CLOSING IN ON SUPERSYMMETRY

V. Barger\textsuperscript{a} and R.J.N. Phillips\textsuperscript{b}

\textsuperscript{a}Physics Department, University of Wisconsin, Madison, WI 53706, USA
\textsuperscript{b}Rutherford Appleton Laboratory, Chilton, Didcot, Oxon OX11 0QX, UK

Abstract

A survey is made of some recent ideas and progress in the phenomenological applications of Supersymmetry (SUSY). We describe the success of SUSY-GUT models, the expected experimental signatures and present limits on SUSY partner particles, and the phenomenology of Higgs bosons in the minimal SUSY model.

1 Introduction

The Standard Model (SM) has been very successful so far; all its predictions that have been tested have been verified to high precision. Important checks remain to be made, of course: the top quark is not yet discovered, the interactions between gauge bosons are still unmeasured, and the Higgs boson remains a totally unconfirmed hypothesis. In spite of its success, however, the apparent arbitrariness and various theoretical limitations of the SM suggest a need for a deeper theory such as Supersymmetry (SUSY) or Superstrings, implying new physics, new particles and new interactions beyond the SM. In this lecture we review some recent developments in SUSY phenomenology: unification of couplings in SUSY-GUT models, experimental signals from SUSY, and Higgs phenomenology in the minimal SUSY extension of the SM (MSSM).

With SUSY,\textsuperscript{1} each fermion has a boson partner (and vice versa), with all the same quantum numbers but with spin differing by $1/2$. Since no such partners have been found, SUSY is plainly a broken symmetry at presently accessible mass scales but could be restored above some higher scale $M_{\text{SUSY}}$.

\textsuperscript{1}Talk presented by V. Barger at the Symposium in Honor of Tetsuro Kobayashi’s 63rd Birthday, Tokyo, March 1993
The primary theoretical motivation for SUSY is that it stabilizes divergent loop contributions to scalar masses, because fermion and boson loops contribute with opposite signs and largely cancel. This cures the naturalness problem in the SM, so long as \( M_{\text{SUSY}} \lesssim \mathcal{O}(1 \, \text{TeV}) \), where otherwise the Higgs mass would require fine-tuning of parameters. There are also attractive practical features: SUSY-GUT models can be calculated perturbatively and can be tested experimentally at supercolliders, where SUSY partners can be produced and studied. Philosophically, SUSY is the last possible symmetry of the \( S \)-matrix, and there is a predisposition to believe that anything not forbidden is compulsory.

Phenomenological interest has focussed mainly on the MSSM, which introduces just one spartner for each SM particle. The gauge symmetry is \( \text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y \); the corresponding spin-1 gauge bosons \( g, W, Z, \gamma \) have spin-1/2 “gaugino” partners \( \tilde{g}, \tilde{W}, \tilde{Z}, \tilde{\gamma} \). The three generations of spin-1/2 quarks \( q \) and leptons \( \ell \) have spin-0 squark and sleptons partners \( \tilde{q} \) and \( \tilde{\ell} \); the chiral states \( f_L \) and \( f_R \) of any given fermion \( f \) have distinct sfermion partners \( \tilde{f}_L \) and \( \tilde{f}_R \), respectively (that can however mix). For anomaly cancellation the single Higgs doublet must be replaced by two doublets \( H_1 \) and \( H_2 \) that have higgsino partners \( \tilde{H}_1 \) and \( \tilde{H}_2 \). The MSSM also conserves a multiplicative \( R \)-parity, defined by

\[
R = (-1)^{2S+L+3B}
\]

where \( S, L, B \) are spin, lepton number and baryon number. The normal particles of the SM all have \( R = +1 \); their spartners which differ simply by 1/2 unit of \( S \), therefore have \( R = -1 \). \( R \)-conservation has important physical implications:

(a) sparticles must be produced in pairs,

(b) heavy sparticles decay to lighter sparticles,

(c) the lightest sparticle (LSP) is stable.

If this LSP has zero charge and only interacts weakly, as seems likely since it has defied detection so far, it will carry off undetected energy and momentum in high-energy collisions (providing possible signatures for sparticle production) and will offer a possible source of cosmological dark matter.

In addition to the more general motivations above, there are also several significant phenomenological motivations for SUSY.

(a) Grand Unified Theories (GUTs) with purely SM particle content do not predict a satisfactory convergence of the gauge couplings at some high GUT scale \( M_G \), but convergence can be achieved if SUSY partners are added (see Section 2).

(b) Starting from equal \( b \) and \( \tau \) Yukawa couplings at the GUT scale \( M_G \), the physical masses can be correctly predicted when the evolution equations include SUSY partners, but not with the SM alone (see Section 2).
2 SUSY-GUT models

As the renormalization mass scale \( \mu \) is changed, the evolution of couplings is governed by the Renormalization Group Equations (RGE). For the gauge group \( SU(3) \times SU(2) \times U(1) \), with corresponding gauge couplings \( g_3(= g_s), g_2(= g), g_1(= \sqrt{5/3} g') \), the RGE can be written in terms of the dimensionless variable \( t = \ln(\mu/M_G) \):

\[
\frac{dg_i}{dt} = \frac{g_i}{16\pi^2} \left[ b_i g_i^2 + \frac{1}{16\pi^2} \left( \sum_{j=1}^{3} b_{ij} g_i^2 g_j^2 - \sum a_{ij} g_i^2 \lambda_j^2 \right) \right]
\]

The first term on the right is the one-loop approximation; the second and third terms contain two-loop effects, involving other gauge couplings \( g_j \) and Yukawa couplings \( \lambda_j \). The coefficients \( b_i, b_{ij} \) and \( a_{ij} \) are determined at given scale \( \mu \) by the content of active particles (those with mass \( < \mu \)). If there are no thresholds (i.e. no changes of particle content) between \( \mu \) and \( M_G \), then the coefficients are constants through this range and the one-loop solution is

\[
\alpha_i^{-1}(\mu) = \alpha_i^{-1}(M_G) - t b_i/(2\pi),
\]

where \( \alpha_i = g_i^2/(4\pi) \); thus \( \alpha_i^{-1} \) evolves linearly with \( \ln \mu \) at one-loop order. If there are no new physics thresholds between \( \mu = M_Z \simeq m_t \) and \( M_G \) (i.e. just a “desert” as in the basic SM) then equations of this kind should evolve the observed couplings at the electroweak scale:

\[
\begin{align*}
\alpha_1(M_Z)^{-1} &= 58.89 \pm 0.11, \\
\alpha_2(M_Z)^{-1} &= 29.75 \pm 0.11, \\
\alpha_3(M_Z)^{-1} &= 0.118 \pm 0.007,
\end{align*}
\]

to converge to a common value at some large scale. Figure 1(a) shows that such a SM extrapolation does NOT converge; this figure actually includes two-loop effects but the evolution is still approximately linear versus \( \ln \mu \), as at one-loop order. GUTs do not work, if we assume just SM particles plus a desert up to \( M_G \).

But if we increase the particle content to include the minimum number of SUSY particles, with a threshold not too far above \( M_Z \), then GUT-type convergence can
happen. Figure 1(b) shows an example with SUSY threshold \( M_{\text{SUSY}} = 1 \text{ TeV} \).\(^4\) SUSY-GUTs are plainly more successful; the evolved couplings are consistent with a common intersection at \( M_G \sim 10^{16} \text{ GeV} \). In fact a precise single-point intersection is not strictly necessary, since the exotic GUT gauge, fermion and scalar particles do not have to be precisely degenerate; we may therefore have several non-degenerate thresholds near \( M_G \), to be passed through on the way to GUT unification.

![Fig. 1. Gauge coupling evolution: (a) in the SM; (b) in a SUSY-GUT example.](image)

The Yukawa couplings also evolve. The evolution equations for \( \lambda_t \) and \( \lambda_b/\lambda_\tau \) are

\[
\frac{d \lambda_t}{dt} = \frac{\lambda_t}{16\pi^2} \left[ - \sum c_i g_i^2 + 6\lambda_t^2 + \lambda_b^2 + \text{2-loop terms} \right],
\]

with \( c_1 = 13/15, c_2 = 3, c_3 = 16/3 \), and

\[
\frac{d(\lambda_b/\lambda_\tau)}{dt} = \frac{(\lambda_b/\lambda_\tau)}{16\pi^2} \left[ - \sum d_i g_i^2 + \lambda_t^2 + 3\lambda_b^2 - 3\lambda_\tau^2 + \text{2-loop terms} \right],
\]

with \( d_1 = -4/3, d_2 = 0, d_3 = 16/3 \). The low-energy values at \( \mu = m_t \) are

\[
\lambda_b(m_t) = \frac{\sqrt{2} m_b(m_b)}{\eta_b v \cos \beta}, \quad \lambda_\tau(m_t) = \frac{\sqrt{2} m_\tau(m_\tau)}{\eta_\tau v \cos \beta}, \quad \lambda_t(m_t) = \frac{\sqrt{2} m_t(m_t)}{v \sin \beta},
\]

where \( \eta_f = m_f(m_f)/m_f(m_t) \) gives the running of the masses below \( \mu = m_t \), obtained from 3-loop QCD and 1-loop QED evolution. The \( \eta_f \) values depend on the value of \( \alpha_3(M_Z) \); for \( \alpha_3(M_Z) = 0.118 \), \( \eta_b = 1.5 \), \( \eta_c = 2.1 \), \( \eta_s = 2.4 \). The running mass values are \( m_\tau(m_\tau) = 1.777 \text{ GeV} \) and \( m_b(m_b) = 4.25 \pm 0.15 \text{ GeV} \). The denominator factors in Eq. (9) arise from the two Higgs vevs \( v_1 = v \cos \beta \) and \( v_2 = v \sin \beta \); they are related to the SM vev \( v = 246 \text{ GeV} \) by \( v_1^2 + v_2^2 = v^2 \), while tan \( \beta = v_2/v_1 \) measures their ratio.

It is frequently assumed that the \( b \)-quark and \( \tau \)-lepton Yukawa couplings are equal at the GUT scale:\(^5\)

\[
\lambda_b(M_G) = \lambda_\tau(M_G).
\]

Figure 2 illustrates the running of \( \lambda_t, \lambda_b \) and \( \lambda_\tau \), obtained from solutions to the RGEs with the appropriate low-energy boundary conditions and the GUT-scale condition of Eq. (10). Note that \( \lambda_t(M_G) \) must be large in order to satisfy the boundary condition \( m_b(m_b) = 4.25 \pm 0.15 \).

![Fig. 2. The running of \( \lambda_t, \lambda_b \) and \( \lambda_\tau \) from low energies to the GUT scale.](image)

As \( \mu \rightarrow m_t \), \( \lambda_t \) rapidly approaches a fixed point.\(^6\) The approximate fixed-point solution for \( m_t \) is

\[
- \sum c_i g_i^2 + 6\lambda_t^2 + \lambda_b^2 = 0.
\]

---

\(^1\) The Yukawa couplings also evolve. The evolution equations for \( \lambda_t \) and \( \lambda_b/\lambda_\tau \) are

\[
\frac{d \lambda_t}{dt} = \frac{\lambda_t}{16\pi^2} \left[ - \sum c_i g_i^2 + 6\lambda_t^2 + \lambda_b^2 + \text{2-loop terms} \right],
\]

with \( c_1 = 13/15, c_2 = 3, c_3 = 16/3 \), and

\[
\frac{d(\lambda_b/\lambda_\tau)}{dt} = \frac{(\lambda_b/\lambda_\tau)}{16\pi^2} \left[ - \sum d_i g_i^2 + \lambda_t^2 + 3\lambda_b^2 - 3\lambda_\tau^2 + \text{2-loop terms} \right],
\]

with \( d_1 = -4/3, d_2 = 0, d_3 = 16/3 \). The low-energy values at \( \mu = m_t \) are

\[
\lambda_b(m_t) = \frac{\sqrt{2} m_b(m_b)}{\eta_b v \cos \beta}, \quad \lambda_\tau(m_t) = \frac{\sqrt{2} m_\tau(m_\tau)}{\eta_\tau v \cos \beta}, \quad \lambda_t(m_t) = \frac{\sqrt{2} m_t(m_t)}{v \sin \beta},
\]

where \( \eta_f = m_f(m_f)/m_f(m_t) \) gives the running of the masses below \( \mu = m_t \), obtained from 3-loop QCD and 1-loop QED evolution. The \( \eta_f \) values depend on the value of \( \alpha_3(M_Z) \); for \( \alpha_3(M_Z) = 0.118 \), \( \eta_b = 1.5 \), \( \eta_c = 2.1 \), \( \eta_s = 2.4 \). The running mass values are \( m_\tau(m_\tau) = 1.777 \text{ GeV} \) and \( m_b(m_b) = 4.25 \pm 0.15 \text{ GeV} \). The denominator factors in Eq. (9) arise from the two Higgs vevs \( v_1 = v \cos \beta \) and \( v_2 = v \sin \beta \); they are related to the SM vev \( v = 246 \text{ GeV} \) by \( v_1^2 + v_2^2 = v^2 \), while tan \( \beta = v_2/v_1 \) measures their ratio.

It is frequently assumed that the \( b \)-quark and \( \tau \)-lepton Yukawa couplings are equal at the GUT scale:\(^5\)

\[
\lambda_b(M_G) = \lambda_\tau(M_G).
\]

Figure 2 illustrates the running of \( \lambda_t, \lambda_b \) and \( \lambda_\tau \), obtained from solutions to the RGEs with the appropriate low-energy boundary conditions and the GUT-scale condition of Eq. (10). Note that \( \lambda_t(M_G) \) must be large in order to satisfy the boundary condition \( m_b(m_b) = 4.25 \pm 0.15 \).

![Fig. 2. The running of \( \lambda_t, \lambda_b \) and \( \lambda_\tau \) from low energies to the GUT scale.](image)

As \( \mu \rightarrow m_t \), \( \lambda_t \) rapidly approaches a fixed point.\(^6\) The approximate fixed-point solution for \( m_t \) is

\[
- \sum c_i g_i^2 + 6\lambda_t^2 + \lambda_b^2 = 0.
\]
Neglecting $g_1$, $g_2$ and $\lambda_b$, $m_t$ is predicted in terms of $\alpha_s(m_t)$ and $\beta$:

$$m_t(m_t) \approx \frac{4}{3} \sqrt{2 \pi \alpha_s(m_t)} \frac{v}{\sqrt{2}} \sin \beta \approx (200 \text{ GeV}) \frac{\tan \beta}{\sqrt{1 + \tan^2 \beta}} . \quad (12)$$

Thus the natural scale of the top-quark mass is large in SUSY-GUT models. Note that the propagator-pole mass is related to this running mass by

$$m_t(\text{pole}) = m_t(m_t) \left[ 1 + \frac{4}{3\pi} \alpha_s(m_t) + \cdots \right] . \quad (13)$$

An exact numerical solution for the relation between $m_t$ and $\tan \beta$, obtained from the 2-loop RGEs for $\lambda_t$ and $\lambda_b/\lambda_t$, is shown in Fig. 3 taking $M_{\text{SUSY}} = m_t$. At large $\tan \beta$, $\lambda_b$ becomes large and the above fixed-point solution no longer applies. In fact, the solutions becomes non-perturbative at large $\tan \beta$ and we impose the perturbative requirements $\lambda_t(M_G) \leq 3.3$, $\lambda_b(M_G) \leq 3.1$, based on the requirement that (2-loop)/(1-loop)$\leq 1/4$ giving $\tan \beta \lesssim 65$. At large $\tan \beta$ there is the possibility of $\lambda_t = \lambda_b = \lambda_\tau$ unification. For most $m_t$ values there are two possible solutions for $\tan \beta$; the lower solution is controlled by the $\lambda_t$ fixed point, following Eqs. (12),(13):

$$\sin \beta \simeq m_t(\text{pole})/(200 \text{ GeV}) . \quad (14)$$

An upper limit $m_t(\text{pole}) \lesssim 200 \text{ GeV}$ is found with the RGE solutions.

Fig. 3. Contours of constant $m_b(m_b)$ in the $(m_t(m_t), \tan \beta)$ plane.$^{14}$

Figure 4 shows the dependence of $\lambda_t(M_G)$ on $\alpha_3(M_Z)$. For $\lambda_t$ to remain perturbative, an upper limit $\alpha_3(M_Z) \lesssim 0.125$ is necessary.

Fig. 4. Dependence of $\lambda_t$ at the GUT scale on $\alpha_3(M_Z).^{14}$

Specific GUT models also make predictions for CKM matrix elements. For example, several models$^{16,19}$ give the GUT-scale relation

$$|V_{cb}(\text{GUT})| = \sqrt{\lambda_c(\text{GUT})/\lambda_t(\text{GUT})} . \quad (15)$$

The relevant RGEs are

$$\frac{d|V_{cb}|}{dt} = - \frac{|V_{cb}|}{16\pi^2} \left[ \lambda_t^2 + \lambda_b^2 + \text{2-loop} \right] , \quad (16)$$

$$\frac{d(\lambda_c/\lambda_t)}{dt} = - \frac{(\lambda_c/\lambda_t)}{16\pi^2} \left[ 3\lambda_t^2 + \lambda_b^2 + \text{2-loop} \right] , \quad (17)$$
in addition to Eqs. (7) and (8). Starting from boundary conditions on $m_c$ and $|V_{cb}|$ at scale $\mu = m_t$, the equations can be integrated up to $M_G$ and checked to see if the above GUT-scale constraint is satisfied. The low-energy boundary conditions are

$$0.032 \leq |V_{cb}(m_t)| \leq 0.054, \quad 1.19 \leq m_c(m_t) \leq 1.35 \text{GeV}. \quad (18)$$

The resulting $|V_{cb}|$ solutions at the 2-loop level are shown in Fig. 5. The contours of $m_b(m_b) = 4.1$ and $4.4 \text{ GeV}$, which satisfy $\lambda_b(M_G)/\lambda_t(M_G) = 1$, are also shown in Fig. 5(a). The shaded region in Fig. 5(a) denotes the solutions that satisfy both sets of GUT-scale constraints. A lower limit $m_t \geq 155 \text{ GeV}$ can be inferred, based on values $m_c = 1.19$ and $\alpha_3(M_Z) = 0.118$ used in this illustration; with $\alpha_3(M_Z) = 0.118$ instead, $m_t$ can be as low as $120 \text{ GeV}$ with $|V_{cb}| = 0.054$. One GUT “texture” that leads to the above $|V_{cb}|$ GUT prediction is given by the following up-quark, down-quark and lepton mass-matrix structure at $M_G$.

$$U = \begin{pmatrix} 0 & C & 0 \\ C & 0 & B \\ 0 & B & A \end{pmatrix}, \quad D = \begin{pmatrix} 0 & F e^{i\phi} & 0 \\ F e^{-i\phi} & E & 0 \\ 0 & 0 & D \end{pmatrix}, \quad E = \begin{pmatrix} 0 & F & 0 \\ F & E & 0 \\ 0 & 0 & D \end{pmatrix}. \quad (19)$$

Fig. 5. Contours of constant $m_b(m_b)$ for $\lambda_b/\lambda_t = 1$ at $\mu = M_G$ and contours of constant $|V_{cb}(m_t)|$, (a) in the $(m_t(m_t), \sin \beta)$ plane and (b) in the $(m_t(m_t), \tan \beta)$ plane.

### 3 Experimental Signatures for SUSY

Experimental evidence for SUSY could come in various forms, for example

(a) discovery of one or more superpartners,

(b) discovery of a light neutral Higgs bosons with non-SM properties and/or a charged Higgs boson,

(c) discovery of $p \rightarrow K \nu$ decay: the present lifetime limit is $10^{32}$ years but Super-Kamiokande will be sensitive up to $10^{34}$ years,

(d) discovery that dark matter is made of heavy ($\lesssim 100 \text{ GeV}$) neutral particles.

GUTs are essential for SUSY phenomenology; without GUTs there would be far too many free parameters. A minimal set of GUT parameters with soft SUSY breaking consists of the gauge and Yukawa couplings $g_i$ and $\lambda_i$, the Higgs mixing mass $\mu$, the common gaugino mass at the GUT scale $m_{1/2}$, the common scalar mass at the GUT scale $m_0$, and two parameters $A, B$ that give trilinear and bilinear scalar couplings. At the weak scale, the gauge couplings are experimentally determined. The Higgs potential depends upon $m_0, \mu, B$ (at tree level) and $m_{1/2}, A, \lambda_t, \lambda_b$ (at one loop). After minimizing the Higgs potential and putting in the measured $Z$ and fermion masses, there remain 5 independent parameters that can be taken as
$m_t, \tan \beta, m_0, m_{1/2}, A$, though other independent parameter sets are often used for specific purposes.

The SUSY particle spectrum consists of Higgs bosons ($h, H, A, H^\pm$), gluinos ($\tilde{g}$), squarks ($\tilde{q}$), sleptons ($\tilde{\ell}$), charginos ($\tilde{W}^\pm_1, \tilde{W}^\pm_2, i = 1, 2$; mixtures of winos and charged higgsinos), neutralinos ($\tilde{Z}_j, j = 1, 2, 3, 4$; mixtures of zinos, photinos and neutral higgsinos). An alternate notation is $\tilde{\chi}_i^\pm$ for $\tilde{W}^\pm_i$ and $\tilde{\chi}_j^0$ for $\tilde{Z}_j$. The evolution of the SUSY mass spectrum from the GUT scale is illustrated in Fig. 6. The running masses are plotted versus $\mu$ and the physical value occurs approximately where the running mass $m = m(\mu)$ intersects the curve $m = \mu$. In the case of the Higgs scalar $H_2$, the mass-square becomes negative at low $\mu$ due to coupling to top; in this region we have actually plotted $-|m(\mu)|$. Negative mass-square parameter is essential for spontaneous symmetry-breaking, so this feature of SUSY-GUTs is desirable; it is achieved by radiative effects. The running masses for the gauginos $\tilde{g}, \tilde{W}, \tilde{B}$ are given by

$$M_i(\mu) = m_{1/2} \frac{\alpha_i(\mu)}{\alpha_i(M_G)},$$

where $i$ labels the corresponding gauge symmetry; this applies before we add mixing with higgsinos to obtain the chargino and neutralino mass eigenstates. In the example of Fig. 6 the squarks are heavier than the gluinos, but the opposite ordering $m_\tilde{q} < m_\tilde{g}$ is possible in other scenarios. Sleptons, neutralinos and charginos are lighter than both squarks and gluinos in general. Note that the usual soft SUSY-breaking mechanisms preserve the gauge coupling relations (unification) at $M_G$.

In order that SUSY cancellations shall take effect at low mass scales as required, the SUSY mass parameters are expected to be bounded by

$$m_\tilde{g}, m_\tilde{q}, |\mu|, m_A \lesssim 1-2 \text{ TeV}. \quad (21)$$

The other parameter $\tan \beta$ is effectively bounded by

$$1 \lesssim \tan \beta \lesssim 65, \quad (22)$$

where the lower bound arises from consistency in GUT models and the upper bound is the perturbative limit. Proton decay gives the constraint $\tan \beta < 85$.

Fig. 6. Representative RGE results for spartner masses.

At LEP I, sufficiently light SUSY particles would be produced through their gauge couplings to the $Z$. Direct searches for SUSY particles at LEP give mass lower bounds

$$m_\tilde{q}, m_\tilde{t}, m_\tilde{\tau}, m_{\tilde{W}_1} \gtrsim 40-45 \text{ GeV}. \quad (23)$$

The limitation of LEP is its relatively low CM energy.

Hadron colliders can explore much higher energy ranges. For $m_\tilde{q} = m_\tilde{t}$, about 100 events would be expected for each of the channels $\tilde{g}\tilde{q}$ and $\tilde{q}\tilde{g}$ at mass 200 GeV so the Tevatron clearly reaches well beyond the LEP range.
The most distinctive SUSY signature is the missing energy and momentum carried off by the LSP, usually assumed to be the lightest neutralino $\tilde{Z}_1$, which occurs in all SUSY decay chains with $R$-parity conservation. At hadron colliders it is only possible to do book-keeping on the missing transverse momentum denoted $p_T$. The missing momenta of both LSPs are added vectorially in $p_T$. The LSP momenta and hence the magnitude of $p_T$ depend on the decay chains.

If squarks and gluinos are rather light ($m_{\tilde{g}}, m_{\tilde{q}} \lesssim 50$ GeV), their dominant decay mechanisms are direct strong decays or decays to the LSP:

\[
\begin{align*}
\tilde{g} &\rightarrow q\tilde{g} \quad \text{if } m_{\tilde{g}} < m_{\tilde{q}}, \\
\tilde{g} &\rightarrow q\bar{q}\tilde{Z}_1 \quad \text{if } m_{\tilde{q}} < m_{\tilde{g}}.
\end{align*}
\]  

In such cases the LSPs carry a substantial fraction of the available energy and $p_T$ is correspondingly large. Assuming such decays and small LSP mass, the present 90% CL experimental bounds from UA1 and UA2 at the CERN $p\bar{p}$ collider ($\sqrt{s} = 640$ GeV) and from CDF at the Tevatron ($\sqrt{s} = 1.8$ TeV) are:

| Experiment | $m_{\tilde{g}}$ | $m_{\tilde{q}}$ |
|------------|----------------|----------------|
| UA1 (1987) | > 53 GeV       | > 45 GeV       |
| UA2 (1990) | > 79           | > 74           |
| CDF (1992) | > 141          | > 126          |

The limits become more stringent if $m_{\tilde{g}}$ and $m_{\tilde{q}}$ are assumed to be comparable.

For heavier gluinos and squarks, many new decay channels are open, such as:

\[
\begin{align*}
\tilde{g} &\rightarrow q\bar{q}\tilde{Z}_i \ (i = 1, 2, 3, 4), \ q\bar{q}\tilde{W}_j \ (j = 1, 2), \ g\tilde{Z}_1, \\
\tilde{q}_L &\rightarrow q\tilde{Z}_i \ (i = 1, 2, 3, 4), \ q\tilde{W}_j \ (j = 1, 2), \\
\tilde{q}_R &\rightarrow q\tilde{Z}_i \ (i = 1, 2, 3, 4).
\end{align*}
\]  

Some decays go via loops (e.g. $\tilde{g} \rightarrow g\tilde{Z}_1$); we have not attempted an exhaustive listing here. Figure 7 shows how gluino-to-heavy-gaugino branching fractions increase with $m_{\tilde{g}}$ in a particular example (with $m_{\tilde{g}} < m_{\tilde{q}}$).

Fig. 7. Example of gluino decay branchings versus mass.

The heavier gauginos then decay too:

\[
\begin{align*}
\tilde{W}_j &\rightarrow Z\tilde{W}_k, \ W\tilde{Z}_i, \ H^0_k \tilde{W}_k, \ H^\pm \tilde{Z}_i, \ f\tilde{f} , \\
\tilde{Z}_i &\rightarrow Z\tilde{Z}_k, \ W\tilde{W}_j, \ H^0_k \tilde{Z}_k, \ H^\pm \tilde{W}_k, \ f\tilde{f}'.
\end{align*}
\]  

Here it is understood that final $W$ or $Z$ may be off-shell and materialize as fermion-antifermion pairs; also $Z$ may be replaced by $\gamma$. To combine the complicated production and cascade possibilities systematically, all these channels have been incorporated in the ISAJET 7.0 Monte Carlo package called ISASUSY.
These multibranch cascade decays lead to higher-multiplicity final states in which the LSPs $\tilde{Z}_1$ carry a much smaller share of the available energy, so $p_T$ is smaller and less distinctive, making detection via $p_T$ more difficult. Remember that leptonic $W$ or $Z$ decays, $\tau$ decays, plus semileptonic $b$ and $c$ decays, all give background events with genuine $p_T$; measurement uncertainties also contribute fake $p_T$ backgrounds. Experimental bounds therefore become weaker when we take account of cascade decays. The CDF 90% CL limits on the $(m_{\tilde{g}}, m_{\tilde{q}})$ masses quoted above are reduced by 10–30 GeV when cascade decays are taken into account.

Cascade decays also present new opportunities for SUSY detection. Same-sign dileptons (SSD) are a very interesting signal, which arises naturally from $\tilde{g}\tilde{g}$ and $\tilde{g}\tilde{q}$ decays because of the Majorana character of gluinos, with very little background. Figure 8 gives an example of this signal. Eqs. (25)–(29) show how a heavy gluino or squark can decay to a chargino $\tilde{W}_j$ and hence, via a real or virtual $W$, to an isolated charged lepton. For such squark pair decays the two charginos — and hence the two leptons — are constrained to have opposite signs, but if a gluino is present it can decay equally into either sign of chargino and lepton because it is a Majorana fermion. Hence $\tilde{g}\tilde{g}$ or $\tilde{g}\tilde{q}$ systems can decay to isolated SSD plus jets plus $p_T$. Cascade decays of $\tilde{q}$ via the heavier neutralinos $\tilde{Z}_i$ offer similar possibilities for SSD, since the $\tilde{Z}_i$ are also Majorana fermions. Cross sections for the Tevatron are shown in Fig. 9.

Fig. 8. Example of same-sign dilepton appearance in gluino-pair decay.

Fig. 9. Same-sign dilepton signals at the Tevatron.

Genuinely isolated SSD backgrounds come from the production of $WZ$ or $Wt\bar{t}$ or $W^+W^+$ (e.g. $uu \rightarrow ddW^+W^+$ by gluon exchange), with cross sections of order $\alpha_s^2$ or $\alpha_s^\alpha_s \alpha_s^3$ or $\alpha_s^2 \alpha_s \alpha_s^3$ compared to $\alpha_s^3$ for gluino pair production, so we expect to control them with suitable cuts. Very large $b\bar{b}$ production gives SSD via semileptonic $b$-decays plus $B-\bar{B}$ mixing, and also via combined $b \rightarrow c \rightarrow s\ell^+\nu$ and $b \rightarrow \bar{c}\ell^+\nu$ decays, but both leptons are produced in jets and can be suppressed by isolation cuts. Also $t\bar{t}$ gives SSD via $t \rightarrow b\ell^+\nu$ and $\bar{t} \rightarrow \bar{b} \rightarrow \bar{c}\ell^+\nu$, but the latter lepton is non-isolated. So SSD provide a promising SUSY signature.

Gluino production rates at SSC/LHC are much higher than at the Tevatron. At $\sqrt{s} = 40$ TeV, the cross section is

$$\sigma(\tilde{g}\tilde{g}) = 10^4, 700, 6 \text{ fb} \quad \text{for } m_{\tilde{g}} = 0.3, 1, 2 \text{ TeV}.$$  \hspace{1cm} (30)

Many different SUSY signals have been evaluated, including $p_T + n$ jets, $p_T +$ SSD, $p_T + n$ isolated leptons, $p_T +$ one isolated lepton + $Z$, $p_T + Z$, $p_T + Z + Z$. SSC cross sections for some of these signals from $\tilde{g}\tilde{g}$ production are shown versus $m_{\tilde{g}}$ in Fig. 10 (for two scenarios, after various cuts); the labels 3,4,5 refer to numbers of isolated leptons.

Fig. 10. SSC cross sections for various SUSY signals, after cuts.
Heavy gluinos can also decay copiously to t-quarks:

\[ \tilde{g} \to t\tilde{t}, tb\tilde{W}^-, bt\tilde{W}^+. \]  

(31)

\( t \to bW \) decay then leads to multiple \( W \) production. For example, for a gluino of mass 1.5 TeV, the \( \tilde{g} \to W, WW, WWZ, WWWW \) branching fractions are typically of order 30%, 30%, 6%, 6%, respectively. For \( m_{\tilde{g}} \sim 1 \) TeV the SUSY rate for 4\( W \) production can greatly exceed the dominant SM \( 4t \to 4W \) mode, offering yet another signal for SUSY.

To summarize this section:

(a) Light SUSY particle searches are based largely on \( \not{p}_T \) signals. But for \( m_{\tilde{g}}, m_{\tilde{q}} > 50 \) GeV cascade decays become important; these cascades both weaken the simple \( \not{p}_T \) signals and provide new signals such as same-sign dileptons.

(b) For even heavier squarks and gluinos, the cascade decays dominate completely and provide further exotic (multi-\( W, Z \) and multi-lepton) signatures.

(c) Gluinos and squarks in the expected mass range of Eq. (21) will not escape detection.

4 MSSM Higgs Phenomenology

In minimal SUSY, two Higgs doublets \( H_1 \) and \( H_2 \) are needed to cancel anomalies and at the same time give masses to both up- and down-type quarks. Their vevs are \( v_1 = v \cos \beta \) and \( v_2 = v \sin \beta \). There are 5 physical scalar states: \( h \) and \( H \) (neutral CP-even with \( m_h < m_H \)), \( A \) (neutral CP-odd) and \( H^\pm \). At tree level the scalar masses and couplings and an \( h-H \) mixing angle \( \alpha \) are all determined by two parameters, conveniently chosen to be \( m_A \) and \( \tan \beta \). At tree level the masses obey

\[ m_h \leq M_Z, m_A, m_H \geq M_Z, m_A, m_{H^\pm} \geq M_W, m_A. \]

Radiative corrections are significant, however. The most important new parameters entering here are the \( t \) and \( \bar{t} \) masses; we neglect for simplicity some other parameters related to squark mixing. One-loop corrections give \( h \) and \( H \) mass shifts of order \( \delta m^2 \sim G_F m_t^4 \ln(m_t/m_t) \), arising from incomplete cancellation of \( t \) and \( \bar{t} \) loops.

The \( h \) and \( H \) mass bounds get shifted up and for the typical case \( m_t = 150 \) GeV, \( m_t = 1 \) TeV we get

\[ m_h < 116 \text{ GeV} < m_H. \]  

(32)

There are also corrections to cubic \( hAA, HAA, Hhh \) couplings, to \( h-H \) mixing, and smaller corrections to the \( H^\pm \) mass. Figure 11 illustrates the dependence of \( m_h \) and \( m_H \) on \( m_A \) and \( \tan \beta \), for two different values of \( m_t \) (with \( m_t = 1 \) TeV still). We shall assume \( \tan \beta \) obeys the GUT constraints \( 1 \leq \tan \beta \leq 65 \) of Eq. (22).

At LEP I, the ALEPH, DELPHI, L3 and OPAL collaborations have all searched for the processes

\[ e^+e^- \to Z \to Z^*h, Ah, \]  

(33)
with $Z^* \rightarrow \ell\ell, \nu\nu, jj$ and $h, A \rightarrow \tau\tau, jj$ decay modes. The $ZZh$ and $ZAh$ production vertices have complementary coupling-strength factors $\sin(\beta - \alpha)$ and $\cos(\beta - \alpha)$, respectively, helping to give good coverage. The absence of signals excludes regions of the $(m_A, \tan \beta)$ plane; Fig. 12 shows typical boundaries for various $m_t$ values, deduced from ALEPH results. These results imply lower bounds

$$m_h, m_A \gtrsim 20-45 \text{ GeV (depending on } \tan \beta).$$

Null searches for $e^+e^- \rightarrow H^+H^-$ exclude a region with $\tan \beta < 1$.

Fig. 11. Contours of $h$ and $H$ masses in the $(m_A, \tan \beta)$ plane for (a) $m_t = 150$ GeV, (b) $m_t = 200$ GeV, with $m_{\tilde{t}} = 1$ TeV.

Fig. 12. Limits from ALEPH searches for (a) $Z \rightarrow Z^*h$ and (b) $Z \rightarrow Ah$ at LEP I, for various $m_t$ values with $m_{\tilde{t}} = 1$ TeV.

LEP II will have higher energy and greater reach. Figure 13 shows approximate discovery limits in the $(m_A, \tan \beta)$ plane for various $m_t$ values, based on projected searches for $e^+e^- \rightarrow ZH \rightarrow \ell\ell jj, \nu\nu jj, jjjj$ and for $e^+e^- \rightarrow (Zh, Ah) \rightarrow \tau\tau jj$, assuming energy $\sqrt{s} = 200$ GeV and luminosity $\mathcal{L} = 500 \text{ pb}^{-1}$. $H^\pm$ searches will not extend this reach.

Fig. 13. Projected limits for various LEP II searches, assuming $\sqrt{s} = 200$ GeV and $\mathcal{L} = 500 \text{ pb}^{-1}$.

Searches for neutral scalars at SSC and LHC will primarily be analogous to SM Higgs searches:

(a) untagged $\gamma\gamma$ signals from $pp \rightarrow (h, H, A) \rightarrow \gamma\gamma$ via top quark loops (Fig. 14);

(b) tagged $\gamma\gamma$ signals from $pp \rightarrow (h, H, A) \rightarrow \gamma\gamma$ plus associated $t\bar{t}$ or $W$, permitting lepton tagging via $t \rightarrow W \rightarrow \ell \nu$ or $W \rightarrow \ell \nu$ decays (Fig. 15);

(c) four-lepton signals from $pp \rightarrow (h, H) \rightarrow ZZ$ or $Z^*Z \rightarrow \ell^+\ell^+\ell^-\ell^-$ (Fig. 16).

Though qualitatively similar to SM signals, these will generally be smaller due to the different coupling constants that depend on $\beta$ and $\alpha$.

Fig. 14. Typical diagram for untagged Higgs $\rightarrow \gamma\gamma$ signals.

Fig. 15. Typical diagrams for lepton-tagged Higgs $\rightarrow \gamma\gamma$ signals.

Fig. 16. Typical diagrams for “gold-plated” four-lepton Higgs signals.
For charged Higgs scalars, the only copious hadroproduction source appears to be top production with \( t \rightarrow bH^+ \) decay (that requires \( m_{H^\pm} < m_t - m_b \)). The subsequent \( H^+ \rightarrow cs, \nu\tau \) decays are most readily detected in the \( \tau\nu \) channel (favored for \( \tan \beta > 1 \)), with \( \tau \rightarrow \pi\nu \) decay (Fig. 17).

Fig. 17. Typical diagram for \( \tau \) signals from top decay via charged-Higgs modes.

SM \( t \)-decays give equal probabilities for \( e, \nu, \tau \) leptons via \( t \rightarrow bW \rightarrow b(e, \mu, \tau)\nu \), but the non-standard \( t \rightarrow bH^+ \rightarrow b\tau\nu \) leads to characteristic excess of \( \tau \). The strategy is to tag one top quark via standard \( t \rightarrow bW \rightarrow b\ell\nu \) decay and to study the \( \tau/\ell \) ratio in the associated top quark decay (\( \ell = e \) or \( \mu \)).

Several groups have studied the detectability of these signals at SSC/LHC, and they all reach broadly similar conclusions. Figure 18 shows typical limits of detectability for untagged and lepton-tagged \( \gamma\gamma \) signals at SSC, assuming luminosities \( \mathcal{L} = 20 \text{ fb}^{-1} \) (two years of running) and \( m_t = 150 \text{ GeV} \). Figure 19(a) shows a similar limit for the \( H \rightarrow 4\ell \) search (no \( h \rightarrow 4\ell \) signal is detectable). Figure 19(b) shows typical limits for detecting the \( t \rightarrow H^+ \rightarrow \tau\nu \) signal; here the value of \( m_t \) is critical, since only the range \( m_{H^+} < m_t - m_b \) can contribute at all. Putting all these discovery regions together with the LEP I and LEP II regions, we see that very considerable coverage of the \( (m_A, \tan \beta) \) plane can be expected — but there still remains a small inaccessible region; see Fig. 20. For \( m_t = 120 \text{ GeV} \) the inaccessible region is larger, for \( m_t = 200 \text{ GeV} \) it is smaller.

Fig. 18. Limits of detectability for \( H, h, A \gamma\gamma \) signals at the SSC, for \( \mathcal{L} = 20 \text{ pb}^{-1} \), (a) without tagging, and (b) with lepton tag.

Fig. 19. Detectability limit for (a) \( H \rightarrow 4\ell \) signals and (b) \( t \rightarrow bH^+ \rightarrow b\tau\nu \) signals at the SSC for \( \mathcal{L} = 20 \text{ pb}^{-1} \).

Fig. 20. Combined LEP and SSC discovery regions for \( m_t = 150 \text{ GeV} \) from Ref. 30; similar results are obtained by other groups.

Figure 21 shows how many of the MSSM scalars \( h, H, A, H^\pm \) would be detectable, in various regions of the \( (m_A, \tan \beta) \) plane. In many regions two or more different scalars could be discovered, but for large \( m_A \) only \( h \) would be discoverable; in the latter region, the \( h \) couplings all reduce to SM couplings, the other scalars become very heavy and approximately degenerate, and the MSSM essentially behaves like the SM.

Fig. 21. How many MSSM Higgs bosons may be discovered (from Ref. 30).

An indirect constraint on the MSSM Higgs sector is provided by the CLEO bound on \( b \rightarrow s\gamma \) decays.

\[
B(b \rightarrow s\gamma) < 5.4 \times 10^{-4} \quad (95\% \text{ C.L.})
\]
In the SM this decay proceeds via a $W$ loop process, but in models with more than one Higgs doublet there are charged Higgs contributions too (Fig. 22). In the MSSM both the $W$ and $H$ amplitudes have the same sign and the branching fraction is directly related to $m_{H^+}$ and $\tan \beta$; hence the CLEO result implies a lower bound on $m_{H^+}$ for given $\tan \beta$ [Fig. 23(a)]. It was recently pointed out\cite{37,38} that this CLEO-based constraint falls in a very interesting and sensitive region when translated to the $(m_A, \tan \beta)$ plane; see Fig. 23(b). Taken at face value, it appears to exclude a large part of the LEP II discovery region and even to exclude the otherwise inaccessible region too.

Fig. 22. $W$ and charged-Higgs loop diagrams contributing to $b \to s\gamma$ decays.

Fig. 23. (a) Lower bound on $m_{H^+}$ for given $\tan \beta$, from $b \to s\gamma$ constraint. (b) Comparison of $b \to s\gamma$ bound with other MSSM Higgs constraints in the $(m_A, \tan \beta)$ plane, for $m_t = 150$ GeV. The regions excluded by the CLEO experimental bound are to the left of the $b \to s\gamma$ curves; the curves shown are updated from Ref. 38.

It is premature however to reach firm conclusions yet. The calculations of Ref. 38 are based on the approximation of Ref. 39, but later work indicates further small corrections\cite{40}. More importantly, other SUSY loop diagrams (especially chargino loops) can give additional contributions of either sign, leading to potentially significant changes in the amplitude\cite{41,42}. However, as theoretical constraints on SUSY particles become more extensive, and as the $B(b \to s\gamma)$ bound itself becomes stronger, we may expect this approach to give a valuable constraint in the MSSM Higgs phenomenology.

Finally, what could a future $e^+e^-$ collider do? We have seen that part of the MSSM parameter space is inaccessible to LEP II. But a possible future linear collider with higher energy and luminosity could in principle cover the full parameter space. It is interesting to know what are the minimum $s$ and $L$ requirements for complete coverage, for given $m_t$. This question was answered in Ref. 43, based on the conservative assumption that only the channels $e^+e^-\to (Zh, Ah, ZH, AH)\to \tau\tau jj$ would be searched, with no special tagging. The results are shown in Fig. 24. We have estimated that including all $Z \to \ell\ell, \nu\nu, jj$ and $h, H, A \to bb, \tau\tau$ decay channels plus efficient $b$-tagging could increase the net signal $S$ by a factor 6 and the net background $B$ by a factor 4, approximately; this would increase the statistical significance $S/\sqrt{B}$ by a factor 3 and hence reduce the luminosity requirement by a factor 9 or so. In this optimistic scenario, the luminosity axis in Fig. 24 would be rescaled downward by an order of magnitude.

Fig. 24. Conservative requirements for a “no-lose” MSSM Higgs search at a future $e^+e^-$ collider. Curves of minimal $(\sqrt{s}, L)$ pairings are shown for $m_t = 120, 150, 200$ GeV; the no-lose region for $m_t = 150$ GeV is unshaded\cite{43}.

To summarize this Section:
(a) The MSSM Higgs spectrum is richer but in some ways more elusive than the SM case.

(b) At least one light scalar is expected.

(c) As $m_A \to \infty$ this light scalar behaves like the SM scalar and the others become heavy.

(d) LEP I, LEP II and SSC/LHC will give extensive but not quite complete coverage of the MSSM parameter space.

(e) For some parameter regions, several different scalars are detectable, but generally one or more remain undetectable.

(f) The $b \to s\gamma$ bound has the potential to exclude large areas of parameter space (possibly including the inaccessible region) but is presently subject to some uncertainty.

(g) A higher-energy $e^+e^-$ collider could cover the whole MSSM parameter space, discovering at least the lightest scalar $h$.

**Acknowledgements**

We thank H. Baer, M. Berger, and P. Ohmann for valuable contributions to the contents of this review. This work was supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, in part by the U.S. Department of Energy under contract no. DE-AC02-76ER00881, and in part by the Texas National Laboratory Research Commission under grant no. RGFY9273.

**References**

1. For reviews, see J. Ellis in *Ten Years of SUSY Confronting Experiment*, CERN-TH.6707/92; H. P. Nilles, Phys. Rep. **110**, 1 (1984); P. Nath, R. Arnowitt, and A. Chamseddine, *Applied N=1 Supergravity*, ICTP series in Theoretical Physics, Vol. I, World Scientific (1984); H. Haber and G. Kane, Phys. Rep. **117**, 75 (1985); X. Tata, in *The Standard Model and Beyond*, p. 304, ed. by J. E. Kim, World Scientific (1991).

2. R. Haag, J. Lopuszanski, and M. Sohnius, Nucl. Phys. **B88**, 257 (1975).

3. U. Amaldi *et al.*, Phys. Lett. **B260**, 447 (1991).

4. J. Ellis *et al.*, Phys. Lett. **B260**, 131 (1991).

5. P. Langacker and M. Luo, Phys. Rev. **D44**, 817 (1991).
6. H. Arason et al., Phys. Rev. Lett. **67**, 2933 (1991).
7. A. Giveon et al., Phys. Lett. **B271**, 138 (1991).
8. J. Hisano, H. Murayama, and T. Yanagida, Tohoku preprint TU-400 (1992); P. Nath and D. Arnowitt, NUB-TH-3056/92; J. L. Lopez et al., Phys. Lett. **B299**, 262 (1993).
9. M. Drees and M. Nojiri, Phys. Rev. **D47**, 376 (1993).
10. S. Kelley et al., Phys. Rev. **D47**, 2461 (1993); J. Ellis et al., Nucl. Phys. **B373**, 55 (1992).
11. R.G. Roberts and L. Roszkowski, RAL-93-003.
12. G.G. Ross and R.G. Roberts, Nucl. Phys. **B377**, 571 (1992), and references therein.
13. Particle Data Book, Phys. Rev. **D45**, (1992).
14. V. Barger, M.S. Berger, and P. Ohmann, Phys. Rev. **D47**, 1093 (1993), and unpublished calculations.
15. M. Chanowitz, J. Ellis, and M. K. Gaillard, Nucl. Phys. **B128**, 506 (1977); A. Buras, J. Ellis, M.K. Gaillard, and D.V. Nanopoulos, Nucl. Phys. **B135**, 66 (1978); H. Georgi and D.V. Nanopoulos, Phys. Lett. **82B**, 392 (1979); H. Georgi and C. Jarlskog, *ibid.* **86B**, 297 (1979); J.A. Harvey, P. Ramond, and D.B. Reiss, *ibid.* **92B**, 309 (1980); G. Giudice, Mod. Phys. Lett. **A7**, 2429 (1992).
16. S. Dimopoulos, L. Hall and S. Raby, Phys. Rev. Lett. **68**, 1984 (1992); Phys. Rev. **D45**, 4192 (1992).
17. B. Pendleton and G. G. Ross, Phys. Lett. **98B**, 291 (1981); C. T. Hill, Phys. Rev. **D24**, 691 (1981).
18. V. Barger, M. S. Berger, T. Han, and M. Zralek, Phys. Rev. Lett. **68**, 3394 (1992).
19. J. Harvey, P. Ramond, and D. B. Reiss, Phys. Lett. **92B**, 309 (1980); Nucl. Phys. **B199**, 223 (1982).
20. S. Kelley et al., Texas A&M preprint CTP-TAMU-16-92.
21. H. Baer and X. Tata, FSU-HEP-921222.
22. H. Baer, X. Tata, and J. Woodside, Phys. Rev. **D41**, 906 (1990); **D42**, 1568 (1991); **D45**, 142 (1992); H. Baer et al., Phys. Rev. **D46**, 303 (1992).
23. UA1 collaboration, Phys. Lett. **B198**, 261 (1987); UA2 collaboration, *ibid.* **B235**, 363 (1990); CDF collaboration, Phys. Rev. Lett. **69**, 3439 (1992).
24. H. Baer, V. Barger, D. Karatas, X. Tata, Phys. Rev. **D36**, 96 (1987).
25. H. Baer, F. Paige, S. Protopopescu, and X. Tata (unpublished).
26. V. Barger, W.-Y. Keung, R. J. N. Phillips, Phys. Rev. Lett. **55**, 166 (1985); R. Barnett, J. Gunion, H. Haber, 1988 Snowmass Summer Study, ed. by S. Jensen, World Scientific (1989).
27. H. Baer et al., FSU-HEP-901110, Proc. 1990 DPF Summer Study at Snowmass, ed. by E. Berger, World Scientific (1992).
28. V. Barger, R.J.N. Phillips and A.L. Stange, Phys. Rev. D45, 1484 (1992).
29. S. P. Li and M. Sher, Phys. Lett. B140, 339 (1984); J. F. Gunion and A. Turski, Phys. Rev. D39, 2701, D40, 2325, 2333 (1989); M. S. Berger, Phys. Rev. D41, 225 (1990); Y. Okada et al., Prog. Theor. Phys. 85, 1 (1991); Phys. Lett. B262, 54 (1991); H. Haber and R. Hempfling, Phys. Rev. Lett. 66, 1815 (1991); J. Ellis et al., Phys. Lett. B257, 83 (1991); R. Barbieri et al., Phys. Lett. B258, 167 (1991); J. Lopez and D. V. Nanopoulos, Phys. Lett. B266, 397 (1991); A. Yamada, Phys. Lett. B263, 233 (1991); M. Drees and M. N. Nojiri, Phys. Rev. D41, 225 (1990); Y. Okada et al., Phys. Rev. D41, 2482 (1991); R. Hempfling, SCIPP-91/39; M. A. Diaz and H. E. Haber, Phys. Rev. D45, 4246 (1992); D. M. Pierce, A. Papadopoulos and S. Johnson, Phys. Rev. Lett. 68, 3678 (1992).
30. V. Barger, K. Cheung, R. J. N. Phillips, and A. L. Stange, Phys. Rev. D46, 4914 (1992).
31. ALEPH Collaboration, D. Decamp et al., Phys. Lett. B246, 306 (1990); B265, 475 (1991); DELPHI Collaboration: P. Abreu et al., Phys. Lett. B245, 276 (1990); Nucl. Phys. B373, 3 (1992); L3 Collaboration: B. Adeva et al., Phys. Lett. B251, 311 (1990); B283, 454 (1992); OPAL Collaboration: M. Z. Akrawy et al., Z. Phys. C49, 1 (1991).
32. M. A. Diaz and H. E. Haber, Phys. Rev. D45, 4246 (1992).
33. H. Baer et al., Phys. Rev. D46, 1067 (1992).
34. J. Gunion et al., Phys. Rev. D46, 2040, 2052 (1992); D47, 1030 (1993).
35. Z. Kunszt and F. Zwirner, Nucl. Phys. B46, 4914 (1992).
36. CLEO Collaboration, report to the Washington APS meeting, April 1993.
37. J. L. Hewett, Phys. Rev. Lett. 70, 1045 (1993).
38. V. Barger, M. S. Berger, and R. J. N. Phillips, Phys. Rev. Lett. 70, 1368 (1993).
39. B. Grinstein, R. Springer, and M. B. Wise, Nucl. Phys. B339, 269 (1989).
40. G. Cella et al., Phys. Lett. B248, 181 (1990); M. Misiak, Phys. Lett. B269, 161 (1991), Zurich report ZH-TH-19/22 (1992); M. A. Diaz, Phys. Lett. B304, 278 (1993).
41. S. Bertolini, F. Borzumati, and A. Masiero, Nucl. Phys. B294, 321 (1987); S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. B353, 591 (1991).
42. R. Barbieri and G. F. Giudice, CERN-TH.6830/93.
43. V. Barger, K. Cheung, R. J. N. Phillips, and A. L. Stange, Phys. Rev. D47, 3041 (1993).