A search for sparticles in zero lepton final states

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ABSTRACT

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Acknowledgements

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

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.

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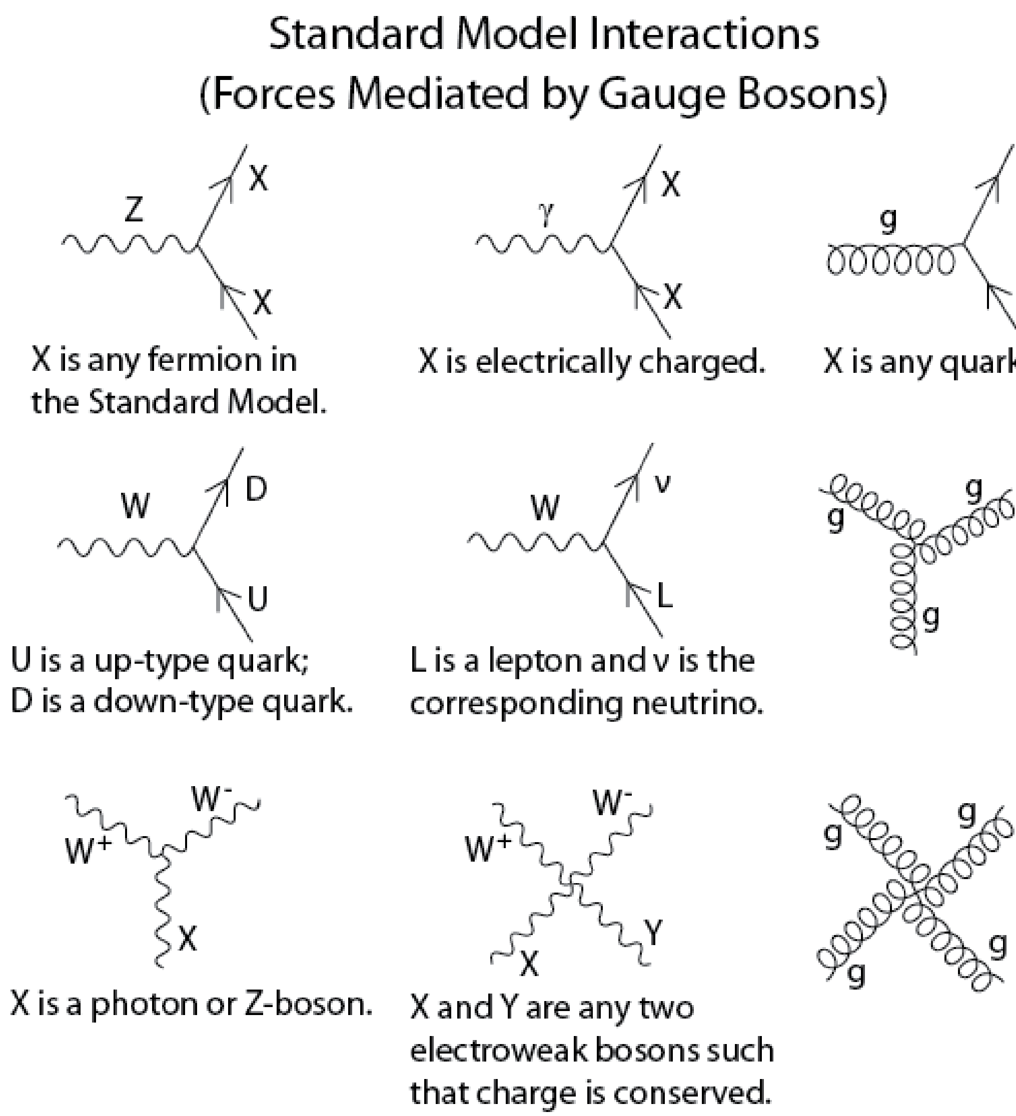
Field Content

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.

Figure 2.1: The interactions of the Standard Model



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; we say there are 8 gluons. The fermion fields $L_L(1,2)_{-1}$ and $E_R(1,1)_{-2}$ are singlets under $SU(3)_C$; we call them *leptons*.

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 so-called "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. The covariant derivative D^μ is given by

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

where T_a are the P

$\mathcal{L}_{kin}$

For each of the vector boson fields, we have the follow field strengths :

$$
\begin{aligned}
G_a^{\mu\nu} &= \partial^\mu G_a^\nu + \partial^\nu G_a^\mu - g_s f_{abc} G_b^\mu G_c^\nu \\
W_a^{\mu\nu} &= \partial^\mu W_a^\nu + \partial^\nu W_a^\mu - g \epsilon_{abc} W_b^\mu W_c^\nu \\
B^{\mu\nu} &= \partial^\mu B^\nu + \partial^\nu B^\mu
\end{aligned}
\tag{2.4}
$$

where g and g_s are the electroweak and strong coupling constant.

We can write the covariant derivative for the Standard Model as

$$
D^\mu = \partial^\mu + i g_s G_a^\mu L_a + i g W_a^\mu T_a + i g' Y B^\mu \tag{2.5}
$$

where L_a and T_a are the generators of $SU(3)_C$ and $SU(2)_L$ respectively for each of the representations. Explicitly, for the $SU(3)_C$ triplets, $L_a = \frac{1}{2}\lambda_a$ and for the $SU(3)_C$ singlets, $L_a = 0.$. For $SU(2)_L$ doublets, $T_a = \frac{1}{2}\sigma_a$ and for $SU(2)_L$ singlets, $T_a = 0$.

GELLMANN and Pauli matrices

The combination of these terms allows us to write the kinetic terms of the Lagrangian as

$$
\begin{aligned}
\mathcal{L}_{kin} &= G^{\mu\nu} G_{\mu\nu} + W^{\mu\nu} W_{\mu\nu} + B^{\mu\nu} B_{\mu\nu} \\
+ D^\mu Q_L D_\mu Q_L + D^\mu U_R D_\mu U_R &+ D^\mu D_R D_\mu D_R + D^\mu L_L D_\mu L_L L + D^\mu E_R D_\mu E_R
\end{aligned}
\tag{2.6}
$$

$\mathcal{L}_{\psi}$

We cannot write down any mass terms for fermions in the Standard Model. Dirac mass terms are forbidden since they are all assigned to "chiral" representations of the gauge symmetry. Majorana mass terms are disallowed since there are no fields with $Y \neq 0$.

$\mathcal{L}_{Yuk}$

We write the Yukawa portion of the Standard Model Lagrangian

$$\mathcal{L}_{Yuk} = Y_{ij} L_{Li} \bar{E}_{Rj} \phi + h.c. \tag{2.7}$$

The Yukawa matrix Y is a general complex 3×3 matrix of dimensionless couplings which can be diagonalized, leading to a diagonal matrix with only three real parameters (y_e, y_μ, y_τ). This reflects the fact that for the electron, muon, and tau lepton, the interaction basis is the same as the mass basis; this is the same as saying an electron has a well-defined mass.

2.1 $\mathcal{L}_\phi$, Electroweak Symmetry breaking and the Higgs Boson

Let us now recall that local gauge invariance means that the vector fields in this theory are *massless*. Näively, it seems this combined with the chirality of the Standard Model, that *none* of the fields have masses. The solution to this seeming conundrum is of course the well-known “Higgs” mechanism, described in Sec. 9.6.

In the Standard Model, the Higgs potential is given by

$$\mathcal{L}_\phi = -\mu^2 \phi^\dagger \phi - \lambda(\phi^\dagger \phi)^2. \tag{2.8}$$

Since λ is dimensionless and real, to have a potential bounded from below, we require $\lambda > 0$. To break the gauge symmetry, we require $\mu^2 < 0$, leading again to the sombrero potential **??**. We define

$$v^2 = -\frac{\mu^2}{\lambda}. \tag{2.9}$$

This allows us to write 2.8 as

$$\mathcal{L}_\phi = -\lambda(\phi^\dagger \phi - \frac{v^2}{2})^2 \tag{2.10}$$

after dropping the constant term.

This means the ϕ field acquires a VEV $| < \phi > | = v/\sqrt{2}$. Choosing the convenient gauge

$$\phi = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}, \tag{2.11}$$

The VEV breaks the $SU(2)_L \otimes U(1)_Y$ symmetry to a $U(1)_{EM}$ subgroup. We can identify the unbroken generator of this $U(1)_{EM}$ subgroup as $Q_{EM} = T_3 + Y/2$, since this vanishes in the down component

$$Q_\gamma \phi = (T_3 + Y/2)\phi = (\frac{1}{2}\begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} + \frac{1}{2}I)\begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}. \tag{2.12}$$

Here we see the indicative γ for the photon, as this unbroken $U(1)_{EM}$ symmetry is of course the symmetry associated to the electromagnetic force mediated by the gauge boson known as the photon.

There are three broken generators : $T_1, T_2, T_3 - Y/2$. These are each associated to one of the massive gauge bosons induced by the symmetry breaking. Choosing a gauge which rotates away the "eaten" Goldstone boson degrees of freedom, we can write the Higgs field as

$$\phi = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \tag{2.13}$$

2.2 Particle Spectrum : Standard Model Lagrangian after Electroweak Symmetry Breaking

We can now return to the Standard Model Lagrangian and use the equation for the Higgs field after EWSB 2.13. This will show us the "physical" particle content of the Standard Model.

Particle content associated to $\mathcal{L}_\phi$

Setting *phi* as in Eq.2.13, we quickly see that we can rewrite Eq.2.10 as

CHECK FACTORS OF TWO

$$\mathcal{L}_\phi = -\lambda(\phi^\dagger\phi - \frac{v^2}{2})^2 = -\lambda(\frac{1}{2}(v+h(x))^2 - \frac{v^2}{2})^2 = -\lambda(h(x)^2 + vh(x))^2 = -\lambda(h(x)^4 + vh(x)^3 + \frac{v^2}{2}h(x)^2). \tag{2.14}$$

Interpreting the Higgs field squared term as the mass term of the Higgs boson, we see that $m_H = \sqrt{2\lambda}v$.

Particle content associated to $\mathcal{L}_{kin}$

Again using Eq.2.13 and $D^\mu = \partial^\mu + ig_s G^\mu_a L_a + igW^\mu_a T_a + ig'YB^\mu$, we can see how the mass terms associated to the three massive gauge bosons, and also see how the photon stays massless. The mass terms for the gauge boson fields come from the kinetic term of the Higgs field :

$$\begin{aligned}\mathcal{L}_{M_V} = D^\mu\phi D_\mu\phi = (igW^\mu_a T_a + ig'YB^\mu)\frac{1}{\sqrt{2}}\begin{pmatrix}0\\v\end{pmatrix}(igW_{\mu,a}T_a + ig'YB_\mu)\frac{1}{\sqrt{2}}\begin{pmatrix}0\\v\end{pmatrix} = \\ \frac{1}{8}|\begin{pmatrix}gW_3 + g'B & g(W_1 - iW_2)\\ g(W_1 + iW_2) & -gW_3 + g'B\end{pmatrix}\begin{pmatrix}0\\v\end{pmatrix}|^2\end{aligned} \tag{2.15}$$

where we have noted that ∂_μ and L_a both disappear when acting on ϕ. Defining the *Weinberg* angle $\tan(\theta_W) = g'/g$ and the following physical fields :

$$\begin{aligned}W^\pm &= \frac{1}{\sqrt{2}}(W_1 \mp iW_2)\\ Z^0 &= \cos\theta_W W_3 - \sin\theta_W B\\ A^0 &= \sin\theta_W W_3 + \cos\theta_W B\end{aligned} \tag{2.16}$$

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

$$\mathcal{L}_{M_V} = \frac{1}{4}g^2v^2W^+W^- + \frac{1}{8}(g^2 + g'^2)v^2Z^0Z^0. \tag{2.17}$$

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

$$
\begin{aligned}
m_W^2 &= \frac{1}{4} g^2 v^2 \\
m_Z^2 &= \frac{1}{4} (g^2 + g'^2) v^2 \\
m_A^2 &= 0
\end{aligned}
\tag{2.18}
$$

2.3 Deficiencies of the Standard Model

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Chapter 3

Supersymmetry

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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

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4.1 Magnets

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Chapter 5

The ATLAS detector

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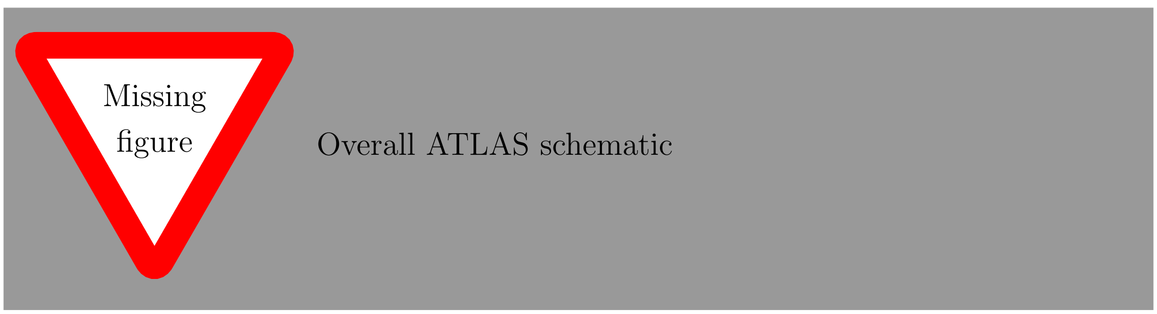

5.1 Inner Detector

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Pixel Detector

Missing figure Pixel detector schematic

Insertable B-Layer

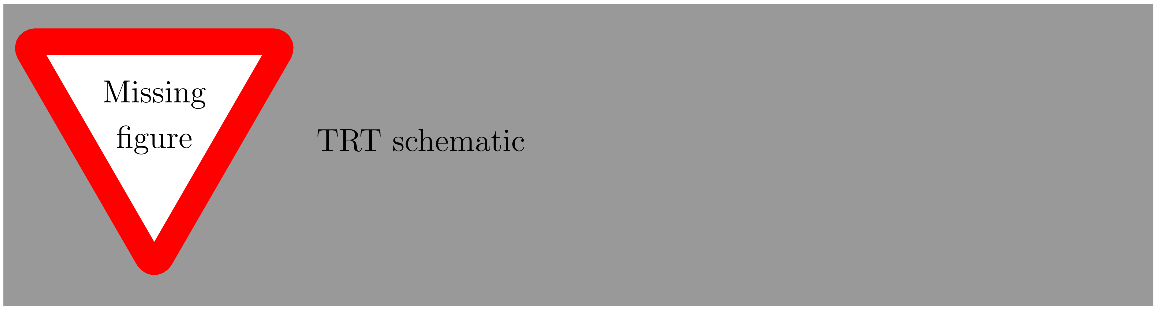
Qualification task, so add a bit more.

Missing figure IBL ROD/BOC schematic

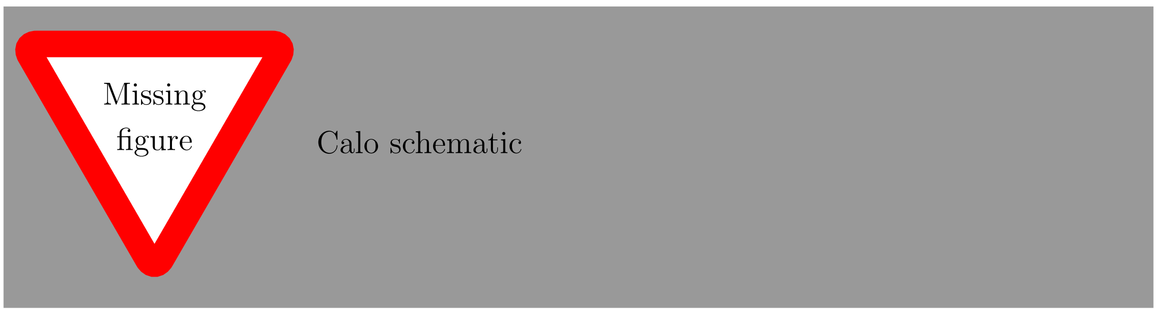
Semiconductor Tracker

Missing figure SCT schematic

Transition Radiation Tracker



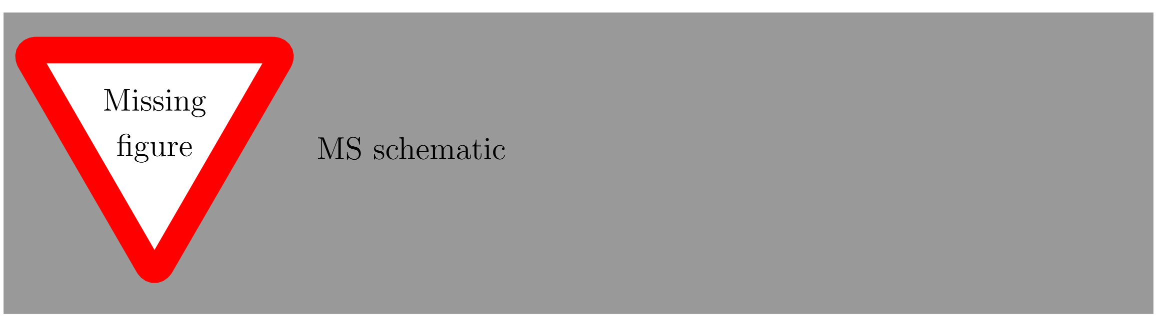
5.2 Calorimeter



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

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9.1 Statistical Analysis

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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

cite Yuval's lectures and notes somehow

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

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

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

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

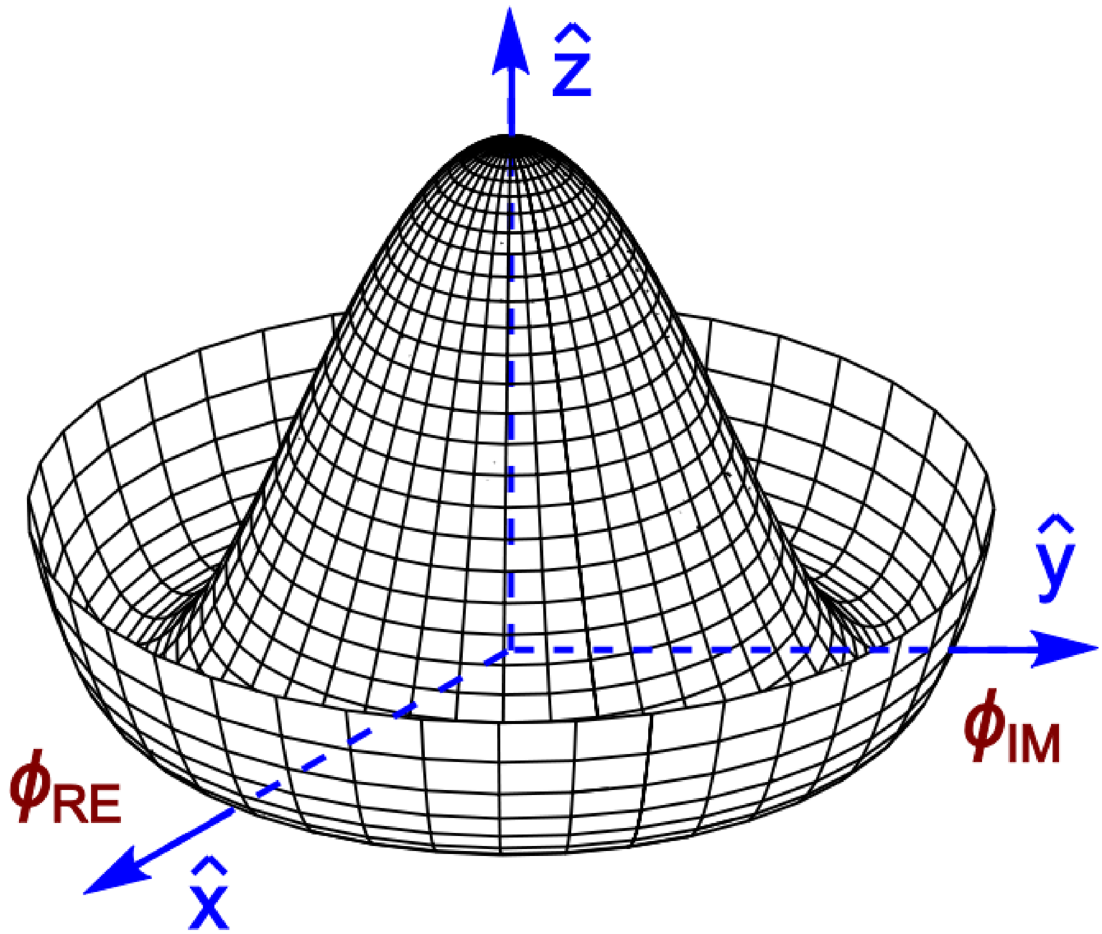
First, note that the theory is only stable when $\lambda > 0$. To understand the effect of SSB, we now enforce that $\mu^2 < 0$, and define $v^2 = -\mu^2/\lambda$. We can then write the scalar potential of this theory as :

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

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

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

Figure 1: Sombrero potential



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

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

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U(1) local symmetry breaking