

Quantum Field Theory

November 30, 2021

1 Preliminaries

1.1 Natural units

$$[c] = LT^{-1} \quad [\hbar] = L^2 MT^{-1} \quad [G] = L^3 M^{-1} T^{-2}$$

We set $c = \hbar = 1$ to give the natural units. All quantities expressed in natural units scale with some power of mass or energy. We use notation below to describe this:

$$XM^\delta, [X] = \delta$$

where δ is the scaling dimension. e.g. $[E] = +1$ and $[L] = -1$. In order to convert back from natural units we need to know what we have calculated. For example, if it is energy then we multiply by c^2 as $E = mc^2$ but in natural units $c = 1$. Basically everything is done in units of energy/mass with c and \hbar giving the scale of choice between length, time and energy.

1.2 Classical Fields

A classical scalar field is one that maps from Minkowski spacetime to the "field space":

$$\phi : \mathbb{R}^{3,1} \rightarrow \mathbb{R}$$

The scalar part of the definition implies that the field is invariant under Lorentz transformations that are defined by keeping the metric ($\eta = +1, -1, -1, -1$) invariant:

$$\Lambda^\mu_\sigma \Lambda^\nu_\tau \eta^{\sigma\tau} = \eta^{\mu\nu}$$

the further restriction of $\Lambda = +1$ removes the possibility of reflection and defines the Lorentz group $G_L = SO(3, 1)$.

The active convention for describing a scalar field is:

$$\phi(x) \rightarrow^\Lambda \phi'(x)$$

where $\phi'(x) = \phi(\Lambda^{-1} \cdot \vec{x})$. It is clear that $\Lambda^{-1} \in SO(3, 1)$ and $(\Lambda^{-1})^\mu_l \Lambda^\nu_\nu = \delta^\mu_\nu$.

1.3 Spacetime derivatives

$$\partial_\mu \phi(x) = \frac{\partial \phi(x)}{\partial x^\mu}$$

$$\partial_\mu \phi(x) \rightarrow^\Lambda \partial_\mu \phi'(x) = \frac{\partial \phi(\Lambda^{-1} \cdot x)}{\partial x^\mu} = (\Lambda^{-1})^\nu_\mu \partial_\nu \phi((\Lambda^{-1} \cdot x))$$

4-Vector Field

$$V : \mathbb{R}^{3,1} \rightarrow \mathbb{R}^{3,1}$$

$$v^\mu(x) \rightarrow^\Lambda \Lambda^\mu_\nu v^\nu(\Lambda^{-1} \cdot x)$$

1.4 Lagrangians

Action

$$S[q] = \int_{t_i}^{t_f} dt L(q(t), \dot{q}(t))$$

Principle of least action: If you vary the path with fixed endpoints then the stationary action ($\partial S = 0$) is equivalent to the Euler-Lagrange equation

1.4.1 Lagrangian for scalar field theory

Demand Lorentz invariant action, locality and at most two time derivatives. Lagrangian is a functional as it must give a number for every field.

$$L(t) = L[\phi, \partial_\mu \phi] = \int d^3x \mathfrak{L}(\phi(x), \partial_{\mu\nu} \phi(x))$$

where \mathfrak{L} is the Lagrangian density. This ensures locality as the field is EOM will only be influenced by the local field.

$$S_{t_i, t_f}[\phi, \partial_\mu \phi] = \int_{t_i}^{t_f} L[\phi, \partial_\mu \phi] dt = \int_{t_i}^{t_f} dt \int d^3x \mathfrak{L}$$

Take infinite time interval (nice interval to take when assuming that before and after interaction particles were very far apart)

$$S = \int_{\mathbb{R}^{3,1}} d^4x \mathfrak{L}(\phi(x), \partial_\mu \phi(x))$$

Consider L.T. $\Lambda \in SO(3, 1)$ and require that \mathfrak{L} is a scalar field:

$$\mathfrak{L}(x) \rightarrow^\Lambda \mathfrak{L}(\Lambda^{-1} \cdot x)$$

then

$$S = \int_{\mathbb{R}^{3,1}} d^4x \mathfrak{L}(x) \rightarrow^\Lambda \int_{\mathbb{R}^{3,1}} d^4x \mathfrak{L}(\Lambda^{-1} \cdot x) = S'$$

choose variables of $y^\mu = \Lambda^{-1\mu}_\nu x^\nu$ giving:

$$S' = \int_{\mathbb{R}^{3,1}} d^4y \mathfrak{L}(y)$$

so the action is Lorentz invariant if the Lagrangian density is a scalar field.

1.4.2 General Lagrangian

$$L = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) \quad (1)$$

general terms you can have if you want at most two time derivatives. Not quite most general as could have an arbitrary function of ϕ as a pre-factor to the kinetic term but is ruled out later by a dimensionality constraint. No need to separately consider $\phi \partial_\mu^\mu \phi$ but once it is inside the action this is the kinetic term by integration by parts (differs only by a surface term).

Everytime two indicies are contracted top and bottom they use the metric:

$$\eta_{\mu\nu} \Lambda_\sigma^\mu = \Lambda_{\mu\sigma}$$

$$\partial_\mu \phi \partial^\mu \phi = \eta_{\mu\nu} \partial^\mu \phi \partial^\nu \phi$$

2 Lecture 3

Principle of Least Action Vary the field configuration: $\phi(x) \rightarrow \phi(x) + \delta\phi(x)$ but fix the boundary condition: $\delta\phi(x) = \phi(t, x) \rightarrow 0$ for $|\mathbf{x}| \rightarrow \infty$ and $t \rightarrow \pm\infty$. Consider variation of action:

$$\begin{aligned} \delta S &= \int d^4x \left[\frac{\partial \mathcal{L}}{\partial \phi} \cdot \delta\phi + \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \delta(\partial_\mu \phi) \right] \\ &= \int d^4x \left[\left(\frac{\partial \mathcal{L}}{\partial \phi} - \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial \partial_\mu \phi} \right) \right) \delta\phi + \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \delta\phi \right) \right] \\ B &= \int_{\mathbb{R}^{3,1}} d^4x \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial \partial_\mu \phi} \delta\phi \right) = \int_{\partial(\mathbb{R}^{3,1})} dS_\mu \frac{\partial \mathcal{L}}{\partial \partial_\mu \phi} \delta\phi \end{aligned}$$

The boundary $\partial\mathbb{R}^{3,1}$ represents where $|\mathbf{x}| \rightarrow \infty$ and $t \rightarrow \pm\infty$. Therefore, $\delta\phi = 0$ and $\partial(\mathbb{R}^{3,1})$. This means we can set this integrand to zero. The above needs modifying to match notes but it is essentially the same derivation we use for Euler-Lagrange in classical mechanics.

2.0.1 Euler-Lagrange equation

$$\frac{\partial \mathcal{L}}{\partial \phi} - \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} \right) = 0 \quad (2)$$

$$\frac{\partial \mathcal{L}}{\partial \phi} = -V'(\phi)$$

$$\frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi)} = \partial^\mu \phi$$

gives general field equation of motion:

$$\partial_\mu \partial^\mu \phi + V'(\phi) = 0 \quad (3)$$

$$\partial_\mu \partial^\mu \phi = \frac{\partial^2}{\partial t^2} - \nabla_{\mathbf{x}}^2$$

This is in general a second order PDE which is hard to solve normally need computer. We want to consider special case that is easy to solve e.g. choose a quadratic potential to give a linear EOM:

$$V(\phi) = \frac{1}{2}m^2\phi^2$$

gives the Klein-Gordon equation:

$$\partial_\mu \partial^\mu \phi + m^2\phi = 0 \quad (4)$$

This is indeed a lorentz invariant equation, which is good as the whole point of the action stuff was to generate EOM that were lorentz invariant. This is a wave equation, and we can immediately say quite a lot about it. It has wave-like solutions, e.g. trial solution:

$$\phi = e^{i\mathbf{x}\cdot\mathbf{p}} = e^{i\omega t - i\mathbf{k}\cdot\mathbf{x}}$$

gives dispersion relation:

$$\omega_k = \sqrt{|k| + m^2}$$

You can see that when we go to the quantum scale this will correspond to free massive particles as, as this is basically the wave-particle dual for a standard particle of mass m . This gives a free particle as we have a linear classical equation so it has a superposition principle so you can add together solutions so you will get non-interacting particles.

The general case is the case when we have a non-linear PDE given by a non-quadratic potential. You won't have localised wave packets they will disperse. If you try to form lumps of fields then they will interact with each other and disperse. No superposition principle. may have exotic solutions like solitons. In notes from last year there is also a short section of applying this to maxwells equation of electromagnetism if his doesn't do this after covering maxwells stuff go read the notes from last year.

2.1 Symmetry

Definition: Variation of the field that leaves the action invariant

Typically has the structure of the group and if it is a continuous variation then it behaves like a manifold (this is a Lie Group). Why are they important in field theory - because they tell us about the conserved quantities in the theory. As every symmetry implies a conservation law by Noether's Theorem. Important in any theory as we will bung some particles in and have some coming out at the but the key constraint on what is allowed and what is not is the conservation laws in the theory e.g. tells us we cant have an electron go in and a positron come out as that would violate a conservation law. More generally than that, they have much richer consequences e.g. non-abelian symmetries then infact the imprint which that leaves on the observables of that theory that is much more than conservation laws. It leaves restrictions on the allowed masses and spectrums etc. Here we are generally talking about global symmetry which is the things we can actually compute (and we will use the field as a proxy for this) change under the application of the symmetry whereas gauge symmetry the observables actually remain unchanged.

We will implemment symmetries as a variation of the field not the coordinates, below are listed some symmetries and their action on coordinates and on fields.

Translation: $\mathbf{x}^\mu \rightarrow \mathbf{x}^\mu + \mathbf{c}^\mu$ for a $\mathbf{c} \in \mathbb{R}^{3,1}$

$$\phi(x) \rightarrow \phi'(x) = \phi(x - c)$$

Automatic that all the actions that are written down are invariant under translation as we integrated over all space so they are translation invariant.

Lorentz: $\mathbf{x}^\mu \rightarrow \Lambda_\nu^\mu \mathbf{x}^\nu$

$$\phi(x) \rightarrow \phi(\Lambda^{-1} \cdot x)$$

If massless ($m = 0$) and $\nu = 0$ then also get a **scale transformation:** $\mathbf{x}^{mu} \rightarrow \lambda \mathbf{x}^\mu$ for $\lambda \in \mathbb{R}^+$

$$\phi(x) \rightarrow \lambda^{-\Delta} \phi(\lambda^{-1} x), \Delta = [\phi]$$

2.2 Internal Symmetry

Does not act on the coordinates, and its generators commute with all the generators of the global symmetry above. These give properties of particles

such as electric charge, flavour and colour.

Example: Complex scalar field

$$\psi : \mathbb{R}^{3,1} \rightarrow \mathbb{C}$$

Need real lagrangian so need to contract with complex conjugate. **When finding the Euler-Lagrange equation of a complex scalar field you treat ϕ and ϕ^* as independent variables so you get two E-L equations.** Below is one possible complex field with a specific form of potential that generates an internal symmetry.

$$\mathcal{L} = \partial_\mu \psi^* \partial^\mu \psi - V(|\psi|^2)$$

This has an internal symmetry as:

$$\psi(x) \rightarrow \psi'(x) = e^{i\alpha} \psi(x)$$

$$\psi^*(x) \rightarrow \psi'^*(x) = e^{-i\alpha} \psi^*(x)$$

which gives both \mathcal{L} and S are invariant. It doesn't need to be invariant in the lagrangian density to be a symmetry but this stronger case often occurs in internal symmetries unlike in the case of Lorentz transformation where \mathcal{L} was not invariant.

Continuous symmetry form matrix Lie groups e.g. Lorentz group

$$G_L = \{\Lambda \in Mat_4(\mathbb{R}), \Lambda \cdot \eta \cdot \Lambda^T = \eta, \det \Lambda = 1\} = SO(3, 1)$$

One can ask about how the different space-time symmetries fit together. The lorentz transformations correspond to $SO(3, 1)$ and when combined with translations (which don't commute with lorentz transformations) you get the Poincare group. Finally, if we specialise to the case where the mass is zero then we get space invariance and this further enhances the group to give the conformal group ($SO(4,2)$).

3 Lecture 4

3.1 Finite vs Infinitesimal transformations

You don't have to understand the full non-linearity of the Lie Group everything can be thought of by just considering the behaviour of the group

elements near the identity. Group elements sufficiently near the identity can be written in the following form:

$$g = \exp(\alpha X), \alpha \in \mathbb{R}, X \in \mathbb{L}(G)$$

For all the examples in the course we are only considering matrix lie groups so all the elements are matrices. $\mathbb{L}(G)$ is called the Lie Algebra of the group. Therefore, $\exp(\alpha X) = \sum_{n=0}^{\infty} \frac{1}{n!} (\alpha X)^n$. We are going to consider this for $|\alpha| \ll 1$ to give $\exp(\alpha X) = I + \alpha X + O(\alpha^2)$, which is called an infinitesimal transformation.

3.1.1 Example

$$\alpha = \partial_\mu \psi^* \partial^\mu \psi - V(|\psi|^2)$$

can undergo finite transformation:

$$\psi \rightarrow g\psi, g = \exp(i\alpha)$$

so has infinitesimal transformation:

$$\psi \rightarrow \psi + \delta\psi, \delta\psi = i\alpha\psi$$

3.1.2 Lorentz Transformations

$$x^\mu \rightarrow x'^\mu = \Lambda^\mu_\nu x^\nu$$

$$\Lambda^\mu_\sigma \Lambda^\nu_\tau \eta^{\sigma\tau} = \eta^{\mu\nu}$$

$$\Lambda = +1$$

Can express lambda as:

$$\Lambda = \exp(s\Omega), \Omega \in \mathbb{L}(SO(3, 1))$$

$$\Lambda = \delta^\mu_\nu + s\Omega^\mu_\nu + O(s^2)$$

Expand in linear order:

$$(\delta^\mu_\alpha + s\Omega^\mu_\alpha) \eta^{\alpha\beta} (\delta^\nu_\beta + s\Omega^\nu_\beta) = \eta^{\mu\nu}$$

Consider $O(s)$:

$$s\Omega^{\mu\nu} + s\Omega^{\nu\mu} = 0$$

so:

$$\Omega_{\mu\nu} = -\Omega_{\nu\mu}$$

As this means the diagonal must be zero and it is symmetric about the diagonal, Ω only has six degrees of freedom. Now linearise the transformation of the scalar field:

$$\begin{aligned}\phi(x) &\rightarrow \phi(\Lambda^{-1} \cdot x) \\ (\Lambda^{-1})^\mu_\nu &= \delta^\mu_\nu - s\Omega^\mu_\nu + O(s^2)\end{aligned}$$

so to $O(s)$

$$\phi(\Lambda^{-1} \cdot x) = \phi(x - s\Omega \cdot x + O(s^2)) = \phi(x) - s\Omega^\mu_{nu} x^{nu} \partial_{mu} \phi(x) + O(s^2)$$

so the infinitesimal transformation is:

$$\begin{aligned}\phi(x) &\rightarrow \phi(x) + \delta\phi(x) \\ \delta\phi(x) &= -s\Omega^\mu_\nu x^\nu \partial_\mu \phi(x)\end{aligned}\tag{5}$$

3.2 Noether's Theorem

A continuous symmetry implies a conserved current.

$$\phi(x) \rightarrow \phi + \delta\phi(x)$$

As the variation is local the transformation will just be in terms of the field and its derivatives at the point:

$$\delta\phi(x) = X(\phi(x), \partial\phi(x))$$

How does \mathfrak{L} vary?

Lorentz transformation

$$\phi(x) \rightarrow \phi(\Lambda^{-1} \cdot x)$$

gives

$$\delta\mathfrak{L} = -s\Omega^\mu_\nu x^\nu \partial_\mu \mathfrak{L}(x) = \partial_\mu (-s\Omega^\mu_\nu x^\nu \mathfrak{L}(x))$$

above we used the fact that Ω is anti-symmetric to vanish the term $-s\Omega^\mu_\nu \partial_\mu (x^\nu) \mathfrak{L}(x)$. So we can rewrite the variation of the lagrangian under the lorentz transformation as a total derivative.

Translation

$$\begin{aligned}\phi(x) &\rightarrow \phi(x - c) \\ \delta\phi(x) &= X = -c^\mu \partial_\mu \phi(x) \\ \delta\mathfrak{L} &= -c^\mu \partial_\mu \mathfrak{L}(x)\end{aligned}$$

If the lagrangian transforms as a total derivative then the action will be invariant.

General symmetry

$$\begin{aligned}\delta\mathfrak{L} &= \partial_\mu F^\mu \\ F^\mu &= F^\mu(\phi(x), \partial\phi(x))\end{aligned}\tag{6}$$

As:

$$S = \int_{\mathbb{R}^{3,1}} d^4x \delta\mathfrak{L} = \int_{\mathbb{R}^{3,1}} d^4x \partial_\mu F^\mu = \int_{\partial(\mathbb{R}^{3,1})} dS_\mu F^\mu$$

So if ϕ , $\partial\phi$ decay fast enough then this goes to 0.

3.2.1 Conserved current

$$j^\mu(x) = j^\mu(\phi(x), \partial\phi(x))$$

is conserved if

$$\partial_\mu j^\mu = 0$$

when ϕ obeys its E-L equations.

$$j^\mu(x) = j^\mu(t, x) = (j^0(t, x), \mathbf{J}(t, x))$$

j^0 is charge density and \mathbf{J} is the current density.

Total charge in region $V \subset \mathbb{R}^3 = Q(t) = \int_V dV j^0(t, x)$.

Rate of change of $Q(t)$ with conserved current:

$$\frac{dQ(t)}{dt} = \int_V dV \frac{\partial}{\partial t} j^0(t, x) = - \int_V dV \nabla \cdot \mathbf{J} = - \int_{\partial V} d\mathbf{S} \cdot \mathbf{J}$$

This means that whenever we have a conserved current we have a conserved charge, and any changes in charge are accounted for by the flux of current across the boundary.

Proof of Noether's theorem

$$\partial\mathfrak{L}(x) = \frac{\partial\mathfrak{L}}{\partial\phi}\delta\phi + \frac{\partial\mathfrak{L}}{\partial(\partial_\mu\phi)}\delta(\partial_\mu\phi)$$

$$\partial\mathfrak{L}(x) = \left[\frac{\partial\mathfrak{L}}{\partial\phi} - \partial_\mu\left(\frac{\partial\mathfrak{L}}{\partial(\partial_\mu\phi)}\right)\right]\delta\phi + \partial_\mu\left(\frac{\partial\mathfrak{L}}{\partial(\partial_\mu\phi)}\right)\delta\phi$$

if ϕ obeys EL then

$$\partial\mathfrak{L}(x) = \partial_\mu\left(\frac{\partial\mathfrak{L}}{\partial(\partial_\mu\phi)}X(\phi, \partial\phi)\right)\tag{7}$$

As for any symmetry we have

$$\delta \mathfrak{L} = \partial_\mu F^\mu$$

so define conserved current to be:

$$j^\mu = \frac{\partial \mathfrak{L}}{\partial(\partial_\mu \phi_a)} X(\phi_a, \partial \phi_a) - F^\mu$$

As both of these are equal to the variation of the lagrangian it will be conserved.

4 Lecture 5

Focus on spacetime translations:

$$x^\mu \rightarrow x^\mu - \epsilon^\mu$$

$$\phi(x) \rightarrow \phi(x + \epsilon) \approx \phi(x) + \epsilon^\mu \partial_\mu \phi(x) + O(\epsilon^2)$$

$$\delta \phi(x) = +\epsilon^\mu \partial_\mu \phi(x)$$

$$\delta \mathfrak{L}(x) = +\epsilon^\mu \partial_\mu \mathfrak{L}(x)$$

$$\delta \phi(x) = \epsilon^\nu X_\nu(\phi)$$

$$\delta \mathfrak{L}(x) = \epsilon^\nu \partial_\mu F_\nu^\mu(\phi)$$

$$X_\nu = \partial_\nu \phi, F_\nu^\mu = \delta_\nu^\mu \mathfrak{L}$$

So this can give us 4 conserved currents(**energy-momentum tensor**):

$$T_\nu^\mu = \frac{\partial \mathfrak{L}}{\partial(\partial_\mu)} \partial_\nu \phi - \delta_\nu^\mu \mathfrak{L} \quad (8)$$

$$\partial_\mu T_\nu^\mu = 0$$

Conserved charge that corresponds to energy comes from the zeroth component:

$$E = \int_{\mathbb{R}^3} d^3 x T^{00}$$

similarly for translations in space:

$$P^i = \int_{\mathbb{R}^3} d^3 x T^{0i}$$

4.1 Example: free scalar field/Klein Gordon

$$\mathfrak{L} = \frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi) - \frac{m^2}{2}\phi^2$$
$$T^{\mu\nu} = \partial^\mu \phi \partial^\nu \phi - \eta^{\mu\nu} \mathfrak{L}$$

Look at conserved energy when $\mu = \nu = 0$.

$$E = \int_{\mathbb{R}^3} d^3 T^{00} = \int_{\mathbb{R}^3} d^3 \left(\frac{1}{2} \dot{\phi}^2 + \frac{1}{2} |\nabla \phi|^2 + \frac{1}{2} m^2 \phi^2 \right)$$

This formula has several nice features as the energy is bounded below which is something we often want as we often have ground states. It also follows the same pattern as the Lagrangian in classical dynamics if the first term is interpreted as being the kinetic energy and the remaining as potential energy (it goes from $L = T - V$ to $E = T + V$).

The energy momentum tensor is clearly a symmetric tensor, now that is not necessarily a general feature but in a system when you have a rotational symmetry you can always didn't really make much sense have a look in the notes. Basically he was trying to say that energy doesn't really make an awful lot of sense until you couple it to gravity and then the energy momentum tensor becomes very important.

4.2 Canonical Quantisation

4.2.1 Quick review of classical hamiltonian mechanics

Consider non-relativistic particle ($m = 1$) in potential $V(q)$:

$$L(q(t), \dot{q}(t)) = \frac{1}{2} \dot{q}^2 - V(q)$$

Define momentum p conjugate to q :

$$p = \frac{\partial L}{\partial \dot{q}} = \dot{q}$$

Define Legendre transform as

$$H = p\dot{q} - L(q, \dot{q})$$

with \dot{q} replaced with p to give:

$$H = \frac{1}{2}p^2 + V(q)$$

Now consider system with N degrees of freedom:

$$L = L(\{q_i(t)\}, \{\dot{q}_i(t)\})$$

$$H = \sum_i p_i \dot{q}_i - L$$

and eliminate \dot{q} .

Poisson Bracket: For $F = F(p, q)$ and $G = G(p, q)$, $\{F, G\} = \sum_i \frac{\partial F}{\partial p_i} \frac{\partial G}{\partial q_i} - \frac{\partial F}{\partial q_i} \frac{\partial G}{\partial p_i}$

$$\{q_i, q_j\} = \{p_i, p_j\} = 0, \{q_i, p_j\} = \delta_{ij}$$

State of system at time $t = 0$ is given by specifying your momentum and position ($q_i(0)$ and $p_i(0)$). So time evolution is completely governed by the hamiltonian. For any function $F(p, q)$,

$$\dot{F} = \{F, H\}$$

In particular, choose $F = p_i$ or q_i gives hamiltons equations.

$$\dot{q}_i = \{q_i, H\} = \frac{\partial F}{\partial p_i}$$

$$\dot{p}_i = \{p_i, H\} = -\frac{\partial F}{\partial q_i}$$

$$\dot{Q} = 0 \iff \{H, Q\} = 0$$

5 Lecture 6

5.0.1 Quantization

States are vectors in a Hilbert space. Quantization means turning classical functions into a hermitian operator. The basic rule as first written down by Dirac is:

$$\{, \} \rightarrow \frac{1}{i\hbar}[,]$$

replace poisson bracket with the commutator. e.g. can apply this recipe to the coordinates and their momentum

$$q_i \rightarrow \hat{q}_i, p_i \rightarrow \hat{p}_i$$

and

$$[\hat{q}_i, \hat{q}_j] = [\hat{p}_i, \hat{p}_j] = 0, [\hat{q}_i, \hat{p}_j] = i\hbar\delta_{ij}$$

In general the problem of quantisation is finding a representation for this algebra which in this case is the $\frac{\partial}{\partial x}$ representation of p .

The time evolution of the quantised system is described by the transformed Hamiltonian $\hat{H}(\hat{p}, \hat{q})$. This often leads to ambiguity as the ordering of the \hat{p} and \hat{q} is important in \hat{H} (most of this we will be able to brush under the carpet as we will mainly deal with free states). If we deal with the Schrödinger picture then the states are thought of as time dependent:

$$-i\hbar \frac{\partial}{\partial t} |\psi(t)\rangle = \hat{H} |\psi(t)\rangle$$

The special case is when $|\psi(t)\rangle$ is an eigenstate of the Hamiltonian where

$$\hat{H} |\psi\rangle = E |\psi\rangle$$

Field Theory is an infinite dimensional dynamical system. We think about this as replacing the dynamical variable $q_i(t)$ with $\phi(\mathbf{x}, t)$. So in the dynamical system we have a discrete label i whereas the field has a continuous label $\mathbf{x} \in \mathbb{R}^3$. It turns out it is not valid to treat these the same but for the purposes of this course we can pretend it is possible. Though in AQFT you will find out you need to take an appropriate limit using something called a regulator. There are two types of infinite in this system: the IR (infrared) divergences associated with the infinite range of x ('large distances'), the UV (ultraviolet) divergences are associated with the fact we have a continuous infinite that we can consider infinitely close together points ('short distances').

As we have a dynamical variable the next step is defining a conjugate momentum:

$$\pi(\mathbf{x}, t) = \frac{\partial \mathcal{L}(\mathbf{x}, t)}{\partial \dot{\phi}(\mathbf{x}, t)}$$

Perform a Legendre transform to get a Hamiltonian density

$$\mathcal{H} = \pi(\mathbf{x}, t) \dot{\phi}(\mathbf{x}, t) - \mathcal{L}(\mathbf{x}, t)$$

Just like in classical dynamics we can now eliminate the time derivative from the hamiltonian density using the definition of π to give $\mathcal{H} = \mathcal{H}(\phi(\mathbf{x}, t), \pi(\mathbf{x}, t))$. Define the hamiltonian as:

$$H = \int_{\mathbb{R}^3} d^3x \mathcal{H}(\mathbf{x}, t)$$

Example

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi - V(\phi))$$

One problem with the hamiltonian treatment is it breaks lorentz invariance so we need to split this up into space and time components.

$$\mathcal{L} = \frac{1}{2}\dot{\phi}^2 - \frac{1}{2}|\nabla\phi|^2 - V(\phi)$$

Conjugate momentum

$$\pi(\mathbf{x}, t) = \frac{\partial \mathcal{L}}{\partial \dot{\phi}} = \dot{\phi}(\mathbf{x}, t)$$

$$\mathcal{H} = \pi\dot{\phi} - \mathcal{L} = \frac{1}{2}\pi^2 + \frac{1}{2}|\nabla\phi|^2 + V(\phi) = T^{00}$$

$$H = \int_{\mathbb{R}^3} d^3x \frac{1}{2}\pi^2 + \frac{1}{2}|\nabla\phi|^2 + V(\phi) = \int_{\mathbb{R}^3} d^3x T^{00} = E$$

Now lets quantise it.

$$\{\phi(\mathbf{x}, t), \pi(\mathbf{y}, t)\} = \delta^{(3)}(\mathbf{x} - \mathbf{y})$$

Suppress/fix the time coordinate for the moment.

Canonical Quantisation

$$\phi(\mathbf{x}) \rightarrow \hat{\phi}(\mathbf{x})$$

$$\pi(\mathbf{x}) \rightarrow \hat{\pi}(\mathbf{x})$$

$$[\hat{\phi}(\mathbf{x}, t), \hat{\pi}(\mathbf{y}, t)] = i\hbar\delta^{(3)}(\mathbf{x} - \mathbf{y})$$

$$H = \int_{\mathbb{R}^3} d^3x \frac{1}{2}\hat{\pi}^2 + \frac{1}{2}|\nabla\hat{\phi}|^2 + V(\hat{\phi})$$

except in the case of free field theory we aren't going to do this as describing the hilbert space is very hard so we will do it for free field theory and perturb around from it. In the above note that we normally work in natural units but we have left the \hbar here to show the parallel with the classical case. The hamiltonian is very hard to understand here except in the case of free field theory as we are taking derivatives of operators and stuff all over the place.

Can solve the free field theory because of the superposition principle allowing us to effectively consider it as an infinite number of simple harmonic oscillators.

5.1 Free Field Theory

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

Equation of motion:

$$\partial_\mu \partial^\mu \phi + m^2 \phi = 0$$

As this is linear we can find general solution by solving with Fourier transform.

$$\phi(\mathbf{x}, t) = \int \frac{d^3 p}{(2\pi)^3} e^{i\mathbf{p} \cdot \mathbf{x}} \tilde{\phi}(\mathbf{p}, t)$$

If we want $\phi(\mathbf{x}, t) \in \mathbb{R}$ then we need $\tilde{\phi}^*(\mathbf{p}, t) = \tilde{\phi}(-\mathbf{p}, t)$. The equation becomes:

$$\left(\frac{\partial^2}{\partial t^2} + |\mathbf{p}|^2 + m^2\right)\tilde{\phi}(\mathbf{p}, t) = 0$$

This is essentially the equation for simple harmonic motion so solutions are:

$$\tilde{\phi}(\mathbf{p}, t) = A_{\mathbf{p}} e^{i\omega_{\mathbf{p}} t} + B_{\mathbf{p}} e^{-i\omega_{\mathbf{p}} t}$$

for $\omega_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$. In reality $A_{\mathbf{p}}^* = A_{-\mathbf{p}}$ to give a real field. Also the action becomes:

$$S = \int dt \int d^3 x \mathcal{L}(\mathbf{x}, t) = \frac{1}{2} \int dt \int d^3 p \tilde{\phi}^*(\mathbf{p}, t) \left(-\frac{\partial^2}{\partial t^2} - |\mathbf{p}|^2 - m^2\right) \tilde{\phi}(\mathbf{p}, t)$$

This is what gives the sense of an infinite set of a decoupled simple harmonic oscillators, as for any fixed value of \mathbf{p} this gives the action of the harmonic oscillator.

6 Lecture 7

6.1 Quantum Simple harmonic oscillator

$$L = \frac{1}{2}\dot{q}^2 - \frac{\omega^2}{2}q^2$$

$$H = \frac{1}{2}\dot{p}^2 + \frac{\omega^2}{2}q^2$$

First off replace $q\hat{q}$ and $p \rightarrow \hat{p}$ with $[\hat{q}, \hat{p}] = i\hbar$. Define creation and annihilation operators:

$$\hat{a} = \sqrt{\frac{\omega}{2}}\hat{q} + \frac{i}{\sqrt{2\omega}}\hat{p}$$

$$\begin{aligned}\hat{a}^\dagger &= \sqrt{\frac{\omega}{2}}\hat{q} - \frac{i}{\sqrt{2\omega}}\hat{p} \\ [\hat{a}, \hat{a}^\dagger] &= \hbar\end{aligned}\tag{9}$$

Rule of ordering always have annihilation operators on the right

$$\begin{aligned}\hat{H} &= \frac{\omega}{2}(\hat{a}^\dagger\hat{a} + \frac{\hbar}{2}) \\ [\hat{H}, \hat{a}] &= -\omega\hbar\hat{a} \\ [\hat{H}, \hat{a}^{dagger}] &= \omega\hbar\hat{a}^{dagger}\end{aligned}$$

Consider:

$$\hat{H}|E\rangle = E|E\rangle$$

Commutation relations tell us

$$\begin{aligned}\hat{H}(\hat{a}^{dagger}|E\rangle) &= (E + \hbar\omega)\hat{a}^\dagger|E\rangle \\ \hat{H}(\hat{a}|E\rangle) &= (E - \hbar\omega)\hat{a}|E\rangle\end{aligned}$$

Either spectrum unbounded below or have $\hat{a}|0\rangle = 0$. Therefore:

$$\hat{H}|0\rangle = \frac{\hbar\omega}{2}|0\rangle$$

and the raising operator gives the infinite tower of excited states

$$\begin{aligned}|n\rangle &= (\hat{a}^\dagger)^n|0\rangle \\ \hat{H}|n\rangle &= (n + \frac{1}{2})\hbar\omega|n\rangle\end{aligned}$$

Define number operator: $\hat{N} = \hat{a}^\dagger\hat{a}$

$$\hat{N}|n\rangle = n|n\rangle, \hat{H} = \omega\hat{N} + \frac{1}{2}\hbar\omega$$

6.2 Quantisation

$$[\hat{\phi}(\mathbf{x}), \hat{\pi}(\mathbf{y})] = i\delta^{(3)}(\mathbf{x} - \mathbf{y})$$

Dirac delta function

$$\int_{(R)^3} f(\mathbf{x})\delta^{(3)}(\mathbf{x} - \mathbf{y}) = f(\mathbf{y})$$

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

$$\hat{H} = \int_{\mathbb{R}^3} d^3x \left(\frac{\hat{\pi}^2}{2} + \frac{|\nabla \hat{\phi}|^2}{2} + \frac{m^2}{2} \hat{\phi}^2 \right)$$

Write field in fourier space. This is a definition that we then check produces the correct commutator relations.

$$\hat{\phi}(\mathbf{x}) = \int \frac{d^3p}{(2\pi)^3} [\hat{a}_{\mathbf{p}} e^{i\mathbf{p}\cdot\mathbf{x}} + \hat{a}_{\mathbf{p}}^\dagger e^{-i\mathbf{p}\cdot\mathbf{x}}] \frac{1}{\sqrt{2\omega_{\mathbf{p}}}}$$

for $\omega_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$.

$$\hat{\pi}(\mathbf{x}) = \int \frac{d^3p}{(2\pi)^3} (-i) \sqrt{\frac{\omega_{\mathbf{p}}}{2}} [\hat{a}_{\mathbf{p}} e^{i\mathbf{p}\cdot\mathbf{x}} - \hat{a}_{\mathbf{p}}^\dagger e^{-i\mathbf{p}\cdot\mathbf{x}}]$$

Make the claim that

$$[\hat{\phi}(\mathbf{x}), \hat{\phi}(\mathbf{y})] = [\hat{\pi}(\mathbf{x}), \hat{\pi}(\mathbf{y})] = 0, [\hat{\phi}(\mathbf{x}), \hat{\pi}(\mathbf{y})] = i\delta^{(3)}(\mathbf{x} - \mathbf{y})$$

is equivalent to:

$$[\hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{q}}] = [\hat{a}_{\mathbf{q}}^\dagger, \hat{a}_{\mathbf{p}}^\dagger] = 0$$

$$[\hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{q}}^\dagger] = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q})$$

Check that:

$$[\hat{\phi}(\mathbf{x}), \hat{\pi}(\mathbf{y})] = \frac{-i}{2} \int \frac{d^3p}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} \sqrt{\frac{\omega_{\mathbf{q}}}{\omega_{\mathbf{p}}}} [\hat{a}_{\mathbf{p}} e^{i\mathbf{p}\cdot\mathbf{x}} + \hat{a}_{\mathbf{p}}^\dagger e^{-i\mathbf{p}\cdot\mathbf{x}}, \hat{a}_{\mathbf{q}} e^{i\mathbf{q}\cdot\mathbf{y}} + \hat{a}_{\mathbf{q}}^\dagger e^{-i\mathbf{q}\cdot\mathbf{y}}]$$

$$[,] = -e^{i\mathbf{p}\cdot\mathbf{x} - i\mathbf{q}\cdot\mathbf{y}} [\hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{q}}^\dagger] + e^{-i\mathbf{p}\cdot\mathbf{x} + i\mathbf{q}\cdot\mathbf{y}} [\hat{a}_{\mathbf{p}}^\dagger, \hat{a}_{\mathbf{q}}] = -(2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q}) (e^{i\mathbf{p}\cdot\mathbf{x} - i\mathbf{q}\cdot\mathbf{y}} + e^{-i\mathbf{p}\cdot\mathbf{x} + i\mathbf{q}\cdot\mathbf{y}})$$

first do $\int d^3q$ which just sets $\mathbf{p} = \mathbf{q}$. Then collect up everything that is left as:

$$[\hat{\phi}(\mathbf{x}), \hat{\pi}(\mathbf{y})] = \frac{1}{2} \int \frac{d^3p}{(2\pi)^3} (e^{i\mathbf{p}\cdot\mathbf{x} - i\mathbf{p}\cdot\mathbf{y}} + e^{-i\mathbf{p}\cdot\mathbf{x} + i\mathbf{p}\cdot\mathbf{y}}) = i\delta^{(3)}(\mathbf{x} - \mathbf{y})$$

Above set is done by inspection using the integral expression for the delta function.

7 Lecture 8

Can substitute these values for $\hat{\phi}$ and $\hat{\pi}$ into H to give:

$$\hat{H} = \frac{1}{4} \int \frac{d^3 p}{(2\pi)^3} \frac{1}{\omega_{\mathbf{p}}} (-\omega_{\mathbf{p}}^2 + |\mathbf{p}|^2 + m^2) (\hat{a}_{\mathbf{p}} \hat{a}_{\mathbf{p}} + \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{-\mathbf{p}}^\dagger) + (\omega_{\mathbf{p}}^2 + |\mathbf{p}|^2 + m^2) (\hat{a}_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger + \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}})$$

$$\hat{H} = \frac{1}{2} \int \frac{d^3 p}{(2\pi)^3} \omega_{\mathbf{p}} (\hat{a}_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger + \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}})$$

7.1 The Vacuum

This is an infinite set of non-interacting quantum simple harmonic oscillators. The state $|0\rangle$ has the lowest energy and is annihilated by all of the lowering operators:

$$\hat{a}_{\mathbf{p}} = 0 \forall \mathbf{p} \in \mathbb{R}^3$$

$$\hat{H} |0\rangle = E_0 |0\rangle$$

E_0 is the vacuum energy. Use the same trick of reordering the operators so the annihilation operator is always to the right.

$$\hat{H} = \frac{1}{2} \int \frac{d^3 p}{(2\pi)^3} \omega_{\mathbf{p}} (\hat{a}_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger + \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}) = \frac{1}{2} \int \frac{d^3 p}{(2\pi)^3} \omega_{\mathbf{p}} (2\hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}} + (2\pi)^3 \delta^{(3)}(0))$$

This is our first sight of the infinities of quantum field theory. First write as normal ordered hamiltonian

$$\hat{H} = \int \frac{d^3 p}{(2\pi)^3} \omega_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}} + \int \frac{d^3 p}{2} \omega_{\mathbf{p}} \delta^{(3)}(0)$$

Therefore vacuum energy is the remaining constant piece:

$$E_0 = \frac{1}{2} \int d^3 p \omega_{\mathbf{p}} \delta^{(3)}(0)$$

As we have two types of infinity:

IR: Large distance (low energy)

$$(2\pi)^3 \delta^{(3)}(\mathbf{p}) = \int_{\mathbb{R}} d^4 x e^{i\mathbf{x} \cdot \mathbf{p}}$$

divergence of $\mathbf{p} = 0$ comes from the infinite volume of space. So the cure for the infrared infinity is to put the theory in a large box of size L and then consider $L \rightarrow \infty$ by taking $x \in [-\frac{L}{2}, \frac{L}{2}]$.

$$(2\pi)^3 \delta^{(3)}(\mathbf{p}) = \lim_{L \rightarrow \infty} \int_{-\frac{L}{2}}^{\frac{L}{2}} d^3 x e^{i\mathbf{x} \cdot \mathbf{p}}$$

$$(2\pi)^3 \delta^{(3)}(0) = \lim_{L \rightarrow \infty} V_L$$

So this infinity is just reflecting the fact that the volume of the box grows to infinity. All this tells us is we should be calculating an energy density rather than an energy as if we have constant energy density it will give us infinite energy when integrated over infinite space. So we should consider the energy density:

$$\epsilon_0 = \lim_{L \rightarrow \infty} \frac{E_0^{(L)}}{V_L}$$

$$E_0^{(L)} = \int_{-\frac{L}{2}}^{\frac{L}{2}} d^4x \int \frac{d^3p}{(2\pi)^3} \frac{1}{2} \omega_{\mathbf{p}} \implies \epsilon_0 = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2} \omega_{\mathbf{p}}$$

with $\omega_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$. Infinity covers from large $|\mathbf{p}|$ with:

$$\epsilon_0 \sim \int |\mathbf{p}|^3 d|\mathbf{p}| = \infty$$

This is giving a UV divergences as we are adding up infinitely small distances (infinity high frequencies and momentum).

Can cure UV divergence by introducing a UV cutoff so short distance rather than large distance. So introduce $\lambda \gg m \gg L^{-1}$ with L is size of box.

$$\epsilon_0^{(\lambda)} = \int_{|\mathbf{p}| < \lambda} \frac{d^3p}{(2\pi)^3} \frac{1}{2} \omega_{\mathbf{p}} = \frac{4\pi}{2(2\pi)^3} \int_0^\lambda |\mathbf{p}|^2 \sqrt{|\mathbf{p}|^2 + m^2} d|\mathbf{p}| = \frac{1}{16\pi^2} \lambda^4 (1 + O(\frac{m^2}{\lambda^2}))$$

Alternatively put theory on a space time lattice with lattice spacing chosen so $a \ll m^{-1}$, find $\epsilon^{(a)} \sim \frac{1}{a^4}$.

One viewpoint is called effective field theory: You need to define QFT as a continuum limit (a limit in which the cutoff goes to infinity/lattice spacing goes to zero). In order to define this limit you need to identify the quantities that hold fixed in that limit. Theory need to be valid up to some maximum scale.

Specific to vacuum energy: This is a very special type of divergence. Just an additive constant we add to the energy. Can remove this additive constant by considering a different quantisation with ordering chosen so it is normal ordering of all products of field operators ($\hat{\phi}$) and momentum ($\hat{\pi}$) operators.

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Normal ordering symbol: $: X :$ place all annihilation operators to the right.
 $: \hat{a} \hat{a}^2 \hat{a} \hat{a}^\dagger := \hat{a}^\dagger \hat{a}^2$

$$: \hat{H} := \frac{1}{2} \int \frac{d^3 p}{(2\pi)^3} \omega_{\mathbf{p}} : (\hat{a}_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger + \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}) := \int \frac{d^3 p}{(2\pi)^3} \omega_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}$$

This implies $(\epsilon_0)_{\text{Normal}} = 0$ as we can see that as $\hat{a} |0\rangle = 0$ so this will annihilate the ground state.

We care about the ground state energy as gravity cares about energy and says that energy gravitates. The energy momentum tensor appears on the RHS of Einsteins equations seen below:

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} = 8\pi G T_{\mu\nu}$$

We might reasonably ask that if we have some QFT coupled to gravity then:

$$\langle 0 | \hat{T}_{\mu\nu} | 0 \rangle = \epsilon_0 g_{\mu\nu}$$

This leads to problems as if you allow the expectation to the RHS of Einsteins equations, then it contributes a cosmological constant to Einsteins equations. Measurement gave a value to the cosmological constant of $\lambda \sim (10^{-3} \text{eV})^2$ which is very small compared to the energy scales of HEP. We might consider QFT with a UV cutoff λ of $\lambda \sim M_{pl} = \sqrt{\frac{\hbar c}{G}}$ then we would get $\epsilon_0 \sim \lambda^4$ which would give $\lambda \sim (10^{28} \text{eV})^2$. So either QFT does not predict the vacuum energy or it does but it is many many orders of magnitude off from the observation.

Casimir Effect: If we put two metal plates a distance d apart, we can repeat the calculation of the energy density in the presence of these plates. Now we need to take into account the boundary conditions of the field that the electric field has to vanish on the plate. So in this context we get $\epsilon_0 = \epsilon_0(d)$ which therefore exerts a force on these plates and can correctly predict the force on these plates (so vacuum energy is a real thing that can be measured).

8.1 Excited states

Define:

$$|\mathbf{p}\rangle = \hat{a}^\dagger |0\rangle \forall \mathbf{p} \in \mathbb{R}^3$$

Energy

$$\hat{H} - E_0 = \int \frac{d^3p}{(2\pi)^3} \omega_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}$$

$$(\hat{H} - E_0) |\mathbf{p}\rangle = (\hat{H} - E_0) \hat{a}^\dagger |0\rangle = \hat{a}^\dagger (\hat{H} - E_0) |0\rangle + [(\hat{H} - E_0), \hat{a}^\dagger] |0\rangle$$

$$[(\hat{H} - E_0), \hat{a}^\dagger] = \int \frac{d^3p}{(2\pi)^3} \omega_{\mathbf{p}} [\hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{q}}^\dagger] = \int \frac{d^3p}{(2\pi)^3} \omega_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger [\hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{q}}^\dagger] = \int \frac{d^3p}{(2\pi)^3} \omega_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger (2\pi)^3 \delta^{(3)}(\mathbf{q} - \mathbf{p})$$

$$(\hat{H} - E_0) |\mathbf{p}\rangle = \omega_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger |0\rangle = \omega_{\mathbf{p}} |\mathbf{p}\rangle$$

$$E_1 = \omega_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$$

Consider classical momentum of the field: $p^i = \int d^3x T^{0i}$ and $T^{\mu\nu} = \partial^\mu \phi \partial^\nu \phi - \eta^{\mu\nu} \mathcal{L}$. This is converted into a quantum operator on ES2 to give result:

$$\hat{\mathbf{p}} = \int \frac{d^3p}{(2\pi)^3} \mathbf{p} \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}$$

$$\hat{\mathbf{p}} |\mathbf{p}\rangle = \mathbf{p} |\mathbf{p}\rangle$$

Therefore, the state $|\mathbf{p}\rangle$ corresponds to a relativistic particle of mass m , spin 0, momentum \mathbf{p} and energy $E_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$. Do the same process of this to angular momentum, and you get that there must be $\mathbf{J} |0\rangle = 0$ for $|0\rangle$ of a particle in its rest frame which confirms it is a spin 0 particle.

8.2 Multiparticle states

We can construct states corresponding to any number of these particles by simply acting by a lot of operators:

$$|\mathbf{p}_1, \mathbf{p}_2, \dots, \mathbf{p}_n\rangle = \hat{a}_{\mathbf{p}_1}^\dagger \hat{a}_{\mathbf{p}_2}^\dagger \dots \hat{a}_{\mathbf{p}_n}^\dagger |0\rangle$$

This gives us energy $E - E_0 = \sum_{a=1}^n E_{\mathbf{p}_a}$ and momentum $\mathbf{p} = \sum_{a=1}^n \mathbf{p}_a$. These particles must be identical particles as:

$$|\mathbf{p}_1, \mathbf{p}_2\rangle = \hat{a}_{\mathbf{p}_1}^\dagger \hat{a}_{\mathbf{p}_2}^\dagger |0\rangle = \hat{a}_{\mathbf{p}_2}^\dagger \hat{a}_{\mathbf{p}_1}^\dagger |0\rangle = |\mathbf{p}_2, \mathbf{p}_1\rangle$$

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Introduce particle number operator:

$$\hat{N} = \int \frac{d^3p}{(2\pi)^3} \hat{a}_{\mathbf{p}}^\dagger \hat{a}_{\mathbf{p}}$$

$$\hat{N} |\mathbf{p}_1, \mathbf{p}_2, \dots, \mathbf{p}_n\rangle = n |\mathbf{p}_1, \mathbf{p}_2, \dots, \mathbf{p}_n\rangle$$

for free field \hat{H} is conserved quantity: $[\hat{H}, \hat{N}] = 0$. This is not the case when we have interactions

9.1 Normalisation

We can choose without any loss of generality to take our vacuum state to have norm 1: $\langle 0 | 0 \rangle = 1$ and

$$\langle p | q \rangle = \langle 0 | \hat{a}_{\mathbf{p}} \hat{a}_{\mathbf{p}}^\dagger | 0 \rangle = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q})$$

We want to change this normalisation system to make it lorentz invariant (above is not lorentz covariant). Instead of considering directly the overlap of the states we will consider the completeness identity:

$$\hat{I} = \int \frac{d^3 p}{(2\pi)^3} \langle \mathbf{p} | \mathbf{p} \rangle$$

This will help us to renormalise the states to a lorentz invariant way as we need to change this measure by changing this completeness identity as the unit operator is by definition lorentz invariant. $d^3 p$ is an integral over the spatial components so is not lorentz invariant, so we could try and reformulate this as the below **relativistic integration measure**:

$$\int d^4 p \delta(p_\mu p^\mu - m^2) \theta(P_0)$$

First note the δ -function and the θ function (picks out the positive root): impose the constraints

$$p^\mu p_\mu = m^2, p^0 > 0 \implies P_0 = +\sqrt{|\mathbf{p}|^2 + m^2} = E_{\mathbf{p}}$$

$$\int d\mu_p = \int d^3 p \int_0^\infty \frac{d(p_0)^2}{2p_0} \delta(p_0^2 - |\mathbf{p}|^2 - m^2) = \int \frac{d^3 p}{2E_{\mathbf{p}}}$$

Rewrite the completeness relation:

$$|\mathbf{p}^4\rangle = \sqrt{2E_{\mathbf{p}}} |\mathbf{p}\rangle$$

$$\hat{I} = \int \frac{d^3 p}{(2\pi)^3} \frac{1}{2E_{\mathbf{p}}} \langle \mathbf{p}^4 | \mathbf{p}^4 \rangle = \int \frac{d\mu_p}{(2\pi)^3} \langle \mathbf{p}^4 | \mathbf{p}^4 \rangle$$

For multiparticle states are renormalised to:

$$\Pi_{a=1}^n (\sqrt{2E_{\mathbf{p}_a}}) \times |\mathbf{p}_1, \mathbf{p}_2, \dots, \mathbf{p}_n\rangle$$

9.2 Complex Scalar Fields

The details are only change in very small places and we will now highlight these main places:

$$L = \frac{1}{2} \partial_\mu \phi^* \partial^\mu \phi - m^2 \psi^* \psi$$

can rewrite interms of real and imaginary parts and what you get is the sum of the lagrangian for two real scalar fields $\phi_1(x)$ and $\phi_2(x)$ where $\psi = \frac{\phi_1(x) + i\phi_2(x)}{\sqrt{2}}$. This makes manifest a global symmetry:

$$\psi \rightarrow e^{i\alpha} \psi, \psi^* \rightarrow e^{-i\alpha} \psi^*$$

This gives rise to a conserved current that you would not have if you just had one field and it is:

$$j^\mu = i(\partial^\mu \psi^*) \psi - i\psi^* (\partial^\mu \psi)$$

We can write down hamiltonian formulation:

$$\pi = \frac{\mathcal{L}}{\partial \dot{\psi}} = \dot{\psi}^*$$

Try quantisation:

$$\psi(\mathbf{x}, t) \rightarrow \hat{\psi}(\mathbf{x}, t)$$

$$\psi^*(\mathbf{x}, t) \rightarrow \hat{\psi}^\dagger(\mathbf{x}, t)$$

$$[\hat{\psi}(\mathbf{x}), \hat{\pi}(\mathbf{y})] = i\delta^{(3)}(\mathbf{x} - \mathbf{y})$$

$$[\hat{\psi}^\dagger(\mathbf{x}), \hat{\pi}^\dagger(\mathbf{y})] = i\delta^{(3)}(\mathbf{x} - \mathbf{y})$$

Now do mode expansion:

$$\hat{\psi}(\mathbf{x}) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}}} (\hat{b}_{\mathbf{p}} e^{i\mathbf{p}\cdot\mathbf{x}} + \hat{c}_{\mathbf{p}}^\dagger e^{-i\mathbf{p}\cdot\mathbf{x}})$$

get e.g.

$$[\hat{b}_{\mathbf{p}}, \hat{b}_{\mathbf{q}}^\dagger] = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q})$$

$$|\mathbf{p}, +\rangle = \hat{c}_{\mathbf{p}}^\dagger |0\rangle$$

$$\hat{b}_{\mathbf{p}}^\dagger |0\rangle = \hat{b}_{\mathbf{p}}^\dagger |\mathbf{p}, -\rangle$$

$$:\hat{H}: = \int \frac{d^3p}{(2\pi)^3} \omega_{\mathbf{p}} (\hat{b}_{\mathbf{p}}^\dagger \hat{b}_{\mathbf{p}} + \hat{c}_{\mathbf{p}}^\dagger \hat{c}_{\mathbf{p}})$$

So we have a $U(1)$ conserved charge by starting from the classical charge we can get:

$$\hat{Q} = i \int d^3x (\hat{\pi} \hat{\psi} - \hat{\psi}^\dagger \hat{\pi}^\dagger) = \int \frac{d^3p}{(2\pi)^3} \omega_{\mathbf{p}} (\hat{b}_{\mathbf{p}}^\dagger \hat{b}_{\mathbf{p}} - \hat{c}_{\mathbf{p}}^\dagger \hat{c}_{\mathbf{p}})$$

To check this is conserved we have to check $[\hat{Q}, \hat{H}] = 0$ which is easy as all the things commute. If we evaluate the conserved charge:

$$\hat{Q} |\mathbf{p}, \pm\rangle = \pm |\mathbf{p}, \pm\rangle$$

So for the complex scalar field we have two particles that in addition to all the quantum numbers we have a charge number which distinguishes between particle and antiparticle. We have operators that conserve the $U(1)$ charge. Even in the interacting theory this symmetry is conserved.

9.3 Time dependence

So far we have worked in the schrodinger picture, so really:

$$\hat{\phi}(\mathbf{x}) = \hat{\phi}_S(\mathbf{x})$$

so there is no time dependence in the field operator the time dependence all lives in the space.

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$$i \frac{d}{dt} |\mathbf{p}(t)\rangle_S = H : |\mathbf{p}(t)\rangle_S = E_{\mathbf{p}} |\mathbf{p}(t)\rangle_S$$

implies

$$|\mathbf{p}(t)\rangle_S = e^{-iE_{\mathbf{p}}t} |\mathbf{p}(0)\rangle_S$$

Heisenberg picture

States (time independent):

$$|\psi\rangle_H = e^{i\hat{H}t} |\psi(t)\rangle_S$$

Operators (time dependent):

$$\hat{O}_H = e^{i\hat{H}t} \hat{O}_S e^{-i\hat{H}t}$$

Define Heisenberg picture field operator:

$$\hat{\phi}(x) = e^{i\hat{H}t} \hat{\phi}(\mathbf{x}) e^{-i\hat{H}t}$$

LHS x is four vector, RHS is three vector (I think not sure)
This obeys the Klein-Gordon operator. As shown in ES2 and below using mode expansion:

$$\hat{\phi}(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_p}} ((\hat{a}_p)_H e^{i\mathbf{p}\cdot\mathbf{x}} + (\hat{a}_p^\dagger)_H e^{-i\mathbf{p}\cdot\mathbf{x}})$$

$$(\hat{a}_p)_H = e^{i\hat{H}t} \hat{a}_p e^{-i\hat{H}t} \implies \frac{d}{dt}(\hat{a}_p)_H = i[\hat{H}, (\hat{a}_p)_H] = -iE_p (\hat{a}_p)_H$$

Therefore:

$$(\hat{a}_p)_H(t) = e^{-iE_p t} (\hat{a}_p)_H(0)$$

similarly

$$(\hat{a}_p^\dagger)_H(t) = e^{+iE_p t} (\hat{a}_p^\dagger)_H(0)$$

so (below equation needs fixing watch the 15 min of lecture 10)

$$\hat{\phi}(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_p}} ((\hat{a}_p)_H e^{i\mathbf{p}\cdot\mathbf{x}} + (\hat{a}_p^\dagger)_H e^{-i\mathbf{p}\cdot\mathbf{x}})$$

$$\mathbf{p} \cdot \mathbf{x} = p_\mu x^\mu = E_p t - \mathbf{p} \cdot \mathbf{x}$$

10.1 Interacting QFT

$$\mathcal{L} = \frac{1}{2} \partial_\mu \tilde{\phi} \partial^\mu \tilde{\phi} - V(\phi)$$

implies hamiltonian

$$H = \int d^3x \frac{1}{2} \dot{\tilde{\phi}}^2 + \frac{1}{2} |\nabla \tilde{\phi}|^2 + V(\tilde{\phi})$$

If we want to have a sensible theory we want it to be bounded below at some point giving the lowest energy:

$$\tilde{\phi}(\mathbf{x}, t) = \tilde{\phi}_0$$

This must be a minimum point so:

$$\frac{\partial V}{\partial \tilde{\phi}} \Big|_{\tilde{\phi}=\tilde{\phi}_0} = 0 \tag{10}$$

$$\frac{\partial^2 V}{\partial \tilde{\phi}^2} \Big|_{\tilde{\phi}=\tilde{\phi}_0} \geq 0 \tag{11}$$

WLOG take $V(\phi_0) = 0$. Now expand the field around the minimum:

$$\phi(x) = \tilde{\phi}(x) - \tilde{\phi}_0$$

Therefore:

$$V(\tilde{\phi}) = \sum_{n=0}^{\infty} \frac{1}{n!} \frac{\partial^n V}{\partial \tilde{\phi}^n} \Big|_{\tilde{\phi}=\tilde{\phi}_0} \phi^n$$

The result is a lagrangian of the following form:

$$\mathfrak{L} = \mathfrak{L}_0 + \mathfrak{L}_{\text{int}}$$

with \mathfrak{L}_0 being the free terms and the $\mathfrak{L}_{\text{int}}$ being the interacting terms (all generated by this expansion around the minimum).

$$\mathfrak{L}_0 = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{m^2}{2} \phi^2$$

$$m^2 = \frac{\partial^2 V}{\partial \tilde{\phi}^2} \Big|_{\tilde{\phi}=\tilde{\phi}_0}$$

$$\mathfrak{L}_{\text{int}} = - \sum_{n=3}^{\infty} \frac{\lambda_n}{n!} \phi^n$$

$$\lambda_n = \frac{\partial^n V}{\partial \tilde{\phi}^n} \Big|_{\tilde{\phi}=\tilde{\phi}_0}$$

If we include general interacting terms there is no way to make progress, we need to work in a perturbation expansion λ_n telling us how far we deviate from the free state. These λ_n are "coupling" constants and are parameters of the theory that we put in. We can only hope to calculate it when it is close to the free field theory (when the coupling constants are "small").

Dimensional analysis

$$[S] = 0, [x] = -1, [\partial_\mu] = +1$$

$$S = \int_{\mathbb{R}^{3,1}} d^4x \mathfrak{L}$$

so $[\mathfrak{L}] = 4$ which implies that:

$$[\partial_\mu \phi \partial^\mu \phi] = 4 \implies [\phi] = +1$$

$$\left[\frac{\lambda_n \phi^n}{n!} \right] = 4 \implies [\lambda_n] = 4 - n$$

so if $n \neq 4$ then $[\lambda_n] \neq 0$.

Consider process with $E \geq m$. We must be expanding in an effective dimensionless parameter, so cannot be just λ, λ^2 as that would have different dimensions so need to add a power of energy into the mix.

$$\tilde{\lambda}_n = \lambda_n (E)^{n-4}$$

This gives $[\tilde{\lambda}] = 0$ so can expand nicely. Consider different cases:

If $n < 4$ then called **relevant coupling**. This will become weakly coupled at high energy as if $\lambda_n \ll (E)^{4-n}$ then $\tilde{\lambda}_n \ll 1$. If we choose $\lambda_n \ll (m)^{4-n}$ then the coupling can be small for all energies, so perturbation theory is good for all energy scales:

If $n = 4$ called **marginal coupling**. Perturbation theory is good for $\lambda_4 \ll 1$.

If $n > 4$ it is called **irrelevant coupling**. Correspondingly it means that compared to the relevant case the high and low energy cases are swapped around. So in general the coupling constant can be small at low energy, but if you go to high enough energy the coupling constant will be large. So perturbation theory is only good at $\lambda \ll (E)^{4-n}$. The key fact is that irrelevant couplings lead to non-renormalisable theories.

Examples:

ϕ^4 -theory:

$$\mathfrak{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{m^2}{2} \phi^2 - \frac{\lambda}{4!} \phi^4$$

As $n = 4$ perturbation theory is good for $\lambda \ll 1$. With the exception of adding a ϕ^3 term this is the only renormalisable Lagrangian of a single scalar field.

10.1.1 Scalar-Yukawa Theory

$$\phi : \mathbb{R}^{3,1} \rightarrow \mathbb{R}$$

$$\psi : \mathbb{R}^{3,1} \rightarrow \mathbb{C}$$

$$\mathfrak{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{m^2}{2} \phi^2 + \frac{1}{2} \partial_\mu \psi \partial^\mu \psi^* - \frac{M^2}{2} \psi \psi^* - g \psi^* \psi \phi$$

This will be weakly coupled if $|g| \ll m, M$.

11 Lecture 11

11.1 Scattering

Looking for the amplitude of the transition from an initial state to a final state: $A_{i \rightarrow f} = \lim_{T \rightarrow \infty} [{}_S \langle f, t = +\frac{T}{2} | e^{-i\hat{H}T} | i, t = -\frac{T}{2} \rangle_S]$. This will be constrained by the overall conservation laws like energy, momentum and charge. We are going to think about our theory perturbatively, so \hat{H} is thought of as:

$$\hat{H} = \hat{H}_0 + \hat{H}_{int}$$

For example in ϕ^4 - theory:

$$\hat{H}_0 = \frac{1}{2} \int_{\mathbb{R}^3} d^3x (\hat{\pi}^2 + |\nabla \hat{\phi}|^2 + m^2 \hat{\phi}^2)$$

$$\hat{H}_{int} = \frac{\lambda}{4!} \int_{\mathbb{R}^3} d^3x \hat{\phi}^4$$

KEY ASSUMPTION: the particles are free in the infinite past and the infinite future (independent of thinking about perturbation theory just a consequence of locality). This implies that $|i\rangle$ and $|f\rangle$ are eigenstates of \hat{H}_0 . You can justify this but we won't as it is a very involved discussion that can be found in textbooks.

We are going to work in the **interaction picture** as it is very convenient for perturbation theory. It is Heisenberg with respect to \hat{H}_0 but Schrodinger with respect to \hat{H}_{int} :

States:

$$|\phi(t)\rangle_I = e^{i\hat{H}_0 t} |\psi(t)\rangle_S$$

Operators:

$$\hat{O}_I(t) = e^{i\hat{H}_0 t} \hat{O}_S e^{-i\hat{H}_0 t}$$

Time evolution:

For schrodinger you get:

$$i \frac{d}{dt} |\psi(t)\rangle_S = \hat{H} |\psi(t)\rangle_S \implies |\psi(t)\rangle_S = e^{-i\hat{H}t} |\psi(0)\rangle_S$$

Interaction picture:

$$i \frac{d}{dt} |\psi(t)\rangle_I = i \frac{d}{dt} (e^{i\hat{H}_0 t} |\psi(t)\rangle_S) = -\hat{H}_0 e^{i\hat{H}_0 t} |\psi(t)\rangle_S + e^{i\hat{H}_0 t} i \frac{d}{dt} |\psi(t)\rangle_S$$

$$i \frac{d}{dt} |\psi(t)\rangle_I = -\hat{H}_0 e^{i\hat{H}_0 t} |\psi(t)\rangle_S + e^{i\hat{H}_0 t} \hat{H} |\psi(t)\rangle_S = e^{i\hat{H}_0 t} \hat{H}_{int} |\psi(t)\rangle_S$$

$$i \frac{d}{dt} |\psi(t)\rangle_I = \hat{H}_I(t) |\psi(t)\rangle_I \quad (12)$$

for $\hat{H}_I = e^{i\hat{H}_0 t} \hat{H}_{int} e^{-i\hat{H}_0 t}$. Compare with ODE:

$$i \frac{d}{dt} F(t) = g(t) F(t) \implies F(t) = e^{\frac{1}{i} \int^t dt' g(t')} F(0)$$

The order of H is important as $[H(t), H(t')] \neq 0$ in general which makes this more complex than the ODE case. As can be seen by guessing the solution:

$$|\psi(t)\rangle = e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')} |\psi(0)\rangle$$

This DOES NOT WORK though as can be seen by substituting it into 12. There is an easy fix though, for any operator valued function $\hat{O}(t)$ define the following time ordering symbol:

$$T[\hat{O}(t_1) \hat{O}(t_2)] = \begin{cases} \hat{O}(t_1) \hat{O}(t_2) & t_1 \geq t_2 \\ \hat{O}(t_2) \hat{O}(t_1) & t_1 \leq t_2 \end{cases}$$

So a valid solution is actually:

$$|\psi(t)\rangle_I = T(e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')}) |\psi(0)\rangle$$

Very slick two line proof in David Tong's notes, now we will check that the first couple terms match:

$$T(e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')}) = T(I + \frac{1}{i} \int^t dt' \hat{H}_I(t') - \frac{1}{2} \int^t \int^t dt' dt'' (\hat{H}_I(t') \hat{H}_I(t'')) + \dots)$$

$$T(e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')}) = I + \frac{1}{i} \int^t dt' \hat{H}_I(t') - \int^t dt' \int^{t'} dt'' (\hat{H}_I(t') \hat{H}_I(t'')) + \dots$$

Check equation 12:

$$i \frac{d}{dt} T(e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')}) |\psi(0)\rangle = (\hat{H}_I(t) + \hat{H}_I(t) \frac{1}{i} \int^t dt' \hat{H}_I(t') + \dots) |\psi(0)\rangle$$

$$i \frac{d}{dt} T(e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')}) |\psi(0)\rangle = \hat{H}_I(t) (I + \frac{1}{i} \int^t dt' \hat{H}_I(t') + \dots) |\psi(0)\rangle \approx \hat{H}_I(t) T(e^{\frac{1}{i} \int^t dt' \hat{H}_I(t')}) |\psi(0)\rangle$$

So the amplitude can be written in the interaction picture:

$$A_{i \rightarrow f} = \lim_{T \rightarrow \infty} [{}_S \langle f, t = +\frac{T}{2} | i, t = \frac{T}{2} \rangle_S] = \lim_{T \rightarrow \infty} [{}_I \langle f, t = \frac{T}{2} | i, t = \frac{T}{2} \rangle_I] = \lim_{T \rightarrow \infty} [{}_I \langle f, t |$$

$$A_{i \rightarrow f} = \lim_{T \rightarrow \infty} [{}_S \langle f, t = +\frac{T}{2} | | i, t = \frac{T}{2} \rangle_S] T \left(e^{\frac{1}{i} \int_{-\frac{T}{2}}^{\frac{T}{2}} dt' \hat{H}_I(t')} \right) | i, t = -\frac{T}{2} \rangle_I] A_{i \rightarrow f} =$$

$$\lim_{T \rightarrow \infty} [{}_S \langle f, t = +\frac{T}{2} | | i, t = \frac{T}{2} \rangle_S] = \frac{T}{2} S | i, t = -\frac{T}{2} \rangle_I] \text{THESE LAST FEW EQUATIONS MESSEDDU}$$

12 Lecture 12

12.1 Scattering in ϕ^4 theory

$$\hat{H}_0 = \frac{1}{2} \int_{\mathbb{R}^3} d^3x (\hat{\pi}^2 + |\nabla \hat{\phi}|^2 + m^2 \hat{\phi}^2)$$

$$\hat{H}_{int} = \frac{\lambda}{4!} \int_{\mathbb{R}^3} d^3x \hat{\phi}^4$$

$$\hat{H}_I(t) = e^{i\hat{H}_0 t} \hat{H}_{int} e^{-i\hat{H}_0 t} = \int d^3x (e^{i\hat{H}_0 t} \hat{\phi}(x) e^{-i\hat{H}_0 t})^4 = \int d^3x \hat{\phi}_I^4$$

$$\hat{\phi}(x) = \hat{\phi}_+(x) + \hat{\phi}_-(x)$$

$$\hat{\phi}_+(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}}} \hat{a}_{\mathbf{p}} e^{-ip \cdot x}$$

$$p^\mu = (E_{\mathbf{p}}, \mathbf{p}0)$$

$$p \cdot x - p_\mu x^\mu = E_{\mathbf{p}} t - \mathbf{p} \cdot \mathbf{x}$$

$$\hat{\phi}_-(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}}} \hat{a}_{\mathbf{p}}^\dagger e^{ip \cdot x}$$

Think about initial state as $|i\rangle = |p_1, \dots, p_{n_i}\rangle = \prod_{a=1}^{n_i} \sqrt{2E_{\mathbf{p}_a}} \hat{a}_{\mathbf{p}_a}^\dagger |0\rangle$ and the final state as $|f\rangle = |p_1, \dots, p_{n_f}\rangle = \prod_{a=1}^{n_f} \sqrt{2E_{\mathbf{p}_a}} \hat{a}_{\mathbf{p}_a}^\dagger |0\rangle$. So:

$$\Lambda_{i \rightarrow f} = \langle f | T[\exp(\frac{1}{i} \frac{\lambda}{4!} \int_{\mathbb{R}^3} d^4x \hat{\phi}_I^4)] | i \rangle$$

To work this out explicitly we will have to expand the exponential and work it out order by order:

$$A_{i \rightarrow f} = \sum_{l=0}^{\infty} \lambda^l A_{i \rightarrow f}^{(l)}$$

$$A_{i \rightarrow f}^{(l)} = \frac{1}{l!} \left(\frac{1}{i(4!)} \right)^l \int d^4x_1 \dots \int d^4x_l \langle f | T[\hat{\phi}^4(x_1) \hat{\phi}^4(x_2) \dots \hat{\phi}^4(x_l)] | i \rangle$$

Need to convert T ordering into normal ordering. Simplest case:

$$T_\alpha = T[\hat{\phi}(x)\hat{\phi}(y)]$$

$$T_\alpha = \hat{\phi}(x)\hat{\phi}(y) = (\hat{\phi}^+(x) + \hat{\phi}^-(x))(\hat{\phi}^+(y) + \hat{\phi}^-(y)) = \hat{\phi}^+(x)\hat{\phi}^+(y) + \hat{\phi}^+(x)\hat{\phi}^-(y) + \hat{\phi}^-(x)\hat{\phi}^+(y) + \hat{\phi}^-(x)\hat{\phi}^-(y)$$

$$T_\alpha = \hat{\phi}^+(x)\hat{\phi}^+(y) + \hat{\phi}^-(x)\hat{\phi}^+(y) + [\hat{\phi}^+(x), \hat{\phi}^-(y)] + \hat{\phi}^-(x)\hat{\phi}^+(y) + \hat{\phi}^-(x)\hat{\phi}^-(y)$$

$$T[\hat{\phi}(x)\hat{\phi}(y)] =: \hat{\phi}(x)\hat{\phi}(y) + D(x-y)$$

where

$$D(x-y) = [\hat{\phi}^+(x), \hat{\phi}^-(y)] = \int \frac{d^3p d^3q}{(2\pi)^6} \frac{1}{\sqrt{2E_p E_q}} [\hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{q}}^\dagger] e^{i(q \cdot y - p \cdot x)} = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2E_p} \hat{a}_{\mathbf{p}}^\dagger e^{ip \cdot (x-y)}$$

for $x^0 < y^0$.

$$T[\hat{\phi}(x)\hat{\phi}(y)] =: \hat{\phi}(x)\hat{\phi}(y) : + \Delta_F(x-y) \quad (13)$$

$\Delta_F(x-y)$ is the Feynman propagator:

$$\Delta_F(x-y) = \begin{cases} D(x-y) & x^0 > y^0 \\ D(y-x) & x^0 < y^0 \end{cases} = \langle 0 | T[\hat{\phi}(x)\hat{\phi}(y)] | 0 \rangle$$

12.2 Wicks Theroem

$$T(\phi_1 \phi_2 \dots \phi_n) =: \phi_1 \dots \phi_n : + : \text{all possible contractions} :$$

12.3 Integral representation of Δ_F

$$I_F(x-y) = \int_{C_F} \frac{d^4p}{(2\pi)^4} \frac{i}{p^2 - m^2} e^{-ip \cdot (x-y)}$$

Split up over p^0 and \mathbf{p} :

$$I_F(x-y) = \int \frac{d^3p}{(2\pi)^3} i e^{-i\mathbf{p} \cdot (\mathbf{x}-\mathbf{y})} \int_{C_F} dp_0 I(p_0)$$

$$I(p_0) = \frac{e^{-ip_0 \cdot (x-y)^0}}{(p_0)^2 - |\mathbf{p}|^2 - m^2}$$

This has poles that lie in the real axis so you cant straight fowardly you need to specify the contour.

$$\frac{1}{p_0^2 - E_{\mathbf{p}}^2} = \frac{1}{2E_{\mathbf{p}}} \left(\frac{1}{p_0 - E_{\mathbf{p}}} - \frac{1}{p_0 + E_{\mathbf{p}}} \right)$$

So the poles are at $p_0 = \pm E_{\mathbf{p}}$ wiht residues:

$$Res_{p^0 = \pm E_{\mathbf{p}}} (I(p_0)) = \pm \frac{1}{2E_{\mathbf{p}}} e^{iE_{\mathbf{p}}(x^0 - y^0)}$$

13 Lecture 13

The contour is defined as going in to the lower half plane for $-E_{\mathbf{p}}$ and into the upper half plane for $E_{\mathbf{p}}$. Now want to evaluate this integral:

$$\int \frac{d^3p}{(2\pi)^4} i e^{i\mathbf{p}(\mathbf{x}-\mathbf{y})} \int_{C_F} dp_0 I(p_0)$$

Set $p_0 = Re(p_0) + iIm(p_0)$ clearly the behaviour of $I(p_0)$ as $p_0 \rightarrow \infty$ will be dominated by the real exponential part:

$$I(p_0) \sim e^{Im(p_0)(x^0-y^0)}$$

this decays rapidly as p^0 in the the LHP for $x^0 > y^0$ and in the UHP for $x^0 < y^0$. Now apply Jordan's Lemma and close in the upper or lower plane. For $x^0 > y^0$ we need to close the contour in the LHP and use cauchy's theorem to only pick up the residue from the positive pole (bearing in mind cauchy's theorem needs a counterclockwise contour and we have a clockwise one so pick up an extra minus sign):

$$\Delta_F(x-y) = \int \frac{d^3p}{(2\pi)^4} i e^{i\mathbf{p}(\mathbf{x}-\mathbf{y})} \times -2\pi i Res_{p_0=E_{\mathbf{p}}}(I(p_0)) = \int \frac{d^3p}{(2\pi)^4} i e^{i\mathbf{p}(\mathbf{x}-\mathbf{y})} \frac{1}{2E_{\mathbf{p}}} e^{-iE_{\mathbf{p}}(x^0-y^0)}$$

$$\Delta_F(x-y) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2E_{\mathbf{p}}} e^{-ip(x-y)} = D(x-y)$$

for $x^0 < y^0$ close the contour in the UHP and pick up pole $p = -E_{\mathbf{p}}$ and reproduces the correct time ordered feyman propagator.

13.0.1 $i\epsilon$ prescription

Consider a contour that goes along the real axis and alter the integral a little bit to push the poles above on the left and down on the right. So poles now at $p^0 = E_{\mathbf{p}} - i\epsilon$ and $p^0 = -E_{\mathbf{p}} + i\epsilon$:

$$\Delta_F(x-y) = \lim_{\epsilon \rightarrow 0^+} \left(\int \frac{d^4p}{(2\pi)^4} \frac{i e^{-ip(x-y)}}{p^2 - m^2 + i\epsilon} \right)$$

This is a more convenient way of automatically making it time ordered.

14 Lecture 14

$$\mathfrak{L} = (\partial_\mu \psi^* \partial^\mu \psi) - M^2 \psi^* \psi + \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{m^2}{2} \phi^2 - g \phi \psi^* \psi$$

We have $[g] = +1$ and $g \ll m, M$. Mode expansions:

$$\begin{aligned}\hat{\phi}(x) &= \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2\epsilon_p}} (\hat{a}_{\mathbf{p}} e^{-ipx} + \hat{a}_{\mathbf{p}}^\dagger e^{ipx}) \\ \hat{\psi}(x) &= \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_p}} (\hat{b}_{\mathbf{p}} e^{-ipx} + \hat{c}_{\mathbf{p}}^\dagger e^{ipx}) \\ \hat{\psi}^\dagger(x) &= \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_p}} (\hat{b}_{\mathbf{p}}^\dagger e^{-ipx} + \hat{c}_{\mathbf{p}} e^{ipx})\end{aligned}$$

with $\epsilon = \sqrt{|\mathbf{p}|^2 + m^2}$ and $E = \sqrt{|\mathbf{p}|^2 + M^2}$. We also have $\hat{a}_{\mathbf{p}} |0\rangle = \hat{b}_{\mathbf{p}} |0\rangle = \hat{c}_{\mathbf{p}} |0\rangle = 0$ and the standard commutation relations apply like $[\hat{a}_{\mathbf{p}}, \hat{a}_{\mathbf{p}}^\dagger] = (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q})$. We will have \hat{a}^\dagger creates a ϕ particle, \hat{b}^\dagger creates a ψ particle and \hat{c}^\dagger creates a anti- ψ particle. Going to try and consider:

$$\psi + \psi \rightarrow \psi + \psi$$

Going to have $p_1 = (E_{\mathbf{p}_1}, \mathbf{p}_1)$ and p_2 going in, and p'_1 and p'_2 going out. We can treat these initial and final states as eigenstates of the hamiltonian so the appropriate states to choose are:

$$|i\rangle = \sqrt{2E_{\mathbf{p}_1}} \sqrt{2E_{\mathbf{p}_2}} \hat{b}_{\mathbf{p}_1}^\dagger \hat{b}_{\mathbf{p}_2}^\dagger |0\rangle, |f\rangle = \sqrt{2E_{\mathbf{p}'_1}} \sqrt{2E_{\mathbf{p}'_2}} \hat{b}_{\mathbf{p}'_1}^\dagger \hat{b}_{\mathbf{p}'_2}^\dagger |0\rangle$$

$$A_{i \rightarrow f} = \langle f | \hat{S} | i \rangle, \hat{S} = T[\exp(\frac{g}{i} \int d^4x \hat{\phi} \hat{\psi}^\dagger \hat{\psi})]$$

Only way to make progress is to expand the exponentials in powers of the coupling constant. Have to think what is the first order to which scattering can occur. Restrict ourselves to the case the $\{p_1, p_2\} \neq \{p'_1, p'_2\}$ so will certainly get zero to the first order in the expansion as they correspond to linearly independent states in our hilbert space as they have different momenta. In order to get a non-zero answer we need to bring down terms with the same operators as in the $|i\rangle$ and $|f\rangle$ states as only contributions will come from comunators. So need to bring down at least two creation operators (\hat{b}^2) and two annihaltion operators ($\hat{b}^{2\dagger}$). So we need to bring down at least two powers of the interaction hamiltonian.

$$A_{i \rightarrow f}^{(2)} = \frac{1}{2} \left(\frac{g}{i}\right)^2 \int d^4x_1 \int d^4x_2 M(x_1, x_2)$$

$$M(x_1, x_2) = \langle f | T[\hat{\phi}(x_1) \hat{\psi}^\dagger(x_1) \hat{\psi}(x_1) \hat{\phi}(x_2) \hat{\psi}^\dagger(x_2) \hat{\psi}(x_2)] | i \rangle$$

Using wicks theorem we can express this as a sum of all possible contractions plus the normal order. We need an explicit power of \hat{b}^2 and $\hat{b}^{2\dagger}$ so if we contract any of the ψ fields they will give zero when stuck back into $|f\rangle$ and $|i\rangle$. However, we are fine to contract the ϕ fields so:

$$M(x_1, x_2) = \Delta_F^{(\phi)}(x_1 - x_2) \langle f | : \psi^\dagger(x_1) \hat{\psi}(x_1) \psi^\dagger(x_2) \hat{\psi}(x_2) : | i \rangle$$

Immediate simplification comes from the fact that only the terms with the \hat{b} in give non zero terms so ψ^{dagger} only consists of creation terms and ψ only consists of annihilation operators. so

$$M(x_1, x_2) = \Delta_F^{(\phi)}(x_1 - x_2) \langle f | \psi^\dagger(x_1) \psi^\dagger(x_2) \hat{\psi}(x_1) \hat{\psi}(x_2) | i \rangle = \Delta_F^{(\phi)}(x_1 - x_2) N$$

By inserting a complete set of free particle eigenstates $I = \sum_\phi |\phi\rangle \langle\phi| = |0\rangle \langle 0| + \sum_{n \geq 1} |n \text{ particle}\rangle \langle n \text{ particle}|$. So can write N as:

$$N = N_i N_f, N_i = \langle 0 | \hat{\psi}(x_1) \hat{\psi}(x_2) | i \rangle = \sqrt{2E_{\mathbf{p}_1}} \sqrt{2E_{\mathbf{p}_2}} \int \frac{d^3 p}{(2\pi)^3} \int \frac{d^3 q}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}} \sqrt{2E_{\mathbf{q}}}} e^{-ipx_1 - iqx_2} \langle 0 | \hat{b}_{\mathbf{p}} \hat{b}_{\mathbf{q}}$$

Use commutator functions:

$$A = (2\pi)^6 (\delta^{(3)}(\mathbf{p}_1 - \mathbf{p}) \delta^{(3)}(\mathbf{p}_2 - \mathbf{q}) + \delta^{(3)}(\mathbf{p}_2 - \mathbf{p}) \delta^{(3)}(\mathbf{p}_1 - \mathbf{q}))$$

So:

$$N_i = (e^{-ip_1 x_1 - ip_2 x_2} + e^{-ip_2 x_1 - ip_1 x_2})$$

Can do same calculation to get answer for N_f :

$$N_f = \langle f | \hat{\psi}^\dagger(x_1) \hat{\psi}^\dagger(x_2) | 0 \rangle = (e^{ip'_1 x_1 + ip'_2 x_2} + e^{ip'_2 x_1 + ip'_1 x_2})$$

So

$$A_{i \rightarrow f}^{(2)} = \frac{1}{2} \left(\frac{g}{i}\right)^2 \int d^4 x_1 \int d^4 x_2 N_f N_i \Delta_F^{(\phi)}(x_1 - x_2) = \frac{1}{2} \left(\frac{g}{i}\right)^2 \int d^4 x_1 \int d^4 x_2 N_f N_i \int \frac{dk^4}{(2\pi)^4} \frac{ie^{ik(x_1 - x_2)}}{k^2 - m^2 + i\epsilon}$$

$$A_{i \rightarrow f}^{(2)} = \frac{1}{2} \left(\frac{g}{i}\right)^2 \int \int \int \frac{d^4 x_1 d^4 x_2 dk^4}{(2\pi)^4} \frac{i}{k^2 - m^2 + i\epsilon} (e^{ix_1(k+p'_1-p_1)} e^{ix_2(k-p'_2+p_2)} + e^{ix_1(k+p'_2-p_1)} e^{ix_2(k-p'_1+p_2)} + x_1$$

Last term above is the same as the first with x_1 and x_2 exchanged which is identical as they are symmetric in the metric. First evaluate by integral over x_1 and x_2 as the integrand depends purely exponential so gives two set of four dimensional delta functions and then can do integral over $d^4 k$. The result from doing this is:

$$A_{i \rightarrow f}^{(2)} = i(-ig)^2 (2\pi)^4 \delta^{(4)}(p_1 + p_2 - p'_1 - p'_2) \times \left(\frac{1}{(p_1 - p'_1)^2 - m^2 + i\epsilon} + \frac{1}{(p_1 - p'_2)^2 - m^2 + i\epsilon} \right)$$

Above section has some incorrect signs that are corrected on the moodle notes

15 Lecture 15

Energy momentum conservation:

$$\mathbf{p}_1 + \mathbf{p}_2 = \mathbf{p}'_1 + \mathbf{p}'_2$$

$$E_{\mathbf{p}_1} + E_{\mathbf{p}_2} = E_{\mathbf{p}'_1} + E_{\mathbf{p}'_2} =$$

In COM frame:

$$\mathbf{p}_1 + \mathbf{p}_2 = \mathbf{p}'_1 + \mathbf{p}'_2 = 0 \implies |\mathbf{p}_1| = |\mathbf{p}_2|, |\mathbf{p}'_1| = |\mathbf{p}'_2|$$

$$E_{\mathbf{p}_1} = E_{\mathbf{p}_2} = E_{\mathbf{p}'_1} = E_{\mathbf{p}'_2}$$

In COM frame: $(p - p')^\mu = (0, \mathbf{p}_1 - \mathbf{p}'_1)$ so $(p - p_1)^2 \leq 0$. As the denominator $(p_1 - p'_1)^2 - m^2 < 0$ and similarly $(p_1 - p'_2)^2 - m^2 < 0$. Therefore it is safe to assume that $\epsilon = 0$ as the other part of the denominator is never zero.

15.1 Feynmann diagrams

Contributions to $\langle f | \hat{S} - I | i \rangle$ can be written in a diagrammatic language of feynamn diagrams. Now need to move to paper!!

15.2 Fermions

So far we have only covered scalar fields that transform under the lorentz transformation as:

$$\phi(x) \rightarrow \phi(\Lambda^{-1}x)$$

We have shown that $\hat{\phi}$ creates a particle of spin 0. We want to make a particle of half-spin.

Non-relativistic description of a spin: $\hat{\mathbf{S}} = (\hat{S}_x, \hat{S}_y, \hat{S}_z)$, \hat{S}^2, \hat{S}_z paricles of spin $s = 0, \frac{1}{2}, \frac{3}{2}$ haave $2s + 1$ spin states:

$$\hat{S}^2 |s, s_z\rangle = s(s + 1)\hbar^2 |s, s_z\rangle$$

$$\hat{S}_z |s, s_z\rangle = s_z\hbar |s, s_z\rangle$$

so get s_z at interval spacings with

$$-s \leq s_z \leq s$$

Need to be able to describe these independant states so this motivates the fact we need multiple components to the field. So need fields that transform in non-trivial representations of the Lorentz group.

16 Lecture 16

Consider the lorentz group: $G_L = \{\Lambda \in Mot_4(\mathbb{R})\}$ where $\Lambda\eta^T = \eta, det = +1$. Consider a representation D of G_L of dimension N which is a way of representing the group elements in terms of matrices:

$$D : G \rightarrow Mat_N(\mathbb{C})$$

(with $det D \neq 0$ I think)

This map is a smooth homomorphism and so preserves the group multiplication, so:

$$D(\Lambda_1)D(\Lambda_2) = D(\Lambda_1\Lambda_2)\forall\Lambda_1, \Lambda_2 \in G$$

Matrices are linearly maps that act on vectors and in this case they act on the representation space. So the representation acts on vectors in the representation space $V = \mathbb{C}^N$. This formalises the variety of ways the same group can act on vectors. Now we want to think about what it means for a field to transform under the lorentz group. So we will write an explicit form of the matrices.

$$D(\Lambda_{AB})$$

A field is a function from spacetime into the field space: $\Psi : \mathbb{R}^{3,1} \rightarrow \mathbb{C}^N$. such a field transforms in representation D of G_L if:

$$\psi_A(x) \rightarrow^{\Lambda \in G_L} \psi'_A(x) = \sum_{B=1}^N D(\Lambda)_{AB} \psi_B(\Lambda^{-1} \cdot x)$$

The field is now a vector and the lorentz group acts on these elements under different representation of the lorentz group.

16.0.1 Representations of the Lorentz group

Scalar field: $\phi(x) \rightarrow \phi(\Lambda^{-1} \cdot x)$, $N = 1$ and $D(\Lambda) = 1 \forall \Lambda \in G_L$. This is the trivial representation.

4-vector field: $A^\mu(x) \rightarrow \Lambda^\mu_\nu A^\nu(\Lambda^{-1} \cdot x)$, this is the fundamental representation which is where you simply use element itself as its representation (only exists for matrix Lie groups). $D(\Lambda) = \Lambda$, $N = 4$.

16.1 Spinor Representation

Going to spend the rest of the lecture deriving this. It is generally hard to construct representations, so we find representations of the Lie Algebra and

then utilise the relationship between them. In physics this means finding a representation of the infinitesimal transformations.

Lie Algebra $\mathbb{L}(G_L)$: Correspond to the tangent vectors at the identity on the Lie group and is useful as there exists an exponential map from the Lie Algebra to the Lie Group:

$$\text{Exp} : \mathbb{L}(G_L) \rightarrow G_L$$

A representation of the Lie Algebra. As $\mathbb{L}(G_L)$ is a linear map, so want matrices:

$$R : \mathbb{L}(G_L) \rightarrow \text{Mat}_N(\mathbb{C})$$

As algebra of a Lie Algebra is the Lie bracket so R must preserve the Lie Bracket so:

$$R([X, Y]) = [R(X), R(Y)] \forall X, Y \in \mathbb{L}(G_L)$$

So it is a representation that preserves all the important structure of the Lie Algebra.

Given a representation R of $\mathbb{L}(G_L)$, we can find a representation D of G_L :

$$\Lambda = \text{Exp}(X), X \in \mathbb{L}(G_L), D(\Lambda) = \text{Exp}(R(X))$$

There is a slight wrinkle here which is the fact that several Lie Groups have the same Lie Algebra like $SU(2)$ and $SO(3)$, so this can't map back to give you all of the Lie Group. So you don't get a representation of the original group but rather the universal cover \tilde{G}_L of G_L . For example with $SO(3)$ has a larger group $\text{spin}(3,1)$ which is a double cover and is exactly a 2:1 map ρ :

$$\rho : \tilde{G}_L \rightarrow G_L$$

Key examples of a double cover is that $SO(3)$ is double covered by $SU(2)$ written: $SO(3) \leftarrow^{2:1} SU(2)$.

Consider proper Lorentz transformations with $G_L = SO(3, 1)$ so $\Lambda = \text{Exp}(\omega)$. For example:

$$\Lambda_{\mu\nu}^\mu = \delta_{\mu\nu}^\mu + \omega_\nu^\mu + \frac{1}{2}\omega_l^\mu \omega_\nu^l + \dots$$

So as $\Lambda \in G_L \implies \Lambda \eta \Lambda^T = \eta$ we get constraint on ω of $\omega^{\mu\nu} + \omega^{\mu\nu} = 0$. This can be shown to be necessary and sufficient. So the Lie Algebra of the Lorentz group is $\mathbb{L}(G_L) = \{\omega \in \text{Mat}_4(\mathbb{R}), \omega + \omega^T = 0\}$. This has 6 independent variables/generators. The natural basis is matrices where we

put a 1 in the (μ, ν) , -1 in (ν, μ) position and zeros everywhere else above the leading diagonal. We can write this very explicitly in terms of the metric tensor:

$$(M^{\rho\sigma})^{\mu\nu} = \eta^{\rho\mu}\eta^{\sigma\nu} - \eta^{\sigma\mu}\eta^{\rho\nu}$$

Above one set of indices labels the set of generators ρ, σ and then another set of indices μ, ν which give you the matrix entries in the generators. E.g.

$$(M^{01})^{\nu\mu} = \begin{pmatrix} 0 & 1 & 0 & 0 \\ -1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}$$

This is the generator of infinitesimal lorentz transformation corresponding to a boost along the x-axis. So a general element of the Lie Algebra lorentz group will look like:

$$\omega_\nu^\mu = \frac{1}{2}\Omega_{\rho\sigma}(M^{\rho\sigma})_\nu^\mu$$

with $\Omega_{\rho\sigma} = -\Omega_{\sigma\rho}$. We have 6 independent constants that appear here that are a way of parametrising the elements of the lorentz group. We need to check the brackets and as everything is linear it is enough to just check the generators:

$$[M^{\rho\sigma}, M^{\tau\nu}] = \eta^{\sigma\tau}M^{\rho\nu} - \eta^{\rho\tau}M^{\sigma\nu} + \eta^{\rho\nu}M^{\sigma\tau} - \eta^{\sigma\nu}M^{\rho\tau} \quad (14)$$

Finally can write G_L in terms of generators:

$$\Lambda = \exp\left(\frac{1}{2}\Omega_{\rho\sigma}M^{\rho\sigma}\right)$$

Clifford Algebra which is four objects $(\gamma^0, \gamma^1, \gamma^2, \gamma^3)$ that obey certain relations. The defining relation is $\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu}\mathbb{I}$. Dirac invented this algebra as a trick. If you can find a representation of this algebra you can then manipulate them into the spinor algebra.

Consider the pauli matrices they obey the nice anticommutator: $\{\sigma^i, \sigma^j\} = 2\delta^{ij}\mathbb{I}_2$. Consider the simple representation of the clifford algebra as:

$$\gamma^0 = \begin{pmatrix} 0 & \mathbb{I}_2 \\ \mathbb{I}_2 & 0 \end{pmatrix}, \gamma^i = \begin{pmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{pmatrix}$$

The trick is to define:

$$S^{\rho\sigma} = \frac{1}{4}[\gamma^\rho, \gamma^\sigma] = \frac{1}{2}\gamma^\rho\gamma^\sigma - \frac{1}{4}\{\gamma^\rho, \gamma^\sigma\} = \frac{1}{2}\gamma^\rho\gamma^\sigma - \frac{1}{2}\eta^{\rho\sigma}\mathbb{I}_2$$

Check that $[S^{\rho\sigma}, S^{\mu\nu}] = \eta^{\sigma\mu}S^{\rho\nu} - \eta^{\rho\mu}S^{\sigma\nu} + \eta^{\rho\nu}S^{\sigma\mu} - \eta^{\sigma\nu}S^{\rho\mu}$

17 Lecture 17

Each Λ is mapped to $S[\Lambda] \in Mat_4(\mathbb{C})$ to give the spinor representation:

$$S[\Lambda] = \exp\left(\frac{1}{2}\Omega_{\rho\sigma}S^{\rho\sigma}\right) \quad (15)$$

$$S^{\mu\nu} = \frac{1}{4}[\gamma^\mu, \gamma^\nu] \quad (16)$$

with γ defined by:

$$\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu}\mathbf{I}_4 \quad (17)$$

This representation is a representation of the covering group (which is larger than the lorentz group) of the lorentz group not the lorentz group itself. e.g. when we rotate about 2π we get a minus sign rather than identity. We can focus on rotations in space by setting our parameters to pick out just the spatial components ($\Omega_{00} = \Omega_{0i} = \Omega_{i0} = 0$)

$$\Omega_{ij} = -\epsilon_{ijk}\phi^k$$

This is a rotation about the vector ϕ by an angle specified by the magnitude of this vector. Exercise: for $\phi^1 = \phi^2 = 0$ show that $\Lambda = \exp(-\phi^3 M^{12})$ acts on 4-vectors as a rotation through an angle ϕ^3 about the z axis.

$$\begin{aligned} S^{ij} &= \frac{1}{4}[\gamma^i, \gamma^j] =_{i \neq j \text{ as antisymmetric}} \frac{1}{2}\gamma^i\gamma^j = \frac{1}{2}\begin{pmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{pmatrix}\begin{pmatrix} 0 & \sigma^j \\ -\sigma^j & 0 \end{pmatrix} \\ S^{ij} &= \frac{1}{2}\begin{pmatrix} -\sigma^i\sigma^j & 0 \\ 0 & -\sigma^i\sigma^j \end{pmatrix} = -\frac{i}{2}\epsilon^{ijk}\begin{pmatrix} \sigma^k & 0 \\ 0 & \sigma^k \end{pmatrix} \\ S[\Lambda] &= \exp\left(-\frac{1}{2}\epsilon_{ijk}\phi^k S^{ij}\right) = \begin{pmatrix} e^{\frac{i\phi\cdot\sigma}{2}} & 0 \\ 0 & e^{\frac{i\phi\cdot\sigma}{2}} \end{pmatrix} \end{aligned}$$

Specialise further to the case $\phi^1, \phi^2 = 0$ then

$$S[\phi^3] = S[\Lambda]|_{\phi_1=\phi_2=0} = \begin{pmatrix} e^{i\frac{\phi^3}{2}} & 0 & 0 & 0 \\ 0 & e^{-i\frac{\phi^3}{2}} & 0 & 0 \\ 0 & 0 & e^{i\frac{\phi^3}{2}} & 0 \\ 0 & 0 & 0 & e^{-i\frac{\phi^3}{2}} \end{pmatrix}$$

$$\Lambda[\phi^3 = 2\pi] = \mathbf{I}_4$$

but

$$S[2\pi] = -\mathbf{I}_4$$

So clearly a double cover. This is also clear that it is a good starting point for talking about fermions as they are antisymmetric under exchange.

17.1 Spinor Fields

We want to have fields that transform in representations of the Lorentz group. Will easily give us singlets for the scalar fields but also lead to particle fields. A field is defined as:

$$\begin{aligned}\psi &: \mathbb{R}^{3,1} \rightarrow \mathbb{C}^4 \\ x &\rightarrow \psi^\alpha(x) \in \mathbb{C}\end{aligned}$$

where $\alpha = \{1, 2, 3, 4\}$ are the spinor indices. No significance to writing a spinor index up or down. Transformation property under a lorentz transformation ($\Lambda \in G_L$):

$$\psi^\alpha(x) \rightarrow \psi^{\alpha'}(x) = S[\Lambda]^\alpha_\beta \psi^\beta(\Lambda^{-1} \cdot x)$$

$$\psi(x) \rightarrow^\Lambda S[\Lambda] \cdot \psi(\Lambda^{-1} \cdot x)$$

Key property of this is that under a 2π rotation that leaves x invariant will change the field:

$$\psi(x) \rightarrow -\psi(x)$$

The spacetime derivative transforms as (follows from before using chain rule):

$$\partial_\mu \psi(x) \rightarrow^\Lambda (\Lambda^{-1})^\nu_\mu S[\Lambda]^\alpha_\beta \partial_\nu \psi^\beta(\Lambda^{-1} \cdot x)$$

Conjugate field:

$$\psi^\dagger_\alpha(x) \rightarrow \psi^\dagger \cdot S[\Lambda]^\dagger$$

Note: S is not unitary:

$$S[\Lambda]^\dagger \neq S[\Lambda]^{-1} \forall \Lambda \in G_L$$

As the lorentz group is a non-compact lorentz group as the boosts are not bounded so the manifold is a hyperboloid and there is a theorem that says any simple non-compact lie group has no unitary representation (missed what he actually said here 36 minute lecture 16/11).

Conjugation

Deduce from definition of clifford algebra that:

$$(\gamma^0)^2 = \mathbf{I}^4$$

So eigenvalues of γ^0 are real and it is hermitian ($(\gamma^0)^\dagger = \gamma^0$)
Can also deduce that for $i \in \{1, 2, 3\}$:

$$(\gamma^i)^2 = -\mathbf{I}_4$$

So γ^i is anti-hermitian $((\gamma^i)^\dagger = -\gamma^i)$.

$$\gamma^0 \gamma^i \gamma^0 = \gamma^0 \{\gamma^i, \gamma^0\} - \gamma^0 \gamma^{0i} = (\gamma^i)^\dagger, \gamma^0 \gamma^0 \gamma^0 = (\gamma^0)^\dagger$$

So

$$(\gamma^\mu)^\dagger = \gamma^0 \gamma^\mu \gamma^0$$

Now we can look at the conjugate properties of the spinor generators of the representation:

$$(S^{\mu\nu})^\dagger = \frac{1}{4}([\gamma^\mu, \gamma^\nu])^\dagger = \frac{1}{4}([\gamma^\nu]^\dagger, [\gamma^\mu]^\dagger) = \frac{1}{4}[\gamma^0 \gamma^\nu \gamma^0, \gamma^0 \gamma^\mu \gamma^0] = \frac{1}{3} \gamma^0 [\gamma^\nu, \gamma^\mu] \gamma^0 = -\gamma^0 S^{\mu\nu} \gamma^0$$

Now lets consider:

$$(S[\Lambda])^\dagger = [\exp(\frac{1}{2} \Omega_{\rho\sigma} S^{\rho\sigma})]^\dagger = \exp(\frac{1}{2} \Omega_{\rho\sigma} (S^{\rho\sigma})^\dagger) = \exp(-\frac{1}{2} \Omega_{\rho\sigma} \gamma^0 S^{\rho\sigma} \gamma^0) = \gamma^0 \exp(-\frac{1}{2} \Omega_{\rho\sigma} S^{\rho\sigma}) \gamma^0$$

Using the fact the $\gamma^0)^2 = 1$ on the last equality. As the inverse of e^P is e^{-P} :

$$(S[\Lambda])^\dagger = \gamma^0 (S[\Lambda])^{-1} \gamma^0 \quad (18)$$

This will allow us to construct real scalar quantities out of these fields.

18 Example class 2

In the Heisenberg picture for the free scalar field:

$$\dot{\phi}(x) = i[H, \phi(x)] = \pi(x), \dot{\pi}(x) = i[H, \pi(x)] = \nabla^2 \phi(x) - m^2 \phi(x)$$

The reason the signs flip when we go into the Heisenberg picture is because $e^{i\mathbf{p} \cdot \mathbf{x}} e^{-iEt} = e^{-ip_\mu x^\mu}$ as $p_\mu x^\mu = p_0 x^0 - \mathbf{p} \cdot \mathbf{x}$ Tips for question 4:

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

$$L = q^k \frac{\partial}{\partial q^j} - q^j \frac{\partial}{\partial q^k}$$

If $f(q) = f(|q|)$ (is spherically symmetric, then $Lf = 0$. **Feynmann diagram notation**

Let Γ be a Feynmann diagram, and let $v(\Gamma)$ be the final value of the diagram (the sum of all the terms it represents in the perturbation series. So $\langle 0|S|0\rangle = \sum_{\Gamma} v(\Gamma)$. $|\Gamma|$ is the value of Γ after applying Feynmann rules: write down $-i\lambda$ at each vertex, write down a propagator at each (internal)

edge, multiply and integrate. $v(\Gamma) = \frac{1}{|Aut(\Gamma)|} |\Gamma|$ with $|Aut(\Gamma)|$ being the symmetry factor of Γ , as $Aut(\Gamma)$ is a map of the vertices to themselves that preserves the graph structure. Count how many ways you can connect the half edges (draw the vertices and half edges without connecting them and then count the number of ways you could connect them up and still generate the same diagram)

Every feynmann diagram only corresponds to one term of the expansion you get addition from having multiple possible feynmann diagram. However, it is actually representing a collection of multiple term e.g $\int \phi(x)\phi^*(x)\phi\phi^*$ and $\int \phi^*(x)\phi(x)\phi\phi^*$ are represented by the same diagram. Number of half edges is equal to the number of fields in the perturbation. In general for a disjoint union of connected diagrams:

$$\Gamma = \Gamma_1^{n_1} \Gamma_2^{n_2} \dots \Gamma_k^{n_k} \implies v(\Gamma) = \frac{1}{n_1! \dots n_k!} v(\Gamma_1)^{n_1} \dots v(\Gamma_k)^{n_k}$$

19 Lecture 18

Define Dirac Adjoint:

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

Under a Lorentz Transformation $\Lambda \in G_L$:

$$\bar{\psi}(x) \rightarrow \bar{\psi}'(x) = \psi^\dagger(\Lambda^{-1} \cdot x) S^\dagger(\Lambda) \gamma^0 = \psi^\dagger(\Lambda^{-1} \cdot x) \gamma^0 S(\Lambda)^{-1} \gamma^0 \gamma^0 = \bar{\psi}(\Lambda^{-1} \cdot x) S(\Lambda)^{-1}$$

Define $\Sigma(x) = \bar{\psi}(x) \psi(x)$. We claim $\Sigma(x)$ is a real scalar field.

Proof it is real: $\Sigma^*(x) = (\psi^\dagger(x) \gamma^0 \psi(x))^\dagger = \psi^\dagger(x) (\gamma^0)^\dagger \psi(x) = \Sigma(x)$

Proof it is scalar: $\Sigma(x) \rightarrow \Sigma'(x) = \bar{\psi}(\Lambda^{-1} \cdot x) S(\Lambda)^{-1} S(\Lambda) \psi(\Lambda^{-1} \cdot x) = \bar{\psi}(\Lambda^{-1} \cdot x) \psi(\Lambda^{-1} \cdot x) = \Sigma(\Lambda^{-1} \cdot x)$

Can also form a vector field: $V^\mu = \bar{\psi}(x) \gamma^\mu \psi(x)$ transforms as

$$V^\mu(x) \rightarrow \Lambda^\mu_\nu V^\nu(\Lambda^{-1} \cdot x) = \Lambda^\mu_{\nu} \bar{\psi}(\Lambda^{-1} \cdot x) \gamma^\nu \psi(\Lambda^{-1} \cdot x)$$

Proof of above:

$$V^\mu \rightarrow \bar{\psi}(\Lambda^{-1} \cdot x) \cdot S(\Lambda)^{-1} \cdot \gamma^\mu \cdot S(\Lambda) \cdot \psi(\Lambda^{-1} \cdot x)$$

Hence this is true if $S[\Lambda]^{-1} \gamma^\mu S[\Lambda] = \Lambda^\mu_\nu \gamma^\nu$. Two spinor acting on gamma matrix is equivalent to one lorentz transfromation. Exercise is to prove this at the lienarised level by expanding out the exponential expressions for these terms and it becomes equivalent to the identity: $[\gamma^\mu, S^{\rho\sigma}] = \eta^{\rho\mu} \gamma^\sigma - \eta^{\sigma\mu} \gamma^\rho$. It is a lot more work to prove the full identity.

19.1 Dirac Action

Action for spinor fields, needs to be lorentz invariant and real. We want the \mathfrak{L} to be a real (at least up to surface terms) scalar field.

$$\mathfrak{L}(x) = \mathfrak{L}_1 + \mathfrak{L}_2 = \bar{\psi}(x)i\gamma^\mu\partial_\mu\psi(x) - m\bar{\psi}(x)\psi(x)$$

We know that $\mathfrak{L}_2 = -m\bar{\psi}\psi$ so is a real scalar field.

$$L_1 \rightarrow^\Lambda \bar{\psi}(\Lambda^{-1}\cdot x)S(\Lambda)^{-1}i\gamma^\mu(\Lambda^{-1})^\nu_\mu S[\Lambda]\partial_\nu\psi(\Lambda^{-1}\cdot x) = \bar{\psi}(\Lambda^{-1}\cdot x)i\Lambda^\mu_\rho\gamma^\rho(\Lambda^{-1})^\nu_\mu\partial_\nu\psi(\Lambda^{-1}\cdot x)$$

So

$$L_1 \rightarrow^\Lambda i\bar{\psi}(\Lambda^{-1}\cdot x)\gamma^\mu\partial_\mu\psi(\Lambda^{-1}\cdot x)$$

So L_1 is a scalar field. now check it is real:

$$L_1 = i\psi^\dagger(x)\gamma^0i\gamma^\mu\psi(x)$$

$$L_1^* = -i_\mu\psi^\dagger(x)\gamma^\mu)^\dagger\gamma^0\psi(x) = -i\partial_\mu\psi^\dagger(x)\gamma^0\gamma^\mu\gamma^0\gamma^0\psi(x) = -i\partial_{mu}\bar{\psi}(x)\gamma^\mu\psi(x) = i\bar{\psi}(x)\gamma^\mu\partial_\mu\psi(x)$$

Dirac action

$$S = \int d^4x (\bar{\psi}(x) \cdot (i\gamma^\mu\partial_\mu - m) \cdot \psi(x)) \quad (19)$$

Having got this action we can derive the field equation (EOM). For these purposes it is convient to think about the two independant fields are $\bar{\psi}$ and ψ rather than real and imaginary parts.

$$(i\gamma_\mu\partial_\mu - m) \cdot \psi(x) = 0 \quad (20)$$

Vary the action with respect to $\psi(x)$ gives $i(\partial_\mu\bar{\psi}) \cdot \gamma^\mu + m\bar{\psi} = 0$. **Notation** on 4 vector $\mathcal{V} = \gamma^\mu V_\mu$ so dirac equation becomes $(i\mathcal{D} - m) \cdot \psi = 0$.

Consider rotations (ϕ): $\Omega_{jk} = -\epsilon_{ijk}\phi^k$ or boosts $\Omega_{i0} = -\phi_i = \chi_i$. We can look at the spinor representation of these and they are all block diagonal:

$$S(\phi) = \begin{pmatrix} e^{i\frac{\phi\cdot\sigma}{2}} & 0 \\ 0 & e^{i\frac{\phi\cdot\sigma}{2}} \end{pmatrix}$$

$$S(\chi) = \begin{pmatrix} e^{\frac{\chi\cdot\sigma}{2}} & 0 \\ 0 & e^{-\frac{\chi\cdot\sigma}{2}} \end{pmatrix}$$

This reflects something non-trivial about the spinor representaion as every lorentz transformation is represented by a block diagonal matrix. So the spinor representation is a reducible representation.

20 Lecture 20

The rules for constructing the field theory are to basically find representations of the Lorentz groups that introduce fields that transform in these representations that produce particles that also live in the corresponding representation of the Lorentz group. So we are talking about the representation theory of the Lorentz group $G_L = SO(3,1)$. We actually we are considering representations of its cover $spin(3,1)$ we normally do this by exponentiating a Lie algebra representation which is much simpler as it is linear and there is no distinction between the covering group and the group. Once we get to classifying the representations of the Lie Algebra there is a useful trick of passing to the complexification of the Lie Algebra, as if you project back from these onto the real Lie algebra they will give you representations. There are many different Lie Algebras that have the same complexification. Therefore:

$$\mathbb{L}(spin(3,1)) = \mathbb{L}(SO(3,1)), \mathbb{L}_{\mathbb{C}}(SO(3,1)) = \mathbb{L}_{\mathbb{C}}(SO(4)) = \mathbb{L}_{\mathbb{C}}(SU(2)) \oplus \mathbb{L}_{\mathbb{C}}(SU(2))$$

We know the finite dimensional representations of $\mathbb{L}_{\mathbb{C}}(SU(2))$ which are classified by $j = 0, \frac{1}{2}, \dots$ and $-j \leq j_z \leq j$. So the representation of spin R_j has dimension $2j+1$. So you can classify the representations of the Lorentz group you can classify by combining these so they are labeled by two spins (j_+, j_-) so the dimension $R_{(j_+, j_-)} = (2j_+ + 1, 2j_- + 1)$. So you can see these dimensions will grow very rapidly with spin.

| | j_+ | j_- | D |
|---------------------------|---------------|---------------|---|
| scalar | 0 | 0 | 1 |
| Left handed Weyl spinors | $\frac{1}{2}$ | 0 | 2 |
| Right handed Weyl spinors | 0 | $\frac{1}{2}$ | 2 |
| 4 vector | $\frac{1}{2}$ | $\frac{1}{2}$ | 4 |

The Dirac spinors term is a combination of left and right handed Weyl spinors $= (\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$.

20.1 Solutions of Dirac Equation

$$(i\gamma^\mu \partial_\mu - m\mathbf{I}_4)\phi(x) = 0$$

act with $(-i\gamma^{\mu\nu}\partial_\nu - m\mathbf{I}_4)$:

$$(\gamma^\mu \gamma^\nu \partial_\nu \partial_\mu + m^2 \mathbf{I}_4)\phi(x) = 0$$

As $\frac{1}{2}\{\gamma^\mu, \gamma^\nu\} = \eta^{\mu\nu} \mathbf{I}_4$ $(\partial_\mu \partial^\mu + m^2)\phi(x) = 0$ So each of the components of ϕ always obeys the Klein Gordon equation, and we know these have wave like solutions with on shell solution:

$$\phi \sim e^{\pm i p \cdot x}$$

So $p^\mu p_\mu = m^2$, $p^0 = E_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$ so we will search for the positive frequency solutions using the ansatz

$$\psi_\alpha(x) = u_\alpha(p) e^{-ip \cdot x}$$

plugging into Dirac:

$$(p_\mu \gamma^\mu - m \mathbf{I}_4) u(p) = 0$$

Use the chiral representation for γ then:

$$\begin{pmatrix} -m \mathbf{I}_2 & p_\mu \sigma^\mu \\ p_\mu \bar{\sigma}^\mu & -m \mathbf{I}_2 \end{pmatrix} u = 0$$

with $\sigma^\mu = (I_2, \sigma^i)$ and $\bar{\sigma}^\mu = (I_2, -\sigma^i)$. Show that the general solution is:

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

where ξ is an arbitrary 2 component complex vector. As you can always diagonalise a hermitian matrix so the sqrt of H is defined as $\sqrt{H} = U^{-1} \sqrt{\Lambda} U$ for U s.t. $H = U^{-1} \Lambda U$ and $\sqrt{\Lambda} = \text{diag}\{\sqrt{\lambda_1}, \sqrt{\lambda_2}\}$.

Choose basis vectors ξ^s with $s = 1, 2$ (spin component) and each one is a 2 component complex vector with

$$(\xi^r)^\dagger \xi^s = \delta^{rs}$$

E.g. Could choose $\xi^1 = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$ $\xi^2 = 0$

1 So the basis of solutions are of the form:

$$\psi^{(r)}(x) = u^{(r)}(p) e^{-ip \cdot x}$$

with $u^{(r)} = \begin{pmatrix} \sqrt{p_\mu \sigma^\mu} \xi^r \\ \sqrt{p_\mu \bar{\sigma}^\mu} \xi^r \end{pmatrix}$ Similarly we have "negative" frequency solutions:

$$\psi^{(r)}(x) = v^{(r)}(p) e^{ip \cdot x}$$

with basis η^s giving:

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

This gives a total of four linearly independent solutions. We are going to need quite a few identities

$u^{r\dagger}(p)u^s(p) = 2p_0\delta^{rs}\bar{u}^r(p)u^s(p) = 2m\delta^{rs}v^{r\dagger}(p)v^s(p) = 2p_0\delta^{rs}\bar{v}^r(p)v^s(p) = -2m\delta^{rs}\bar{u}^s(p)v^r(p) = 0$ *for* $p' = (p_0, -\mathbf{p})$. Instead contracting the spin indices

$$\sum_{s=1}^2 u_{\alpha}^s(p)\bar{u}_{\beta}^s(p) = (\not{p} + m)_{\alpha\beta}$$

$$\sum_{s=1}^2 v_{\alpha}^s(p)\bar{v}_{\beta}^s(p) = (\not{p} - m)_{\alpha\beta}$$

20.2 Quantization of spinor field

$$L = i\bar{\psi}(x)\gamma_{\mu}^{\mu}\psi(x) = m\bar{\phi}(x)\psi(x)$$

Symmetry of dirac action:

$$\psi(x) \rightarrow e^{-i\alpha}\psi(x)$$

$$\bar{\psi}(x) \rightarrow e^{i\alpha}\bar{\psi}(x)$$

Conserved current:

$$j_{\nu}^{\mu} = \bar{\psi}\gamma^{\mu}\psi = V^{\mu}$$

$$\partial_{\mu}j_{\nu}^{\mu}$$

Conserved charge:

$$Q = \int d^3x j_{\nu}^0 = \int d^3x \bar{\psi}\gamma^0\psi = \int \psi^{\dagger}\psi$$

21 Lecture 21

21.1 Hamiltonian Formulation

Conjugate momentum of $\psi(x)$:

$$\pi(x) = \frac{\partial \mathcal{L}}{\partial \dot{\psi}(x)} = i\bar{\psi}(x)\gamma^0 = i\gamma^{\dagger}(x)$$

Hamiltonian:

$$H = \int d^3x \mathcal{H}$$

$$\mathcal{H}(x) = \pi(x)\dot{\psi}(x) - \mathcal{L}(x) = \bar{\psi}(x)(-i\gamma^i\partial_i + m)\psi(x)$$

Expand the field in a complete set of solutions of the Dirac equation, these are labelled by p_μ , $s = 1, 2$ (the momentum is on shell which means $p^\mu p_\mu = m^2$).

$$\psi(x) = u^s(p)e^{-ip \cdot x}, v^s(p)e^{ip \cdot x}$$

$$\psi_\alpha(x) = \sum_{s=1}^2 \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_p}} (b_p^s u_\alpha^s(p)e^{-ip \cdot x} + c_p^{s\dagger} v_\alpha^s(p)e^{ip \cdot x})$$

21.1.1 Spin-Statistics Theorem

Correlates in a relativistic field theory how we quantise a field with its spin. In particular it says that if we have particles with half integer spin we must quantise these particles as fermions. If we don't do this then the energy of our system will be unbounded below.

Quantization in the Schrodinger picture (at fixed time): $\psi(x) = \psi(\mathbf{x}, t) \rightarrow \psi(\hat{x})$ and $\psi^\dagger(x) = \psi^\dagger(\mathbf{x}, t) \rightarrow \psi^\dagger(\hat{x})$. Impose canonical anti-commutator relations:

$$\{\hat{\psi}_\alpha(\mathbf{x}), \hat{\psi}_\beta^\dagger(\mathbf{y})\} = \delta_{\alpha\beta} \delta^{(3)}(\mathbf{x} - \mathbf{y})$$

$$\{\hat{\psi}_\alpha(\mathbf{x}), \hat{\psi}_\beta(\mathbf{y})\} = \{\hat{\psi}_\alpha^\dagger(\mathbf{x}), \hat{\psi}_\beta^\dagger(\mathbf{y})\} = 0,$$

Mode expansion:

$$\hat{\psi}_\alpha(\mathbf{x}) = \sum_{s=1}^2 \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_p}} (\hat{b}_p^s u_\alpha^s(p)e^{ip \cdot x} + \hat{c}_p^{s\dagger} v_\alpha^s(p)e^{-ip \cdot x})$$

This implies that all the anticommutators vanish except:

$$\{b_p^s, b_q^{r\dagger}\} = \{c_p^s, c_q^{r\dagger}\} = (2\pi)^3 \delta^{rs} \delta^{(3)}(\mathbf{p} - \mathbf{q})$$

If we consider this in the context of the Fermionic Simple Harmonic Oscillator then:

We introduce \hat{b} and \hat{b}^\dagger but rather than posing their commutator relation we posit an anticommutator:

$$\{\hat{b}, \hat{b}^\dagger\} = 1, \{\hat{b}, \hat{b}\} = \{\hat{b}^\dagger, \hat{b}^\dagger\} = 0$$

This implies that $\hat{b}^2 = \hat{b}^{\dagger 2} = 0$ so they are nilpotent operators. We can try and define a Hamiltonian for our fermions:

$$H = \frac{1}{2} \omega (\hat{b}^\dagger \hat{b} - \hat{b} \hat{b}^\dagger) = \omega (\hat{N}_F - \frac{1}{2})$$

with $N_F = \hat{b}^\dagger \hat{b}$. With respect to this operator \hat{b} and \hat{b}^\dagger are raising and lower operators. Can check that:

$$[\hat{H}_F, \hat{b}] = -\omega \hat{b}, [\hat{H}_F, \hat{b}^\dagger] = \omega \hat{b}^\dagger$$

Ground state $|0\rangle$ has $\hat{H}_F |0\rangle = -\frac{1}{2}\omega |0\rangle$. Then there is also a $|1\rangle = \hat{b}^\dagger |0\rangle = \frac{1}{2}\omega 1$ but no more states as $(\hat{b}^\dagger)^2 |0\rangle = 0$. This means you can only fill a fermionic system once so this imposes the pauli exclusion principle.

Starting with our hamiltonian for a free field after normal ordering we get:

$$H = \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} \sum_{s=1}^2 (\hat{b}_{\mathbf{p}}^{s\dagger} \hat{b}_{\mathbf{p}}^s - \hat{c}_{\mathbf{p}}^s \hat{c}_{\mathbf{p}}^{s\dagger}) = \int \frac{d^3p}{(2\pi)^3} E_{\mathbf{p}} (\sum_{s=1}^2 (\hat{b}_{\mathbf{p}}^{s\dagger} \hat{b}_{\mathbf{p}}^s + \hat{c}_{\mathbf{p}}^{s\dagger} \hat{c}_{\mathbf{p}}^s) - (2\pi)^3 \delta^{(3)}(0))$$

So we can clearly see we have a hamiltonian that is just made up of number operators for all the possible momentum and particles so the field is again just an infinite number of harmonic oscillators. If we had used commutators instead of anticommutators we would have got a minus sign between the bs and cs, which would mean we could have an arbitrarily negative energy by occupying the c states. So have to use fermion quantisation description rather than the bosonic (anticommutator means they are antisymmetric under exchange whereas commutator means they are symmetric under exchange).

$$\hat{Q} = \int \frac{d^3p}{(2\pi)^3} \sum_{s=1}^2 (\hat{b}_{\mathbf{p}}^{s\dagger} \hat{b}_{\mathbf{p}}^s - \hat{c}_{\mathbf{p}}^{s\dagger} \hat{c}_{\mathbf{p}}^s)$$

So b particles and c particles have opposite charge.

Particle Spectrum:

$$\hat{b}_{\mathbf{p}}^r |0\rangle = \hat{c}_{\mathbf{p}}^r |0\rangle = 0, r = 1, 2 \forall \mathbf{p} \in \mathbb{R}^3$$

1-particle state:

$$\hat{b}_{\mathbf{p}}^{s\dagger} = |p, s\rangle$$

As b have positive charge $Q = +1$ and energy $E_{\mathbf{p}} = \sqrt{|\mathbf{p}|^2 + m^2}$ but we also have the spin index which can take two values (as appropriate for a spin 1/2 particle). These states are associated with the two different spinor indices which are then associated with our arbitrary choice of ξ basis so we could pick zbasis? (more to be done to show that these represent the spin of these particles need to act on it with angular momentum etc.).

1-anti particle state:

$$\hat{c}_{\mathbf{p}}^{s\dagger} = |p, s\rangle$$

so $Q = -1$ and $E = E_{\mathbf{p}}$. The fact we only get four states rather than eight states is because the hermitian conjugate is also the canonical conjugate of the field (lol not sure what this means).

Multi-particle states:

$$|p, r; q, s\rangle = \hat{b}_{\mathbf{p}}^{r\dagger} \hat{b}_{\mathbf{q}}^{s\dagger} |0\rangle = -|q, s; p, r\rangle$$

simply because $\{b_{\mathbf{p}}^{r\dagger}, b_{\mathbf{q}}^{s\dagger}\} = 0$

22 Lecture 22

22.1 Interacting Fermions

$$\mathfrak{L} = \mathfrak{L}_0 + \mathfrak{L}_{int}$$

$$\mathfrak{L}_0 = \bar{\psi}(i\gamma^\mu \partial_\mu - m)\psi$$

Dimensional analysis: $[\mathfrak{L}] = 4$, $[\partial] = +1 \implies [\psi] = [\bar{\psi}] = \frac{3}{2}$. Consider the simplest possible self-coupling that is consistent with lorentz invariance (needs to be powers of $\psi\bar{\psi}$):

$$\mathfrak{L}_{int} = -g_F(\psi\bar{\psi})^2$$

as $[g_F] = -2$ so this is an irrelevant coupling (leads to non-renormalisable field theory). So we need to consider fermions coupling to other fields e.g. Yukawa theory or coupling to bosons:

22.2 Yukawa theory

Real scalar field $\phi(x)$:

$$\mathfrak{L} = \mathfrak{L}_0 + \frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi) - \frac{\mu^2}{2}\phi^2 - \lambda\phi\psi\bar{\psi}$$

$[\lambda] = 0$ so this is a marginal interaction.

22.2.1 Calculate scattering amplitudes

Dyson's formula which will give rise to Time ordered products of the field operators, and then we will need to contract them with a version of wicks theorem. In doing so you get an answer in terms of a set of Feynman rules. So the rest of the course we will fill in these theories for Yukawa theory and QED (which uses gauge fields).

Feynmann Propagator

$$S_F(x - y)_{\alpha\beta} = \langle 0 | T[\hat{\psi}_\alpha(x) \hat{\bar{\psi}}_\beta(y)] | 0 \rangle$$

Like in the scalar case we are thinking about the Heisenberg picture, so the fields above are the Heisenberg picture field operators. We need to worry about the order that fields appear in the product in all cases as they have a vanishing anticommutator. So:

$$T[\hat{\psi}_\alpha(x) \hat{\bar{\psi}}_\beta(y)] = \begin{cases} \hat{\psi}_\alpha(x) \hat{\bar{\psi}}_\beta(y) & x^0 > y^0 \\ -\hat{\bar{\psi}}_\beta(y) \hat{\psi}_\alpha(x) & x^0 < y^0 \end{cases}$$

In the case where $x^0 = y^0$ they still anticommute so in order for the expression to be unambiguous it needs this minus sign. Lets start with the first ordering above ($x^0 > y^0$) call it M_+ . Consider the mode expansion:

$$\hat{\psi}_\alpha(x) = \sum_{s=1}^2 \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}}} (\hat{b}_{\mathbf{p}}^s u_\alpha^s(p) e^{-ip \cdot x} + \hat{c}_{\mathbf{p}}^{s\dagger} v_\alpha^s(p) e^{ip \cdot x})$$

$$\hat{\bar{\psi}}_\beta(x) = (\hat{\psi}^\dagger(x) \gamma^0)_\beta$$

So

$$\hat{\bar{\psi}}_\beta(x) = \sum_{s=1}^2 \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}}} (\hat{b}_{\mathbf{p}}^{s\dagger} \bar{u}_\beta^s(p) e^{ip \cdot x} + \hat{c}_{\mathbf{p}}^s \bar{v}_\beta^s(p) e^{-ip \cdot x})$$

So

$$M_+ = |0\rangle \hat{\psi}_\alpha(x) \hat{\bar{\psi}}_\beta(y) \langle 0| = \sum_{r,s=1}^2 \int \frac{d^3p}{(2\pi)^3} \int \frac{d^3q}{(2\pi)^3} \frac{1}{\sqrt{2E_{\mathbf{p}}2E_{\mathbf{q}}}} (u_\alpha^s(p) \bar{u}_\beta^r(q) e^{-i(p \cdot x - q \cdot y)} \langle 0 | \hat{b}_{\mathbf{p}}^r \hat{b}_{\mathbf{q}}^{s\dagger} | 0 \rangle$$

We can replace $\hat{b}_{\mathbf{p}}^r \hat{b}_{\mathbf{q}}^{s\dagger}$ with $\{\hat{b}_{\mathbf{p}}^r, \hat{b}_{\mathbf{q}}^{s\dagger}\} = (2\pi)^3 \delta^{rs} \delta^{(3)}(\mathbf{p} - \mathbf{q})$ so contract out the p and q:

$$M_+ = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2E_{\mathbf{p}}} e^{-ip \cdot (x-y)} \sum_{s=1}^2 u_\alpha^s(p) \langle u |_\beta^s(p)$$

and as $\sum_{s=1}^2 u_\alpha^s(p) \langle u|_\beta^s(p) = (\not{p} + m\mathbf{I}_4)_{\alpha\beta}$:

$$M_+ = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2E_p} e^{-ip \cdot (x-y)} (\not{p} + m\mathbf{I}_4)_{\alpha\beta}$$

Similarly:

$$M_- = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2E_p} e^{ip \cdot (x-y)} (\not{p} - m\mathbf{I}_4)_{\alpha\beta}$$

and we have:

$$S_F(x-y)_{\alpha\beta} = \begin{cases} M_+ & x^0 > y^0 \\ -M_- & x^0 < y^0 \end{cases}$$

Use the same technique as for the scalar case to turn this into a contour integral that picks up the minus signs by depending on the plane you close it in:

$$S_F(x-y)_{\alpha\beta} = \int \frac{d^4p}{(2\pi)^4} i \frac{(\not{p} + m)_{\alpha\beta}}{p^2 - m^2 + i\epsilon} e^{-ip \cdot (x-y)}$$

Wicks theorem for fermions

Now the individual creation operators and annihilation operators do not commute with themselves so you also have to keep track of the ordering of the sets of the same operators:

$$T[\hat{\psi}_\alpha \hat{\bar{\psi}}_\beta(y)] =: \hat{\psi}_\alpha \hat{\bar{\psi}}_\beta(y) : + \hat{\bar{\psi}}_\alpha \hat{\psi}_\beta(y)$$

The operators still anticommute within the normal ordering! Take the vacuum expectation value:

$$\hat{\bar{\psi}}_\alpha \hat{\psi}_\beta(y) = \langle 0 | T[\hat{\bar{\psi}}_\alpha(x) \hat{\psi}_\beta(y)] | 0 \rangle = S_F(x-y)_{\alpha\beta}$$

The contraction between $\psi\psi$ and $\bar{\psi}\bar{\psi}$ would vanish. The fermionic operators anticommute inside normal ordering:

$$: \hat{\psi}_1 \hat{\psi}_2 : = - : \hat{\psi}_2 \hat{\psi}_1 :$$

Wicks Theorem

Let $\hat{\pi}$ be an arbitrary product $\hat{\psi}, \hat{\bar{\psi}}$ with $\hat{\psi}(x_1) = \hat{\psi}_1$ and let

$$\hat{\pi} = \hat{\psi}_1 \hat{\psi}_2 \hat{\bar{\psi}}_3 \hat{\bar{\psi}}_4$$

For any such product, the time ordered product is normal order of the sum over all the contractions:

$$T[\hat{\pi}] =: G[\hat{\pi}] :$$

22.3 Gauge Fields

4-vector fields associated with the electromagnetic fields, as the observed fields E and B can be expressed in terms of the 4-vector potential A^μ which is a 4-vector field meaning it transforms as:

$$A^\mu(x) \rightarrow \Lambda^\mu_\nu A^\nu(\Lambda^{-1}x)$$

This field is not in of its self a physical thing there are transformations you can perform on A that leave the physically observable things E and B unchanged which is what makes this a gauge field. There is the observable field strength tensor which contains \mathbf{E} and \mathbf{B} :

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu$$

This field strength tensor is invariant under gauge transformations of A^μ :

$$A^\mu \rightarrow A^\mu + \partial^\mu \lambda, : \mathbb{R}^{3,1} \rightarrow \mathbb{R}$$

We need to write an action for this field theory with Lagrangian density (which is the only thing it could be to produce linear EOM whilst being Lorentz invariant and gauge invariant):

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu}$$

Conjugate momenta:

$$\pi_\mu = \frac{\partial \mathcal{L}}{\partial \dot{A}_\mu} = \begin{cases} 0 & \mu = 0 \\ -F^{0i} = E^i & \mu = \{1, 2, 3\} \end{cases}$$

This won't work with our normal canonical commutation relations as the zeroth term is 0. This is because we need to treat the gauge transformation properly

Configurations of A^μ that are related by a gauge transformation are physically indistinguishable.

The configuration space is all configurations of the field at some fixed time, and it is covered by sets of gauge orbits. A is foliated by gauge orbits and we need to fix a gauge.

Lorentz gauge

$$\partial_\mu A^\mu = 0$$

This condition is lorentz invariant, and also doesn't completely fix the gauge. Now think about dynamics with this gauge, by modifying the Lagrangian:

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

This reduces to Maxwell when the lorentz gauge is imposed and now we will have conjugate momentum for all μ . We can further specify to the Feynmann gauge with $\epsilon = 1$ but you don't need to. It is a good exercise to leave ϵ unset and check that it makes no difference to the final equations.

$$\partial_\mu \left(\frac{\partial \mathcal{L}}{\partial (\partial_\mu A^\mu)} \right) = -\partial_\mu F^{\mu\nu} - \partial_\mu (\eta^{\mu\nu} (\partial_\rho A^\rho)) = -\partial_\mu \partial^\mu A^\nu + \partial_\mu \partial^\nu A^\mu - \partial^\nu \partial_\mu A^\mu = -\partial^\nu \partial_\mu A^\mu$$

So the canonical momenta is:

$$\pi^0 = -(\partial_\mu A^\mu), \pi^i = \partial^i A^0 - \dot{A}^i$$

Strategy (Gupta-Bleuler procedure) is to first quantise and then constrain to the lorentz gauge. Consider in the schrodinger picture:

$$A^\mu \rightarrow \hat{A}^\mu, \phi^\mu \rightarrow \hat{\pi}^\mu$$

with Non-vanishing commutators:

$$[\hat{A}^\mu(\mathbf{x}), \hat{\pi}^\mu(\mathbf{y})] = i\eta^{\mu\nu} \delta^{(3)}(\mathbf{x} - \mathbf{y})$$

We have a spin 1 field so we expect to be quantised in a bosonic way hence why we are using commutators.

Consider mode expansion of field theory:

$$\partial^\nu A^\mu = 0$$

So this has solutions of the form $\epsilon^\mu e^{\pm i p \cdot x}$ with. These are labeled by on shell 4-momentum ($p_\mu p^\mu = 0, p^0 \geq 0 \implies p^\mu = (E_{\mathbf{p}}, \mathbf{p})$ and the ϵ (the polarisation 4-vector).

Chose an orthonormal basis $\epsilon_\mu^{(\lambda)}(p)$ ($\lambda \in \{0, 1, 2, 3\}$) with $\eta^{\mu\nu} \epsilon_\mu^{(\lambda)}(p) \epsilon_\nu^{(\rho)}(p) = \eta^{\lambda\rho}$ this tells us that $\epsilon^{(0)}$ is timelike and $\epsilon^{(i)}$ is spacelike. can show completeness relation: $\eta_{\lambda\rho} \epsilon_\mu^{(\lambda)}(p) \epsilon_\nu^{(\rho)}(p) = \eta_{\mu\nu}$.

It is useful to choose $\epsilon^{(1)}, \epsilon^{(2)}$ to be transverse ($p^\mu \epsilon_\mu^{(i)} = 0$) but $p^\mu \epsilon_\mu^{(3)} \neq 0$. Now we can mode expand the field operator:

$$\hat{A}_\mu(x) = \int \frac{d^3 p}{(2\pi)^3} \frac{1}{\sqrt{2|\mathbf{p}|}} \sum_{\lambda=0}^3 \epsilon_\mu^{(\lambda)}(p) (\hat{a}_{\mathbf{p}} e^{i p \cdot x} + \hat{a}_{\mathbf{p}}^\dagger e^{-i p \cdot x})$$

$$\hat{\pi}_\mu(x) = i \int \frac{d^3p}{(2\pi)^3} \sqrt{\frac{|\mathbf{p}|}{2}} \sum_{\lambda=0}^3 \epsilon_\mu^{(\lambda)}(p) (\hat{a}_\mathbf{p} e^{i\mathbf{p}\cdot\mathbf{x}} - \hat{a}_\mathbf{p}^\dagger e^{-i\mathbf{p}\cdot\mathbf{x}})$$

Now should check that the field commutation relations imply the oscillator commutation relations:

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

This creates a hilbert space with a non-positive norm so is not positive semidefinite (this is because we aren't talking about a physical space, when we restrict to the physical subspace this will disappear). We can calculate the hamiltonian:

$$\hat{H} = \int \frac{d^3p}{(2\pi)^3} |\mathbf{p}| \left(\sum_{i=1}^3 \hat{a}_\mathbf{p}^{i\dagger} \hat{a}_\mathbf{p}^i - \hat{a}_\mathbf{p}^{0\dagger} \hat{a}_\mathbf{p}^0 \right)$$

So the vacuum states are defined such that:

$$\hat{a}_\mathbf{p} |0\rangle = 0$$

The "Big" "Hilbert" space:

$$|\mathbf{p}_1, \dots, \mathbf{p}_n, \lambda_1, \dots, \lambda_n\rangle = \hat{a}_{\mathbf{p}_1}^{\dagger\lambda_1} \dots \hat{a}_{\mathbf{p}_n}^{\dagger\lambda_n} |0\rangle$$

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Heisenberg field operator

$$A_\mu(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2|\mathbf{p}|}} \sum_{\lambda=0}^3 \epsilon_\mu^{(\lambda)}(p) (\hat{a}_\mathbf{p}^\lambda e^{-ip\cdot x} + \hat{a}_\mathbf{p}^{\lambda\dagger} e^{ip\cdot x})$$

$$[\hat{a}_\mathbf{p}^\lambda, \hat{a}_\mathbf{q}^\mu] = -\eta^{\lambda\mu} (2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q})$$

We can get particles which we call photons by acting in different components on the vacuum state as above. We know this is wrong as there should only be two degrees of freedom corresponding to a photon as the polarisation is aligned transverse in 2d whereas this gives us 4 degrees of freedom. There is also a mathematical problem as the big "Hilbert space" spanned by this states contain negative norm states and 0 norm states so it is not really a hilbert space.

$$\langle \mathbf{p}, 0 | \mathbf{q}, 0 \rangle = \langle 0 | \hat{a}_\mathbf{p}^0 \hat{a}_\mathbf{q}^{0\dagger} | 0 \rangle = -(2\pi)^3 \delta^{(3)}(\mathbf{p} - \mathbf{q})$$

These problems go away when we impose the lorentz gauge constraint which classically is:

$$\partial_\mu A^\mu = 0$$

To impose this quantumly we use the Gupta-Bleuler strategy. We define a physical subspace s.t.

$$\langle \Psi' | \partial_\mu \hat{A}(x) | \Psi \rangle = 0 \forall | \Psi' \rangle, | \Psi \rangle \in \mathcal{H}_p$$

$$\partial_\mu \hat{A}^\mu = \int \frac{d^3 p}{(2\pi)^3} \frac{1}{\sqrt{2|\mathbf{p}|}} \sum_{\lambda=0}^3 (-ip^\mu) \epsilon_\mu^{(\lambda)}(p) (\hat{a}_{\mathbf{p}}^\lambda e^{-ip \cdot x} - \hat{a}_{\mathbf{p}}^{\lambda\dagger} e^{ip \cdot x})$$

This won't constrain transverse polarisations for which $p^\mu \epsilon_\mu$ as it will vanish always. Choose a canonical lorentz frame with $p^\mu = (p, 0, 0, p)$, and orthonormal basis $\epsilon_\mu^{(\lambda)}(p) = \delta_\mu^\lambda$. The point is that the GB condition then allows only photons with polarisation obeying $p^\mu \epsilon_\mu = 0$ so we have transverse photons: $|\psi\rangle = |\mathbf{p}, 1, 2\rangle$ and also certain linearly combinations of unphysical photons: $|\phi\rangle = |\mathbf{p}, 0\rangle - |\mathbf{p}, 3\rangle$ ("time like photon and longitudinal photon"). So we haven't completely removed the extra degrees of freedom but this is good enough, as these states have zero norm and the expectation of the Hamiltonian is also zero. This means you can get rid of them by taking a quotient, by identifying states which differ by a state with zero norm $|\psi\rangle \sim |\psi'\rangle + |\phi\rangle$. The upshot after all of this is we can restrict our attention to the two transverse operations.

23.0.1 Photon propagator

$$\Delta_{\mu\nu}(x-y) = \langle 0 | T[\hat{A}_\mu(x) \hat{A}_\nu(y)] | 0 \rangle$$

Evaluate using the mode expansion in exactly the same way and we will get an integral representation of the form:

$$\int \frac{d^4 p}{(2\pi)^4} \frac{-i\eta_{\mu\nu}}{p^2 + i\epsilon} e^{-ip \cdot (x-y)}$$

23.0.2 Coupling to matter

Maxwell equation with source

$$\partial_\mu F^{\mu\nu} = J^\nu$$

$$J^\nu = (\rho(\mathbf{x}, t), \mathbf{J}(\mathbf{x}, t))$$

The charge conservation equation means we must have a conserved current:

$$\partial_\mu J^\mu = 0$$

Comes from the action:

$$S = \int d^4x \mathfrak{L}$$

$$\mathfrak{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - A_\mu J^\mu$$

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$$

F is invariant under the lorentz gauge transformation, and so is $A_\mu J^\mu$ as

$$A_\mu \rightarrow A_\mu + \partial_\mu \lambda \implies A_\mu J^\mu \rightarrow A_\mu J^\mu + \partial_\mu \lambda J^\mu = A_\mu J^\mu$$

Source is provided by the spinor field so write down the dirac lagrangian:

$$\mathfrak{L}_F = \bar{\psi}(x)(i\gamma^\mu \partial_\mu - m)\psi(x)$$

This has a U(1) symmetry that gives rise to a conserved current:

$$\psi(x) \rightarrow e^{-i\alpha}\psi(x), \bar{\psi}(x) \rightarrow e^{i\alpha}\bar{\psi}(x) \implies \partial_\mu J^\mu_\nu = 0, J^\mu_\nu = \bar{\psi}\gamma^\mu\psi$$

23.0.3 Quantum Electrodynamics

$$\mathfrak{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \bar{\psi}(i\gamma^\mu \partial_\mu - m)\psi - e(\bar{\psi}\gamma^\mu\psi)A_\mu$$

Check that this is gauge invariant:

It is obviously a lorentz scalar as this was already demonstrated for the dirac lagrangian and the FF term. We have shown that $\bar{\psi}\gamma^\mu\psi$ is a lorentz vector so contracts to a lorentz scalar with A_μ . Now we can keep the symmetry if we turn it into a local symmetry:

$$\psi(x) \rightarrow e^{-ie\alpha(x)}\psi(x), \bar{\psi}(x) \rightarrow e^{ie\alpha(x)}\bar{\psi}$$

Want to check this is equivalent to $A^\mu \rightarrow A^\mu + \partial^\mu \alpha(x)$.

Define a covariant derivative:

$$D_\mu\psi(x) = \partial_\mu\psi(x) + ieA_\mu(x)\psi(x)$$

$$D_\mu\psi(x) \rightarrow \partial_\mu(e^{-ie\alpha(x)}\psi(x)) + ie(A_\mu(x) + \partial_\mu\alpha(x))e^{-ie\alpha(x)}\psi(x)$$

$$D_\mu\psi(x) \rightarrow e^{-ie\alpha(x)}(\partial_\mu\psi(x) - ie\partial_\mu\alpha(x)\psi(x) + ieA_\mu\psi(x) + ie\partial_\mu\alpha(x)\psi(x)) = e^{-ie\alpha(x)}D_\mu\psi(x)$$

Now we write

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \bar{\psi}(i\gamma^\mu D_\mu - m)\psi$$

$$\mathcal{L}_{int} = -e\bar{\psi}\gamma^\mu\psi A_\mu$$

Dimensional analysis: $[\mathcal{L}] = 4$ $[\psi] = [\bar{\psi}] = \frac{1}{2}$ and $[A^\mu] = 1$ so $[e] = 0$. So e is a marginal coupling. The true expansion parameter $\alpha = \frac{e^2}{4\pi}$. If we identify e with the charge of the electron and restoring units by dividing through by $\hbar c$ and this comes out to $\alpha \approx \frac{1}{137}$. It is the fact that this number being small that allows perturbation theory to work as each term differs from the next by powers of this number.

24 Examples Class 3

To show that $Tr\gamma^\mu = 0$ we observe that $\{\gamma^\mu, \gamma^\nu\} = 0$ for $\mu \neq \nu$ so $\gamma^\nu\gamma^\mu(\gamma^\nu)^{-1} + \gamma^\nu = 0$ and then take trace of both sides. Also neat trick:

$$\gamma^\nu\gamma^\mu = \frac{1}{2}\{\gamma^\nu, \gamma^\mu\} + \frac{1}{2}[\gamma^\nu, \gamma^\mu]$$

Q5:

$$(p \cdot \sigma)(p \cdot \bar{\sigma}) = p^2 \neq 0$$

$$\bar{\psi}u\bar{b}^\dagger + \bar{v}c$$

I forgot the bars on the u and v

Fermionic sign in feynmann rules: if we consider $\langle 0|T(\psi_1\psi_2\psi_3\psi_4)|0\rangle = \langle 0|T(\psi_1\psi_2\psi_3\psi_4)|0\rangle + \langle 0|T(\psi_1\psi_2\psi_3\psi_4)|0\rangle + \langle 0|T(\psi_1\psi_2\psi_3\psi_4)|0\rangle$ to make sure that the contractions are between adjacent operators we need to count how many times we need to swap fermionic operators and this gives us the parity. No way of doing this from feynmann diagrams (as it depends on the order of the external operators the ones defining the final and initial state).

When looking for relative signs just need to think about how many times you need to swap positions to go from one contraction to another

The reason you cant contract out any more terms as you will either get a loop or one particle will pass throughout without interacting. Also in feynmann rules look at the back of q8 to see the rules for a fermionic paying particular attention to the u and v associated with ingoing and outgoing lines.