Chapter 2

STANDARD MODEL (AND BEYOND!)

The Standard Model (SM) is a set of theories that describe fundamental particle physics and the interactions of all known elementary particles, except gravity[1]. Kickstarted by Sheldon Glashow's [1] discovery of combining electromagnetic and weak interactions in 1961, it has evolved since then into its current form that we know today. Many precision analyses have been performed at many particle physics experiments such as AGS, E288, PETRA, UA1, D0, DONUT, and the experiments at the Large Hadron Collider (LHC) (see table 2.2). All of these experimnts have measured the cross-section for various processes and show good agreement to the predictions of the Standard Model. For example, fig. 2.1 describes the production cross-section measurements measured by the ATLAS detector compared to theoretical expectations for common decay processes. The data/theory ratio are shown on the right side of the figure and serve to show how successful the Standard Model has been. However, it had posed a few problems such as requiring spontaneous symmetry breaking in order to explain the heavy masses of the bosons that mediate the weak interactions (see section 2.1.2). The Higgs boson, discovered in 2012 [2], explained this missing piece.

Section 2.1 will provide background information about the Standard Model, the theories, and its particles. There are still many other puzzles of the Standard Model that need to be reconciled and will be discussed in section 2.2, thanks in large part due to [3, 4].

[1] As far as we know, gravity is too weak to play any significant role in ordinary particle processes.

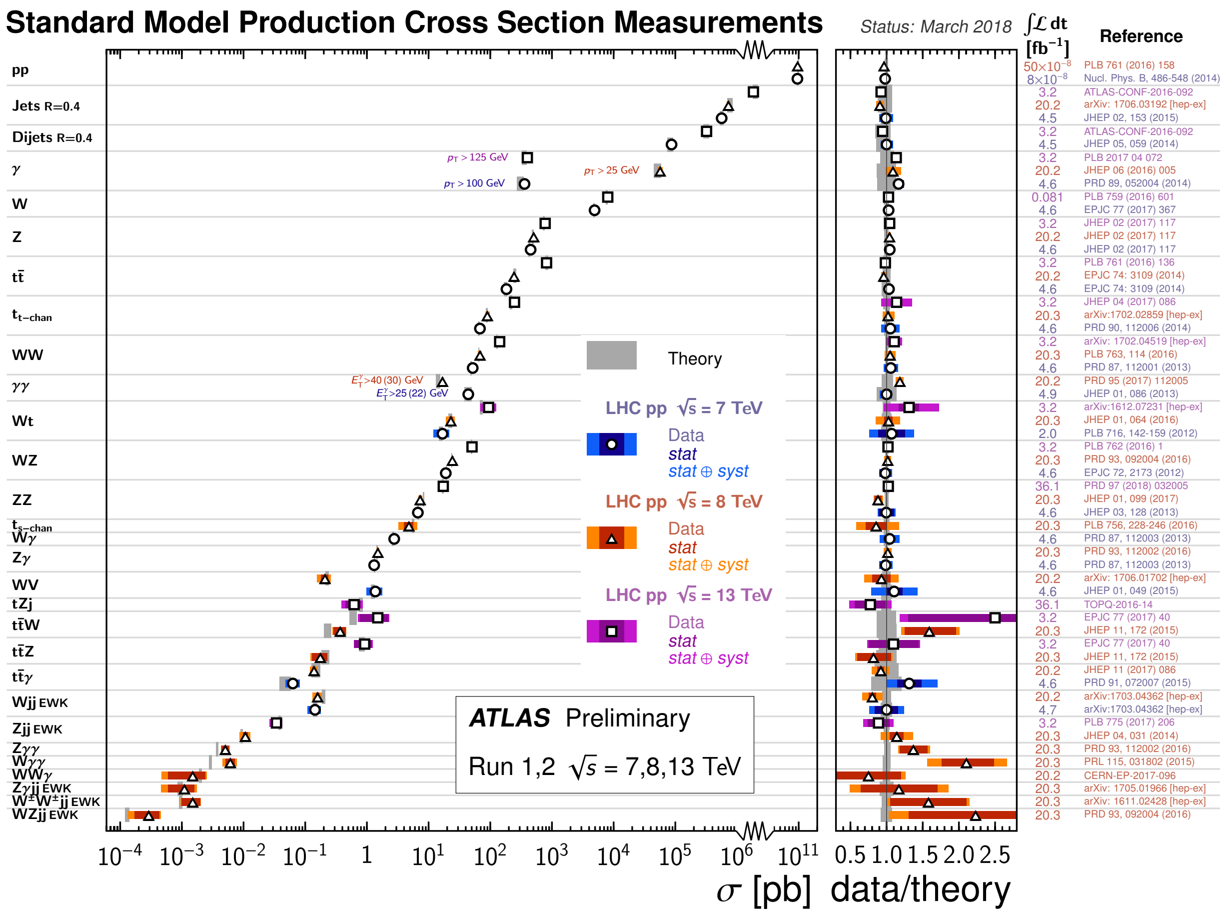


Figure 2.1: [5] Summary of several Standard Model total and fiducial production cross section measurements, corrected for leptonic branching fractions, compared to the corresponding theoretical expectations. All theoretical expectations were calculated at next-to-leading-order (NLO) or higher. The dark-color error bar represents the statistical uncertainty. The lighter-color error bar represents the full uncertainty, including systematics and luminosity uncertainties. The data/theory ratio, luminosity used and reference for each measurement are also shown. Uncertainties for the theoretical predictions are quoted from the original ATLAS papers. They were not always evaluated using the same prescriptions for PDFs and scales. The Wgamma and Zgamma theoretical cross-sections have non-perturbative corrections applied to the next-to-next-to-leading-order (NNLO) fixed order calculations [6]. Not all measurements are statistically significant yet.

2.1 The Standard Model

The SM is the most comprehensive quantum field theory of particle physics today. It encompasses a single, concise model made of up two theories: the Glashow-Weinberg-Salam theory of QED (section 2.1.2) which describes the electromagnetic and weak nuclear forces and QCD (section 2.1.3) which describes the strong nuclear force; with two classes of particles: fermions and bosons. These two theories form the symmetry group of the Standard Model [7, 8, 9, 10]

$$\underbrace{\mathrm{SU_C}(3)}_{\text{QCD}} \otimes \underbrace{\mathrm{SU_L}(2) \otimes \mathrm{U_Y}(1)}_{\text{QED}}. \tag{2.1}$$

$\mathrm{SU_C}(3)$ is denoted with a subscript to ensure it is not confused with the non-gauge theory flavor SU(3), $\mathrm{SU_L}(2)$ represents the weak gauge vectors in the theory, and $\mathrm{SU_Y}(1)$ denoting the gauge group of weak hypercharge. As the SM is a quantum field theory, the fundamental objects are quantum fields. These are:

- fermionic fields ψ^L, ψ^R (for left/right chirality),
- electroweak boson fields W_1, W_2, W_3, B,
- gluon field G,
- and the Higgs field ϕ.

For example, the massless electroweak boson fields are given mass due to the Higgs mechanism through mixing, to create physically observable particles[11]

$$Z = \cos\theta_W W_3 - \sin\theta_W B, \tag{2.2}$$

$$A = \sin\theta_W W_3 + \cos\theta_W B, \tag{2.3}$$

$$W^{\pm} = \frac{1}{\sqrt{2}}(W_1 \mp iW_2). \tag{2.4}$$

In eq. (2.2), θ_W is the Weinberg angle or weak mixing angle $\sin^2\theta_W = 0.2223(21)$ [12].

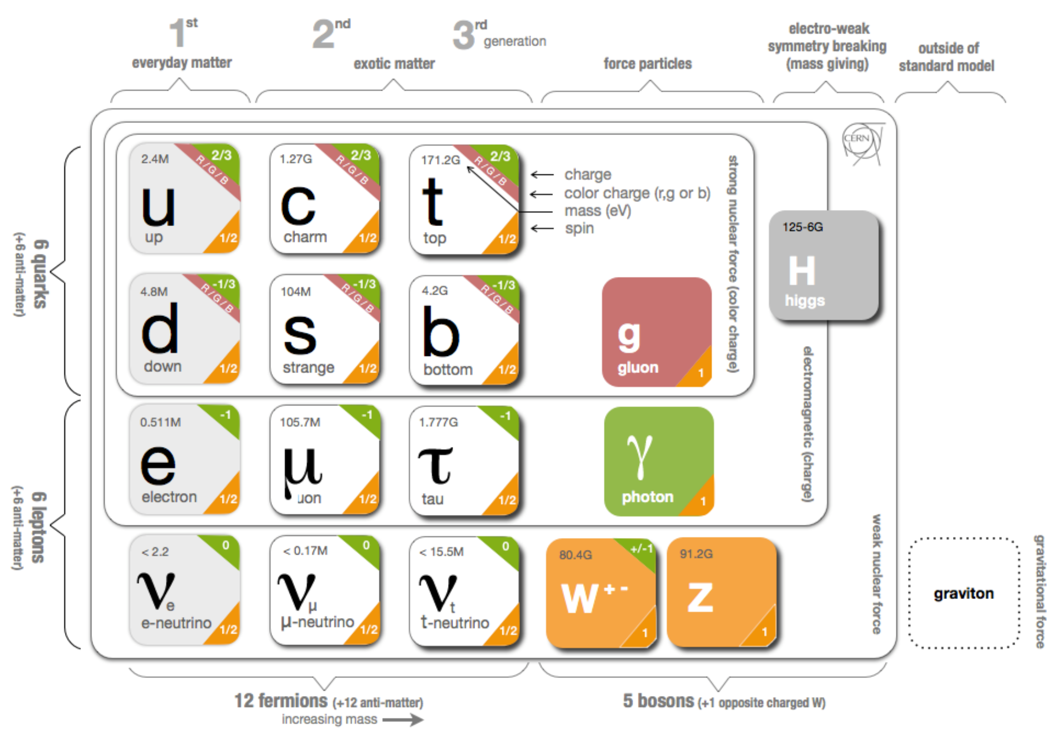


Figure 2.2: [13] A diagram of the Standard Model of particles. Shown are three generations of twelve fermions (quarks and leptons), all with spin $\frac{1}{2}$. The five force carriers (bosons) are shown: gluon, photon, W/Z bosons, and the Higgs boson. Also depicated is the graviton, a theoretical mediator of the gravitational force which is not currently in the Standard Model. All gauge bosons, except for the Higgs boson.

Figure 2.2 summarizes the fermions and bosons known today with table 2.2 providing a brief timeline of the discoveries. Table 2.1 summarizes the fundamental forces and how they interact with the different particles of the SM.

Property	Interaction			
	Gravitational	Weak	Electromagnetic	Strong
Acts On	Mass-Energy	Flavor	Electric Charge	Color Charge
Particles Experiencing	All	Quarks, Leptons	Charged	Quarks, Gluons
Particles Mediating	Graviton	W/Z bosons	Photons	Gluons
Strength at 10^{-18}m	10^{-41}	0.8	1	25

Table 2.1: [14] The strengths of the interactions (forces) are shown relative to the strength of the electromagnetic force for two u quarks separated by 10^{-18}m, the scale of quarks.

Fermions are spin-$\frac{1}{2}$ particles and follow Fermi-Dirac statistics, **bosons** are integer spin and follow Bose-Einstein statistics. These particles are the result of enforcing the symmetry in eq. (2.1) by introducing fields and interactions as mentioned in table 2.1. Fermions and bosons also have anti-particles of the same mass but opposite quantum charge. The **photon** is a mediator of the electromagnetic force and couples to all fermions with a non-zero electromagnetic charge; itself being massless, neutrally charged, and with spin 1. The electrically-neutral **gluon** is the mediator of the strong force and couples to all fermions with a color[2] charge. The gluon also carries color charge, color/anti-color pair, so it also participates in strong interactions[3] unlike the photon. The color flavor is SU(3) which means given the three colors red-green-blue, there are actually nine possible combinations of color/anti-color but only eight gluons in reality. The ninth possibility is a colorless singlet that is unobservable via strong interaction $\left(r\bar{r} + g\bar{g} + b\bar{b}\right)/\sqrt{3}$ and does not exist. Gluons are massless with spin 1. The other nice thing about the color terminology is that all naturally occurring particles are colorless[4]. It's a nice rule that helps to explain why you cannot make a particle out of two quarks qq or four quarks $qqqq$, but instead see particles like mesons $q\bar{q}$,

[2]Color does not actually mean "color" as if a quark actually appears red. Physicists would say that a quark has one unit of red-ness, for example.

[3]Quarks and gluons have different strong coupling strengths, with quark-gluon color factor $C_F = 4/3$ and gluon-gluon color factor $C_A = 3$ [15].

[4]Total amount of each color is zero or all three colors are present in equal amounts.

baryons qqq, and the antibaryons $\bar{q}\bar{q}\bar{q}$. The **W/Z bosons** are mediators of the weak force and couple to all fermions. The W bosons have electromagnetic charges of ± 1 while the Z boson is electromagnetically neutral, all with spin 1.

What	When	Who	Paper
Photon	1895	Wilhelm Röntgen	[16]
Electron	1897	J.J. Thomson	[17]
Proton	1919	Ernest Rutherford	[18]
Neutron	1932	James Chadwick	[19]
Muon	1937	Seh Neddermeyer, Carl Anderson	[20]
Electron neutrino	1956	Clyde Cowan, Frederick Reines	[21]
Muon neutrino	1962	BNL (AGS)	[22]
Up Quark Down Quark Strange Quark	1969	SLAC	[23, 24]
Charm Quark	1974	SLAC and MIT	[25, 26]
Tau	1975	SLAC-LBL	[27]
Bottom Quark	1977	Fermilab (E288)	[28]
Gluon	1979	DESY (PETRA)	[29]
W/Z Bosons	1983	CERN (UA1)	[30, 31]
Top Quark	1995	Fermilab (D0, CDF)	[32, 33]
Tau Neutrino	2000	Fermilab (DONUT)	[34]
Higgs Boson	2012	CERN LHC (ATLAS, CMS)	[2, 35]

Table 2.2: The abridged timeline of particle physics discoveries of the fermions and bosons that make up the Standard Model known today.

Finally, each force has an associated radiation where a real or virtual particle can be emitted. A photon can be radiated through the electromagnetic force, and this is known as bremsstrahlung. A quark can radiate a gluon[5] through the strong force. A similar process can also occur through the weak force where a quark can radiate a W/Z boson [36]. The search presented in this thesis focuses primarily on the strong interaction.

2.1.1 Spontaneous Symmetry Breaking

Spontaneous Symmetry Breaking (SSB) is the process in which a symmetry of a theory is not realized[6] in the lowest energy configuration (the vacuum expectation value, v.e.v. or vev). The classical example of describing such a situation is to imagine a pencil standing straight up on a table. The pencil is in a state of maximum energy with infinitely many ground states when it is lying horizontal on the table. The high energy state has a symmetry of rotation about the z-axis, but none of the ground states have this symmetry! So a physicist will say that when the pencil falls over, the rotational symmetry about the z-axis is "spontaneously broken".

To explain this with a toy model [11], consider a complex scalar field $\Phi = (\phi_1 + i\phi_2)/\sqrt{2}$. The Lagrangian density for this is

$$\mathcal{L} = \partial_\mu \Phi^\dagger \partial^\mu \Phi - m^2 \Phi^\dagger \Phi. \tag{2.5}$$

If Φ is constant, independent of space and time, only the $m^2\Phi^\dagger\Phi$ term contributes to the energy. Since the mass, m, is real, m^2 is positive and the energy is a minimum with the trivial solution $\phi_1 = \phi_2 = 0$. So $\Phi = 0$ is the ground state. Now, take the same equation but flip the sign in front of m^2 and now the Lagrangian is unstable as it is not bounded from

[5] Gluon showers

[6] I say realized, and not "broken", because I believe the phrase "broken" confuses people. There's nothing that is broken, but simply "transformed".

below. One can make this stable again by introducing a term $(m^2/2\phi_0^2)(\Phi^\dagger\Phi)^2$, and then the Lagrangian density is

$$\mathcal{L} = \partial_\mu \Phi^\dagger \partial^\mu \Phi - V(\Phi^\dagger\Phi), \qquad V(\Phi^\dagger\Phi) = \frac{1}{2\phi_0^2} m^2 \left[\Phi^\dagger\Phi - \phi_0\right]^2 + \text{constant}. \tag{2.6}$$

Just like in eq. (2.5), eq. (2.6) has minimum energy when Φ is constant (independent of space and time) where $\Phi^\dagger\Phi = \phi_0$. Instead of a unique field Φ, there is an infinite number of vacuum states described by $|\Phi|^2 = \phi_0$. In eq. (2.6), there is a global U(1) symmetry $\Phi \to \Phi' = e^{-i\theta}\Phi$ such that $\mathcal{L} \to \mathcal{L}' = \mathcal{L}$. If one picks out a particular direction in (ϕ_1, ϕ_2) space for which Φ is real, and take the vacuum state to be $(\phi_0, 0)$, the U(1) symmetry is lost. That is, the Lagrangian has some "global" symmetry that appears to have been lost when a ground state is picked out for the field.

So what does SSB give us? Well, you need to reinterpret the new fields after the loss of the symmetry. To expand about the ground state, the procedure is to put in $\Phi = \phi_0 + (\chi + i\phi)/\sqrt{2}$ for two real scalar fields χ, ψ, so the Lagrangian is now written in two terms: $\mathcal{L} = \mathcal{L}_{\text{free}} + \mathcal{L}_{\text{int}}$ with a free component and an interacting component corresponding to interactions between the free particles. Here

$$\mathcal{L}_{\text{free}} = \frac{1}{2}\partial_\mu\chi\partial^\mu\chi - m^2\chi^2 + \frac{1}{2}\partial_\mu\psi\partial^\mu\psi, \tag{2.7}$$

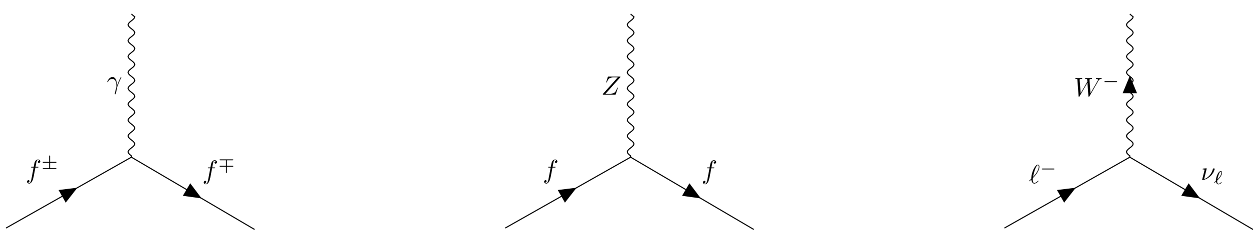
represents the free particle fields and contains terms quadratic in the fields. Notice that in eq. (2.7), there is a $-m^2\chi^2$ term which implies that the χ field has a scalar spin-zero particle of mass $m\sqrt{2}$. For the ψ field, there is no corresponding term so it is a massless, scalar, spin-zero particle. ψ is known as a Nambu-Goldstone[7] boson which are massless

[7]Yoichiro Nambu was a professor here at the University of Chicago.

particles that always arise as a result of the loss of a global symmetry [37].

2.1.2 Quantum Electrodynamics (QED)

QED is the oldest and perhaps simplest of the SM theories and has influenced the design of other theories. The QED theory corresponds to the $\mathrm{SU_W}(2) \otimes \mathrm{U_Y}(1)$ symmetry that is spontaneously broken by the Higgs mechanism providing mass-eigenstates corresponding to the $Z^0, W^{\pm}$ bosons, and the photon. All electromagnetic and weak phenomena are reducible to fundamental processes in fig. 2.3. To describe more complicated processes, you simply combine two or more replicas of this vertex. Each vertex introduces a factor of $\alpha = 1/137$ which is a small number, so only needs to sum over a smaller number of Feynman diagrams to get a reasonable approximation of the probability amplitude.



(a) Electromagnetic (b) Weak Neutral (c) Weak Charged

Figure 2.3: The elementary processes of Quantum Electrodynamics. Note that time is horizontal (a convention in ATLAS). In (a), a charged particle, f, enters, emits (or absorbs) a photon, γ, and exits. In (b), the Z boson mediates such processes. In (c), a lepton converts into corresponding neutrino with emission or absorption of $W^{\pm}$. These diagrams were made with TikZ-Feynman [38].

In order to describe QED, it will be sufficient to describe the process by which the masses of the electroweak bosons arise through the loss of global symmetries

$$\mathrm{SU}(2)_\mathrm{L} \times \mathrm{U}(1)_\mathrm{Y} \to \mathrm{U}(1)_\mathrm{EM} \tag{2.8}$$

This idea was explored by Yang and Mills in 1954 [39] and will be re-explored here. First, introduce a two-component field $\Phi = (\Phi_A, \Phi_B)$ where $\Phi_A = \phi_1 + i\phi_2$ and $\Phi_B = \phi_3 + i\phi_4$. In this case, a simple Lagrangian density that has global $\mathrm{U}(1) \otimes \mathrm{SU}(2)$ symmetry is described by eq. (2.6). If $V(\Phi^\dagger\Phi) = m^2\Phi^\dagger\Phi$, this Lagrangian density would correspond to four independent free scalar fields with the same mass m. In the SM, we need to describe the local symmetries from the global symmetries. Defining τ^k as the generators of SU(2), which are identical to the Pauli spin matrices in eq. (2.9)

$$\tau^0 = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}, \tau^1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \tau^2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \tau^3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \tag{2.9}$$

The U(1) transformation, $\Phi \to \Phi' = e^{-i\theta\tau^0}\Phi$ requires the introduction of a vector gauge field $B_\mu(x)\tau^0$ to become a local symmetry

$$B_\mu(x) \to B'_\mu(x)0 = B_\mu(x) + \frac{2}{g_1}\partial_\mu\theta, \qquad i\partial_\mu \to i\partial_\mu - \frac{g_1}{2}B_\mu, \tag{2.10}$$

where g_1 is a dimensionless parameter of the theory. For SU(2) where $U = e^{-i\alpha^k\tau^k}$ for three real numbers α^k and τ^k are the generators in eq. (2.9), a vector gauge field $W^k_\mu(x)$ is introduced:

$$\boldsymbol{W}_\mu(x) = W^k_\mu(x)\tau^k, \qquad \boldsymbol{W}_\mu(x) \to \boldsymbol{W}'_\mu(x) = \boldsymbol{U}(x)\boldsymbol{W}_\mu(x)\boldsymbol{U}^\dagger(x) + \frac{2i}{g_2}(\partial_\mu\boldsymbol{U}(x))\boldsymbol{U}^\dagger(x), \tag{2.11}$$

where g_2 is another dimensionless parameter of the theory. Finally, one needs to define the covariant derivative D_μ as

$$D_\mu \Phi = \left[\partial_\mu + \frac{ig_1}{2} B_\mu + \frac{ig_2}{2} \boldsymbol{W}_\mu\right] \Phi, \qquad D'_\mu \Phi' = e^{-i\theta} \boldsymbol{U} D_\mu \Phi. \tag{2.12}$$

So the locally gauge invariant Lagrangian density corresponding to eq. (2.6) is

$$\mathcal{L}_\Phi = (D_\mu \Phi)^\dagger D^\mu \Phi - V(\Phi^\dagger \Phi) \tag{2.13}$$

So now we're ready to write out the dynamical contribution to the Lagrangian density associated with the gauge fields:

$$\mathcal{L}_{\text{dyn}} = -\frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{8} \text{Tr}\left(\boldsymbol{W}_{\mu\nu} \boldsymbol{W}^{\mu\nu}\right), \tag{2.14}$$

with the field strength tensor for $B_\mu(x)$ straightforward to write out. As the SU(2) group is non-Abelian, $\boldsymbol{W}_\mu(x)$ is trickier

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \tag{2.15a}$$

$$\boldsymbol{W}_{\mu\nu} = \left[\partial_\mu + \frac{ig_2}{2} \boldsymbol{W}_\mu\right] \boldsymbol{W}_\nu - \text{same, but } \mu \leftrightarrow \nu. \tag{2.15b}$$

Now, because of the nice features of the Pauli matrices, specifically that $\text{Tr}(\tau^i)^2 = 2$ and $\text{Tr}(\tau^i \tau^j) = 0, i \neq j$, eq. (2.14) can be written more simply as

$$\mathcal{L}_{\text{dyn}} = -\frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{4} W^3_{\mu\nu} W^{3\mu\nu} - \frac{1}{2} W^-_{\mu\nu} W^{+\mu\nu}, \tag{2.16}$$

where the W field has defined complex mixing for convenience

$$W_\mu^\pm = \frac{1}{\sqrt{2}}\left(W_\mu^1 - iW_\mu^2\right), \qquad W_{\mu\nu}^\pm \text{ written similarly, and} \tag{2.17a}$$

$$W_{\mu\nu}^3 = \partial_\mu W_\nu^3 - \partial_\nu W_\mu^3 - ig_2\left(W_\mu^- W_\nu^+ - W_\nu^- W_\mu^+\right) \tag{2.17b}$$

Now, we are at the point with eq. (2.16) to apply the methodology of losing the symmetry as described previously in section 2.1.1. Since there are three real parameters $\alpha^k(x)$ in SU(2), a gauge is chosen such that $\Phi_A = 0$ (two conditions) and $\Phi_B = \phi_0$ is real (one condition). The ground and excited states are then of the form

$$\Phi_{\text{ground}} = \begin{pmatrix} 0 \\ \phi_0 \end{pmatrix}, \tag{2.18a}$$

$$\Phi_{\text{excited}} = \begin{pmatrix} 0 \\ \phi_0 + h(x)/\sqrt{2} \end{pmatrix}, \qquad h(x) \text{ is real} \tag{2.18b}$$

so plugging this into eq. (2.13), one obtains

$$\begin{aligned}\mathcal{L}_\Phi =& \frac{1}{2}\partial_\mu h \partial^\mu h + \frac{g_2^2}{2} W_\mu^- W^{+\mu}\left(\phi_0 + \frac{h}{\sqrt{2}}\right)^2 + \underbrace{\frac{1}{4}\left(g_1^2 + g_2^2\right) Z_\mu Z^\mu \left(\phi_0 + \frac{h}{\sqrt{2}}\right)^2}_{Z_\mu = W_\mu^3 \cos\theta_w - B_\mu \sin\theta_w} \\ & \underbrace{- m^2 h^2 + \frac{m^2}{\phi_0\sqrt{2}} h^3 + \frac{m^2}{8\phi_0^2} h^4}_{V(h)}.\end{aligned} \tag{2.19}$$

$V(h)$ here is the Higgs potential which takes on the shape of a mexican hat; a local maxima at the origin and the potential drops off before rising up again with a local minima along a circle around the origin. $A_\mu = W^3_\mu \sin\theta_w + B_\mu \cos\theta_w$ (the orthogonal complement to Z_μ) with

$$\cos\theta_w = \frac{g_2}{\sqrt{g_1^2+g_2^2}} \qquad \text{and} \tag{2.20a}$$

$$\sin\theta_w = \frac{g_1}{\sqrt{g_1^2+g_2^2}}, \tag{2.20b}$$

where w stands for the Weinberg angle. So we have $\mathcal{L} = \mathcal{L}_{\text{dyn}} + \mathcal{L}_\Phi$ from eqs. (2.16) and (2.19). Putting it all together and rewriting a little bit[8]

$$\begin{aligned}
\mathcal{L} =& \frac{1}{2}\partial_\mu h \partial^\mu h - m^2 h^2 \\
& - \frac{1}{4} Z_{\mu\nu} Z^{\mu\nu} + \frac{1}{4}\phi_0^2 (g_1^2 + g_2^2) Z_\mu Z^\mu \\
& - \frac{1}{4} A_{\mu\nu} A^{\mu\nu} \\
& - \frac{1}{2}\left[(D_\mu W^+_\nu)^* - (D_\nu W^+_\mu)^*\right]\left[D^\mu W^{+\nu} - D^\nu W^{+\mu}\right] + \frac{1}{2} g_2^2 \phi_0^2 W^-_\mu W^{+\mu} \\
& + \mathcal{L}_{\text{int}}
\end{aligned} \tag{2.21}$$

where $Z_{\mu\nu} = \partial_\mu Z_\nu - \partial_\nu Z_\mu$ ($A_{\mu\nu}$ is written similarly) and $D_\mu W^+_\nu = (\partial_\mu i g_2 \sin\theta_w A_\mu) W^+_\nu$.

[8] It helps to have a really, *really* big chalkboard here.

Looking at this, one can extract out the masses of the particles associated with the fields:

$$m_A = 0, \tag{2.22a}$$

$$m_W = \phi_0 \frac{g_2}{\sqrt{2}} = 80.385 \pm 0.015\,\text{GeV}, \tag{2.22b}$$

$$m_Z = \phi_0 \sqrt{\frac{g_1^2 + g_2^2}{2}} = 91.1876 \pm 0.0021\,\text{GeV}, \tag{2.22c}$$

$$m_h = m\sqrt{2} = 125.09 \pm 0.24\,\text{GeV}. \tag{2.22d}$$

From experimental observations [40], we know all of these masses experimentally, including the mass of the Higgs boson found on July 4th, 2012 [2, 35] by the ATLAS and CMS collaborations. Finally, notice that $\cos\theta_w = m_W/m_Z$ is a reported ratio in PDG [40] as well. So what we've seen from basic principles of QED is that starting with a two-component complex field (composed of four real fields), one can find the global symmetry of SU(2)⊗SU(1), lose that symmetry locally as in eq. (2.8), trigger the Higgs mechanism, and find a Nambu-Goldstone boson instead. The real, initially-massless fields now gain mass[9] through their interaction with $h(x)$ and we can write out the interacting portion of the Lagrangian $\mathcal{L}_{\text{int}}$.

2.1.3 Quantum Chromodynamics (QCD)

QCD is a quantum field theory describing the strong force, governed by the symmetry $\text{SU}_\text{C}(3)$[41, 42]. I will state the Lagrangian density for this theory to illuminate how this compares to eq. (2.21) but the procedure is very similar to QED. In QCD, there are three fields[10] for each flavor of quark and are put into color triplets. The top quark, for example,

[9] Well, except for the massless photon of course. Technically, the photon interacts with the "Higgs doublet" but this is not the component of the Higgs field whose excitations are the Higgs bosons.

[10] Read: color.

looks like eq. (2.23).

$$\boldsymbol{t} = \begin{pmatrix} t_r \\ t_g \\ t_b \end{pmatrix} \tag{2.23}$$

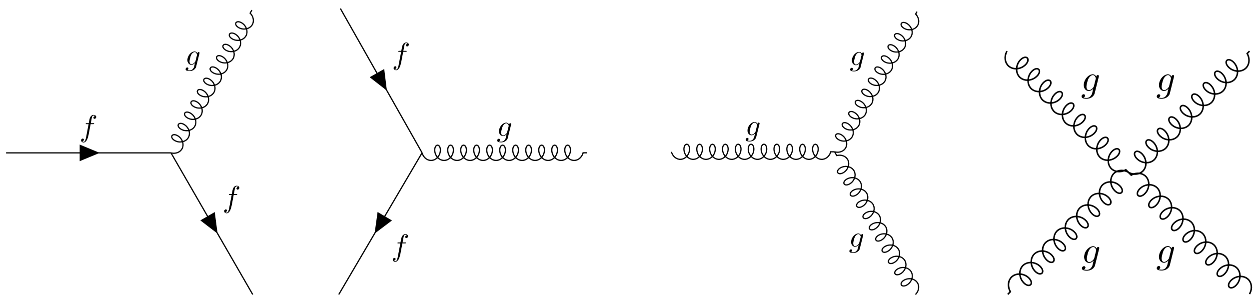
where $t_c, \{c \mid r, g, b\}$ represents the four-component Dirac spinors. You state a local SU(3) transformation under which the theory is invariant $\boldsymbol{q} \to \boldsymbol{q'} = \boldsymbol{U}\boldsymbol{q}$. This lets us write down eq. (2.24) where the gluon gauge fields, $\boldsymbol{G}_\mu$, are similar to the weak gauge fields, $\boldsymbol{W}_\mu$, and the covariant derivative $D_\mu = \partial_\mu + ig\boldsymbol{G}_\mu$ as by Yang-Mills construction [39].

$$\mathcal{L}_{\text{QCD}} = \underbrace{-\frac{1}{4}\sum_{a=1}^{8} G^a_{\mu\nu}G^{a\mu\nu}}_{\mathcal{L}_{\text{gluon}}} + \underbrace{\sum_{f=1}^{6}\left[\bar{\boldsymbol{q}}_f i\gamma^\mu(\partial_\mu ig\boldsymbol{G}_\mu)\boldsymbol{q}_f - m_f\bar{\boldsymbol{q}}_f\boldsymbol{q}_f\right]}_{\mathcal{L}_{\text{quark}}} \tag{2.24}$$

$\mathcal{L}_{\text{gluon}}$ has a sum over the eight gluons of $\text{SU}_\text{C}(3)$ and provides the kinetic terms for gluons and their self-interactions as in fig. 2.4c. $\mathcal{L}_{\text{quark}}$ has a sum over the six flavors of quarks with m_f the "true" masses given to the quarks by coupling to the Higgs field; and provides the kinetic terms for the quarks and their interactions with gluons.

Any number of interactions may follow from a single initial state, but the probability of a final state occurring decreases as the complexity of the final state increases. A set of Feynman diagrams representing basic strong force interactions is shown in fig. 2.4. The probability of a given Feynman diagram is determined by many factors, including the probabilities of each interaction point, all proportional to the strong coupling constant.

All lowest order QCD diagrams are on the order of $O(\alpha_S^2)$, such as for example, the Feynman diagram representing the strong force binding two quarks together to make hadrons such as neutrons and protons is shown in fig. 2.5.



(a) $q \to qg$ (b) $q\bar{q} \to g$ (c) $g \to gg$ (d) $gg \to gg$

Figure 2.4: A few Feynman diagrams of basic QCD interactions to lowest order, including (a) gluon radiation, (b) quark/anti-quark annihilation, (c) gluon splitting, and (d) gluon self-coupling. These diagrams were made with TikZ-Feynman [38].

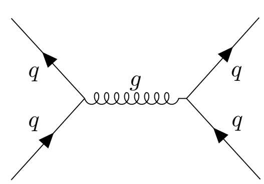


Figure 2.5: A Feynman diagram at leading order (LO) with probability amplitude proportional to the square of the strong coupling constant. This particular Feynman diagram represents the interaction between quarks that, for example, binds them into hadrons. These diagrams were made with TikZ-Feynman [38].

So it seems that at least up to this point, QCD looks suspiciously like QED, and that's not an accident. There are some important differences, such as the size of the coupling constants where QED introduces a factor $\alpha_{\text{QED}} = 1/137$ but for QCD, the factor $\alpha_{\text{QCD}} > 1$ is larger[11]. This was initially a headache as calculations[12] beyond NLO would contribute just as equally, if not more. It was found that the strong coupling constant decreases at higher energy scales (or probing smaller distances) and is called a "running coupling constant"[13]. This discovery by Gross, Wilczek, and Politzer won the Nobel Prize in 2004 [43, 44, 42, 45]. This phenomenon is known as **asymptotic freedom**, and allows the Feynman diagrams as a legitimate tool for QCD calculations in the high-energy regime. As the energy scale goes up, the strength of the strong force goes down to zero, allowing for perturbative calculations. In the other direction, as the energy scale goes down, a non-perturbative approach needs to be taken as the coupling constant blows up [**doi:10.1146**]. This will be discussed briefly later in this section. As mentioned in [42, 45], there is a kind of competition between the quark loops and gluon loops in the Feynman diagrams that determines whether the effective coupling constant increases or decreases at short distances. It turns out, compared the coupling constants as a function of the energy scale between QED and QCD, it's clear to see why the running coupling is different [46] in eq. (2.25).

$$\alpha_{\text{QED}}(Q^2) = \frac{e^2}{4\pi - \frac{e^2}{3\pi}\ln\left(\frac{Q^2}{4m^2}\right)} \qquad \text{and} \tag{2.25a}$$

$$\alpha_{\text{QCD}}(Q^2) = \frac{g^2}{4\pi - \frac{g^2}{4\pi}\frac{1}{n_c}\left[2n_f - 11n_c\right]\ln\left(\frac{Q^2}{\lambda^2}\right)} \tag{2.25b}$$

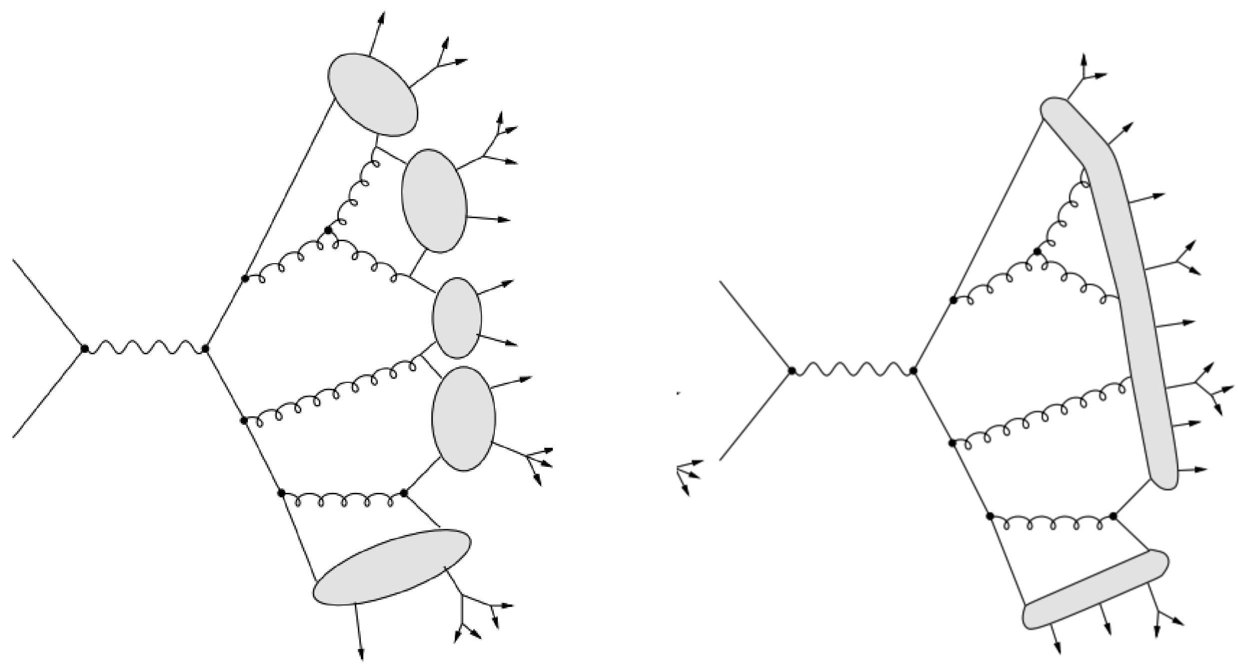
The special piece to notice is in the denominator of α_{QCD} in eq. (2.25) where n_c, n_f are for the number of colors and number of flavors in the theory. If this piece is negative, that is,

[11]Hence the theories are said to be strongly coupled or weakly coupled.

[12]These calculations would involve infinitely more loops.

[13]This also happens for α_{QED} too.

$f(n_f, n_c) \equiv 2n_f - 11n_c < 0$, then the α_{QCD} decreases at short distances (large Q^2). For SM QCD there are 6 flavors (quarks) and 3 colors, so $f(n_f, n_c) < 0$ and this is the basis of asymptotic freedom where color-charged particles barely interact with each other at small distances.



(a) Cluster

(b) Lund

Figure 2.6: [47] Cartoon of the (a) cluster hadronization model which treats individual color singlets separately and the (b) lund string hadronization model which propagates field lines of color flux. These two leading models approximate the non-perturbative process of hadronization to map colored partons onto stable, colorless hadrons.

On the opposite end of the energy scale, for low Q^2 and large distances on the order of femtometers, a non-perturbative approach needs to be taken to evaluate the interactions. Physics simulators, which try to approximate the non-perturbative behavior of QCD, pick from two different popular options shown in fig. 2.6. The cluster model starts with gluon splitting into $q\bar{q}$ to form clusters that are used to predict final state hadrons. The Lund string model, on the other hand, uses the $q\bar{q}$ pair to estimate the intensity of the color flux string and generates gluons and hadrons based on kinks in this flux.

There currently is no analytic proof of this behavior (or the transition to this behavior) known as **color confinement**, but it can be observed experimentally at a particle detector. To describe it in a qualitative manner, as quarks and gluons separate, the strong force increases in strength. At a certain point, it is energetically favorable to produce a quark/anti-quark pair rather than put in more work to separate the two particles. In other words, separating two particles with color will produce bound states which are colorless. What this means for detectors like ATLAS is that non-colorless particles cannot be directly detected due to color confinement. All physically and directly observable particles are colorless.

The boosted partons (quarks and gluons) that come flying out of the proton-proton collisions with large amounts of energy will create colorless bound states. This process is called **hadronization** and refers to the transition of colored partons to colorless hadrons. Partons can also radiate collinear gluons which in turn radiate $q\bar{q}$ collimated pairs, through a process known as **showering**. These steps are shown in fig. 2.7, a partonic representation of the process of a single colored parton generating multiple, colorless, measurable hadron showers. The green arrows on either side of the event are the proton bunches which have gluons radiating which form two different groups of interactions. The hard scatter[14] (large red circle) of a proton-proton collision is the highest energy interaction in the event. The secondary interactions (purple blob) form the underlying event, involve smaller momentum transfers. From the hard scatter, the high energy partons shower according to perturbative QCD (red showers). At a low enough energy level where perturbation theory becomes invalid and color confinement takes over, the partons hadronize (green blobs) into various colorless hadrons.

[14] At the large energies of the LHC, the "core process" here are gluon-gluon scattering.

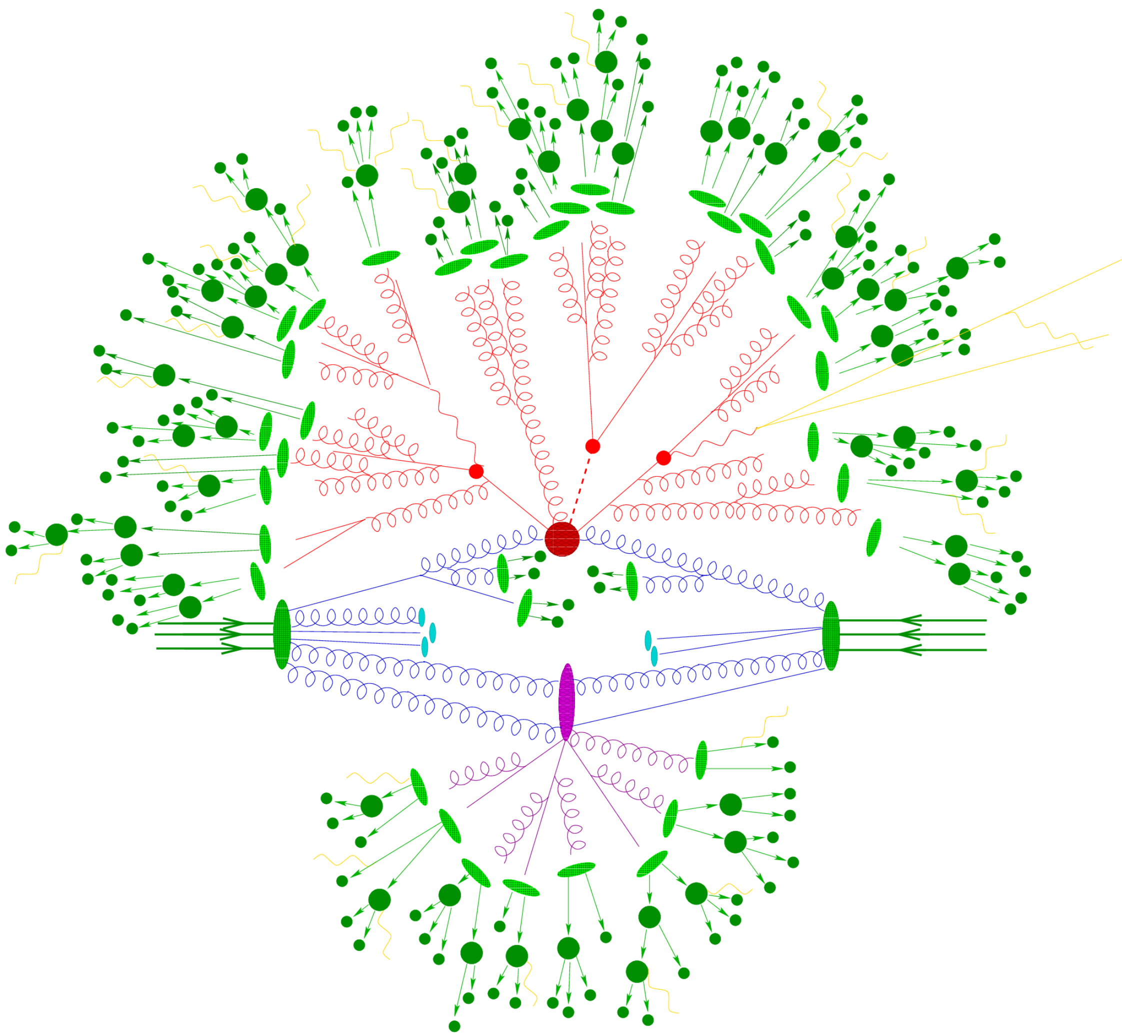
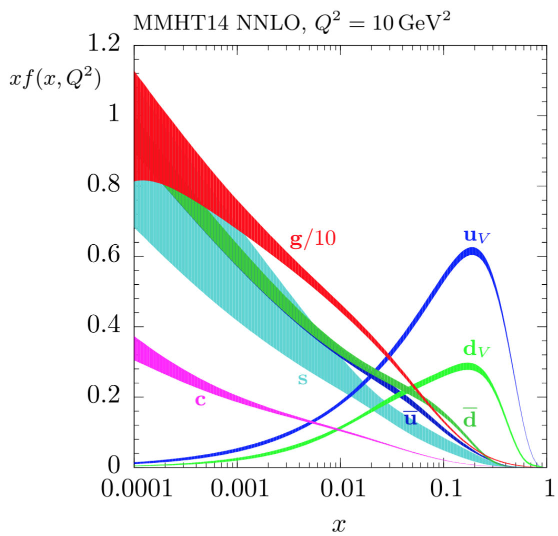


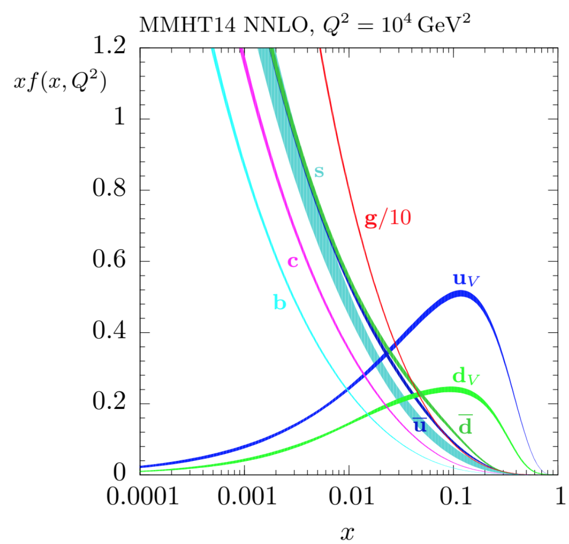
Figure 2.7: [48] Sketch of a hadron-hadron collision as simulated by a Monte-Carlo event generator. The red blob in the center represents the hard collision, surrounded by a tree-like structure representing Bremsstrahlung as simulated by parton showers. The purple blob indicates a secondary hard scattering event. Parton-to-hadron transitions are represented by light green blobs, dark green blobs indicate hadron decays, while yellow lines signal soft photon radiation.

2.1.4 Parton Distribution Function

The name "parton" was proposed by Richard Feynman in 1969 [49] as a generic description for any particle constituent within the proton, neutron, and other hadrons. At first, the hadrons were thought to consist of doublets and triplets of quarks ($q\bar{q}$ and qqq). However, through high-energy proton-proton collisions at the LHC and the interaction processes, these "valence" quarks and gluons can also produce an arbitrary number of lower-energy virtual partons, "sea"[15] quarks and gluons. These were first observed by James Bjorken and Emmanuel Paschos in 1969 [50]. Now, we know that protons (neutrons) are made up of two (one) up quarks, u, and one (two) down quark, d, along with the gluons, g, that hold them together.



(a) $Q^2 = 10\,\mathrm{GeV}^2$



(b) $Q^2 = 10^4\mathrm{GeV}^2$

Figure 2.8: [51] MMHT2014 NNLO PDFs at $Q^2 = 10\,\mathrm{GeV}^2$ and $Q^2 = 10^4\mathrm{GeV}^2$ with associated 68% confidence-level uncertainty bands. The function $xf(x,Q^2)$ is plotted versus x for different flavors: $u, \bar{u}, d, \bar{d}, s = \bar{s}$, and g.

While the Standard Model enables the calculations of cross-sections as a function of the

[15] Personally, as a Deaf person, I propose that we call them "ocean" or "plum" quarks, to reduce confusion with the existing "c" quarks.

energies of colliding partons (quarks and gluons), the LHC is a proton-proton collider. It is also useful to know the cross-sections for a given proton energy, the parton distribution function (PDF) [52]. This is because the collisions at the LHC are really between partons inside the protons. The PDF is a function that provides the probability density of finding the given parton in the given hadron with the given momentum. PDFs are parameterized by Q^2 and "Bjorken x" (or just x). The Q^2 corresponds to the energy scale of the collision[16] process and x represents the momentum fraction of the proton that the interacting parton holds. For proton-proton collider experiments like LHC, proton PDFs are the most interesting, reconstructed using data from proton scattering experiments. Figure 2.8 shows one such example of a PDF: the u, d quark, and the g gluon generally dominate at low energies (low Q^2), while other virtual partons are more likely to participate in the interaction processes at high energies (high Q^2). Many other PDFs exist [53, 54, 51, 55], and for LHC Run 2, the global PDFs NNPDF3.0, MMHT14 and CT14 are the latest used which use the data from LHC Run 1 for further constraints. There is some uncertainty in these PDFs which contribute to uncertainties in the predicted proton-proton cross-sections and are often one of the dominant sources of uncertainty for many important searches and analyses at the LHC, especially for precision cross-section measurements.

Factorization [56] is a concept that was implicit in the discussion about hadronization and PDFs. In particular, what factorization allows us to do is define a cutoff scale Q_F above which collinear radiation is directly treated and below which it is absorbed into the PDF definition. Effectively, this allows us to separate the calculation of phenomena which are perturbatively calculable from phenomena which are not. The total cross section σ for a collision process [57], $ab \to n$ may be derived by integrating over all possible initial state momentums for partons a and b, hadron h, the parton p_{T} fraction x_a^h and x_b^h, and weighting

[16] The center-of-mass energy at a proton-proton collider, $\sqrt{s}$, is not related to Q^2, but is instead the upper bound on Q^2.

them by their PDF f_a^h and f_b^h can be written as shown in eq. (2.26).

$$\sigma(Q_F, Q_R) = \sum_{a,b} \int_0^1 \mathrm{d}x_a \mathrm{d}x_b \int_{\text{hadrons}} f_a^h(x_a, Q^2) f_b^h(x_b, Q^2) \mathrm{d}\sigma_{ab \to n} \tag{2.26}$$

The other scale involved is the renormalization scale Q_R [58, 59] which accounts for the logarithmically divergent contributions of the Lagrangian through the process of renormalization. Unlike Q_F which represents the scale at which the hadron is being probed, Q_R is a non-physical effect that accounts for the limited knowledge used in lower order calculations.

2.1.5 Top Quark Decays

The top quark is a 3rd generation parton and is fundamental to this thesis analysis as we search for four Lorentz-boosted top quarks in the final state. As such, it is appropriate to provide a little bit more detail about the top quark and its decay. The timescale for strong force interactions is on the order of 10^{-24}s. The top quark has a lifetime of 10^{-25}s which is due to its large mass. Thus, the top quark is a unique parton in that it decays before it can hadronize, allowing physicists to measure the "bare mass" of the top quark [60]. Figure 2.9 shows the two dominant decay modes of a top quark, through the weak interaction, producing a W-boson and a down-type quark (down, strange, or bottom).

The W-boson branching ratios are listed in eq. (2.27) [40]. The top quark will decay hadronically (to two quarks) approximately 70% of the time and leptonically approximately 30% of the time. As well as top quark decays, I also discuss $t\bar{t}$[17] which has three different kinds of decays: hadronic (both W bosons decay hadronically), semi-leptonic (one W boson decays leptonically), and fully leptonic (both W bosons decay leptonically). $t\bar{t}$ decays hadronically

[17] Colloquially "tee-tee-bar" or written "ttbar".

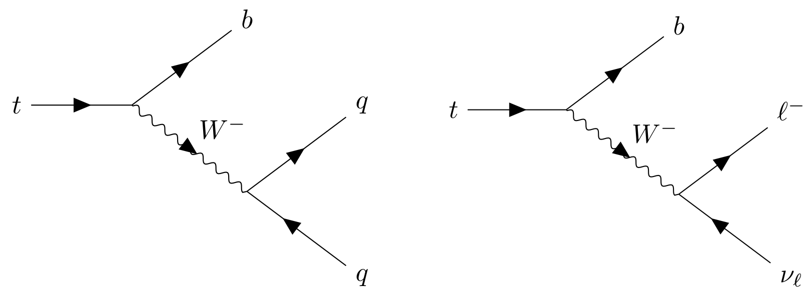


Figure 2.9: Feynman diagrams showing the top quark decays for (a) hadronic and (b) leptonic. These diagrams were made with TikZ-Feynman [38].

about 50% of the time, semi-leptonically about 40% of the time, and full-leptonically about 10% of the time.

$$\mathrm{BR}(W \to e\bar{\nu}_e) = 0.1046 \pm 0.0042(\mathrm{stat}) \pm 0.0014(\mathrm{syst}), \tag{2.27a}$$

$$\mathrm{BR}(W \to \mu\bar{\nu}_\mu) = 0.1050 \pm 0.0041(\mathrm{stat}) \pm 0.0012(\mathrm{syst}), \tag{2.27b}$$

$$\mathrm{BR}(W \to \tau\bar{\nu}_\tau) = 0.1075 \pm 0.0052(\mathrm{stat}) \pm 0.0021(\mathrm{syst}), \tag{2.27c}$$

$$\mathrm{BR}(W \to q\bar{q}) = 0.6832 \pm 0.0061(\mathrm{stat}) \pm 0.0028(\mathrm{syst}). \tag{2.27d}$$

2.2 Beyond the Standard Model

The Standard Model has been tested over the last few decades by many experiments and shown to be robust. The fermion fields of leptons and quarks interact through the mediation of vector bosons. The renormalizability of the SM requires that the vector boson fields be introduced through the requirement of local gauge symmetry as in sections 2.1.1 to 2.1.3. However, we know this is not a complete model given the success so far, as certain assumptions are still made that need to be reconciled, motivated by naturalness, such as:

- the matter/anti-matter asymmetry not observed in the detector [61],

- the fine-tuning required to the quantum corrections to keep the Higgs mass around the electroweak scale [62],
- the lack of inclusion of gravity, and the lack of dark matter candidates [63] even though it is largely agreed upon that dark matter exists [64],
- the scale difference between the Planck scale and the Electroweak scale (the so-called Hierarchy problem) [65],
- and many more [66, 67, 68, 69]

Many Beyond the Standard Model (BSM) theories have been proposed, all with a variety of testable signatures. Attempts have been made to carry unification further, by combining the electroweak and strong interactions in a higher, unified symmetry, which could only manifest at extremely high energies of order 10e13 TeV. The Higgs boson interacts with all other gauge fields in QED and QCD so that quantum loops in the Feynman diagrams are created to correct the Higgs mass. It seems absurd, and incredibly coincidental, that the Higgs mass, which was expected to be around the Planck scale, receives gigantic corrections on the order of 10^{17} to be on the electroweak scale. In addition, the coupling of the Higgs to some quarks covers two orders of magnitudes which does not seem *natural*. Hence, **naturalness** is a strong motivation for many physicists, myself included.

If you look at SM and understand the corrections to the Higgs mass, the Feynman diagram in fig. 2.10 shows an example of the loop correction to the Higgs mass that requires such precise fine-tuning that it doesn't seem natural for the Higgs mass to be as light as observed in 2012. The top mass has the largest coupling, and therefore the largest correction to the Higgs mass, which means this quantum-level correction is roughly described by eq. (2.28), where λ_t is the Yukawa coupling of the top and Λ_{UV} is the ultraviolet momentum cutoff used to regulate the loop integral of the theory [4], which for the SM is the Planck mass.

This naturalness motivation is the strongest motivation for trying to find supersymmetry.

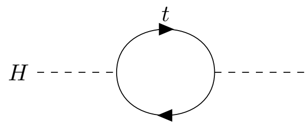


Figure 2.10: [4] An example of a loop diagram which corrects the Higgs mass.

$$\Delta m_H^2 = -\frac{\lambda_t|^2}{8\pi^2}\Lambda_{\mathrm{UV}}^2 + \dots \tag{2.28}$$

For the rest of this section, I discuss the theoretical framework of supersymmetry, the supersymmetric particles, and introduce the simplified models that I studied.

2.2.1 Supersymmetry

Supersymmetry (SUSY) [70, 71, 72, 73, 74, 75] is a generalization of space-time symmetries that predicts new bosonic partners for the fermions and new fermionic partners for the bosons of the SM. If R-parity[18] is conserved [76], SUSY particles are produced in pairs and the lightest supersymmetric particle (LSP) is stable. The scalar partners of the left-handed and right-handed quarks, the squarks $\tilde{q}_{\mathrm{L}}$ and $\tilde{q}_{\mathrm{R}}$, can mix to form two mass eigenstates $\tilde{q}_1$ and $\tilde{q}_2$, ordered by increasing mass. SUSY can solve the hierarchy problem [77, 78, 79, 80] reducing unnatural tuning in the Higgs sector by orders of magnitude, provided that the superpartners of the top quark have masses not too far above the weak scale. The large top Yukawa coupling results in significant $\tilde{t}_{\mathrm{L}}$–$\tilde{t}_{\mathrm{R}}$ mixing so that the mass eigenstate $\tilde{t}_1$ is typically lighter than the other squarks [81, 82].

If supersymmetry exists, it should contain SSB. From a theoretical perspective, there should

[18] Also known as Matter parity. All SM particles and Higgs boson have even R-parity $P_R = +1$ while the squarks, sleptons, gauginos, and higgsinos have odd R-parity $P_R = -1$.

be a Lagrangian density that is invariant under supersymmetry but a ground state that is not. This is analogous to what has been discussed before in section 2.1.1. On top of this, the theory should be renormalizable to compensate for the effects of self-interactions, infinities arising in calculated quantities, and the differences in descriptions between small-distance-scale physics and large-distance-scale physics [4]. In a supersymmetric extension of the SM [83, 84], each of the known fermions (bosons) is therefore either in a chiral or gauge supermultiplet and must have a superpartner boson (fermion) with spin differing by $\frac{1}{2}$ unit. The names, while appearing somewhat humorous, serve to make the connection from the superpartner to their physical SM partner more obvious are generated as follows:

- the names for the spin-0 partners of the quarks and leptons are constructed by preprending an "s" for scalar (or superpartner) to be called squarks, sleptons, and sfermions,
- the symbols for the squarks and sleptons are the same as for the corresponding fermion, but with a tilde added such as $\tilde{e}_L, \tilde{e}_R$[19],
- the supersymmetric fermions take the name of their superpartner bosons, but with an "ino" appended, such as "wino", "gluino"

Given this fermion-boson symmetry, as well as R-parity, the Lagrangian density for an unbroken symmetric theory can be written down. Using a similar mechanism to break this symmetry, gauge fields are introduced, and mass states arise as a mixing of the gauge states. Unlike electroweak which is slightly easier to break, there are a few gotchas this time:

- In SM, there is one Higgs boson; in SUSY, there are two complex Higgs doublets. The reason for this is that the fermionic partner of a Higgs must be able to cancel gauge

[19] The leptons and quarks have left/right handedness and superpartners for each version, as the superpartners are spin-0. SM neutrinos ν_ℓ are always left-handed, so superpartners are just $\tilde{\nu}_\ell$.

anomalies which are usually the traces of hypercharge matrices. In the SM this works out because $Y = 0$, but for SUSY, $Y = \pm 1$. So there must be two complex Higgs doublets to account for each hypercharge variation. This is a heuristic motivation.

- In the SM, the Yukawa couplings between the Higgs and the 3rd generation fermions (t, b, τ) are much larger than the first and second generations. Normally, it is not very easy to diagonalize the gauge eigenstates for the fermions, however in the minimal supersymmetric model (MSSM) that is being considered, an approximate can be made to treat the Yukawa couplings for first and second generation as negligible. An example is shown in eq. (2.29) [4].

- The higgsinos and electroweak gauginos mix with each other because of the effects of electroweak symmetry breaking. The neutral higgsinos and electroweak gauginos mix to form **neutral**inos[20], and the charged versions mix to form **charg**inos.

$$\boldsymbol{M}_{\chi^0} = \begin{pmatrix} M_1 & 0 & -m_Z s_W c_\beta & m_Z s_W s_\beta \\ 0 & M_2 & m_Z c_W c_\beta & -m_Z c_W s_\beta \\ -m_Z c_\beta s_W & m_Z c_\beta c_W & 0 & -\mu \\ m_Z s_\beta s_W & -m_Z s_\beta c_W & -\mu & 0 \end{pmatrix}. \qquad s_k \equiv \sin\theta_k, c_k \equiv \cos\theta_k \tag{2.29}$$

In eq. (2.29), the M_i terms come from the soft component of the supersymmetric Lagrangian [4, eq. 6.3.1]. Diagonalizing this matrix allows us to form the neutralinos χ_k^0 as a mixture of the wino, bino, and higgsinos gauge eigenstates. Different mixtures of these

[20]Not, as I sometimes mistakenly think, the superpartners of the neutrinos which are the **s**neutrinos.

gauge eigenstates correspond to different decay products in the final state. A similar procedure exists for the charginos $\chi_k^{\pm}$. Both neutralinos and charginos are conventionally ordered (and labeled) in k in terms of increasing mass, such that $\tilde{\chi}_1^{\pm} < \tilde{\chi}_2^{\pm}$ and $\tilde{\chi}_1^0 < \tilde{\chi}_2^0 < \tilde{\chi}_3^0 < \tilde{\chi}_4^0$. Table 2.3 shows a summary of the various SUSY particles and their mass eigenstates.

Names	Spin	P_R	Gauge Eigenstates	Mass Eigenstates
Higgs bosons	0	+1	$H_u^0, H_d^0, H_u^+, H_d^-$	$h^0, H^0, A^0, H^{\pm}$
squarks	0	-1	$\tilde{u}_L, \tilde{u}_R, \tilde{d}_L, \tilde{d}_R$	(same)
			$\tilde{s}_L, \tilde{s}_R, \tilde{c}_L, \tilde{c}_R$	(same)
			$\tilde{t}_L, \tilde{t}_R, \tilde{b}_L, \tilde{b}_R$	$\tilde{t}_1, \tilde{t}_2, \tilde{b}_1, \tilde{b}_2$
sleptons	0	-1	$\tilde{e}_L, \tilde{e}_R, \tilde{\nu}_e$	(same)
			$\tilde{\mu}_L, \tilde{\mu}_R, \tilde{\nu}_\mu$	(same)
			$\tilde{\tau}_L, \tilde{\tau}_R, \tilde{\nu}_\tau$	$\tilde{\tau}_1, \tilde{\tau}_2, \tilde{\nu}_\tau$
neutralinos	$\frac{1}{2}$	-1	$\tilde{B}^0, \tilde{W}^0, \tilde{H}_u^0, \tilde{H}_d^0$	$\tilde{\chi}_1^0\ \tilde{\chi}_2^0\ \tilde{\chi}_3^0\ \tilde{\chi}_4^0$
charginos	$\frac{1}{2}$	-1	$\tilde{W}^{\pm}, \tilde{H}_u^+, \tilde{H}_d^-$	$\tilde{\chi}_1^{\pm}, \tilde{\chi}_2^{\pm}$
gluino	$\frac{1}{2}$	-1	$\tilde{g}$	(same)
goldstino (gravitino)	$\frac{1}{2}$ $\frac{3}{2}$	-1	$\tilde{G}$	(same)

Table 2.3: [4] The undiscovered particles in the Minimal Supersymmetric Standard Model (with sfermion mixing for the first two families assumed to be negligible).

Now, as shown in fig. 2.11 for SUSY compared to fig. 2.10 for SM, the stop squark is the bosonic superpartner to the top quark which provides an equal and opposite contribution[21] to the correction of the Higgs mass. In the limit of top-squark masses $m_{\tilde{t}_k} \gg m_t$ much greater than the top quark mass, the largest **finite** correction to the higgs mass m_{h^0} is [4] shown in eq. (2.30), where $\Delta_{\text{threshold}}$ is a small correction based on the top-squark mixing angle and the Higgs quartic coupling, α is a mixing angle of the Higgs couplets, and λ_t is

[21]Spin-statistics theorem states that fermions have a negative contribution and bosons have a positive contribution.

the top Yukawa coupling. So in addition to having a light stop [85, 86, 87], there is also a strong motivation[22] for a light gluino [88], as the gluino couples to the stop squark and pulls the stop mass up. And finally, since the Higgsinos also contribute, and the Higgs and Higgsinos should have similar masses, and the Higgsinos mix with the Wino and Bino to form neutralinos, we can motivate a light neutralino particle.

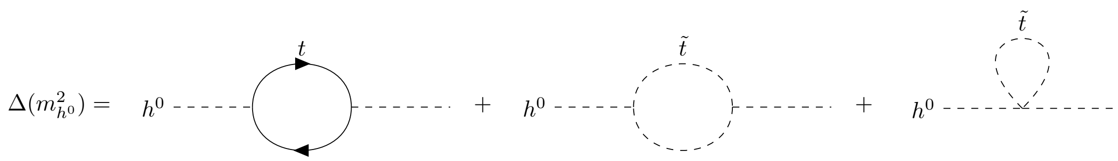


Figure 2.11: [4] An updated version of fig. 2.10 with the inclusion of the stop squark, top quark one-loop diagrams. The stop squark is a bosonic superpartner of the fermionic top quark, provides equal and opposite contribution to the top quark loop, cancelling out the contribution. There are two loops because there are two bosonic partners for the top quark, a fermion with spin.

$$\Delta(m^2_{h^0}) = \frac{3}{4\pi^2}\cos^2\alpha\lambda_t^2 m_t^2\left[\ln(m_{\tilde{t}_1} m_{\tilde{t}_2}/m_t^2) + \Delta_{\text{threshold}}\right]. \tag{2.30}$$

Figure 2.12 shows the theoretical cross-sections of the supersymmetric particles at the LHC assuming a center-of-mass collision energy $\sqrt{s} = 13\,\text{TeV}$. Compared to Run-I, the rate of gluino production has increased by a factor of 50. Since naturalness [89] is a strong motivator for the gluinos ($\tilde{g}$) to have a mass around the TeV scale in order to limit their contributions to the radiative corrections to the top squark masses, also at the TeV scale, and the lightest supersymmetric partner $\tilde{\chi}_1^0$ is also motivated to be light as well, one expects these particles to be produced copiously during Run 2 operation of the LHC at 13 TeV. For these reasons, the search for gluino production with decays via top squarks is a highly motivated search

[22]Warning: slightly heuristic argument ahead.

to perform. In section 2.2.2, I introduce the simplified SUSY model that is the crux of the search presented in this thesis.

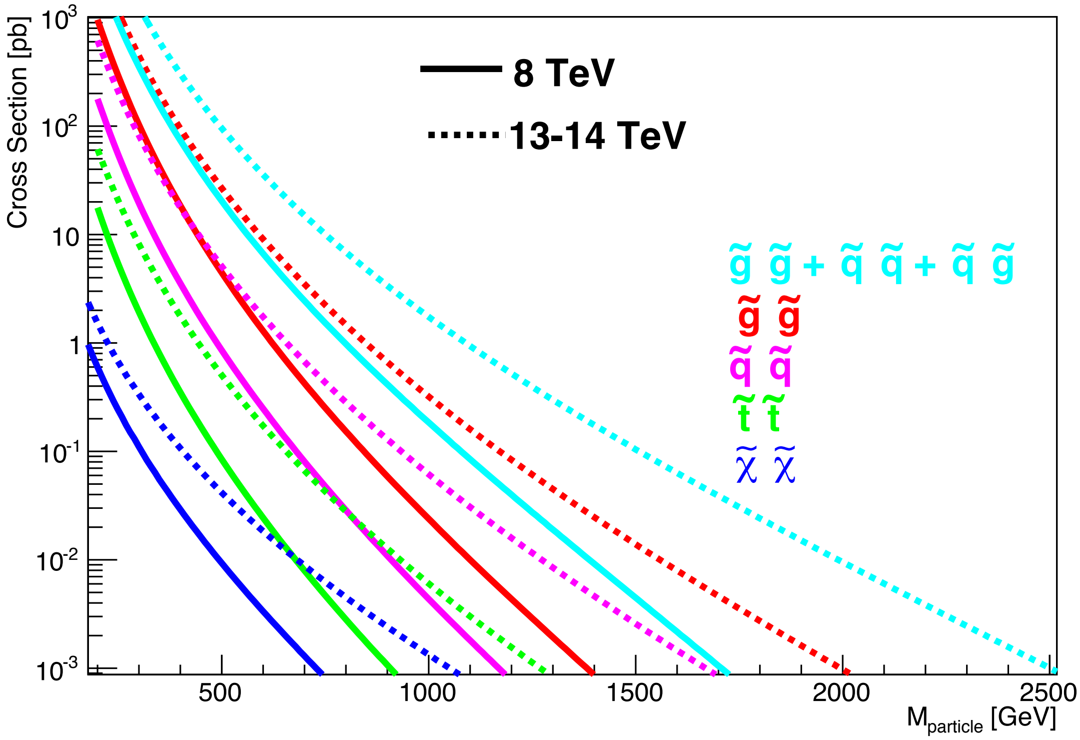
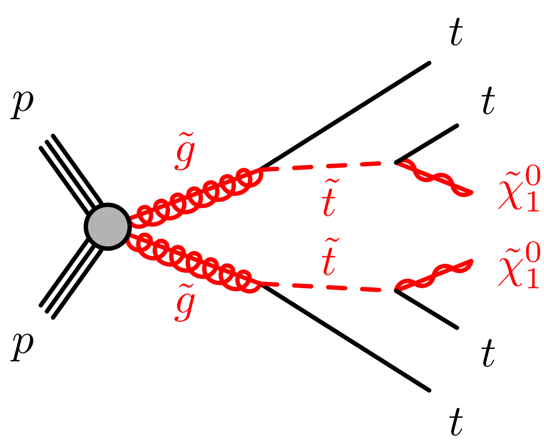


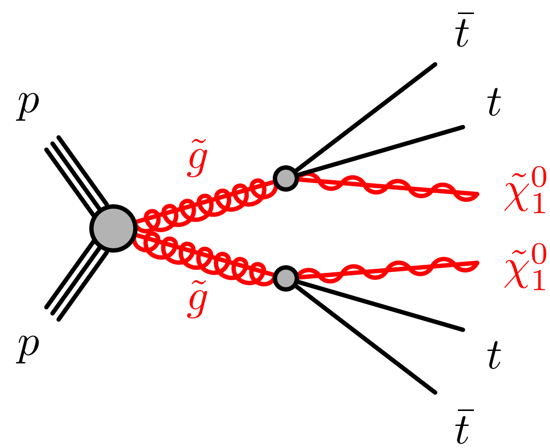
Figure 2.12: [90, 91] Theoretical cross-sections of gluino pair production are shown in the model of equal degenerate squark masses, as a function of gluino mass at $\sqrt{s} = 13\,\text{TeV}$. Gluinos, because of their strong color coupling, have the highest theoretical cross section of the sparticles that could be found at the LHC.

2.2.2 Searching for New Physics using Simplified Models

A model of new physics is defined by a TeV-scale effective Lagrangian describing its particles and their interactions. The efforts so far have focused around motivating the lightness of stops, gluinos, and neutralinos; but have we neglected the other sparticles? For the sake of experimental physicists, a simplified model is generally a limit of a more general new-physics scenario with all but a few particles integrated out [92, 93, 94] by setting them to very large mass scales. Simplified models are useful to provide topology-based limits on searches to identify the boundaries of search sensitivity and derive limits on more general models by reinterpreting [95] the limits in the context of a different signal topology. Another particular reason a simplified model helps in the search for new physics is to understand the sensitivity of the detector technology and analysis design. Experimentalists and theorists alike can identify kinematic ranges for which existing searches are not efficient or sensitive, and then define new search strategies to attempt to cover the gaps in the exploration of phase-space. An example of two simplified models are shown in fig. 2.13 for gluino production with final states consisting of four top quarks and a large missing transverse energy[23] from the neutralinos $\tilde{\chi}_1^0$. The 1st and 2nd generation squarks are assumed to be much larger than the gluino mass. The gluino can produce $t\bar{t} + \tilde{\chi}_1^0$ by decaying either off-shell through a heavier stop squark or on-shell through a lighter stop squark.



(a) on-shell



(b) off-shell

Figure 2.13: The decay topology of the simplified model for $\tilde{g} \to \bar{t}\tilde{t}_1 \to t\bar{t}\tilde{\chi}_1^0$ for both (a) on-shell and (b) off-shell stops, $\tilde{t}_1$. The difference between the two diagrams is that in the off-shell diagram, the stops are integrated out of the simplified model by setting the mass of the stop to 5 TeV.

These simplified models can be parameterized allowing for projected views in phase-space.

[23]This is an assumption, for R-parity conserving (RPC) scenarios such that $\tilde{\chi}_1^0$ is stable, does not decay, and escapes the detector unseen. In R-parity violating (RPV) scenarios [83, 96, 97, 98], the lightest supersymmetric particle is unstable and decays to SM particles.

The cross-section of gluinos (fig. 2.12), $\sigma(pp \to \tilde{g}\tilde{g} \to X)$ is one parameter. In both the on-shell and off-shell models in fig. 2.13, there are two parameters for the gluino mass $m_{\tilde{g}}$ and the neutralino mass $m_{\tilde{\chi}_1^0}$. In the on-shell model in fig. 2.13a, there is an extra parameter for the mass of the stop squark $m_{\tilde{t}_1}$, but the off-shell model does not have this parameter, setting the mass of the stop squark $m_{\tilde{t}_1} = 5\,\mathrm{TeV}$. Finally, the branching ratio for $\tilde{g}$ to decay to $\tilde{t}_1\ \bar{t}$ is assumed to be 100% in this simplified model to reduce the number of parameters. This is clearly not physical (if we find SUSY) but models with multiple decay modes can be studied by taking linear combinations of the results of simplified models for 100% branching ratios.

At the end of the day, one needs to remember that all of these are theories and we, as experimentalists, make many assumptions to simplify the theories into a set of reduced observables to search for. In particular, the search presented in this analysis has assumed that the $\tilde{t}$ have higher masses than the rest of the squarks, but it could be possible that the $\tilde{t}$ has a lower mass. If there is indeed a sign of new physics observed through the search designed around a simplified model, further studies and analysis reinterpretations need to be performed to determine what that new physics is. In the next chapter, I'll discuss how we can leverage the world's most powerful collider to search for new physics.

Glossary

ATLAS a general-purpose detector at the LHC. 1, 2, 19

BSM Beyond the Standard Model. 25

LHC Large Hadron Collider. 1, 19, 21, 22, 31, 32

LO leading order. 16

LSP lightest supersymmetric particle. 26

NLO next-to-leading-order. 2

NNLO next-to-next-to-leading-order. 2

PDF parton distribution function. 22, 23

QCD Quantum Chromodynamics. A theory describing the strong interactions of SM particles.. 3, 14, 16, 17, 25

QED Quantum Electrodynamics. A theory describing the electromagnetic and weak interactions of SM particles.. 3, 9, 14, 16, 25

SM Standard Model. 1, 3, 5, 9, 10, 17, 24–29, 33

SSB Spontaneous Symmetry Breaking. 7, 8, 27

SUSY Supersymmetry. 26, 28, 29, 31, 34

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