Phenomenology of Extra-Dimensional Higgsless Models

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Master thesis in High Energy Physics for the University of Twente, performed at the NIKHEF institute in Amsterdam. Under supervision of Prof. Bob van Eijk
Abstract

Since the minimal Higgs mechanism used in the SM suffers from several problems a wide range of Higgs sector expansions and alternatives exist. The MSSM predicts a non-minimal Higgs scenario with a charged Higgs, which is one subject of this thesis. The charged Higgs study serves as a test for the main topic: the phenomenology of Higgsless models. Higgsless models provide an alternative for the Higgs mechanism; they are a group of extra dimensional theories that have been expanded to include electroweak symmetry breaking. These models contain an infinite tower of Kaluza-Klein excitations of the SM vector bosons. Although the theoretical interest in these models is growing, the phenomenological and experimental aspects of the new, hypothetical particles have not extensively been studied. This thesis contains a first step towards a full detector study of the fermiophobic Z' particle. A Monte Carlo implementation of the two production processes is presented along with a scheme for a correct treatment of polarisations. The results show that the production cross sections are of the order $100\,\text{fb}$. This in combination with several distinguishing signal characteristics make a search for the Z' particle at the LHC promising. The charged Higgs study illustrates the use of spin effects for scalar/vector boson discrimination, which is also of interest for the Higgsless models. No detailed phenomenological spin analysis could be performed on the Z' simulations, however, a study of several of its basic features is presented. The results of these studies show a future dedicated detector study to be useful.
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Chapter 1

Elementary particle physics

The atom, even though its name suggests so, is not an elementary particle. It consists of a nucleus surrounded by a cloud of electrons. This was already discovered in the early 20th century by Rutherford [1]. After his discovery the nucleus itself was proven to consist of protons and neutrons, which were then again proven to consist of particles called quarks. Up till now there is no experimental proof that the electrons and quarks themselves contain even smaller building blocks. We can therefore say that all ordinary matter can be broken down to quarks and electrons. However, these are not the only particles present in nature, they make up only part of one of at least three generations of particles. Apart from these three families there also exist anti-matter versions of every particle, the anti-particles, and the particles that carry the different forces, the gauge bosons. Particle physics concentrates on all these particles and the forces through which they interact.

1.1 The Standard Model

The Standard Model is the theoretical model that most accurately describes the behaviour of elementary particles. It contains all of the experimentally discovered particles together with the forces with which these particles interact. The complete Standard Model can be summarised with the two diagrams presented in figures 1.1 and figure 1.2. Figure 1.1 shows all the elementary particles of the SM, while figure 1.2 presents how the different forces of nature act upon which groups of particles. In this section the three different interactions will be discussed with a special emphasis on the weak interaction, since it plays an important role in this thesis. The Higgs mechanism, used to break electroweak symmetry in the SM, will be discussed in a separate section.

1.1.1 Gauge bosons

Within elementary particle physics forces are described by particles, the gauge bosons. The way this is done is by quantising the force fields. The forces acting between two objects are replaced by a continuous flow of force particles. A way to imagine this is
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Figure 1.1: The particles that make up the Standard Model.

Figure 1.2: Diagram showing the interactions between the particles of the Standard Model.
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with the often used example of two people throwing a ball to each other. Every time the ball is thrown the two people exert a force, which pushes them apart, this is a repulsive force. In a similar way one could, with some imagination, think of a force acting the other way around, an attractive force. One example of an attractive force would be that between a positive and a negatively charged particle. The force acting there is electromagnetism, carried by photons. So for electromagnetism the photon is the gauge boson, the quantised version of the electromagnetic force field.

Next to the electromagnetic interaction, there is the strong interaction, carried by gluons, the weak interaction, carried by vector bosons called W and Z and gravity. The latter has not yet been incorporated in the Standard Model. This means that we do not know how to quantise gravity in a satisfactory way and therefore do not know which, if any, particles are responsible for this force.

1.1.2 Electromagnetism

After gravity electromagnetism (EM) is the most encountered force in every day life and is responsible for phenomena such as light and electricity. At the atomic scale this force keeps different atoms apart because the electrons from the different atoms exert a repulsive force. Photons are exchanged between particles with electric charge, which, as can be seen in figure 1.1, are most of the fermions of the SM and the W bosons.

1.1.3 Strong interaction

The strong interaction, exchanged between quarks by gluons, is the interaction responsible for keeping nuclei together. The relative strength of 100 times that of the electromagnetic force, makes it possible for the strong to overcome the repulsive electromagnetic force between the protons in the nucleus. The property that causes the quarks to bond to the gluons is called colour, it can be seen as the equivalent of charge in electromagnetism. Where every positive electric charge attracts a negative charge using photons, quarks with different colour attract or bind through gluons. In total there are the three primary colours, red, blue and green. Every quark has one of these colours and is attracted to a quark with one of the other colours. When three of the different quarks are combined they make up a so called white or colour neutral nucleon. Every particle found in nature is colour neutral. They always have to either be part of a baryon, a combination of 3 quarks like the proton, or a meson. Where a meson is a combination of a quark and an antiquark carrying anti-colour.

One important feature of the strong interaction is its asymptotic freedom, this is the property which causes the strong force to become weaker with decreasing distance. The other extreme is that for increasing distances the force goes to infinity. This is the reason why coloured particles can never be found in nature. When one for example creates a quark antiquark pair, both with a very high energy, they will not move away from each other simply because the attracting force becomes stronger and stronger. This continues until the energy stored in the gluon between the two is high enough for it to be able to form another quark antiquark pair. Resulting in 2 mesons instead of
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...the highly energetic one we started with. This in combination with gluon radiation can cause one high energetic quark to create a whole range of hadronic particles, called a jet.

1.1.4 Weak force

In the process of $\beta$-decay a neutron decays into a proton, an electron and a neutrino. For this an interaction different from the electromagnetic force and the strong force is needed, because electromagnetism does not couple to the neutron and gluons do not couple to leptons. The responsible interaction here is the weak force. It has a strength of $10^{-3}$ relative to the electromagnetic force and is mediated by the vector bosons. There are two sorts of vector bosons, the electrically neutral Z bosons and the W’s, charged gauge bosons with a charge of $\pm 1$. The last particle in the $\beta$-decay process, the neutrino, only interacts through the weak force, for it has charge nor colour and is therefore hard to detect in experiments. The reason why the weak force is not experienced in every day life is that its gauge bosons are, unlike those from the electromagnetic and strong force, massive. Therefore, Z’s and W’s mainly occur at high energies and the chance of having a virtual one at low energies is much smaller than for example having a photon.

One especially important feature of the weak force is that it violates parity. The result of parity violation is that couplings to right-handed particles (particles with spin and momentum antiparallel) are absent for the massive gauge boson. Where the EM and the strong force do not distinguish between left or right-handed fermions, the weak force fully violates the symmetry in parity. It never couples to a right-handed particle, as a result only left-handed neutrinos exist in the SM. The importance of this will be seen in the following chapters of this thesis.

In the electron positron scattering process, both the photon and the Z boson can be the force carrier since both couple to the particles involved. However, when this happens at low energies, $E << M_Z$, the chances of having a Z boson in the the process are small. For higher energetic incoming particles the chances of creating a Z boson increase and discrimination between photons and Z’s becomes impossible. Here the weak and electromagnetic force combine into the electroweak force (EW). The two forces unify and processes evolve through the exchange of a superposition of the photon and the Z. The unification is illustrated in figure 1.3, which shows the theoretically extrapolated strength of the different forces as a function of the energy.

1.2 The Higgs mechanism

One of the problems with the SM as explained up to this point is that it can not account for the masses of the weak gauge bosons. For it to remain a renormalisable theory, a theory in which everything stays finite, all fermions and bosons must be massless, which is not what is observed in nature. A second problem arises when calculating scattering processes between W and Z bosons. Exponential growth of the amplitude with increasing energy results in violation of unitarity. This basically means that the
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Figure 1.3: Relative strength of the different forces as a function of the energy (or temperature).

probability of the scattering process becomes larger than 1, an unphysical result. The mechanism incorporated in the SM that cures both these problems in an elegant way is the Higgs mechanism [2]. The first part of this section contains an introduction to gauge symmetries to understand the problems inherent to a Higgsless Standard Model. This is followed by a description of the Higgs mechanism and its consequences. This section will end with a short discussion of the several problems with the Higgs mechanism.

1.2.1 Gauge symmetry

In quantum mechanics the amplitude for a particle to move from say A to B is described by a path integral. This is an integral over all possible paths going from A to B where each path is weighted by its action $S$, the space time integral of the Lagrangian. One example of such a Lagrangian in field theory is:

$$\mathcal{L} = \frac{1}{2} (\partial_{\mu} \phi)(\partial^{\mu} \phi) - \frac{1}{2} m^2 \phi^2$$

which after substitution into the Lagrange equation leads to the Klein-Gordon equa-
tion that describes the motion of spinless particles. The Dirac equation, which describes the motion of particles with spin \( \frac{1}{2} \), can be obtained from the Lagrangian:

\[
\mathcal{L} = i \bar{\psi} \gamma_\mu \partial^\mu \psi - m \bar{\psi} \psi
\]

Where \( \psi \) is the wavefunction for a spin \( \frac{1}{2} \) particle. Since the physics should not depend on the phase of the wave function the Lagrangian must be invariant when adding a phase factor of the form \( e^{i\alpha} \) to \( \psi \). In other words, the Lagrangian is invariant under phase transitions or, since these phase transitions form a so called unitary group, invariant under global \( U(\alpha) \) transformations. A symmetry implies a conserved quantity and vice versa, a connection first shown by Noether in [3]. For example physical invariance under spatial translations can be shown to lead to conservation of momentum. Requiring \( \mathcal{L} \) to obey \( U(1) \) symmetry leads to conservation of charge.

It is also possible to construct Lagrangians that are invariant under local transformations, that is transformations containing a space-time dependent part. The Lagrangian should then not change under for example a local \( U(1) \) transformation:

\[
\psi(x) \rightarrow e^{i\alpha(x)} \psi(x)
\]

Examining:

\[
\mathcal{L} = i \bar{\psi} \gamma_\mu \partial^\mu \psi - m \bar{\psi} \psi
\]

it is straightforward to see that this Lagrangian is not invariant under these \( U(1) \) transformations due to the presence of a derivative. One way to solve this is by replacing the normal derivative by a covariant derivative, which transforms as:

\[
D_\mu \psi \rightarrow e^{i\alpha(x)} D_\mu \psi(x)
\]

This is accomplished by choosing \( D_\mu \equiv \partial_\mu - ieA_\mu \), where \( A_\mu \) transforms as \( A_\mu \rightarrow A_\mu - ieA_\mu \). Substituting this covariant derivative into the Lagrangian results in:

\[
\mathcal{L} = i \bar{\psi} \gamma_\mu D^\mu \psi - m \bar{\psi} \psi = \bar{\psi}(i\gamma_\mu \partial^\mu - m) \psi + e \bar{\psi} \gamma^\mu \psi A_\mu
\]

By demanding the Lagrangian to also obey local symmetries, the Lagrangian gains an additional term. In this new term \( \psi \) could be a fermion wave function, the term \( \bar{\psi} \) an antifermion wave function, the two are connected by a new field \( A_\mu \). \( A_\mu \) is known as the gauge field. For this particular example \( A_\mu \) should be interpreted as the photon field. Therefore, imposing local gauge symmetry has lead to a modified Lagrangian containing the kinetic terms of the electron (or fermions in general) and a new term describing the coupling between the fermions and gauge field. The addition of a kinetic term for the gauge bosons is also allowed as long as it is invariant under local \( U(1) \) symmetry. For the electromagnetic Lagrangian this extra term would be \(-\frac{1}{4} F_{\mu\nu} F^{\mu\nu}\) where \( F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \). Combining all of these elements we arrive at the total Lagrangian of quantum electrodynamics (QED):
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\[ \mathcal{L} = \bar{\psi}(i\gamma_\mu \partial^\mu - m)\psi + e\bar{\psi}\gamma^\mu \psi A_\mu - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} \]

From this QED Lagrangian one can derive the so called Feynman rules for quantum electrodynamics. Each term in the Lagrangian can be associated with a propagator or a vertex factor (an interaction between different particles). Propagators are represented by all terms containing two field terms. For the QED Lagrangian above \( \bar{\psi}(i\gamma_\mu \partial^\mu - m)\psi \) would be the propagator for the fermions. In a Feynman diagram this would be represented as the left part of figure 1.4. The terms with more than two field elements can be associated with a vertex (figure 1.4 right). In the QED Lagrangian, this is \( e\bar{\psi}\gamma^\mu \psi A_\mu \), which describes an interaction, or coupling, between two fermions and a photon \( A_\mu \). One can now combine these in all possible ways to create Feynman diagrams for interactions, such as the electron-positron scattering diagram shown in figure 1.5.

Figure 1.4: Feynman diagram representations of a propagator (left) and a vertex (middle).

Figure 1.5: One Feynman diagram representation of electron electron scattering.

For quantum chromo dynamics (QCD), the theory of the strong interaction, we can derive a Lagrangian in a similar way. The major difference between QED and QCD is that while QED has to obey \( U(1) \) symmetry QCD obeys \( SU(3) \) symmetry. The transformations are part of the special (S), unitary (U) group. Transformations from the \( U(1) \) group can be represented by 1x1 unitary matrices, those for \( SU(3) \) can be represented as unitary 3x3 matrices with determinant 1. The group is non-Abelian, implying self interactions of the gauge bosons. The Lagrangian for QCD is:

\[ \mathcal{L} = \bar{q}(i\gamma_\mu \partial^\mu - m)q + g(\bar{q}\gamma^\mu T_a q)G^a_\mu - \frac{1}{4}G^a_\mu G^a_\mu \]

Where \( T_a \), with \( a \) going from 1 to 8, are unitary traceless matrices, stemming from the property that there are 8 different gluons (differing in colour anti-colour combination). \( G^a_\mu \) is the gauge boson field that transforms as \( G^a_\mu \rightarrow G^a_\mu - \frac{1}{4}G^a_\mu \partial_\mu, \) - \( f_{abc} \alpha_\beta G^a_\mu. \)
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Of course we can now try to apply a similar procedure for the electroweak force, the combination of the electromagnetic and weak interaction, which should obey $SU(2)_L \otimes U(1)_Y$. Where the subscript L indicates the left-handedness of the weak interaction. When unifying these two interactions the four separate gauge bosons are written as two doublets. This is done by mixing the two W’s and writing the mixtures as $W^1_\mu$ and $W^2_\mu$, these are defined as:

$$W^\pm_\mu = \frac{1}{2}(W^1_\mu \pm W^2_\mu)$$

and the Z and the photon as a combination of the two fields $W^3_\mu$ and $B_\mu$:

$$Z_\mu = \frac{gW^3_\mu - g'B_\mu}{\sqrt{g^2 + g'^2}}$$

$$A_\mu = \frac{g'W^3_\mu + gB_\mu}{\sqrt{g^2 + g'^2}}$$

The covariant derivative of the electroweak interaction can now be written as:

$$D_\mu = \partial_\mu - \frac{1}{2}ig\tau W_\mu - \frac{1}{2}ig'YB_\mu$$

Where the $\tau$ are the $SU(2)$ generators, the Pauli matrices multiplied by $\frac{1}{2}$ and the $Y$, called hypercharge, is the generator of the $U(1)$ part.

Just like for the QED and QCD theories a Lagrangian can be created using this covariant derivative. However, when looking at the Lagrangians of QED and QCD there are no mass terms present for the gauge bosons. For these interactions this is of course not necessary since both gluons and photons are massless. The gauge bosons of the weak interaction are not, they are in fact very heavy. An additional term to account for this mass can be added by hand and would be of the form $MB_\mu B^\mu$ where $B_\mu$ is the gauge boson field. However, under transformations this new term will never be invariant and lead to problems like divergences. Ignoring this problem is a solution, however, when using this Lagrangian to calculate processes containing loops as in figure 1.6 one will end up with terms diverging with increasing energy. So a solution needs to be found that keeps these processes renormalisable for high energies. A solution which does this is called spontaneous symmetry breaking, opposed to adding the term by hand, which is called explicit symmetry breaking. One way to break the symmetry spontaneously is with the Higgs mechanism. This mechanism introduces at least one new scalar particle, the Higgs particle. This additional particle is not only needed for spontaneous symmetry breaking but also prevents the EW force from becoming strong at high energies, due to additional interfering diagrams it adds to the process of vector boson scattering, as will be discussed later.
1.2.2 Electroweak Symmetry Breaking

For the introduction of mass terms in the Lagrangian in a satisfying way, EW symmetry needs to be broken spontaneously. The Higgs mechanism accomplishes this is by adding a new field, a complex weak $SU(2)$ scalar (spin 0 particle) doublet $\Phi$:

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}$$

which can be introduced in the Lagrangian with the covariant derivative from $SU(2)_L \otimes U(1)_Y$ gauge symmetry discussed in the previous section:

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

The important part of this Lagrangian is the potential $V(\Phi)$, which, whilst remaining gauge invariant, may have the form:

$$V(\Phi) = \mu^2 |\Phi|^2 + \lambda |\Phi|^4$$

with $\lambda > 0$. Depending on the sign of $\mu^2$ this can have two different shapes. The first, with $\mu^2$ positive, has its minimum at $\Phi = 0$, the other has an infinite amount of minima at $\Phi = \pm \nu$ with $\nu = \sqrt{-\mu^2/\lambda}$ as can be seen in figure 1.7. By choosing the solution with an infinite amount of minima the symmetry is broken. The fields are then no longer expanded around 0 but rather around one of the possible minima that lie at a distance $\nu$. One now has to choose one of these minima, resulting in a choice of the form of the scalar doublet. For the Higgs mechanism this is the following:

$$\phi_1 = \phi_2 = \phi_4 = 0$$

and

$$\phi_3 = \nu$$

This results in $\frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \nu \end{pmatrix}$, which, when expanding around this minimum becomes $\Phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \nu + h(x) \end{pmatrix}$. Used in the Lagrangian this leads to:

$$\mathcal{L} = \left| -\frac{1}{2} i g \vec{W}_\mu - \frac{1}{2} i g' B_\mu \right| \Phi^2 - V(\Phi)$$

$$= \frac{1}{8} \nu^2 \mu^2 \left( (W_\mu^1)^2 + (W_\mu^2)^2 \right) + \frac{1}{8} \nu^2 \left[ (g W_\mu^3 - g' B_\mu)^2 \right] - V(\Phi)$$
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This can be rewritten in terms of the actual gauge bosons resulting in a Lagrangian with a mass term for $W^\pm$ of the form $\frac{1}{2}\nu g$, a mass term for $Z$, $\frac{1}{2}\nu\sqrt{g^2 + g'^2}$ and no mass term for the photon. Complying with the requirements of the SM. So by adding a scalar doublet with its own potential, the mass terms for the gauge bosons appear spontaneously. A subsequent result of adding the complex scalar doublet is a new particle from the expansion around $\phi_3$, the $h(x)$ field. By giving the gauge bosons mass they also gain a possible polarisation state, longitudinal polarisation. This is the reason why there is only one new scalar in the Lagrangian while a complex doublet was introduced. The other 3 fields, the Goldstone bosons, are absorbed in the gauge bosons, they are converted into the the polarisation states of the $W^\pm$ and the $Z$.

The one surviving field introduced with this mechanism, the scalar Higgs boson, couples to every massive SM particle and to itself via a triple and a quadruple coupling. Its mass follows from the Lagrangian, $M_H = \nu\sqrt{2\lambda}$. In this term the value of $\nu$ is known from measurements of muon decay, however, the value of $\lambda$ is completely unknown and cannot be derived from any other SM constant. This means that the mass of the hypothetical Higgs particle is not known and that future searches for this particle have to scan a large mass region.

1.2.3 Vector Boson Scattering

Apart from breaking the electroweak symmetry the Higgs mechanism solves a second problem of a Higgsless SM, the problem that the electroweak force becomes strong at high energies. In the scattering processes of longitudinally polarised gauge bosons, for example $W_LW_L \rightarrow W_LW_L$, the cross section suffers from terms diverging with increasing energy. The different first order Feynman diagrams contributing to this
The amplitudes corresponding to these different diagrams are [4]:

\[(a) \ A_{s\gamma} = -\frac{1}{16}ie^2s^2\beta^2(3-\beta^2)^2\cos\theta,\]

\[A_{sZ} = -\frac{1}{16}ig_{WWZ}\frac{s^3}{s - \frac{M_Z^2}{M_W^2}}\beta^2(3-\beta^2)^2\cos\theta,\]

\[(b) \ A_{t\gamma} = -\frac{ie^2s^3}{32t}\left[\beta^2(4-2\beta^2+\beta^4)+\beta^2(4-10\beta^2+\beta^4)\cos\theta+(2-11\beta^2+10\beta^4)\cos^2\theta+\beta^2\cos^3\theta\right],\]

\[A_{tZ} = -\frac{ig_{WWZ}s^3}{32(t - \frac{M_Z^2}{M_W^2})}\left[\beta^2(4-2\beta^2+\beta^4)+\beta^2(4-10\beta^2+\beta^4)\cos\theta\right.\]

\[+\left.(2-11\beta^2+10\beta^4)\cos^2\theta + \beta^2\cos^3\theta\right]\]

\[(c) \ A_4 = \frac{1}{16}ig_{WWWW}s^2(1 + 2\beta^2 - 6\beta^2\cos\theta - \cos^2\theta),\]

Where \(t = -\frac{1}{2}s\beta^2(1 - \cos\theta)\) and \(\beta = \sqrt{1 - 4/s}\) and the labels \(s\) and \(t\) refer to the \(s, (s = [p_1 + p_2]^2)\) and the \(t\) (\(t = [p_1 - p_3]^2\)) channel respectively. When adding these amplitudes the total amplitude can be approximated for \(s >> 1\) by:

\[A_1 \approx \frac{1}{4}ig^2 \left(-s - t + 2\frac{M_Z^2}{M_W^2} + \frac{2t}{s}\left(\frac{M_Z^2}{M_W^2} - 4\right) + 8\sin\theta_W s \left(\frac{1}{t} - \frac{1}{t - \frac{M_Z^2}{M_W^2}}\right) + \frac{2\frac{M_Z^2}{M_W^2} s}{t - \frac{M_Z^2}{M_W^2}}\right)\]

As can be seen this is proportional to \(s\). This means that the cross section, proportional to \(|A|^2\), will grow with \(s^2 \approx E^4\) leading to problems at high energies. Typically
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Figure 1.9: The first order Feynman diagrams contributing to WW scattering with a Higgs propagator [4].

unitary is violated above 1.8 TeV. The solution comes from the Higgs mechanism, which adds two extra diagrams to the process, diagrams (a) and (b) of figure 1.9. The amplitudes corresponding to these two diagrams have to be added to the total amplitude, they are:

\[
(a) \quad A_{sH} = -\frac{1}{16} i g^2 s^2 (1 + \beta^2)^2 \frac{1}{s - m_H^2 / m_W^2 + i M_H \Gamma_H / M_W^2},
\]

\[
(b) \quad A_{tH} = -\frac{1}{16} i g^2 s^2 (\beta^2 - \cos \theta)^2 \frac{1}{t - m_H^2 / m_W^2 + i M_H \Gamma_H / M_W^2},
\]

Resulting in:

\[
A_{\text{total}} = A_1 + A_{sH} + A_{tH} \approx -\frac{1}{2} i g^2 \left[ \frac{M_Z^2}{M_W^2} \left( 1 + \frac{s}{t} + \frac{t}{s} \right) + \frac{M_H^4}{M_W^2} - i \frac{M_H \Gamma_H}{M_W^2} \right]
\]

For large values of s, the terms growing with s (or t) are now cancelled. In this way the Higgs mechanism is used to keep the process of longitudinal vector boson scattering well behaved for all energies. So this relatively simple mechanism solves both of the major problems of the Higgsless SM. However as will be shown next the Higgs Mechanism itself also suffers from several problems.

1.2.4 Problems with the Higgs mechanism

The Higgs particle has not yet been discovered and as mentioned before, there is no direct theoretical prediction of its mass. There are however indirect measurements of its mass. For example measurements of the contributions the Higgs makes through loop diagrams to the propagators of particles have been performed. Calculations of these contributions together with results from earlier direct searches by the LEP experiment at CERN have lead to the mass prediction shown in figure 1.10. This figure shows that there is a region up to 114 GeV that is, with a 95% certainty, excluded by the LEP experiments. However, the indirect measurements indicate that the mass is most likely not much higher than this.
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Figure 1.10: Probability predictions of $m_H$ from the LEP experiment. The full line is the result of the fit using all high energy data, together with an error band representing an estimate of the theoretical error as a result of missing higher order corrections. The vertical band shows the 95% confidence exclusion limit on $m_H$ from the direct search. The dashed line represents the results using a different value of $\delta\alpha_{\text{had}}$, the dotted line gives the results when low energy data is also included. [5].

These indirect measurements do not proof that there has to be a Higgs particle, they only indicate that there are contributions to certain processes. This could be a Higgs particle, ideally one with a mass around 120 GeV, but it is also plausible that the explanation comes from other unknown particles.

The indirect influences of particles on one another also leads to a problem with the Higgs itself, the so called hierarchy problem. If the Higgs boson exists with a mass below one TeV then, according to some, an unnatural situation occurs. The measured mass is then many order sizes smaller than the "bare" mass. The 120 GeV is only the apparent mass. This is the real or bare mass together with all the contributions to it from SM particles.

In the SM every fermion makes the following contribution to the Higgs mass [6]:

$$\Delta m_H^2 = -\frac{|\lambda_f|^2}{8\pi^2}[\Lambda_{UV}^2 + \ldots]$$

Where the $\Lambda_{UV}^2$ is the Ultra Violet cutoff, which can be seen as the energy up to which the SM should be valid. Usually this value is taken to be the Planck Scale, $M_{PL}$. 

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1.2. THE HIGGS MECHANISM

Here $\lambda_f$ is the Yukawa coupling of the Higgs boson to the fermion, which is proportional to the fermion mass. For a boson the contribution is [6]:

$$\Delta m_H^2 = \frac{|\lambda_s|}{16\pi^2}[\Lambda_{UV}^2 + ...]$$

Here the contribution to the squared Higgs mass only goes with $\lambda_s$ and is positive instead of negative. The result is that the main contribution to the measurable Higgs mass comes from the heaviest SM fermion, the top quark. If Planck Scale is taken as the Ultra Violet cutoff, the contribution of the top quark is an order 30 larger than this squared mass itself. A difference of this order can seem unnatural, leaving many physicists to question the validity of the SM.

The most popular solutions to this problem are supersymmetric models. These models, of which the MSSM (which is discussed in the next chapter) is the simplest, make use of the difference in sign between the contributions of bosons and fermions. The model introduces a bosonic super partner for every SM fermion and a fermionic super partner for every SM boson. This is done in such a way that all the different contributions cancel each other. These models however are not without their own short comings. The most obvious being that the new particles do not have the same mass as their SM partners. If the masses were identical supersymmetry would already have been detected in past experiments. So one immediately has to assume that this symmetry is broken. The energy scale where this breaking occurs cannot be too high for this mechanism to work. If this scale becomes too high, the mass differences become too large and sufficient cancellations no longer occur. Therefore, the scale at which this breaking takes place creates its very own hierarchy problem, this scale is itself much smaller than the Planck scale.

A second problem with the Higgs mechanism is of a cosmological origin [7]. The problem arises from the non negative vacuum expectation value associated with the Higgs. The result is a non zero energy-momentum density in the universe that is theorised to interact with gravity causing a curvature of space-time. Using the expected Higgs mass from figure 1.10 to calculate the expected curvature one obtains a result that differs greatly from the Euclidean universe observed by experiments. One way to get around this problem is by assuming that the universe was counter curved right after its creation, before electroweak symmetry was broken. This justification implies that the symmetry breaking would exactly cancel this counter curvature. Of course this sounds like a contrived and therefore unsatisfactory solution to the problem.

Both of these problems together with the fact that the Higgs boson has yet to be detected, leave room for competing theories. Examples of this are the already mentioned MSSM, studied in the next chapter, complex Technicolour models, top quark condensate models and the group of theories that form the main subject of this thesis, extra dimensional Higgsless models.
Chapter 2

Charged Higgs

2.1 Charged Higgs from the MSSM

The SM Higgs mechanism introduced in the previous chapter is, as mentioned, only one of many solutions for EWSB. In the mechanism used in the SM only one boson is introduced, similar results can however be achieved using more bosons. For models that introduce two or more doublets of bosons one of the new scalars needs to have charge $\pm 1$ [6]. One of the more popular theories using a non-minimal Higgs scenario is the Minimal Super Symmetrical Model (MSSM). It uses 4 Higgs bosons to break EWS, the details of which can be found in [6]. The theoretical limits on this model indicate that if MSSM exists the LHC should be able to find its first signatures.

The new scalar with electric charge $\pm 1$ of the MSSM is the charged Higgs, $(H^\pm)$. This $H^\pm$ is allowed to have a mass of approximately $m_{W^\pm}$ [8], making a discrimination between the two difficult. If $m_{H^\pm}$ is indeed low, the only major differences between $H^\pm$ and $W^\pm$ are their branching ratios and their spin. The scalar $H^\pm$ has spin 0, whereas the $W^\pm$ is a vector boson with spin 1. This will result in a different distribution of their decay products, as was shown by [9]. The Higgsless models, which are the main subject of this thesis, also contain particles only differing in their spin from the SM Higgs. It is interesting to investigate the results of these spin differences for the MSSM first since theoretical predictions of this already exist. In this chapter we will test these theoretical predictions and try to exploit the spin differences to make a faster discovery of the charged Higgs at the LHC possible. For this a detailed study was performed with the theory department of the DESY institute in Hamburg. Since the results for the MSSM study and those for the later discussed Higgsless models are acquired using Monte Carlo simulations, a short introduction the used software packages will be presented in the next section. This will be followed by a discussion of the important parameters for a charged Higgs study and the results of the performed simulations.
2.2 Monte Carlo Techniques

2.2.1 PYTHIA

For the study presented in this chapter, as for those in the next chapters, Monte Carlo simulations were performed. The software package used for this throughout this thesis is PYTHIA [10]. PYTHIA has the ability to simulate a wide range of both SM and Beyond the Standard Model (BSM) processes using the Monte Carlo method. It factorises each collision event in several parts. Most importantly there is the hard process, for example $q\bar{q} \rightarrow ZW$, for which PYTHIA contains the expressions of the squared amplitude. It also handles all the other processes within a collision: the structure functions of the incoming particles, the colour flow, the initial and final state radiation and the decay chains of the final state particles. By handling each part separately it calculates complete events both accurately and fast. The output of PYTHIA can be used as input for a detector simulation with which predictions can be made for experiments. One limitation of PYTHIA is that spin information is lost for the decay products. PYTHIA loses this information because it uses the spin averaged amplitude. Final state particle spin effects are accounted for by assigning a weight to each event. This weight is calculated separately using the individual spin dependent amplitudes. This is however only done for the final state particles so for further generations of decay products the spin information is lost. Since the interesting effects of the charged Higgs polarisation carry through beyond its first generation of decay products, PYTHIA can not be used on its own to produce the simulations presented in this chapter.

2.2.2 TAUOLA

For decays including $\tau$ leptons the spin problem can be solved with the TAUOLA package [11] that can be coupled to PYTHIA. TAUOLA takes the $\tau$’s from PYTHIA as input and then handles the complete decay chains up to the stable particles using the spin dependent amplitudes. For each event the TAUOLA data is then returned to PYTHIA. Since for majority of the studied parameter space the $H^\pm$ mainly decays through $\tau$, this package was required for the simulations discussed in this chapter.

2.2.3 CompHEP

A second Monte Carlo software package is CompHEP [12]. This was used in the study of the Higgsless models presented in the next chapters. The main difference between this package and PYTHIA is that CompHEP takes the Lagrangian of a theory as input and calculates the squared amplitudes of every possible hard process from this Lagrangian. This way a wide range of processes is directly accessible, making this package preferable to PYTHIA for background simulations. The study of new theories is also made easier since one only needs its Lagrangian, whereas for PYTHIA every single new process must be implemented individually. The main limitation of CompHEP is that it only handles everything up to the hard process. When a full collision simulation is required PYTHIA can be used to further process the results of CompHEP.
CHAPTER 2. CHARGED HIGGS

2.3 Production and Decay of the Charged Higgs

![Diagram of Feynman diagrams for charged Higgs production and decay](image)

Figure 2.1: The first order Feynman diagrams for the production mechanisms of $t\bar{t}$ through gluon fusion and quark antiquark fusion.

Using LEP and Tevatron data an experimental lower bound of approximately $m_{W^\pm}$ [8] has been set on $m_{H^\pm}$. However, there are no tight upper constraint on its mass or its couplings to the SM particles. Therefore, different mass regions need to be probed to find it, all of which require a different search strategy. In this chapter the strategy for finding a charged Higgs with a mass below that of the top quark will be discussed first. In this region the branching ratio of the $H^\pm$ into $\tau$ is almost a 100% and the dominant production process is top decay ($t \to H^+ + b$). Therefore, the LHC process of interest is top pair production via either quark antiquark or gluon-gluon fusion, figure 2.1, with the subsequent decay channel $t \to H^+ b \to \tau \nu + b$. The other top quark will be used for triggering and is therefore required to decay via $W^+ b \to \mu \nu + b$. The main background for this is the process where both partners of the top quark pair decay via the $W^\pm$ channel. Especially for $m_{H^\pm} \approx m_{W^\pm}$ these processes are difficult to distinguish. Discovery is however possible using an excess of $\tau$ decay events which originates from the almost 100% BR of the $H^\pm$, for the $W^\pm$ this is only 10.8%.

The goal of this study is to see if the difference in spin between the charged Higgs and the vector bosons can be used to make a discovery of the charged Higgs easier. The scalar nature of the charged Higgs makes the $\tau$'s produced from its decay left-handed whereas a $\tau$ from $W^\pm$ decay will always be right-handed (see figure 2.2). Subsequent decay of the $\tau$ will then result in a difference in kinematical distributions for both processes. The channels of interest for this research are the one-prong decay channels of the $\tau$. These are: $\tau^+ \to \pi^+ \pi^0$, $\tau^+ \to \rho^+ \pi^0 \to \pi^+ \pi^0 \pi^0$, $\tau^+ \to a^+ \pi^0 \to \pi^+ \pi^0 \pi^0$. The result of prior studies [9], [13] show that the $\pi^+$ coming from the charged Higgs on average acquire a higher portion of the energy than in the case of a $W$ decay. With the use of this characteristic it might be possible to improve the signal to background ratio
by applying a lower cutoff in the charged pion energy spectrum. In case of a charged Higgs with a mass above that of the W boson this effect will be increased by the extra momentum boost the $\tau$ gets from the available extra energy.

Figure 2.2: Decay differences between charged Higgs and W bosons, the charged Higgs will always result in a left-handed lepton whereas the W will result in a right-handed lepton.

The polarisation strategy will be extended to masses above $m_t$. In this region a relatively high value of $\tan \beta$, the ratio between the vacuum expectation values of the two Higgs doublets, is required. This parameter influences the branching ratios of $H^\pm$ as can be seen from equation (2.1). Here $V_{ij}$ are the CKM matrix elements. As shown in equation (2.1), the branching ratio of $H^\pm$ to $\tau$ increases with increasing $\tan \beta$ while that to the quarks decreases. For low values of $\tan \beta$ ($\tan \beta < 10$) the BR to $\tau \bar{\nu}$ will be very small for $m_{H^+} > m_t$. Both for the high and the low mass region 4 different values of $\tan \beta$ ($\tan \beta = 10, 20, 40, 60$) were investigated.

\[
\mathcal{L} = \frac{g}{\sqrt{2} m_W} H^+ [\cot \beta V_{ij} m_{u_i} \bar{u}_i d_j L + \tan \beta V_{ij} m_{d_j} \bar{u}_i L + \tan \beta m_{L_j} \bar{\nu}_j L + H.c.] \quad (2.1)
\]

In the next section the polarisation strategy will be applied for a charged Higgs search in the low mass region ($m_{H^+} < m_t$). Here the main background comes from $W^+ \to \tau$, with the $W^\pm$ produced by top quark decay. The relevant simulated process is therefore top quark pair production. Since both signal and background are produced by the same hard interaction process an important parameter for the signal to background ratio is the BR of the top quark. This BR has been a research subject in [14]. The values used in the simulations were taken from this paper.

In the second section $m_{H^\pm} > m_t$ will be studied. Production by top decay becomes impossible there and a different production process has to be investigated. The process with the highest cross section in this area is $bg \to H^+ t$ [15]. The main background process is again top quark pair production with one top quark decaying into $\tau$ via a $W^\pm$. This process is not only interesting for $m_{H^\pm} > m_t$ but already for $m_{H^\pm} > 140GeV$. It is therefore also used in the $m_{H^\pm} < m_t$ simulations. One important parameter for the cross section here is the proton bottom quark density, which is dependent on the beam
energy. Therefore, all simulations were performed not only with the LHC design $\sqrt{s}$ of 14 TeV. The now planned start up energy of 7 TeV and also 10 TeV, which is the planned $\sqrt{s}$ for the second year of operation were used in separate simulations.

2.4 Light Charged Higgs

2.4.1 $m_{H^\pm} = m_{W^\pm}$

When setting $m_{H^\pm} = m_{W^\pm}$ the only difference in the production mechanism of the $\tau$ is the spin of its mother particle, making this parameter choice ideal for studying polarisation effects. The simulations presented here were performed using PYTHIA 6.4 [10] in combination with TAUOLA [11] that handles the $\tau$ decays. The branching ratio of the top quark from PYTHIA was modified using figure 2.1, which is taken from [14]. Both the production process via top quark decay and $bg \rightarrow H^\pm t$ were used for this simulation, though for this parameter choice the fast majority of the charged Higgs events comes from top quark decay.

![Figure 2.3: Branching ratios of $H^\pm$ to $\tau$ from [14].](image)

To investigate the effect of the polarisation on the measurable charged pion a high statistics run was performed for $m_{H^\pm} = m_{W^\pm}$ and $\tan \beta = 20$. The detector efficiency for this process was later applied in the analysis by multiplying the number of events by a factor of 0.5 accounting for the $b$-tagging efficiency and a factor of 0.108 for the requirement that the partner of the top quark decays via a $W^\pm$ into a muon for triggering. The result of this first polarisation simulation can be seen in figure 2.4. It shows the histograms of the ratio of $E_{\pi^\pm}$ over $E_\tau$, both measured in the labframe. The
2.4. LIGHT CHARGED HIGGS

<table>
<thead>
<tr>
<th>$\tan \beta$</th>
<th>$m_{H^+} = 80$</th>
<th>$m_{H^+} = 100$</th>
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<td>5.0</td>
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</table>

Table 2.1: *Used branching ratios (in %) from [14]*

ratio of the two processes as a function of this variable clearly shows that, as predicted by [9], the $\pi^\pm$ production via charged Higgs results in more energetic charged pions than the $W^\pm$ channel.

Figure 2.4: *Histograms of the $\pi$ over $\tau$ energy ratio for $\pi$ coming from: $H^+$ (left), $W^+$ (middle) and the normalised ratio of these two, $Z'/H$ (right) for $m_{H^\pm} = 80$, $\tan \beta = 20$ and $CM=14\text{TeV}$.***

To see if this effect can be used to improve the signal to background ratio simulations were performed for $\tan \beta = 10, 20, 40, 60$. From these simulations the integrated luminosity at the LHC required for a $5\sigma$ discovery was calculated for the three relevant centre of mass energies. The results can be seen in tables 2.2, 2.4 and 2.6. The integrated luminosities were once calculated using the complete pion energy spectrum and a second time after applying the ideal lower energy cutoff in this spectrum. The ideal lower energy cutoff for $m_{H^\pm} = m_{W^\pm}$ was however calculated to be 0, implying that the polarisation effect is not large enough to make a lower cutoff favourable. The reason is that even though with increasing charged pion energy the cross section of $H^\pm$ increases with respect to $W^\pm$, the cross section itself remains low for high energies. Therefore, a cutoff in this energy spectrum results in removing the fast majority of the signal events.
2.4.2 \( m_{W^\pm} < m_{H^\pm} < m_t \)

For \( m_{W^\pm} < m_{H^\pm} \) the pions from charged Higgs decay will get an extra mass boost compared to \( W^\pm \) events. This will increase the rise in the relative cross section that was seen before. When this effect is large enough, cutoff application could become advantageous.

For these simulations the same settings were used as for the \( m_{H^\pm} = m_{W^\pm} \) case. Simulations were performed for charged Higgs masses of 100, 120, 140, 150 and 160 GeV. For all of these masses 4 simulations were done with different values of \( \tan \beta \) (10, 20, 40 and 60) and again for the three different centre of mass energies. The amount of \( t\bar{t} \) events needed to claim a 5\( \sigma \) discovery before applying a cut for any of these combinations can be found in tables 2.2, 2.4 and 2.6. The integrated luminosity required for 5\( \sigma \) after applying the ideal cutoff can be found in tables 2.3, 2.5 and 2.7. As can be seen from these tables the combined effects of the mass boost and the polarisation make a lower cutoff favourable for masses above 150 GeV.

When lowering the centre of mass energy the most important change is the \( t\bar{t} \) production cross section. At 10 TeV this cross section is reduced by a factor of \( \approx 2 \) and at 7 TeV by a factor of \( \approx 5 \). Since this results in a decrease of both the signal and background the result on the signal to background ratio, which is defined as \( \sigma = \frac{\text{signal}}{\sqrt{\text{background}}} \), will be a reduced with \( \sqrt{2} \) and \( \sqrt{5} \) respectively. So even though a lower centre of mass energy is not favourable, a charged Higgs search remains possible in most of the parameter space. The cross section of the other simulated production process, \( bg \rightarrow H^+t \), get reduced more severely for lower collision energies. This is a result of the lower proton bottom quark density. This, however, only becomes relevant for the higher charged Higgs masses.
2.4. LIGHT CHARGED HIGGS

<table>
<thead>
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<th>m_{H^+} = 80</th>
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Table 2.2: Needed data (in fb$^{-1}$) without cutoff in charged pion energy spectrum for $H^\pm$ production at 14 TeV centre of mass energy.

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Table 2.3: Needed data (in fb$^{-1}$) with cutoff in charged pion energy spectrum for $H^\pm$ production at 14 TeV centre of mass energy.

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Table 2.4: Needed data (in fb$^{-1}$) without cutoff in charged pion energy spectrum for $H^+$ production at 10 TeV centre of mass energy.
### CHAPTER 2. CHARGED HIGGS

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Table 2.5: Needed data (in fb$^{-1}$) with cutoff in charged pion energy spectrum for $H^+$ production at 10 TeV centre of mass energy.

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Table 2.6: Needed data (in fb$^{-1}$) without cutoff in charged pion energy spectrum for $H^+$ production at 7 TeV centre of mass energy.

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</table>

Table 2.7: Needed data (in fb$^{-1}$) with cutoff in charged pion energy spectrum for $H^+$ production at 7 TeV centre of mass energy.
2.5 Heavy Charged Higgs

For $m_{H^±} \approx m_t$ and $m_{H^±} > m_t$ the dominant production process is $bg \rightarrow H^±t$. The background for this process again comes from top quark pair production, with subsequent decay into $W \rightarrow \tau$ by one of the top partners. The $H^±$ branching ratio to $\tau$ can no longer be considered 100% in this region. These branching ratios were calculated by PYTHIA itself and the used values can be found in figure 2.5. This figure shows that the BR to $\tau$ becomes small for higher charged higgs mass values. This is a result of the decay into a neutral higgs ($m_{h^0} = 115$) $W^±$ pair that starts to take over in this region. For higher values of $\tan \beta$ this effect gets less prominent, however, here the BR into $t\bar{b}$ becomes large and the decay in $\tau$ can be considered negligible for $m_{H^±} > 300$. Therefore, no simulations were performed for those values of $m_{H^±}$.

The same efficiency factors of 0.5 and 0.108 were applied in the data analysis accounting for the b-tagging efficiency and the needed $\mu$ from a $W^±$ decay.

![Figure 2.5: Branching ratios of the charged Higgs from PYTHIA for $\tan \beta = 10$ (upper-left), $\tan \beta = 20$ (upper-right), $\tan \beta = 40$ (lower-left) and $\tan \beta = 60$ (lower-right). With the three main decay channels, $\tau\nu_\tau$ in blue (slash-dot), $h^0W^±$ in green (dotted), $t\bar{b}$ in black (full).](image)

The cross sections of both the $H^±$ production processes become comparable around $m_{H^±} > 150\text{GeV}$. For masses above $m_t$ the $bg \rightarrow H^±t$ will completely take over. Since the cross section of the background stays the same more and more data will be needed for the $5\sigma$ discovery. The result is that only high $\tan \beta$ values are interesting especially
for $m_{H^\pm} > m_t$. Also the dependence on the collision energy gets more important in this region. For example for $m_{H^\pm} = 200$ and $\tan \beta = 10$ the cross section is reduced by a factor of $\approx 9$ when decreasing the centre of mass energy from 14 TeV to 7 TeV. The reduction of the background is a factor of $\approx 5$ when doing this. The result is that especially for the 7 TeV simulations only the parameter space with high $\tan \beta$ can be probed within the planned LHC run time.

As was already seen for the low mass charged Higgs case, the combined effect of the polarisation together with the momentum boost make a low energy cutoff in the pion energy favourable. As can be expected the ideal lower cutoff value becomes higher when increasing the charged Higgs mass. The results for all the high mass charged Higgs simulations without cutoff application can be found in tables 2.8, 2.10 and 2.12 and in 2.9, 2.11 and 2.13 after ideal cutoff was applied.
2.5. HEAVY CHARGED HIGGS

<table>
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Table 2.8: Needed data (in fb$^{-1}$) without cutoff in charged pion energy spectrum for $H^+$ production at 14 TeV centre of mass energy.

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<td>1.73</td>
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Table 2.9: Needed data (in fb$^{-1}$) with cutoff in charged pion energy spectrum for $H^+$ production at 14 TeV centre of mass energy.

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Table 2.10: Needed data (in fb$^{-1}$) without cutoff in charged pion energy spectrum for $H^+$ production at 10 TeV centre of mass energy.
CHAPTER 2. CHARGED HIGGS

\[ \tan \beta m_{H^+} = 180 \quad m_{H^+} = 200 \quad m_{H^+} = 250 \quad m_{H^+} = 300 \]

<table>
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<tr>
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Table 2.11: Needed data (in $fb^{-1}$) with cutoff in charged pion energy spectrum for $H^+$ production at 10 TeV centre of mass energy.

\[ \tan \beta m_{H^+} = 180 \quad m_{H^+} = 200 \quad m_{H^+} = 250 \quad m_{H^+} = 300 \]

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Table 2.12: Needed data (in $fb^{-1}$) without cutoff in charged pion energy spectrum for $H^+$ production at 7 TeV centre of mass energy.

\[ \tan \beta m_{H^+} = 180 \quad m_{H^+} = 200 \quad m_{H^+} = 250 \quad m_{H^+} = 300 \]

<table>
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Table 2.13: Needed data (in $fb^{-1}$) with cutoff in charged pion energy spectrum for $H^+$ production at 7 TeV centre of mass energy.

27
2.6 Conclusions

For the light charged Higgs scenario the use of the polarisation effect was shown not to be advantageous. The high BR of the top quark into $H^\pm$ in this mass region makes the use of the excess of $\tau$ events alone sufficient for a discovery $H^\pm$ within the LHC running time. The polarisation effect combined with the momentum boost start to make cutting favourable in the intermediate region where production by top quark decay and $bg \rightarrow H^\pm t$ become equally important. For masses above 150 GeV the needed integrated luminosity can be reduced by a few percent to almost 50% for $m_H^\pm = 300$. The possible amount of reduction for each parameter combination at 14 TeV can be seen in table 2.14. Going to higher masses would result in an even higher reduction, however, the production cross section becomes smaller and for $m_{H^+} > 300$ GeV the $H^+ \rightarrow \tau \nu_\tau$ BR becomes negligible. Therefore, other search methods are required in this part of parameter space.

Operating the LHC at energies of 7 or 10 TeV has, as expected, a negative effect. This is especially true for high $m_{H^\pm}$ because the production process involves a b-quark from a proton. Since lower beam energies result in a lower b-quark density the signal process gets reduced more than the background process. In the light charged Higgs scenario the signal to background ratio does not suffer badly from the beam energy reduction.

From the study presented in this chapter one can conclude that the polarisation of the resonance particle can be used to improve search methods. The MSSM is not the only model introducing new particles that differ in spin from a SM particle. The Higgsless models, which are the subject of the remainder of this thesis, suffer from the same problem. The results from this chapter indicate that a similar search strategy could therefore also be of interest for these models.

<table>
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<td>30.8</td>
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Table 2.14: Calculated reduction factor from cutoff in percent for 14 TeV, parameter combinations for which there is insufficient data were left blank.
Figure 2.6: Integrated luminosity after cuts required for a $5\sigma$ discovery. For $\tan \beta = 10$ (upper-left), $\tan \beta = 20$ (upper-right), $\tan \beta = 40$ (lower-left) and $\tan \beta = 60$ (lower-right) for the three different CMS energies, 14 TeV in black (full), 10 TeV in blue (slash-dotted), 7 TeV in green (dotted).
Chapter 3

Extra Dimensional Higgsless Theories

As shown in chapter 1, the Higgs mechanism used in the SM suffers from several problems. It is therefore interesting to look for alternatives. One such alternative including more Higgs bosons was discussed in the previous chapter. There are however also models in which no Higgs is needed at all. A group of such models are the extra dimensional Higgsless theories. They will be the subject of the remainder of this thesis. This chapter contains a theoretical introduction to these models. Two chapters will follow containing the results of a phenomenological study. These will include the results of a polarisation study similar to that performed for the MSSM charged Higgs.

3.1 Extra dimensional theories

Before explaining the idea behind the extra dimensional Higgsless models a description of extra dimensional models is required. Extra dimensional models do not all result in Higgsless models. Using the idea that their properties can be used to break EWS without a Higgs was later applied to a subgroup of them. The original motive for adding extra dimensions to the SM was to overcome the hierarchy problem. Supersymmetric theories do this by making the low energy regime of EW, or the low Higgs mass, look natural. Extra dimensional models lower the energy scale at which gravity becomes important, or make the Planck scale we measure look natural.

The gravitational force is calculated using:

\[ g(r) = -\frac{Gm_e r}{r^2} \] (3.1)

which is Newton’s law of gravitation. It can be rewritten using the Planck Mass as:

\[ g(r) = -\frac{m_e}{r^2} \frac{1}{M_{Pl}^2} \] (3.2)

If more dimensions are added to the model equation (3.2) has to be modified to:
CHAPTER 3. EXTRA DIMENSIONAL HIGGSLESS THEORIES

\[ g(r) = -m \frac{e_r}{M_{Pl}^{2+\delta} r^{2+\delta}} + 3 + 1 + \delta \]

This is true when the extra dimensions have an infinite size, which is obviously not correct. However, when assuming that the extra dimensions have a finite volume, say \( n \), the gravitational force is described by:

\[ g(r) = -m \frac{e_r}{M_{Pl}^{2+\delta} r^{2+\delta} V_n} \]

When \( r >> n \) this should behave as Newton’s law, so we get:

\[ -m \frac{e_r}{M_{Pl}^2 r^2} = -m \frac{e_r}{M_{Pl}^{2+\delta} r^{2+\delta} V_n} \]

which gives \( M_{Pl}^2 = M_{Pl}^{2+\delta} V_n \).

When requiring that all the SM particles do not probe the extra dimensions, its effects will not be noticeable for energies below the inverse size of these dimensions. If gravitons, the force carriers of gravity, are able to probe the extra dimensions this could account for the weakness of gravity. When the gravitons move into the higher dimensions, they cause the flux in the 4 dimensional world to become lower, causing the force to appear weaker. The lower flux of gravitons raises the apparent value of the Planck Scale to what we measure, \( M_{PL} = 2 \times 10^{18} GeV \).

The reason these extra dimensions have not shown up in experiments is that all the SM particles are restricted to the 4 dimensional subspace. The SM particles ”live” on the 3+1 dimensional sides of the 4+n dimensional space, called 3-branes.

The idea of using extra dimensional models to break EWS comes from warped extra dimensions. This is a subgroup of the extra dimensional models in which there is usually just one extra dimension, so there is a total of 5 space-time dimensions. In this extra dimension the space is warped to prevent the need for fine tunings. The warped space means that there is an exponential term in the metric:

\[ ds^2 = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2 \]

Here x are the coordinates of the normal dimension and y is the coordinate of the extra dimension. The two sides or 3 branes reside at \( y = 0 \) and \( y = \pi R \) where R is the size of the interval on which the extra dimension is compactified. For more detail on the mathematical frame work the reader is referred to for example [16]. This metric can be created in such a way that excitations of the SM gauge bosons are expected to exist. Simply speaking the SM particles can be seen as a standing wave within the extra dimension, this wave can then be excited into higher states or Kaluza-Klein (KK) excitations of the particles. These KK excitations will appear as heavier versions of the SM particles. Again the mathematics of this exceeds the topic of this thesis but can be found in for example [17], [18], [19]. By choosing correct boundary conditions for these
3.2 Higgsless Vector Boson Scattering

Waves a model can be created in which the gauge bosons acquire mass without the need for a Higgs particle, so EWSB is accomplished. A second consequence is that the SM gauge bosons, the W and Z are only the lowest of a whole KK tower of excitations. Their first excitations, the Z’ and the W’, may be detectable at the LHC.

### 3.2 Higgsless Vector Boson Scattering

The reason why these W’ and Z’ particles are of interest in a Higgsless model is that a SM without a Higgs boson suffers from two problems. First of all EWS needs to be broken to create masses for the gauge bosons. In the Higgsless models these masses come from a correct choice of the boundary conditions for these particles on the 3 branes, this can be found in for example [20] or [21]. Since the description of this is purely mathematical the interested reader is referred to these citations for further reading.

The second problem is the divergence of the gauge boson scattering amplitudes. In the Higgsless models, as the name suggests, there is no Higgs present and one would expect problems since the diagrams containing a Higgs are now lost. The way this problem is solved in extra dimensional Higgsless models is with the KK excitations of the vector bosons. These particles add new diagrams to the earlier discussed W_L W_L → W_L W_L process. These new diagrams are simply the SM diagrams containing a Z, now replaced by a Z’. They add the following terms to the total amplitude:

\[
A_{sZ'} = -\frac{1}{16} i g_{WWZ'} \frac{s^3}{s - \frac{M_Z^2}{M_W^2}} \beta^2 (3 - \beta^2)^2 \cos \theta, \\
A_{tZ'} = -\frac{i g_{WWZ'}^2 s^3}{32 (t - \frac{M_Z^2}{M_W^2})} [\beta^2 (4 - 2 \beta^2 + \beta^4) + \beta^2 (4 - 10 \beta^2 + \beta^4) \cos \theta \\
+ (2 - 11 \beta^2 + 10 \beta^4) \cos^2 \theta + \beta^2 \cos^3 \theta]
\]

One can now imagine the contributions from further excitations. Adding just the two amplitudes coming from the s and t channel mediated by Z’ the following term arises in the total amplitude [22]:

\[
A^{(4)} = i (g_{WWW}^2 - g_{WW\gamma}^2 - g_{WWZ}^2 - g_{WWZ'}^2) [(3 + 6 \cos \theta - \cos^2 \theta) + 2(3 - \cos^2 \theta)]
\]

This is the part of the amplitude that grows with \(s^2\) or \(E^4\). For the process to remain unitary through the entire energy spectrum this has to be 0. Therefore, \(g_{WWW}^2 - g_{WW\gamma}^2 - g_{WWZ}^2 - g_{WWZ'}^2\) must be equal to 0, or \(g_{WWW} = g_{WW\gamma} + g_{WWZ} + g_{WWZ'}\). So, the quadruple coupling of W to itself has to be equal to the coupling of a Z and photon to two W’s plus that of the Z’ to the two W’s. There is however also a term growing with \(s\) or \(E^2\):
CHAPTER 3. EXTRA DIMENSIONAL HIGGSLESS THEORIES

\[ A^{(2)} = \frac{1}{M_W^2} (4g_{WWWW}^2 M_W^2 - 3(g_{WWZ}^2 M_Z^2 + g_{WWZ'}^2 M_{Z'}^2)) \]

\[ - \frac{1}{2M_W^2} [4g_{WWWW}^2 M_W^2 - 3(g_{WWZ}^2 M_Z^2 + g_{WWZ'}^2 M_{Z'}^2) + \]

\[ (12g_{WWWW}^2 M_W^2 + g_{WWZ}^2 (M_Z^2 - 16M_W^2) + g_{WWZ'}^2 (M_{Z'}^2 - 16M_W^2))] \]

This term also has to be 0 if the process is to be well behaved for all energies, using the result obtained from the \(A^{(4)}\) term this expression becomes:

\[ A^{(2)} = \frac{1}{M_W^2} (4g_{WWWW}^2 M_W^2 - 3(g_{WWZ}^2 M_Z^2 + g_{WWZ'}^2 M_{Z'}^2))(1 - \sin^2 \theta) \]

Requiring this to be 0 results in a second relation between the masses and couplings of the different excitations:

\[ 4g_{WWWW}^2 M_W^2 = 3(g_{WWZ}^2 M_Z^2 + g_{WWZ'}^2 M_{Z'}^2) \]

As long as the mass and couplings of the first excitation follow from these requirements longitudinal WW scattering stays well behaved. The obvious new problem is the scattering of the excitations of the W, the W' bosons. This process runs into exactly the same problem as the normal WW scattering process, albeit at a higher energy. For the W'W' process all that needs to be added to solve this is the second excitation of the Z, the Z'. This can then be extended to a Z''' for W''' scattering and so on. So to keep everything unitary eventually the whole KK towers of both the gauge bosons have to be added to the models, resulting in an infinite number of excitations. The two new requirements for the whole towers can now be derived in a similar way as for WW scattering:

\[ g_{nnnn}^2 = \sum_k g_{nnk}^2 \]  \hspace{1cm} (3.3)

\[ 4g_{nnnn}^2 M_n^2 = 3 \sum_k g_{nnk}^2 M_k^2 \]  \hspace{1cm} (3.4)

Where n can be the W boson and k the Z.

3.3 W' and Z'

For the remainder of this thesis only the first excitations will be added to the SM. This approximation of the model is called the saturation limit [23]. That this is an appropriate approximation for physics below the TeV scale and physics at the LHC can be shown using the rules that ensure vector boson scattering unitarity. They can be summarised as (3.3) and (3.4), where n can for example be a W boson and k is the Kaluza-Klein tower of Z excitations.
3.3. W' AND Z'

When only using the Z and the Z', equation (3.4) can be rewritten into (3.5) for the WWZ coupling:

\[ 4g_{WWW}^2 M_W^2 = 3 \sum_k g_{WVk}^2 M_k^2 \]

\[ = 4g_{WWW}^2 M_W^2 = 3g_{WWZ}^2 M_Z^2 + 3g_{WWZ'}^2 M_{Z'}^2 \]

\[ = 4g_{WWZ}^2 \cos \Omega_W M_W^2 = 3g_{WWZ}^2 M_Z^2 + 3g_{WWZ}^2 M_{Z'}^2 \]

\[ = 4g_{WWZ}^2 M_Z^2 = 3g_{WWW}^2 M_W^2 + 3g_{WWZ}^2 M_{Z'}^2 \]

\[ \frac{4}{3}g_{WWZ}^2 M_Z^2 - g_{WWZ}^2 M_Z^2 = g_{WWZ}^2 M_{Z'}^2 \]

\[ \frac{1}{3}g_{WWZ}^2 M_Z^2 = g_{WWZ}^2 M_{Z'}^2 \]

So we finally get,

\[ \frac{1}{3} M_Z^2 = \frac{g_{WWZ}}{g_{WWZ'}} \] (3.5)

From the experimental constraint on \( \frac{g_{WWZ}}{g_{WWZ'}} \), which bounds this value from above to about 0.03 [22], it follows that \( m_{Z'} \) must be at least several hundred GeV. As can be derived from equations (3.3) and (3.4). A second excitation will be even heavier and will, in realistic models, probably not be lighter than 1 TeV. Using only the first excitation with an appropriate mass results in a model that stays unitary up to \( \approx 10 \) TeV [22]. Therefore, such a model will be sufficient for a phenomenological study of Higgsless models at the LHC.

Using equations (3.3) and (3.4) a full spectrum of the couplings and masses for the KK excitations of the gauge bosons can be derived. However, these equations say nothing about the coupling of the new excitations to the SM fermions. Experimental data from LEP does give limits to these couplings. As already explained in the Higgs section of this thesis indirect measurements of particles were done in these experiments. For the Higgs this was an indirect measurement of its mass, for the Z' and the W' it can be measurements of its couplings to fermions. Its values can be obtained from the Precision Electroweak Constraints (PEC).

These constraints are experimentally obtained values with which every Beyond the Standard Model (BSM) theory can be tested. When present, the effects of BSM particles on SM processes should result in parameter changes in the current experimentally probed region. The important parameters for new physics are the Peskin Takeuchi parameters called S,T and U. They are constructed in such a way that they are all equal to zero at a reference point in the Standard Model, with a particular value chosen for the Higgs boson mass. Their values are derived from the oblique corrections, which are the loop diagrams present in the propagators of four fermion interactions. The
parameters depend on the vacuum polarisation functions of the 3 gauge bosons, $\Pi_{GG}$, and that of the mixing between $\gamma$ and $Z$, $\Pi_{Z\gamma}$. They are defined as:

$$\alpha_S = 4s_w^2c_w^2 \left[ \Pi'_{ZZ}(0) - \frac{c_w^2 - s_w^2}{s_w c_w} \Pi'_{Z\gamma}(0) - \Pi'_{\gamma\gamma}(0) \right]$$

$$\alpha_T = \frac{\Pi_{WW}(0)}{M_W^2} - \frac{\Pi_{ZZ}(0)}{M_Z^2}$$

$$\alpha_U = 4s_w^2 \left[ \Pi'_{WW}(0) - c_w^2 \Pi'_{ZZ}(0) - 2s_w c_w \Pi'_{Z\gamma}(0) - s_w^2 \Pi'_{\gamma\gamma}(0) \right]$$

Without going in to their precise meaning one can for example see that $T$ is a parameter for isospin violation. $T$ depends on the loop diagrams in the $W$ and $Z$ boson propagator. Where the $Z$ couples to either 2 up-type or 2 down-type fermions, the $W$ couples to a pair consisting of one up and one down-type. If one now calculates the isospin violation of a certain model one can check the consistency of this model with the experimental value for $T$.

These parameters can be used to put experimental constraints on every BSM theory. For example new physics models with an extra group of vector bosons will change the values of the Peskin Takeuchi parameters since new vector bosons will cause new loop diagrams in the gauge boson propagators (if they do couple to the SM gauge bosons). The masses and couplings of these new vector bosons therefore need values that do not alter the values of $S, T$ and $U$ outside of their current experimentally verified areas. Since most electroweak parameters, which can be calculated from the $S, T$ and $U$ values, are known through experiments to be close to the SM predicted values the effect of new physics on these parameters must be small.

Whereas the EW parameters like $g_{VWVW}$ can be varied by a few percent from the SM predicted value, the couplings of the new gauge bosons to fermions are much restricted by the PEC. With the help of the parameters a $Z'$ and $W'$ with strong couplings to fermions can be be excluded. We therefore have to assume that the new gauge bosons have a weak or non-existing coupling to the SM fermions. That this is theoretically possible by using fermion delocalisation was shown by for example [22]. The interested reader can find the principle behind this in literature such as [22],[24] and [25]. In this thesis we will set the fermion couplings to 0, making the new gauge bosons fermiophobic. There are two important consequences of this. The first is that the number of diagrams including a $Z'$ or $W'$ is restricted. So unlike the Higgs boson only several channels stay available for both production and decay. Secondly, this lack of decay channels in combination with the weak couplings to the decay products causes the lifetimes to be relatively long. This results in a narrow Breit-Wigner resonance. The width of the resonances is a few GeV [23], whereas that of an equally heavy Higgs boson is in the order of 10 to a 100 GeV. A possible consequence of this is that the higher order excitations can be dark matter candidates. However, to my knowledge this has never been studied in detail.

The extra dimensional Higgsless models can schematically be presented as in figure 3.1.
As mentioned earlier, without couplings to fermions the number of diagrams with a Z’ or W’ is restricted. For the Z’, which will be the main focus if the next two chapters, the two remaining production processes at the LHC are Vector Boson Fusion (VBF) and Z’-strahlung. They can be seen in figure 3.2 and will both be studied in this thesis. The VBF process has been studied in detail before for heavy SM Higgs bosons. A similar search method will be developed for the Z’ with an extra emphasis on its spin 1 nature. The Z’-strahlung process will have 3 W’s in its final state of which one is highly energetic, making it interesting for phenomenological studies. These characteristics in combination with the narrow Breit-Wigner and heavy mass make the Z’ a promising candidate for a search at the LHC.
Figure 3.2: The two production mechanisms of a fully fermiophobic $Z'$. The diagram on the left depicts $Z'$-strahlung, the right diagram shows the VBF process.
Chapter 4

Z’-strahlung

Z’-strahlung is the process where a quark and an antiquark fuse into a $W^\pm$ that later radiates off a $Z'$. This chapter contains a description of the Monte Carlo implementation of this process and a discussion of its results. The process will be shown to have a relatively high cross section and several distinctive phenomenological signatures, making both its discovery at the LHC and a $Z'/Higgs$ discrimination probable. Signatures resulting from its spin 1 nature will especially be interesting for future studies. Therefore, a strategy for correct polarisation handling in the Monte Carlo simulations is also presented in this chapter.

4.1 Monte Carlo Implementation

For an implementation of this process in Monte Carlo simulation software an expression for the differential cross section or amplitude squared is required. Such an expression can be derived from that of the $WZ$ production process calculated in [27]. The expression for $WZ$ production contains three different processes, namely that of Z-strahlung and two processes with quark propagators. The Feynman diagrams of which can be seen in figure 4.1. Since the $Z'$ is fermiophobic the contributions from the two processes with quark propagators need to be removed for $WZ'$ production. This reduces the expression from [27] to equation (4.1).
\[ \frac{d\sigma}{dt} = \frac{\pi \alpha^2}{6s^2 \sin^4 \theta_W} \times \left( \frac{1}{s - M_W} \right)^2 \left[ \left( \frac{9 - 8 \sin^2 \theta_W}{4} \right) (ut - M_W^2 M_Z^2) \right. \]
\[
\left. + \left( 8 \sin^2 \theta_W - 6 \right) s(M_W^2 + M_Z^2) \right] \]

(4.1)

Since the coupling of the Z' to a W pair differs from the ZWW coupling by a factor \( g_{WZZ'} / g_{WWZ} \), the expression in (4.1) needs to be multiplied by this factor squared. The then obtained expression was implemented in PYTHIA 8.108. However, plotting the differential cross section as a function of the \( p_T \) of the W boson one gets figure 4.2.

![Figure 4.2: Transverse momentum of the W after radiating the Z' with \( m_{Z'} = 500 \text{GeV} \). An unphysical cutoff at \( p_T = 400 \text{GeV} \) is present in the distribution.](image)

This indicates a 0 or possibly negative cross section for values of the W boson energy below 400 GeV. This is unphysical and therefore the implementation must be incorrect. The expression presented in [27] is a rewritten version of that presented in [26]. The two were compared and as proven in Appendix A an error was introduced by the authors of [27] during rewriting. Therefore, the expression from [26] was used as a starting point to obtain the following expression:

\[ \frac{d\sigma}{dt} = \frac{4\pi \alpha^2}{s^2} \left( \frac{G_{V-A}^{ij}}{e} \right)^2 \left( \frac{se_{Z}/e}{s - M_W^2} \right)^2 A(s, t, u) \]  

(4.2)

With:

\[ A(s, t, u) = \left( \frac{ut}{M_W^2 M_Z^2} - 1 \right) \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2 + 8M_W^2 M_Z^2}{4s^2} \right] \]

\[ + \left( \frac{M_W^2 + M_Z^2}{M_W^2 M_Z^2} \right) \left[ s/2 - M_W^2 - M_Z^2 + (M_W^2 - M_Z^2)^2 / 2s \right] \]
4.1. MONTE CARLO IMPLEMENTATION

\[ G_{V-A}^{ij} = \frac{M_W \sqrt{G_F}}{2^4} V_{ud} \]

\[ e_Z = e \cot \theta_W \]

Using equation (4.2), multiplied by \( \frac{g_{WWZ'}}{g_{WWZ}} \), in the implementation results in the distribution shown in figure 4.3. The sharp cutoff at 400 GeV is now absent. Several more tests were performed to see if this new implementation is indeed correct. These test were performed with the help of the CompHEP 4.5.1 [12] software package. CompHEP can be used to calculate distributions like that shown in 4.3 for several models. The Higgsless model, however, is not present in the package. A simplified version can be produced by adding only the \( Z' \) particle along with its coupling to WW to the SM. This is a copy of the ZWW coupling with an added constant. Setting \( m_{Z'} = 25 \text{GeV} \) several parameters were plotted using both the corrected PYTHIA implementation and CompHEP. The results of these tests can be found in figures 4.4 and 4.5. The low \( Z' \) mass was chosen since it makes the characteristics of the distributions more prominent. It must also be noted that CompHEP does not allow the propagating resonance to be stable whereas PYTHIA does allow this. This causes a resonance peak according to a Breit Wigner distribution to be present at \( \sqrt{s} = m_W \) in the PYTHIA results. For the CompHEP results this peak is missing. Figures 4.4 and 4.5 do show that the first PYTHIA implementation gives unphysical results. An example of this is the rise of the cross section at \( \sqrt{s} = 170 \text{GeV} \). This should be at \( m_H + m_{Z'} - \Gamma_W - \Gamma_{Z'} \approx 100 \). The \( p_T \) plot also contains an unexplainable resonance peak. Figure 4.6 shows the ratio of the results from the corrected PYTHIA implementation and those from CompHEP, since the ratio is stable this implementation was chosen for further simulations.

Since the expression from [27], used for the WZ production process in PYTHIA 8.108, was proven to be incorrect, a new version of the WZ production process was written for future incorporation in PYTHIA. The cross section expression from [26] was used for this. A full explanation of this can be found in Appendix B.

\[ ^1 \text{Only qualitative comparisons could be made since the CompHEP package calculates using a specific quark input, for example } ud, \text{ whereas PYTHIA averages over all quark combinations.} \]
Figure 4.3: $p_T$ of the W after radiating the $Z'$ of 500 GeV. A sharp cutoff is now absent.
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Figure 4.4: The differential cross section as a function of $p_T$ for a $Z'$ of 25 GeV. The upper right figure shows the results from CompHEP, the upper left for the new PYTHIA implementation. The result from the original PYTHIA implementation is shown in the lower figure. Only a qualitative comparison can be made.
Figure 4.5: The differential cross section as a function of $\sqrt{s}$ for a $Z'$ of 25 GeV from CompHEP (upper left), new PYTHIA implementation (upper right) and original implementation (lower middle). Only a qualitative comparison can be made.
Figure 4.6: Ratio of the $\sqrt{s}$ distribution from the corrected PYTHIA implementation over those from CompHEP. The ratio is a constant, showing the results to be qualitatively equal.
4.2 Cross Sections

Using PYTHIA the cross section of the Z'-strahlung process was calculated for different Z' masses. A comparison of these cross sections with those of the Higgs strahlung process can be found in figure 4.7. The coupling of the Z' to the WW pair used here was calculated using equation (3.5). The assumption was used that for $m_{Z'} = 350\text{GeV}$ the value of $g_{WWZ'}/g_{WWZ} = 0.02$. This is well within the experimentally allowed limits of about 0.03 [22]. With this starting assumption the rest of the values of $g_{WWZ'}/g_{WWZ}$ were calculated. Furthermore the cross sections of the processes were calculated for the LHC with a centre of mass energy of 14 TeV using the default parton distribution function from PYTHIA 8.108, CTEQ5.1. The results for Z'-strahlung are in agreement with the W' production cross sections presented in [23], those for Higgs strahlung are in agreement with [28].

![Cross sections of both Higgs strahlung and Z'-strahlung as a function of $M_{\text{resonance}}$.](image)

Figure 4.7: Cross sections of both Higgs strahlung and Z'-strahlung as a function of $M_{\text{resonance}}$.

Figure 4.7 shows that the two cross sections are not of the same order size. The cross section of the Z' is almost a factor of 100 larger than that of Higgs strahlung and therefore not much smaller than the SM process of WZ creation. The reason for this relatively large cross section is the disappearance of the interfering quark diagrams. For the WZ creation process the different diagrams interfere negatively with one another resulting in a lower cross section. The cross section of Higgs strahlung makes this process of no interest an LHC search. The main reason for this is that there are other production processes with a much higher cross sections. On top of this the Higgs width, which increases rapidly with increasing mass, makes it more difficult to spot. For the Z' the width stays in the GeV range even when its mass is 700 GeV [23]. This results
4.3. BACKGROUND SIMULATION

in a narrow resonance peak for the the $Z'$, making a $5\sigma$ discovery easier.

Since the LHC will not start at the planned $\sqrt{s}$ of 14 TeV the cross sections were also calculated for the earlier mentioned 7 and 10 TeV centre of mass energies. The results of these calculations can be seen in figure 4.8.

![Figure 4.8: Cross sections $Z'$-strahlung as a function of $M_{Z'}$ for the three different planned centre of mass energies of the LHC.](image)

Figure 4.8 shows that the cross section for this process decreases with a factor of $\approx 3$ for 10 TeV and a factor of $\approx 8$ for 7 TeV. This drop in cross section is partly the result of the lower gluon density in the protons at these energies. The gluon density is directly connected to the antiquark densities and therefore influences the process' cross section.

4.3 Background simulation

Background simulations were performed using both CompHEP 4.5.1 and PYTHIA 8.108. PYTHIA was used to simulate the $Z'$-strahlung process using the implementation discussed before. CompHEP is capable of creating events of all the possible SM processes with three vector boson final states. This background was chosen since $W$ and $Z$ bosons can hardly be distinguished from one another when decaying hadronically. For a leptonic decay of these particles a distinction is possible, however, only $\approx \frac{1}{1000}$ of these events decays fully leptonically. A combination of leptonic and hadronic decay will therefore be of interest for LHC searches. Processes involving the Higgs were also incorporated in the background simulations. The Higgs was chosen to have a mass equal to that of the $Z'$ to simulate the most difficult scenario.
CHAPTER 4. Z’-STRAHLUNG

All processes that can be mistaken for a 3 vector boson signal were not taken into account here. One of the most important of these is $t\bar{t}$ production, which results in 2 jets from W-decay and 2 jets from b-decay. Resulting in a similar signature to Z’-strahlung when one b-quark is misidentified as a W jet and the other as QCD radiation. The production cross section of this process is at least a factor $10^3$ higher than that of Z’ production making it a potential problem. For now these kind of background processes were left out since a dedicated detector simulation is needed for a signal and background discrimination study. This will be the subject of a follow up study, [29].

The production of a simulation with $M_{Z'} = 500$ GeV can be seen in figure 4.9. The invariant mass was calculated using the 2 decay products of the Z’ or in the case of a background event the 2 vector bosons with the highest energy. Different methods have been tested for the calculation of the background invariant mass, the differences between these methods were however small. They all resulted in a clearly visible resonance peak.

![Figure 4.9: The Z’ signal for $M_{Z'} = 500$ GeV with the SM background for $\approx 100 fb^{-1}$. A clear resonance from the Z’ can be seen as well as a wide resonance from the SM Higgs boson of 500 GeV.](image)

The results shown in 4.9 agree with those presented for W’-strahlung in [30]. The high resonance peak in the figure is partly a result of its long life time. This small resonance width makes it, as discussed before, easier to find than an equally massed Higgs, the effects of which can hardly be seen in this plot. However, as mentioned before, processes resembling the final state of 3 gauge bosons like $t\bar{t}$ production have a much larger cross section making a Z’ discovery more difficult than discussed here. A dedicated detector study is needed to investigate this. Also the final 3 W state will make event reconstruction difficult. This also requires continued research using detector analysis software.
4.4 Polarisation effects

As shown in the previous chapter polarisation measurements can be a helpful tool to distinguish between different resonance particles. For Z’-strahlung discrimination from a heavy SM Higgs boson or neutral MSSM Higgs boson is required. This can be achieved by studying the differences in kinematics of the decay products. This section contains an explanation of the expected differences, a proposed method for a full polarisation study and the first results of such a study.

4.4.1 Z’ vs. Higgs: Vector vs. Scalar

Since we are interested in the fermiophobic Z’ we will only look at its decay into WW. When a SM Higgs decays into a W pair both W’s will have opposite spin or both will be longitudinally polarised, indicated in figure 4.10. Due to parity violation, which forces all neutrino’s to be left-handed, the decay will, in the situation with W’s with opposite spin, result in two left-handed neutrino’s and 2 right-handed leptons. For two W’s with spin parallel to their direction of motion the final state leptons will all be left-handed, see figure 4.10. The first situation occurs in 2/3 of the decays and the lower one in the remaining 1/3 of all the Higgs to WW decays.

\[ \begin{array}{c}
\text{Ratio:} \\
2/3 \\
1/3
\end{array} \]

Figure 4.10: Illustration of a Higgs decay into a W pair with subsequent leptonic decay. The red arrows illustrate the direction of spin.

Z’ decay can occur in 3 different ways. Firstly there is the decay as depicted in the upper part of figure 4.11. The spin 1 nature of the Z’ forces one of the W’s to be polarised like the Z’, while the other is longitudinally polarised. The W with equal polarisation to the Z’ will result in both the lepton and the neutrino being left-handed. The other W will decay into a left-handed neutrino and a right-handed lepton. The situation depicted in the middle is similar to the decay depicted in the upper part of figure 4.10. The last situation will result in 2 left-handed leptons. The first situation occurs in 2/3 of all the decays since a Z’ with spin ±1 will result in the depicted decay channel. The second possibility occurs for a Z’ of spin 0 when one of the W’s has spin
+1 or -1. So an extra factor of $2/3$ has to be added, resulting in a fraction of $2/9$ for this decay channel. The last channel then has the remaining fraction of $1/9$. This leads to the result that for a $Z'$ decay $5/9$ of all outgoing leptons is right-handed and $4/9$ is left-handed.

![Figure 4.11: Illustration of a $Z'$ decay into a $W$ pair with subsequent leptonic decay. The red arrows illustrate the direction of spin.](image)

When adding up all the possibilities for the final handedness of the outgoing measurable leptons a difference is found between $Z'$ and Higgs decay. However, this is only true when taking randomly polarised resonance particles. The polarisation of the resonance is, however, dependent on the production process since the cross sections differ per polarisation configuration. The above mentioned ratios will therefore not be measured in a detector. More complicated, process dependent polarisation distributions are needed to calculate these ratios correctly. To accomplish this an interface between PYTHIA and another program called CAMORRA was imagined.

### 4.4.2 PYTHIA-CAMORRA Interface

Within PYTHIA polarisations are accounted for in the following way. The hard process is first calculated using the spin averaged cross section. After the particles are given momentum in this first step the decay of the resonances is handled. Here a $Z'$ can for example decay into $WW$, however, their momenta and angles depend on the $Z'$ polarisation. Therefore, these momenta and angles are calculated together with a weight that accounts for the spin of the resonance. This weight can be seen as account-
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Figure 4.12: Schematic representation of the intended PYTHIA-CAMORRA interface.

The weights for the probability of the resonance having a certain spin. The way these weights are calculated is however non trivial, they are derived from complex distributions, the calculations of which require an in depth study. A solution to this problem for the charged Higgs search was the use of the TAUOLA package. No similar package exists for the decays of the $Z'$, therefore, a different solution had to be found.

A solution was developed that involves the CAMORRA program currently under development at the University of Nijmegen [31]. The CAMORRA program is able to calculate the cross sections of the production processes, including the spin correlations, using spin dependent matrix elements. However, PYTHIA is needed to create the
CHAPTER 4. Z'-STRAHLUNG

phase space for CAMORRA. An interface between the two programs was therefore planned that uses PYTHIA to produce the incoming partons. CAMORRA then handles the events and calculates the polarisation dependent weight. The same event is also processed by PYTHIA including the decay and hadronisation. The results of PYTHIA are then multiplied by the weights from CAMORRA. This way the CAMORRA weights are used to reshape the isotropical distributions produced by PYTHIA. A schematic representation of this can be seen in 4.12.

This interface can be used to calculate the whole process, including the effects of polarisation, together with full decay and hadronisation. Although this method would work for any process it could not yet be used for the work presented here since it is still under development. Only a restricted study can therefore be presented in this thesis. However, the work is ongoing and more results are assumed to be presented in [29].

4.4.3 Characteristics of the Radiating W

Parameters that can be studied without the CAMORRA-PYTHIA interface are those of the radiating W and the Z’ itself. The momentum and production angle of the radiating W are calculated in the part of the code that handles the hard process.

The simulations for such a study were performed for three different resonance masses, 350 GeV, 500 GeV and 700 GeV. Next to the self implemented Z’-strahlung process the Higgs strahlung process was used here to compare the characteristics of the two. Different resonance masses are interesting since higher masses are expected to wash out the polarisation effects. Apart from the different resonance masses all simulations were once performed for high precision, to see the ideal results, and once for the planned yearly amount of integrated luminosity of 10 fb$^{-1}$. This was done to see if boson discrimination can be performed at the LHC.

The $p_T$ distribution of the W from Z’-strahlung differs greatly from that of the Higgs strahlung process as can be seen in 4.13, 4.14 and 4.15. In the Z’ process the W has a much higher $p_T$ than in the case of Higgs strahlung. This is a result of the higher average $\sqrt{s}$. The difference is large enough for it to still be clearly visible even for a yearly amount of LHC data, as can be seen in 4.19 and 4.21. For 700 GeV it becomes more difficult, this is due to both the mass boost effects and the lower cross section. Within the LHC run time of about 10 years, however, here too the polarisation effect will become clear. The differences in the pseudorapidity of the resulting W are less striking but also clearly visible. The results of this can be seen in figures 4.16, 4.17 and 4.18 for high precision simulations and figures 4.22, 4.23 and 4.24 for the integrated luminosity of 10 fb$^{-1}$.

As expected, the higher mass reduces the differences in distributions. For example the ratio of the peaks in the pseudorapidity plots is reduced from 1.6 to 1.2 when going from 350 GeV to 700 GeV. A larger problem for these higher masses is the cross section reduction. However, the $p_T$ signature in Z’-strahlung does make this a promising channel for boson discrimination. In this $p_T$ distributions the differences between Higgs and Z’ become smaller with increasing resonance mass and the cross section does go down with it. However, the average value of the distributions shifts to higher values for
increasing mass making it easier to distinguish the events from the background. This can compensate some of the loss in production cross section.

Figure 4.13: $p_T$ distribution of the 3rd $W$ for $Z'$-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 350\text{GeV}$.

Figure 4.14: $p_T$ distribution of the 3rd $W$ for $Z'$-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 500\text{GeV}$.
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Figure 4.15: $p_T$ distribution of the 3rd $W$ for Z'-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 700\text{GeV}$.

Figure 4.16: $\eta$ distribution of the 3rd $W$ for Z'-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 350\text{GeV}$. 
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Figure 4.17: $\eta$ distribution of the 3rd $W$ for $Z'$-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 500\text{GeV}$.

Figure 4.18: $\eta$ distribution of the 3rd $W$ for $Z'$-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 700\text{GeV}$.
Figure 4.19: $p_T$ distribution of the 3rd W for Z$'$-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 350\text{GeV}$. For $10 \text{fb}^{-1}$.

Figure 4.20: $p_T$ distribution of the 3rd W for Z$'$-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 500\text{GeV}$. For $10 \text{fb}^{-1}$.
4.4. POLARISATION EFFECTS

Figure 4.21: $p_T$ distribution of the 3rd W for Z'-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 700$GeV. For $10 fb^{-1}$.

Figure 4.22: $\eta$ distribution of the 3rd W for Z'-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 350$GeV. For $10 fb^{-1}$.
Figure 4.23: η distribution of the 3rd W for Z'-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 500$ GeV. For $10 \text{fb}^{-1}$.

Figure 4.24: η distribution of the 3rd W for Z'-strahlung (left), Higgs strahlung (middle) and the normalised ratio of the two (right) for $m_{Z'} = m_H = 700$ GeV. For $10 \text{fb}^{-1}$. 
4.5 Conclusions

The cross section for Z'-strahlung is relatively large compared to Higgs strahlung. However it does drop faster than that of Higgs strahlung with increasing resonance mass. This is a result of the maximum allowed value of $\frac{g_{WWZ'}}{g_{WWZ}}$, which decreases with increasing Z' mass. The choice of $\frac{g_{WWZ'}}{g_{WWZ}}$ of 0.02 for $m_{Z'} = 350$ GeV was however close to the maximum allowed value and there is no real lower limit on this value. The cross sections could therefore be much smaller. The small width of the Z', which is partly due to this low value of $\frac{g_{WWZ'}}{g_{WWZ}}$, causes a clear resonance to be present in the background simulations. Also the distinctive high $p_T$ distribution of both the Z' and the radiating W make this channel interesting. Not only could these help to distinguish it from a similar process with a scalar boson, like the SM Higgs, it could also help in background reduction. Both a trigger on a high $p_T$ W jet and a trigger on the high $p_T$ decay products of the Z' would greatly reduce the background. Unfortunately no study of these decay products could be performed yet. One can however assume that the high $p_T$ of the Z' will cause the two W's from its decay to be close together. Spin effects could then either increase or decrease this effect on the W decay products. Future studies can show this easily using the presented idea of the PYTHIA-CAMORRA interface. With the help of this interface both the detectable decay products of the 3rd W and the Z' can be studied for the application of kinematical cuts. Another interesting subject of future research is a detector study in which a more detailed signal and background study can be performed. Processes like $t\bar{t}$ production could for example proof to be a problematic background for this process. The importance of this process as a background depends on the particle recognition efficiency of the detector software. These detector studies are also important for testing which decay channels of the W's are of interest for LHC searches. These studies are currently planned for [29].
Chapter 5

Vector Boson fusion

The second production process of the fully fermiophobic $Z'$ is Vector Boson Fusion (VBF). In this process the $Z'$ is created by the fusion of two W-bosons that are radiated off two incoming quarks or antiquarks. The Feynman diagram is presented in figure 5.1. Since the fermiophobic $Z'$ only couples to W’s the final state of this process will consist of two W’s and the two jets resulting from the outgoing quarks. A similar process exists for Higgs production. This has been studied in detail in the past, for example by [32]. This chapter contains a description of the process implementation in PYTHIA, the results of a signal to background simulation and a phenomenological study with an emphasis on the polarisation effects.

Figure 5.1: Feynman diagram of the $Z'$ VBF production mechanism.

5.1 Monte Carlo Implementation

Since the VBF process is of little interest for non-fermiophobic $Z'$ models it is not present in PYTHIA 8.108. Therefore, just like for the $Z'$-strahlung process, it had to be implemented. CompHEP was chosen to produce the required expression of the amplitude squared, the main reason for this choice is that no explicit cross section calculations of this process have been performed up to date. Also unlike the $Z'$-strahlung
5.1. MONTE CARLO IMPLEMENTATION

process no equivalent can be found amongst the processes already included within PYTHIA.

The process implementation in PYTHIA is based on the VBF process for Higgs production \(^1\) since this already contains the complicated basis of a \(2 \rightarrow 3\) process. Also this process can serve as a test for the imagined method of process implementation. For this test CompHEP was used to calculate the amplitude squared. This expression was than used to replace that already present in PYTHIA. Both implementations gave equal results, proving that the method can be used.

To change the Higgs production process into a \(Z'\) production process two things need to be changed. The first is the code which handles the decay, however, as already explained for the \(Z'\)-strahlung process this is not achievable within the work of this thesis. Therefore, this code was simply set to always give a weight of 1, resulting in isotropic distributions of the decay parameters. In the future this can be fixed using the CAMORRA program. The second change required is the expression of the amplitude squared. In CompHEP the process was calculated with the model used before for the \(Z'\)-strahlung tests, the SM with an additional \(Z'\) particle. The only couplings of the \(Z'\) are those to \(W\)'s, which are copies of the \(Z\) boson couplings with an extra factor of \(g_{WWZ'}/g_{WWZ}\). The calculated CompHEP expression was implemented in the PYTHIA code. Again several tests were performed to compare the results from PYTHIA and those acquired with separate CompHEP simulations. The results can be seen in figures 5.2, 5.3 and 5.4. These results show to be qualitatively similar. Again the ratio between the two is a constant, like that shown for \(Z'\)-strahlung.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{distribution}
\caption{Distribution of the invariant mass of \(q'\) and \(q''\) for the VBF process with \(m_{Z'} = 500\text{GeV}\) for CompHEP (left) and PYTHIA (right)}
\end{figure}

\(^1\)Coming from HiggsSM:ff2Hff (t:WW)
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Figure 5.3: Distribution of $p_T$ of the $Z'$ for the VBF process with $m_{Z'} = 500\text{GeV}$ for CompHEP (left) and PYTHIA (right)

Figure 5.4: Distribution of $\sqrt{s}$ for the VBF process with $m_{Z'} = 500\text{GeV}$ for CompHEP (left) and PYTHIA (right)

5.2 Cross Sections

The cross sections for several masses were calculated using PYTHIA. These were then compared with those of Higgs production through VBF and $Z'$ production through strahlung. The same values of $g_{WWZ'}$ as for the $Z'$-strahlung study were used here. The results of the cross section calculations are presented in figure 5.5.

As can be seen from figure 5.5 the cross sections for production through VBF are of the same order size as those of $Z'$-strahlung. The VBF process only takes over as the dominant production channel for $m_{Z'} = 800\text{GeV}$. These results are similar to those presented in [23] where both cross sections are calculated for the processes of fermiophobic $W'$ creation. In [23] the cross sections are also of the same order size, however, $W'$ production through VBF takes over around $m_{W'} = 400\text{GeV}$. The cause of this, apart from the difference in couplings, lies in the different initial quark combinations that can lead to the processes. For example, $Z'$ production through VBF
5.3 Background

The main backgrounds that need to be considered for this process are the same as those studied for Higgs production through VBF [32]. These are all the states with 2 W’s and 2 jets in the final state, but also those with 2 Z’s and 2 jets and the combination of a W a Z and 2 jets. The combined cross sections of these processes is however small compared to the process of $t\bar{t}$ production, which has a cross section of the order
CHAPTER 5. VECTOR BOSON FUSION

Figure 5.6: Cross section (fb) of VBF $Z'$ production, $Z'$-strahlung and Higgs production through VBF for the three centre of mass energies.

$6 \times 10^5 fb$ at the LHC. Its final state of 2 W’s and 2 b-jets makes it looks similar to the signal of $Z'$ production. Unlike for strahlung a simple signal to $t \bar{t}$ background study can be performed here since the final states are now equivalent. Other backgrounds with direct WW production need to be studied in more detail in combination with detector studies. The reason for this are the numerous processes that resemble this final state in a detector. Jet reconstruction, W tagging efficiency and initial state radiation are important features of this and require a dedicated studies.

Using only the $t \bar{t}$ process should give a useful insight in the discovery possibilities through this process. Both the VBF and the top pair production processes were simulated in PYTHIA. The result of a simulation with $m_{Z'} = 500 GeV$ at 14 TeV can be seen in figure 5.7.

Figure 5.7 shows that the high cross section of $t \bar{t}$ production makes the $Z'$ resonance hard to spot. Several features can however be exploited to reduce the background. The first of these is the absence of bottom quarks in the signal process. There will of course be several VBF events with a b-quark in the final state but the number of these will be negligible. The ATLAS detector has a b-tagging efficiency of approximately 50%, which means it is able to identify half of all the b-jets. Since the presence of one b-jet is enough to exclude an events as signal, the background process considered here can be reduced by 75%.

A second characteristic that can be used to differentiate between signal and back-
5.3. BACKGROUND

Figure 5.7: Invariant mass distribution of the W pair for the processes of \( t\bar{t} \) production and \( Z' \) production (red) through VBF with \( m_{Z'} = 500 \text{GeV} \) for \( \approx 3 \text{fb}^{-1} \). The peak resulting from the \( Z' \) has been enlarged in the middle.

Background is jet direction. The signal process will have two jets in opposite directions. This is a result of the high momentum of the partons in the beam direction. In the interaction the quark will acquire \( p_T \) which is small compared to the momentum along the beam pipe, resulting in a forward jet. For the b-jets of the background process this is not the case. To illustrate this the difference in pseudorapidity of the two jets for both signal and background is shown in figure 5.8.

Figure 5.8: \( |\Delta \eta| \) distribution of the two jets for the processes of \( t\bar{t} \) (right) production and \( Z' \) production (left) through VBF with \( m_{Z'} = 500 \text{GeV} \).
Figure 5.8 shows that the difference between the two processes can be used to greatly enhance the signal to background ratio. This was tested by removing all events where $|\Delta \eta|$ was smaller than 4. Plotting the invariant mass of the W pair again after application of this simple cut gives the result shown in figure 5.9. This figure shows a dramatic increase in the signal to background ratio when compared to figure 5.7.

More cuts are able to further increase the signal to background ratio. These cuts can for example be on the $p_T$ or direction of the leptons coming from W decay. For the application of these cuts a correct handling of the W decay, including polarisation, is needed. A study of these cuts was therefore not possible here. However, it can be imagined that a higher $p_T$ is expected for the signal leptons since these come from the heavy $Z'$. Using these extra cuts and the one studied in this subsection the problem of $t\bar{t}$ production as a background can be greatly reduced. Again a follow up study is needed to confirm this.

### 5.4 Polarisation effects

Just like for the $Z'$-strahlung process only a basic study of the polarisation effects can be performed here. Whereas a study of the hard process particles is possible, their decay products can not be studied. The earlier proposed CAMORRA-PYTHIA interface is needed for these purposes. The particles that can be studied for this process are the $Z'$ itself and the two final state quarks. Their properties will be discussed in the first part of this section. Since this process becomes important for the higher $Z'$ masses,
the effects of these high masses on a prominent feature of the decay products was also studied. This study was not performed for the VBF process but rather for the fermion annihilation production process, so a fermiophile Z' is considered there. The reason for this choice was that this process is already implemented in PYTHIA including a correct weight calculation. Although this process is absent in the fermophobic Higgsless model, this study is useful for gaining insight in the effect of the high resonance masses.

5.4.1 Z' vs. Higgs: Vector vs. Scalar

Similar to the strahlungs processes a comparison will be made between the process of Higgs and Z' production through VBF. The parameters that can be studied here are those of the particles of the hard interaction, the Z' and the jets. Surprisingly the differences between the processes in the $\eta$ distributions of the quarks are quite large, as can be seen in figures 5.10, 5.12 and 5.14. Also a similar difference as for the strahlungs process can be found in the $p_T$ spectrum. The on average higher $p_T$ of the Z', when compared to the Higgs, should lead to an easier discovery. As expected, the differences again become smaller with increasing resonance mass. Since the cross sections are similar to those of Z'-strahlung the amount of data acquired within the LHC runtime will in theory be enough to distinguish the two processes. Detector capabilities might however make this more difficult.
Figure 5.10: $|\Delta\eta|$ distribution of the two hadronic jets for $Z'$ production (left), Higgs production (middle) and the normalised ratio, $Z'/H$ (right). For $m_{Z'} = m_H = 350\text{GeV}$

Figure 5.11: $p_T$ distribution of the for $Z'$ (left), Higgs (middle) and the normalised ratio, $Z'/H$ (right). For $m_{Z'} = m_H = 350\text{GeV}$
5.4. POLARISATION EFFECTS

Figure 5.12: $|\Delta \eta|$ distribution of the two hadronic jets for $Z'$ production (left), Higgs production (middle) and the normalised ratio, $Z'/H$ (right). For $m_{Z'} = m_H = 500\text{GeV}$

Figure 5.13: $p_T$ distribution of the for $Z'$ (left), Higgs (middle) and the normalised ratio, $Z'/H$ (right). For $m_{Z'} = m_H = 500\text{GeV}$
CHAPTER 5. VECTOR BOSON FUSION

Figure 5.14: $|\Delta \eta|$ distribution of the two hadronic jets for $Z'$ production (left), Higgs production (middle) and the normalised ratio, $Z'/H$ (right). For $m_{Z'} = m_H = 700$GeV.

Figure 5.15: $p_T$ distribution of the for $Z'$ (left), Higgs (middle) and the normalised ratio, $Z'/H$ (right). For $m_{Z'} = m_H = 700$GeV.
5.4. POLARISATION EFFECTS

5.4.2 Polarisation searches through $\Delta \Phi_{ll}$

One of the most striking signatures of Higgs production through VBF is the equal
direction of the charged leptons. The reason for this feature is that this process occurs
predominantly through the exchange of longitudinally polarised vector bosons. The
same holds for its subsequent decay. So, as shown in figure 5.16, both of the leptons
from W decay have to travel in the same direction. This makes the difference in $\phi$
of the two oppositely charged leptons a distinctive signal for this process. This was shown
in [32] for a Higgs mass of 170 GeV. The distribution, resulting from the scalar nature
of the Higgs, was used there to reduce the signal to background ratio. This can be seen
from figure 5.17, taken from [32].

![Figure 5.16: Schematic representation of Higgs boson decay through vector bosons and
their polarisations [32].](image)

![Figure 5.17: $\Delta \Phi$ distribution of the two leptons for the Higgs signal on the left and that
of its main background channels on the right [32].](image)

For a $Z'$ decay one would expect a distribution in $\Delta \Phi$ different from that of the
Higgs. This difference can be used to discriminate between the two signals. However,
the study performed in [32] was done for a Higgs with a mass of 170 GeV, much lower
than a realistic mass of the $Z'$. A resonance with a higher mass will cause a higher
momentum boosts of the W’s, resulting in different distributions. For high enough
boosts this will result in the two leptons travelling in opposite directions despite the
polarisation effects. To test this, simulations were performed in PYTHIA for Higgs
production through VBF with the subsequent decay chain of $H \rightarrow WW \rightarrow \mu^+\mu^-\nu\bar{\nu}$. These simulations were performed for several Higgs boson masses starting with 170 GeV
to see if the results match those from [32]. Similar simulations were then performed for
$m_H = 200\text{GeV}$, $m_H = 250\text{GeV}$, $m_H = 350\text{GeV}$, $m_H = 500\text{GeV}$, $m_H = 700\text{GeV}$. The
results of which can be seen in figure 5.18.
Figure 5.18: $\Delta \Phi$ distributions for several values of $m_H$: 170 GeV left upper corner, 200 GeV right upper corner, 250 GeV left middle, 350 GeV right middle, 500 GeV left bottom and 700 GeV right bottom.

Figure 5.18 shows that for 170 GeV a distribution similar to that of [32] is found. Also, as expected, the distribution changes to one similar to the background when increasing the mass of the Higgs boson. For higher mass values the effect of the polarisation is lost. To see how this influences the possibilities of $Z'$ Higgs discrimination a second set of simulations was performed. For this both the processes $f \bar{f} \rightarrow H$ and $f \bar{f} \rightarrow Z'$ were simulated first with $m_H = m_{Z'} = 170$ GeV. This is not a realistic mass for a $Z'$ boson, however, for this mass the momentum boost effect should still be
small. For a realistic \( Z' \) mass, \( m_H = m_{Z'} = 600 \text{GeV} \) the polarisation effect will be much smaller, this situation was therefore also simulated. Note that the \( Z' \) is not taken to be fermiophobic here because this production mechanism is not possible for a fully fermiophobic \( Z' \). The reason why this production mechanism was chosen is that these processes are both already implemented in PYTHIA, they already include a correct handling of the polarisations.

The result of the first simulations with \( m = 170 \text{GeV} \) can be seen in figure 5.19. This figure shows that whereas the polarisation of the Higgs forces the two leptons in the same direction, the distribution of the \( Z' \) looks more similar to the background processes from figure 5.17. This is the result of its vector boson character, which allows for several different final state polarisation combinations. This distribution could therefore be used to discriminate between the two resonances in the fermiophile scenario.

![Figure 5.19: \( \Delta \Phi_H \) distributions from \( f \bar{f} \rightarrow Z' \) on the left, \( f \bar{f} \rightarrow H \) in the middle and the ratio of the two on the right, for \( m_H = m_{Z'} = 170 \text{GeV} \).](image)

The distributions for \( m_H = m_{Z'} = 600 \text{GeV} \) can be found in figure 5.20. Here the Higgs signal looks much more like that of the \( Z' \), as can be seen from the ratio. The ratio between the two processes drops for higher \( m_{\text{resonance}} \), thus requiring more data to discriminate between the two. The drop in cross sections with increasing mass adds to this problem. It can therefore be concluded that whereas this feature can easily be used for low mass resonance discrimination, for high masses it will lose its use. Since the here studied production channel of the Higgs and the VBF process show similar results in the directions of the leptons the same can be expected for the \( Z' \). Therefore, the result of this study can be extended for the VBF process. However, for confirmation this has to be studied using the PYTHIA-CAMORRA interface in the future.
Figure 5.20: $\Delta \Phi_{ll}$ distributions from $f \bar{f} \to Z'$ on the left, $f \bar{f} \to H$ in the middle and the ratio of the two on the right, for $m_H = m_{Z'} = 600$ GeV.

5.5 Conclusions

For most allowed masses the cross section for $Z'$ production through VBF is smaller, but of the same order as the strahlung process. It does however decrease less with increasing $m_{Z'}$, making the process interesting for higher resonance masses. The signal also offers several features that make it distinguishable from the background, the most promising being the two forward jets. However, as mentioned before a dedicated detector study is needed to study the background and the ways in which it can be reduced. The cut on jet separation presented in this thesis is a first step towards this, it already shows to be of great use for $t \bar{t}$ background reduction.

Using this process for vector scalar discrimination is not as easy as for the strahlungs case. The first reason for this is that the $Z'$-strahlungs process has the prominent high $p_T$ W, which distinguishes it from Higgs strahlung. The VBF process shows differences between the two resonances, however these are not as large as that of the strahlungs process. A second reason is that the VBF process becomes more interesting for higher resonance masses. There the effects coming from polarisation get washed out by the momentum boosts. However, the lack of information on the polarisations of the W pairs means that many features could not be studied in this thesis. An indirect study of the lepton direction did show that the differences suffer greatly with increasing mass.

So although the process has a lower cross section, VBF does seem promising for future studies. Especially for lower LHC centre of mass energies this process becomes at least of equal interest since it suffers less from such a decrease in energy. Last of all VBF is the only of the two processes that can occur at a future electron positron collider. Making it the only channel suitable for precision measurements of the $Z'$. 
Chapter 6

Conclusions

The Higgs mechanism of the SM is both simple and elegant, however, it suffers from several problems and has not yet been confirmed experimentally. Therefore, it is interesting to look for alternative ways to break electroweak symmetry. The extra-dimensional Higgsless models do not only offer such an alternative but also solve the hierarchy problem and could even introduce dark matter candidates. Although theorists are still working on the details of the models, a phenomenological study for an LHC search can already be performed. Some simple studies were already performed in the past however these mostly focused on the W' or on specific models in which the KK to fermion couplings were well defined. Therefore, it is interesting to study the Z' in a general way by setting these fermion couplings to 0.

Since the Z’ differs from heavy Higgs bosons in spin, the effects of this are an interesting research topic. A study of the better understood MSSM model showed that effects resulting from polarisation can be measured. For MSSM, however, the presented polarisation effect were over shadowed by the larger effect of $\tau$ BR differences. This makes polarisation effects of little use for a low mass charged Higgs search. In the high mass region, where the effects are increased due to mass boosts, utilising them becomes advantageous. However, with increasing mass the cross sections go down and only for high values of $\tan \beta$ a discovery remains possible.

Using Monte Carlo software the process of Z'-strahlung was shown to have a relatively large cross section when compared to for example Higgs strahlung. In combination with the narrow resonance width this could lead to a possible discovery at the LHC. The main problems will however lie in the reconstruction because the final state contains 3 W’s. The event reconstruction along with a study of the background of this signal will have to be performed in the future to show how promising this production channel really is. Some signatures of this process already show this channel to be interesting, especially for scalar vector discrimination. The high momentum of both the Z’ and the radiating W are the most striking features that could be studied here. This high momentum along with the large cross section make this process an interesting candidate for continued research.

The second production process of the Z’ is through VBF. In the lower Z’ mass
region this has a lower cross section than \( Z' \)-strahlung. However, its final state, which contains only 2 \( W \)'s and two distinctive forward jets, make it interesting. The forward jets can help with a large background reduction, this was shown for the \( t\bar{t} \) background. The angular separation of the two forward jets for this process also differs from that of the same process with a Higgs. This makes resonance discrimination possible, however, the differences are not as large as for example the high \( W \) \( p_T \) of the strahlungs process. More differences are expected from the kinematics of the decay products, these need to be studied in the future. One striking feature of the decay products was already studied indirectly and showed to be of great value for the low mass region, for higher realistic masses however, the effect becomes negligible. Continued studies of this process using the PYTHIA-CAMORRA interface are therefore strongly recommended. Especially since this process becomes the dominant production process for high resonance masses.

In this mass region the exploitation of the spin effects becomes more important due to low cross sections resulting in low statistics. The last important feature of this process is that its cross section suffers less than that of the strahlungs process when decreasing the CMS energy, making it more interesting for these lower collision energies.

In total it can be concluded that this appealing theory has two \( Z' \) production processes that are both interesting for future studies. Further investigation is needed in the field of detector simulations and polarisation effects for which a method is proposed in this thesis. The studied processes do have a relatively high cross section, which combined with their phenomenological distinctive signatures make a discovery within the LHC runtime probable.
Appendix

Appendix A

The differential cross section derived in [27] from [26] is:

\[
\frac{d\sigma}{dt} = \frac{\pi \alpha^2}{6s^2 \sin^4 \theta_W} \times \left( \frac{1}{s - M_W} \right)^2 \left[ \left( \frac{9 - 8 \sin^2 \theta_W}{4} \right) \left( ut - M_W^2 M_Z^2 \right) + \frac{(8 \sin^2 \theta_W - 6)}{4} s(M_W^2 + M_Z^2) \right]
\]

The original from [26] is:

\[
\frac{d\sigma}{dt} = \frac{4\pi \alpha^2}{s^2} \left( \frac{G_{ij}^A V}{e} \right)^2 \left( \frac{se_Z/e}{s - M_W^2} \right)^2 A(s, t, u)
\]

With:

\[
A(s, t, u) = \left( \frac{ut}{M_W^2 M_Z^2} - 1 \right) \left[ \frac{1}{4} - M_W^2 + M_Z^2 \right] + \left( \frac{M_W^2 + M_Z^2}{2s} \right) \left[ s/2 - M_W^2 - M_Z^2 + (M_W^2 - M_Z^2)^2/2s \right]
\]

\[
G_{ij}^A = \frac{M_W \sqrt{G_F}}{2} V_{ud}
\]

\[
e_Z = e \cot \theta_W
\]

It can be shown that the two cross section terms however differ by writing the derived expression as:

\[
\frac{d\sigma}{dt} = \frac{\pi \alpha^2}{6s^2 \sin^4 \theta_W} \left( \frac{1}{s - M_W} \right)^2 B(s, t, u)
\]

And the original as:

\[
\frac{d\sigma}{dt} = \frac{4\pi \alpha^2}{s^2} \left( \frac{G_{ij}^A V}{e} \right)^2 \left( \frac{se_Z/e}{s - M_W^2} \right)^2 A(s, t, u)
\]

This gives:

\[
\frac{4\pi \alpha^2}{s^2} \left( \frac{G_{ij}^A V}{e} \right)^2 \left( \frac{se_Z/e}{s - M_W^2} \right)^2 A(s, t, u) = \frac{\pi \alpha^2}{6s^2 \sin^4 \theta_W} \left( \frac{1}{s - M_W} \right)^2 B(s, t, u)
\]

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It is now clear that both sides have the common terms \( A(s, t, u) = \frac{1}{6 \sin^4 \theta_W} \left( \frac{1}{s - M_W} \right)^2 B(s, t, u) \)

\[
4 \left( \frac{G^i_j}{e} \right)^2 \left( \frac{e z/e}{s - M_W^2} \right)^2 A(s, t, u) = \frac{1}{6 \sin^4 \theta_W} \left( \frac{1}{s - M_W} \right)^2 B(s, t, u)
\]

\[
24 \left( \frac{1}{e^2 \sqrt{2} \sin \theta_W} \right)^2 \left( \frac{e z/e}{s - M_W^2} \right)^2 A(s, t, u) = \frac{1}{6 \sin^4 \theta_W} \left( \frac{1}{s - M_W} \right)^2 B(s, t, u)
\]

\[
24s^2 \left( \frac{1}{8 \sin^2 \theta_W} \right)^2 \left( \frac{e \cot \theta_W/e}{s - M_W^2} \right)^2 A(s, t, u) = \frac{1}{6 \sin^4 \theta_W} \left( \frac{1}{s - M_W} \right)^2 B(s, t, u)
\]

Filling in \( A(s, t, u) \) and \( B(s, t, u) \) we get:

\[
3s^2 \frac{M_W^2}{M_Z^2} \left( \left( \frac{ut}{M_W^2 M_Z^2} - 1 \right) \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2 + 8M_W^2 M_Z^2}{4s^2} \right] 
+ \frac{M_W^2 + M_Z^2}{2s} \left[ s/2 - M_W^2 - M_Z^2 + (M_W^2 - M_Z^2)^2 / 2s \right] \right)
= \frac{9 - 8 \sin^2 \theta_W}{4} \left( ut - M_W^2 M_Z^2 \right) + \frac{(8 \sin^2 \theta_W - 6)}{4} s (M_W^2 + M_Z^2)
\]

\[
3s^2 \frac{1}{M_Z^2} \left( \left( ut - M_W^2 M_Z^2 \right) \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2 + 8M_W^2 M_Z^2}{4s^2} \right] 
+ (M_W^2 + M_Z^2) \left[ s/2 - M_W^2 - M_Z^2 + (M_W^2 - M_Z^2)^2 / 2s \right] \right)
= \left( \frac{9 - 8 \sin^2 \theta_W}{4} ( ut - M_W^2 M_Z^2 ) + \frac{(8 \sin^2 \theta_W - 6)}{4} s (M_W^2 + M_Z^2) \right)
\]

It is now clear that both sides have the common terms \((ut - M_W^2 M_Z^2) X_1\) and \((M_W^2 + M_Z^2) X_2\) and that we can divide these expressions in two parts. The first part is:

\[
\frac{9 - 8 \sin^2 \theta_W}{4} ( ut - M_W^2 M_Z^2 ) M_W^2 M_Z^2
\]

\[
= 3s^2 \frac{M_W^2}{M_Z^2} \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2 + 8M_W^2 M_Z^2}{4s^2} \right]
\]

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\[
\frac{9 - 8 \sin^2 \theta_W}{4} M_W^4 M_Z^2 = 3 s^2 \frac{M_W^2}{M_Z^2} \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2 + 8 M_W^2 M_Z^2}{4s^2} \right]
\]

\[
\left( \frac{1}{4} + 2 \frac{M_W^2}{M_Z^2} \right) M_W^2 M_Z^2 = 3 s^2 \frac{M_W^2}{M_Z^2} \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2 + 8 M_W^2 M_Z^2}{4s^2} \right]
\]

\[
\left( \frac{1}{4} + 2 \frac{M_W^2}{M_Z^2} \right) M_Z^4 = 3 \left[ \frac{1}{4} s^2 - \frac{M_W^2 + M_Z^2}{2} + \frac{(M_W^2 + M_Z^2)^2 + 8 M_W^2 M_Z^2}{4} - \frac{1}{12} M_Z^2 - \frac{2}{3} M_W^2 M_Z^2 \right] = 0
\]

Resulting in:
\[
s^2 - 2s(M_W^2 + M_Z^2) + (M_W^2 + M_Z^2)^2 + 5 \frac{1}{3} M_W^2 M_Z^2 - \frac{1}{3} M_Z^2 = 0
\]

The second part is:
\[
3 s^2 \frac{1}{M_Z^2} (M_W^2 + M_Z^2) \left[ s/2 - M_W^2 - M_Z^2 + (M_W^2 - M_Z^2)^2/2s \right] = \frac{(8 \sin^2 \theta_W - 6)}{4} s(M_W^2 + M_Z^2)
\]

\[
3 \frac{1}{M_Z^4} \left[ s^2/2 - (M_W^2 + M_Z^2)s + (M_W^2 - M_Z^2)^2/2 \right] = \frac{(8 \sin^2 \theta_W - 6)}{4}
\]

\[
12 \left[ s^2/2 - (M_W^2 + M_Z^2)s + (M_W^2 - M_Z^2)^2/2 \right] = M_Z^2 (2 - \frac{8 M_W^2}{3 M_Z^2})
\]

\[
s^2 - 2(M_W^2 + M_Z^2)s + (M_W^2 - M_Z^2)^2 = M_Z^2 (1/3 - \frac{4 M_W^2}{3 M_Z^2})
\]

\[
s^2 - 2(M_W^2 + M_Z^2)s + (M_W^2 - M_Z^2)^2 + \frac{4}{3} M_W^2 M_Z^2 - \frac{1}{3} M_Z^2 = 0
\]

\[
s^2 - 2(M_W^2 + M_Z^2)s + M_W^4 - \frac{2}{3} M_W^2 M_Z^2 + \frac{2}{3} M_Z^2 = 0
\]

The two expressions differ indicating that there is an error somewhere in the derivation. Apart from this error an assumption must also have been made to end up with an almost constant expression for s. The expression chosen for s, can be derived from:
\[
s^2 - 2s(M_W^2 + M_Z^2) + (M_W^2 + M_Z^2)^2 + 5 \frac{1}{3} M_W^2 M_Z^2 - \frac{1}{3} M_Z^2 = 0
\]
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or from:

\[ s^2 - 2(M_W^0 + M_Z^0)s + M_W^4 - \frac{2}{3}M_W^2M_Z^2 + \frac{2}{3}M_Z^4 = 0 \]

The first of these however results in \( s \) having an imaginary part. The second expression gives a different result:

\[ s = 2(M_W^0 + M_Z^0) \pm \sqrt{4(M_W^4 + M_Z^4)^2 - 4(M_W^4 - \frac{2}{3}M_W^2M_Z^2 + \frac{2}{3}M_Z^4)} \]

\[ s = 2(M_W^0 + M_Z^0) \pm \sqrt{4(M_W^4 + M_Z^4 + 2M_W^2M_Z^2) - 4(M_W^4 - \frac{2}{3}M_W^2M_Z^2 + \frac{2}{3}M_Z^4)} \]

\[ s = 2(M_W^0 + M_Z^0) \pm \sqrt{\frac{32}{3}M_W^2M_Z^2 + \frac{4}{3}M_Z^4} \]

\[ s = 2(M_W^0 + M_Z^0) \pm \sqrt{(1 + 8\cos^2 \theta_W)\frac{4}{3}M_Z^4} \]

indicating that an error has been made in the first expression. From the above we can conclude that when going from [26] to [27] the assumption was made that \( s \) was on a resonance in the cross section, which is of course not correct for this process. The value of \( s \) and \( 2(M_W^0 + M_Z^0) \pm \sqrt{(1 + 8\cos^2 \theta_W)\frac{4}{3}M_Z^4} \) where both printed for this process (with the PYTHIA implementation). This shows that the two values are equal and that this assumption was made when deriving this expression.

After some calculation it can be seen that when using the same expression for \( s \) in the other part of the equation from [27] one gets:

\[ 3s^2 \frac{M_W^2}{M_Z^2} (ut - M_W^2M_Z^2) \left[ \frac{1}{4} - \frac{M_W^4 + M_Z^4}{2s} + \frac{(M_W^2 + M_Z^2)^2}{4s^2} \right] \]

So it seems that the factor of \( 8M_W^2M_Z^2 \) in the last part was left out. This is not a justifiable simplification and it therefore seems that it was forgotten.
Appendix B

To implement the differential cross section from the paper of Brown et all. [26] the function that needs to be changed is the `sigmaKin()` function. The new cross section from [26] is:

\[
\frac{d\sigma}{dt} = \frac{4\pi\alpha^2}{s^2} \left( \frac{G_{V-A}^{ij}}{e} \right)^2 \left[ \frac{(se_Z/e)^2}{s - M_W^2} A(s, t, u) + \frac{2se_Z/e}{s - M_W^2} \left( -\frac{g^i}{e} I(s, t, u) + \frac{g^j}{e} I(s, u, t) \right) + 
\frac{g^i - g^j}{e} E(s, t, u) + \frac{g^i g^j}{e^2} \left( \frac{ut - M_W^2 M_Z^2}{u^2} \right) + \frac{g^i g^j}{e^2} \left( \frac{ut - M_W^2 M_Z^2}{t^2} \right) + 2 \frac{g^i g^j}{e^2} s(M_W^2 + M_Z^2)/ut \right]
\]

With:

\[
G_{V-A}^{ij} = \frac{M_W \sqrt{G_F}}{2^4} \delta_{ju} \delta_{id}
\]

\[
e_Z = e \cot \theta_W
\]

\[
A(s, t, u) = \left( \frac{ut}{M_W^2 M_Z^2} - 1 \right) \left[ \frac{1}{4} - \frac{M_W^2 + M_Z^2}{2s} + \frac{(M_W^2 + M_Z^2)^2}{4s^2} \right] + \left( \frac{M_W^2 + M_Z^2}{M_W^2 M_Z^2} \right) [s/2 - M_W^2 - M_Z^2 + (M_W^2 - M_Z^2)^2/2s]
\]

\[
I(s, t, u) = \frac{1}{4} \left( \frac{ut}{M_W^2 M_Z^2} - 1 \right) \left[ 1 - \frac{M_W^2 + M_Z^2}{s} - \frac{(M_W^2 + M_Z^2)^2}{st} \right] + \left( \frac{M_W^2 + M_Z^2}{M_W^2 M_Z^2} \right) [s - M_W^2 - M_Z^2 + (2M_W^2 - M_Z^2)^2/t]
\]

\[
E(s, t, u) = \frac{1}{4} \left( \frac{ut}{M_W^2 M_Z^2} - 1 \right) + \frac{1}{4} \frac{s(M_W^2 + M_Z^2)}{M_W^2 M_Z^2}
\]

Using the definition of $G_F$ the common prefactor was rewritten as follows:

\[
\frac{4\pi\alpha^2}{s^2} \left( \frac{G_{V-A}^{ij}}{e} \right)^2 = \frac{4\pi\alpha^2}{s^2} \left( \frac{M_W^2 G_F}{\sqrt{2} e^2} \right)
\]

\[
= \frac{4\pi\alpha^2}{s^2} \left( \frac{M_W^2}{e^2 \sqrt{2}} \right) \left( \frac{\sqrt{2} e^2}{8 \sin^2 \theta_W M_W^2} \right)
\]

\[
= \frac{4\pi\alpha^2}{s^2} \left( \frac{1}{8 \sin^2 \theta_W} \right)
\]
With this prefactor the implementation in PYTHIA looks as follows:

```c
double resBW = 1. / (pow2(sH - mWS) + mwWS);

double cotT = sqrt(cos2thetaW / sin2thetaW);

double prefac = 12.0 * M_PI * pow2(alpEM) / (sH2 * 8. * sin2thetaW);

double A = ((uH*tH)/(s3*s4) - 1.) * (0.25 - ((s3+s4)/(2.*sH)) + (pow2(s3+s4) + 8.*s3*s4/(4.*sH2)) + (s3+s4)/(s3*s4) * (sH/2. - s3 - s4 + pow2(s3-s4)/(2.*sH));

double B1 = lun * ( 0.25 * ((uH*tH)/(s3*s4) - 1.) * (1.- (s3+s4)/sH - (4.*s3*s4)/(sH * tH)) + ((s3+s4)/(2.*s3*s4)) * (sH - s3 - s4 + 2.*s3*s4/tH));

double B2 = lde * ( 0.25 * ((uH*tH)/(s3*s4) - 1.) * (1.- (s3+s4)/sH - (4.*s3*s4)/(sH * uH)) + ((s3+s4)/(2.*s3*s4)) * (sH - s3 - s4 + 2.*s3*s4/uH));

double E = 0.25 * ((uH*tH)/(s3*s4) -1.) + 0.5 * sH*(s3+s4)/(s3*s4);

sigma0 = prefac* ( pow2(cotT) * sh2 * resBW * A + 2.* sh* cotT * resBW * (sH - mWS) * (B2-B1) + pow2(lun-lde) * E + pow2(lde) * (uH*tH-s3*s4)/uH2 + pow2(lun) * (uH*tH-s3*s4)/tH2 + 2.*lun*lde * sH *(s3+s4)/(uH*tH)) ;
```

Notice that an extra factor of 3 was added to the prefactor to correct for the factor 1/3 in the CKM element calculation in \texttt{sigmaHat()}. Apart from this the \texttt{lun} and \texttt{lde} variables were simply taken to be those from the \texttt{initProc} function.

The result of this new implementation is that the cross section does not drop below 0 like the old implementation. A second result is that when setting the coupling of the Z to fermions to 0, \texttt{lun=0, lde=0}, the function behaves as it should according to tests preformed with CompHEP. Whereas the old implementation gave results with an unnatural lower cutoff in the \(p_T\) of the W and Z. The last result is a change in the cross section of the process. For the LHC energy this cross section used to be \(2.580 \times 10^{-08} \text{mb}\), it is now increased to \(9.075 \times 10^{-08} \text{mb}\).
Appendix C

A problem occurs when setting all the couplings of the Z' to fermions to 0. The reason for this is that the photon, Z and Z' are treated as a mixture. During production of one of the three the mixture is calculated according to the individual couplings to up and down type quarks. This is required for the flavour of in- and outcoming quarks to be treated correctly. Problems in this mixing procedure do however arise when setting all the couplings of the Z' to 0. This problem was pointed out to the authors of PYTHIA and it will be fixed in one of the future distributions. For now a simple solution [34], which fixes this problem is by adding the following to the function `ResonanceDecays.cc`

```cpp
// Mother flavour - relevant for gamma*/Z0 mixing.
int idIn = process[decayer.mother1()].id();
// Debug.
if (id0 == 32) idIn = 0;
```
Bibliography


[31] G. van Oord, in preperation.


[34] Private communications with T. Sjostrand.
Acknowledgements

After about 16 months I finally managed to finish this thesis, of course there are several people who helped to make this possible. First of all I’d like to thank Bob van Eijk. Most importantly for introducing me to particle physics, but also for getting me into NIKHEF, CERN, The University of Geneva and DESY, and of course for providing me with all the help required for this thesis. I enjoyed the large amount of freedom I got in my work while always being able to walk into your office whenever there were problems of either a scientific or bureaucratic nature. Secondly I’d like to thank Ahmed Ali and the whole DESY theory department for their hospitality and providing me with ideas and help for the polarisation studies. For always quickly providing solutions whenever I encountered problems with PYTHIA I would like to thank Torbjorn Sjostrand. Both the very complicated and the extremely stupid questions were always answered within one day. For the work on the CAMORRA-PYTHIA interface I would like to thank both Gijs van Oord and Ronald Kleiss for their help and assistance.

I also need to thank Pieter van der Deyl who spend a lot of his time listening to me complaining, but still helped me out with several problems, of which several were relevant for my work. And for saying:” moet je is kijken wat er gebeurd als je de Z’ weer fermiofiel maakt”. Then of course Lars Beemster deserves a big thank you for several reasons. First of all for helping me out whenever there root/c++/latex/printer/unix related problems, but also for all the useful and useless discussions we had. But most importantly for providing the ”master kamer” with a refrigerator and the beer to fill it, oh and for making me feel less guilty about showing up late by always showing up even later.

For this report I also need to thank the very Australian Tessa Charles for proof reading this report, the (useful) discussions and for teaching me the correct pronunciation of the words ”Higgs” and ”data”, this thesis wouldn’t have been the same without it. I also need to thank the NS since I’ve spend about 400 hours (≈ 40.000 km) on their trains in the last couple of months.

Last of all I want to thank my parents who provided me with food and shelter in the many periods I’ve been homeless, I wouldn’t have survived without it.

If you have read this thesis and think you contributed somehow but have not been acknowledged, just fill in the following sentence: I would like to thank ............ for .................................................................