Supersymmetry Phenomenology*

V. Barger

Physics Department, University of Wisconsin, Madison, WI 53706, USA

The phenomenological implications of a low-energy supersymmetry are surveyed, with particular attention given to unification constraints and the role of a large top quark Yukawa couplings. Generic expectations for sparticle mass spectra are presented along with prospects for their discovery and study at present and future colliders.

1. Introduction

The Standard Model (SM) is a pillar of success as an effective theory. Precision experiments agree with SM radiative corrections to an accuracy $\lesssim 0.1\%$. However, the SM Higgs sector is problematic. Longitudinal $W$-boson scattering, $W_L W_L \to W_L W_L$, violates unitarity if the Higgs mass exceed about 1 TeV, but the quadratic divergences in radiative corrections to the Higgs mass give a Higgs mass that is naturally of the order of the Planck mass. A way out of this contradiction is a low-energy fermion-boson supersymmetry (SUSY) in which each SM fermion (boson) has a boson (fermion) superpartner; see Table 1. Two Higgs doublets are required in the minimal supersymmetric standard model (MSSM), one ($H_u$) to give mass to the up-type quarks and leptons and the other ($H_d$) to give mass to the down-type fermions. The $W^\pm$, $H^\pm$ sparticle states mix and their spin-1/2 mass eigenstates are the charginos, denoted by $\tilde{\chi}^{\pm}_1,2$. Similarly, the neutral sparticles states $\tilde{B}^0, \tilde{W}^0, \tilde{H}_u^0, \tilde{H}_d^0$ mix to give the neutralino mass eigenstates $\tilde{\chi}^0_{1,2,3,4}$. In exact SUSY, a particle and its sparticle companion have the same mass and couplings. In broken SUSY the sparticles have higher masses than the particles but the exact SUSY coupling relationships are maintained. The additional radiative contributions to the Higgs mass from sparticle loops cancel the quadratic divergence SM loop contribution, solving the naturalness and gauge hierarchy problems. There is a vast literature on supersymmetry phenomenology and the reader may consult recent reviews[4] and textbooks[5, 6] for references.

2. Gauge Coupling Unification

The Renormalization Group Equations (RGE), found by analyzing loop corrections, predict the evolution of the couplings with energy scale. The $\alpha_i = g^2_i/(4\pi)$, with label $i=1,2,3$ for U(1), SU(2), SU(3) satisfy the RGEs,

$$\frac{d\alpha_i}{dt} = \frac{1}{2} \left[ b_i \alpha_i^2 + \frac{1}{4\pi} \sum_{j=1}^{3} b_{ij} \alpha_i^2 \alpha_j + \cdots \right]$$

Table 1

| Gauge multiplets | Boson fields | Fermionic partners |
|-----------------|--------------|--------------------|
| $SU(3)$         | $g^a$        | $\tilde{g}^a$     |
| $SU(2)$         | $W^i$        | $\tilde{W}^i$     |
| $U(1)$          | $B$          | $\tilde{B}$       |

| Matter multiplets | $\tilde{L}^i = (\tilde{\nu}_i, \tilde{e}^{-}_L)$ | $(\nu_i, e^-)_L$ |
|-----------------|-------------------------------------------------|-----------------|
|                  | $\tilde{\nu}^i = (\tilde{\nu}_L^i, \tilde{\nu}_R^i)$ | $(\nu, d)_L$ |
|                  | $\tilde{d}^i = (\tilde{d}_L^i, \tilde{d}_R^i)$ | $u^+_L$          |
|                  | $\tilde{u}^i = (\tilde{u}_L^i, \tilde{u}_R^i)$ | $d^+_L$          |
|                  | $H_u^0$ | $(\tilde{H}_u^0, \tilde{H}_d^0)_L$ |
|                  | $H_d^0$ | $(\tilde{H}_d^0, \tilde{H}_u^0)_L$ |

where $t = \ln(\mu/M_\text{G})$ and $\mu$ is the running mass parameter. The $b_i$ are known constants from the particle content of the loops. The running of the strong coupling constant $\alpha_s(t) = \alpha_3$ is now convincingly established; see Fig. 1[4].

Starting from the measured $\alpha_i$ at $\mu = M_Z$, the couplings can be extrapolated to higher scales, assuming that only particles of low mass contribute to the loops (i.e., there is a particle desert between the TeV scale and the unification scale). If a Grand Unified Theory (GUT) exists at a high scale $M_G$, then the three couplings should evolve to a common point of intersection. Unification of the gauge couplings is remarkably realized if supersymmetric particle masses are between $\sim 100$ GeV and $\sim 1$ TeV, as illustrated in Fig. 2 for a 1 TeV SUSY mass scale. Unification occurs for $\alpha_s(M_Z) = 0.13 \pm 0.01$ (see e.g. Ref. [3]), where the uncertainty is associated with the SUSY mass spectrum. This value is remarkably consistent with the LEP measurement $\alpha_s(M_Z) = 0.120 \pm 0.005$[3]. No such common intersection of the couplings occurs in the SM.

3. Yukawa Coupling Evolution

The fermion masses are generated by the vacuum expectation values (vevs) $v_u$ of the $H_u^0$ and $v_d$ of the
Then the third generation masses are related to the Higgs Yukawa couplings $\lambda_f$ as

$$m_t(m_t) = \lambda_t(m_t) v_u , \quad m_b(m_t) = \lambda_b(m_b) v_d ,$$

$$m_\tau(m_\tau) = \lambda_\tau(m_\tau) v_d .$$  \hspace{1cm} (3)

The Yukawa couplings also evolve with scale. In terms of the $Y_f \equiv \lambda_f^2/(4\pi)$, the RGEs for the Yukawa couplings at one-loop order are

$$\frac{dY_t}{dt} = \frac{1}{2\pi} Y_t \left( 6Y_t + Y_b - \frac{16}{3} \alpha_3 - 3\alpha_2^2 \right) ,$$

$$\frac{dY_b}{dt} = \frac{1}{2\pi} Y_b \left( Y_t + 6Y_b + Y_\tau - \frac{16}{3\alpha_3} - 3\alpha_2 \right) ,$$

$$\frac{dY_\tau}{dt} = \frac{1}{2\pi} Y_\tau \left( 3Y_b + 4Y_\tau + 3\alpha_2 \right) .$$ \hspace{1cm} (4)

Here the small contributions from $\alpha_1$ have been ignored. To correctly predict $m_b/m_\tau$ with $Y_b = Y_\tau$ unification\cite{6}, large values of $Y_t$ (close to the perturbative bound) at the GUT scale are required\cite{5}. In this circumstance evolution drives $\lambda_t$ towards a quasi-infrared fixed point ($d\lambda_t/dt \simeq 0$)\cite{8},\cite{9},\cite{10},\cite{11} independent of the precise value of $\lambda_t$ at the GUT scale.

Figure 3 shows the $\lambda_t$ infrared fixed point band in the $\tan\beta$ vs. $m_t$ plane. For $m_t = 175$ GeV there is a low $\tan\beta$ solution with $\tan\beta \simeq 1.8$ (\beta \simeq 60\degree), given by

$$m_t = 200 \mathrm{ GeV} \sin\beta .$$ \hspace{1cm} (7)

There is also a large $\tan\beta$ fixed point solution with $\tan\beta \simeq 56$ for which triple Yukawa coupling unification ($\lambda_t = \lambda_b = \lambda_\tau$) is approximately realized at the GUT scale\cite{5}. These fixed point solutions are very attractive theoretically, though it is too soon to rule out values of $\tan\beta$ between the two fixed points, or even below.

4. Soft SUSY Breaking

Supersymmetry is usually presumed to be broken in a hidden sector with this breaking transmitted to the observable sector via a) gravitational interactions ($N = 1$ supergravity) or b) gauge sector interactions of a messenger sector with effective scale $\lesssim 100$ TeV. Both categories of model contain “soft” breaking mass terms that do not reintroduce quadratic divergences. All gauge-invariant soft mass terms are included in the Lagrangian

$$\mathcal{L}_{\text{soft}} = - \sum_{\text{scalars}} m^2 A^2 - \sum_{\text{gauginos}} M(\lambda \lambda + \bar{\lambda} \bar{\lambda})$$

$$+ B_{\mu \nu \lambda \bar{\lambda}} \tilde{H}_d^i \tilde{H}_d^j + A_{\lambda \mu} \bar{Q} \tilde{U} \tilde{U}$$

$$+ A_{\beta \lambda} \bar{Q} \tilde{D} \tilde{D} + A_{\tau \lambda} \bar{L} \tilde{E} \tilde{H}_d .$$ \hspace{1cm} (8)

Subsequently we concentrate on the phenomenology of the minimal supergravity model (mSUGRA), which assumes universal soft SUSY breaking parameters at the GUT scale:
Figure 3. Contours of fixed $b$-quark mass in the plane of $\tan \beta$ versus the top-quark mass, along with contours of constant GUT scale Yukawa couplings. The top Yukawa infrared fixed point region is given by the line shading. From Ref. [7].

$m_0$ common scalar mass
$m_{1/2}$ common gaugino mass
$A_0$ common trilinear coupling
$B_0$ common bilinear coupling
$\mu_0$ Higgs mixing mass

Starting from these universal parameters at the GUT scale, the soft parameters at the electroweak scale are obtained from RGE evolution. Figure 4 shows typical results of the evolution. The Higgs miracle is explained by the evolution: With a large $\lambda_t$ at the GUT scale the mass-squared of $H_u$ is driven negative at the electroweak scale[11]. The values of $|\mu(M_Z)|$ and $B(M_Z)$ are then determined by the minimization of the Higgs potential. At tree level the minimization conditions are:

$$\mu^2 = \frac{m_{H_u}^2 - m_{H_u}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \frac{1}{2} M_Z^2 ,$$

$$B\mu = \frac{1}{2} (m_{H_u}^2 + m_{H_u}^2 + 2\mu^2) \sin^2 \beta .$$

For the $\lambda_t$ fixed point solution at low $\tan \beta$ the $A_t$ parameter also approaches a fixed point, independent of its GUT scale value[10]. Thus the low-energy phenomenology of the mSUGRA models is given in terms of the parameters $m_0, m_{1/2}, \text{sign} \mu, \text{and} \tan \beta$.

The evolution of the gaugino masses $M_1, M_2, M_3$ depends only on the gauge couplings

$$\frac{dM_i}{dt} = -b_i \alpha_i M_t ,$$

where $b_1 = 33/5, b_2 = 1, b_3 = -3$, so the $M_i$ at any scale are simply related,

$$M_3 / \alpha_3 = M_2 / \alpha_2 = M_1 / \alpha_1 .$$

The values of the $M_i$ at the $M_Z$ scale are given in terms of the GUT scale gaugino mass by

$$M_3 \simeq 3.2 m_{1/2} , \ M_2 \simeq 0.88 m_{1/2} , \ M_1 \simeq 0.44 m_{1/2} .$$

The gluino mass is $M_3$. In the following we discuss the qualitative features of sparticle masses for the low $\tan \beta \lambda_t$ fixed point, for which $|\mu| \gg M_Z$ is obtained from RGE evolution.

5. Chargino Sector

In a first approximation the mass matrix in the $(\tilde{W}^\pm, \tilde{H}^\pm)$ basis is diagonal,

$$\mathcal{M} = \begin{pmatrix} M_2 & 0 \\ 0 & -\mu \end{pmatrix} .$$

Thus the chargino mass eigenstates are

$$\chi_1^\pm \sim \tilde{W}^\pm , \ \chi_2^\pm \sim \tilde{H}^\pm .$$

6. Neutralino Sector

Here the approximate mass matrix in the $(\tilde{B}^0, \tilde{W}^0, \tilde{H}_u^0, \tilde{H}_d^0)$ basis is

$$\mathcal{M} = \begin{pmatrix} M_1 & 0 & 0 & \mu \\ 0 & M_2 & 0 & 0 \\ 0 & 0 & \mu & 0 \\ \mu & 0 & 0 & M_3 \end{pmatrix} .$$

Thus the two lightest neutralinos are approximately gauginos,

$$\chi_1^0 \sim \tilde{B}^0 , \ \chi_2^0 \sim \tilde{W}^0 .$$
The masses of the lightest color singlet ino states and the gluino are approximately related as
\[ \chi_0^0 : \chi_1^0 : \chi_1^\pm : \tilde{g} = 1 : 2 : 2 : 7. \] (18)

7. Stop Sector

The stop mass-squared matrix in the \( \tilde{t}_L, \tilde{t}_R \) basis is
\[ m^2 = \begin{pmatrix} L^2 & a m_t \\ a m_t & R^2 \end{pmatrix}, \] (19)
where
\[ a = A_t + \mu \cot \beta, \] (20)
\[ L^2 = m_t^2 + M_Q^2 - 0.35 M_Z^2 |\cos 2\beta|, \] (21)
\[ R^2 = m_t^2 + M_U^2 - 0.15 M_Z^2 |\cos 2\beta|. \] (22)

The off-diagonal terms are proportional to \( m_t \). Consequently there may be large mixing. Diagonalization of \( m^2 \) leads to two stop mass eigenstates \( \tilde{t}_1, \tilde{t}_2 \). The \( \tilde{t}_1 \) may be light if the mixing is large or if \( M_Q^2 \lesssim 0 \). The stop masses and mixings determine the precise value of the light Higgs boson mass through the radiative corrections to \( m_h \). For example, in the limit of large mass of the CP-odd Higgs state \( A \), the mass of the lightest CP-even Higgs boson \( h \) is given by [12]

\[ m_h^2 = M_Z^2 \cos 2\beta \left( 1 - \frac{3 m_t^2 t}{8 \pi^2 v^2} \right) \]
\[ + \frac{3 m_t^2}{4 \pi^2 v^2} \left[ t + \frac{\kappa \tilde{t}}{2} + \left( \frac{3}{32 \pi^2} \right) \frac{m_t^2}{v^2} \right] (\kappa \tilde{t} + t^2), \] (23)
with
\[ t = \ln \left( \frac{M_Z^2}{m_t^2} \right) \]
\[ \kappa = \frac{2 \tilde{A}_t}{M_S^2} \left( 1 - \frac{\tilde{A}_t^2}{12 M_S^2} \right) \]
\[ \tilde{A}_t = A_t + \mu \cot \beta, \]
where \( M_S^2 = (m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2)/2 \). The radiative corrections can substantially increase the tree-level bound \( m_{h^0} \lesssim M_Z |\cos 2\beta| \) to \( m_{h^0} \lesssim 130 \text{ GeV} \).

8. Generic SUSY Mass Spectra

Representative results for the SUSY mass spectra for the low tan \( \beta \) fixed point scenario are shown in Fig. 5, where the predicted masses for \( m_{1/2} = 150 \text{ GeV} \) are given versus \( m_0 \). Typical superpartner masses are
- selectron \( \gtrsim 70 \text{ GeV} \)
- sneutrino \( \gtrsim 100 \text{ GeV} \)
- stop \( \gtrsim 90 \text{ GeV} \)
- chargino \( \gtrsim 100 \text{ GeV} \)
- LSP \( \gtrsim 50 \text{ GeV} \)
- gluino \( \gtrsim 350 \text{ GeV} \)

The production and decays of the sparticles offer many interesting possibilities for experimental searches.

![Figure 5. Characteristic mass spectra for sparticles in the mSUGRA model (from Ref. [7].)](image)

9. Light Higgs Search

Ongoing searches for Higgs bosons at LEP-2 are based on the processes
\[ e^+ e^- \rightarrow (h \text{ or } H) Z \rightarrow b \bar{b} q \bar{q}, b \bar{b} \ell \ell, b \bar{b} \nu \nu, \tau \bar{\tau} q \bar{q}, \] (24)
\[ e^+ e^- \rightarrow h A \rightarrow b b b, \tau \bar{\tau} b b. \] (25)

Figure 6 shows the regions of the \((\tan \beta, m_h)\) plane presently excluded \((M_S = 1 \text{ TeV is assumed})\) [13].

Upgrades at the Tevatron collider (Main Injector,
Using vertex detectors to identify $h$ will allow for a SM Higgs search up to $120$ GeV and possibly higher through the process $(26)\, q\bar{q} \to Wh,\, q\bar{q} \to t\bar{t}h$.

Using vertex detectors to identify $h \to b\bar{b}$ decays. If $m_A$ is large, then the couplings of the lightest MSSM Higgs boson are essentially the same as the SM Higgs, and the mass reach for the MSSM Higgs is comparable to that above for the SM Higgs.

At the LHC at least one of the MSSM Higgs bosons should be found in any region of the $(\tan \beta, m_A)$ parameter space. For the lightest MSSM Higgs, the decays $h \to \gamma\gamma$ and $h \to ZZ^* \to 4$ leptons $(l = e, \mu)$ are important search modes, with the $h$ produced by gluon-gluon fusion ($gg \to h$).

10. Neutralino Dark Matter

A discrete quantum number known as $R$-parity, $R = (-)^{3B+L+2S}$, is commonly introduced in supersymmetry models to keep the proton sufficiently stable: $R = +1$ for particles and $R = -1$ for sparticles. Then sparticles are produced only in pairs and the lightest supersymmetric particle (LSP) is stable. In mSUGRA models the LSP is the lightest neutralino, $\chi_1^0$.

In the early Universe, when the temperature exceeded $m_{\chi_1}$, LSPs would have existed abundantly in thermal equilibrium, with the annihilation rate balanced by pair production. However, after the temperature dropped below $m_{\chi_1}$ and the annihilation rate dropped below the expansion rate of the Universe, a relic cosmological abundance of the LSP would remain; the LSP is an attractive candidate for the observed dark matter. In terms of $\Omega \equiv \rho/\rho_c$, where $\rho_c$ is the critical mass density to close the universe, clusters and large scale structure indicate $\Omega > 0.2$. However, the nucleosynthesis of the heavy elements constrains $\Omega h^2 \leq 0.03$, where $h$ is the Hubble constant in units of $100$ km/s/Mpc. Recent determinations find $h \simeq 0.65$. Thus the cosmologically interesting region for neutralino dark matter is $0.1 \leq \Omega_{\chi_1} h^2 \leq 0.5$.

Figure 7 illustrates regions of the mSUGRA $m_0$ and $m_{1/2}$ parameters that are compatible with the cosmological constraint, for the low $\tan \beta$ fixed point solution.[13] The neutralino explanation of dark matter indicates a low mass scale for supersymmetry in this fixed point scenario.

11. Finding Sparticles at Colliders

At $e^+e^-$ or $\mu^+\mu^-$ colliders, the production of sparticles may give very clean signatures. Two such examples are

$$e^+e^- \to \chi_1^+ \chi_1^- \to (W^+\chi_1^0)(W^-\chi_1^0),$$

$$e^+e^- \to \tilde{\mu}^+\tilde{\mu}^- \to (\mu^+\chi_1^0)(\mu^-\chi_1^0).$$

The $W$-bosons from the chargino decay can be real or virtual. From LEP-2 searches at c.m. energies $\sqrt{s} = 161–172$ GeV, the ALEPH collaboration has placed the limits[13]

$$M_{\chi_1^0} > 85 \text{ GeV}; \ \ M_{\tilde{\mu}^\pm} > 70 \text{ GeV}; \ \ M_{\tilde{t}_1} > 63 \text{ GeV}.$$ (30)

Post discovery, the next step will be to determine the sparticle spins, masses and couplings. The two-body decay $\tilde{\mu}_R \to \chi_1^0\mu^-$ is a particularly simple example. The two endpoints of the flat energy spectrum for the decay muon give the relations

$$\frac{m_{\tilde{\mu}_R}^2 - m_{\chi_1}^2}{2m_{\tilde{\mu}_R}} = \sqrt{\frac{1 - \beta_{\tilde{\mu}}}{1 + \beta_{\tilde{\mu}}} (E_{\mu}^{\text{max}}_{\text{lab}})},$$

$$\frac{m_{\tilde{\mu}_R}^2 - m_{\chi_1}^2}{2m_{\tilde{\mu}_R}} = \sqrt{\frac{1 + \beta_{\tilde{\mu}}}{1 + \beta_{\tilde{\mu}}} (E_{\mu}^{\text{min}}_{\text{lab}})},$$

where $\beta_{\tilde{\mu}} = (1 - 4m_{\mu}^2/s)^{1/2}$ is the smuon velocity. Thus the measured muon energy endpoints determine both $m_{\tilde{\mu}_R}$ and $m_{\chi_1}$.[19] Figure 8 shows the results of a realistic simulation. Similar mass determinations from kinematic endpoints of other decay processes are possible at hadron colliders[20].
The signatures of pair-produced heavy sparticles are isolated leptons, missing transverse energy (LSPs and neutrinos), and jets. Missing $E_T$ searches at the Tevatron exclude regions of gluino and squark masses as shown in Fig. 9. The trilepton signal from $ud \to W^{+*} \to \chi_1^0 \chi_2^0$ with $\chi_1^+ \to \chi_1^0 \ell^+ \nu$ and $\chi_2^0 \to \chi_1^0 \ell^+ \ell^-$ decays ($\ell = e, \mu$) gives the highest future SUSY mass reach at the Tevatron collider[23]. At the LHC many SUSY channels are accessible and there is a good safety margin for discovery up to the TeV scale. The comparative reach of the Tevatron, LHC and NLC is illustrated in Fig. 10 in the space of mSUGRA scalar and gaugino masses[24]. Contours of 1 TeV gluino and squark masses are given in the figure for reference.

12. Conclusion

Supersymmetry is a compelling extension of the Standard Model which solves the quadratic divergence problem, gives unification of the gauge couplings, and accounts for the dark matter in the Universe. The heavy top quark plays a pivotal role in unified SUSY models. A large top-quark Yukawa coupling, needed for $b$-$t$ unification and to explain electroweak symmetry breaking as a radiative effect, leads to an infrared fixed point prediction of the top quark mass. The stop sector gives important radiative contributions to the mass of the lightest Higgs boson.

Heavy sparticles will decay through multistep cascades. An example is gluino decay through the following chain

$$\tilde{g} \to q\tilde{q}, \quad \tilde{q} \to q\chi_i, \quad \chi_i \to \chi_j W, \quad \chi_j \to f\tilde{f}\chi_1^0.$$  (32)
Search strategies for sparticles are in place for the LEP-2, Tevatron, LHC, NLC and FMC colliders. Experiments at higher energies will soon reveal whether a weak scale supersymmetry exists. A SUSY revolution would rival the excitement of the last three decades when the quark structure of matter was uncovered and the SM put in place.

Acknowledgments

I am grateful to Chung Kao for helpful comments in the preparation of this report. This research was supported in part by the U.S. Department of Energy under Grant No. DE-FG02-95ER40896 and in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation.

REFERENCES

1. X. Tata, Proc. of the IX Jorge A. Swieca Summer School, Campos do Jordão, Brazil (in press), hep-ph/9706307; R. Arnowitt and P. Nath, Lectures presented at the VII J. A. Swieca Summer School, Campos do Jordao, Brazil, 1993 CTP-TAMU-52/93; J. Bagger in QCD and Beyond, Proceedings of the 1995 TASI, D. Soper, Editor (World Scientific, 1996); M. Drees, KEK-TH-501, hep-ph/9611409 (1996); S. Dawson, hep-ph/9612229 (1996); M. Dine, hep-ph/9612389 (1996); W. de Boer, Prog. Part. & Nucl. Phys. 33, 201 (1994).

2. J. Wess and J. Bagger, Supersymmetry and Supergravity, (Princeton University Press, revised edition 1991); G. G. Ross, Grand Unified Theories (Benjamin/Cummings, 1985); P. West, Introduction to Supersymmetry and Supergravity, (World Scientific, 1986); R. N. Mohapatra, Unification and Supersymmetry, (Springer-Verlag, 1986); H. J. Müller-Kirsten and A. Wiedemann, Supersymmetry: An Introduction with Conceptual and Computational Details (World Scientific, 1987).

3. J. Gunion, H. Haber, G. Kane and S. Dawson, The Higgs Hunter’s Guide, (Addison-Wesley, 2nd printing, 1991).

4. S. Martí i García, hep-ex/9704016 (1997).

5. P. Langacker and J. Erler, hep-ph/9703428 (1997); P. Langacker and N. Polonsky, Phys. Rev. D52, 3081 (1995).

6. M.S. Chanowitz, J. Ellis and M.K. Gaillard, Nucl. Phys. B128, 506 (1977).

7. V. Barger, M.S. Berger, P. Ohmann, Phys. Rev. D 47, 1093 (1993); Phys. Rev. D 49, 4908 (1994); V. Barger, M.S. Berger, P. Ohmann and R.J.N. Phillips, Phys. Lett. B 314, 351 (1993).

8. B. Pendleton and G.G. Ross, Phys. Lett. B 98, 291 (1981); C.T. Hill, Phys. Rev. D 24, 691 (1981).

9. C.D. Froggatt, R.G. Moorhouse and I.G. Knowles, Phys. Lett. B 298, 356 (1993).

10. J. Bagger, S. Dimopoulos and E. Masso, Phys. Rev. Lett. 55, 920 (1985); H. Arason, et al., Phys. Rev. Lett. 67 2933 (1991); and Phys. Rev. D 46, 3945 (1992); P. Langacker, N. Polonsky, Phys. Rev. D 50, 2199 (1994); W.A. Bardeen, M. Carena, S.
11. K. Inoue, A. Kakuto, H. Kamatsu, and S. Takeshita, Prog. Theor. Phys. 68, 927 (1982); L. Alvarez-Gaume, J. Polchinski, and M.B. Wise, Nucl. Phys. B221, 495 (1983); J. Ellis, J.S. Hagelin, D.V. Nanopoulos, and K. Tamvakis, Phys. Lett. B125, 275 (1983); L.E. Ibáñez and C. Lopez, Nucl. Phys. B233, 511 (1984); L.E. Ibáñez, C. Lopez, and C. Munuños, Nucl. Phys. 256, 218 (1985).

12. M. Carena et al., hep-ph/9602250, in Vol. 1, Report of the Workshop on Physics at LEP2, ed. by G. Altarelli, T. Sjostrand and F. Zwirner, CERN yellow report 96-01; M. Carena, M. Quiros, and C.E.M. Wagner, Nucl. Phys. B461, 407 (1996); H.E. Haber, R. Hempfling, and A.H. Hoang, Z. Phys. CC75, 539 (1997).

13. J. Carr et al., ALEPH Collaboration, CERN-PPE/97-071 (1997).

14. A. Stange, W. Marciano, and S. Willenbrock, Phys. Rev. D49, 1354 (1994); D50, 4491 (1992); S. Mrenna and G.L. Kane, hep-ph/9406337 (1994); S. Kim, S. Kuhlmann, and W.M. Yao, CDF/ANAL/EXOTIC/PUBLIC/3904, in Proceedings of the 1996 DPF/DPB Summer Study on New Directions for High-Energy Physics, Snowmass, CO.

15. E. Richter-Was, D. Froidevaux, F. Gianotti, L. Poggiali, D. Cavalli, S. Resconi, CERN-TH-96-111 (1996); CMS Technical Proposal, CERN/LHCC 94-38 (1994); Atlas Technical Proposal, CERN/LHCC 94-43 (1994).

16. G. Jungman, M. Kamionkowski and K. Griest, Phys. Rep. 267, 195 (1996); P. Nath and R. Arnowitt, presented at the International Workshop on Aspects of Dark Matter in Astrophysics and Particle Physics, Heidelberg, Germany, Sept. 1996, hep-ph/9610460; M. Drees, invited talk at 12th DAE Symposium on High-Energy Physics, Guahati, India, Dec. 1996, hep-ph/9703260; R. Barbieri, M. Frigeri and G.F. Giudice, Nucl. Phys. B 313, 725 (1989); J.L. Lopez, D.V. Nanopoulos and K.-J. Yuan, Phys. Lett. B 267, 219 (1991) and Nucl. Phys. B370, 445 (1992); J. Ellis and L. Roszkowski, Phys. Lett. B 283, 252 (1992); L. Roszkowski and R. Roberts, Phys. Lett. B 309, 329 (1993); M. Kawasaki and S. Mizuta, Phys. Rev. D 46, 1634 (1992); G.L. Kane, C. Kolda, L. Roszkowski and J.D. Wells, Phys. Rev. D 49, 6173 (1994); E. Diehl, G.L. Kane, C. Kolda and J.D. Wells, Phys. Rev. D 52, 4223 (1995). M. Drees and M.M. Nojiri, Phys. Rev. D 47, 376 (1993). R. Arnowitt and P. Nath, Phys. Lett. B 299, 58 (1993); B 307, 403(E) (1993); Phys. Rev. Lett. 70, 3696 (1993); Phys. Rev. D 54, 2374 (1996); J.L. Lopez, D.V. Nanopoulos and K. Yuan, Phys. Rev. D 48, 2766 (1993); M. Drees and A. Yamada, Phys. Rev. D 53, 1586 (1996). H. Baer and M. Brhlik, Phys. Rev. D 53, 597 (1996). J. Ellis, T. Falk, K.A. Olive and M. Schmitt, Phys. Lett. B 388, 97 (1996); CERN Report No. CERN-TH-97-105, hep-ph/9705444 and references therein.

17. V. Barger and C. Kao, UW-Madison report no. MADPH-97-992, hep-ph/9704063 (1997).

18. G. Cowan, CERN Particle Physics Seminar, Feb. 25, 1997.

19. S. Kuhlman et al., the NCD ZDR Design Group and NLC Physics Working Group, Physics and Technology of the Next Linear Collider, report submitted to Snowmass 96, SLAC-R-0485 [hep-ex/9605011] (1996).

20. I. Hinchliffe, F. Paige, M. Shapiro, J. Soderqvist, and W. Yao, Phys. Rev. D55, 5520 (1997).