A search for sparticles in zero lepton final states

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Submitted in partial fulfillment of the
requirements for the degree of
Doctor of Philosophy
in the Graduate School of Arts and Sciences

COLUMBIA UNIVERSITY

2016

ABSTRACT

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Dedication

Chapter 1

Introduction

Particle physics is a remarkably successful field of scientific inquiry. The ability to precisely predict the properties of a exceedingly wide range of physical phenomena, such as the description of the cosmic microwave background [1, 2], the understanding of the anomalous magnetic dipole moment of the electron [3, 4], and the measurement of the number of weakly-interacting neutrino flavors [5] is truly amazing.

The theory that has allowed this range of predictions is the *Standard Model* of particle physics (SM). The Standard Model combines the electroweak theory of Glashow, Weinberg, and Salam [6–8] with the theory of the strong interactions, as first envisioned by Gell-Mann and Zweig [9, 10]. This quantum field theory (QFT) contains a tiny number of particles, whose interactions describe phenomena up to at least the TeV scale. These particles are manifestations of the fields of the Standard Model, after application of the Higgs Mechanism. The particle content of the SM consists only of the six quarks, the six leptons, the four gauge bosons, and the scalar Higgs boson.

Despite its impressive range of described phenomena, the Standard Model has some theoretical and experimental deficiencies. The SM contains 26 free parameters [1]. It would be more theoretically pleasing to understand these free parameters in terms of a more fundamental theory. The major theoretical concern of the Standard Model, as it pertains to this thesis, is the *hierachy problem*[11–15]. The light mass

[1] This is the Standard Model corrected to include neutrino masses. These parameters are the fermion masses (6 leptons, 6 quarks), CKM and PMNS mixing angles (8 angles, 2 CP-violating phases), W/Z/Higgs masses (3), the Higgs field expectation value, and the couplings of the strong, weak, and electromagnetic forces (3 α_{force}) .

of the Higgs boson (125 GeV) should be quadratically dependent on the scale of UV physics, due to the quantum corrections from high-energy physics processes. The most perplexing experimental issue is the existence of *dark matter*, as demonstrated by galactic rotation curves [16–22]. This data has shown that there exists additional matter which has not yet been seen interacting with the particles of the Standard Model. There is no particle in the SM which can act as a candidate for dark matter.

Both of these major issues, as well as numerous others, can be solved by the introduction of *supersymmetry* (SUSY) [15, 23–35]. In supersymmetric theories, each SM particles has a so-called *superpartner*, or sparticle partner, differing from given SM particle by 1/2 in spin. These theories solve the hierachy problem, since the quantum corrections induced from the superpartners exactly cancel those induced by the SM particles. In addition, these theories are usually constructed assuming $R-$parity, which can be thought of as the "charge" of supersymmetry, with SM particles having $R = 1$ and sparticles having $R = -1$. In collider experiments, since the incoming SM particles have total $R = 1$, the resulting sparticles are produced in pairs. This produces a rich phenomenology, which is characterized by significant hadronic activity and large missing transverse energy ($E_{\mathrm{T}}^{\mathrm{miss}}$), which provide significant discrimination against SM backgrounds [36].

Despite the power of searches for supersymmetry where $E_{\mathrm{T}}^{\mathrm{miss}}$ is a primary discriminating variable, there has been significant interest in the use of other variables to discriminate against SM backgrounds. These include searches employing variables such as αT, $M_{T,2}$, and the razor variables (M_R, R^2) [37–47]. In this thesis, we will present the first search for supersymmetry using the novel Recursive Jigsaw Reconstruction (RJR) technique. RJR can be considered the conceptual successor of the razor variables. We impose a particular final state "decay tree" on an events, which roughly corresponds to a simplified Feynmann diagram in decays containing weakly-interacting particles. We account for the missing degrees of freedom associated to the weakly-interacting particles by a series of simplifying assumptions, which allow us to calculate our variables of interest at each step in the decay tree. This allows an unprecedented understanding of the internal structure of the decay and the ability to construct additional variables to reject Standard Model backgrounds.

This thesis details a search for the superpartners of the gluon and quarks, the gluino and squarks, in final states with zero leptons, with 13.3 $\mathrm{fb}^{-1}$of data using the ATLAS detector. We organzie the thesis as follows. The theoretical foundations of the Standard Model and supersymmetry are described in Chapters 2 and 3. The Large Hadron Collider and the ATLAS detector are presented in Chapters 4 and 5. Chapter 5 provides a detailed description of Recursive Jigsaw Reconstruction and a description of the variables used for the particular search presented in this thesis. Chapter 6 presents the details of the analysis, including details of the dataset, object reconstruction, and selections used. In Chapter 7, the final results are presented; since there is no evidence of a supersymmetric signal in the analysis, we present the final exclusion curves in simplified supersymmetric models.

Chapter 2

The Standard Model

2.1 Overview

A Standard Model is another name for a theory of the internal symmetry group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, with its associated set of parameters. *The* Standard Model refers specifically to a Standard Model with the proper parameters to describe the universe. The SM is the culmination of years of work in both theoretical and experimental particle physics. In this thesis, we take the view that theorists cite construct a model with the field content and symmetries as inputs, and write down the most general Lagrangian consistent with those symmetries. Assuming this model is compatible with nature (in particular, the predictions of the model are consistent with previous experiments), experimentalists are responsible measuring the parameters of this model This will be applicable for this chapter and the following one.

Additional theoretical background is in 9.6. The philosophy and notations are inspired by [48, 49].

2.2 Field Content

The Standard Model field content is

$$\begin{aligned}
\text{Fermions} \;&: Q_L(3,2)_{+1/3}, U_R(3,1)_{+4/3}, D_R(3,1)_{-2/3}, L_L(1,2)_{-1}, E_R(1,1)_{-2} \\
\text{Scalar (Higgs)} \;&: \phi(1,2)_{+1} \\
\text{Vector Fields} \;&: G^\mu(8,1)_0, W^\mu(1,3)_0, B^\mu(1,1)_0
\end{aligned} \tag{2.1}$$

where the $(A, B)_Y$ notation represents the irreducible representation under $SU(3)$ and $SU(2)$, with Y being the electroweak hypercharge. Each of these fermion fields has an additional index, representing the three generation of fermions.

We observed that Q_L, U_R, and D_R are triplets under $SU(3)_C$; these are the *quark* fields. The *color* group, $SU(3)_C$ is mediated by the *gluon* field $G^\mu(8,1)_0$, which has 8 degrees of freedom. The fermion fields $L_L(1,2)_{-1}$ and $E_R(1,1)_{-2}$ are singlets under $SU(3)_C$; we call them the *lepton* fields.

Next, we note the "left-handed" ("right-handed") fermion fields, denoted by L (R) subscript, The left-handed fields form doublets under $SU(2)_L$. These are mediated by the three degrees of freedom of the "W" fields $W^\mu(1,3)_0$. These fields only act on the left-handed particles of the Standard Model. This is the reflection of the "chirality" of the Standard Model; the left-handed and right-handed particles are treated differently by the electroweak forces. The right-handed fields, U_R, D_R, and E_R, are singlets under $SU(2)_L$.

The $U(1)_Y$ symmetry is associated to the $B^\mu(1,1)_0$ boson with one degree of freedom. The charge Y is known as the electroweak hypercharge.

To better understand the phenomenology of the Standard Model, let us investigate each of the *sectors* of the Standard Model separately.

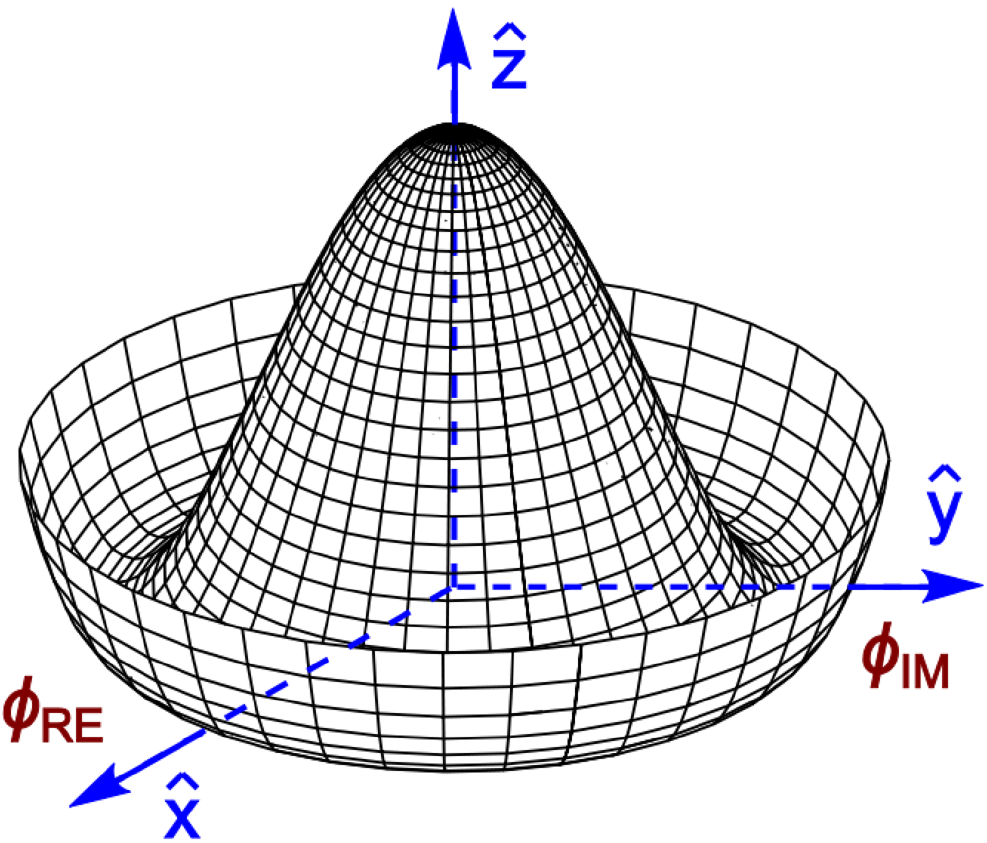
Electroweak sector

The electroweak sector refers to the $SU(2)_L \otimes U(1)_Y$ portion of the Standard Model gauge group. Following our philosophy of writing all gauge-invariant and renormalizable terms, the electroweak Lagrangian can be written as

$$\mathcal{L} = W_a^{\mu\nu} W_{\mu\nu}^a + B^{\mu\nu} B_{\mu\nu} + (D^\mu \phi)^\dagger D_\mu \phi - \mu^2 \phi^\dagger \phi - \lambda(\phi^\dagger \phi)^2. \tag{2.2}$$

where $W_a^{\mu\nu}$ are the three ($a = 1, 2, 3$) gauge bosons associated to the $SU(2)_L$ gauge group, $B^{\mu\nu}$ is the one gauge boson of the $U(1)_Y$ gauge group, and ϕ is the complex

Figure 2.1: Sombrero potential



Higgs multiplet. The covariant derivative D^μ is given by

$$D^\mu = \partial^\mu + \frac{ig}{2} W_a^\mu \sigma_a + \frac{ig'}{2} B^\mu \tag{2.3}$$

where $i\sigma_a$ are the Pauli matrices times the imaginary constant, which are the generators for $SU(2)_L$, and g and g' are the $SU(2)_L$ and $U(1)_Y$ coupling constants, respectively. The field strength tensors $W_a^{\mu\nu}$ and $B^{\mu\nu}$ are given by the commutator of the covariant derivative associated to each field

$$\begin{aligned} B^{\mu\nu} &= \partial^\mu B^\nu - \partial^\nu B^\mu \\ W_a^{\mu\nu} &= \partial^\mu W_a^\nu - \partial^\nu W_a^\mu - g\epsilon_{abc} W_a^\mu W_b^\nu, \qquad\qquad i = 1, 2, 3 \end{aligned} \tag{2.4}$$

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The terms in the Lagrangian 2.2 proportional to μ^2 and λ make up the "Higgs potential" [50]. As normal (see Appendix 9.6), we restrict $\lambda > 0$ to guarantee our potential is bounded from below, and we also require $\mu^2 < 0$, which gives us the standard "sombrero" potential shown in 2.1.

This potential has infinitely many minima at $<\phi>=\sqrt{2m/\lambda}$; the ground state is *spontaneously* broken by the choice of ground state, which induces a vacuum expection value (VEV). Without loss of generality, we can choose the Higgs field ϕ to point in the real direction, and write the Higgs field ϕ in the following form :

$$\phi = \frac{1}{\sqrt{2}} \exp(\frac{i}{v}\sigma_a\theta_a) \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \tag{2.5}$$

We choose a gauge to rotate away the dependence on θ_a, such that we can write simply

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \tag{2.6}$$

Now, we can see how the masses of the vector bosons are generated from the application of the Higgs mechanism. We plug Eq.2.6 back into the electroweak Lagrangian, and only showing the relevant mass terms in the vacuum state where $h(x) = 0$ see that (dropping the Lorentz indices) :

$$\begin{aligned} \mathcal{L}_M &= \frac{1}{8} \left| \begin{pmatrix} gW_3 + g'B & g(W_1 - iW_2) \\ g(W_1 + iW_2) & -gW_3 + g'B \end{pmatrix} \begin{pmatrix} 0 \\ v \end{pmatrix} \right|^2 \\ &= \frac{g^2v^2}{8} \left[W_1^2 + W_2^2 + (\tfrac{g'}{g}B - W_3)^2 \right] \end{aligned} \tag{2.7}$$

Defining the *Weinberg* angle $\tan(\theta_W) = g'/g$ and the following *physical* fields :

$$\begin{aligned} W^{\pm} &= \frac{1}{\sqrt{2}}(W_1 \mp iW_2) \\ Z^0 &= \cos\theta_W W_3 - \sin\theta_W B \\ A^0 &= \sin\theta_W W_3 + \cos\theta_W B \end{aligned} \tag{2.8}$$

we can write the piece of the Lagrangian associated to the vector boson masses as

$$\mathcal{L}_{M_V} = \frac{1}{4}g^2v^2W^+W^- + \frac{1}{8}(g^2 + g'^2)v^2Z^0Z^0. \tag{2.9}$$

and we have the following values of the masses for the vector bosons :

$$\begin{aligned} m_W^2 &= \frac{1}{4}v^2g^2 \\ m_Z^2 &= \frac{1}{4}v^2(g^2 + g'^2) \\ m_A^2 &= 0 \end{aligned} \tag{2.10}$$

We thus see how the Higgs mechnanism gives rise to the masses of the $W^\pm$ and Z boson in the Standard Model; the mass of the photon is zero, as expected. The $SU(2)_L \otimes U(1)_Y$ symmetry of the initially massless $W_{1,2,3}$ and B fields is broken to the $U(1)_{EM}$. Of the four degrees of freedom in the complex Higgs doublet, three are "eaten" when we give mass to the $W^\pm$ and Z_0, while the other degree of freedom is the Higgs particle, as found in 2012 by the ATLAS and CMS collaborations [51, 52].

Quantum Chromodynamics

Quantum chromodynamics (or the theory of the *strong* force) characterizes the behavior of *colored* particles, collectively known as *partons*. The partons of the Standard Model are the (fermionic) quarks, and the (bosonic) gluons. The strong force is governed by $SU(3)_C$, an unbroken symmetry in the Standard Model, which implies the gluon remains massless. Defining the covariant derivative for QCD as

$$D^\mu = \partial^\mu + ig_sG_a^\mu L_a, a = 1, ..., 8 \tag{2.11}$$

where L_a are the generators of $SU(3)_C$, and g_s is the coupling constant of the strong force. The QCD Lagrangian then is given by

$$\mathcal{L}_{\text{QCD}} = i\bar{\psi}_f D_\mu \gamma^\mu \psi_f - \frac{1}{4}G_{a,\mu\nu}G_a^{\mu\nu} \tag{2.12}$$

where the summation over f is for quarks *families*, and $G_a^{\mu\nu}$ is the gluon field strength tensor, given by

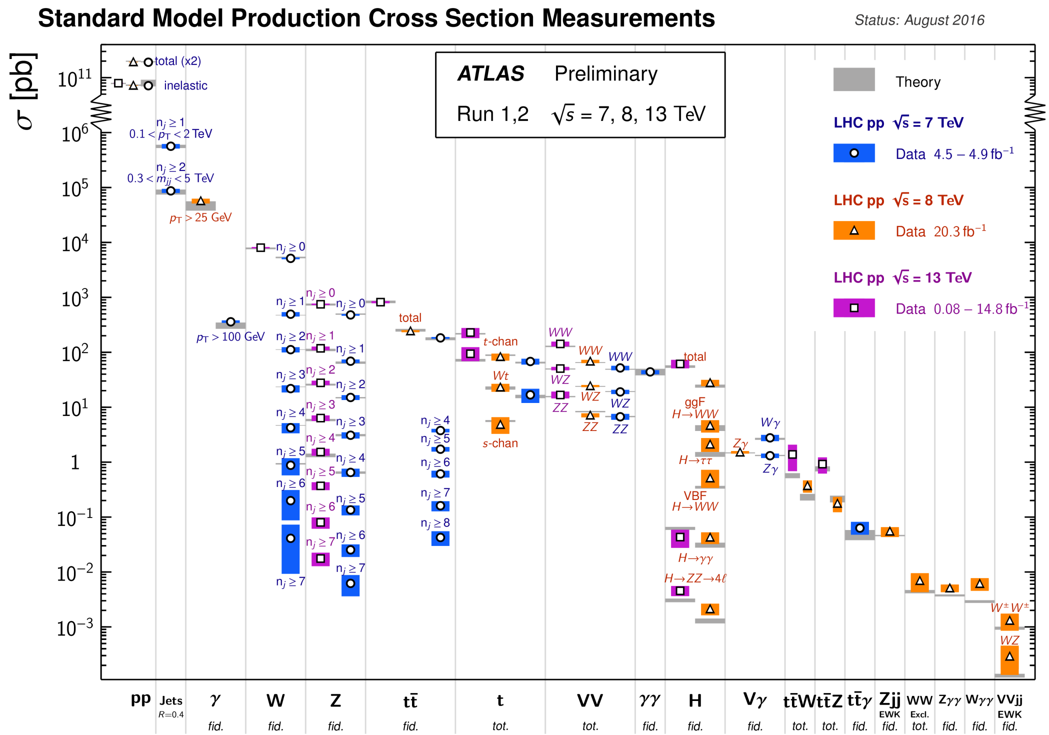
$$G_a^{\mu\nu} = \partial^\mu G_a^\nu - \partial^\nu G_a^\mu - g_s f^{abc} G_b^\mu G_c^\nu, a, b, c = 1, ..., 8 \tag{2.13}$$

where f^{abc} are the structure constants of $SU(3)_C$, which are analogous to ϵ_{abc} for $SU(2)_L$. The kinetic term for the quarks is contained in the standard ∂_μ term, while the field strength term contains the interactions between the quarks and gluons, as well as the gluon self-interactions.

Written down in this simple form, the QCD Lagrangian does not seem much different from the QED Lagrangian, with the proper adjustments for the different group structures. The gluon is massless, like the photon, so one could näively expect an infinite range force, and it pays to understand why this is not the case. The reason for this fundamental difference is the gluon self-interactions arising in the field strength tensor term of the Lagrangian. This leads to the phenomena of *color confinement*, which describes how one only observes color-neutral particles alone in nature. In contrast to the electromagnetic force, particles which interact via the strong force experience a *greater* force as the distance between the particles increases. At long distances, the potential is given by $V(r) = -kr$. At some point, it is more energetically favorable to create additional partons out of the vacuum than continue pulling apart the existing partons, and the colored particles undergo *fragmentation*. This leads to *hadronization*. Bare quarks and gluons are actually observed as sprays of hadrons (primarly kaons and pions); these sprays are known as *jets*, which are what are observed by experiments.

It is important to recognize the importance of understanding these QCD interactions in high-energy hadron colliders such as the LHC. Since protons are hadrons, proton-proton collisions such as those produced by the LHC are primarily governed by the processes of QCD. In particular, by far the most frequent process observed in LHC experiments is dijet production from gluon-gluon interactions (see Fig.2.2). These

Figure 2.2: Cross-sections of various Standard Model processes



gluons that interact are part of the *sea* particles inside the proton; the simple $p = uud$ model does not apply. The main *valence uud* quarks are constantly interacting via gluons, which can themselves radiate gluons or split into quarks, and so on. A more useful understanding is given by the colloquially-known *bag* model [53, 54], where the proton is seen as a "bag" of (in principle) infinitely many partons, each with energy $E < \sqrt{s} = 6.5$ TeV. One then collides this (proton) bag with another, and views the products of this very complicated collision, where calculations include many loops in nonpertubative QCD calculations.

Fortunately, we are generally saved by the QCD factorization theorems [55]. This allows one to understand the hard (i.e. short distance or high energy) $2 \to 2$ parton process using the tools of perturbative QCD, while making series of approximations known as a *parton shower* model to understand the additional corrections from nonpertubative QCD. We will discuss the reconstruction of jets by experiments in Ch.5.

Fermions

We will now look more closely at the fermions in the Standard Model [56].

As noted earlier in Sec.2.2, the fermions of the Standard Model can be first distinguished between those that interact via the strong force (quarks) and those which do not (leptons).

There are six leptons in the Standard Model, which can be placed into three *generations*.

$$\begin{pmatrix} e \\ \nu_e \end{pmatrix}, \begin{pmatrix} \mu \\ \nu_\mu \end{pmatrix}, \begin{pmatrix} \tau \\ \nu_\tau \end{pmatrix} \tag{2.14}$$

There is the electron (e), muon (μ), and tau (τ), each of which has an associated neutrino (ν_e, ν_μ, ν_τ). Each of the so-called charged ("electron-like") leptons has electromagnetic charge -1, while the neutrinos all have $q_{EM} = 0$.

Often in an experimental context, lepton is used to denote the stable electron and metastable muon, due to their striking experimental signatures. Taus are often treated separately, due to their much shorter lifetime of $\tau_\tau \sim 10^{-13}s$; these decay through hadrons or the other leptons, so often physics analyses at the LHC treat them as jets or leptons, as will be done in this thesis.

As the neutrinos are electrically neutral, nearly massless, and only interact via the weak force, it is quite difficult to observe them directly. Since LHC experiments rely overwhelmingly on electromagnetic interactions to observe particles, the presence of neutrinos is not observed directly. Neutrinos are instead observed by the conservation of four-momentum in the plane transverse to the proton-proton collisions, known as *missing transverse energy*.

There are six quarks in the Standard Model : up, down, charm, strange, top, and bottom. Quarks are similar organized into three generations :

$$\begin{pmatrix} u \\ d \end{pmatrix}, \begin{pmatrix} c \\ s \end{pmatrix}, \begin{pmatrix} t \\ b \end{pmatrix} \tag{2.15}$$

where we speak of “up-like” quarks and “down-like” quarks.

Each up-like quark has charge $q_{up} = 2/3$, while the down-like quarks have $q_{down} = -1/3$. At the high energies of the LHC, one often makes the distinction between the light quarks (u, d, c, s), the bottom quark, and top quark. In general, due to the hadronization process described above, the light quarks, with masses $m_q <\sim 1.5 GeV$ are indistinguishable by LHC experiments. Their hadronic decay products generally have long lifetimes and they are reconstructed as jets.[1]. The bottom quark hadronizes primarily through the B-mesons, which generally travels a short distance before decaying to other hadrons. This allows one to distinguish decays via $b-$quarks form other jets; this procedure is known as *b-tagging* and will be discussed more in Ch.5. Due to its large mass, the top quark decays before it can hadronize; there are no bound states associated to the top quark. The top is of particular interest at the LHC; it has a striking signature through its most common decay mode $t \to Wb$. Decays via tops, especially $t\bar{t}$ are frequently an important signal decay mode, or an important background process.

Interactions in the Standard Model

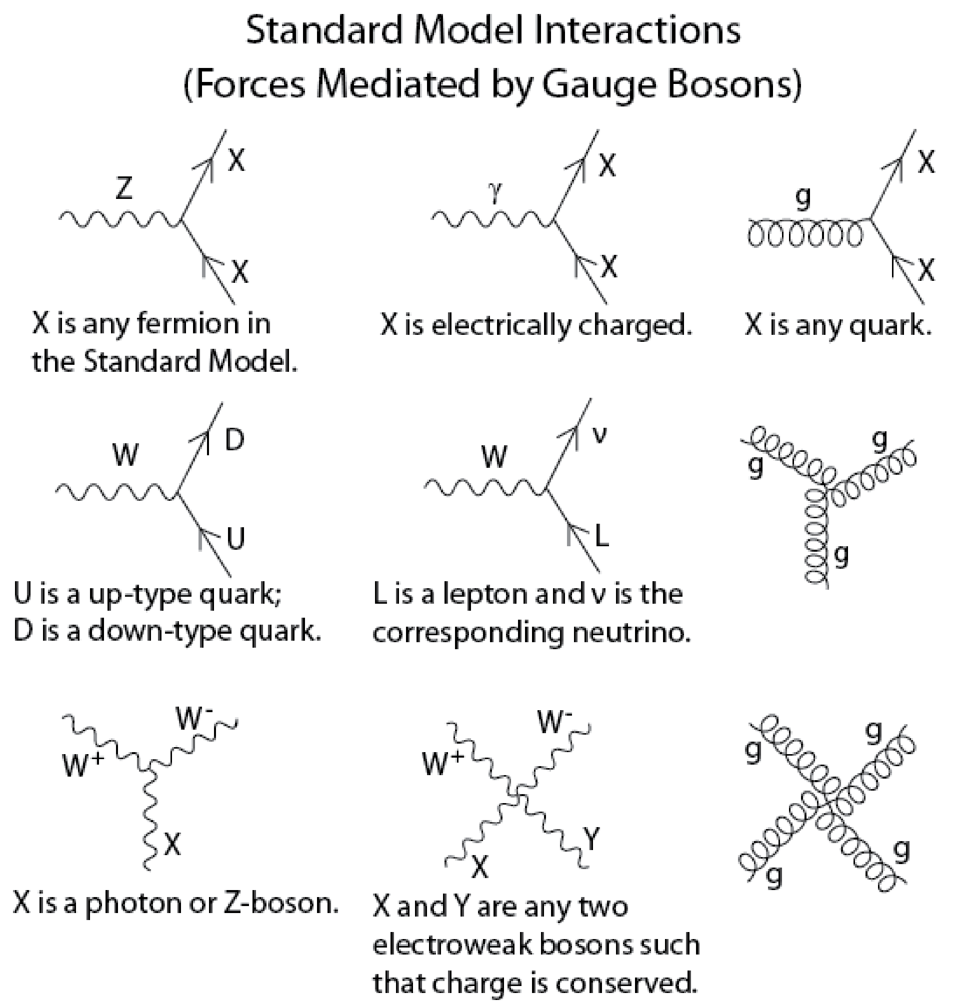
We briefly overview the entirety of the fundamental interactions of the Standard Model; these can also be found in 2.3.

The electromagnetic force, mediated by the photon, interacts with via a three-point coupling all charged particles in the Standard Model. The photon thus interacts with all the quarks, the charged leptons, and the charged $W^{\pm}$ bosons.

The weak force is mediated by three particles : the $W^{\pm}$ and the Z^0. The Z^0 can interacts with all fermions via a three-point coupling. A real Z_0 can thus decay to a fermion-antifermion pair of all SM fermions except the top quark, due to its large

[1] In some contexts, charm quarks are also treated as a separate category, although it is quite difficult to distinguish charm quarks from the other light quarks.

Figure 2.3: The interactions of the Standard Model



mass. The $W^{\pm}$ has two important three-point interactions with fermions. First, the $W^{\pm}$ can interact with an up-like quark and a down-like quark; an important example in LHC experiments is $t \to Wb$ The coupling constants for these interactions are encoded in the unitary matrix known as the Cabibbo–Kobayashi–Maskawa (CKM) matrix [57, 58], and are generally known as flavor-changing interactions. Secondly, the $W^{\pm}$ interacts with a charged lepton and its corresponding neutrino. In this case, the unitary matrix that corresponds to CKM matrix for quarks is the identity matrix, which forbids (fundamental) vertices such as $\mu \to We$. For leptons, instead this is a two-step process : $\mu \to \nu_m uW \to \nu_m u\bar{\nu}_e e$. Finally, there are the self-interactions of the weak gauge bosons. There is a three-point and four-point interaction; all combinations are allowed which conserve electric charge.

The strong force is mediated by the gluon, which as discussed above also carries the strong color charge. There is the fundamental three-point interaction, where a quark radiates a gluon. Additionally, there are the three-point and four-point gluon-only interactions.

2.3 Deficiencies of the Standard Model

At this point, it is quite easy to simply rest on our laurels. This relatively simple theory is capable of explaining a very wide range of phenomenom, which ultimately break down only to combinations of nine diagrams shown in Fig.2.3. Unfortunately, there are some unexplained problems with the Standard Model. We cannot go through all of the potential issues in this thesis, but we will motivate the primary issues which naturally lead one to *supersymmetry*, as we will see in Ch.3.

The Standard Model has many free paramaters; see Table 2.1 In general, we prefer models with less free parameters. A great example of this fact, and the primary experimental evidence for EWSB, is the relationship between the couplings of the weak force and the masses of the gauge bosons of the weak force :

$$\rho \equiv \frac{m_W^2}{m_Z^2 \cos^2\theta_W} \stackrel{?}{=} 1 \tag{2.16}$$

where ? indicates that this is a testable prediction of the Standard Model (in particular, that the gauge bosons gain mass through EWSB). This relationship has been measured within experimental and theoretical predictions. We would like to produce additional such relationships, which would exist if the Standard Model is a low-energy approximation of some other theory.

An additional issue is the lack of *gauge coupling unification*. The couplings of any quantum field theory "run" as a function of the distance scales (or inversely,

Table 2.1: Parameters of the Standard Model. For values dependent on the renormalization scheme, we use a combination of the on-shell normalization scheme [59–62] and modified minimal subtraction scheme with $m_{\bar{M}S}$ as indicated in the table[63]

.

m_e	Electron mass	511 keV
m_μ	Muon mass	105.7 MeV
m_τ	Tau mass	1.78 GeV
m_u	Up quark mass	1.9 MeV ($m_{\bar{M}S} = 2GeV$)
m_d	Down quark mass	4.4 MeV ($m_{\bar{M}S} = 2GeV$)
m_s	Strange quark mass	87 MeV ($m_{\bar{M}S} = 2GeV$)
m_c	Charm quark mass	1.32 GeV ($m_{\bar{M}S} = m_c$)
m_b	Bottom quark mass	4.24 GeV ($m_{\bar{M}S} = m_b$)
m_t	Top quark mass	172.7 GeV (on-shell renormalization)
θ_{12} CKM	12-mixing angle	13.1°
θ_{23} CKM	23-mixing angle	2.4°
θ_{13} CKM	13-mixing angle	0.2°
δ CKM	CP-violating Phase	0.995
g'	U(1) gauge coupling	0.357 ($m_{\bar{M}S} = m_Z$)
g	SU(2) gauge coupling	0.652 ($m_{\bar{M}S} = m_Z$)
g_s	SU(3) gauge coupling	1.221 ($m_{\bar{M}S} = m_Z$)
θQCD	QCD vacuum angle	∼0
VEV	Higgs vacuum expectation value	246 GeV
m_H	Higgs mass	125 GeV

energy scales) of the theory. The idea is closely related to the unification of the electromagnetic and weak forces at the so-called *electroweak scale* of $O(100 \text{ GeV})$. One would hope this behavior was repeated between the electroweak forces and the strong force at some suitable energy scale. The Standard Model does automatically not exhibit this behavior, as we can see in Fig.2.4.

The most significant problem with the Standard Model is the *hierarchy problem*. In its most straightforward incarnation, the Higgs scalar field is subject to quantum corrections through loop diagrams, as shown in Fig.2.5. For demonstration, we use the contributions from the top quark, since the top quark has the largest Higgs Yukawa coupling due to its large mass. In general, we should expect these corrections to quadratically dependent on the scale of the ultraviolet physics, Λ. Briefly assume there is no new physics before the Planck scale of gravity, $\Lambda_{\text{Planck}} = 10^{19}$ GeV. In this

Figure 2.4: The running of Standard Model gauge couplings. The Standard Model couplings do not unify at high energies, which indicates it cannot completely describe nature through the Planck scale.

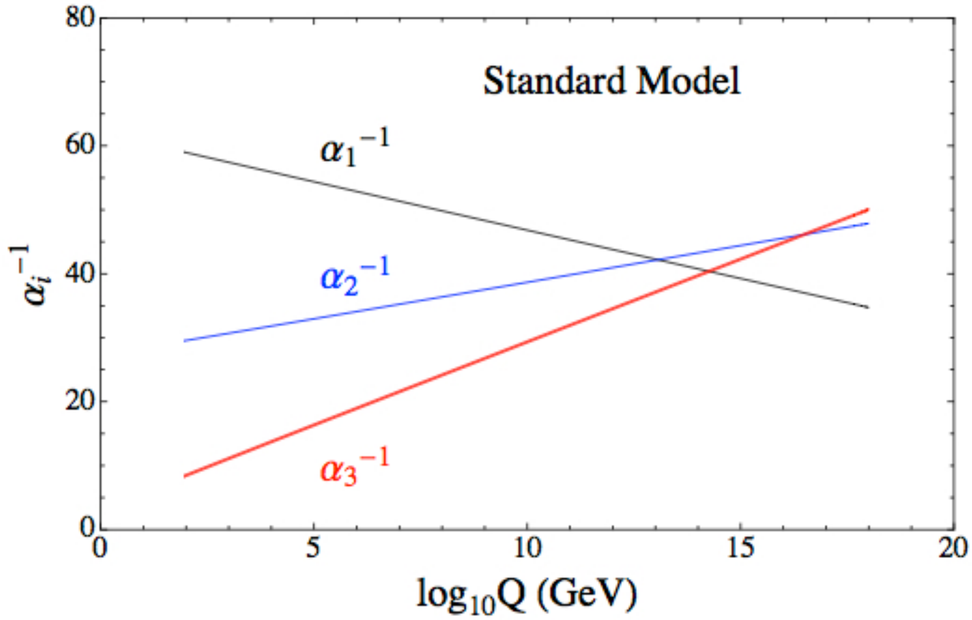
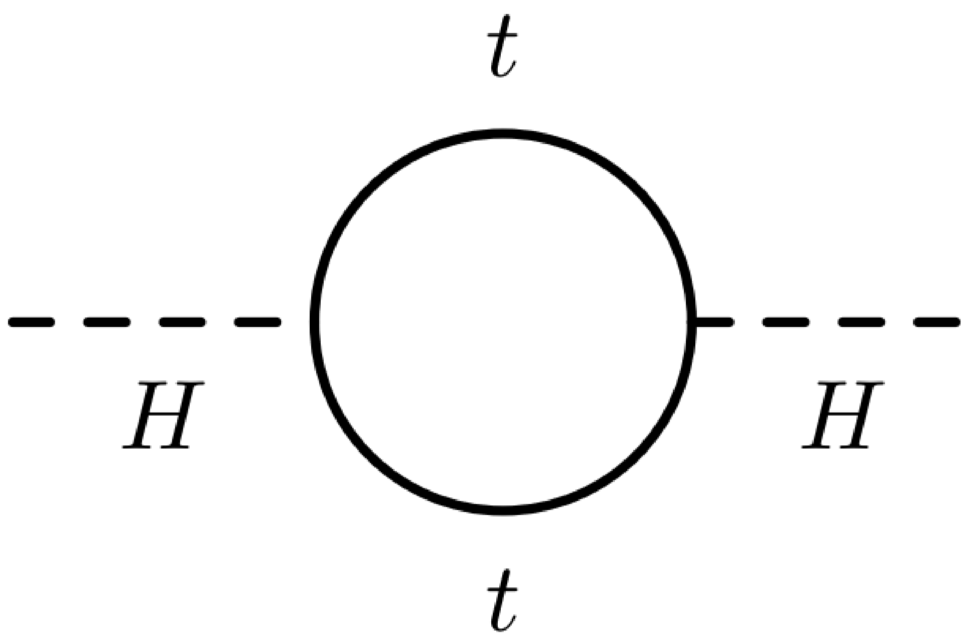


Figure 2.5: The dominant quantum loop correction to the Higgs mass in the Standard Model.



case, we expect the corrections to the Higgs mass like

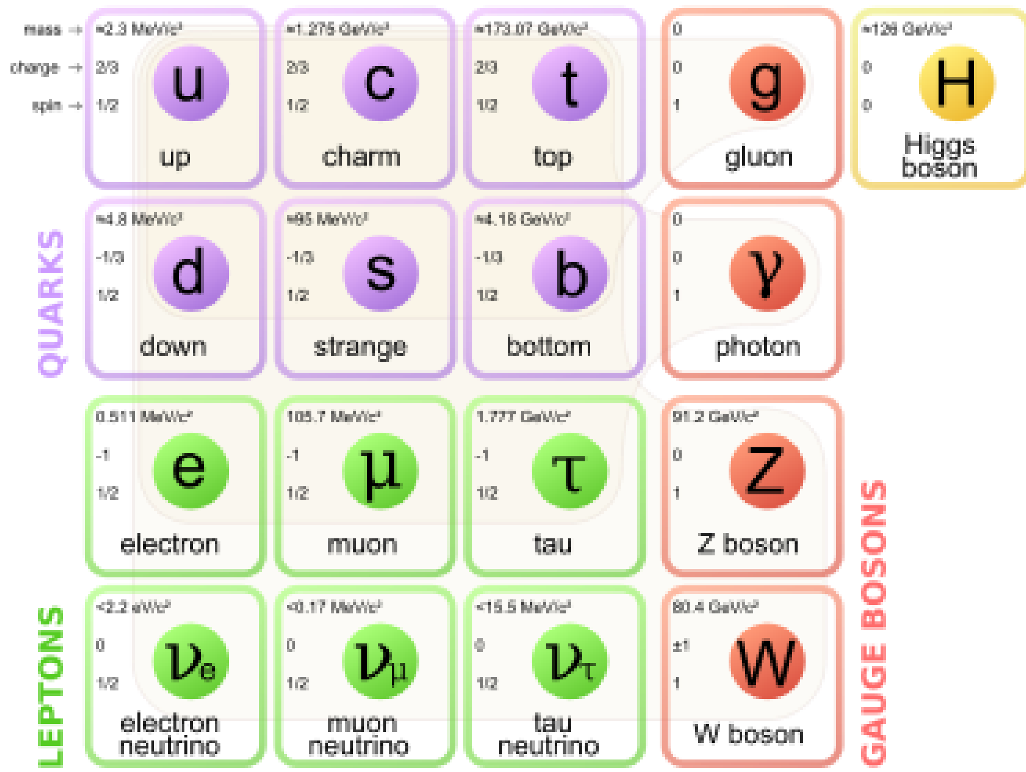
$$\delta m_H^2 \approx \left(\frac{m_t}{8\pi^2 {<}\phi{>}_{VEV}} \right)^2 \Lambda_{Planck}^2. \tag{2.17}$$

To achieve the miraculous cancellation required to get the observed Higgs mass of 125 GeV, one needs to then set the bare Higgs mass m_0, our input to the Standard Model Lagrangian, itself to a *precise* value $\sim 10^{19}$ GeV. This extraordinary level of parameter finetuning is quite undesirable, and within the framework of the Standard Model, there is little that can be done to alleviate this issue.

An additional concern, of a different nature, is the lack of a *dark matter* candidate in the Standard Model. Dark matter was discovered by observing galactic rotation curves, which showed that much of the matter that interacted gravitionally was invisible to our (electromagnetic) telescopes [16–22]. The postulation of the existence of dark matter, which interacts at least through gravity, allows one to understand these galatic rotation curves. Unfortunately, no particle in the Standard Model could possibly be the dark matter particle. The only candidate truly worth another look is the neutrino, but it has been shown that the neutrino content of the universe is simply too small to explain the galatic rotation curves [22, 64]. The experimental evidence from the galactic rotations curves thus show there *must* be additional physics beyond the Standard Model, which is yet to be understood.

In the next chapter, we will see how these problems can be alleviated by the theory of supersymmetry.

Figure 2.6: Particles of the Standard Model



Chapter 3

Supersymmetry

This chapter will introduce supersymmetry (SUSY) [15, 23–35]. We will begin by introducing the concept of a *superspace*, and discuss some general ingredients of supersymmetric theories. This will include a discussion of how the problems with the Standard Model described in Ch.2 are naturally fixed by these theories.

The next step is to discuss the particle content of the *Minimally Supersymmetric Standard Model* (MSSM). As its name implies, this theory contains the minimal additional particle content to make Standard Model supersymmetric. We then discuss the important phenomonological consequences of this theory, especially as it would be observed in experiments at the LHC.

3.1 Supersymmetric theories : from space to superspace

Coleman-Mandula "no-go" theorm

We begin the theoretical motivation for supersymmetry by citing the "no-go" theorem of Coleman and Mandula [65]. This theorem forbids *spin-charge unification*; it states that all quantum field theories which contain nontrivial interactions must be a direct product of the Poincarégroup of Lorentz symmetries, the internal product from of gauge symmetries, and the discrete symmetries of parity, charge conjugation, and time reversal. The assumptions which go into building the Coleman-Mandula theorem are quite restrictive, but there is one unique way out, which has become known as *supersymmetry* [26, 66]. In particular, we must introduce a *spinorial* group generator Q. Alternatively, and equivalently, this can be viewed as the addition of anti-commuting coordinates; space plus these new anti-commuting coordinates is then called *superspace* [67]. We will not investiage this view in detail, but it is also a quite intuitive and beautiful way to construct supersymmetry[15].

Supersymmetry transformations

A *supersymmetric* transformation Q transforms a bosonic state into a fermionic state, and vice versa :

$$Q\,|\mathrm{Fermion}\rangle = |\mathrm{Boson}\rangle \tag{3.1}$$

$$Q\,|\mathrm{Boson}\rangle = |\mathrm{Fermion}\rangle \tag{3.2}$$

To ensure this relation holds, Q must be an anticommuting spinor. Additionally, since spinors are inherently complex, $Q^\dagger$ must also be a generator of the supersymmetry transformation. Since Q and $Q^\dagger$ are spinor objects (with $s = 1/2$), we can see that supersymmetry must be a spacetime symmetry. The Haag-Lopuszanski-Sohnius extension [66] of the Coleman-Mandula theorem [65] is quite restrictive about the forms of such a symmetry. Here, we simply write the (anti-) commutation relations [15] :

$$Q_\alpha, Q^\dagger_{\dot\alpha} = -2\sigma_{\alpha\dot\alpha^\mu} P_\mu \tag{3.3}$$

$$Q_\alpha, Q_{\dot\beta} = Q^\dagger_{\dot\alpha}, Q^\dagger_{\dot\beta} = 0 \tag{3.4}$$

$$[P^\mu, Q_\alpha] = [P^\mu, Q^\dagger_{\dot\alpha}] = 0 \tag{3.5}$$

Supermultiplets

In a supersymmetric theory, we organize single-particle states into irreducible representations of the supersymmetric algebra which are known as *supermultiplets*. Each supermultiplet contains a fermion state $|\mathrm{F}\rangle$ and a boson state $|\mathrm{B}\rangle$; these two states are the known as *superpartners*. These are related by some combination of Q and $Q^\dagger$, up to a spacetime transformation. Q and $Q^\dagger$ commute with the mass-squared operator $-P^2$ and the operators corresponding to the gauge transformations [15]; in particular, the gauge interactions of the Standard Model. In an unbroken supersymmetric theory, this means the states $|\mathrm{F}\rangle$ and $|\mathrm{B}\rangle$ have exactly the same mass, electromagnetic charge, electroweak isospin, and color charges. One can also prove [15] that each supermultiplet contains the exact same number of bosonic (n_B) and fermion (n_F) degrees of freedom. We now explore the possible types of supermultiples one can find in a renormalizable supersymmetric theory.

Since each supermultiplet must contain a fermion state, the simplest type of supermultiplet contains a single Weyl fermion state ($n_F = 2$) which is paired with $n_B = 2$ scalar bosonic degrees of freedom. This is most conveniently constructed as single complex scalar field. We call this construction a *scalar supermultiplet* or *chiral supermultiplet*. The second name is indicative; only chiral supermultiplets can contain fermions whose right-handed and left-handed components transform differently under the gauge interactions (as of course happens in the Standard Model).

The second type of supermultiplet we construct is known as a *gauge* supermultiplet. We take a spin-1 gauge boson (which must be massless due to the gauge symmetry, so $n_B = 2$) and pair this with a single massless Weyl spinor[1]. The gauge bosons transform as the adjoint representation of the their respective gauge groups; their fermionic partners, which are known as gauginos, must also. In particular, the left-handed and right-handed components of the gaugino fermions have the same

[1] Choosing an $s = 3/2$ massless fermion leads to nonrenormalizable interactions.

gauge transformation properties.

Excluding gravity, this is the entire list of supermultiplets which can participate in renormalizable interactions in what is known as $N = 1$ supersymmetry. This means there is only one copy of the supersymmetry generators Q and $Q^\dagger$. This is essentially the only "easy" phenomenological choice, since it is the only choice in four dimensions which allows for the chiral fermions and parity violations built into the Standard Model, and we will not look further into $N > 1$ supersymmtry in this thesis.

The primary goal, after understanding the possible structures of the multiplets above, is to fit the Standard Model particles into a multiplet, and therefore make predictions about their supersymmetric partners. We explore this in the next section.

3.2 Minimally Supersymmetric Standard Model

To construct what is known as the MSSM [**susyPrimer** , 68–71], we need a few ingredients and assumptions. First, we match the Standard Model particles with their corresponding superpartners of the MSSM. We will also introduce the naming of the superpartners (also known as *sparticles*). We discuss a very common additional restraint imposed on the MSSM, known as $R-$parity. We also discuss the concept of soft supersymmetry breaking and how it manifests itself in the MSSM.

Chiral supermultiplets

The first thing we deduce is directly from Sec.**??**. The bosonic superpartners associated to the quarks and leptons *must* be spin 0, since the quarks and leptons must be arranged in a chiral supermultiplet. This is essentially the note above, since the chiral supermultiplet is the only one which can distinguish between the left-handed and right-handed components of the Standard Model particles. The superpartners of the quarks and leptons are known as *squarks* and *sleptons*, or *sfermions* in aggregate.

(for “scalar quarks”, “scalar leptons”, and “scalar fermion”[2]). The “s-” prefix can also be added to the individual quarks i.e. *selectron*, *sneutrino*, and *stop*. The notation is to add a $\sim$ over the corresponding Standard Model particle i.e. $\tilde{e}$, the selectron is the superpartner of the electron. The two-component Weyl spinors of the Standard Model must each have their own (complex scalar) partner i.e. e_L, e_R have two distinct partners : $\tilde{e_L}, \tilde{e_R}$. As noted above, the gauge interactions of any of the sfermions are identical to those of their Standard Model partners.

Due to the scalar nature of the Higgs, it must obviously lie in a chiral supermultiplet. To avoid gauge anomolies and ensure the correct Yukawa couplings to the quarks and leptons[15], we must add additional Higgs bosons to any supersymmetric theory. In the MSSM, we have two chiral supermultiplets. The SM (SUSY) parts of the multiplets are denoted $H_u(\tilde{H}_u)$ and $H_d(\tilde{H}_d)$. Writing out H_u and H_d explicitly:

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix} \tag{3.6}$$

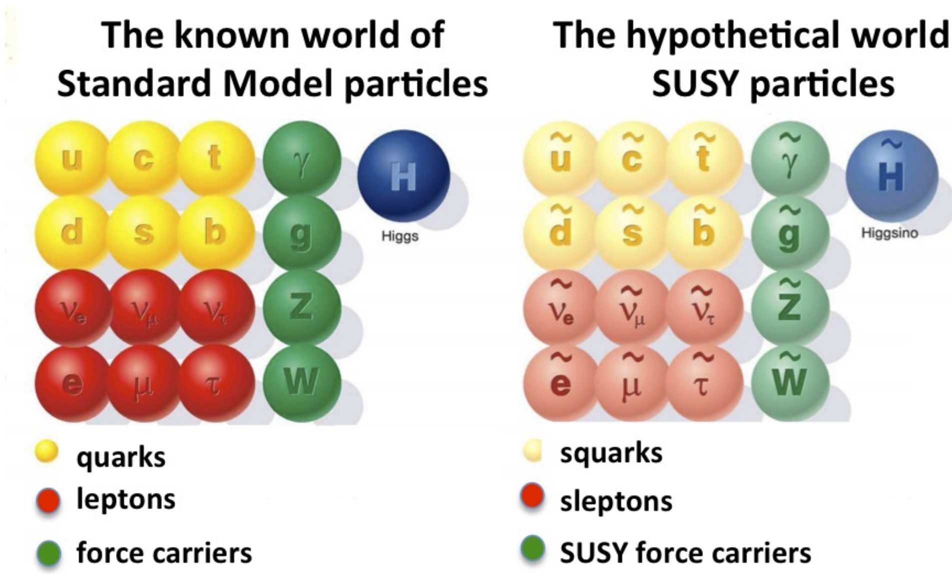
$$H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix} \tag{3.7}$$

$$\tag{3.8}$$

we see that H_u looks very similar to the SM Higgs with $Y = 1$, and H_d is symmetric to this with $+ \to -$, with $Y = -1$. The SM Higgs boson, h_0, is a linear superposition of the neutral components of these two doublets. The SUSY parts of the Higgs multiplets, $\tilde{H}_u$ and $\tilde{H}_d$, are each left-handed Weyl spinors. For generic spin-1/2 sparticles, we add the “-ino” suffix. We then call the partners of the two Higgs collectively the *Higgsinos*.

[2] The last one should probably have bigger scare quotes.

Figure 3.1: Particles of the MSSM



Gauge supermultiplets

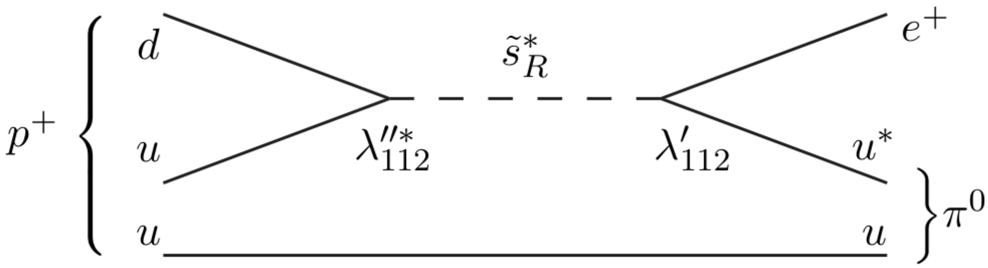
The superpartners of the gauge bosons must all be in gauge supermultiplets since they contain a spin-1 particle. Collectively, we refer to the superpartners of the gauge bosons as the gauginos.

The first gauge supermultiplet contains the gluon, and its superpartner, which is known as the *gluino*, denoted $\tilde{g}$. The gluon is of course the SM mediator of $SU(3)_C$; the gluino is also a colored particle, subject to $SU(3)_C$. From the SM before EWSB, we have the four gauge bosons of the electroweak symmetry group $SU(2)_L \otimes U(1)_Y$: $W^{1,2,3}$ and B^0. The superpartners of these particles are thus the *winos* $W^{\tilde{1},2,3}$ and *bino* $\tilde{B}^0$, where each is placed in another gauge supermultiplet with its corresponding SM particle. After EWSB, without breaking supersymmetry, we would also have the zino $\tilde{Z}^0$ and photino $\tilde{\gamma}$.

The entire particle content of the MSSM can be seen in Fig.3.1.

At this point, it's important to take a step back. Where are these particles? As stated above, supersymmetric theories require that the masses and all quantum

Figure 3.2: This Feynmann diagram shows how proton decay is induced in the MSSM, if one does not impose $R-$parity.



numbers of the SM particle and its corresponding sparticle are the same. Of course, we have not observed a selectron, squark, or wino. The answer, as it often is, is that supersymmetry is *broken* by the vacuum state of nature [15].

$R-$parity

This section is a quick aside to the general story. $R-parity$ refers to an additional discrete symmetry which is often imposed on supersymmetric models. For a given particle state, we define

$$R = (-1)^{3(B-L)+2s} \tag{3.9}$$

where B, L is the baryon (lepton) number and s is the spin. The imposition of this symmetry forbids certain terms from the MSSM Lagrangian that would violate baryon and/or lepton number. This is required in order to prevent proton decay, as shown in Fig.3.2[3]. .

In supersymmetric models, this is a $\mathbb{Z}_2$ symmetry, where SM particles have $R = 1$ and sparticles have $R = -1$. We will take $R-parity$ as part of the definition of the MSSM. We will discuss later the *drastic* consequences of this symmetry on SUSY phenomenology

[3]Proton decay can actually be prevented by allowing only one of the four potential R-parity violating terms to survive.

Soft supersymmetry breaking

The fundamental idea of *soft* supersymmetry breaking[15, 34, 35, 72, 73] is that we would like to break supersymmetry without reintroducing the quadratic divergences we discussed at the end of Chapter 2. We write the Lagrangian in a form :

$$\mathcal{L}_{\text{MSSM}} = \mathcal{L}_{\text{SUSY}} + \mathcal{L}_{\text{soft}} \tag{3.10}$$

In this sense, the symmetry breaking is "soft", since we have separated out the completely symmetric terms from those soft terms which will not allow the quadratic divergences to the Higgs mass.

The explicitly allowed terms in the soft-breaking Lagrangian are [35].

- Mass terms for the scalar components of the chiral supermultipletss
- Mass terms for the Weyl spinor components of the gauge supermultipletss
- Trilinear couplings of scalar components of chiral supermultiplets

In particular, using the field content described above for the MSSM, the softly-broken portion of the MSSM Lagrangian can be writen

$$\begin{aligned}
\mathcal{L}_{\text{soft}} = &-\frac{1}{2}\left(M_3\tilde{g}\tilde{g} + M_2\tilde{W}\tilde{W} + M_1\tilde{B}\tilde{B} + c.c.\right) &(3.11)\\
&-\left(\tilde{u}a_u\tilde{Q}H_u - \tilde{d}a_d\tilde{Q}H_d - \tilde{e}a_e\tilde{L}H_d + c.c.\right) &(3.12)\\
&-\tilde{Q}^\dagger m_Q^2\tilde{Q} - \tilde{L}^\dagger m_L^2\tilde{L} - \tilde{u}m_u^2\tilde{u}^\dagger - \tilde{d}m_d^2\tilde{d}^\dagger - \tilde{e}m_e^2\tilde{e}^\dagger &(3.13)\\
&-m_{H_u}^2 H_u^* H_u - m_{H_d}^2 H_d^* H_d - (bH_uH_d + cc). &(3.14)
\end{aligned}$$

where we have introduced the following notations :

1. M_3, M_2, M_1 are the gluino, wino, and bino masses.
2. a_u, a_d, a_e are complex 3×3 matrices in family space.
3. $m_Q^2, m_u^2, m_d^2, m_L^2, m_e^2$ are hermitian 3×3 matrices in family space.

4. $m^2_{H_u}, m^2_{H_u}, b$ are the SUSY-breaking contributions to the Higgs potential.

We have written matrix terms without any sort of additional notational decoration to indicate their matrix nature, and we now show why. The first term 1 are straightforward; these are just the straightforward mass terms for these fields. There are strong constraints on the off-diagonal terms for the matrices of 2 [74, 75]; for simplicity, we will assume that each $a_i, i = u, d, e$ is proportional to the Yukawa coupling matrix : $a_i = A_{i0} y_i$. The matrices in **??** can be similarly constrained by experiments [68, 75–82] Finally, we assume that the elements 4 contributing to the Higgs potential as well as all of the 1 terms must be real, which limits the possible CP-violating interactions to those of the Standard Model. We thus only consider flavor-blind, CP-conserving interactions within the MSSM.

The important mixing for mass and gauge interaction eigenstates in the MSSM occurs within electroweak sector, in a process akin to EWSB in the Standard Model. The neutral portions of the Higgsinos doublets and the neutral gauginos $(\tilde{H}^0_u, \tilde{H}^0_d, \tilde{B}^0, \tilde{W}^0$ of the gauge interaction basis mix to form what are known as the *neutralinos* of mass basis :

$$M_{\tilde{\chi}} = \begin{pmatrix} M_1 & 0 & -c_\beta s_W m_Z & s_\beta s_W m_Z \\ 0 & M_2 & c_\beta c_W m_Z & -s_\beta c_W m_Z \\ -c_\beta s_W m_Z & c_\beta c_W m_Z & 0 & -\mu \\ s_\beta s_W m_Z & -s_\beta c_W m_Z & -\mu & 0 \end{pmatrix} \tag{3.15}$$

where $s(c)$ are the sine and cosine of angles related to EWSB, which introduced masses to the gauginos and higgsinos. Diagonalization of this matrix gives the four neutralino mass states, listed without loss of generality in order of increasing mass : $\tilde{\chi^0_{1,2,3,4}}$.

The neutralinos, especially the lightest neutralino $\tilde{chi}^0_1$, are important ingredients in SUSY phenomenology.

The same process can be done for the electrically charged gauginos with the charged portions of the Higgsino doublets along with the charged winos $(\tilde{H_u^+}, \tilde{H_d^+}, \tilde{W^+}, \tilde{W^-}$. This leads to the *charginos*, again in order of increasing mass : $\tilde{\chi}^{\pm}_{1,2}$.

3.3 Phenomenology

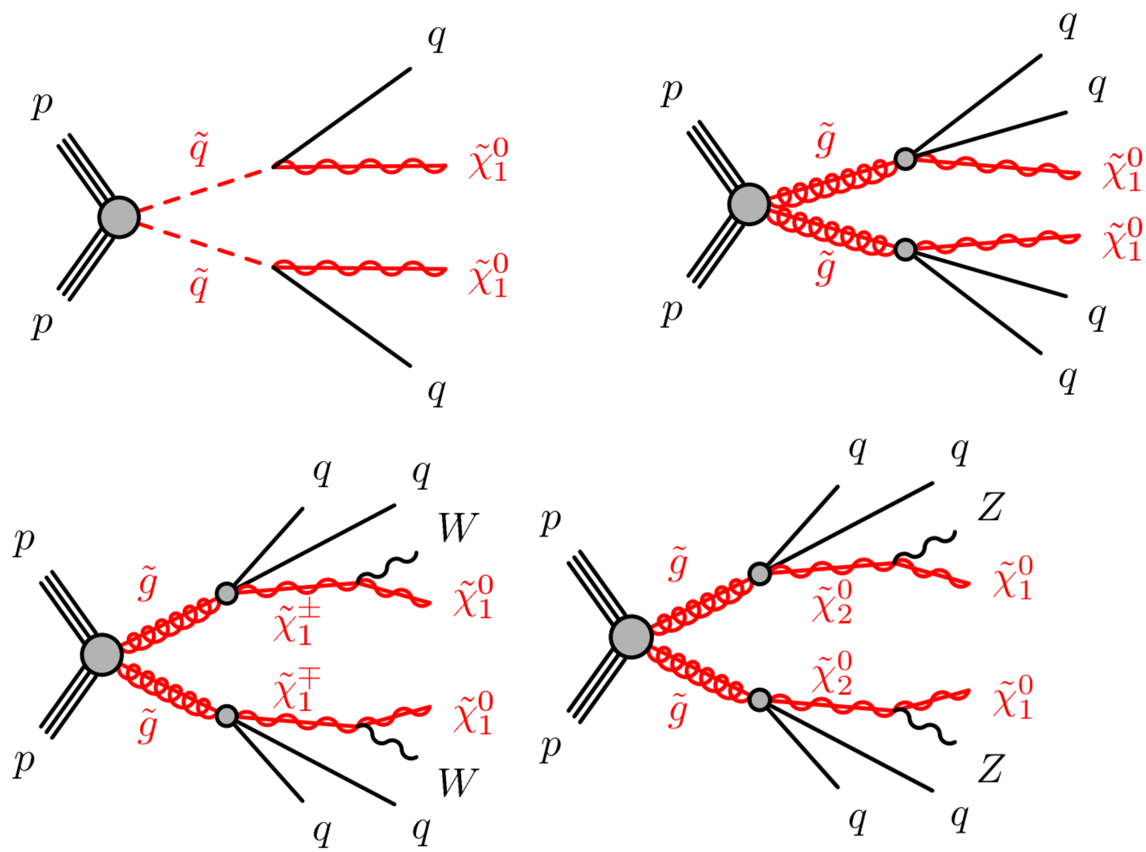
We are finally at the point where we can discuss the phenomenology of the MSSM, in particular as it manifests itself at the energy scales of the LHC.

As noted above in Sec.3.2, the assumption of $R-$parity has important consequences for MSSM phenomenology. The SM particles have $R = 1$, while the sparticles all have $R = -1$. Simply, this is the "charge" of supersymmetry. Since the particles of LHC collisions (pp) have total incoming $R = 1$, we must expect that all sparticles will be produced in *pairs*. An additional consequence of this symmetry is the fact that the lightest supersymmetric particle (LSP) is *stable*. Off each branch of the Feynmann diagram shown in Fig., we have $R = -1$, and this can only decay to another sparticle and a SM particle. Once we reach the lightest sparticle in the decay, it is absolutely stable. This leads to the common signature $E_{\mathrm{T}}^{\mathrm{miss}}$ for a generic SUSY signal.

For this thesis, we will be presenting an inclusive search for squarks and gluinos with zero leptons in the final state. This is a very interesting decay channel[4], due to the high cross-sections of $\tilde{g}\tilde{g}$ and $\tilde{q}\tilde{q}$ decays, as can be seen in Fig.**??** [83]. This is a direct consequence of the fact that these are the colored particles of the MSSM. Since the sparticles interact with the gauge groups of the SM in the same way as their SM partners, the colored sparticles, the squarks and gluinos, are produced and decay as governed by the color group $SU(3)_C$ with the strong coupling g_S. The digluino production is particularly copious, due to color factor corresponding to the color octet

[4] Prior to Run1, probably the most *most* interesting SUSY decay channel.

Figure 3.3: SUSY signals considered in this thesis



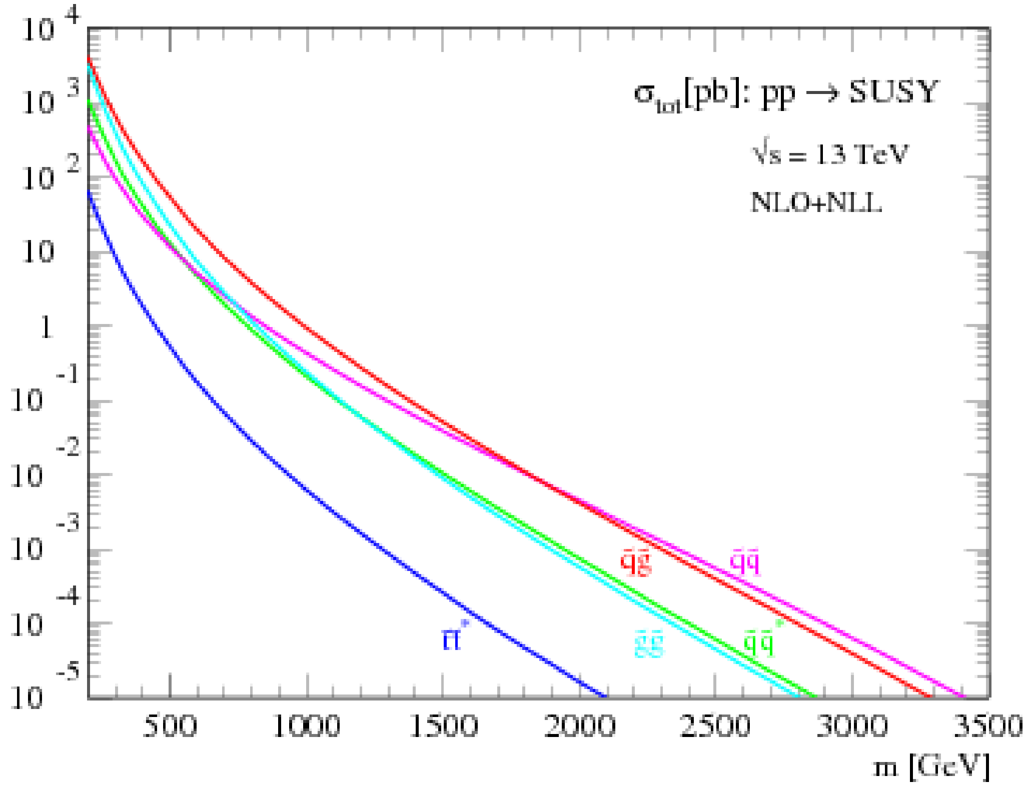
of $SU(3)C$.

In the case of disquark production, the most common decay mode of the squark in the MSSM is a decay directly to the LSP plus a single SM quark [15]. This means the basic search strategy of disquark production is two jets from the final state quarks, plus missing transverse energy for the LSPs. There are also cascade decays, the most common of which, and the only one considered in this thesis, is $\tilde{q} \to q\tilde{\chi}^{\pm} \to qW^{\pm}\chi^0$.

For digluino production, the most common decay is $\tilde{g} \to g\tilde{q}$, due to the large g_S coupling. The squark then decays as listed above. In this case, we generically search for four jets and missing transverse energy from the LSPs. We can also have the squark decay in association with a $W^{\pm}$ or Z^0; in this thesis, we are interested in those cases where this vector boson goes hadronically.

In the context of experimental searches for SUSY, we often consider *simplified*

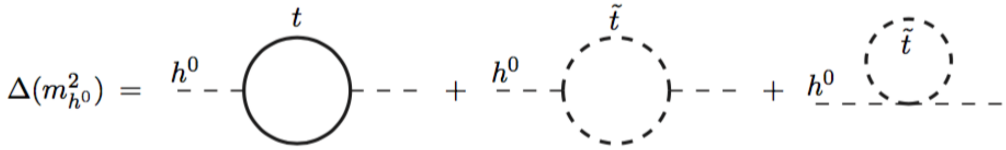
Figure 3.4: SUSY production cross-sections as a function of sparticle mass at $\sqrt{s} =$ 13 TeV.



models. These models make certain assumptions which allow easy comparisons of results by theorists and rival experimentalists. In the context of this thesis, the simplified models will make assumptions about the branching ratios described in the preceding paragraphs. In particular, we will often choose a model where the decay of interest occurs with 100% branching ratio. This is entirely for ease of interpretation by other physicists[5], but it is important to recognize that these are more a useful comparison tool, especially with limits, than a strict statement about the potential masses of sought-after beyond the Standard Model particle.

[5] In the author's opionion, this often leads to more confusion than comprehension. We will revisit the shortcomings of simplified models in the Conclusion to this thesis.

Figure 3.5: Loop diagrams correct the Higgs mass in the MSSM



3.4 How SUSY solves the problems with the SM

We now return to the issues with the Standard Model as described in Ch.2 to see how these issues are solved by supersymmetry.

Quadratic divergences to the Higgs mass

The quadratic divergences induced by the loop corrections to the Higgs mass, for example from the top Yukawa coupling, goes as

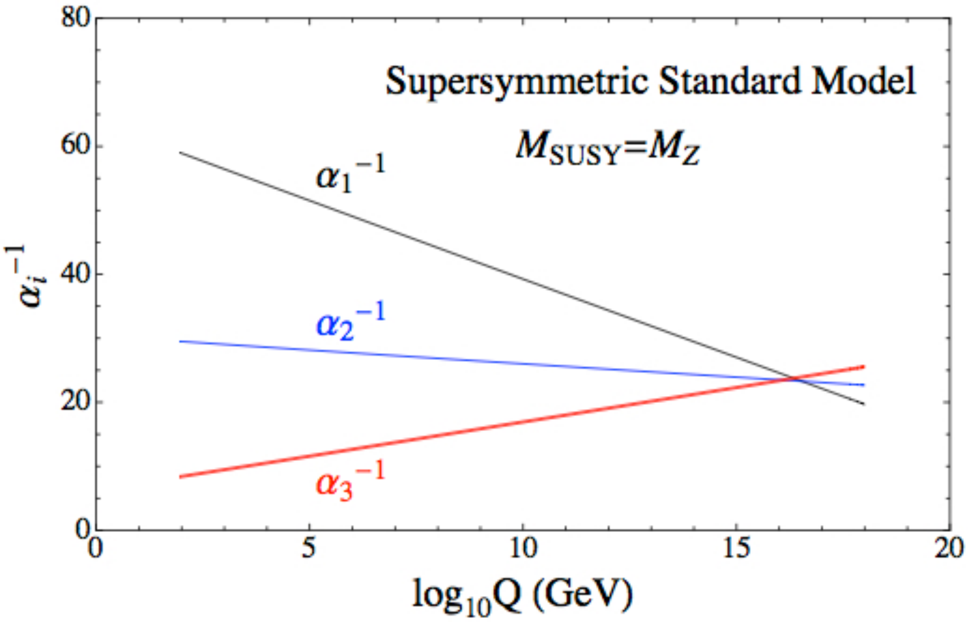
$$\delta m_H^2 \approx \left(\frac{m_t}{8\pi^2 {<}\phi{>}_{VEV}}\right)^2 \Lambda^2_{Planck}. \tag{3.16}$$

The miraculous thing about SUSY is each of these terms *automatically* comes with a term which exactly cancels this contribution[15]. The fermions and bosons have opposite signs in this loop diagram to all orders in pertubation theory, which completely solves the hierarchy problem. This is the most well-motivated reason for supersymmetry.

Gauge coupling unification

An additional motivation for supersymmetry is seen by the gauge coupling unification high scales. In the Standard Model, as we saw the gauge couplings fail to unify at high energies. In the MSSM and many other forms of supersymmetry, the gauge couplings unify at high energy, as can be seen in Fig.??. This provides additional aesthetic motivation for supersymmetric theories.

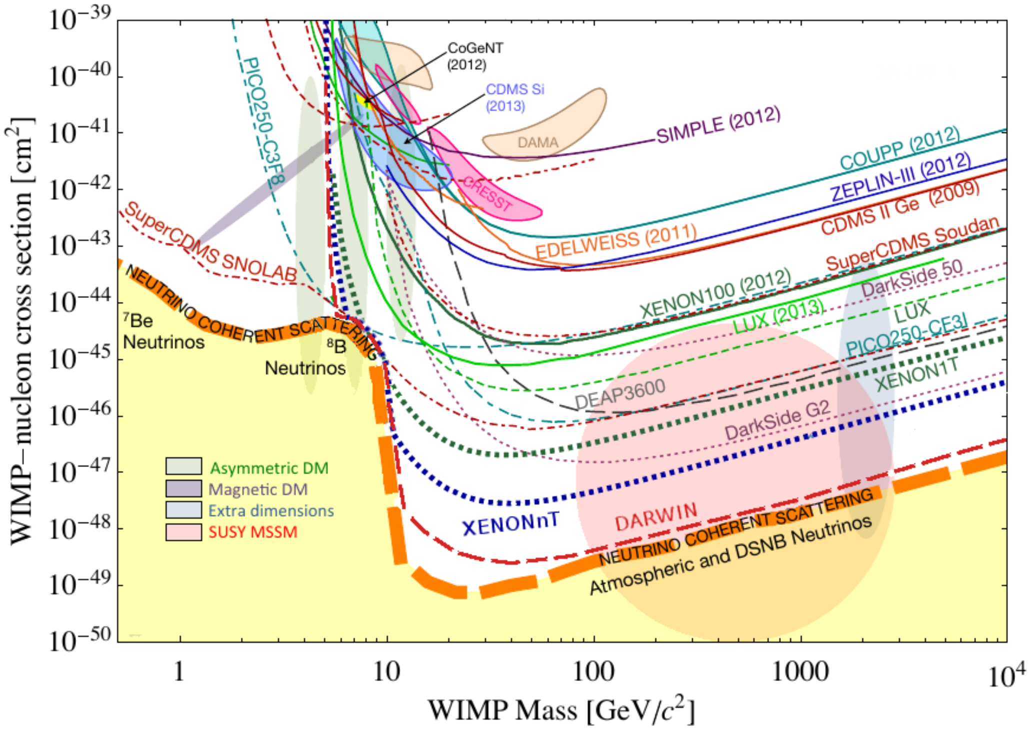
Figure 3.6: The running of Standard Model gauge couplings; compare to Fig.2.4. The MSSM gauge couplings nearly intersect at high energies.



Dark matter

As we discussed previously, the lack of any dark matter candidate in the Standard Model naturally leads to beyond the Standard Model theories. In the Standard Model, there is a natural dark matter candidate in the lightest supersymmetric particle[15] The LSP would in dark matter experiments be called a *weakly-interacting massive particle* (WIMP), which is a type of cold dark matter [22, 84]. These WIMPS would only interact through the weak force and gravity, which is exactly as a model like the MSSM predicts for the neutralino. In Fig.3.7, we can see the current WIMP exclusions for a given mass. The range of allowed masses which have not been excluded for LSPs and WIMPs have significant overlap. This provides additional motivation outside of the context of theoretical details.

Figure 3.7: WIMP exclusions from direct dark matter detection experiments.



3.5 Conclusions

Supersymmetry is the most well-motivated theory for physics beyond the Standard Model. It provides a solution to the hierarchy problem, leads to gauge coupling unification, and provides a dark matter candidate consistent with galatic rotation curves. As noted in this chapter, due to the LSPs in the final state, most SUSY searches require a significant amount of missing transverse energy in combination with jets of high transverse momentum. However, there is some opportunity to do better than this, especially in final states where one has two weakly-interacting LSPs on opposite sides of some potentially complicated decay tree. We will see how this is done in Ch.??.

Chapter 4

The Large Hadron Collider

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4.1 Magnets

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Chapter 5

The ATLAS detector

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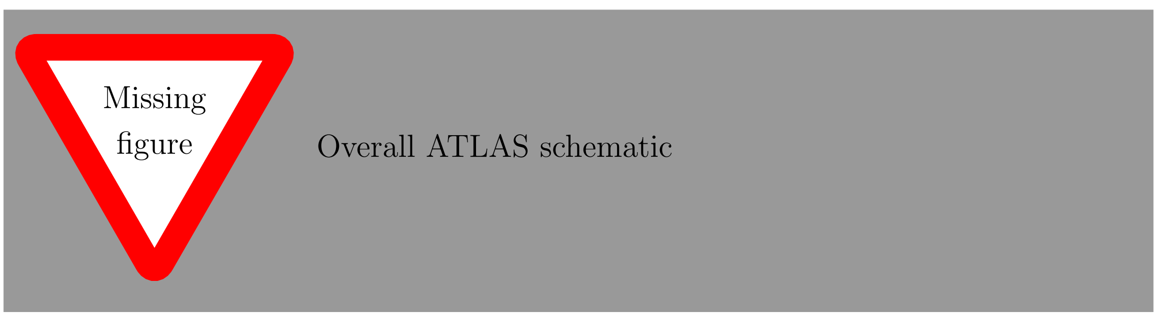

5.1 Inner Detector

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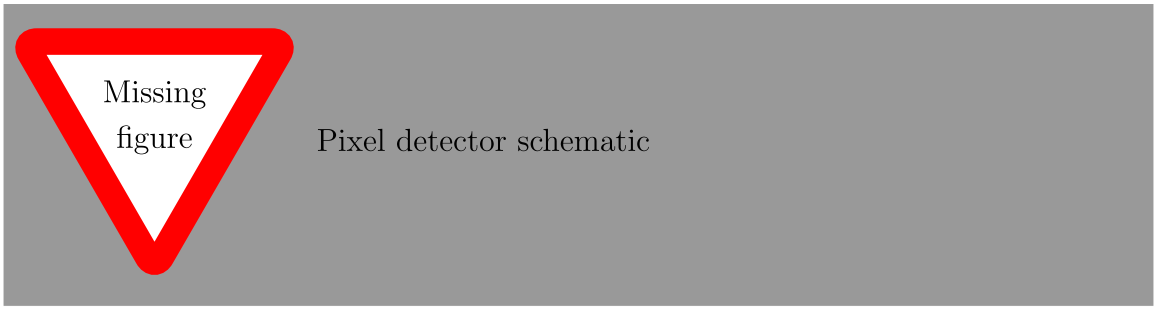
Pixel Detector



Insertable B-Layer

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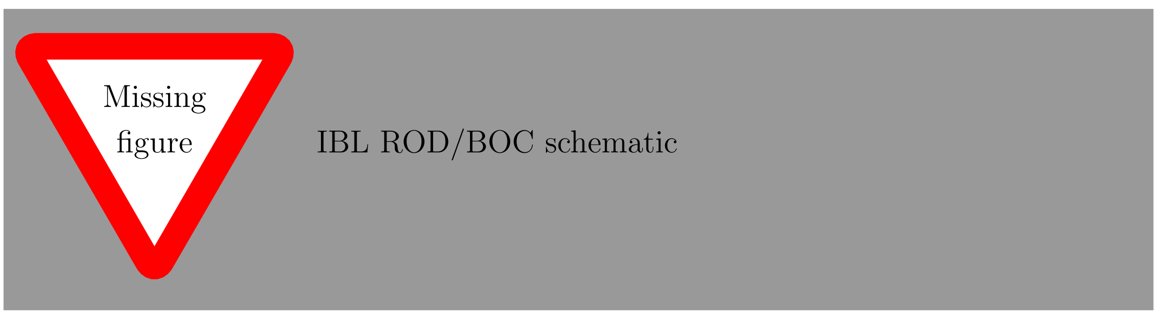
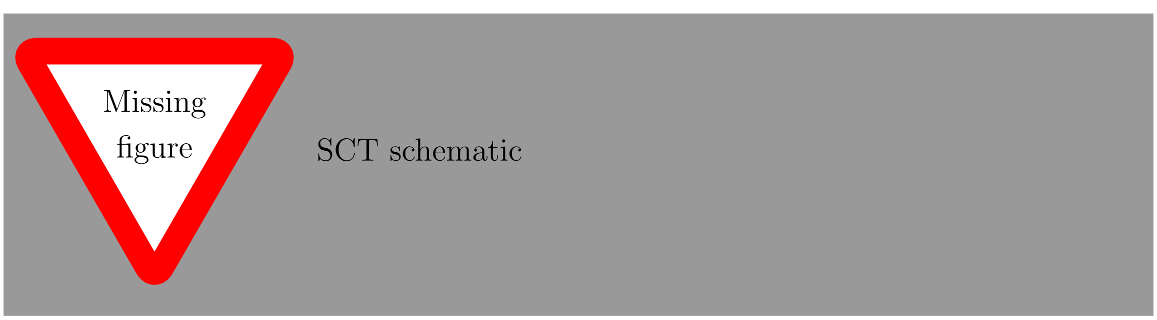
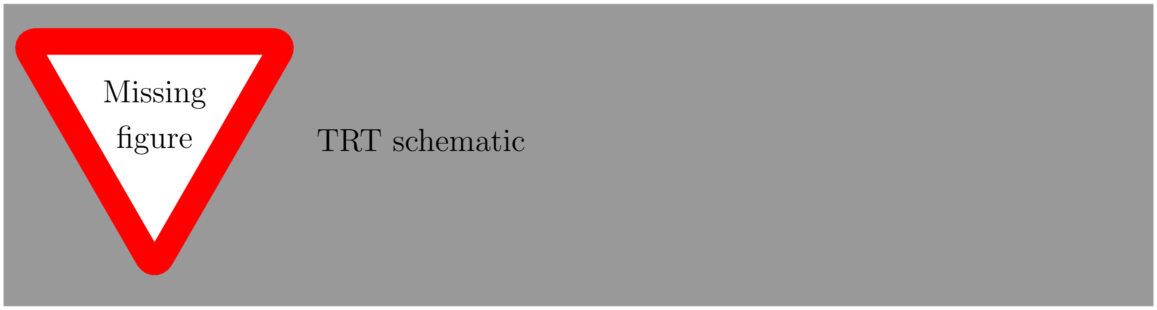


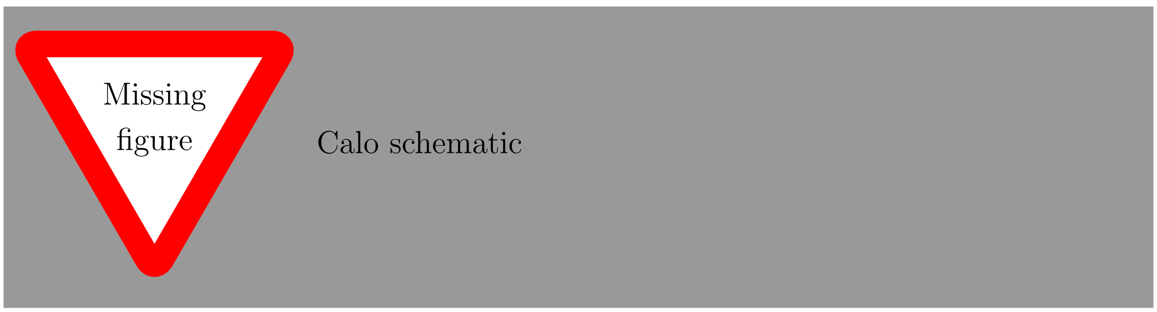
Semiconductor Tracker



Transition Radiation Tracker



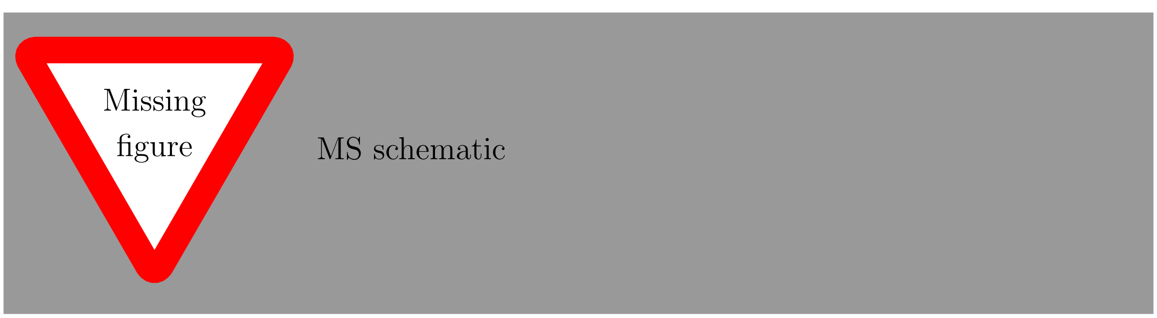
5.2 Calorimeter



Electromagnetic Calorimeter

Hadronic Calorimeter

5.3 Muon Spectrometer



Chapter 6

The Recursive Jigsaw Technique

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6.1 Razor variables

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6.2 SuperRazor variables

6.3 The Recursive Jigsaw Technique

6.4 Variables used in the search for zero lepton SUSY

Chapter 7

Title of Chapter 1

Chapter 8

Title of Chapter 1

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8.1 Object reconstruction

Photons, Muons, and Electrons

Jets

Missing transverse momentum

Probably longer, show some plots from the PUB note that we worked on

8.2 Signal regions

Gluino signal regions

Squark signal regions

Compressed signal regions

8.3 Background estimation

Z vv

W ev

ttbar

Chapter 9

Title of Chapter 1

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9.1 Statistical Analysis

maybe to be moved to an appendix

9.2 Signal Region distributions

9.3 Pull Plots

9.4 Systematic Uncertainties

9.5 Exclusion plots

Conclusion

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9.6 New Section

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The Standard Model

In this appendix, we provide a brief overview of the basic ingredients involved in construction of the Standard Model Lagrangian : quantum field theory, symmetries, and symmetry breaking.

Quantum Field Theory

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In this section, we provide a brief overview of the necessary concepts from Quantum Field Theory (QFT).

In modern physics, the laws of nature are described by the "action" S, with the imposition of the principle of minimum action. The action is the integral over the spacetime coordinates of the "Lagrangian density" $\mathcal{L}$, or Lagrangian for short. The Lagrangian is a function of "fields"; general fields will be called $\phi(x^\mu)$, where the indices μ run over the space-time coordinates. We can then write the action S as

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$$S = \int d^4x \mathcal{L}[\phi_i(x^\mu), \partial_\mu \phi_i(x^\mu)] \tag{9.1}$$

where we have an additional summation over i (of the different fields). Generally, we impose the following constraints on the Lagrangian :

1. Translational invariance - The Lagrangian is only a function of the fields ϕ and their derivatives $\partial_\mu \phi$
2. Locality - The Lagrangian is only a function of one point x_μ in spacetime.

3. Reality condition - The Lagrangian is real to conserve probability.

4. Lorentz invariance - The Lagrangian is invariant under the Poincarégroup of spacetime.

5. Analyticity - The Lagrangian is an analytical function of the fields; this is to allow the use of pertubation theory.

6. Invariance and Naturalness - The Lagrangian is invariant under some internal symmetry groups; in fact, the Lagrangian will have *all* terms allowed by the imposed symmetry groups.

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7. Renormalizabilty - The Lagrangian will be renormalizable - in practice, this means there will not be terms with more than power 4 in the fields.

The key item from the point of view of this thesis is that of "Invariance and Natural". We impose a set of "symmetries" and then our Lagragian is the most general which is allowed by those symmetries.

Symmetries

Symmetries can be seen as the fundamental guiding concept of modern physics. Symmetries are described by "groups". . To illustrate the importance of symmetries and their mathematical description, groups, we start here with two of the simplest and most useful examples : $\mathbb{Z}_2$and $U(1)$.

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$\mathbb{Z}_2$symmetry

$\mathbb{Z}_2$symmetry is the simplest example of a "discrete" symmetry. Consider the most general Lagrangian of a single real scalar field $\phi(x_\mu)$

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 - \frac{\mu}{2\sqrt{2}}\phi^3 - \lambda\phi^4 \tag{9.2}$$

Now we *impose* the symmetry

$$\mathcal{L}(\phi) = \mathcal{L}(-\phi) \tag{9.3}$$

This has the effect of restricting the allowed terms of the Lagrangian. In particular, we can see the term $\phi^3 \to -\phi^3$ under the symmetry transformation, and thus must be disallowed by this symmetry. This means under the imposition of this particular symmetry, our Lagrangian should be rewritten as

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{m^2}{2}\phi^2 - \lambda\phi^4 \tag{9.4}$$

The effect of this symmetry is that the total number of ϕ particles can only change by even numbers, since the only interaction term $\lambda\phi^4$ is an even power of the field. This symmetry is often imposed in supersymmetric theories, as we will see in Chapter 3.

$U(1)$ symmetry

$U(1)$ is the simplest example of a continuous (or *Lie*) group. Now consider a theory with a single complex scalar field $\phi = \operatorname{Re}\phi + i\operatorname{Im}\phi$

$$\mathcal{L}_\phi = \delta_{i,j}\frac{1}{2}\partial_\mu\phi_i\partial^\mu\phi_j - \frac{m^2}{2}\phi_i\phi_j - \frac{\mu}{2\sqrt{2}}\phi_i\phi_j\phi_k - \lambda\phi_i\phi_j\phi_k\phi_l \tag{9.5}$$

where $i, j, k, l = Re, Im$. In this case, we impose the following $U(1)$ symmetry : $\phi \to e^{i\theta}, \phi^* \to e^{-i\theta}$. We see immediately that this again disallows the third-order terms, and we can write a theory of a complex scalar field with $U(1)$ symmetry as

$$\mathcal{L}_\phi = \partial_\mu\phi\partial^\mu\phi^* - \frac{m^2}{2}\phi\phi^* - -\lambda(\phi\phi^*)^2 \tag{9.6}$$

Local symmetries

The two examples considered above are "global" symmetries in the sense that the symmetry transformation does not depends on the spacetime coordinate x_μ. We know look at local symmetries; in this case, for example with a local $U(1)$ symmetry, the transformation has the form $\phi(x_\mu) \to e^{i\theta(x_m u)}\phi(x_\mu)$. These symmetries are also known as "gauge" symmetries; all symmetries of the Standard Model are gauge symmetries.

There are wide-ranging consequences to the imposition of local symmetries. To begin, we note that the derivative terms of the Lagrangian 9.2 are *not* invariant under a local symmetry transformation

$$\partial_\mu\phi(x_\mu) \to \partial_\mu(e^(i\theta(x_\mu)\phi(x_\mu)) = (1 + i\theta(x_\mu))e^(i\theta(x_\mu)\phi(x_\mu) \tag{9.7}$$

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This leads us to note that the kinetic terms of the Lagrangian are also not invariant under a gauge symmetry. This would lead to a model with no dynamics, which is clearly unsatisfactory.

Let us take inspiration from the case of global symmetries. We need to define a so-called "covariant" derivative D^μ such that

$$D^\mu\phi \to e^{iq\theta(x^\mu)D^\mu\phi} \tag{9.8}$$

$$D^\mu\phi^* \to e^{-iq\theta(x^\mu)D^\mu\phi} \tag{9.9}$$

$$\tag{9.10}$$

Since ϕ andϕ^* transforms with the opposite phase, this will lead the invariance of the Lagrangian under our local gauge transformation. This D^μ is of the following form

$$D^\mu = \partial_\mu - igqA^\mu \tag{9.11}$$

where A^μ is a vector field we introduce with the transformation law

$$A^\mu \to A^\mu - \frac{1}{g}\partial_\mu \theta \tag{9.12}$$

and g is the coupling constant associated to vector field. This vector field A^μ is also known as a "gauge" field.

Since we need to add all allowed terms to the Lagrangian, we define

$$F^{\mu\nu} = A^\mu A^\nu - A^\nu A^\mu \tag{9.13}$$

and then we must also add the kinetic term :

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} \tag{9.14}$$

The most general renormalizable Lagrangian with fermion and scalar fields can be written in the following form

$$\mathcal{L} = \mathcal{L}_{kin} + \mathcal{L}_\phi + \mathcal{L}_\psi + \mathcal{L}Yukawa \tag{9.15}$$

Symmetry breaking and the Higgs mechanism

Here we view some examples of symmetry breaking. We investigate breaking of a global $U(1)$ symmetry and a local $U(1)$ symmetry. The SM will break the electroweak symmetry $SU(2)xU(1)$, and in Chapter 3 we will see how supersymmetry must also be broken.

There are two ideas of symmetry breaking

- Explicit symmetry breaking by a small parameter - in this case, we have a small parameter which breaks an "approximate" symmetry of our Lagrangian. An example would be the theory of the single scalar field 9.2, when $\mu << m^2$ and

$\mu << \lambda$. In this case, we can often ignore the small term when considering low-energy processes.

- Spontaneous symmetry breaking (SSB) - spontaneous symmetry breaking occurs when the Lagrangian is symmetric with respect to a given symmetry transformation, but the ground state of the theory is *not* symmetric with respect to that transformation. This can have some fascintating consequences, as we will see in the following examples

Symmetry breaking a

U(1) global symmetry breaking

Consider the theory of a complex scalar field under the $U(1)$ symmetry, or the transformation

$$\phi \to e^{i\theta}\phi \tag{9.16}$$

The Lagrangian for this theory is

$$\mathcal{L} = \partial^{\mu}\phi^{\dagger}\partial_{\mu}\phi + \frac{\mu^2}{2}\phi^{\dagger}\phi + \frac{\lambda}{4}(\phi^{\dagger}\phi)^2 \tag{9.17}$$

Let us write this theory in terms of two scalar fields, h and ξ : $\phi = (h+i\xi)/\sqrt{(}2)$. The Lagrangian can then be written as

$$\mathcal{L} = \partial^{\mu}h\partial_{\mu}h + \partial^{\mu}\xi dmu\xi - \frac{\mu^2}{2}(h^2+\xi^2) - \frac{\lambda}{4}(h^2+\xi^2)^2 \tag{9.18}$$

First, note that the theory is only stable when $\lambda > 0$. To understand the effect of SSB, we now enforce that $\mu^2 < 0$, and define $v^2 = -\mu^2/\lambda$. We can then write the scalar potential of this theory as :

$$V(\phi) = \lambda(\phi^{\dagger}\phi - v^2/2)^2 \tag{9.19}$$

Minimizing this equation with respect to ϕ, we can see that the "vacuum expectation value" of the theory is

$$2 < \phi^{\dagger}\phi > = < h^2 + \xi^2 > = v^2 \tag{9.20}$$

We now reach the "breaking" point of this procedure. In the (h, ξ) plane, the minima form a circle of radius v. We are free to choose any of these minima to expand our Lagrangian around; the physics is not affected by this choice. For convenience, choose $< h >= v, < \xi^2 >= 0$.

Now, let us define $h' = h - v, \xi' = \xi$ with VEVs $< h' >= 0, < \xi' >= 0$. We can then write our spontaneously broken Lagrangian in the form

$$\mathcal{L} = \frac{1}{2}\partial_\mu h' \partial^\mu h' + \frac{1}{2}\partial_\mu \xi' \partial^\mu \xi' - \lambda v^2 h'^2 - \lambda v h'(h'^2 + \xi'^2) - \lambda(h'^2 + \xi'^2)^2 \tag{9.21}$$