# Properties and Discovery Prospects of Higgs Bosons from the Minimal Supersymmetric Standard Model

Candidacy Paper

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**Abstract**

Supersymmetric theories provide the only known models of the fundamental interactions which have the potential to resolve the problems of naturalness and heirarchy while retaining the Higgs boson as a fundamental spin zero particle. The Minimal Supersymmetric Standard Model (MSSM) is a supersymmetric theory which introduces the minimum number of new parameters beyond the Standard Model. This theory requires two Higgs doublets to generate masses for all of the fermions and gauge bosons, and predicts a spectrum of three neutral Higgs boson and two charged Higgs bosons. The properties and coupling of these bosons to fermions and gauge bosons are summarized. Limits on the masses of some of these bosons have been obtained at the LEP $e^+e^-$ collider at CERN and are presented here. Higgs decay channels which may be observable at future hadron colliders such as the LHC are also discussed.

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## 1. Introduction

The Standard Model (SM)[1] has proven to be capable of precisely predicting a broad spectrum of particle physics phenomena, however it is not expected to stand up to scrutiny at the TeV energy scale. In particular, the mechanism by which the weak gauge bosons and the fermions acquire mass has not been tested experimentally. Extensions of the minimal SM, such as the Minimal Supersymmetric Standard Model (MSSM),[2] have been proposed in an attempt to expand the range of physics theory beyond the electroweak scale. These theories provide a more natural basis for constructing Grand Unified Theories, in which the strong force is united with the weak and electromagnetic (EM) forces, and typically predict extensive new phenomenology beyond the SM. This paper summarizes the properties of the MSSM Higgs sector, which is responsible for electroweak symmetry breaking and the generation of particle masses within this theory. It is hoped that future high energy hadron colliders such as the LHC[3] will provide a better understanding of symmetry breaking and particle masses.

The remainder of section 1 provides a summary of the SM and illustrates how spontaneous breaking of a local gauge symmetry and the Higgs mechanism can combine to produce gauge boson masses. The MSSM is introduced in section 2, and the properties and couplings of the Higgs bosons predicted by the MSSM are described in section 3. Current experimental limits on Higgs masses and the prospects for their discovery at future hadron colliders are discussed in section 4.

### *1.1. The Standard Model*

The SM is a highly successful theory describing the interactions between fundamental particles. In this theory, matter is composed of point-like spin 1/2 fermions which interact via the strong, weak and EM forces. These forces arise through the exchange of spin 1 particles called gauge bosons. Some of the properties of the SM fermions and bosons are given in table 1. Fermions are classified as either leptons or quarks based on their ability to interact via the strong force. The leptons consist of the electron ($e$), muon ($\mu$), tau ($\tau$) and their associated neutrinos. These particles possess integer electric charge and do not interact via the strong force. The six quarks (u, d, s, c, t, and b) each possess fractional charge and can interact by the strong force as well as the weak and EM forces. Each fermion is associated with an antiparticle of the opposite electric charge.

Gauge Bosons (spin = 1)

| Name | Charge | Mass (GeV) |
|---|---|---|
| photon ($\gamma$) | 0 | 0 |
| $Z^0$ | 0 | 91.16 |
| $W^+$ | +1 | 80.6 |
| $W^-$ | −1 | 80.6 |
| gluon ($g$) | 0 | 0 |

Higgs Boson (spin = 0)

| Name | Charge | Mass (GeV) |
|---|---|---|
| $H$ | 0 | ? |

Fermions (spin = 1/2)

| $T^3$ | Charge (Q) | Flavour | Mass (GeV) | Flavour | Mass (GeV) | Flavour | Mass (GeV) |
|---|---|---|---|---|---|---|---|
| +1/2<br>−1/2 | 0<br>−1 | $\begin{pmatrix} \nu_e \\ e \end{pmatrix}$ | $< 2 \times 10^{-8}$<br>$5.1 \times 10^{-4}$ | $\begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}$ | $< 3 \times 10^{-4}$<br>0.106 | $\begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}$ | $< 4 \times 10^{-2}$<br>1.784 |
| +1/2<br>−1/2 | +2/3<br>−1/3 | $\begin{pmatrix} u \\ d \end{pmatrix}$ | $\sim 4 \times 10^{-3}$<br>$\sim 7 \times 10^{-3}$ | $\begin{pmatrix} c \\ s \end{pmatrix}$ | $\sim 1.5$<br>$\sim 0.15$ | $\begin{pmatrix} t \\ b \end{pmatrix}$ | $\sim 174$<br>$\sim 4.7$ |

Table 1: Some properties of the fermions and bosons in the Standard Model.[4] The discovery of the top ($t$) quark has not yet been confirmed.[5]

The fermions can be further catagorized into three families, each consisting of a lepton and a quark weak isospin doublet. Each family therefore contains a pair of quarks with charges $+2/3$ and $-1/3$, a charged lepton and a neutrino. The weak force is able to couple members of each doublet to one another by charged current interactions (mediated by a $W^{\pm}$). For example, an electron can be converted to an electron neutrino by coupling to a $W^{\pm}$. Charged current interactions can also couple quarks from one doublet to quarks in another doublet, but with a lower probability than interactions within the same doublet. These probabilites are parameterized by the Cabbibo-Kobayashi-Maskawa (CKM) matrix.[6] Processes which couple two different fermions of the same electric charge (flavour changing neutral currents or FCNCs) are not permitted in the SM.

The strong, weak and EM forces are mediated by the exchange of particles known as gauge bosons, which arise due to the invariance of the SM Lagrangian under $SU(3)_C \times SU(2)_L \times U(1)_Y$ local gauge transformations. The $SU(3)_C$ group determines the couplings between strongly interacting particles by the exchange of colour carrying gauge bosons called gluons. The $SU(2)_L \times U(1)_Y$ gauge group produces the unified electroweak interaction described by Glashow, Weinberg and Salam.[7] The subscript $L$ indicates that the weak force only couples to left-handed particles. The weak hypercharge $Y$ is related to the electric charge $Q$ and the weak isospin $T^3$ by $Q = T^3 + Y/2$.

The masses of the gauge bosons and fermions are the result of couplings between the gauge or fermion fields and a scalar field called a Higgs field. These couplings are required in order to generate particle masses in a gauge invariant, Lorentz invariant and renormalizable way. The Higgs field spontaneously breaks the local $SU(2)_L \times U(1)_Y$ gauge symmetry to $\mathrm{U}(1)_Q$, to produce the separate EM and weak forces. The resulting massive gauge bosons, $W^{\pm}$ and $Z^0$, are associated with the weak interaction, while the photon ($\gamma$), which is associated with the unbroken $U(1)_Q$ symmetry, remains massless. As a consequence of this mechanism, the SM predicts the existence of a single scalar particle of unknown mass called a Higgs boson.[8]

### *1.2. Spontaneous Symmetry Breaking*

The theory describing the weak and electromagnetic forces is a local gauge theory. The essential properties of these theories can be illustrated using the simple case of an abelian gauge symmetry.[9] Suppose there exists a complex scalar field $\phi(x)$

$$\phi(x) \equiv \frac{\phi_1 \pm i\phi_2}{\sqrt{2}} \tag{1}$$

which is invariant under a local gauge transformation

$$\phi \rightarrow \phi' = \ e^{ig\chi(x)}\phi \quad , \tag{2}$$

where $\chi(x)$ is a function of the spatial coordinates. Invariance under such a a transformation requires the existence of a gauge field $A^\mu$ which compensates for the effect of the transformation on $\phi$. This field must transform as

$$A^\mu \rightarrow A'^\mu = A^\mu - \partial^\mu \chi^\mu(x) \quad . \tag{3}$$

The Lagrangian (density) describing the kinetic energy of the fields $\phi$ and $A^\mu$ is

$$\mathcal{L}_{free} = |D^\mu \phi|^2 - \frac{1}{4} F^{\mu\nu} F_{\mu\nu} \tag{4}$$

where $D^\mu \equiv \partial^\mu + igA^\mu$ is the covariant derivative and $F^{\mu\nu} \equiv \partial^\mu A^\nu - \partial^\nu A^\mu$ is the field strength tensor for $A^\mu$. The gauge boson associated with the field $A^\mu$ must be massless because if one attempts to introduce a mass term (of the form $1/2m^2 A^\mu A_\mu$) to this Lagrangian "by hand", the gauge invariance of the theory is destroyed. However, the gauge bosons $W^\pm$ and $Z^0$ which are associated with the weak interaction are known to be massive.[10] Any theory describing the weak interaction must somehow incorporate gauge boson masses while retaining the desired symmetry properties of the Lagrangian. This can be achieved by "spontaneously" breaking the symmetry by introducing a potential $V(\phi)$

$$V(\phi) = \mu^2 |\phi|^2 - \lambda |\phi|^4 \tag{5}$$

(where $\lambda$ is real and positive) which is an even function of the field $\phi$. The ground state (or vacuum expectation) $< \phi >_0$ of the system corresponds to the absolute minimum of this potential. If $\mu^2 > 0$ then the potential has a unique minima at $\phi = 0$ as shown in figure 1a, so the vacuum expectation is given by $< \phi >_0 = 0$. In this case the potential is symmetric in $\phi$, so the symmetry of $\mathcal{L}$ is unbroken. If $\mu^2 < 0$ the potential possesses degenerate local minima at values of $\phi \neq 0$ (figure 1b). The possible ground states of $\phi$ are given by

$$< \phi >_0^2 = -\mu^2/2|\lambda| \equiv \nu^2/2 \quad . \tag{6}$$

Because the ground state of the system does not possess the symmetry of the Lagrangian, the symmetry is said to be spontaneously broken.

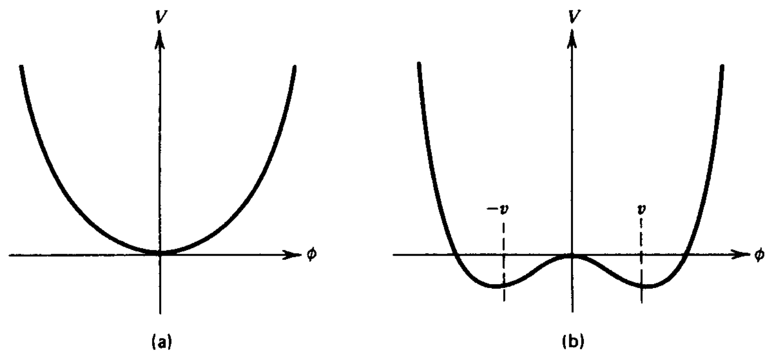


Figure 1: The potential $V(\phi)$ for the two cases $\mu^2 > 0$ (a) and $\mu^2 < 0$ (b)

This situation can be illustrated by considering an infinite array of magnetic dipoles (a Heisenberg ferromagnet). Above the Curie temperature the individual dipoles are randomly oriented, so that the array has no net magnetic dipole and the system is therefore rotationally invariant. As the temperature decreases to the Curie point, interactions between neighbouring dipoles cause the array to become spontaneously aligned, producing a net magnetic dipole oriented in some arbitrary direction. The rotational invariance of the system is therefore broken, however the underlying rotational symmetry of the Lagrangian can be observed by repeatedly heating and cooling the array and noting that the dipole moment changes direction each time.

### *1.3. The Higgs Mechanism*

In order to see how mass terms for the gauge bosons arise through spontaneous symmetry breaking, it is necessary to expand the field $\phi$ about the ground state as is appropriate in the case of small perturbations. This requires the selection of a particular ground state out of the continuum of degenerate minima. $\phi$ can be choosen to be real ($\phi_1 = \nu/\sqrt{2}$, $\phi_2 = 0$) without loss of generality, and can therefore be expanded about the ground state in terms of two real fields $\eta$ and $\xi$;

$$\phi(x) = e^{i\xi(x)/\nu}[\nu + \eta(x)]/\sqrt{2}$$

Substituting this expression into the Lagrangian from equations 4 and 5 gives

$$\mathcal{L} = \frac{1}{2}|\partial_\mu \eta|^2 + 2\mu^2\eta^2 + \frac{1}{2}|\partial_\mu \xi|^2 - \frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{g^2\nu^2}{2}A_\mu A^\mu + g\nu A_\mu(\partial^\mu \xi) + \text{ h.o.}$$

where $\mathcal{L} = \mathcal{L}_{free} + V(\phi)$. This expression can be expressed in a simpler form by making an appropriate gauge transformation

$$\phi \to \phi' = e^{-i\xi/\nu}\phi(x) = \frac{\nu + \eta(x)}{\sqrt{2}} \quad \text{so that} \quad A_\mu \to A'_\mu = A_\mu + \frac{1}{g\nu}\partial_\mu \xi \quad .$$

This leads to an effective Lagrangian containing the desired mass term for the gauge boson field .

$$\mathcal{L} \approx \frac{1}{2}[|\partial_\mu \eta|^2 + 2\mu^2\eta^2] - [\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{g^2\nu^2}{2}A_\mu A^\mu] \tag{7}$$

The last two terms in this expression correspond to a gauge boson with mass $m = g\nu$. The $\xi$ field has dissapeared, having been "eaten" by the gauge field in order to produce the new longitudinal polarization state of the massive gauge boson. The first expression in square brackets describes a scalar field with mass $m_H = \sqrt{-2\mu^2}$ associated with a particle called a Higgs boson. Because $\mu$ is a free parameter in this theory, $m_H$ is not known. Scalar bosons arise in this way anytime a continuous symmetry is spontaneously broken.[11] The interaction between the massless gauge fields and the massless scalar field which results in gauge boson masses and no massless scalar particles is known as the Higgs mechanism.[8]

## 2. The Minimal Supersymmetric Model

There are a number problems associated with the Higgs sector of the minimal SM which seem to suggest that it may be only a part of a larger structure. For example, quadratic divergences arise in the calculation of first order corrections to the Higgs boson mass which cannot be regulated within the SM. A solution to this problem is to postulate additional particle loops which cancel the divergent terms. For example, in supersymmetry (SUSY) theories this cancellation occurs by the contributions of loop graphs involving the supersymmetric partners of SM particles.[12]

Extensions of the minimal SM are constrained by two important experimental observations. The first is the fact that the ratio

$$\rho \equiv \frac{m_W^2}{m_Z^2 \cos^2\theta_W} \approx 1 \tag{8}$$

where $m_W$ and $m_Z$ are the masses of the $W^\pm$ and $Z^0$ respectively and $\theta_W$ is the weak mixing angle. This condition is determined by the Higgs sector of the theory, and is satisfied by any model containing an arbitrary number of Higgs singlets and doublets. In particular, $\rho = 1$ occurs automatically in the SM since it possesses only a single

$Y = +1$ Higgs doublet. Models containing more complex Higgs representations (such as Higgs triplets) do not in general satisfy this criteria, although Higgs representations with $SU(2)_L$ isospin $T$ and hypercharge $Y$ satisfying

$$(2T+1)^2 - 3Y^2 = 1 \tag{9}$$

will have $\rho = 1$. Many models not satisfying this condition can be artificially made to have $\rho \approx 1$ at tree level by unnatural fine tuning of the model parameters.[2]

The second major constraint on non-minimal models is the experimental non-observation of flavour changing neutral currents (FCNCs). In the minimal SM these are automatically absent at tree level, however this is not generally true for models containing more than one Higgs doublet. It has been proven that tree level FCNCs (mediated by Higgs bosons) do not occur as long as all fermions of a given electric charge couple to only a single Higgs doublet.[13]

### *2.1. Supersymmetry*

SUSY theories postulate the existence of an additional symmetry between particles of different spins.[12] Each fermion is associated with a boson of the same electric charge, lepton number and baryon number. Because supersymmetry is not observed to be an exact symmetry of nature, it must be broken in such a way that supersymmetric particles acquire masses which are different than their SM partners, and presumably larger than the energy scale which has been probed at existing particle accelerators. If SUSY breaking occurs at an energy scale much less than the Planck scale then it is reasonable to assume that it is broken only by very small non-perturbative corrections rather than at tree level, since otherwise the breaking strength would be expected to be of the same order as the energy scale of the theory (ie. the Planck scale). This could provide a solution to the so called "heirarchy problem", which addresses the reason why the electroweak energy scale ($\sim 10^2$ GeV) should be so much less than the grand unification scale or the Planck scale. If electroweak symmetry breaking is to occur at such low energies, then the SM Higgs doublet must be massless at energies of the order of the Planck scale. It is however "unnatural" to have massless elementary charged scalars in non-supersymmetric theories.[14] However, in SUSY theories massless scalars must be associated with massless fermions and therefore must exist at sufficiently high energies that SUSY is unbroken. Once SUSY breaking has occurred, these scalars no longer are required to be massless. If SUSY is spontaneously broken by non-perturbative effects at some energy scale much less than the Planck scale, the Higgs doublets could acquire a (negative) mass squared by these same effects, thus initiating $SU(2)_L \times U(1)_Y$ symmetry breaking. SUSY therefore provides a possible explanation for the "low" energy scale of the electroweak interaction. The requirement of perturbative unitarity at the energy scale $\sim 10^2$ GeV suggests that evidence of SUSY should be observable at energies of the order of 1 TeV [2](although the supersymmetry breaking scale may be somewhat higher). In particular, at least one Higgs boson must have a mass less than about one TeV.

SUSY theories are the only known theoretical framework in which the problems of heirarchy and naturalness can be resolved without requiring that the Higgs bosons be composite in nature.[15] It is hoped that these theories will permit the eventual unification of gravity and particle physics at the Planck scale ($10^{19}$ GeV).

### *2.2. The MSSM*

The minimal supersymmetric extension of the Standard Model (MSSM) is minimal in the sense that it contains the simplest possible Higgs structure and the fewest new parameters permitted in a supersymmetric theory.[15] It does however predict a wealth of new physics phenomenology beyond the SM in both the Higgs and fermion sectors. Each SM particle is associated with a supersymmetric partner with spin ($S$) differing by a half integer (see table 2). For example, the leptons and quarks ($q$ and $l$) are associated with supersymmetric bosons called squarks ($\tilde{q}$) and sleptons ($\tilde{l}$). An exact discrete symmetry ensures lepton ($L$) and baryon ($B$) number conservation. This results in a new multiplicative quantum number called R-parity, which is defined as

$$R \equiv (-1)^{3B+L+2S} \tag{10}$$

All SM particles have $R = +1$ while supersymmetric particles have $R = -1$.

Before symmetry breaking the MSSM gauge group is $SU(3)_C \times SU(2)_L \times U(1)_Y$ , as in the SM. This symmetry is spontaneously broken to $U(1)_Q$ by introducing a pair of Higgs doublets $\phi_1, \phi_2$ with opposite hypercharge along with a potential $V(\phi_1, \phi_2)$. In a general two Higgs doublet model,[2] the Higgs fields can be written as

$$\phi_1 = \begin{pmatrix} \phi_1^0 \\ \phi_1^- \end{pmatrix} \quad , \quad \phi_2 = \begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix} \quad , \tag{11}$$

where the components of $\phi_1$ and $\phi_2$ are complex scalar fields. Two doublets are required in order to generate masses for all of the particles in the MSSM while avoiding tree-level FCNC's. An additional reason is that SUSY associates scalar fields with

| Before SUSY Breaking | | | After Symmetry Breaking | | |
|---|---|---|---|---|---|
| $S=0$ | $S=1/2$ | $S=1$ | $S=0$ | $S=1/2$ | $S=1$ |
| $\tilde{q}_{\mathrm{L,R}}$ | $q$ | | $\tilde{q}_{\mathrm{L,R}}$ | $q$ | |
| $\tilde{l}_{\mathrm{L,R}}$ | $l$ | | $\tilde{l}_{\mathrm{L,R}}$ | $l$ | |
| | $\tilde{g}$ | $g$ | | $\tilde{g}$ | $g$ |
| | $\tilde{\gamma}$ | $\gamma$ | | $\tilde{\chi}_4^0$ | $\gamma$ |
| | $\tilde{z}^0$ | $Z^0$ | | $\tilde{\chi}_3^0$ | $Z^0$ |
| $h^0$, $A^0$ | $\tilde{h}^0$ | | $h^0$, $A^0$ | $\tilde{\chi}_2^0$ | |
| $H^0$ | $\tilde{H}^0$ | | $H^0$ | $\tilde{\chi}_1^0$ | |
| $H^\pm$ | $\tilde{h}^\pm$ | | $H^\pm$ | $\tilde{\chi}_2^\pm$ | |
| | $\tilde{w}^\pm$ | $W^\pm$ | | $\tilde{\chi}_1^\pm$ | $W^\pm$ |

Table 2: Spectrum of particles in the MSSM before and after SUSY and $SU(2)\times U(1)$ symmetry breaking. After symmetry breaking, mixing occurs between the neutral fermion states to produce the neutralinos ($\tilde{\chi}_i^0$) and between the charged Higgsinos and the $W$-inos to produce the charginos ($\tilde{\chi}_i^\pm$).

fermions of a given helicity, so a Higgs doublet is required for each fermion helicity state. The doublet with hypercharge $Y = -1$ ($\phi_1$) couples to the down-type quarks and the charged leptons, while the $Y = +1$ doublet ($\phi_2$) couples only to the up-type quarks. The general form of the potential which possesses the required gauge symmetry is

$$\begin{aligned} V(\phi_1, \phi_2) = \; & \lambda_1(\phi_1^\dagger\phi_1 - \nu_1^2)^2 + \lambda_2(\phi_2^\dagger\phi_2 - \nu_2^2)^2 \\ & +\lambda_3[(\phi_1^\dagger\phi_1 - \nu_1^2) + (\phi_2^\dagger\phi_2 - \nu_2^2)]^2 + \lambda_4[(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) - (\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1)] \\ & +\lambda_5[\mathrm{Re}(\phi_1^\dagger\phi_2) - \nu_1\nu_2\cos\xi]^2 + \lambda_6[\mathrm{Im}(\phi_1^\dagger\phi_2) - \nu_1\nu_2\sin\xi]^2 \end{aligned} \tag{12}$$

where $\nu_1$ and $\nu_2$ are real parameters. This potential also possesses a discrete symmetry $\phi_1 \to -\phi_1$ which is only "softly" violated, ensuring that FCNC's due to loop graphs are not too large. If all of the $\lambda_i$ are non-negative the symmetry is broken to $U(1)_Q$ as required. The vacuum expectation values of the two Higgs doublets can then be choosen without loss of generality to be

$$< \phi_1 >= \begin{pmatrix} 0 \\ \nu_1 \end{pmatrix} \quad , \quad < \phi_2 >= \begin{pmatrix} 0 \\ \nu_2 \end{pmatrix} \quad . \tag{13}$$

In the MSSM, this potential must also spontaneously break SUSY, placing further constraints on the form of $V(\phi_1, \phi_2)$ which are not discussed here.[16] Restricting $\lambda_5 = \lambda_6$ allows the phase $\xi$ to be removed by a suitable redefinition of the fields $\phi_1$ and $\phi_2$. This ensures that there is no CP violation in the Higgs sector of the MSSM, and that $\nu_1$ and $\nu_2$ are real and non-negative. Additional relations between the $\lambda_i$'s are obtained from the requirement that $\nu_1$ and $\nu_2$ are non-zero;

$$\begin{aligned} & \lambda_2 = \lambda_1 \\ & \lambda_3 = 1/8\,(g^2 + g'^2) - \lambda_1 \\ & \lambda_4 = 2\lambda_1 1/2\,g'^2 \\ & \lambda_5 = \lambda_6 = 2\lambda_1 - 1/2\,(g^2 + g'^2) \end{aligned} \tag{14}$$

where $g$ and $g'$ are the $U(1)$ and $SU(2)$ gauge group couplings respectively.

After symmetry breaking, the neutral Higgsinos ($\tilde{h}^0, \tilde{H}^0$) and the neutral gauginos ($\tilde{\gamma}, \tilde{z}^0$) mix to produce the four neutralinos ($\tilde{\chi}_i^0$), and a similar mixing of the $\tilde{h}^\pm$ and $\tilde{w}^\pm$ produces the charginos ($\tilde{\chi}_i^\pm$). As a consequence of R-parity conservation supersymmetric particles must be produced in pairs, and will generally decay rapidly to lighter SUSY particles. The lightest supersymmetric particle (LSP), which is generally taken to be $\tilde{\chi}_1^0$, cannot decay to lighter particles and therefore is absolutely stable.

## 3. Properties of the MSSM Higgs Sector

By counting the number of independent fields before and after symmetry breaking one can determine the number of physical Higgs bosons in the MSSM. Before symmetry breaking there are two doublets of complex scalar fields corresponding to eight real Higgs fields. There are also three massless gauge bosons with two polarization states each, for a total of fourteen degrees of freedom. After symmetry breaking, the three gauge bosons become massive and so each possess a third longitudinal polarization state (9 degrees of freedom). The remaining five fields correspond to physical Higgs bosons. Three of these are the neutral bosons $h^0, H^0$, and $A^0$, and the remaining two are charged bosons ($H^{\pm}$). The $h^0$ and $H^0$ are massive scalar particles (where $m_h < m_H$ by definition) similar to the SM Higgs boson. The $A^0$ is frequently referred to as a pseudoscalar, although this is technically incorrect. In the absence of fermions, C (charge conjugation) and P (parity) are individually conserved. In this case $J^{PC} = 0^{+-}$ for the $A^0$, indicating that it is a CP-odd scalar rather than a pseudoscalar. However, when fermions are included in the theory C and P are no longer good quantum numbers although CP is still (approximately) conserved. In this case the $A^0$ can be treated as having $J^{PC} = 0^{-+}$, making it a a CP-odd pseudoscalar. $J^{PC}$ and $J^P$ for the gauge and Higgs bosons in the MSSM are listed in table 3.

Because $\phi_1$ couples only to down-type fermions, the masses of these particles are proportional to $\nu_1$. Similarily, the masses of up-type fermions are proportional to $\nu_2$. The values of $\nu_1$ and $\nu_2$ are not individually known, however because the gauge bosons acquire mass by coupling to both Higgs doublets, the quadratic sum of $\nu_1$ and $\nu_2$ is given by

$$\sqrt{\nu_1^2 + \nu_2^2} = \frac{2m_W}{g} = 246 \text{ GeV} \tag{15}$$

in the low-energy limit of electroweak theory.

At tree level, the MSSM can be completely specified by five parameters. These are frequently choosen to be the squark and gluino masses ($m_{\tilde{g}}, m_{\tilde{q}}$), the Higgsino mass parameter ($\mu$), the mass of any one of the Higgs bosons (denoted $m_\phi$), and the ratio of the two vacuum expectation values

$$\tan\beta \equiv \frac{\nu_2}{\nu_1} \quad . \tag{16}$$

Of these, only $\tan\beta$ and $m_\phi$ are necessary to fully specify the Higgs sector. Once these values have been fixed, the masses of the remaining gauge bosons are given by

| | $J^{PC}$ | | | $J^{P}$ | |
|---|---|---|---|---|---|
| $\gamma$ | $1^{--}$ | $(1^{--})$ | $W^{\pm}$ | $1^{-}$ | $(1^{-},1^{+})$ |
| $Z^0$ | $1^{--}$ | $(1^{--},1^{++})$ | | | |
| $H^0$ | $0^{++}$ | $(0^{++},0^{--})$ | $H^{\pm}$ | $0^{+}$ | $(0^{+},0^{-})$ |
| $h^0$ | $0^{++}$ | $(0^{++},1^{--})$ | | | |
| $A^0$ | $0^{+-}$ | $(0^{+-},0^{-+})$ | | | |

Table 3: $J^{PC}$ and $J^{P}$ for the bosons in the MSSM in the absence of fermions. The quantities in parethesis are the allowable $J^{PC}$ and $J^{P}$ assignments when fermions are included in the theory.

the following tree level relations:

$$\begin{aligned} m^2_{H^0,h^0} &= \frac{1}{2}\left(m_A^2+m_Z^2 \pm \sqrt{(m_A^2+m_Z^2)^2-(2m_A m_Z\cos 2\beta)^2}\right) \\ m^2_{H^\pm} &= m_W^2+m_A^2 \\ m_h^2+m_H^2 &= m_A^2+m_Z^2 \end{aligned} \tag{17}$$

The mixing angle $\alpha$ which diagonalizes the $h^0, H^0$ mass matrix is given by

$$\cos 2\alpha = -\cos 2\beta \frac{m_A^2-m_Z^2}{m_H^2-m_h^2} \quad , \qquad -\frac{\pi}{2} < \alpha \le 0 \quad . \tag{18}$$

These expressions lead to tree level mass relations between the gauge and Higgs bosons which are true for arbitrary MSSM parameterizations

$$m_h \le m_A \le m_H \quad , \qquad m_W \le m_{H^\pm} \quad \text{and} \quad m_h \le m_Z \le m_H \quad . \tag{19}$$

These relations are not strictly true if radiative corrections are included in the calculation of the Higgs masses.[17,18] In particular, the recent evidence for a heavy top quark $(t)$[5,19] indicates that radiative corrections may increase the mass of the $h^0$ well beyond that of the $Z^0$. Assuming that $m_A > m_Z$ and that $m_Z \ll m_t \ll m_{\tilde{t}}$ (where $m_{\tilde{t}}$ is the top squark mass), one-loop radiative corrections give an approximate upper bound on $m_h$ of[20]

$$m_h^2 \le m_Z^2 + \frac{3g^2 m_Z^4}{16\pi^2 m_W^2} \times \left[\ln\left(\frac{m_{\tilde{t}}^2}{m_t^2}\right)\left(\frac{2m_t^4-m_t^2 m_Z^2}{m_Z^4}\right)+\frac{m_t^2}{3m_Z^2}\right] \quad . \tag{20}$$

Given a top quark mass of $\sim 170$ GeV, $m_h$ can exceed $m_Z$ by $\sim 30$ GeV for $m_{\tilde{t}}$ of the order of 1 TeV. Radiative corrections to the masses of the remaining Higgs bosons can also be significant, and generally increase the tree level masses.

### *3.1. Couplings to Gauge Bosons*

In addition to generating gauge boson and fermion masses, spontaneous symmetry breaking also produces coupling terms in the MSSM Lagrangian between the Higgs bosons and the vector bosons ($V$). The requirement that $V_L V_L \rightarrow V_L V_L$ does not exceed unitarity limits restricts the couplings of the two neutral scalar Higgs to be

$$g^2_{H^0VV} + g^2_{h^0VV} = g^2_{H_{SM}VV} \quad , \tag{21}$$

where $g_{H_{SM}VV}$ represents the coupling of the SM Higgs boson a pair of vector bosons $W^\pm$ and $Z^0$. More specifically

$$\begin{aligned} g_{h^0VV} &= g_{H_{SM}VV}\sin(\beta-\alpha) \\ g_{H^0VV} &= g_{H_{SM}VV}\cos(\beta-\alpha) \end{aligned} \tag{22}$$

where $\alpha$ and $\beta$ are the angles defined in equations 16 and 18. The couplings of the neutral scalar Higgs bosons of the MSSM are therefore suppressed relative to the equivalent SM Higgs couplings to pairs of vector bosons by the factors in table 4. The degree of this suppression is dependent on the values of the parameters $m_\phi$ and $\tan\beta$. In most realistic scenarios $\cos(\beta-\alpha) \ll 1$ and $\sin(\beta-\alpha) \approx 1$, so that the $H^0$ is highly suppressed while the $h^0$ couplings to the gauge bosons are similar to the SM Higgs boson couplings. This suppression is illustrated as a function of $m_{H^+}$ in figure 2. In contrast to the $h^0$ and $H^0$, couplings between the pseudoscalar $A^0$ and pairs of vector bosons are not permitted at tree level due to CP conservation.

A feature of two Higgs doublet models which is not present in the minimal SM is the existence of couplings involving a gauge boson and two Higgs bosons. These couplings only occur between non-identical Higg bosons such as $W^+H^-A^0$ and $Z^0H^+H^-$. Unitarity constraints uniquely determine the strength of four of these couplings:

$$\begin{aligned} g_{h^0A^0Z^0} &= \tfrac{g}{2\cos\theta_W}\cos(\beta-\alpha) & \qquad g_{H^0A^0Z^0} &= \tfrac{g}{2\cos\theta_W}\sin(\beta-\alpha) \\ g_{H^+W^-h^0} &= \tfrac{g}{2}\cos(\beta-\alpha) & \qquad g_{H^+W^-H^0} &= \tfrac{g}{2}\sin(\beta-\alpha) \end{aligned} \tag{23}$$

Although the $H^+H^-\gamma$ vertex is permitted at tree level, couplings between the neutral Higgs bosons and one or two photons vanish. These couplings do however occur at the one-loop level, generated radiatively through fermion loops. The decay $h^0 \rightarrow \gamma\gamma$ which occurs by this mechanism is of great importance in the detection of a light neutral Higgs in both the SM and the MSSM (see figure 3a). The couplings $h^0gg$, $H^0gg$ and $A^0gg$ arise in a similar fashion and are expected to be an important Higgs production mechanism (known as gluon-gluon fusion) at hadron colliders such as the LHC (figure 3b).

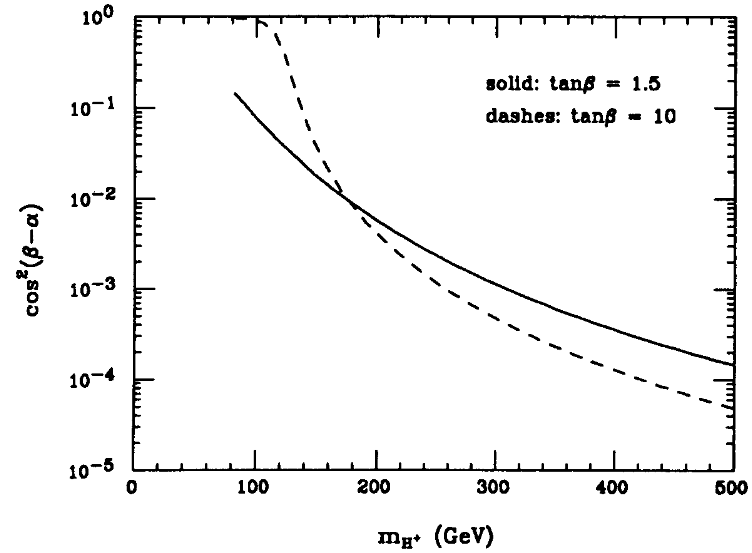


Figure 2: Suppression factor $\cos^2(\beta - \alpha)$ as a function of $m_{H^+}$ for two different values of $\tan\beta$.

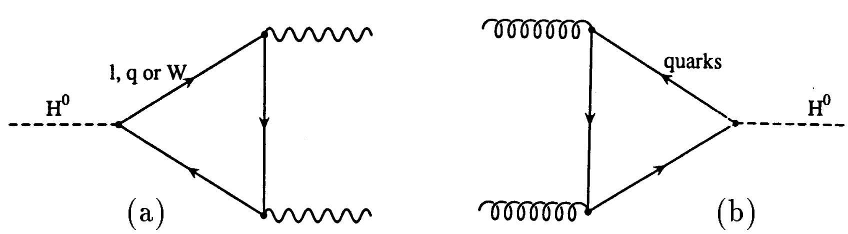


Figure 3: Phenomenologically important Higgs boson couplings which arise through fermion loops. a) The $H^0 \to \gamma\gamma$ decay mode. b) Higgs production by gluon-gluon fusion. $H^0$ here means any of the three neutral bosons $h^0, A^0$ or $H^0$.[21]

| | $h^0$ | $H^0$ | $A^0$ |
|---|---|---|---|
| $Z^0Z^0$ | $\sin(\beta-\alpha)$ | $\cos(\beta-\alpha)$ | - |
| $W^+W^-$ | $\sin(\beta-\alpha)$ | $\cos(\beta-\alpha)$ | - |
| $u\bar{u}$ | $\frac{\cos\alpha}{\sin\beta}$ | $\frac{\sin\alpha}{\sin\beta}$ | $-i\gamma_5\cot\beta$ |
| $d\bar{d}$ | $-\frac{\sin\alpha}{\cos\beta}$ | $\frac{\cos\alpha}{\cos\beta}$ | $-i\gamma_5\tan\beta$ |

Table 4: Suppression factors for neutral Higgs boson couplings to fermion and gauge boson pairs relative to SM Higgs boson couplings.

### *3.2. Couplings to Fermions*

Fermion masses are generated by introducing Yukawa-like couplings to the Higgs field as in the SM.[9] After symmetry breaking, these terms produce couplings between the fermions and the MSSM Higgs bosons with strengths which are proportional to the fermion masses. As a result, the dominant Higgs boson couplings are to the fermions of the third generation $(t, b, \tau)$, which are more massive than the other two generations. The interaction terms in the MSSM Lagrangian are

$$\begin{aligned}\mathcal{L}_{Hf\bar{f}} = &-\frac{g}{2m_W\cos\beta}\,\bar{d}\,M_d\,d\,(H^0\cos\alpha - h^0\sin\alpha) + \frac{ig\tan\beta}{2m_W}\,\bar{d}\,M_d\,\gamma_5\,d\,A^0 \\ &-\frac{g}{2m_W\sin\beta}\,\bar{u}\,M_u\,u\,(H^0\sin\alpha + h^0\cos\alpha) + \frac{ig\cot\beta}{2m_W}\,\bar{u}\,M_u\,\gamma_5\,u\,A^0 \\ &+\frac{g}{2\sqrt{2}m_W}(H^+\bar{u}\,[\cot\beta\,M_uU_{ckm}(1-\gamma_5) + \tan\beta\,U_{ckm}M_d(1+\gamma_5)]\,d)\end{aligned} \tag{24}$$

where $u$ and $d$ represent the three generations of up-type and down-type quarks respectively and $M_u$ and $M_d$ are the corresponding diagonal mass matrices. $U_{ckm}$ is the CKM mixing matrix.[6] The interactions between leptons and Higgs bosons can be obtained by substituting the appropriate lepton fields and mass matrices for $u$, $d$, $M_u$ and $M_d$, and replacing $U_{ckm}$ with 1.

The first and third terms of equation 24 describe interactions with the neutral scalar Higgs bosons $h^0$ and $H^0$, which are identical in form to SM Higgs interactions with fermions apart from a factor which depends on $\beta$ and the Higgs mass mixing angle $\alpha$. Because $\phi_1$ couples only to down-type quarks and leptons, these fermions have couplings proportional to $1/\nu_1$ ($\propto 1/\cos\beta$) while up-type fermions have couplings proportional to $1/\sin\beta$. From the second and fourth terms of equation 24 it is apparent that the $A^0$ couples to fermion pairs via the $\gamma_5$ matrix as expected for a pseudoscalar particle. The vertex factors for the couplings of neutral Higgs bosons to fermions can be obtained from the SM Higgs couplings by including the factors given in table 4. In the case $\tan\beta > 1$ the $h^0$ and $H^0$ couplings to $u\bar{u}$ will be suppressed relative to SM couplings, while $d\bar{d}$ couplings will be enhanced. For $\tan\beta < 1$ the opposite is true. It should be noted that the presence of diagonal quark mass matrices $M_u$ and $M_d$ ensures the absence of FCNCs in interactions involving the three neutral Higgs bosons.

The charged Higgs $H^\pm$ couples up-type fermions to down-type fermions in a manner similar to the charged current interactions mediated by the $W^\pm$. The presence of $U_{ckm}$ in the last two terms of equation 24 permits mixing between quarks of different generations. For interactions between the $H^+$ and leptons, the neutrino mass matrix is zero so that the interaction takes the form

$$\mathcal{L}_{H^+f\bar{f}} \propto H^+ \tan\beta\, M_l\, \bar{\nu}\,(1+\gamma_5)\, l = H^+ \tan\beta\, M_l\, \bar{\nu}\, l_R$$

so right-handed charged leptons are coupled to right-handed antineutrinos as expected.

## 4. Limits and Discovery Potential

In this section, SUSY Higgs production mechanisms and decay channels are discussed in the context of $e^+e^-$ collisions at the $Z^0$ resonance such as occur at LEP I,[22] and $pp$ collisions at the Large Hadron Collider (LHC) proposed by CERN.[3] The LHC will collide two proton beams at a centre of mass energy of 14 TeV with a luminosity of $10^{34}$ cm$^{-2}$s$^{-1}$. Two general purpose experiments will search for signatures of SM and non-SM Higgs decays. This paper concentrates on Higgs discovery prospects using the ATLAS experiment[23] at the LHC. It is assumed here that the masses of supersymmetric particles are greater than MSSM Higgs boson masses, so that only Higgs boson decays to SM particles need to be considered.

### *4.1. MSSM Higgs searches at LEP*

It is clear from the tree level mass relations (equation 19) that only the $h^0$ and $A^0$ may be light enough to be found at LEP I, although radiative corrections do not permit the MSSM to be ruled out if no Higgs is found. The dominant Higgs bosons production mechanisms at $e^+e^-$ colliders proceed via couplings to the intermediate vector bosons. The lightest neutral Higgs boson $h^0$ could be produced at LEP I through the bremsstrahlung process

$$e^+e^- \to Z^0 \to Z^{0*}h^0$$

provided $m_{h^0} < m_Z$. The decay products of the $Z^{0*}$ can be used for event tagging. In practice only leptonic $Z^0$ decays are considered because of the large multijet backgrounds. Clearly the $A^0$ cannot be produced in this way since it does not couple to pairs of gauge bosons. Given a branching ratio for $(Z^0 \to l^+l^-)$ of $\sim 3.3\%$ where $l$ is either $\mu$ or $e$, the SM event rate for

$$Z^0 \to Z^{0*}H \to Hl^+l^-$$

where $H$ is a SM Higgs boson, decreases from $\sim 150$ to less than one event per $10^6$ $Z^0$s over the Higgs mass range 10 GeV to 60 GeV.[24] The equivalent MSSM process $Z^0 \to h^0l^+l^-$ is suppressed relative to the SM process by a factor $\sin^2(\beta-\alpha)$. The effect of this suppression is illustrated as a function of $\tan\beta$ in figure 4. For large values of $\tan\beta$, $h^0$ production is strongly suppressed, making its discovery at LEP I very difficult. For small $\tan\beta$, $h^0$ production rates are similar to SM Higgs production, in which case the $h^0$ may be observable, but it may not be possible to distinguish it from a SM Higgs boson.

The ideal situation from an experimental point of view is if both $m_h$ and $m_A$ are less than $m_Z$. In this case, as $\tan\beta$ goes to unity, $m_h$ approaches zero, whereas if $\tan\beta$ is large then the $h^0$ and $A^0$ are approximately degenerate in mass. If the two Higgs bosons are sufficiently light, the mechanism

$$e^+e^- \to Z^0 \to A^0h^0 \qquad .$$

could permit the production of $h^0A^0$ pairs at LEP I. The $Z^0$ partial width to Higgs boson pairs is proportional to $\cos^2(\beta-\alpha)$, so this process complement the Higgs bremsstrahlung process. If either the bremsstrahlung or pair production process happens to be strongly suppressed, then the other process will be relatively unsuppressed. The branching ratio for $h^0A^0$ pair production compared with the SM process $Z^0 \to \nu\bar{\nu}$ is shown as a function of $\tan\beta$ in figure 5.

As noted in the previous section the couplings of the neutral Higgs bosons to fermions can be either suppressed or enhanced relative to SM Higgs couplings depending on the value of $\tan\beta$. The dominant decay channels of the $h^0$ and $A^0$ are

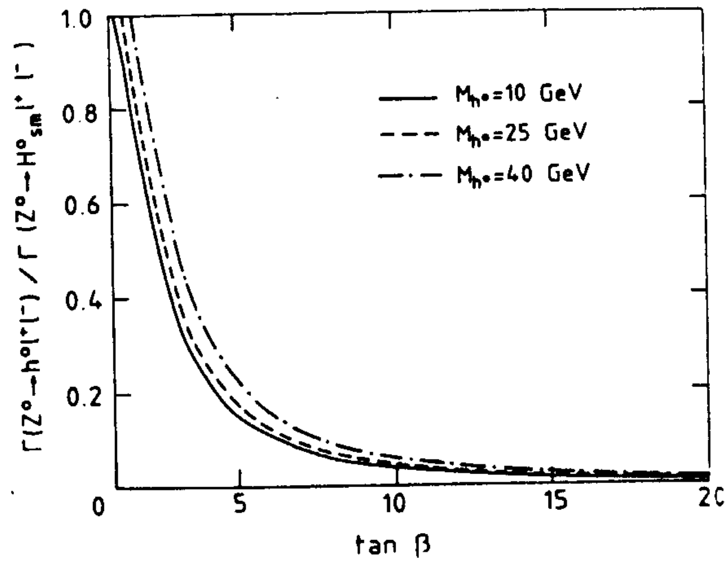


Figure 4: Suppression factor for $Z^0 \to h^0 l^+ l^-$ compared to the SM $Z^0 \to H l^+ l^-$ channel.[25]

expected to be $(h^0, A^0 \to b\bar{b})$ and $(h^0, A^0 \to \tau\bar{\tau})$. Limits on the mass of the $h^0$ have been obtained at LEP I by considering both Higgs pair production and the bremsstrahlung processes in which the virtual $Z^0$ decays to $\mu^+\mu^-$, $e^+e^-$ or $\nu\bar{\nu}$. Both $\tau^+\tau^-$ and $b\bar{b}$ Higgs final states have been considered. Virtual $Z^0$ decays to $\mu^+\mu^-$ or $e^+e^-$ in the bremsstrahlung process can be used for event tagging and energy reconstruction, whereas large jet backgrounds to Higgs pair production from $Z^0 \to q\bar{q}$ must be reduced by considering jet multiplicity and event kinematics. In this case three signatures are possible; $h^0A^0 \to b\bar{b}b\bar{b}$, $\tau^+\tau^-b\bar{b}$ or $\tau^+\tau^-\tau^+\tau^-$. b-quarks are identified by requiring a high transverse momentum electron or muon be associated with a jet, while $\tau$'s produce a high thrust jet with only a single charge track. Selection efficiencies for these processes range from 3% to 15%.[26] Because of the complimentary nature of the bremsstrahlung and $h^0A^0$ pair production processes, by combining the results of the two searches, a region in the $(m_h, m_A)$ plane can be excluded.[26,27] This can be interpreted as as lower limit on the mass of the $h^0$ of 44 GeV, and 22 GeV for $A^0$ assuming that $\tan\beta > 1$.[4] Combined analyses of precision electroweak data from LEP, SLC and CDF have shown that current estimates of the top quark mass are consistent with a MSSM $h^0$ (or SM $H$) of mass $\sim$ 100 GeV.[28]

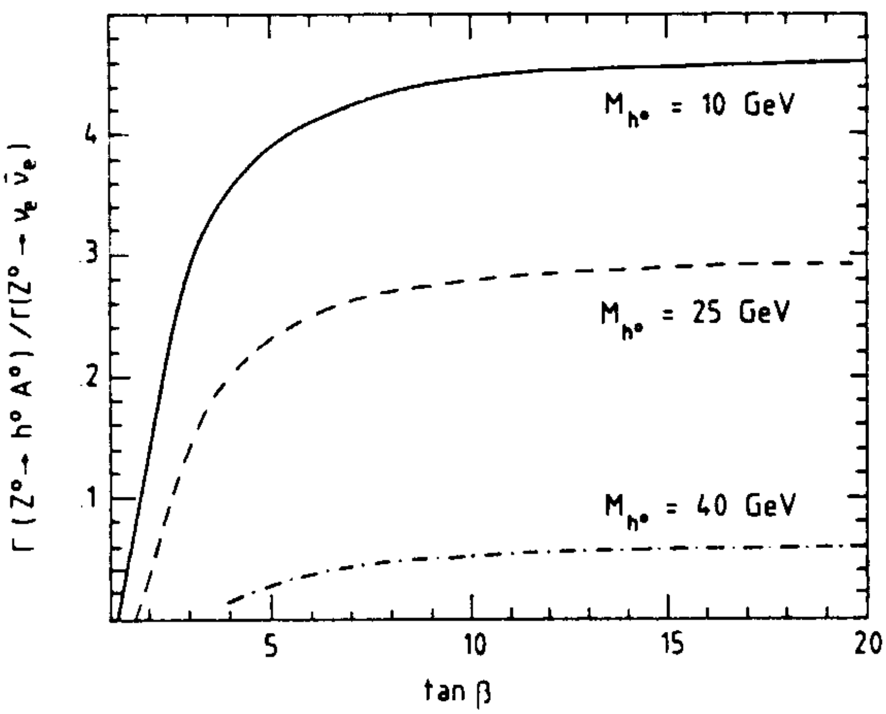


Figure 5: Branching ratio for the MSSM production process $Z^0 \to h^0 A^0$ relative to the SM process $Z^0 \to \nu\bar{\nu}$ as a function of the parameter $\tan\beta$.[25]

Charged Higgs searches have also been performed at LEP I. Although a MSSM $H^+$ must have $m_{H^+} \geq m_W$ and therefore is not expected to be produced at LEP I, these limits are included here because the discovery of a light charged Higgs could potentially rule out the MSSM. If $2m_{H^+} < m_{Z^0}$, $H^\pm$ pairs could be produced from decays of the $Z^0$ in a manner analogous to neutral Higgs pair production. The charged Higgs boson decays to the heaviest available $l\nu$ or $q\bar{q}$ pairs, which in this scenario means $H^+ \to \tau^+\nu_\tau$ or $c\bar{s}$. The three possible final state signature are therefore $Z^0 \to H^+H^- \to \tau^+\nu_\tau\tau^-\bar{\nu}_\tau$, $\tau\nu_\tau cs$ and $c\bar{s}\bar{c}s$. The dominant sources of background to these signatures are from $e^+e^- \to \tau^+\tau^-$ and $e^+e^- \to q\bar{q}$, but can be be substantially reduced by considering event kinematics. For example, the signature $\tau^+\nu_\tau\tau^-\bar{\nu}_\tau$ is characterized by acoplanar taus and significant missing energy. The scalar nature of the decaying Higgs boson can also be used. Limits from the ALEPH and L3 experiments of $m_{H^+} >$ 41 GeV at the 95% confidence level[26,29] essentially rule out the possibility $2m_{H^+} < m_Z$. The region $m_W \leq m_{H^\pm}$ predicted by the MSSM has not yet been tested by experiment.

These limits are expected to increase as a result of the upgrade of LEP to a centre of mass energy of $\sim$ 180 GeV (LEP II).[30] The dominant Higgs production mechanisms will again be due to couplings to $Z^0$ pairs:

$$\begin{aligned} e^+e^- \to \quad Z^{0*} \quad & \to h^0 Z^0 \\ & \to h^0 A^0 \\ & \to H^+H^- \end{aligned}$$

The mass limit on the SM Higgs is expected to improve to about 80 GeV, which can also be interpreted as an absolute upper limit for $h^0$ searches. The charged Higgs and heavy neutral scalar $H^0$ are not likely to be observable, therefore it is doubtful that the MSSM could either be substantiated or ruled out by LEP II.

### *4.2. Neutral Higgs discovery prospects at the LHC*

At hadron colliders such as the LHC, neutral Higgs bosons are produced predominantly by gluon-gluon fusion $gg \to h^0, A^0, H^0$ occuring through fermion loops (see figure 3b). Because the Higgs couplings are proportional to fermion masses, the dominant contribution to these couplings is due to top quark loops, and are therefore sensitive to the top quark mass. Higgs production will be further suppressed (or enhanced) by the factors in table 4 depending on the value of $\tan\beta$. Intermediate vector boson fusion $W^+W^- \to h^0, H^0$ contributes to a lesser degree to the $h^0$ and $H^0$ production cross-section.[31]

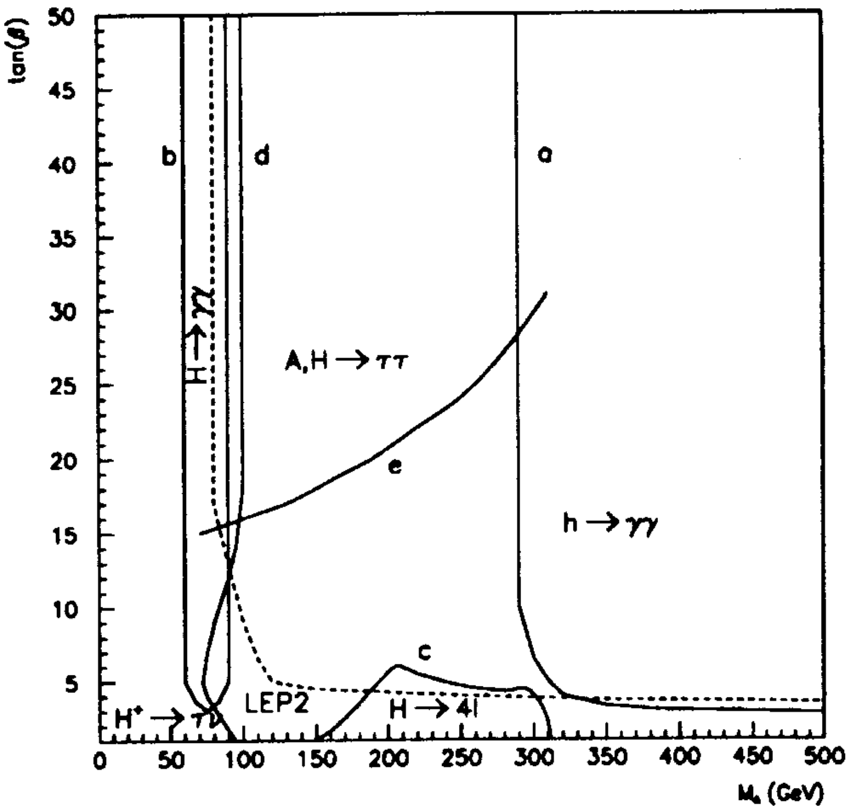


Figure 6: MSSM Higgs boson discovery contours in the $(\tan\beta, m_A)$ plane for the ATLAS experiment at the LHC. The region to the right of curve (a) and inside curve (b) can be tested via $h^0, H^0 \to \gamma\gamma$. The small region below curve (c) can be tested by searching for the $H^0$ decay to four leptons. Charged Higgs searches via $H^+ \to \tau\nu$ test the region to the left of curve (d). This curve shifts to the right for a top quark mass greater than 140 GeV. The large region above curve (e) can be tested using $\tau\tau$ decays of the $A^0$ and $H^0$.

Large QCD backgrounds at $pp$ and $p\bar{p}$ colliders make it impossible to search for neutral Higgs bosons through hadronic decay modes, therefore the rarer leptonic decay modes must be considered. For the $H^0$ and the $A^0$ the preferred search channel is the decay to $\tau\bar{\tau}$, where at least one of the $\tau$'s decays to leptons. This channel is expected to be observable in the case that $\tan\beta$ is large, so that $H^0$ and $A^0$ couplings to down-type fermions are enhanced relative to up-type. Because the cross-sections for $A^0 \rightarrow \tau\bar{\tau}$ and $H^0 \rightarrow \tau\bar{\tau}$ can both be large,[32] and for $m_{A^0} > 150$ GeV the masses of the two Higgs bosons are similar, the two signals can be combined to give a cross-section of typically 10 to 100 pb[33,34] in the region above curve (e) in figure 6. A second channel which may be available in the search for the $H^0$ is the decay $H^0 \rightarrow Z^0Z^0$ where the $Z^0$'s decay to either $\mu^+\mu^-$ or $e^+e^-$. This is the preferred channel to search for an intermediate mass SM Higgs boson, but in the MSSM this mode is strongly suppressed for large values of $\tan\beta$. As a result, it is only useful over a small region of the available parameter space (curve (c) in figure 6).

The four lepton channel is unavailable to the $h^0$ since its mass is expected to be less than about 130 GeV. The dominant $h^0$ decay is expected to be $h^0 \rightarrow b\bar{b}$, however this channel will not be observable above QCD backgrounds. The decay to two photons via charged fermion or gauge boson loops (as shown in figure 3a) is considered to be the most promising channel for $h^0$ discovery.[35] Studies have indicated that the SM $H \rightarrow \gamma\gamma$ decay[36,37] may be observable above the two photon continuum background due to $gg \rightarrow \gamma\gamma$ and $q\bar{q} \rightarrow \gamma\gamma$, provided that the detector possesses electromagnetic calorimetry capable of photon energy measurements at the level of 1 - 2%.[23] The branching ratio of $h^0 \rightarrow \gamma\gamma$ is reduced compared to the SM equivalent process by the contributions of loops due to supersymmetric fermions and bosons which have the opposite sign to the SM contributions. The suppression of $h^0$ couplings to $W^{\pm}$ (table 4) can further reduce the observed rate. As a result, the rate is only significant if $m_{A^0}$ is large so that the $h^0WW$ coupling is close to the SM coupling. In this case it is unlikely that the $h^0$ could be distinguished from a SM Higgs. The channel $H^0 \rightarrow \gamma\gamma$ could also be examined in the case that the $H^0$ is light ($<$ 100 GeV). It should be noted that there is a large region of parameter space in which the $h^0$ and $H^0$ both fall into the intermediate mass range (100 GeV $< m_h, m_H <$ 300 GeV) for which there are no adequate channels for neutral Higgs searches. However, it may be possible to cover some of this parameter space using searches for a charged Higgs boson.

### 4.3. *Charged Higgs boson searches at the LHC*

The production mechanisms and decay channels of significance in the search for a charged Higgs are strongly dependent on the relative masses of the $H^{\pm}$, the top quark and the three neutral Higgs bosons. If the mass of the charged Higgs is larger than $m_t+m_b$ then the dominant decay of the $H^+$ would be to $t\bar{b}$, however this channel is not likely to be observable at future hadron colliders due to large QCD backgrounds. The channel $H^+ \to W^+h^0$ where the $h^0$ decays predominantly to $b\bar{b}$, may be observable in this mass range if it is permitted by kinematics, although the branching ratio is small.[2] If $m_W + m_h < m_{H^+} < m_t + m_b$, then this becomes the dominant $H^+$ decay. The large jet background to $h^0 \to b\bar{b}$ can be substantially reduced by tagging the lepton (and missing transverse energy) or jets produced by the decay of the $W$.

Because of the expected high rate of $t\bar{t}$ production from gluon-gluon fusion at the LHC, the $t \to H^+b$ charged Higgs production mode is extremely important in the case that $m_{H^+} < (m_t - m_b)$.[38] $t\bar{t}$ events can be easily tagged using a single lepton trigger from the dominant $t \to Wb$ decay, where $W \to e\nu_e$ or $\mu\nu_\mu$. The decays of the second top quark can then be examined for evidence of charged Higgs production. If $\tan\beta < 1$ then the $H^+$ decays predominantly to $c\bar{s}$, so it may be seen as a peak in the $t\bar{t}$ dijet mass spectrum. Observation of this channel would require efficient b-tagging and good jet energy measurement and is therefore only likely to be observable during "low" luminosity ($\sim 10^{33}$ cm$^{-2}$ s$^{-1}$) running at the LHC.

Over the range $1 < \tan\beta < 50$, the channel $t \to H^+b$ competes with the dominant SM decay mode $t \to W^+b$, with a branching ratio ranging from 4% to 50%. In this case the $H^+ \to \tau\nu_\tau$ channel is enhanced relative to hadronic $H^+$ decays. Because of the mass dependence of the Higgs couplings, the $t \to H^+b \to \tau\nu_\tau b$ is preferred over $e\nu_e b$ and $\mu\nu_\mu b$ final states. This could result in an apparent deviation from lepton universality in top quark couplings to the $W$.[39] Specifically,

$$\frac{\mathrm{Br}(t \to \tau\nu_\tau b)}{\mathrm{Br}(t \to \mu\nu_\mu b)} = 1 + \frac{\mathrm{Br}(t \to H^+b)\ \mathrm{Br}(H^+ \to \tau\nu_\tau)}{\mathrm{Br}(t \to W^+b)\ \mathrm{Br}(W^+ \to \tau\nu_\tau)} \neq 1 \tag{25}$$

would suggest the existence of the $H^{\pm}$, although this measurement does not provide a measure of the $H^+$ mass. The presence of this alternate top decay mechanism favouring $\tau\nu_\tau$ over other leptonic final states has been proposed to explain the (until recently) non-observation of the top quark. Because these analyses have neglected $\tau\nu_\tau$ final states, the observed rates for top production may be lower than the actual rate. Present bounds on the top mass from electroweak measurements at LEP[19] and from

Fermilab[5] are not necessarily valid if the MSSM is correct. Bounds on the $H^+$ have been obtained at the CERN SPS for the case of a relatively light top quark. Searches for $t \to H^+ b$, where $H^+ \to \tau\nu_\tau$, exclude the region $m_{H^+} < 65$ GeV if $(m_t - m_{H^+})$ is less than about 15 GeV under the assumption that $\mathrm{Br}(H^+ \to \tau\nu_\tau) = 1$.[40] It should be noted that this mass range is not permitted for a charged Higgs boson from the MSSM.

## 5. Conclusion

Although there is at present no experimental evidence which contradicts the SM, there are theoretical difficulties such as naturalness and heirarchy which suggest that other possibilities should be considered. The MSSM represents the simplest case of a class of SUSY theories which utilize more complex Higgs representations to generate gauge boson and fermion masses. The couplings of these bosons to fermions and gauge bosons are similar in form to SM Higgs couplings, but are generally suppressed relative to these by factors which depend on the parameters $\tan\beta$ and $m_\phi$. Because MSSM Higgs boson production and decay cross sections are sensitive to these couplings, the channels available for Higgs searches are strongly dependent on these unknown values.

Searches for MSSM Higgs bosons at LEP I have placed limits the $h^0$ and $A^0$, but large radiative corrections to $m_h$ do not permit the MSSM to be ruled out if no Higgs bosons are found at LEP I or II. Even if a neutral Higgs boson is found at LEP, it is unlikely that it would be possible to determine if it was from the SM or the MSSM. The parameter space available to remaining Higgs bosons $(H^0, H^\pm)$ has not yet been probed at all by experiments. The LHC will be the first major accelerator at which the MSSM can be tested over a large range of the available parameters. Higgs boson searches at the LHC will be difficult, due to small cross sections for leptonic decay modes and large backgrounds to hadronic decays. Leptonic signatures such as $H^+ \to \tau\nu$ and $A^0, H^0 \to \tau^+\tau^-$ can be used to restrict large regions of the available parameter space. Rarer decays such as $h^0 \to \gamma\gamma$ and $H^0 \to l^+l^-l^+l^-$ may also be observable. The search for Higgs bosons from the MSSM will require high performance standards for LHC detectors, and a good understanding of SUSY Higgs phenomenology. The discovery of a non-SM Higgs boson would provide valuable insight into electroweak symmetry breaking and the origin of fermion and gauge boson mass.

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