A search for sparticles in zero lepton final states

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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, 23–33]. 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 [34].

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) [35–45]. 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 $\text{fb}^{-1}$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

cite Yuval's lectures and notes somehow

The Standard Model is another name for the theory of the internal symmetry group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$. This quantum field theory is the culmination of years of work in both theoretical and particle physics. In this thesis, we take the view one constructs a model with the field content and symmetries as inputs, and then writes down the most general Lagrangian consistent with those symmetries. This will be applicable for this chapter and the following one. Additional theoretical background is in 9.6.

CHECK

cite

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 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 (maximum degree 4 in the mass) terms, the electroweak Lagrangian can be written as

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

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 Higgs multiplet with weak hypercharge $Y = 1$. The covariant derivative D^μ is given by

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

where σ_a are the Pauli matrices, 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

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

The terms proportional in the Lagrangian to μ^2 and λ make up the "Higgs potential" . As normal (i.e. in 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 ??.

Cite

footnote how we probably shouldn't say mexican hat?

CITE THIS PICTURE

This potential has a minimum at $VALUE$; 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 :

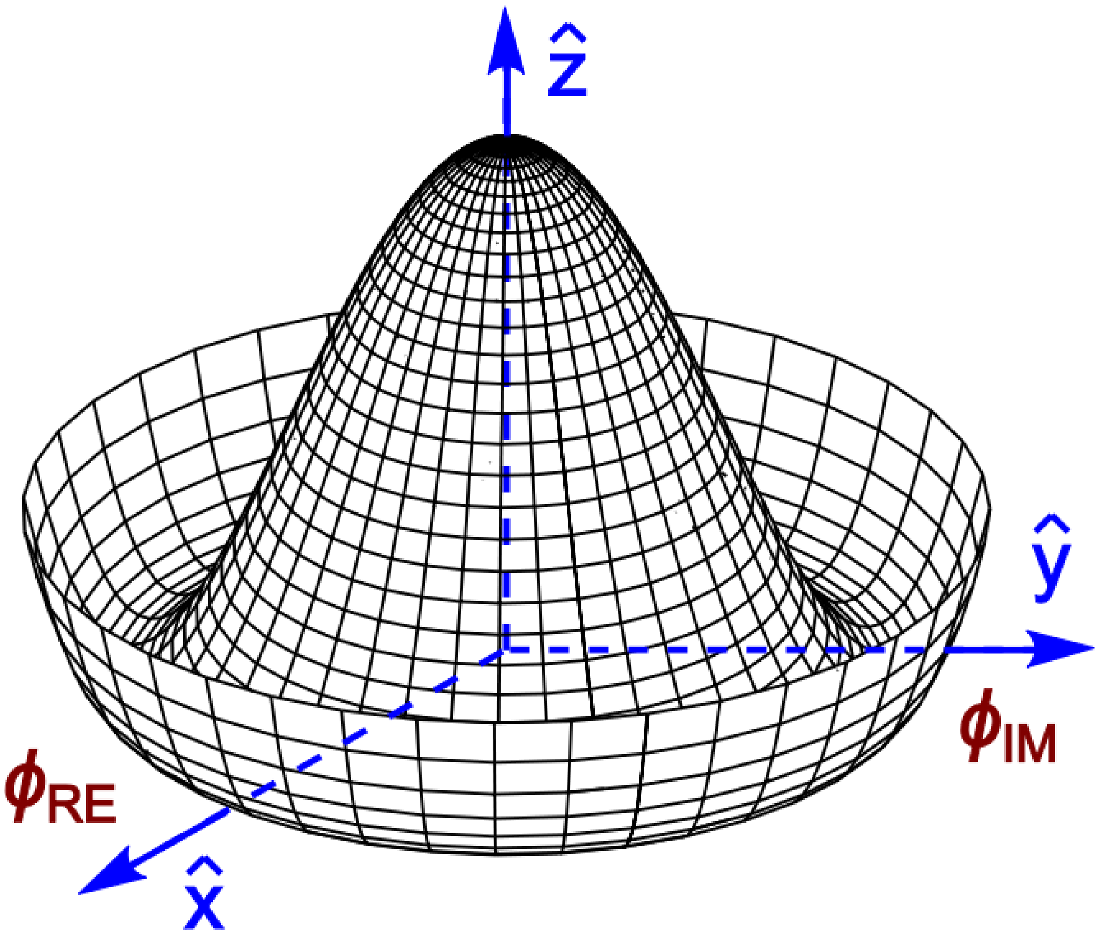
$$\phi = \frac{1}{\sqrt{2}} \exp(\frac{i}{v}\sigma_a\theta_a) \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \tag{2.5}$$

We can choose a gauge to rotate away the dependent 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

Figure 2.1: Sombrero potential



$h(x) = 0$ see that (dropping the Lorentz indices) :

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

CHECK THIS EQUATION, especially the squaring?, I'm a bit confused 162

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 see 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 :

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

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 . cite

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 $SU(3)_C$ which is an unbroken symmetry in the Standard Model; this implies the gluon remains massless. Defining the covariant derivative for QCD as

$$D^\mu = \partial^\mu + ig_sG_a^\mu L_a \tag{2.11}$$

where L_a are the generators of $SU(3)_C$, often represented by the Gell-Mann matrices, and g_s is the coupling constant of the strong force. . The QCD Lagrangian then is given by check the logic here? are tehy actually them?

$$\mathcal{L}_{\mathrm{QCD}} = i\bar{\psi}_f D_\mu \gamma^\mu \psi_f - \frac{1}{4}G^a_{\mu\nu}G_a^{\mu\nu} \tag{2.12}$$

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

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

sp

where f^{abc} are the structure constants of $SU(3)_C$, which are analogus 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.

maybe cite the QED lagrangian?

reference figure

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 (primarily 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. These 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, 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.

get that cross section picture

bag model?

Fortunately, we are generally saved by the QCD factorization theorem . This allows one to understand the hard (i.e. short distance or highest 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.**??**.

cite QCD factorization

.

show ratio of ee to qq?

Fermions

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

As noted earlier with regards to the field content, 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*. 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 $q_{EM} = 0$.

cite pdg probably

Often in an experimental context, lepton is used to denote the electron (stable) and muon (metastable), due to their striking experimental signatures. Taus are often treated separately, due to their much shorter lifetime of τ_τ; 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*.

footnote?

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

$$\begin{pmatrix} u & d \end{pmatrix}, \begin{pmatrix} c & s \end{pmatrix}, \begin{pmatrix} t & b \end{pmatrix} \tag{2.14}$$

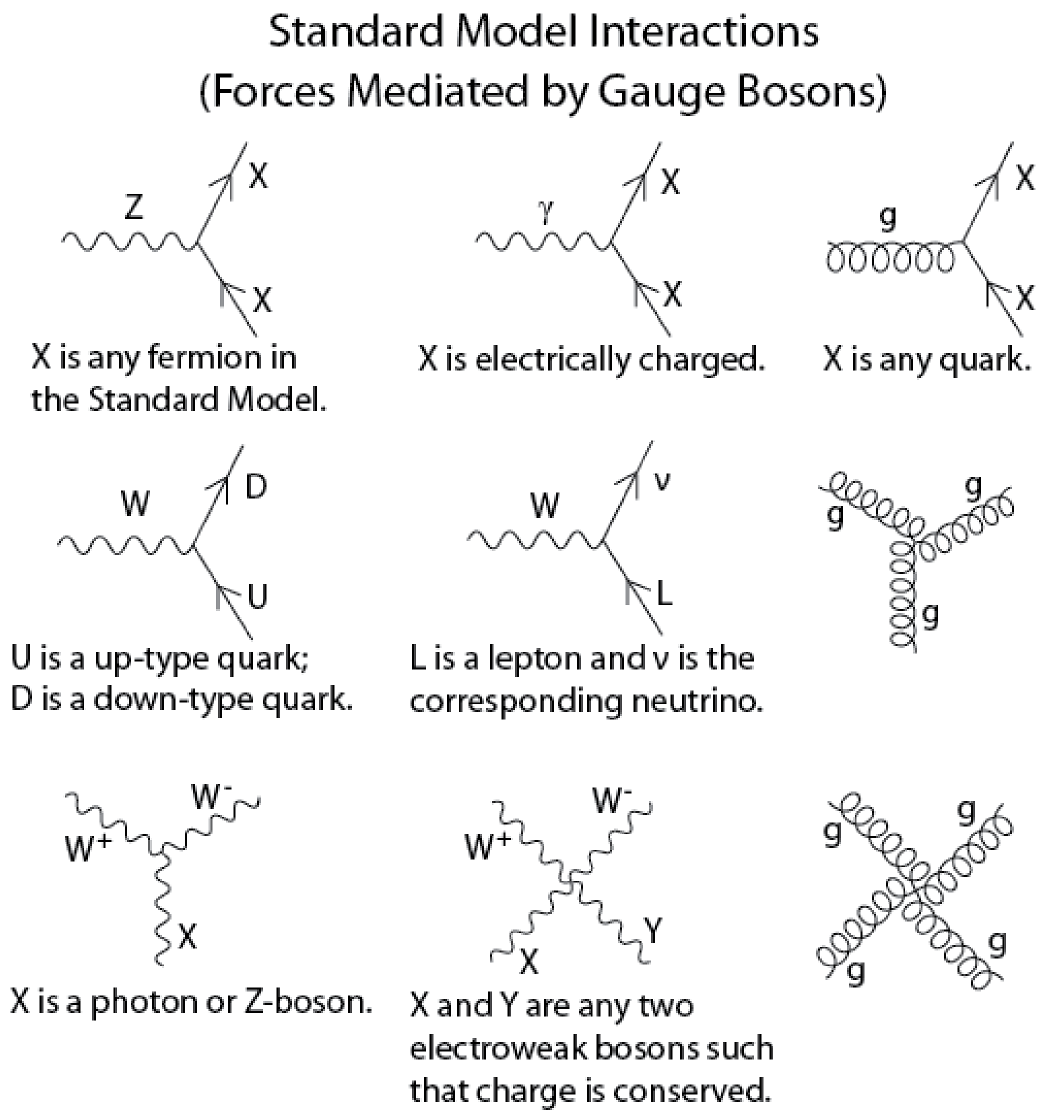
where we speak of "up-like" quarks and "down-like" quarks.

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 are indistinguishable by LHC experiments, and reconstructed as jets.. The bottom quark hadronizes primarly through a relatively long-lived particle known as the B (name), which generally travels a short distance before decay. This feature allows what is known as *b-tagging*; this will be further discussed in Ch.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 with a large cross-section, which can be used to distinguish signal processes with decays to top quarks, or understand top production as a background process.

footnote about charm tagging

refCh ATLAS.

Figure 2.2: The interactions of the Standard Model



Interactions in the Standard Model

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

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, governed by the coupling CHECK constant g'. A real Z_0 can thus decay to two of each fermion in the Standard Model except for the top quark, due to its large 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. The coupling constants for these interactions are encoded in the cite CKM matrix Secondly, the $W^{\pm}$ interacts with a charged lepton and its corresponding neutrino. 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 gluon interacts with any quark. Additionally, there are the gluon-only interactions, which occur in a three-point and four-point interaction.

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 Eq.**??**. 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.**??**.

add list of them from PDG

only?

The Standard Model has many free paramaters , especially when corrected for neutrino masses. In general, we prefer models with less free parameters. A great example of this fact, and additionally some of the strongest experimental proof of

EWSB, is the relationship between the couplings of the weak force and the masses of the gauge bosons of the weak force :

$$\rho m_Z^2 \cos^2\theta_W \stackrel{?}{=} 1 \tag{2.15}$$

where ? indicates that this is a testable prediction of the Standard Model (in particular, that the gauge bosons gain mass through EWSB). This relation has been shown to be true 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.

cite pdg

An additional issue, although not strictly fundamental, is the lack of gauge coupling unification. The couplings of any quantum field theory "run" as a function of the distance scales (or inversely, 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 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, without some additional theoretical gymnastics.

maybe show the lack of unification for the Standard Model?

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.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 scale of Planck scale of gravity, $\Lambda_{\text{Planck}} = 10^{19}$ GeV. In this case, we expect the corrections to the Higgs mass like

loop figure!!

$$\delta m_H^2 \approx \left(\frac{m_t}{8\pi^2 < \phi >_{VEV}}\right)^2 \Lambda^2_{Planck}. \tag{2.16}$$

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 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.

cite

show one

cite lectures from summer school on arxiv?

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. 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* be this 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 (maybe say more). 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.

2.4 Conclusions

The Standard Model is an extraordinary theory. It is a culmination of decades of work in both theoretical and experimental physics. blah some more

Chapter 3

Supersymmetry

Here you can write some introductory remarks about your chapter. I like to give each sentence its own line.

When you need a new paragraph, just skip an extra line.

3.1 Motivation

Only Additional allowed Lorentz invariant symmetry

Dark Matter

Cancellation of quadratic divergences in corrections to the Higgs Mass

3.2 Supersymmetry

3.3 Additional particle content

3.4 Phenomenology

R parity Consequences for sq/gl decays

Chapter 4

The Large Hadron Collider

Here you can write some introductory remarks about your chapter. I like to give each sentence its own line.

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4.1 Magnets

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Chapter 5

The ATLAS detector

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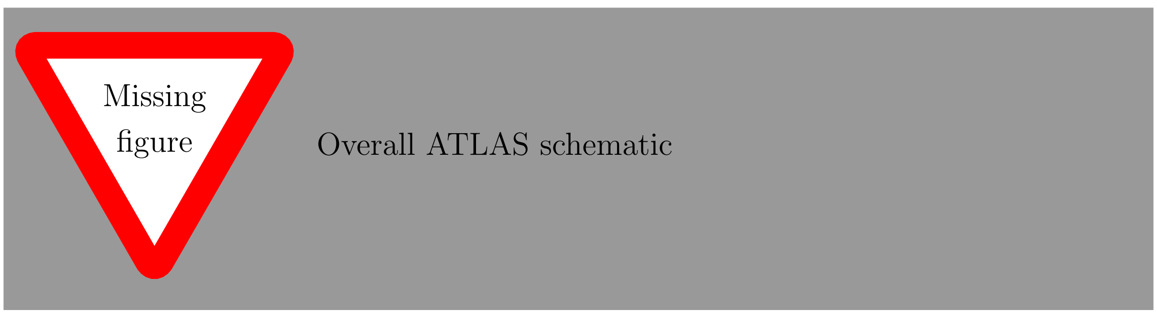

5.1 Inner Detector

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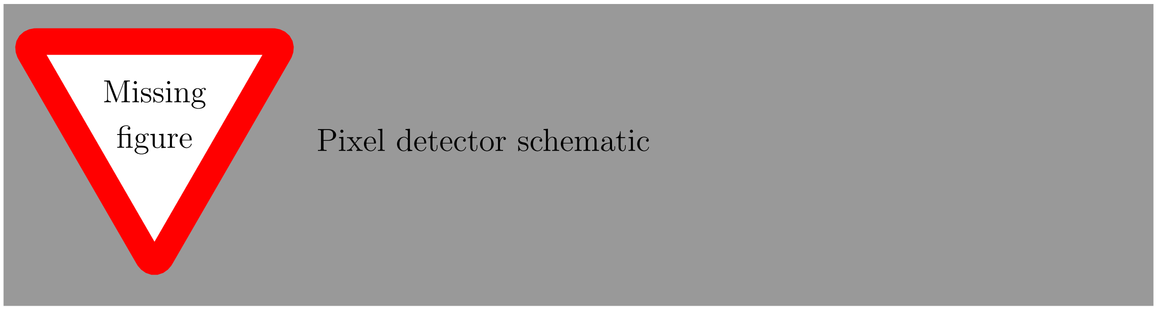
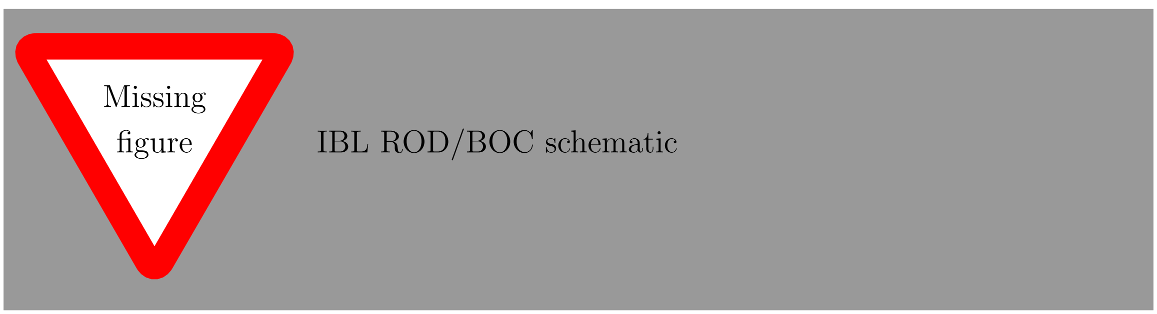
Pixel Detector

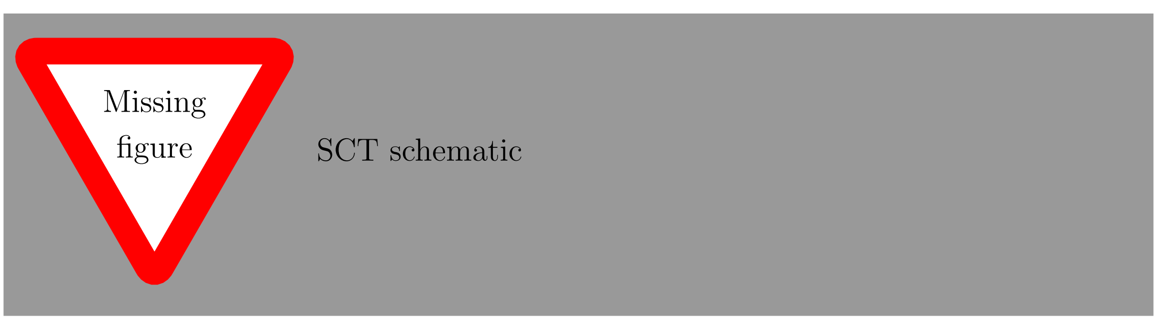


Insertable B-Layer

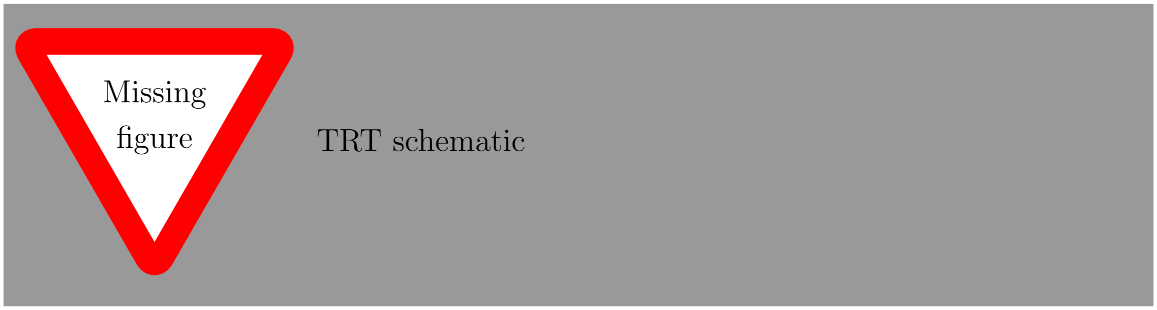
Qualification task, so add a bit more.



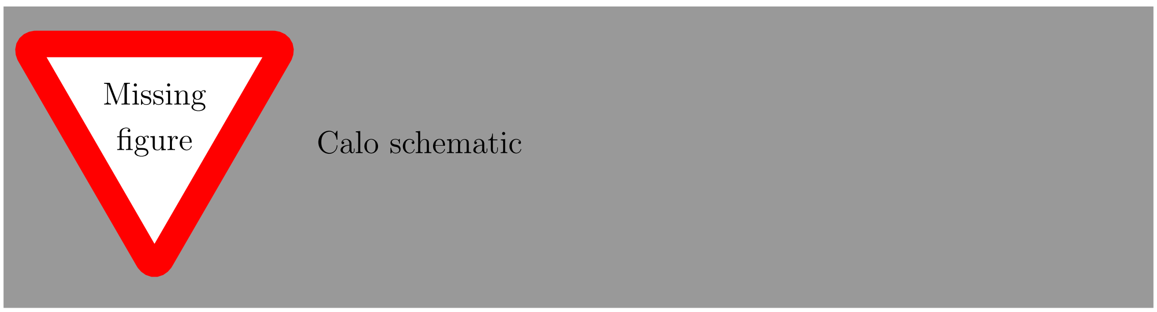
Semiconductor Tracker



Transition Radiation Tracker



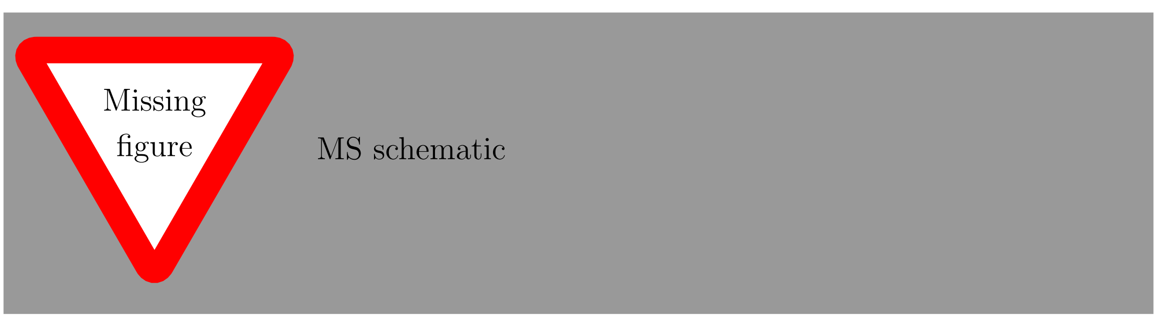
5.2 Calorimeter



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

Here you can write some introductory remarks about your chapter. I like to give each sentence its own line.

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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

Here you can write some introductory remarks about your chapter. I like to give each sentence its own line.

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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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When you need a new paragraph, just skip an extra line.

9.6 New Section

By using the asterisk to start a new section, I keep the section from appearing in the table of contents. If you want your sections to be numbered and to appear in the table of contents, remove the asterisk.

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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 $U(1)$.

cite?

$\mathbb{Z}_2$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

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

GET THIS RIGHT

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

$$\begin{aligned} D^\mu \phi &\to e^{iq\theta(x^\mu)D^\mu \phi} \\ D^\mu \phi^* &\to e^{-iq\theta(x^\mu)D^\mu \phi} \end{aligned} \tag{9.8}$$

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

$$D^\mu = \partial_\mu - igqA^\mu \tag{9.9}$$

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

$$A^\mu \to A^\mu - \frac{1}{g}\partial_\mu \theta \tag{9.10}$$

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

$$F^{\mu\nu} = A^\mu A^\nu - A^\nu A^\mu \tag{9.11}$$

and then we must also add the kinetic term :

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} \tag{9.12}$$

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

$$\mathcal{L} = \mathcal{L}_{kin} + \mathcal{L}_\phi + \mathcal{L}_\psi + \mathcal{L}Yukawa \tag{9.13}$$

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

$$\phi \to e^{i\theta}\phi \tag{9.14}$$

The Lagrangian for this theory is

$$\mathcal{L} = \partial^\mu \phi^\dagger \partial_\mu \phi + \frac{\mu^2}{2}\phi^\dagger \phi + \frac{\lambda}{4}(\phi^\dagger \phi)^2 \tag{9.15}$$

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

$$\mathcal{L} = \partial^\mu h \partial_\mu h + \partial^\mu \xi dmu \xi - \frac{\mu^2}{2}(h^2 + \xi^2) - \frac{\lambda}{4}(h^2 + \xi^2)^2 \tag{9.16}$$

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 :

$$V(\phi) = \lambda(\phi^\dagger \phi - v^2/2)^2 \tag{9.17}$$

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

$$2 < \phi^\dagger \phi > = < h^2 + \xi^2 > = v^2 \tag{9.18}$$

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

$$\mathcal{L} = \frac{1}{2}\partial_\mu h' \partial^\mu h' + \frac{1}{2}\partial_\mu \xi' \partial^\mu \xi' - \lambda v^2 h'^2 - \lambda v h'(h'^2 + \xi'^2) - \lambda (h'^2 + \xi'^2)^2 \tag{9.19}$$