16\pi^2} \left[ \frac{5}{2} c_{\text{adj}} + ? c(f) \right] \frac{1}{\varepsilon} + \text{finite}$$

**Vertex.** The divergent contribution is

$$\Gamma_{\rho ij}^a(p, q)|_{1\text{L}} = -igk^\varepsilon \gamma_\rho (T_f^a)_{ij} \frac{g^2}{16\pi^2} [c_{\text{adj}} + c_2(f)] \frac{1}{\varepsilon} + o(\varepsilon^{-1})$$

The counter term gives a contribution of

$$-iK_1 gk^\varepsilon \gamma_\rho T_f^a$$

Therefore one obtains

$$-igk^\varepsilon \gamma_\rho T_f^a \left[ K_1 + \frac{g^2}{16\pi^2} (c_{\text{adj}} + c_2(f)) \frac{1}{\varepsilon} \right] = \text{finite} \implies Z_1 = 1 - \frac{g^2}{16\pi^2} [c_{\text{adj}} + c_2(f)] \frac{1}{\varepsilon} + \text{finite}$$

**Remark.** The ward identity of QED,  $Z_1 = Z_2$ , does not hold due to  $c_{\text{adj}}$ .

## Lecture 19

### 12.5 Slavnov–Taylor identities

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**Remark.** One may not renormalize the primed coupling constants because, for non-abelian gauge theories, there is only one coupling constant. The gauge-fixing term breaks the gauge invariance of the Lagrangian, but keeps BRST invariance. Gauge invariance implies the Slavnov–Taylor identities, which state that the coupling constants renormalize in the same way. One need not compute  $Z_4$ ,  $Z_5$ ,  $Z_7$  and  $Z_8$ .

**Remark.** The renormalization function  $Z_6$  appears in

$$c_0^a = Z_6^{\frac{1}{2}} c^a$$

To be renormalized, one has to compute the diagram contributing to the ghost propagator. The integral reduces to one already computed.

**Exercise.** Find  $Z_6$  at one-loop.

**Exercise.** Find the explicit form of the Slavnov–Taylor identity for non-abelian gauge theories. See Ramond, eq. 8.7.40.

**Remark.** The renormalization functions  $Z_1$ ,  $Z_2$ ,  $Z_3$  and  $Z_6$  are gauge dependent. The computations differ based on the gauge. The structure of the ghost action depends on the gauge-fixing condition used. In the axial gauge,  $A_3^a = 0$ , ghosts decouple and one does not have a ghost-loop gauge field diagram.

## 12.6 Asymptotic freedom

Consider the one-loop renormalization functions

$$\begin{aligned} Z_2 &= 1 - \frac{g^2}{16\pi^2} c_2(f) \frac{1}{\varepsilon} + \text{finite} \\ Z_3 &= 1 + \frac{g^2}{16\pi^2} \left( \frac{5}{3} c_{\text{adj}} - \frac{4}{3} c(f) \right) \frac{1}{\varepsilon} + \text{finite} \\ Z_1 &= 1 - \frac{g^2}{16\pi^2} [c_{\text{adj}} + c_2(f)] \frac{1}{\varepsilon} + \text{finite} \end{aligned}$$

The bare coupling constant is

$$\begin{aligned} g_0 &= Z_1 Z_2^{-1} Z_3^{-\frac{1}{2}} g k^\varepsilon \\ &= g k^\varepsilon \left[ 1 - \frac{g^2}{16\pi^2} \left( c_{\text{adj}} + \frac{N}{d_f} c(f) \right) \frac{1}{\varepsilon} \right] \left[ 1 + \frac{g^2}{16\pi^2} \frac{N}{d_f} c(f) \frac{1}{\varepsilon} \right] \\ &\quad \times \left[ 1 - \frac{g^2}{16\pi^2} \left( \frac{5}{6} c_{\text{adj}} - \frac{2}{3} c(f) \right) \frac{1}{\varepsilon} \right] + \text{finite} \\ &= k^\varepsilon \left[ g - \frac{g^3}{16\pi^2} \left( \frac{11}{6} c_{\text{adj}} - \frac{2}{3} c(f) \right) \frac{1}{\varepsilon} + \text{finite} \right] \end{aligned}$$

where  $N$  is ? and  $d_f$  is the dimension of the representation. In dimensional regularization, the beta function is obtained from the residue of the simple pole in  $\varepsilon$ . In general

$$g_0 = k^{n\varepsilon} \left[ g + \frac{a_1}{\varepsilon} + \dots \right]$$

from which the beta function is

$$\beta(g) = k \, d_k g = -n [a_1 - g \, d_g a_1]$$

In this case,  $n = 1$  and

$$a_1 = -\frac{g^3}{16\pi^2} \left( \frac{11}{6} c_{\text{adj}} - \frac{2}{3} c(f) \right)$$

The beta function is

$$\beta(g)|_{1L} = 2a_1 = -\frac{g^3}{16\pi^2} \left( \frac{11}{3} c_{\text{adj}} - \frac{4}{3} c(f) \right)$$

**Abelian gauge theory.** For an abelian gauge theory, like QED, one sets  $c_{\text{adj}} = 0$  and obtains

$$\beta(g)|_{1L} = \frac{g^3}{12\pi^2} c(f)$$

For QED, one takes the fundamental representation, so  $c(f) = 1$ . One notices that the beta function is non-negative. Since the beta function is related to the running of the coupling constant

$$k \, d_k g = \beta(g)$$

To higher energies, the coupling constant gets stronger.

**Non-abelian gauge theory.** For non-abelian gauge theories, a positive beta function means that the behaviour of the theory is similar to QED. Though, one may obtain a negative beta function for

$$\frac{11}{3}c_{\text{adj}} > \frac{4}{3}c(f)$$

For higher energies, the coupling weakens. This is the asymptotic freedom. This means that perturbation theory is only applicable in the ultraviolet.

**Example.** Consider  $G = \text{SU}(n)$  with the fundamental and adjoint representations giving

$$c(f) = \frac{1}{2}, \quad c_{\text{adj}} = n$$

Then

$$\beta(g) = -\frac{g^3}{16\pi^2} \left[ \frac{11}{3}n - \frac{2}{3} \right]$$

The bracket is positive for  $n > \frac{2}{11}$ . In general, for more than one flavour of fermions

$$\mathcal{L} \propto \bar{\psi}^A \not{\partial} \psi^A + i m_A \bar{\psi}^A \psi^A + i g A_\mu^a \bar{\psi}^A T^a \psi^A$$

where  $A = 1, \dots, n_f$  is the flavour index, one has a  $n_f$  multiplying  $c(f)$ . One may recall that the gauge fields live in the adjoint representation? of the group, while the fermions live in the fundamental representation [r]

$$\beta(g)|_{1\text{L}} = 2a_1 = -\frac{g^3}{16\pi^2} \left( \frac{11}{3}c_{\text{adj}} - \frac{4}{3}c(f)n_f \right)$$

In order to make the bracket positive, one may play with the choice of the gauge group or the number of flavours  $n_f$ .

For QCD, one has  $G = \text{SU}(3)$  and  $n_f = 6$ . Therefore

$$\frac{11}{3} \cdot 3 - \frac{4}{3} \cdot \frac{1}{2} \cdot 6 = 11 - 4 = 7 > 0$$

The theory exhibits asymptotic freedom. It is sufficient to look at one-loop to determine the presence of asymptotic freedom, since higher loops give smaller contributions [r].

**Absence of a Landau pole.** Integrating the beta function of  $g^2$

$$d_{\ln k} g^2 = 2g d_{\ln k} g = -\frac{g^4}{8\pi^2} \left[ \frac{11}{3}c_{\text{adj}} - \frac{4}{3}c(f)n_f \right]$$

gives a simpler result

$$g^2(k) = g^2(k_0) \left[ 1 + \frac{g^2(k_0)}{8\pi^2} \left( \frac{11}{3}c_{\text{adj}} - \frac{4}{3}c(f)n_f \right) \ln \frac{k}{k_0} \right]^{-1}$$

If the parenthesis positive, then there is no Landau pole due to the presence also of a plus sign. Asymptotic freedom and a Landau pole are mutually exclusive.

For small values of the coupling constant, one obtains

$$g^2(k) = g^2(k_0) - \frac{g^4(k_0)}{8\pi^2} \left( \frac{11}{3}c_{\text{adj}} - \frac{4}{3}c(f) \right) \ln \frac{k}{k_0} + o(\ln k/k_0)$$

The first term is the leading contribution, the second term is the next-to-leading contribution and so on.

The non-perturbative region can be described either with computations on the lattice or with the AdS/CFT correspondence. The computations can be done in supersymmetric theories giving a superconformal boundary theory. Recent work has been done on theories that break both supersymmetry and conformal symmetries.



**Confinement.** At low energies, meaning great distances, one may define free asymptotic states. For a theory with asymptotic freedom, far away particles can only be composite particles. One cannot define free asymptotic states for elementary particles. The free asymptotic states can only be defined for composite particles that are neutral with respect to the gauge group (singlets).

For example in QCD, quarks are in the fundamental representation of SU(3) while anti-quarks are in the anti-fundamental representation. One may obtain a singlet and an octet

$$3 \otimes \bar{3} = 1 \oplus 8 = \text{trivial} + \text{adjoint}$$

The singlet are the mesons. With three quarks, one has

$$3 \otimes 3 \otimes 3 = 1 \oplus \dots$$

The singlet are the baryons.

**Intuitive interpretation of asymptotic freedom.** See Peskin. In QED, the coupling constant gets weaker for greater distances, because the charges get screened by the polarization of the vacuum. A lone charge self-interacts generating virtual particle–anti-particle pairs.

From Maxwell's equations

$$\partial_\mu F^{\mu\nu} = J^\nu$$

one finds

$$E^i = F^{0i} \implies \partial_i E^i = e \delta^{(3)}(\mathbf{x}), \quad E^i \sim \frac{1}{r^2}$$

For non-abelian gauge theories, one has

$$D_\mu F^{\mu\nu a} = J^{\nu a}$$

therefore

$$\partial_i E^{ai} = g \delta^{(3)}(\mathbf{x}) \delta^{a1} + ig c^{abc} A_i^b E^{ci}$$

The extra term matters: it can contribute positively or negatively to the first term. This term acts as a source. [r] One may have a counter-screening: a charge may get stronger. For example,  $a = 1$  gives

$$\partial_i E^{1i} = g \delta^{(3)}(\mathbf{x}) + ig c^{1bc} A_i^b E^{ci}$$

## 13 Spontaneous symmetry breaking

See Cheng. One may prove that gauge theories with spontaneous symmetry breaking are renormalizable.

### 13.1 Review of abelian gauge theories

Consider an abelian gauge theory, in particular a complex scalar theory with U(1) gauge invariance. The matter fields transform as

$$\varphi' = e^{-i\alpha(x)} \varphi, \quad A'_\mu = A_\mu - \frac{1}{g} \partial_\mu \alpha$$

The covariant derivative is

$$D_\mu = \partial_\mu - ig A_\mu$$

The Lagrangian is

$$\mathcal{L} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} + (D_\mu \varphi)^\dagger (D_\mu \varphi) - m^2 \varphi^\dagger \varphi - \lambda (\varphi^\dagger \varphi)^2$$

The potential is

$$V(\varphi) = m^2 \varphi^\dagger \varphi + \lambda (\varphi^\dagger \varphi)^2$$

Its extrema are given by

$$0 = \frac{\partial V}{\partial |\varphi|^2} = -m^2 + 2\lambda|\varphi|^2 \implies |\varphi|^2 = \frac{m^2}{2\lambda} \implies |\varphi| = \frac{1}{\sqrt{2}}\sqrt{\frac{m^2}{\lambda}} = \frac{v}{\sqrt{2}} \neq 0$$

The potential is the Mexican hat. The minima are related by a  $U(1)$  transformation. One chooses a particular minimum by fixing the expectation value

$$\langle 0 | \varphi | 0 \rangle = \frac{v}{2}$$

One may rewrite the field in terms of its real and imaginary components

$$\varphi = \frac{1}{\sqrt{2}}(\varphi_1 + i\varphi_2)$$

so that one chooses

$$\langle 0 | \varphi_1 | 0 \rangle = v_1, \quad \langle 0 | \varphi_2 | 0 \rangle = v_2$$

The vacua are related by an  $SO(2)$  transformation

$$\begin{bmatrix} \varphi'_1 \\ \varphi'_2 \end{bmatrix} = \begin{bmatrix} \cos \alpha & \sin \alpha \\ -\sin \alpha & \cos \alpha \end{bmatrix} \begin{bmatrix} \varphi_1 \\ \varphi_2 \end{bmatrix}$$

The kinetic term of the field becomes

$$(D_\mu \varphi)^\dagger (D_\mu \varphi) = \frac{1}{2}(\partial_\mu \varphi_1 \partial^\mu \varphi_1 + \partial_\mu \varphi_2 \partial^\mu \varphi_2) - ig(\partial_\mu \varphi_1 + i \partial_\mu \varphi_2) A^\mu (\varphi_1 + i\varphi_2) + \frac{g^2}{2} A_\mu A^\mu (\varphi_1^2 + \varphi_2^2)$$

The potential is

$$m^2 \varphi^\dagger \varphi - \lambda (\varphi^\dagger \varphi)^2 = \frac{1}{2} m^2 (\varphi_1^2 + \varphi_2^2) - \frac{\lambda}{4} (\varphi_1^2 + \varphi_2^2)$$

Since it is easier working with a field with zero vacuum expectation value, one parameterizes the fluctuations around the vacuum as

$$\varphi'_1 = \varphi_1 - v \quad \varphi'_2 = \varphi_2$$

Substituting this into the Lagrangian, one finds a massless scalar boson  $\varphi'_2$ . For a global symmetry this is a Goldstone boson, while for a gauge symmetry it is a would-be Goldstone boson. Also, the gauge field becomes massive

$$\frac{1}{2} g^2 v^2 A_\mu A^\mu \implies m = gv$$

Counting the degrees of freedom, one notices that the would-be Goldstone boson should become the third degree of freedom of the gauge field.

The Lagrangian after spontaneous symmetry breaking is still gauge invariant under

$$\begin{aligned} (\varphi'_1 + v) &\rightarrow \cos \alpha (\varphi'_1 + v) + \sin \alpha \varphi'_2 \\ \varphi'_2 &\rightarrow -\sin \alpha (\varphi'_1 + v) + \cos \alpha \varphi'_2 \end{aligned}$$

but gauge invariance is broken in the equations of motion.

## Lecture 20

A better way to counter physical degrees of freedom is parametrize the field in this way

$$\varphi(x) = \frac{1}{\sqrt{2}} \eta(x) e^{i \frac{1}{v} \xi(x)}$$

with  $\eta$  and  $\xi$  real scalar fields. The vacuum expectation values of the fields are

$$\langle 0 | \eta | 0 \rangle = v, \quad \langle 0 | \xi | 0 \rangle = 0$$

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One parametrizes the fluctuations around the vacuum to have

$$\varphi(x) = \frac{1}{\sqrt{2}}[\eta(x) + v]e^{\frac{i}{v}\xi(x)}$$

For small fluctuations, one may expand the above to get

$$\varphi = \frac{1}{\sqrt{2}}(\eta + v) \left[ 1 + \frac{i}{v}\xi + o(\xi) \right] = \frac{1}{\sqrt{2}}[\eta + v + i\xi + o(\eta, \xi)]$$

At the lowest order, this expansion corresponds to the previous one

$$\eta \sim \varphi'_1, \quad \xi \sim \varphi'_2$$

The free Lagrangian with these fields is

$$\mathcal{L}_0 = \frac{1}{2}(\partial_\mu \eta \partial^\mu \eta + \partial_\mu \xi \partial^\mu \xi)$$

while the interaction Lagrangian is

$$\mathcal{L}_{\text{int}} = \frac{1}{2}(\eta^2 + \xi^2) - \frac{\lambda}{4}(\eta^2 + v^2 + 2v\eta + \xi^2)^2 + \dots$$

One may combine quadratic terms in  $\xi$  to see that they give zero: the field  $\xi$  has no mass term. In fact

$$\xi^2 \left[ \frac{1}{2}m^2 - \frac{\lambda}{4}2v^2 \right] = \xi^2 \left[ \frac{1}{2}m^2 - \frac{\lambda}{4}2\frac{m^2}{\lambda} \right] = 0$$

Therefore, the field  $\xi$  is a would-be Goldstone mode. This mode can be removed by performing a gauge transformation that removes the exponential in the parametrization of the field  $\varphi$ :

$$\varphi'(x) = e^{-\frac{i}{v}\xi(x)}\varphi(x) = \frac{1}{\sqrt{2}}[\eta(x) + v]$$

This particular choice of the gauge parameter of the U(1) transformation is called unitary gauge. This gauge transformation has to be balanced by a transformation of the gauge field

$$A'_\mu(x) = A_\mu(x) - \frac{1}{gv} \partial_\mu \xi(x)$$

The Lagrangian in terms of these fields is

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} = -\frac{1}{4}F'_{\mu\nu}F'^{\mu\nu}$$

also

$$\begin{aligned} \mathcal{L}_{\text{Higgs}} &= (D_\mu \varphi)^\dagger (D_\mu \varphi) + m^2 \varphi^\dagger \varphi - \lambda(\varphi^\dagger \varphi)^2 = (D_\mu \varphi)'^\dagger (D_\mu \varphi)' + m^2 \varphi'^\dagger \varphi' - \lambda(\varphi'^\dagger \varphi')^2 \\ &= \frac{1}{2}[\partial_\mu \eta - igA'_\mu(v + \eta)][\partial^\mu \eta + igA'^\mu(v + \eta)] + \frac{1}{2}m^2(v + \eta)^2 - \frac{\lambda}{4}(v + \eta)^2 \end{aligned}$$

The field  $\xi$  disappears and the gauge field gets a mass term. One loses the degree of freedom of  $\xi$ , but gains one degree from the mass term

$$\frac{1}{2}g^2v^2A'_\mu A'^\mu$$

**Renormalizability.** In the unitary gauge, unitarity is manifest, but renormalizability is not manifest. One has to change gauge where the latter becomes manifest.

Consider the kinetic term of the gauge field

$$\mathcal{L}_{\text{kin}} = -\frac{1}{4}(\partial_\mu A_\nu - \partial_\nu A_\mu)^2 + \frac{1}{2}g^2v^2A_\mu A^\mu = \dots = \frac{1}{2}A_\mu[(\Box + M^2)\eta^{\mu\nu} - \partial^\mu \partial^\nu]A_\nu$$

where one has  $M = gv$ . In momentum space, the kinetic term is

$$\frac{1}{2}A_\mu[(-p^2 + M^2)\eta^{\mu\nu} + p^\mu p^\nu]A_\nu = -\frac{1}{2}A_\mu[(p^2 - M^2)\eta^{\mu\nu} - p^\mu p^\nu]A_\nu$$

The kinetic term is therefore

$$K^{\mu\nu} = (p^2 - M^2)\eta^{\mu\nu} - p^\mu p^\nu$$

The propagator is invertible

$$\Delta_{\mu\nu} = \frac{1}{p^2 - M^2 + i\varepsilon} \left[ -\eta_{\mu\nu} + \frac{p_\mu p_\nu}{M^2} \right]$$

**Exercise.** Check that it is the inverse

$$K^{\mu\nu} \Delta_{\mu\nu} = \delta_\rho^\mu$$

One has found that gauge theories present a finite number of topologies of divergent diagrams [r]. While in the ordinary massless case, the propagator is damped in the ultraviolet, while in this case the propagator is constant. By revisiting of the superficial degree of divergence of a diagram, one finds that the degree appears not to be renormalizable: there are an infinite number of topologies of apparently diverging diagrams. The symmetry-broken theory seems to be non-renormalizable.

This cannot be the end. One has only chosen the classical solution of the Lagrangian which is still gauge invariant. One still expects renormalizability to be realized due to non-trivial cancellations of ultraviolet-divergent diagrams. Power counting only gives the worst possible divergence of a diagram, but a diagram can still converge.

### 13.2 Renormalizability

't Hooft proved the renormalizability of Yang-Mills gauge theories with spontaneous symmetry breaking. One has to choose a different group of gauges, called  $R_\xi$  gauges.

Consider the parametrization of the field with real and imaginary components

$$\varphi(x) = \frac{1}{\sqrt{2}}(\varphi_1 + v + i\varphi_2)$$

Recall the kinetic term

$$(D_\mu \varphi)^\dagger (D_\mu \varphi) = \frac{1}{2}(\partial_\mu \varphi_1 \partial^\mu \varphi_1 + \partial_\mu \varphi_2 \partial^\mu \varphi_2) - ig(\partial_\mu \varphi_1 + i \partial_\mu \varphi_2)A^\mu(\varphi_1 + i\varphi_2) + \frac{g^2}{2}A_\mu A^\mu(\varphi_1^2 + \varphi_2^2)$$

One has to substitute

$$\varphi_1 \rightarrow \varphi_1 + v$$

One generates an unclear term

$$-gv \partial_\mu \varphi_2 A^\mu$$

This looks like a two point vertex of a mass term. 't Hooft observes that, since the gauge field is still massless, the Lagrangian has to include the gauge-fixing term in a gauge different from Lorenz's. The original gauge-fixing term

$$-\frac{1}{2\xi}(\partial_\mu A^\mu)^2$$

is replaced by a general gauge-fixing function

$$\mathcal{L}_{\text{gf}} = -\frac{1}{2\xi}(\partial_\mu A^\mu + \xi M \varphi_2)^2 = -\frac{1}{2\xi}(\partial_\mu A^\mu)^2 - M \partial_\mu A^\mu \varphi_2 - \frac{1}{2}\xi M^2 \varphi_2^2$$

The second addendum cancels the unclear term when integrating by parts. One also notices that the field  $\varphi_2$  gets a mass term.

**Propagators.** The free Lagrangian then becomes

$$\mathcal{L}_0 = \frac{1}{2}(\partial_\mu \varphi_1 \partial^\mu \varphi_1 - 2m^2 \varphi_1^2) + \frac{1}{2}(\partial_\mu \varphi_2 \partial^\mu \varphi_2 - \xi M^2 \varphi_2^2) - \frac{1}{4}(\partial_\mu A_\nu - \partial_\nu A_\mu)^2 - \frac{1}{2\xi}(\partial_\mu A^\mu)^2 + \frac{1}{2}M^2 A_\mu A^\mu$$

From this term one reads the propagators. The two scalar propagators

$$\Delta_1(p) = \frac{1}{p^2 - 2m^2 + i\varepsilon}, \quad \Delta_2(p) = \frac{1}{p^2 - \xi M^2 + i\varepsilon}$$

One should be careful about the second expression, because the pole of the propagator is the mass of a particle: one is propagating some particles. The gauge field propagator is

$$\Delta_{\mu\nu}(p) = \frac{1}{p^2 - M^2 + i\varepsilon} \left[ -\eta_{\mu\nu} + (1 - \xi) \frac{p_\mu p_\nu}{p^2 - \xi M^2} \right] \sim \frac{1}{p^2}, \quad p^2 \rightarrow \infty$$

In the ultraviolet, the gauge field propagator is damped and power counting works correctly: spontaneous-symmetry-broken gauge theories are renormalizable.

One must notice that this is just the Higgs sector: one should also add matter fields.

**Exercise.** Check that  $\Delta_{\mu\nu}$  is the inverse of the kinetic term.

**Vertices.** The interaction Lagrangian is given by

$$\mathcal{L}_{\text{int}} = g A_\mu (\partial^\mu \varphi_1) \varphi_2 - g A_\mu (\partial^\mu \varphi_2) \varphi_1 + \frac{1}{2} g^2 A_\mu A^\mu (\varphi_1^2 + \varphi_2^2) - \frac{\lambda}{4} (\varphi_1^2 + \varphi_2^2)^2 + g^2 v A_\mu A^\mu \varphi_1 - \lambda v (\varphi_1^3 + \varphi_1 \varphi_2^2)$$

**Exercise.** Derive this expression from the original Lagrangian.

**Remark.** In the  $R_\xi$  gauges, the field  $\varphi_2$  does not decouple and describes peculiar particles which are also found in the gauge field propagator. One expects that unitarity is restored due to the balance between non-physical particles in  $\varphi_2$  and the non-physical contributions in the gauge field  $A_\mu$ .

Renormalizability is manifest, but not unitarity. There is a propagation of spurious degrees of freedom. Unitarity is restored if the contributions coming from spurious degrees of freedom cancel.

**Remark.** To check that the cancellation happens, one may rewrite the gauge field propagator in a different way

$$\Delta_{\mu\nu}(p) = \frac{1}{p^2 - M^2 + i\varepsilon} \left[ -\eta_{\mu\nu} + (1 - \xi) \frac{p_\mu p_\nu}{p^2 - \xi M^2} \right] = \frac{-\eta_{\mu\nu} + \frac{p_\mu p_\nu}{M^2}}{p^2 - M^2 + i\varepsilon} - \frac{p_\mu p_\nu}{M^2} [p^2 - \xi M^2 + i\varepsilon]^{-1}$$

One sees a division from the typical massive propagator and the unphysical part.

**Remark.** In the unitary gauge, the spurious scalar field disappears and the gauge field contains only the physical massive propagator. One may notice that the unitary gauge is recovered from the  $R_\xi$  gauge in the limit  $\xi \rightarrow \infty$ .

**Example.** Consider the tree-level contributions to the scattering process

$$\varphi_1(k_1) + \varphi_1(k_2) \rightarrow A(k_3) + A(k_4)$$

Looking at the vertices of the theory, computing the Feynman amplitudes, one finds that the gauge field propagator cancels the spurious contribution from the  $\varphi_2$  field propagator.

**Exercise.** Check.

**Remark.** The whole discussion also applies for non-abelian gauge theories.

**Remark.** This implies that the Standard Model is a renormalizable theory.

## 14 Anomalies

See Ramond.

**Intuition.** Consider a classical theory with a global symmetry. Let  $\phi$  be the set of fields. The symmetry formally corresponds to the fact that the classical action is invariant under the global transformation

$$S[\phi] = S[\phi']$$

For a quantum theory, one considers the effective action

$$e^{-\Gamma[\phi]} = \int [\mathcal{D}\phi] e^{-S[\phi]}$$

Under the symmetry transformation, the classical action is invariant, but the path integral measure may not be. If the measure is not invariant, then the symmetry is not preserved in the quantum theory and it exhibits an anomaly.

If the measure is invariant, the effective action  $\Gamma[\phi]$  is invariant: the symmetry is preserved in the quantum theory. If the measure is not invariant, then the symmetry is broken in the quantum theory: the symmetry is anomalous.

The breaking of the symmetry is due to the regularization procedure of the theory. For example, massless QED is a scale-invariant field theory. In the regularization, one introduces an energy cutoff and the theory is no longer scale-invariant. This is the reason why the coupling constant runs. At the fixed points of the RG flow, one recovers scale-invariance by also seeing that the stress-energy tensor is traceless since

$$T^\mu_\mu \sim \beta(g)$$

At the fixed points, the parameters no longer run.

**Example.** A few examples of anomalous symmetries are the following.

- For space-time [r] symmetries, there are scale symmetry and conformal symmetry. Poincaré symmetry is always preserved (apart from putting a sharp cutoff).
- For internal symmetries, in four space-time dimensions there is only chiral symmetry acting on fermions.

## Lecture 21

### 14.1 Classical chiral symmetry

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Consider Euclidean formalism. Consider the Dirac Lagrangian

$$\mathcal{L} = \bar{\psi}(\not{D} + im)\psi = \bar{\psi}(\not{\partial} + im)\psi + i\bar{\psi}\not{A}\psi, \quad D_\mu = \partial_\mu + iqA_\mu$$

**Phase transformation.** The gauge U(1) transformations of the fields are

$$\psi' = e^{i\alpha(x)}\psi, \quad A'_\mu = A_\mu - \partial_\mu\alpha$$

The global symmetry corresponds to a conserved Noether current

$$J^\mu \propto \bar{\psi}\gamma^\mu\psi$$

The conserved current can be found in a simple way. Consider the free action

$$S = \int d^4x \bar{\psi}(\not{\partial} + im)\psi$$

The variation of the action under the gauge transformation is

$$\begin{aligned} \delta S &= \int d^4x [\delta\bar{\psi}(\not{\partial} + im)\psi + \bar{\psi}(\not{\partial} + im)\delta\psi] \\ &= \int d^4x i\partial_\mu\alpha(x) \bar{\psi}\gamma^\mu\psi = - \int d^4x \alpha(x) \partial_\mu(\bar{\psi}\gamma^\mu\psi) \end{aligned}$$

The symmetry of the action under any  $\alpha(x)$  implies

$$\delta S = 0 \implies \partial_\mu(\bar{\psi}\gamma^\mu\psi) = 0 \implies \partial_\mu J^\mu = 0$$

**Chiral transformation.** Consider chiral transformations of the fields

$$\psi' = e^{i\beta\gamma_5}\psi, \quad \bar{\psi}' = \bar{\psi}e^{i\beta\gamma_5}$$

Consider the kinetic term transforms as

$$\begin{aligned} \bar{\psi}' \not{\partial} \psi' &= \bar{\psi} e^{i\beta\gamma_5} \not{\partial} (e^{i\beta\gamma_5} \psi) = \bar{\psi} \not{\partial} \psi + \bar{\psi} e^{i\beta\gamma_5} \partial_\mu (e^{-i\beta\gamma_5}) \gamma^\mu \psi \\ &= \bar{\psi} \not{\partial} \psi - i \partial_\mu \beta(x) \bar{\psi} \gamma_5 \gamma^\mu \psi = \bar{\psi} \not{\partial} \psi + i \partial_\mu \beta(x) \bar{\psi} \gamma^\mu \gamma_5 \psi \end{aligned}$$

The mass term transforms as

$$\bar{\psi}' \psi' = \bar{\psi} e^{i\beta\gamma_5} e^{i\beta\gamma_5} \psi = \bar{\psi} e^{2i\beta\gamma_5} \psi = \bar{\psi} \psi + 2i\beta \bar{\psi} \gamma_5 \psi + o(|\beta|)$$

Therefore, the variation of the action is

$$\begin{aligned} \delta S &= \int d^4x [i \partial_\mu \beta(x) \bar{\psi} \gamma^\mu \gamma_5 \psi - 2m\beta(x) \bar{\psi} \gamma_5 \psi] \\ &= - \int d^4x \beta(x) [\partial_\mu (i \bar{\psi} \gamma^\mu \gamma_5 \psi) + 2m \bar{\psi} \gamma_5 \psi] = 0 \end{aligned}$$

This implies that

$$\partial_\mu (i \bar{\psi} \gamma^\mu \gamma_5 \psi) = -2m \bar{\psi} \gamma_5 \psi$$

One defines the chiral current to be

$$J_\mu^5 = i \bar{\psi} \gamma_\mu \gamma_5 \psi$$

For massive fields, the chiral current is not conserved

$$\partial^\mu J_\mu^5 = -2m \bar{\psi} \gamma_5 \psi$$

and chiral transformations are not a symmetry of the chiral theory. For massless fields, the theory possesses chiral symmetry

$$\partial^\mu J_\mu^5 = 0$$

## 14.2 Quantum chiral symmetry

One would like to study whether chiral symmetry is kept in the quantum theory. One finds that the symmetry is not kept. In massless QED, the ABJ anomaly appears.

One may study the anomaly in two ways. The first is perturbative: at one-loop, one finds that the current is not conserved. This result is non-perturbative because anomalies typically do not get higher order corrections, like in this case. The second way is studying how the path integral measure is not invariant under chiral transformations.

Consider the action

$$S = \int d^4x \left[ \bar{\psi} (\not{\partial} + iq\not{A}) \psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} \right]$$

The gauge field term contains the global U(1) current

$$\bar{\psi} iq\not{A} \psi = A_\mu J^\mu$$

If one defines a functional integral where one integrates only the spinor fields, then the gauge field is a background field and acts a source for the current. [r] One defines

$$e^{-Z[A_\mu, S_\mu^5]} = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] e^{-S[\psi, \bar{\psi}, A_\mu, S_\mu^5]}$$

where one has

$$S[\psi, \bar{\psi}, A_\mu, S_\mu^5] = \int d^4x [\bar{\psi} \not{\partial} \psi + A_\mu i \bar{\psi} \gamma^\mu \psi + S_\mu^5 i \bar{\psi} \gamma^\mu \gamma_5 \psi]$$

The last two terms are currents and they are composite operators. This is the generating functional of such composite operators.

One would like to see whether

$$\partial^\mu J_\mu^5 = 0$$

in the quantum theory. One computes the three-point correlation function

$$\langle J_\mu(x) J_\nu(y) J_\rho^5(z) \rangle$$

If the chiral symmetry is preserved, one should find

$$\frac{\partial}{\partial z^\rho} \langle J_\mu(x) J_\nu(y) J_\rho^5(z) \rangle = \langle J_\mu(x) J_\nu(y) \partial_\rho J_\rho^5(z) \rangle = 0$$

**Remark.** Since the chiral current  $J^5$  is even under charge conjugation, the only non-zero  $n$ -point correlation functions are the ones with an even number of currents  $J$ .

**Remark.** Integrating in four dimensions, one shall find integrals with  $\gamma_5$ . Typically, one brings the Dirac matrices outside the integral and computes a pure momentum integral analytically continuing the integral in  $D = 4 - 2\varepsilon$  dimensions. The fifth gamma matrix is not defined outside four dimensions. One write the matrix as product of the other matrices [r].

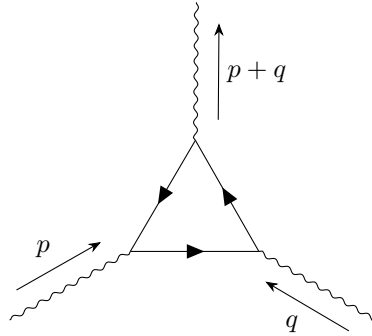
One may combine dimensional regularization with dimensional reduction. One brings the external momenta outside the integral, one analytically continues only the integration variable. When computing the gamma structure, the Dirac matrices related to the external momenta are in four dimensions, while the ones coming from the integration momentum is in more dimensions. In this way one also keeps  $\gamma_5$  in four dimensions. One has to first compute the integral then apply the identities of the gamma matrices.

#### 14.2.1 One-loop contributions

The three-point correlation function is given by

$$\langle J_\mu(x) J_\nu(y) J_\rho^5(z) \rangle = - \frac{\delta^3 e^{-Z}}{\delta A_\mu(x) \delta A_\nu(y) \delta S_\rho^5(z)} \Big|_{A=S^5=0}$$

One needs to compute the diagram [r] diagr



Checking the anomaly in momentum space, means

$$(p+q)_\rho T_{\mu\nu\rho}(p, q, p+q) \stackrel{?}{=} 0$$

**Feynman rules.** The fermion propagators is

$$-\frac{i}{\not{p}}$$

The vector and axial vertices are

$$i q k^\varepsilon, \quad i k^\varepsilon$$



**Triangle diagram.** The integral associated to the diagram is

$$\begin{aligned} T_{\mu\nu\rho}(p, q) &= -iq^2 k^{3\varepsilon} \int \frac{d^4 l}{(2\pi)^4} \text{Tr} \left[ \gamma_\rho \gamma_5 \frac{-i}{\not{l} - \not{p}} \gamma_\nu \frac{-i}{\not{l}} \gamma_\mu \frac{-i}{\not{l} + \not{p}} \right] \\ &= q^2 k^{3\varepsilon} \int \frac{d^4 l}{(2\pi)^4} \frac{\text{Tr}[\gamma_5 \gamma_\rho (\not{l} - \not{q}) \gamma_\nu \not{l} \gamma_\mu (\not{l} + \not{p})]}{(l - q)^2 l^2 (l + p)^2} \end{aligned}$$

At the second line, one has applied

$$-\frac{i}{\not{l}} = i \frac{\not{l}}{l^2}$$

It is easier to compute the difference

$$T_{\mu\nu\rho}(p, q) - T_{\nu\mu\rho}(q, p)$$

If one sends  $l \rightarrow -l$  in the second term, one finds

$$\begin{aligned} T &= T_{\mu\nu\rho}(p, q) - T_{\nu\mu\rho}(q, p) = e^2 k^{3\varepsilon} \int \frac{d^4 l}{(2\pi)^4} \frac{1}{(l - q)^2 l^2 (l + q)^2} \left\{ \text{Tr}[\gamma_5 \gamma_\rho (\not{l} - \not{q}) \gamma_\nu \not{l} \gamma_\mu (\not{l} + \not{p})] \right. \\ &\quad \left. - \text{Tr}[\gamma_5 \gamma_\rho (\not{l} + \not{p}) \gamma_\mu \not{l} \gamma_\nu (\not{l} - \not{q})] \right\} \\ &= q^2 k^{3\varepsilon} [\text{Tr}(\gamma_5 \gamma_\rho \gamma_\alpha \gamma_\nu \gamma_\beta \gamma_\mu \gamma_\gamma) - \text{Tr}(\gamma_5 \gamma_\rho \gamma_\gamma \gamma_\mu \gamma_\beta \gamma_\nu \gamma_\alpha)] \int \frac{d^4 l}{(2\pi)^4} \frac{(l - q)_\alpha l_\beta (l + p)_\gamma}{(l - q)^2 l^2 (l + p)^2} \\ &= \text{Traces} \times 2 \int_0^1 dx \int_0^{1-x} dy \int \frac{d^d l}{(2\pi)^d} \left[ \frac{l_\alpha l_\beta l_\gamma}{(\text{Den})^3} + \frac{p_\gamma l_\alpha l_\beta - q_\alpha l_\beta l_\gamma}{(\text{Den})^3} - q_\alpha p_\gamma \frac{l_\beta}{(\text{Den})^3} \right] \end{aligned}$$

At the third line, one has applied Feynman combining with

$$\text{Den} = x(l - q)^2 + y(l + p)^2 + (1 - x - y)l^2 = l^2 + 2lP + M^2$$

where

$$P_\mu = yp_\mu - xq_\mu, \quad M^2 = xq^2 + yp^2$$

and integrals

$$\int_0^1 dx \int_0^{1-x} dy$$

At the fourth line, one has applied B.17, B.18, B.19 of Ramond.

Consider the first addendum

$$I_1 = 2 \int_0^1 dx \int_0^{1-x} dy \frac{(-1)}{(4\pi)^{2-\varepsilon} \Gamma(3)} \left[ P_\alpha P_\beta P_\gamma \frac{\Gamma(1+\varepsilon)}{D^{1+\varepsilon}} + \frac{1}{2} (\delta_{\alpha\gamma} P_\beta + \delta_{\beta\gamma} P_\alpha + \delta_{\alpha\beta} P_\gamma) \frac{\Gamma(\varepsilon)}{D^\varepsilon} \right]$$

where

$$D \equiv M^2 - P^2$$

One has to compute the integral with the traces. Consider the first addendum

$$\begin{aligned} T_{1a} &= [\text{Tr}(\gamma_5 \gamma_\rho \gamma_\alpha \gamma_\nu \gamma_\beta \gamma_\mu \gamma_\gamma) - \text{Tr}(\gamma_5 \gamma_\rho \gamma_\gamma \gamma_\mu \gamma_\beta \gamma_\nu \gamma_\alpha)] P_\alpha P_\beta P_\gamma \\ &= \text{Tr}(\gamma_5 \gamma_\rho \not{P} \gamma_\nu \not{P} \gamma_\mu \not{P}) - (\mu \leftrightarrow \nu) \\ &= P^2 \text{Tr}(\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu \not{P}) - 2 \text{Tr}(\gamma_5 \gamma_\rho \not{P} \gamma_\mu \not{P}) P_\nu - (\mu \leftrightarrow \nu) \\ &= P^2 \text{Tr}(\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu \not{P}) + 0 - (\mu \leftrightarrow \nu) \\ &= 2P^2 \text{Tr}(\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu \not{P}) \end{aligned}$$

At the third line and beyond, one has applied the following gamma technology

$$\begin{aligned} \not{P} \gamma_\mu \not{P} &= P^2 \gamma_\mu - 2P_\mu \not{P} \\ \text{Tr}[\gamma_5 \gamma_\mu \gamma_\nu] &= 0, \quad \text{in 4D} \\ \gamma_\alpha \gamma_\mu \gamma^\alpha &= 2(1 - \varepsilon) \gamma_\mu \end{aligned}$$

The second addendum is

$$\begin{aligned} T_{1b} &= [\text{Tr}(\gamma_5 \gamma_\rho \gamma_\alpha \gamma_\nu \gamma_\beta \gamma_\mu \gamma_\gamma) - \text{Tr}(\gamma_5 \gamma_\rho \gamma_\gamma \gamma_\mu \gamma_\beta \gamma_\nu \gamma_\alpha)] (\delta_{\alpha\gamma} P_\beta + \delta_{\beta\gamma} P_\alpha + \delta_{\alpha\beta} P_\gamma) \\ &= [8(1 - \varepsilon) - 4(1 - \varepsilon) + 8] \text{Tr}(\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu \not{P}) \end{aligned}$$

Therefore, the total contribution is

$$T_1 = T_{1a} + T_{1b} = 4(3 - \varepsilon) \text{Tr}(\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu \not{P})$$

The integral is

$$I_1 = -2 \int_0^1 dx \int_0^{1-x} dy \frac{\Gamma(1 + \varepsilon)}{(4\pi)^{2-\varepsilon} D^{1+\varepsilon}} \left[ P^2 + \frac{3-\varepsilon}{\varepsilon} D \right] \text{Tr}(\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu \not{P})$$

[r] not same  $I_1$  as above, as the above is missing the traces.

The second integral is

$$\begin{aligned} I_2 &= \int_0^1 dx \int_0^{1-x} dy \frac{\Gamma(1 + \varepsilon)}{(4\pi)^{2-\varepsilon} D^{1+\varepsilon}} \left[ 2P^2 + 2 \frac{1-\varepsilon}{\varepsilon} D \right] \\ &\quad \times \{ \text{Tr}[\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu (\not{p} - \not{q})] - 2P_\nu \text{Tr}[\gamma_5 \gamma_\rho \not{P} \gamma_\mu (\not{p} - \not{q})] + 2P_\mu \text{Tr}[\gamma_5 \gamma_\rho \not{P} \gamma_\nu (\not{p} - \not{q})] \} \end{aligned}$$

The third integral is

$$I_3 = - \int_0^1 dx \int_0^{1-x} dy \frac{\Gamma(1 + \varepsilon)}{(4\pi)^{2-\varepsilon} D^{1+\varepsilon}} [\text{Tr}(\gamma_5 \gamma_\rho \not{p} \gamma_\mu \not{P} \gamma_\nu \not{q}) - \text{Tr}(\gamma_5 \gamma_\rho \not{q} \gamma_\mu \not{P} \gamma_\nu \not{p})]$$

## Lecture 22

The first two integrals contain diverging contributions. The divergent terms are proportional to  $M^2 - P^2$  which simplifies one factor at the denominator: the expansion in  $\varepsilon$  gives a constant term.

The divergent term from the first integral becomes

$$J_1 = -2 \int_0^1 dx \int_0^{1-x} dy \frac{\Gamma(1 + \varepsilon)}{(4\pi)^{2-\varepsilon} D^\varepsilon} \frac{1}{2} \frac{2-\varepsilon}{\varepsilon} \text{Tr}[\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu (\not{p} - \not{q})]$$

The divergent term from the second integral becomes

$$J_2 = 2 \int_0^1 dx \int_0^{1-x} dy \frac{\Gamma(1 + \varepsilon)}{(2\pi)^{2+\varepsilon} D^\varepsilon} \frac{1-\varepsilon}{\varepsilon} \text{Tr}[\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu (\not{p} - \not{q})]$$

[r] The sum of the two is

$$J_1 + J_2 = \left[ -2 - \frac{1}{2} \frac{2-\varepsilon}{\varepsilon} + 2 \frac{1-\varepsilon}{\varepsilon} \right] \int \dots = - \int \dots$$

Therefore, the triangle diagram gives

$$\begin{aligned} T_{\mu\nu\rho}(p, q) &= -q^2 \int_0^1 dx \int_0^{1-x} dy \frac{\Gamma(1 + \varepsilon)}{(2\pi)^{2-\varepsilon} D^\varepsilon} \text{Tr}[\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu (\not{p} - \not{q})] + \text{finite} \\ &= -\frac{q^2}{16\pi^2} \int_0^1 dx \int_0^{1-x} dy \text{Tr}[\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu (\not{p} - \not{q})] + \text{finite} \\ &= -\frac{q^2}{32\pi^2} \text{Tr}[\gamma_5 \gamma_\rho \gamma_\nu \gamma_\mu (\not{p} - \not{q})] + \text{finite} \end{aligned}$$

At the second line, one may take the limit  $\varepsilon \rightarrow 0$  and perform the integrals.

One may then find

$$(p + q)_\rho T_{\mu\nu\rho}(p, q) = -\frac{q^2}{32\pi^2} \text{Tr}[\gamma_5 (\not{p} + \not{q}) \gamma_\nu \gamma_\mu (\not{p} - \not{q})]$$

The finite terms do not contribute when contracting with  $p + q$ .

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**Trace.** One may evaluate the trace to see if it is zero

$$\begin{aligned}
 T &= \text{Tr}[\gamma_5(\not{p} + \not{q})\gamma_\nu\gamma_\mu(\not{p} - \not{q})] = (p+q)_\rho(p-q)_\sigma \text{Tr}(\gamma_5\gamma_\rho\gamma_\nu\gamma_\mu\gamma_\sigma) \\
 &= (p_\rho p_\sigma - q_\rho q_\sigma - p_\rho q_\sigma + q_\rho p_\sigma) \text{Tr}(\gamma_5\gamma_\rho\gamma_\nu\gamma_\mu\gamma_\sigma) \\
 &= (p_\rho p_\sigma - q_\rho q_\sigma - p_\rho q_\sigma + q_\rho p_\sigma) \text{Tr}(\gamma_5\gamma_\sigma\gamma_\rho\gamma_\nu\gamma_\mu) \\
 &= -(p_\rho p_\sigma - q_\rho q_\sigma + q_\rho p_\sigma - p_\rho q_\sigma) \text{Tr}(\gamma_5\gamma_\sigma\gamma_\rho\gamma_\nu\gamma_\mu) \\
 &= (p_\rho q_\sigma - q_\rho p_\sigma) \text{Tr}(\gamma_5\gamma_\sigma\gamma_\rho\gamma_\nu\gamma_\mu) \\
 &= p_\rho q_\sigma \text{Tr}(\gamma_5[\gamma_\sigma, \gamma_\rho]\gamma_\nu\gamma_\mu) \\
 &= p_\rho q_\sigma \text{Tr}(\gamma_5\{\gamma_\sigma, \gamma_\rho\}\gamma_\nu\gamma_\mu) - 2p_\rho q_\sigma \text{Tr}(\gamma_5\gamma_\rho\gamma_\sigma\gamma_\nu\gamma_\mu) \\
 &= -2p_\rho q_\sigma \text{Tr}(\gamma_5\gamma_\rho\gamma_\sigma\gamma_\nu\gamma_\mu) \\
 &= -2p_\rho q_\sigma(-4\varepsilon_{\rho\sigma\nu\mu}) = 8\varepsilon_{\nu\mu\rho\sigma}p_\rho q_\sigma = -8\varepsilon_{\mu\nu\rho\sigma}p_\rho q_\sigma \neq 0
 \end{aligned}$$

At the fifth line, one has noticed that the expression is symmetric in  $\rho\sigma$  and one may introduce the anti-commutator of gamma matrices in the trace. At the sixth line, one has noticed that the expression is anti-symmetric in  $\rho\sigma$ , so one introduces a commutator inside the trace.

In the end, one finds the ABJ anomaly

$$(p+q)_\rho T_{\mu\nu\rho}(p, q) = \frac{q^2}{4\pi^2} \varepsilon_{\mu\nu\rho\sigma} p_\rho q_\sigma$$

#### 14.2.2 Non-perturbative contribution

See Ramond, Srednicki. One may wonder if the anomaly gets contribution from higher orders. One may compute the anomaly using Fujikawa's non-perturbative method.

The [r] is

$$e^{-Z[A]} = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] e^{-S[\psi, \bar{\psi}, A]}$$

where the gauge field is a background field which acts as a source for the electromagnetic current. The action is

$$S[\psi, \bar{\psi}, A] = \bar{\psi} \not{D} \psi = \bar{\psi} \not{\partial} \psi + iq A_\mu \bar{\psi} \gamma^\mu \psi$$

A local chiral transformation of the spinor fields is

$$\psi' = e^{i\beta(x)\gamma_5} \psi \sim \psi + i\beta(x)\gamma_5 \psi, \quad \bar{\psi}' = \bar{\psi} e^{i\beta(x)\gamma_5} \sim \bar{\psi} + \bar{\psi} i\beta(x)\gamma_5$$

For an infinitesimal transformation, the action becomes

$$S'[\psi', \bar{\psi}', A] = S[\psi, \bar{\psi}, A] - \int d^4x \beta(x) \partial_\mu (i\bar{\psi} \gamma_\mu \gamma_5 \psi)$$

One would like to understand whether the consistency of the procedure requires that

$$\partial_\mu J_\mu^5 = 0$$

One has to impose the chiral transformation also on the measure. In the functional integral, the chiral transformation is a change of variables. Therefore, one may write

$$\int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] e^{-S[\psi, \bar{\psi}, A]} = \int [\mathcal{D}\psi' \mathcal{D}\bar{\psi}'] (\det J)^2 e^{-S'[\psi', \bar{\psi}', A]}$$

If one substitutes the action  $S'$ , then one obtains the term on the left and a term proportional to  $\partial_\mu J_\mu^5$ . If the Jacobian is non-trivial [r].

**Jacobian.** The spinor fields are Grassmann-odd fields and the Jacobian of the change of variables is the reciprocal of typical integration

$$[\mathcal{D}\psi' \mathcal{D}\bar{\psi}'] = (\det J)^{-2} [\mathcal{D}\psi \mathcal{D}\bar{\psi}]$$

Therefore, one obtains

$$e^{-Z[A]} = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] e^{-S[\psi, \bar{\psi}, A]} = \int [\mathcal{D}\psi \mathcal{D}\bar{\psi}] (\det J)^{-2} e^{-S[\psi, \bar{\psi}, A]} \exp \left[ - \int d^4x \beta(x) \partial_\mu J_\mu^5 \right]$$

[r] One may find the Jacobian

$$\psi'(x) = e^{i\beta(x)\gamma_5} \psi(x) = e^{i\beta(x)\gamma_5 \delta^{(4)}(x-y)} \psi(y) = J(x, y) \psi(y)$$

[r] ??

Therefore

$$\begin{aligned} (\det J)^{-2} &= e^{-2 \ln \det J} = e^{-2 \text{Tr} \ln J} \\ &= \exp \left[ -2 \text{Tr} \left( i\beta(x) \gamma_5 \delta^{(4)}(x-y) \right) \right] \\ &= \exp \left[ -2i \text{Tr} \int d^4x \beta(x) \delta^{(4)}(x-x) \gamma_5 \right] \end{aligned}$$

[r] ??

The trace is done on both the spinor discrete indices and the continuous position indices. One has to regularize the delta function

$$\delta^{(4)}(x-y) = \int \frac{d^4k}{(2\pi)^4} e^{ik(x-y)} \rightarrow \int \frac{d^4k}{(2\pi)^4} e^{-\frac{k^2}{M^2}} e^{ik(x-y)}$$

where  $M$  is a regularization parameter that is removed in the limit  $M \rightarrow \infty$ . The momentum in the exponent can be replaced by the derivative in the position

$$\delta^{(4)}(x-y) \rightarrow \int \frac{d^4k}{(2\pi)^4} e^{M^{-2} \partial_x^2} e^{ik(x-y)} = e^{M^{-2} \partial_x^2} \delta^{(4)}(x-y)$$

One would like not to break gauge covariant even during regularization: one has to modify the derivative. Therefore

$$\partial_x \rightarrow i\mathcal{D}_x$$

where the imaginary unit is needed to keep the self-adjoint nature of the operator. Thus

$$\delta^{(4)}(x-y) \rightarrow e^{M^{-2} (i\mathcal{D}_x)^2} \delta^{(4)}(x-y) = \int \frac{d^4k}{(2\pi)^4} e^{-M^{-2} (\mathcal{D}_x)^2} e^{ik(x-y)}$$

Consider the following general identity

$$f(\partial_x) e^{ikx} = e^{ikx} f(\partial_x + ik)$$

which holds also for the covariant derivative. One may rewrite

$$\delta^{(4)}(x-y) \rightarrow \int \frac{d^4k}{(2\pi)^4} e^{ik(x-y)} e^{-M^{-2} (\mathcal{D} + i\mathbf{k})^2}$$

The exponent is

$$\begin{aligned} (\mathcal{D} + i\mathbf{k})^2 &= \mathcal{D}\mathcal{D} - \mathbf{k}\mathbf{k} + i\mathcal{D}\mathbf{k} + i\mathbf{k}\mathcal{D} \\ &= \gamma_\mu \gamma_\nu D_\mu D_\nu - \gamma_\mu \gamma_\nu k_\mu k_\nu + i\gamma_\mu \gamma_\nu (D_\mu k_\nu + k_\mu D_\nu) \\ &= \frac{1}{2} (\{\gamma_\mu, \gamma_\nu\} + [\gamma_\mu, \gamma_\nu]) D_\mu D_\nu - \frac{1}{2} \{\gamma_\mu, \gamma_\nu\} k_\mu k_\nu + i\{\gamma_\mu, \gamma_\nu\} D_\mu k_\nu \\ &= \frac{1}{2} (\{\gamma_\mu, \gamma_\nu\} + [\gamma_\mu, \gamma_\nu]) D_\mu D_\nu + k^2 + i\{\gamma_\mu, \gamma_\nu\} D_\mu k_\nu \\ &= -D^2 - 2i\frac{1}{4} [\gamma_\mu, \gamma_\nu] D_\mu D_\nu + k^2 - 2iD_\mu k_\mu \\ &= -D^2 + k^2 - 2iD_\mu k_\mu - 2i\gamma_{\mu\nu} D_\mu D_\nu \\ &= -D^2 + k^2 - 2iD_\mu k_\mu - i\gamma_{\mu\nu} [D_\mu, D_\nu] \\ &= -D^2 + k^2 - 2iD_\mu k_\mu - i\gamma_{\mu\nu} igF_{\mu\nu} \\ &= -D^2 + k^2 - 2iD_\mu k_\mu + \gamma_{\mu\nu} gF_{\mu\nu} \end{aligned}$$

The regularized delta function is

$$\begin{aligned}\delta^{(4)}(x-y)|_{\text{reg}} &= \int \frac{d^4 k}{(2\pi)^4} e^{ik(x-y)} \exp \left[ -\frac{1}{M^2} (-D^2 + k^2 - 2iD_\mu k_\mu + \gamma_{\mu\nu} g F_{\mu\nu}) \right] \\ &= M^4 \int \frac{d^4 k'}{(2\pi)^4} e^{iMk'(x-y)} e^{-k'^2} \exp \left[ \frac{D^2}{M^2} + \frac{2i}{M} k'_\mu D_\mu - \frac{g}{M^2} \gamma_{\mu\nu} F_{\mu\nu} \right]\end{aligned}$$

At the second line, one changes variables

$$k_\mu = Mk'_\mu$$

One may compute the trace with the regularized Delta function

$$\begin{aligned}\text{Tr} \left[ \delta_{\text{reg}}^{(4)}(x-x) \gamma_5 \right] &= M^4 \int \frac{d^4 k}{(2\pi)^4} e^{-k^2} \text{Tr} \left[ \exp \left( \frac{D^2}{M^2} + \frac{2i}{M} k_\mu D_\mu - \frac{g}{M^2} \gamma_{\mu\nu} F_{\mu\nu} \right) \gamma_5 \right] \\ &= M^4 \int \frac{d^4 k}{(2\pi)^4} e^{-k^2} \text{Tr} \left[ \gamma_5 + \left( \frac{D^2}{M^2} + \frac{2i}{M} k_\mu D_\mu - \frac{g}{M^2} \gamma_{\mu\nu} F_{\mu\nu} \right) \gamma_5 + \dots \right] \\ &= \int \frac{d^4 k}{(2\pi)^4} e^{-k^2} \frac{1}{2} q^2 \text{Tr} [\gamma_{\mu\nu} F_{\mu\nu} \gamma_{\rho\sigma} F_{\rho\sigma} \gamma_5] \\ &= \frac{q^2}{2} \left[ \frac{i}{4} \right]^2 \int \frac{d^4 k}{(2\pi)^4} e^{-k^2} F_{\mu\nu} F_{\rho\sigma} \text{Tr} ([\gamma_\mu, \gamma_\nu] [\gamma_\rho, \gamma_\sigma] \gamma_5) \\ &= 4q^2 \left[ \frac{i}{4} \right]^2 \int \frac{d^4 k}{(2\pi)^4} e^{-k^2} F_{\mu\nu} F_{\rho\sigma} \text{Tr} (\gamma_\mu \gamma_\nu \gamma_\rho \gamma_\sigma \gamma_5) \\ &= \dots \\ &= -\frac{q^2}{2^5} \frac{4\pi^2}{(2\pi)^4} F_{\mu\nu} F_{\rho\sigma} (-4) \varepsilon_{\mu\nu\rho\sigma} \\ &= \frac{q^2}{32\pi^2} \varepsilon_{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma}\end{aligned}$$

At the second line, one Taylor expands the exponential in the limit  $M \rightarrow \infty$ . At the third line, one has applied

$$\text{Tr} \gamma_5 = 0, \quad \text{Tr}(\gamma_5 \gamma_\alpha \gamma_\beta) = 0$$

[r] and noted that only one term survives which comes for the quadratic term in the expansion.

The Jacobian is thus

$$\begin{aligned}(\det J)^{-2} &= \exp \left[ -2i \text{Tr} \int d^4 x \beta(x) \delta^{(4)}(x-x) \gamma_5 \right] \\ &= \exp \left[ -2i \frac{q^2}{32\pi^2} \int d^4 x \beta(x) \varepsilon_{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma} \right]\end{aligned}$$

This exponential has to balance the one of the exponential of the chiral current. Therefore

$$2i \frac{q^2}{32\pi^2} \int d^4 x \beta(x) \varepsilon_{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma} + \int d^4 x \beta(x) \partial_\mu J_\mu^5 = 0$$

From which the chiral anomaly

$$\partial_\mu J_\mu^5 = i \frac{q^2}{16\pi^2} \varepsilon_{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma} = i \frac{q^2}{8\pi^2} F_{\mu\nu} \tilde{F}_{\rho\sigma}$$

[r] sign

where  $\tilde{F}$  is the dual tensor. The generalization to  $N$  fermions is

$$\partial_\mu J_\mu^5 = i \frac{q^2}{8\pi^2} N F_{\mu\nu} \tilde{F}_{\rho\sigma}$$

**Result from triangle diagram.** Consider the anomaly written differently

$$\partial_\rho J_\rho^5 = i \frac{q^2}{16\pi^2} \varepsilon_{\mu\nu\eta\sigma} F_{\mu\nu} F_{\eta\sigma} = \frac{iq^2}{4\pi^2} \varepsilon_{\mu\nu\eta\sigma} \partial_\mu A_\nu \partial_\eta A_\sigma$$

Going to momentum space and setting the momenta like in the triangle diagram

$$\partial_\rho \rightarrow i(p+q)_\rho, \quad \partial_\mu \rightarrow -i\rho_\mu, \quad \partial_\eta \rightarrow -iq_\eta$$

Therefore

$$\begin{aligned} i(p+q)_\rho T_{\nu\sigma\rho}(p, q) &= \frac{iq^2}{4\pi^2} (-i)^2 p_\mu q_\eta \varepsilon_{\mu\nu\eta\sigma} \\ (p+q)_\rho T_{\nu\sigma\rho}(p, q) &= \frac{q^2}{4\pi^2} \varepsilon_{\nu\sigma\mu\eta} p_\mu q_\eta \end{aligned}$$

This is the same result as the one-loop computation. The chiral anomaly is one-loop exact.