

# Standard Model

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”If you can’t explain it simply enough you don’t understand it well enough”

- Albert Einstein

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# 1 The Quark Model

- Long ago it was realized that the proton and the neutron have very similar masses and thus in regions where the Electromagnetic is weak compared to the strong force, there is an approximate symmetry between the neutron. With this, it was postulated that these two particles were two states of the same particle, *the nucleon*. With this in analogous to spin states, we can write the proton as  $|p\rangle = (1, 0)^T$  and the neutron as  $|n\rangle = (0, 1)^T$ . This lead to the idea of *Isospin*, as the proton and the neutron can be considered to form an isospin doublet, with total Isospin  $1/2$  and a third component of  $I_3 = \pm 1/2$ . Just like spin, our Lagrangian (if we ignore the electromagnetic terms) should be invariant under unitary transformations of these states (which will be  $\in U(2)$ ), meaning there is a conserved charge associate with this transformation. We can do this exact same procedure for the up and down quarks. The strong force treats all the quarks equally, and seeing as the up and down quark have approximately the same masses, we can treat them as spin states just like the nucleon. In this case the conserved quantity associated with this symmetry is known as *flavor*.

## 1.1 Isospin

- $U(2)$  has 4 degrees of freedom, and thus 4 generators, one of these can be chosen to be a scaling by a phase factor of the identity,  $e^{i\theta}\mathbb{1}$ , since overall phase factors of  $U(1)$  do not change our states, this can be ignored, leaving us with the 3 Pauli matrices  $\sigma^i$ , which are the generators of  $SU(2)$ , the main symmetry group here. From this we can proceed in the exact same manner as we do with spin, recognizing that these matrices form a non-Abelian Lie algebra based on their commutators, where we can define raising and lowering operators, to enable us to write down states analogous to  $|lm\rangle$  for spin. This quantity that is like spin is called *Isospin*. This is 3-vector and is defined as:

$$\mathbf{T} = \frac{1}{2}\boldsymbol{\sigma}$$

- The components obey the following commutations relations:

$$[T_i, T_j] = i\epsilon_{ijk}T_k$$

Where we have sum of the index  $k$ . The measurable quantity from this system is the total isospin,  $\mathbf{T}^2 = T_1^2 + T_2^2 + T_3^2$ . We will label states by their total isospin and the third component of isospin  $I_3$ , i.e.  $\phi(I, I_3)$ . The up quark is then  $|u\rangle = \phi(\frac{1}{2}, \frac{1}{2})$  and the down quark is  $|d\rangle = \phi(\frac{1}{2}, -\frac{1}{2})$ .

## 1.2 Anti-quark Doublet

- The above treatment of up and down quarks is called a *quark doublet*, which we write as:

$$q = \begin{pmatrix} u \\ d \end{pmatrix}$$

We would like to have the same treatment of anti-quarks. We know that the complex conjugate of any quark will give us the anti-quark (eg.  $u^* = \bar{u}$ ), but we don't want to write down something like  $\bar{q} = (\bar{u}, \bar{d})^T$  as then this will transform via  $U^*$  instead of  $U$  and will not follow the same symmetries. Instead we should find some combination that does transform via  $U$ . We write this as  $\bar{q} = S(\bar{u}, \bar{d})^T$  and then impose that  $SU^* = US$ . Solving this equation results in the matrix  $S = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}$ . Meaning the quark anti-state can be written as:

$$\bar{q} = \begin{pmatrix} -\bar{d} \\ \bar{u} \end{pmatrix}$$

### 1.3 Mesons

- Mesons are bound states of a quark and anti-quark pair, since quarks are spin half, this makes mesons bosons. Since Mesons are comprised of two quarks, we can think of adding their Isospin in exactly the same way we add the spin of two particles together. Since the quarks have isospin  $\frac{1}{2}$ , this will create the familiar  $\frac{1}{2} \otimes \frac{1}{2} = 1 \oplus 0$ . Meaning there will be four possible states, a triplet with  $I_3 = 1$  and a singlet with  $I_3 = 0$ . These state correspond to meson particles! Each state will correspond to more than one particle as we are not considering different combinations of spins (which affect the mass!). These particles are:

$$\begin{aligned}
 \phi(1, 1) &= -|u\bar{d}\rangle & &= |\pi^+\rangle, |\rho^+\rangle \\
 \phi(1, 0) &= \frac{1}{\sqrt{2}} (|u\bar{u}\rangle - |d\bar{d}\rangle) & &= |\pi^0\rangle, |\rho^0\rangle \\
 \phi(1, -1) &= |d\bar{u}\rangle & &= |\pi^-\rangle, |\rho^-\rangle \\
 \phi(0, 0) &= \frac{1}{\sqrt{2}} (|u\bar{u}\rangle + |d\bar{d}\rangle) & &= |\eta\rangle, |\omega\rangle
 \end{aligned}$$

### 1.4 $SU(3)$ Flavour

- The above described  $SU(2)$  symmetry is almost exact as the up and down quarks have almost the same mass. What we can also do is consider extending this symmetry to the strange quark. This makes it a  $SU(3)$  symmetry, as the quark doublet now becomes a quark triplet  $q = (u, d, s)^T$ , for which we will need a new basis of generators. The standard choice of generators are the *Gell Mann matrices*:

$$\begin{aligned}
 \lambda_1 &= \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda_2 &= \begin{pmatrix} 0 & -i & 0 \\ i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & \lambda_3 &= \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, & u &\leftrightarrow d \\
 \lambda_4 &= \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}, & \lambda_5 &= \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}, & & & u &\leftrightarrow s \\
 \lambda_6 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, & \lambda_7 &= \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}, & & & d &\leftrightarrow s \\
 \lambda_8 &= \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}, & & & & & \text{equal treatment of u,d}
 \end{aligned}$$

- This particular choice is normalized such that  $\text{tr}(\lambda_i \lambda_j) = 2\delta_{ij}$ . With this we can define the  $SU(3)$  Isospin in a similar manner to  $SU(2)$  by identifying:

$$\hat{T}_i = \frac{1}{2} \lambda_i$$

The total Isospin is then  $\sum_i T_i^2 = \frac{1}{4} \lambda_i^2 = \frac{4}{3} \mathbb{1}$ .  $SU(3)$  is different to  $SU(2)$  in the fact that it has two mutually commuting generators instead of 1. We can call once again the component along the direction of the third generator  $T_3$  the third component of the isospin  $I_3$  and identify states by this number, but we then also need to consider the component along the direction of  $\hat{T}_8$  as  $\hat{T}_8$  commutes with  $\hat{T}_3$ . The component along this direction is known as the *Hyper-charge* and is denoted with a  $Y$ . Strictly we define the hyper charge to be  $Y = \frac{1}{\sqrt{3}} \lambda_8$ .

- From the definitions of the Gell Mann matrices we can check the values of the  $I_3$  and  $Y$ , for the 3 quarks:

$$\begin{aligned} \hat{T}_3 u &= +\frac{1}{2}u & \hat{Y}u &= +\frac{1}{3}u \\ \hat{T}_3 d &= -\frac{1}{2}d & \hat{Y}d &= +\frac{1}{3}d \\ \hat{T}_3 s &= 0 & \hat{Y}s &= -\frac{2}{3}s \end{aligned}$$

We can plot  $Y$  and  $I_3$  for the three quarks and their anti-particles:

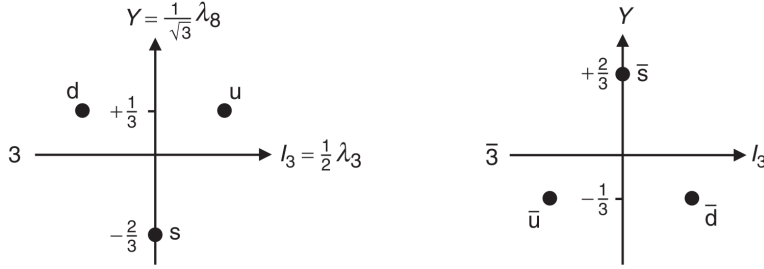


Figure 1: *Quark Isospin and Hyper-charge*

## 1.5 Light Mesons

- Since we have used  $\lambda_3$  and  $\lambda_8$  to form what is known as the Cartan sub-algebra, we can take the remaining  $\lambda_i$  and form raising and lowering operators. There will be three pairs, that which step respectively between the  $d \leftrightarrow u$ ,  $s \leftrightarrow u$  and  $d \leftrightarrow s$ :

$$\hat{T}_{\pm} = \frac{1}{2}(\lambda_1 \pm i\lambda_2) \quad (1.1)$$

$$\hat{V}_{\pm} = \frac{1}{2}(\lambda_4 \pm i\lambda_5) \quad (1.2)$$

$$\hat{U}_{\pm} = \frac{1}{2}(\lambda_6 \pm i\lambda_7) \quad (1.3)$$

We can use these to find all possible Mesons made out of these 3 quarks. This is done by identifying the extreme states (states with maximal  $I_3$  or  $Y$ ), then apply the raising and lowering operators to exhaust all other possible states. Since we are combining 3 possible quarks with 3 possible anti-quarks<sup>1</sup> then this is a case of  $3 \otimes \bar{3} = 8 \oplus 1$ <sup>2</sup>. This decomposition is into an octet and a singlet and takes the below visual form:

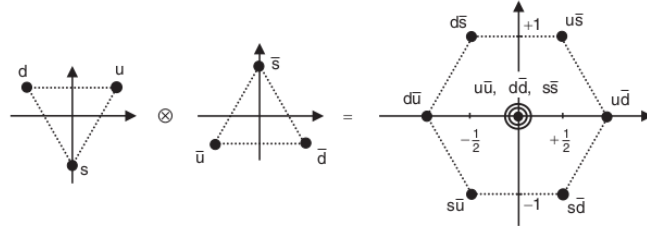


Figure 2: *Decomposition of quark-anti-quark combinations*

<sup>1</sup>Here we are only looking at 2 quark combinations here, i.e. Mesons

<sup>2</sup>The bar on the 3 just indicates that this is the anti-quark triplet

Where the singlet is plotted along with the two octet elements that have  $I_3 = Y = 0$ . These quark combinations are as we mentioned before Mesons! We now have generated more of these having considered the strange quark as well. We are not how-ever considering spin. Quarks are spin 1/2 particles meaning it is possible to form spin 0 or spin 1 combinations of them. It turns out that these different combinations affect the mass of the resulting bound system, meaning these are different particles. We thus have two sets of 9 particles:



Figure 3: All Mesons, graphed by Isospin and Hyper-charge

- In terms of the quarks these are for spin-0:

Pions:  $|\pi^+\rangle = |u\bar{d}\rangle, \quad |\pi^0\rangle = \frac{1}{\sqrt{2}}(|u\bar{u}\rangle - |d\bar{d}\rangle), \quad |\pi^-\rangle = |d\bar{u}\rangle$

Kaons:  $|K^+\rangle = |u\bar{s}\rangle, \quad |K^0\rangle = |d\bar{s}\rangle, \quad |\bar{K}^0\rangle = |s\bar{d}\rangle, \quad |K^-\rangle = |s\bar{u}\rangle$

Eta and Eta Prime:  $|\eta\rangle = \frac{1}{\sqrt{6}}(|u\bar{u}\rangle + |d\bar{d}\rangle - 2|s\bar{s}\rangle), \quad |\eta'\rangle = \frac{1}{\sqrt{3}}(|u\bar{u}\rangle + |d\bar{d}\rangle + |s\bar{s}\rangle)$

- And for the spin-1 mesons:

Rho Mesons:  $|\rho^+\rangle = |u\bar{d}\rangle, \quad |\rho^0\rangle = \frac{1}{\sqrt{2}}(|u\bar{u}\rangle - |d\bar{d}\rangle), \quad |\rho^-\rangle = |d\bar{u}\rangle$

Kaon\* Mesons:  $|K^{*+}\rangle = |u\bar{s}\rangle, \quad |K^{*0}\rangle = |d\bar{s}\rangle, \quad |\bar{K}^{*0}\rangle = |s\bar{d}\rangle, \quad |K^{*-}\rangle = |s\bar{u}\rangle$

Omega and Phi Mesons:  $|\omega\rangle = \frac{1}{\sqrt{2}}(|u\bar{u}\rangle + |d\bar{d}\rangle), \quad |\phi\rangle = |s\bar{s}\rangle$

- The Heavy Mesons are constructed from the bottom and charm quarks in a similar way.

## 1.6 Baryons

- Baryons are combinations of 3 quarks/ anti-quarks. This makes them fermions of spin  $\frac{1}{2}$  or  $\frac{3}{2}$ . This means we need to calculate  $3 \otimes 3 \otimes 3$ . It turns out that the calculation of  $3 \otimes 3$  is a little different to that of  $3 \otimes \bar{3}$ . Since we don't have anti-quarks we can't properly form a state that has  $I_3 = Y_3 = 0$ . This means the decomposition is  $3 \otimes 3 = 6 \oplus 3$ . This can be proved by the standard ladder operator calculations. We are then left with  $3 \otimes (6 \oplus 3)$  which breaks down into the standard  $3 \otimes \bar{3} = 8 \oplus 1$  and the new  $3 \otimes 6 = 10 \oplus 8$ . Overall this means the full decomposition is:

$$3 \otimes 3 \otimes 3 = 10 \oplus 8 \oplus 8 \oplus 1$$

- We can visualize this nicely with the following:

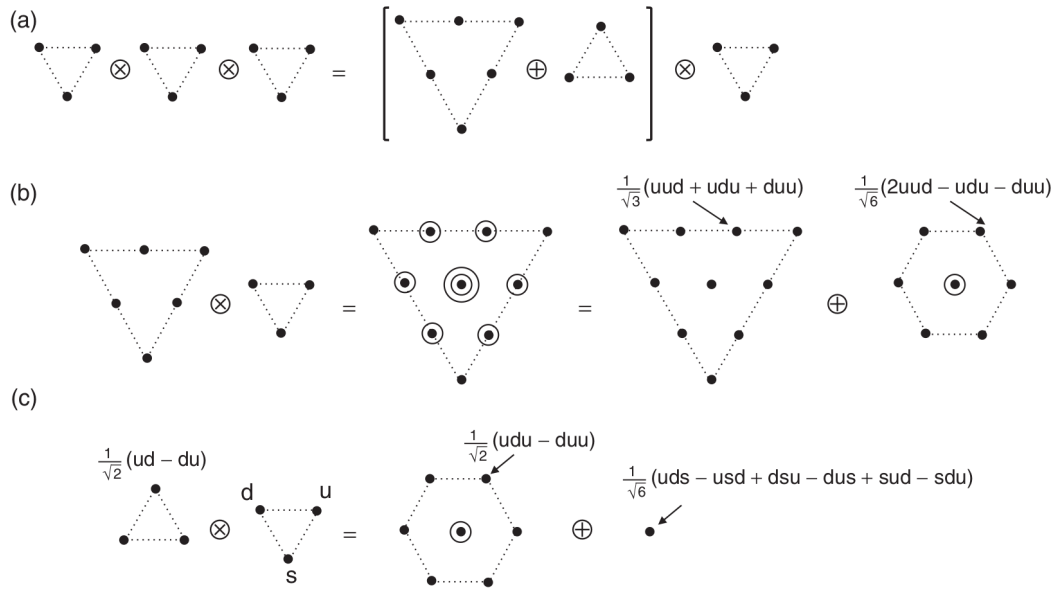


Figure 4: *All Mesons, graphed by Isospin and Hyper-charge*

- The quark composition of the individual Baryons is then:

Nucleons:  $|p\rangle = |uud\rangle, \quad |n\rangle = |udd\rangle$

$$\text{Delta Baryons:} \quad |\Delta^{++}\rangle = |uuu\rangle, \quad |\Delta^+\rangle = |uud\rangle, \quad |\Delta^0\rangle = |udd\rangle, \quad |\Delta^-\rangle = |ddd\rangle$$

Sigma Baryons:  $|\Sigma^+\rangle = |uus\rangle, \quad |\Sigma^0\rangle = |uds\rangle, \quad |\Sigma^-\rangle = |dds\rangle$

Xi Baryons:  $|\Xi^0\rangle = |uss\rangle, \quad |\Xi^-\rangle = |dss\rangle$

Omega Baryon:  $|\Omega^{-}\rangle = |sss\rangle$

## 1.7 Total Wavefunction

- There are many possible values of flavour, spin and colour (which we will encounter later) that a Baryon can have:

$$\psi = \phi_{\text{flavour}} \chi_{\text{spin}} \xi_{\text{colour}} \eta_{\text{space}}$$

However, not all of these states are valid as since baryons are fermions, their total wave-functions need to be anti-symmetric under exchange of any two quarks. We will see later that the colour wavefunction  $\xi_{\text{colour}}$  is totally anti-symmetric. We are discussing the quarks with  $l = 0$ , so zero spatial angular momentum, and since the spatial wavefunction transforms by  $(-1)^l$  under parity, then  $\eta_{\text{space}}$  is symmetric. This means that  $\phi_{\text{flavour}}\chi_{\text{spin}}$  must be symmetric. This allows us to determine the wave-function superpositions of the quarks in terms of their flavour and spin.



## 2 Scattering & Decay

- We will need to develop some mathematical tools for studying scattering of particles and Decay.

### 2.1 Decay

- Unstable particles have an equal probability of decaying at every instant in time. This means the probability of survival to time  $t$  for a particle  $A$  is governed by:

$$\begin{aligned}\frac{dP}{dt} &= -\Gamma_A P \\ \implies P &\propto e^{-\Gamma_A t}\end{aligned}$$

The constant  $\Gamma_A$  is known as the *total width*. If there are multiple decay processes  $A \rightarrow f$ , then the total width is  $\Gamma_A = \sum_f \Gamma(A \rightarrow f)$ . The probability ratio of the decay into an individual particle  $f$  is known as the *branching ratio*:

$$BR(A \rightarrow f) = \frac{\Gamma(A \rightarrow f)}{\Gamma_A}$$

### 2.2 Cross Section

- The main object to interest to us in calculations will be the cross section, as this is measurable. Our scattering setup to calculate the cross section will be as follows. We consider a beam of  $A$  particles of density  $n_A$  at a velocity  $v_A$  at a target  $B$ :



Figure 5: *Scattering setup*

- The rate of scattering is then given by:

$$\frac{\text{events}}{\text{sec}} = n_A v_A \sigma$$

Where  $\sigma$  is the cross section. It is useful for our calculations to consider all possible momenta the  $n$  particles created in the process could have. The probability of finding each given momentum configuration of the final particles is given by the *differential cross section*:  $d\sigma/(d^3p_1 \cdots d^3p_n)$ , which is related to the cross section by:

$$\sigma = \int d^3p_1 \cdots d^3p_n \frac{d\sigma}{d^3p_1 \cdots d^3p_n}$$

### 2.3 Master Formula

- We know from our study of Quantum Field Theory <sup>3</sup>, that the quantum mechanical transition matrix element is given by:

$$\langle 12 \cdots n | T | A(p_A) \rangle = \mathcal{M}(A \rightarrow 1 + 2 + \cdots n) (2\pi)^4 \delta(p_A - \sum_j p_j)$$

Where  $T$  is the time evolution operator. The factor of  $\mathcal{M}$  is called the *invariant Matrix element*.

- To find the total rate of decay we need to integrate over the total phase, i.e. all momenta that the particles could have subject to the constraint of momentum conservation. This will be written as:

$$\int d\Pi_n = \int \frac{d^3 p_1}{(2\pi)^3 2E_1} \cdots \frac{d^3 p_n}{(2\pi)^3 2E_n} (2\pi)^4 \delta(p_A - \sum_j p_j)$$

Where we also for convenience have included in this the factors of  $1/2E$ , so as to satisfy Lorentz invariance.

- Another formula we will then need from QFT is a version of Fermi's Golden rule, which tells us the particle width for a given transition to an  $n$ -particle final state  $f$ :

$$\Gamma(A \rightarrow f) = \frac{1}{2m_A} \int d\Pi_n |\mathcal{M}(A \rightarrow f)|^2$$

If the final products have spin we will need to sum over these spins. Similarly we need to do the same for the initial state, but often this state is unknown so we simply average over all possible spins instead. This expression is actually a simplification of the following more general expression for the cross section in terms of the invariant matrix element:

$$\sigma(A + B \rightarrow f) = \frac{1}{2E_A 2E_B |v_A - v_B|} \int d\Pi_n |\mathcal{M}(A \rightarrow f)|^2$$

### 2.4 2 Particle Collision

- We can now use the above equations to analyze the problem of a two particle collision. We will impossibly do this for the case where  $\mathcal{M}$  is constant, as this represents the case where each point in the phase space is equally likely. We will consider an initial state of momentum  $P$  and a final state of two particles with momentum  $\mathbf{p}_1$  and  $\mathbf{p}_2$ . In this scenario the phase space integral is just:

$$\int \Pi_2 = \int \frac{d^3 p_1}{(2\pi)^3 2E_1} \frac{d^3 p_2}{(2\pi)^3 2E_2} (2\pi)^4 \delta^{(4)}(P - p_1 - p_2) \quad (2.1)$$

If we work in Center of Mass (CoM) frame, where  $\mathbf{P} = 0$ , then the three spatial parts of the delta function enforce that:  $\mathbf{p}_1 = -\mathbf{p}_2$  and thus if  $P = (E_{\text{CM}}, 0) \implies p_1 = (E_1, \mathbf{p})$  &  $p_2 = (E_2, -\mathbf{p})$ , where the particles are on shell so  $E_1 = \sqrt{p^2 + m_1^2}$  and  $E_2 = \sqrt{p^2 + m_2^2}$ . With this 2.1 becomes:

$$\int \Pi_2 = \int \frac{d^3 p}{(2\pi)^3} \frac{2\pi}{2E_1 2E_2} \delta(E_{\text{CM}} - E_1 - E_2) \quad (2.2)$$

In spherical co-ords this volume element is just  $d^3 p = p^2 dp d\Omega$ . We can then deal with this integral by using the property of the delta function for a function with only one zero at  $x_0$ :

$$\begin{aligned} \delta(f(x)) &= \frac{1}{|f'(x_0)|} \delta(x - x_0) \\ \implies \delta(E_{\text{CM}} - E_1(p) - E_2(p)) &= \frac{\delta(p - \tilde{p})}{|dE_1/dp + dE_2/dp|} \end{aligned}$$

<sup>3</sup>See my notes on QFT here.

We can evaluate from the on-shell conditions that  $dE_1/dp = p/E_1$  and  $dE_2/dp = p/E_2$ , so we can write 2.2 as:

$$\int \Pi_2 = \int \frac{\tilde{p}^2 d\Omega}{16\pi^2 E_1 E_2} \frac{1}{|\tilde{p}/E_1 + \tilde{p}/E_2|} = \frac{1}{E_1 + E_2} \int \frac{\tilde{p} d\Omega}{16\pi^2}$$

$$= \frac{1}{8\pi} \left( \frac{2\tilde{p}}{E_{\text{CM}}} \right) \int \frac{d\Omega}{4\pi}$$

(2.3)

We can note that in the relativistic limit  $E \gg m_1, m_2 \implies E_{\text{CM}} = E_1 + E_2 \approx \tilde{p} + \tilde{p} = 2\tilde{p}$ , which makes the factor in brackets in 2.3 unity. This means 2.3 becomes constant in the relativistic limit.

## 3 Electron Positron Annihilation

- We would now like to discuss the annihilation of an electron  $e^-$  and its anti-particle the positron  $e^+$ . Firstly we will calculate the cross section for the process  $e^+e^- \rightarrow \mu^+\mu^-$  ( $\mu$  being muons) as this calculation will serve as a basis for many more complicated calculations. In doing this we will once again rely heavily on my QFT notes as we need to work with spinors for the fermions.

### 3.1 Feynman Diagrams

- Despite their pretty form, Feynman diagrams do not represent physical process that actually happen. Despite this we will still use them for the basis of figuring out which terms in the expansion of the interaction part of the scattering which (as per section 8.4 of my QFT notes) takes the form:

$${}_{\text{in}} \langle \mathbf{p}_1 \mathbf{p}_2 | \mathbf{k}_A \mathbf{k}_B \rangle_{\text{out}} \propto \lim_{t \rightarrow \infty (1-i\epsilon)} {}_0 \langle \mathbf{p}_1 \mathbf{p}_2 | T \left\{ \exp \left( -i \int_{-t}^t dt' H_I(t') \right) \right\} | \mathbf{k}_A \mathbf{k}_B \rangle_0 \quad (3.1)$$

We will expand in powers of the interaction parameter. The Lagrangian for QED (also per my QFT notes, see section 9.1) is given by<sup>4</sup>:

$$\mathcal{L}_{\text{QED}} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \bar{\psi} (i\gamma^\mu \partial_\mu - m) \psi - e \bar{\psi} \gamma^\mu \psi A_\mu \quad (3.2)$$

- We will be working, for simplicity, in the high energy regime, where all particles in the  $e^+e^- \rightarrow \mu^+\mu^-$  reaction can be treated as massless. Recall that the  $\psi$ 's here are Dirac fermions, which are comprised of two 2-component spinors known as Weyl spinors. These spinors are called left and right handed spinors and the  $\psi$ 's which are 4-component spinors can be written as the following<sup>5</sup>:

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

Note that for this calculation it is most useful to use the Chiral representation of the Dirac gamma matrices. This means<sup>6</sup>:

$$\gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{pmatrix} \implies \gamma^0 = \begin{pmatrix} 0 & \mathbb{I} \\ \mathbb{I} & 0 \end{pmatrix}, \quad \gamma^i = \begin{pmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{pmatrix}$$

- With this we can see that the QED Lagrangian 3.2 can be expanded in terms of the left and right component spinors:

$$\begin{aligned} \mathcal{L}_{\text{QED}} = & -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + i\psi_R^\dagger \sigma^\mu \partial_\mu \psi_R + i\psi_L^\dagger \bar{\sigma}^\mu \partial_\mu \psi_L - e\psi_R^\dagger \sigma^\mu \psi_R A_\mu - e\psi_L^\dagger \bar{\sigma}^\mu \psi_L A_\mu \\ & + m \left( \psi_R^\dagger \psi_L + \psi_L^\dagger \psi_R \right) \end{aligned}$$

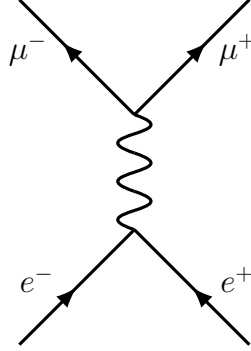
From this we can see that the last term, which is proportional to the mass, couples the left handed spinors to the right handed ones. In fact if we have  $m = 0$ , which is indeed the situation we will be considering, then the two left and right handed particles completely decouple and act as separate particles. This means in the high energy case a right handed electron cannot turn into a left handed one. Note that both species of particles also have their own interaction terms so they still interact amongst themselves. This isolation is essentially *Conservation of helicity*.

<sup>4</sup>Note here the Dirac adjoint which is the bar on some of the  $\psi$ 's is defined as  $\bar{\psi} \equiv \psi^\dagger \gamma^0$ .

<sup>5</sup>For a full discussion of Weyl spinors see section 4.6 of my QFT notes.

<sup>6</sup>Note here  $\bar{\sigma} = (\mathbb{I}, -\sigma)$

- We are now ready to consider the interactions. Recall that in the calculation of terms in the power series of 3.1 we contract ingoing and outgoing states with intermediate states. Since the interaction terms contain 3 fields, in order to have an even number of fields to contract we need to look at the quadratic terms of the interaction Hamiltonian  $H_I$ . This allows us to immediately write down the Feynman diagram that correspond to the first order interaction terms. We can reason that the following diagram is the only first order term as the fields available for contraction are  $\bar{\psi}, A, \psi, \bar{\psi}, A, \psi$ . Two of the  $\psi$ 's contract with the ingoing momenta and two with the outgoing in order to achieve  $e^+e^- \rightarrow \mu^+\mu^-$ , the final  $A$ 's must contract with themselves. This takes the following diagrammatic form:



### 3.1.1 Massless Fermions

- In the case of  $m = 0$  fermions The Dirac equation  $(i\gamma^\mu \partial_\mu - m)\psi = 0$  reduces to the following equations on the two spinors:

$$\begin{aligned} i\sigma^\mu \partial_\mu \psi_L &= 0 \\ i\bar{\sigma}^\mu \partial_\mu \psi_R &= 0 \end{aligned}$$

This means the solutions will take the following form (there are two types of solutions):

$$\psi_R = u_R(p)e^{-ip \cdot x}, \quad \psi_L = u_L(p)e^{-ip \cdot x}$$

Here the  $u_R(p)$  and  $u_L(p)$  are two component spinors that must satisfy their own versions of the Dirac equation. If we consider only plane waves propagating in the  $\hat{z}$  direction, so  $p = (E, 0, 0, p_z)$  then this equation is:

$$\begin{aligned} (E - p_z \sigma_3)u_R &= \begin{pmatrix} E - p_z & 0 \\ 0 & E + p_z \end{pmatrix} u_R = 0 \\ (E + p_z \sigma_3)u_L &= \begin{pmatrix} E + p_z & 0 \\ 0 & E - p_z \end{pmatrix} u_L = 0 \end{aligned}$$

Both of these have the same solutions  $\begin{pmatrix} 1 \\ 0 \end{pmatrix}$  and  $\begin{pmatrix} 0 \\ 1 \end{pmatrix}$ . For  $u_R$ ,  $\begin{pmatrix} 1 \\ 0 \end{pmatrix}$  has eigenvalue  $+E$  and  $\begin{pmatrix} 0 \\ 1 \end{pmatrix}$  has eigenvalue  $-E$  in the case of  $u_L$  it is the other way round. It is convenient (see section 4.8 of my QFT notes) to normalize these spinors by  $\sqrt{2E}$  and write them in the following:

$$\begin{aligned} u_R &= \sqrt{2E}\xi_+, & u_L &= \sqrt{2E}\xi_- \\ v_R &= \sqrt{2E}\xi_-, & v_L &= \sqrt{2E}\xi_+ \end{aligned}$$

Where we have change  $u \rightarrow v$  for the negative energy solutions.

- Big disclaimer,  $\xi_+$  and  $\xi_-$  are defined to be one of the two solution  $(1, 0)^T$  and  $(0, 1)^T$ , how-ever the choice of which one they are depends on the direction of motion.  $\xi_+$  is defined to be the spinor with spin up (that means positive eigenvalue) along the direction of motion and  $\xi_-$  as the spinor with spin down (negative eigenvalue). However, we need to be careful in annihilation processes as then anti-particles are travelling the opposite direction and thus which spinor  $(1, 0)^T$  or  $(0, 1)^T$  has positive eigenvalue can swap.

### 3.2 Matrix Element

- The setup for the scattering in the Center of Mass (CM) frame will be the following. The electron and the positron will be moving in the opposite directions along the  $\hat{z}$  axis with momenta  $p_-$  and  $p_+$ . This means  $p_- = (E, 0, 0, E)$  and  $p_+ = (E, 0, 0, -E)$ . The outgoing muon and anti-muon can be scattered at angles in the  $\hat{x} - \hat{z}$  plane. Hence their outgoing momenta are given by  $p'_- = (E, E \sin \theta, 0, E \cos \theta)$  and  $p'_+ = (E, -E \sin \theta, 0, -E \cos \theta)$ .
- We can then use Feynman rules for QED (see section 9.2 of my QFT notes). Let us first deal with the annihilation of the electron and positron vertex. The Feynman rules tell us to add to our Matrix element a  $-e\gamma^\mu$  for every vertex, note that in our case since we have split our Lagrangian into right and left handed spinors we instead add a  $-e\sigma^\mu$  for right handed spinors and  $-ie\bar{\sigma}^\mu$  for left handed spinors. We then have the contraction of the fields with the incoming momenta. One of the fields is a  $\psi^\dagger$ <sup>7</sup> this contracts with the incoming momenta of the positron to give us a  $v^\dagger(p_+)$  as per the Feynman rules. Similarly the  $\psi$  contracts with the incoming momenta of the electron to give us a  $u(p_-)$ . Mathematically this looks like the following. We will pick to do this for the case of a right handed electron and left handed positron<sup>8</sup>:

$$\begin{aligned} \langle 0 | -e\psi_R^\dagger \sigma^\mu \psi_R | e_R^-(p_-) e_L^+(p_+) \rangle &= \langle 0 | -e\overbrace{\psi_R^\dagger \sigma^\mu \psi_R}^{e_R^-(p_-) e_L^+(p_+)} | e_R^-(p_-) e_L^+(p_+) \rangle = -e v_L^\dagger(p_+) e^{ip_+ \cdot x} \sigma^\mu u_R(p_-) e^{-ip_- \cdot x} \\ &= \sqrt{2E}(0 \ 1)(1, \sigma) \sqrt{2E} \begin{pmatrix} 1 \\ 0 \end{pmatrix} e^{i(p_- - p_+) \cdot x} = 2E(0, 1, i, 0)^\mu e^{i(p_- - p_+) \cdot x} \end{aligned}$$

This vector we see here is nothing more then a polarization vector that we use to describe photons. We will call it  $\epsilon_+^\mu = \frac{1}{\sqrt{2}}(0, 1, i, 0)^\mu$ . Note that the left handed current  $-e\psi_L^\dagger \bar{\sigma}^\mu \psi_L$  when acted on this electron positron ket is zero as in the massless case the right handed electron does not couple to the left handed fermions.

- Next up is the propagator, which corresponds to the contraction of the two  $A_\mu$  fields. This propagator in momentum space is given simply by <sup>9</sup> (see section 4.8 of my QFT notes):

$$\langle 0 | A_\nu A_\mu | 0 \rangle = \overbrace{A_\nu A_\mu} = \frac{-i\eta_{\mu\nu}}{q^2}$$

- Finally we have the contraction of the fields with the final muon momenta. This takes the form, we will choose the muon to have :

$$\langle \mu_R^-(p_-) \mu_L^+(p_+) | (-e)\psi_R^\dagger \sigma^\nu \psi_R | 0 \rangle = -e \langle \mu_R^-(p_-) \mu_L^+(p_+) | \overbrace{\psi_R^\dagger \sigma^\nu \psi_R}^{e_R^-(p_-) e_L^+(p_+)} | 0 \rangle = -e$$

<sup>7</sup>Note this is not a  $\bar{\psi}$  here as we used the  $\gamma^0$  from the  $\bar{\psi}$  to make  $\gamma^\mu$  diagonal so that the vertex cleanly gave us  $\sigma^\mu$  and  $\bar{\sigma}^\mu$ .

<sup>8</sup>Note jsut as we mentioned earlier since the positron is going in the opposite direction  $v_L^\dagger(p_+) = (0 \ 1)^T$

<sup>9</sup>Note there is implicit time ordering operators on all quantities in this section.

