SEARCH FOR PHYSICS BEYOND THE STANDARD MODEL

V. Barger$^a$ and R.J.N. Phillips$^b$

$^a$Physics Department, University of Wisconsin, Madison, WI 53706, USA
$^b$Rutherford Appleton Laboratory, Chilton, Didcot, Oxon OX11 0QX, UK

Abstract

We survey some recent ideas and progress in looking for particle physics beyond the Standard Model, connected by the theme of Supersymmetry (SUSY). We review the success of SUSY-GUT models, the expected experimental signatures and present limits on SUSY partner particles, and Higgs phenomenology in the minimal SUSY model.

1 Introduction

As we stand at the beginning of 1993, the Standard Model (SM) is in excellent shape; all its predictions that have been tested have been verified to high precision. Important checks remain to be made, however: the top quark is not yet discovered, the interactions between gauge bosons are still unmeasured, and the Higgs boson remains a totally unconfirmed hypothesis. There may still be some surprises here, especially in the Higgs sector.

But even if all these checks give SM results, the apparent arbitrariness and the theoretical limitations of the SM suggest the workings of some deeper principles, embodied perhaps in Supersymmetry (SUSY) or Superstrings. Such ideas imply new physics, new particles and new interactions beyond the SM. The present review covers a small number of selected topics related to SUSY: unification of couplings in SUSY-GUT models, experimental signals from SUSY, and Higgs phenomenology in the minimal SUSY extension of the SM (MSSM).

SUSY requires each fermion to have 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 explored mass scales but could hold at a higher scale; we denote the typical scale of the superpartners as $M_{\text{SUSY}}$.

*Talk presented by V. Barger at the International Symposium on 30 Years of Neutral Currents, Santa Monica, California (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 $M_{\text{SUSY}} \lesssim O(1 \text{ TeV})$, 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 [1], and there is a predisposition to believe that anything not forbidden is compulsory.

The development of SUSY ideas in recent years is briefly as follows [2].

| Year | Event | Authors |
|------|-------|---------|
| 1966–68 | SUSY for baryon-meson system; SUSY algebra (non-relativistic) | Miyazawa |
| 1971 | Two-dimensional supersymmetry: dual string models with fermions | Ramond, Neveu + Schwarz |
| 1973 | Four-dimensional supersymmetric field theories | Wess + Zumino |
| 1974 | Absences of many divergences | Wess + Zumino, Iliopoulos, Ferrara |
| 1976 | Supergravity (local supersymmetry) | Friedman + van Nieuwenhuizen + Ferrara, Deser + Zumino |
| 1977 | Model building | Fayet + . . . |
| 1981 | Naturalness/hierarchy | Maiani, 't Hooft, Witten |
| 1981 | SUSY SU(5) GUT | Dimopoulos + Raby + Wilczek, Dimopoulos + Georgi, Sakai |
| 1984 | Strings | Green + Schwarz |
| 1987–present | SUSY RGEs + data | Amaldi et al. + . . . |
The high publication rate for SUSY papers has now reached a plateau, showing a steady continuation of interest (Fig. 1).

Fig.1. SUSY papers in SPIRES, 1980–1992, containing “supersymmetry” or “supersymmetric” in the title (figure provided by A. L. Stange).

Phenomenological interest has focussed mainly on the Minimal Supersymmetric extension of the SM (MSSM), which introduces just one spartner for each SM particle. The gauge symmetry is SU(3)_c × SU(2)_L × U(1)_Y; the corresponding spin-1 gauge bosons g, W, Z, γ have spin-1/2 “gaugino” partners ˜g, ˜W, ˜Z, ˜γ. The three generations of spin-1/2 quarks q and leptons ℓ have spin-0 squark and slepton partners ˜q and ˜ℓ; the chiral states ˜f_L and ˜f_R of any given fermion f have distinct sfermion partners ˜f_L and ˜f_R, respectively. For anomaly cancellation the single Higgs doublet must be replaced by two doublets H_1 and H_2 that have higgsino partners ˜H_1 and ˜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. R distinguishes the normal particles of the SM, which all have R = +1, from their spartners which differ simply by 1/2 unit of S and therefore have R = −1. R-conservation comes from restricting the types of coupling that are allowed. It has immediate and 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 no candidates are yet discovered, 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.

As work has proceeded, several significant phenomenological motivations for SUSY have emerged, in addition to the more general motivations above.

(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) [3, 4, 5].

(b) Starting from equal b and τ 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) [6, 7].

(c) Proton decay is too rapid in a SM GUT but can be acceptable in SUSY-GUT models where M_G is higher [8].

(d) Assuming R-parity conservation, the lightest SUSY partner (LSP) is stable and provides a plausible candidate for the origin of dark matter making Ω ∼ 1 [9, 10, 11].

(e) SUSY-GUT models lead naturally to the Higgs field developing a vacuum expectation value, when the top mass is larger than M_W [12].
2 Unification of couplings in SUSY-GUT models

The evolution of couplings, as the renormalization mass scale $\mu$ is changed, 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 $t = \ln(\mu/M_G)$:

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

(2)

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

$$\alpha_i^{-1}(\mu) = \alpha_i^{-1}(M_G) - tb_i/(2\pi)$$

(3)

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. nothing but a “desert” as in the basic SM) then equations of this kind should evolve the observed couplings at the electroweak scale $[13]$

$$\begin{align*}
\alpha_1(M_Z)^{-1} &= 58.89 \pm 0.11, \quad (4) \\
\alpha_2(M_Z)^{-1} &= 29.75 \pm 0.11, \quad (5) \\
\alpha_3(M_Z)^{-1} &= 0.118 \pm 0.007, \quad (6)
\end{align*}$$

to converge to a common value at some large scale. Figure 2 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 $M_G$.

Fig. 2. Evolution of gauge couplings in the SM.

If however 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 3 shows two examples with SUSY threshold $M_{\text{SUSY}} = m_t = 150$ GeV or $M_{\text{SUSY}} = 1$ TeV $[14]$, the threshold difference being compensated by a small change in $\alpha_2(M_Z)$. SUSY-GUTs are plainly more successful: the evolved couplings are consistent with a common intersection at $M_G \sim 10^{16}$ 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. 3. Evolution of gauge couplings in two SUSY-GUT examples, with SUSY thresholds at $M_{\text{SUSY}} = m_t = 150$ GeV or $M_{\text{SUSY}} = 1$ TeV $[14]$. 

4
The Yukawa couplings also evolve. The evolution equations for $\lambda_t$ and $\lambda_b/\lambda_t$ are

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

(7)

with $c_1 = 13/15$, $c_2 = 3$, $c_3 = 16/3$, and

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

(8)

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

$$\lambda_b(m_t) = \frac{\sqrt{2} m_b(m_b)}{\eta_b v \cos \beta}$$  

(9)

$$\lambda_\tau(m_t) = \frac{\sqrt{2} m_\tau(m_\tau)}{\eta_\tau v \cos \beta}$$  

(10)

$$\lambda_t(m_t) = \frac{\sqrt{2} m_t(m_t)}{v \sin \beta}$$  

(11)

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 $\eta$ values depend on the value of $\alpha_3(M_Z)$, as illustrated in Fig. 4. The running mass values are $m_\tau(m_\tau) = 1.777$ GeV and $m_b(m_b) = 4.25 \pm 0.15$ GeV. The denominator factors in Eqs. (9)–(11) arise from the two Higgs vevs $v_1 = v \cos \beta$ and $v_2 = v \sin \beta$; they are related to the SM vev $v = 246$ GeV by $v_1^2 + v_2^2 = v^2$, while $\tan \beta = v_2/v_1$ measures their ratio.

Fig. 4. The scaling factors $\eta_q$ for the running masses $m_q(\mu)$ [14].

A common boundary condition assumed at the GUT scale is that the $b$-quark and $\tau$-lepton Yukawa couplings are equal there [15, 16]:

$$\lambda_b(M_G) = \lambda_\tau(M_G).$$  

(12)

Figure 5 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 (12). 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. 5. The running of $\lambda_t$, $\lambda_b$ and $\lambda_\tau$ from low energies to the GUT scale [14].

As $\mu \to m_t$, $\lambda_t$ rapidly approaches a fixed point [17]. The approximate fixed-point solution for $m_t$ is

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

(13)

Neglecting $g_1$, $g_2$ and $\lambda_b$, $m_t$ is predicted in terms of $\alpha_s(m_t)$ and $\beta$ [14, 16, 18]:

$$m_t(m_t) \approx \frac{4}{3} \sqrt{2\pi \alpha_s(m_t)} \frac{v}{\sqrt{2}} \sin \beta$$  

$$\approx \frac{v}{\sqrt{2}} \sin \beta$$  

$$\approx (200 \text{ GeV}) \frac{\tan \beta}{\sqrt{1 + \tan^2 \beta}}.$$  

(14)
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
\[
m_t(\text{pole}) = m_t(m_t) \left[ 1 + \frac{4}{3\pi} \alpha_s(m_t) + \cdots \right]. \tag*{(15)}
\]

An exact numerical solution for the relation between \(m_t\) and \(\tan \beta\), obtained from the
2-loop RGEs for \(\lambda_t\) and \(\lambda_b/\lambda_r\), is shown in Fig. 6 [14] 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
\(\lambda_t(M_G) \lesssim 3.3, \lambda_b(M_G) \lesssim 3.1\), based on the requirement that \((2\text{-loop})/(1\text{-loop}) \lesssim 1/4\). At
large \(\tan \beta\) there is the possibility of \(\lambda_t = \lambda_b = \lambda_r\) unification. For most \(m_t\) values there are
two possible solutions for \(\tan \beta\); the lower solution is
\[
\sin \beta \simeq m_t(\text{pole})/200 \text{ GeV}. \tag*{(16)}
\]
An upper limit \(m_t(\text{pole}) \lesssim 200 \text{ GeV}\) is found with the RGE solutions.

Fig. 6. Contours of constant \(m_b(m_b)\) in the \((m_t(m_t), \tan \beta)\) plane [14].

Figure 7 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. 7. 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
\[
|V_{cb}(\text{GUT})| = \sqrt{\lambda_c(\text{GUT})/\lambda_t(\text{GUT})}. \tag*{(17)}
\]
The relevant RGEs are
\[
\frac{d|V_{cb}|}{dt} = -\frac{|V_{cb}|}{16\pi^2} \left[ \lambda_t^2 + \lambda_b^2 + 2\text{-loop} \right], \tag*{(18)}
\]
\[
\frac{d(\lambda_c/\lambda_t)}{dt} = -\frac{(\lambda_c/\lambda_t)}{16\pi^2} \left[ 3\lambda_t^2 + \lambda_b^2 + 2\text{-loop} \right], \tag*{(19)}
\]
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
\[
0.032 \leq |V_{cb}(m_t)| \leq 0.054, \quad 1.19 \leq m_c(m_t) \leq 1.35 \text{ GeV}. \tag*{(20)}
\]
The resulting solutions at the 2-loop level are shown by the dashed curves in Fig. 8. The
contours of \(m_b(m_b) = 4.1 \text{ and } 4.4 \text{ GeV}\), which satisfy \(\lambda_b(M_G)/\lambda_r(M_G) = 1\), are also shown.
The shaded region in Fig. 8(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 $\alpha_3(M_Z) = 0.110$ in this illustration; with $\alpha_3(M_Z) = 0.118$ instead, $m_t$ can be as low as 120 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 matrices at $M_G$ [10]:

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

Fig. 8. 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 [14, 18].

### 3 Experimental SUSY Signatures

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 boson 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 ($\lesssim 100$ GeV) neutral particles.

GUTs are essential for SUSY phenomenology, since otherwise 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 ($\tilde{\ell}^\pm$), charginos ($\tilde{W}_i^\pm$, $i = 1, 2$; mixtures of winos and charged higgsinos), neutralinos ($\tilde{\chi}_j^0, j = 1, 2, 3, 4$; mixtures of zinos, photinos and neutral higgsinos). An alternate notation is $\tilde{\chi}_i^+$ for $\tilde{W}_i^+$ and $\tilde{\chi}_j^0$ for $\tilde{Z}_j$. The evolution of the SUSY mass spectrum from the GUT scale [22, 23] is illustrated in Fig. 9. The running masses are plotted versus $\mu$ and the physical value occurs 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; here it is achieved by radiative effects. The running masses for the gauginos $\tilde{g}, \tilde{W}, \tilde{B}$ are given by

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

(22)
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. 9 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 \( M_G \).

Fig. 9. Representative RGE results for spartner masses [12].

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

\[
m_{\tilde{g}}, m_{\tilde{q}}, |\mu|, m_A \lesssim 1-2 \text{ TeV}.
\]

The other parameter \( \tan \beta \) is effectively bounded by

\[
1 \lesssim \tan \beta \lesssim 65,
\]

where the lower bound arises from consistency in GUT models and the upper bound is the perturbative limit. Proton decay gives the constraint \( \tan \beta < 85 \) [8].

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

\[
m_{\tilde{q}}, m_{\tilde{\ell}}, m_{\tilde{\nu}}, m_{\tilde{W}_1} > \sim 40-45 \text{ GeV}.
\]

The limitation of LEP is its relatively low CM energy.

Hadron colliders can explore much higher energy ranges. Figure 10 shows lowest-order gluon-gluon, gluon-quark and quark-antiquark subprocesses for SUSY particle hadroproduction. Figure 11 shows squark and gluino predictions for the Tevatron \( p-\bar{p} \) collider [21, 22], assuming degenerate masses \( m_{\tilde{q}} = m_{\tilde{g}} \) (summing \( L \) and \( R \) squarks plus antisquarks of all flavors). The right-hand vertical axis shows the number of events for the luminosity 25 pb\(^{-1}\) expected in 1993; we see that about 100 events would be expected for each of the channels \( \tilde{g}\tilde{g} \) and \( \tilde{g}\tilde{q} \) at mass 200 GeV, so the Tevatron clearly reaches well beyond the LEP range.

Fig. 10. Typical SUSY production subprocesses in hadron collisions.

Fig. 11. Tevatron cross sections for \( \tilde{g}\tilde{g} \), \( \tilde{g}\tilde{q} \) and \( \tilde{q}\tilde{q} \) production, versus squark/gluino mass.

The most distinctive signature of SUSY production is the missing energy and momentum carried off by the undetected 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 \((m_{\tilde{g}}, m_{\tilde{q}} \lesssim 50 \text{ GeV})\), their dominant decay mechanisms are direct strong decays or decays to the LSP:

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

(26)

(27)

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 \((\sqrt{s} = 640 \text{ GeV})\) and from CDF at the Tevatron \((\sqrt{s} = 1.8 \text{ TeV})\) are [23]

| 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 the squark and gluino masses are assumed to be comparable.

For heavier gluinos and squarks, many new decay channels are open, such as decays into the heavier gauginos:

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

(28)

(29)

(30)

Some decays go via loops (e.g. \(\tilde{g} \rightarrow g\tilde{Z}_1\)); we have not attempted an exhaustive listing here. Figure 12 shows how gluino-to-heavy-gaugino branching fractions increase with \(m_{\tilde{g}}\) in a particular example (with \(m_{\tilde{g}} < m_{\tilde{q}}\)) [24].

The heavier gauginos then decay too:

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

(31)

(32)

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 \(\gamma\). In practice, chargino decays are usually dominated by \(W\)-exchange transitions (Fig. 13a); neutralino decays are often dominated by sfermion exchanges (Fig. 13b) because the \(\tilde{Z}_2\tilde{Z}_1 Z\) coupling is small. 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 [25].

Fig. 12. Example of gluino decay branchings versus mass [24].

Fig. 13. Examples of (a) chargino decay by \(W\)-exchange, (b) neutralino decay by sfermion exchange.
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 $\not{p}_T$ is smaller and less distinctive (Fig. 14), making detection via $\not{p}_T$ more difficult. (Leptonic $W$ or $Z$ decays, $\tau$ decays, plus semileptonic $b$ and $c$ decays, all give background events with genuine $\not{p}_T$; measurement uncertainties also contribute fake $\not{p}_T$ backgrounds.) Experimental bounds therefore become weaker when we take account of cascade decays. Figure 15 shows typical CDF 90% CL limits in the $(m_{\tilde{g}}, m_{\tilde{q}})$ plane; the dashed curves are limits assuming only direct decays (26)–(27), while solid curves are less restrictive limits including cascade decays (28)–(32).

The cascade decays also present new opportunities for SUSY detection. Same-sign dileptons (SSD) are a very interesting signal [26], 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 16 gives an example of this signal. Eqs. (28)–(32) 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 $\not{p}_T$. The cascade decays of $\tilde{g}\tilde{g}$ 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 illustrated in Fig. 17.

![Fig. 14. Typical $\not{p}_T$ distributions from direct and cascade decays of gluino pairs at the Tevatron](image)

![Fig. 15. 1992 CDF limits in the $(m_{\tilde{g}}, m_{\tilde{q}})$ plane, with or without cascade decays, for a typical choice of parameters](image)

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 $\alpha^2_2$ or $\alpha_2\alpha^2_3$ or $\alpha^2_2\alpha^2_3$ compared to $\alpha^2_3$ for gluino pair production, so we expect to control them with suitable cuts. Very large $bb$ production gives SSD via semileptonic $b$-decays plus $B\bar{B}$ mixing, and also via combined $b \rightarrow c \rightarrow s\ell^+\nu$ and $\bar{b} \rightarrow \bar{c}\ell^+\nu$ decays, but both leptons are produced in jets and can be suppressed by stringent isolation criteria. 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.

![Fig. 16. Example of same-sign dilepton appearance in gluino-pair decay](image)

![Fig. 17. Same-sign dilepton signals at the Tevatron](image)

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} (33)

Many different SUSY signals have been evaluated, including $\not{p}_T + n$ jets, $\not{p}_T + \text{SSD}$, $\not{p}_T + n$ isolated leptons, $\not{p}_T + \text{one isolated lepton} + Z$, $\not{p}_T + Z$, $\not{p}_T + Z + Z$. SSC cross sections for
some of these signals from $\tilde{g}\tilde{g}$ production are shown versus $m_\tilde{g}$ in Fig. 18 (for two scenarios, after various cuts); the labels 3,4,5 refer to numbers of isolated leptons \cite{22}.

Heavy gluinos can also decay copiously to $t$-quarks \cite{22, 28}:

$$\tilde{g} \rightarrow t\bar{t}Z_i, tbW^-, bt\bar{W}^+ .$$

(34)

$t \rightarrow bW$ decay then leads to multiple $W$ production. For example, for a gluino of mass 1.5 TeV, the $\tilde{g} \rightarrow W, WW, WWZ, WWWW$ branching fractions are typically of order 30%, 30%, 6%, 6%, respectively. Figure 19 illustrates SSC cross sections for multi-$W$ production via gluino pair decays (assuming $m_\tilde{g} < m_\tilde{q}$). We see that for $m_\tilde{g} \sim 1$ TeV the SUSY rate for $4W$ production can greatly exceed the dominant SM $4t \rightarrow 4W$ mode, offering yet another signal for SUSY \cite{28}.

To summarize this section:

(a) Experimental SUSY particle searches have hitherto been based largely on $p_T$ signals. But for $m_\tilde{g}, m_\tilde{q} > 50$ GeV cascade decays become important; these cascades both weaken the simple $p_T$ signals and provide new signals such as same-sign dileptons, which will be pursued at the Tevatron.

(b) For even heavier squarks and gluinos, the cascade decays dominate completely and provide further exotic (multi-$W, Z$ and multi-lepton) signatures, which will be pursued at the SSC and LHC.

(c) Gluinos and squarks in the expected mass range of Eq. (23) will not escape detection.

4 SUSY 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 $v_2 = v \sin \beta$ as mentioned previously. There are therefore 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 important, however \cite{28}. The most important new parameters entering here are the $t$ and $\tilde{t}$ 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 $\tilde{t}$ loops. The $h$ and $H$ mass bounds get shifted up and for the typical case $m_t = 150$ GeV, $m_\tilde{t} = 1$ TeV we get

$$m_h < 116 \text{ GeV} < m_H.$$  

(35)
There are also corrections to cubic \( hAA, HAA, Hhh \) couplings, to \( h-H \) mixing, and smaller corrections to the \( H^\pm \) mass. Figure 20 illustrates the dependence of \( m_h \) and \( m_H \) on \( m_A \) and \( \tan \beta \), for two different values of \( m_t \) (with \( m_{\tilde{t}} = 1 \) TeV still). We shall assume \( \tan \beta \) obeys the GUT constraints \( 1 \leq \tan \beta \leq 65 \) of Eq. (24).

At LEP I, the ALEPH, DELPHI, L3 and OPAL collaborations [31] have all searched for the processes

\[
e^+e^- \to Z \to Z'h, Ah, \tag{36}
\]

with \( Z^* \to \ell\ell, \nu\nu, jj \) and \( h, A \to \tau\tau, jj \) 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. 21 shows typical boundaries for various \( m_t \) values, deduced from ALEPH results [30, 31]. These results imply lower bounds

\[
m_h, m_A \gtrsim 20-45 \text{ GeV (depending on } \tan \beta \text{).} \tag{37}
\]

Null searches for \( e^+e^- \to H^+H^- \) also exclude a region with \( \tan \beta < 1 \) [32].

Fig. 20. 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. 21. Limits from ALEPH searches for (a) \( Z \to Z'h \) and (b) \( Z \to Ah \) at LEP I, for various \( m_t \) values with \( m_{\tilde{t}} = 1 \) TeV [30, 31].

LEP II will have higher energy and greater reach. Figure 22 shows approximate discovery limits in the \((m_A, \tan \beta)\) plane for various \( m_t \) values, based on projected searches for \( e^+e^- \to ZH \to \ell\ell jj, \nu\nu jj, jjjj \) and for \( e^+e^- \to (Zh, Ah) \to \tau\tau jj \), assuming energy \( \sqrt{s} = 200 \) GeV and luminosity \( \mathcal{L} = 500 \text{ pb}^{-1} \). \( H^\pm \) searches will not extend this reach.

Fig. 22. Projected limits for various LEP II searches, assuming \( \sqrt{s} = 200 \) GeV and \( \mathcal{L} = 500 \text{ pb}^{-1} \) [30].

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

(a) untagged \( \gamma\gamma \) signals from \( pp \to (h, H, A) \to \gamma\gamma \) via top quark loops (Fig. 23);

(b) tagged \( \gamma\gamma \) signals from \( pp \to (h, H, A) \to \gamma\gamma \) plus associated \( t\bar{t} \) or \( W \), permitting lepton tagging via \( t \to W \to \ell\nu \) or \( W \to \ell\nu \) decays (Fig. 24);

(c) “gold-plated” four-lepton signals from \( pp \to (h, H) \to ZZ \) or \( Z^*Z \to \ell^+\ell^-\ell^+\ell^- \) (Fig. 25).

Though qualitatively similar to SM signals, these will generally be smaller due to the different coupling constants that depend on \( \beta \) and \( \alpha \).
For charged Higgs scalars, the only copious hadroproduction source appears to be top production with $t \to bH^+$ decay (that requires $m_{H^\pm} < m_t - m_b$). The subsequent $H^+ \to c\bar{s}, \nu\bar{\tau}$ decays are most readily detected in the $\tau\nu$ channel (favored for $\tan\beta > 1$), with $\tau \to \pi\nu$ decay (Fig. 26).

For charged Higgs scalars, the only copious hadroproduction source appears to be top production with $t \to bH^+$ decay (that requires $m_{H^\pm} < m_t - m_b$). The subsequent $H^+ \to c\bar{s}, \nu\bar{\tau}$ decays are most readily detected in the $\tau\nu$ channel (favored for $\tan\beta > 1$), with $\tau \to \pi\nu$ decay (Fig. 26).

SM $t$-decays give equal probabilities for $e, \nu, \tau$ leptons via $t \to bW \to b(e, \mu, \tau)\nu$, but the non-standard $t \to bH^+ \to b\tau\nu$ leads to characteristic excess of $\tau$. The strategy is to tag one top quark via standard $t \to bW \to b\ell\nu$ 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 various signals at SSC/LHC, and they all reach broadly similar conclusions [30, 33, 34, 35]. Figures 27 and 28 show typical limits of detectability for untagged and lepton-tagged $\gamma\gamma$ signals at SSC, assuming luminosities $\mathcal{L} = 20$ fb$^{-1}$ (two years of running) and $m_t = 150$ GeV. Figure 29 shows a similar limit for the $H \to 4\ell$ search (no $h \to 4\ell$ signal is detectable). Figure 30 shows typical limits for detecting the $t \to H^+ \to excess \tau$ 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. 31. For $m_t = 120$ GeV the inaccessible region is larger, for $m_t = 200$ GeV it is smaller.

Figure 32 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.

Figure 32 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.
An indirect constraint on the MSSM Higgs sector is provided by the CLEO bound on $b \rightarrow s\gamma$ decays [36],
\[ B(b \rightarrow s\gamma) < 8.4 \times 10^{-4} \quad (95\% \text{ C.L.}) . \] 
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. 33). In the MSSM both the $W$ and $H$ amplitudes have the same sign and the branching fraction is directly related to $m_{H^+}$ and $\tan \beta$ (Fig. 34); hence the CLEO result implies a lower bound on $m_{H^+}$ for given $\tan \beta$ (Fig. 35). It was recently pointed out [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. 36. Taken at face value, it appears to exclude a large part of the LEP II discovery region and furthermore to exclude much of the otherwise inaccessible region too; with future improvements in the CLEO bound, perhaps the whole of the inaccessible region could be excluded.

Fig. 33. $W$ and charged-Higgs loop diagrams contributing to $b \rightarrow s\gamma$ decays.

Fig. 34. Dependence of $B(b \rightarrow s\gamma)$ on $m_{H^\pm}$ and $\tan \beta$ in the MSSM (from Ref. [38]), neglecting other SUSY loops.

Fig. 35. Lower bound on $m_{H^+}$ for given $\tan \beta$, from $b \rightarrow s\gamma$ constraint [38]. The region excluded by the CLEO experimental bound is to the left of the $b \rightarrow s\gamma$ curve.

Fig. 36. Comparison of $b \rightarrow s\gamma$ bound with other MSSM Higgs constraints in the $(m_A, \tan \beta)$ plane, for $m_t = 150$ GeV [38].

It is premature however to reach any firm conclusions from the results above. The calculations of Ref. [38] are based on the approximation of Ref. [39], but later work indicates possible further small corrections [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 [41, 42]. However, as theoretical constraints on SUSY particles become more extensive, and as the $B(b \rightarrow s\gamma)$ bound itself becomes stronger, we may expect this approach to give a valuable constraint in the MSSM Higgs phenomenology.

[Postscript: at the Washington APS meeting April 1993, CLEO reported an improvement in the bound of Eq. (38) to $5.4 \times 10^{-4}$].

Finally we may ask what a future $e^+e^-$ collider could do. We have seen that part of the MSSM parameter space is inaccessible to $e^+e^-$ collisions at $\sqrt{s} = 200$ GeV, $\mathcal{L} = 500\text{ pb}^{-1}$, for $m_t = 150$ GeV and $m_{\tilde{t}} = 1$ TeV. But a possible future linear collider with higher energy and luminosity could in principle cover the full parameter space. In is interesting to know what are the minimum $s$ and $\mathcal{L}$ requirements for complete coverage, for given $m_t$. This question was answered in Ref. [43], based on the conservative assumption that only the channels $e^+e^- \rightarrow (Zh, Ah, ZH, AH) \rightarrow \tau\tau jj$ would be searched, with no special tagging. The results are shown in Fig. 37. We have estimated that including all $Z \rightarrow \ell\ell, \nu\bar{\nu}, jj$ and $h, H, A \rightarrow bb, \tau\tau$ 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. 37 would be rescaled downward by an order of magnitude.
Fig. 37. Minimal 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 [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 usually 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] R. Haag, J. Lopuszanski, and M. Sohnius, Nucl. Phys. B88, 257 (1975).

[2] For reviews, see J. Ellis, 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).

[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 R. Arnowitt, NUB-TH-3056/92; J. 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. D45, 2482 (1992); 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 by D. Kreinick at the Carleton University Beyond the Standard Model Conference, Ottawa, Canada, June 1992.

[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, MAD/PH/704 (1992), to be published in Phys. Rev. D.