A Simplified Summary of
Supersymmetry

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Abstract. I give an overview of the motivations for and theory/phenomenology
of supersymmetry.

INTRODUCTION

The overview consists of three parts. Namely,

• WHY do theorists find supersymmetry so attractive?

• WHAT do we expect to see experimentally?

• HOW do we go about observing what is predicted?

In each part, I will give only the barest outline of the relevant issues and
discussions. In order to keep the presentation simple, I will often be less
than precise in the technicalities. Well-known results will not be referenced in
detail. Further discussion and references can be found, for example, in [1–4].

WHY SUPERSYMMETRY IS ATTRACTIVE

First, there are some very general aesthetic considerations.

1) To appear in Future High Energy Colliders, proceedings of the ITP Symposium, U.C.
Santa Barbara, October 21-25, 1996, AIP Press. Also presented at the Aspen Winter
Conference on High Energy Physics, January 1997, Aspen, CO.
- SUSY is the only non-trivial extension of the Lorentz group which lies at the heart of quantum field theory; the simplest such extension is referred to as $N = 1$ supersymmetry and requires the introduction of a single (two-component) spinorial (anti-commuting) dimension to space, the extra dimension(s) being denoted \( \theta \). Taylor expansions of a superfield $\hat{\Phi}(x, \theta)$ then take the form $\hat{\Phi}(x, \theta) = \phi(x) + \sqrt{2}(\theta \psi(x)) + (\theta \theta) F(x)$ (higher orders in $\theta$ being zero), implying an automatic association of a spin-1/2 field $\psi(x)$ with every spin-0 field $\phi(x)$ and vice-versa ($F(x)$ is an auxiliary field, i.e. it does not represent a dynamically independent field degree of freedom).

- If SUSY is formulated as a local symmetry, then a spin-2 (graviton) field must be introduced, thereby leading automatically to (SUGRA) models in which gravity is unified with the other interactions. Further, SUGRA reduces to general relativity in the appropriate limit.

- SUSY appears in superstrings.

- Historically, adding the Lorentz group to quantum mechanics required introducing an antiparticle for every particle; why should history not repeat — SUSY + quantum field theory requires a ‘sparticle’ for every particle.

Of course, since we have not detected any of the superpartners of the known Standard Model (SM) particles, it is clear that supersymmetry is a broken symmetry. Thus, it is possible that supersymmetry could be irrelevant at the energy scales where experiments can currently be performed. However, there are many reasons to suppose that the superpartners have masses below a TeV and, therefore, could be discovered anytime now.

- String theory solutions with a non-supersymmetric ground state are quite problematic, implying that all the physical states of the theory should have similar masses aside from the effects of supersymmetry breaking (which should be a perturbation on the basic string theory solution).

- SUSY solves the hierarchy problem, e.g. $m_{\text{Higgs}}^2 < (1 \text{ TeV})^2$ (as required to avoid a non-perturbative WW sector), via spin-1/2 loop cancellation of spin-0 loop quadratic divergences:

\[
m_H^2 \sim [m_H^0]^2 + \lambda^2 (m_{\text{boson}}^2 - m_{\text{fermion}}^2) \ln \frac{\Lambda^2}{(m_f^2 + m_f^2)}, \quad (m_f \sim m_{SUSY}); \quad (1)
\]

$m_H^2$ can be small if $[m_H^0]^2$ is, provided $m_f^2 \lesssim (1 \text{ TeV})^2$.

- SUSY implies gauge coupling unification (at $M_U \sim \text{few} \times 10^{16}$ GeV) if $m_{SUSY} \lesssim 1 - 10 \text{ TeV}$ and there is nothing (other than complete SU(5) representations) between $m_{SUSY}$ and $M_U$. Of course there are some qualifications to this statement.
a) Perturbative unification requires that the number of families (only complete families lead to unification) must be $\leq 4$ [5].

b) There is a possible difficulty in the string theory context in that $M_U < M_S$ if we accept the perturbative result that $M_S = M_{\text{Planck}}/\sqrt{8\pi} \sim 2 \times 10^{18}$ GeV. However, it has been emphasized [6] that in non-perturbative approaches it may be possible for $M_S$ to be substantially lower, in principle even as low as a TeV.

c) Unification takes place only if there are exactly 2 Higgs doublets (+ singlets) below $M_U$. This is also the minimal Higgs field content required to give both up and down quarks masses and to guarantee anomaly cancellation.

d) The precise scale $m_{\text{SUSY}}$ preferred for exact unification depends upon the precise value of $\alpha_s(m_Z)$. For currently accepted values of $\alpha_s(m_Z) < 0.118$, an effective $m_{\text{SUSY}}$ value above 1 TeV is seemingly preferred although other subtle issues could alter this preference [7].

- Electroweak symmetry breaking (EWSB) occurs automatically for the simplest universal boundary conditions at $M_U$ by virtue of the fact that the mass-squared parameter for the Higgs field coupled to the top-quark ($m_{H_2}^2$) is driven negative (during evolution from $M_U$ down to $m_Z$) by the large top-quark Yukawa contribution to its renormalization group equation (RGE). The associated symmetry breaking occurs very naturally at an energy scale in the vicinity of $m_Z$ if $m_{H_2}^2 < (1 - 2 \text{ TeV})^2$ at $M_U$.

- If the supersymmetric partner of the QCD gluon has mass below $\sim 1$ TeV, then most GUT boundary conditions imply that the lightest supersymmetric particle (LSP) will have mass on the 100 GeV scale and would interact weakly and have other properties that make it a natural candidate for the cold dark matter of the universe. However, this is only true if this LSP is essentially stable. In certain variants of supersymmetry, to be discussed later, this is not the case.

**WHAT WE EXPECT TO SEE**

A) Sparticles: in the minimal supersymmetric model (MSSM) every normal SM particle has its supersymmetric counterpart.

$$\begin{align*}
[u,d,c,s,t,b]_{L,R} & [e,\mu,\tau]_{L,R} & [\nu_{e,\mu,\tau}]_L & g \quad & W^\pm, H^\pm & \gamma, Z, H^0_1, H^0_2 \\
[\bar{u},\bar{d},\bar{c},\bar{s},\bar{t},\bar{b}]_{L,R} & [\bar{e},\bar{\mu},\bar{\tau}]_{L,R} & [\bar{\nu}_{e,\mu,\tau}]_L & \bar{g} \quad & \bar{\chi}^\pm_{1,2} & \bar{\chi}^0_{1,2,3,4}
\end{align*}$$

The quark, lepton and neutrino partners are the spin-0 squarks, sleptons and sneutrinos; the partner of the gluon is the spin-1/2 gluino; and the partners...
of the charged (neutral) vector bosons and Higgs bosons are the spin-1/2 charginos (neutralinos). Often the latter can be approximately separated into the spin-1/2 bino and wino gaugino partners of the U(1) $B$ and SU(2) $W$ gauge fields and the higgsino partners of the Higgs fields. In other cases, these states are strongly intermixed.

There is a possibly exact discrete symmetry of the theory, called R-parity, such that SM particles have $R = +$ while the sparticle partners have $R = -$. If R-parity is an exact symmetry then any physical process must always involve an even number of sparticles, and the LSP, normally the $\tilde{\chi}^0_1$, will be stable against decay to SM particles.

B) A very special Higgs sector [8].

- For the minimal two-doublet Higgs sector, the physical eigenstates comprise two CP-even scalars ($h^0, H^0$), a CP-odd scalar ($A^0$), and a charged Higgs pair ($H^\pm$).

- SUSY implies that the Higgs self couplings have strength of order $g$, the SU(2) SM coupling, which has the consequence that $m_{A^0} \leq 130$ GeV ($\leq 150$ GeV, if singlet Higgs fields are added).

- For boundary conditions such that EWSB is an automatic consequence of the RGE's, $m_{H^0} \sim m_{A^0} \sim m_{H^\pm} > 200$ GeV is very probable, in which case the $h^0$ will have properties very much like those predicted for the SM Higgs ($h_{SM}$) in the minimal one-doublet SM, while the $H^0, A^0$ decouple from $WW, ZZ$.

Thus, there is little question as to what we should see, but the very uncertain nature of supersymmetry breaking implies that there is a great deal of uncertainty as to the exact mass scale at which we should see the new sparticles and as to the new experimental signatures that will appear when sparticles are produced.

There is one important general point. If R-parity is exact, the sparticles must be produced in pairs. (Limits on R-parity violation suggest that single sparticle production is at best very weak in any case.) Thus, in order to observe supersymmetric particles at hadron machines we must have significant $gg$ and/or $q\bar{q}$ luminosity at $\sqrt{s} > 2m_{\text{SUSY}}$; at an $e^+e^-$ or $\mu^+\mu^-$ collider the $\sqrt{s}$ must exceed the sum of the masses of the the two sparticles one hopes to observe. Large masses ($\gtrsim 1$ TeV) for some sparticles are certainly possible, in which case a lepton collider with $\sqrt{s} \sim 3 - 4$ TeV will be needed.

With regard to the Higgs sector, the $h^0$ is guaranteed to be light and should be easily detected, perhaps even at LEP2 or the Tevatron. If a light $h^0$ is not found, then we must abandon the possibility of low-energy supersymmetry as we now understand it. In the $m_{A^0} \sim m_{H^0} \sim m_{H^\pm} > 200$ GeV decoupling limit, the heavier Higgs must be pair produced, e.g. $e^+e^-, \mu^+\mu^- \rightarrow H^0 A^0, H^+H^-$. Energy reach could again prove to be crucial.
Supersymmetry breaking

The couplings of the complete complex of sparticles and particles are simultaneously fixed by the superpotential, denoted $W$, which involves products of superfields (with their particle and sparticle component fields). As a result, almost all couplings involving sparticles are related by ‘Clebsch-Gordon’ factors to the couplings of their SM counterparts. The only exception is the possible presence in $W$ of R-parity violating ($\tilde{R}$) couplings. I shall temporarily ignore such couplings, but will return later to this subject.

Most of the uncertainty in phenomenology is related to the many possible scenarios for supersymmetry breaking, which lead to many different predictions for the masses (especially relative masses) of the sparticles and for the detailed experimental signatures that will be present when they are produced. The main constraint on supersymmetry breaking is that it should be ‘soft’ in the sense that it should not destroy the very attractive SUSY solution to the naturalness and hierarchy problems. The possible supersymmetry breaking terms in the Lagrangian can then be enumerated.

- gaugino masses: $M_i \lambda_i \lambda_i$ where $i = 3, 2, 1$ for SU(3), SU(2), U(1) and $\lambda_i$ denotes the spin-1/2 partner of the corresponding gauge field.
- scalar masses: e.g.
  \[ m_{H_1}^2 |H_1|^2 + m_{H_2}^2 |H_2|^2 + m_{\tilde{t}, \tilde{b}}^2 (\tilde{t}_L \tilde{t}_L + \tilde{b}_L \tilde{b}_L) + m_{\tilde{t}_R}^2 \tilde{t}_R \tilde{t}_R + m_{\tilde{b}_R}^2 \tilde{b}_R \tilde{b}_R \] (2)
- ‘A’ terms: e.g. $A_t \lambda_i (\tilde{t}_L H_2^0 - \tilde{b}_L H_2^+ ) \tilde{t}_R^*$.
- ‘B’: $B\mu (H_1^0 H_2^0 - H_1^+ H_2^-)$, where $\mu$ is the parameter appearing in the superpotential term $W \ni \mu \tilde{H}_1 \tilde{H}_2$.

Altogether, including CP-violating phases, the above comprise 105 independent and unknown (although limited in magnitude if we are to maintain the naturalness of the model and avoid a charge and/or color breaking ground state) parameters beyond the SM. The obvious question is whether the possibility of so many a priori unknown parameters is a good or bad thing.

- Bad: If we want to know ahead of time exactly what to look for, then it is bad in that the uncertainties associated with so many parameters imply a very large range of phenomenological possibilities.
- Good: If we want to be confident that we will learn something from what we observe, then the existence of so many a priori unknown parameters is good. In particular, by evolving the low-energy parameters up to $M_U$, one can hope to uncover $M_U$-scale boundary conditions that imply an underlying organization for the 105 parameters that can be associated with an attractive and theoretically compelling GUT/string model.
The lesson for machine builders and experimentalists is that being prepared for a wide range of possibilities as to HOW? we see and fully explore SUSY is a necessary evil. The rest of the talk reviews some of the many possibilities discussed to date.

**HOW TO LOOK FOR SUPERSYMMETRY**

It will be important to check that the Higgs sector fits within the supersymmetric model constraints and to learn whether it contains more than the minimal two doublet fields, in particular whether or not there are additional singlet Higgs fields. Direct discovery of sparticles will be even more important. We would hope to eventually observe all the sparticles of the theory. I give a short description of Higgs phenomenology and then turn to sparticle phenomenology.

**Detecting the SUSY Higgs bosons**

It has been clearly established that for a SUSY Higgs sector consisting of only the minimal two-doublets at least one of the SUSY Higgs bosons will be detectable at the large hadron collider (LHC) and at the planned $\sqrt{s} \sim 500$ GeV next linear lepton collider (NLC). In the $m_{A^0} > 2m_Z$ decoupling limit, it is always the $h^0$ whose observability is guaranteed. Detection of the $H^0, A^0, H^\pm$ is not guaranteed. In particular, at the NLC $H^0 A^0$ and $H^+ H^-$ pair production is not kinematically allowed if $m_{A^0} \gtrsim \sqrt{s}/2$. Also, at the LHC there is a region, see Fig. 1, in the standard $(m_{A^0}, \tan \beta)$ parameter space (which specifies the tree-level properties of the Higgs bosons) characterized by moderate $\tan \beta \gtrsim 3$ and $m_{A^0} > 200$ GeV such that only the $h^0$ will be detectable.

The two most urgent and best-motivated questions are:

- • Is discovery of one Higgs boson guaranteed if singlet Higgs fields are added to the minimal two doublet fields? This question is particularly important in light of the fact that string theories typically lead to a Higgs sector with extra singlets.

- • What is required in order to discover the $H^0, A^0, H^\pm$ of the MSSM? Here, the region of concern is the $m_{A^0} > 200$ GeV parameter region that is natural when EWSB is automatically broken by virtue of the RGE’s.

Both questions have been explored in the literature and I summarize the conclusions to date.

- • It can be demonstrated [10–13] that at least one of the light CP-even Higgs bosons ($h_{1,2,3}$) of a two-doublet plus one singlet Higgs sector will be observed at the NLC. This result follows from the fact that if the
lightest \(h_1\) has weak \(ZZ\) coupling then (one of) the heavier ones must actually be almost as light and have substantial \(ZZ\) coupling (and thus be discoverable in the \(e^+e^-\rightarrow Z^*\rightarrow Zh\) production mode). This result probably extends to the inclusion of several singlet Higgs fields.

- At a \(\mu^+\mu^-\) collider, direct \(s\)-channel \(\mu^+\mu^-\rightarrow h_1\) production is guaranteed to be visible [14]. Only the (predicted) \((\mu^+\mu^-h_1)\) coupling is needed; the \(h_1\) can be decoupled from \(ZZ,WW\) without affecting its detectability. A scan search in the \(\leq 150\) GeV region is required with \(\Delta E_{\text{beam}} \lesssim 0.01\%\).

- At the LHC, discovery of at least one Higgs boson is no longer guaranteed if a singlet Higgs field is present in addition to the minimal two doublet fields [15]. The no-discovery “holes” in parameter space are never terribly large, but are certainly not insignificant when \(\tan\beta\) is in the moderate \(\tan\beta \sim 3 - 10\) range.

- At an \(e^+e^-\) collider, the only means for detecting any of the \(H^0, A^0, H^\pm\) when \(m_{A^0} \gtrsim \sqrt{s}/2\) (so that \(H^0A^0, H^+H^-\) pair production is forbidden) is to run the collider in the \(\gamma\gamma\) collider mode [16,17]. In this way, single \(\gamma\gamma\rightarrow H^0, A^0\) production can be observed for \(m_{A^0} \lesssim 0.8\sqrt{s}\) provided high integrated luminosity (\(e.g. L \sim 200\) fb\(^{-1}\)) is accumulated. There is still
no certainty that a $\gamma\gamma$ collider facility will be included in the final NLC plans, nor regarding the luminosity that would be available.

- At a $\mu^+\mu^-$ collider, direct $\mu^+\mu^- \rightarrow H^0, A^0$ production will be observable for any $m_{A^0}$ between $\sqrt{s}/2$ (the pair production limit) and $\sqrt{s}$ provided that $\tan \beta \gtrsim 3$ (below this the LHC will find the $H^0, A^0$, see Fig. 1) and that $L = 100$ fb$^{-1}$ (for $\sqrt{s} \sim 500$ GeV) is available for the $\sqrt{s}/2 \rightarrow \sqrt{s}$ scan [18]. It must be possible to maintain high instantaneous luminosity throughout the scan region (perhaps requiring two cheap storage rings designed for optimal luminosity in different $\sqrt{s}$ ranges).

- At both the $e^+e^-$ and $\mu^+\mu^-$ collider, it is envisioned that the machine will be upgraded to increasingly large $\sqrt{s}$. Once $\sqrt{s} \gtrsim 2m_{A^0}$, $H^0A^0$ and $H^+H^-$ pair production will be observable. It has been shown that this will be true even if these Higgs bosons have substantial decays to sparticles [19,20]. The hierarchy motivations for supersymmetry suggest that $m_{A^0}$ will certainly lie below $1–1.5$ TeV. A strong form of naturalness suggests $m_{A^0} \lesssim 500$ GeV [21].

If a SM-like $h^0$ is observed, the most precise determination of all its properties requires both $\sqrt{s} = 500$ GeV running at the NLC and an $s$-channel scan of the resonance peak in $\mu^+\mu^- \rightarrow h^0$ collisions at $\sqrt{s} \sim m_{h^0}$ [22]. This argues strongly for having both types of machine, especially since very substantial $L = 200$ fb$^{-1}$ accumulated luminosity is needed at both machines in order to achieve good precision for a model-independent determination of all the Higgs couplings and its total width. Once the $H^0, A^0$ and/or $H^\pm$ have been detected (whether in pair production or single production at the muon collider), the relative branching ratios for $H^0, A^0, H^\pm$ decays to different types of channels can be measured and very strong constraints will be placed on a host of the important parameters of the SUSY model, and thus on the GUT boundary conditions [19,20]. At a muon collider, $\mu^+\mu^- \rightarrow H^0, A^0$ studies would require substantial $L$ at a $\sqrt{s}$ that is almost certain to be significantly different than that employed for $\mu^+\mu^- \rightarrow h^0$ studies. Perhaps more than one muon collider will turn out to be needed.

**Detecting the sparticles**

The phenomenology of sparticles is determined by the nature and source of supersymmetry breaking and whether or not R-parity is broken. A selection of possibilities is discussed below. A useful recent review covering some of the following theoretical material is [23].
1 Minimal Supergravity (mSUGRA)

This is the simplest and most fully investigated model. It assumes that the boundary conditions are set at $M_U$, that there is a desert in between $M_U$ and $m_{SUSY}$, and that R-parity is exact. Further, at $M_U$ the supersymmetry breaking parameters listed earlier are taken to be universal and to have no CP-violating relative phases. Universal boundary conditions are natural in the picture where SUSY-breaking (SUSY) arises in a hidden sector and is only communicated to the visible sector at the GUT/string scale via interactions involving gravity (which of course knows nothing about quantum numbers). The result is a model specified by just five parameters and one sign. The GUT boundary conditions are:

$$m_{H_1} = m_{H_2} = m_{\tilde{q}_i} = m_{\tilde{\ell}_i} \ldots = m_0, \quad M_3 = M_2 = M_1 = m_{1/2},$$

$$A_{ijk} = A_0, \quad |\mu|, \quad \text{sign}(\mu), \quad B_0.$$  

Beginning with these boundary conditions at $M_U$, evolution to low energies yields some simple and important results.

- $M_3 : M_2 : M_1 \sim \alpha_3 : \alpha_2 : \alpha_3 \sim 7 : 2 : 1$, leading to a similar ratio of the (MS) gluino to wino to bino masses.  
- Approximate squark degeneracy is maintained for the first two generations, implying acceptably small FCNC.
- The third family stops are normally significantly mixed and split, and the lightest squark will be the lighter $\tilde{t}_1$ eigenstate.
- EWSB is automatic, as described earlier — it is convenient to trade the parameters $|\mu|$ and $B$ for $\tan \beta$ and the (known) value of $m_Z$.
- $|\mu|$ and $m_{A_0} \sim m_{H^0} \sim m_{H^\pm}$ are typically large ($> 200$ GeV) unless $\tan \beta \gg 1$;
- In the usual case where $\mu > M_2$, one finds the following: the LSP is the (stable) $\tilde{\chi}^0_1 \simeq \tilde{B}$ with $m_{\tilde{\chi}^0_1} \sim M_1$ and is a good (cold) dark matter candidate; the lightest chargino ($\tilde{\chi}^\pm_1$) and 2nd lightest neutralino ($\tilde{\chi}^0_2$) are approximately SU(2) charged and neutral winos and have similar mass $\sim M_2$; the heavier chargino ($\tilde{\chi}^\pm_2$) and neutralinos ($\tilde{\chi}^0_{3,4}$) are charged and neutral higgsinos with masses $\sim |\mu|$.
- If $m_0^2 \gg m_{1/2}^2$, then squarks, sleptons, etc. are heavier than the lighter (gaugino) $\tilde{\chi}$’s and are approximately degenerate;
- If $m_0^2 \ll m_{1/2}^2$, then $m_{\tilde{\tau}} < m_{\tilde{\tau}}$ and $m_{\tilde{t}_1}$ is distinctly smaller than the other $m_{\tilde{q}_i}$; charged sleptons can be lighter than most gauginos and it is possible (not included in later discussions) that the LSP could be the $\tilde{\nu}_\tau$.

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2) This assumes a sufficiently simple form for the Kahler potential.

3) The gluino pole mass is substantially larger than its $\overline{\text{MS}}$ mass.
The experimental signatures are determined by how the sparticles decay. In m\text{SUGRA}, the above-outlined mass hierarchy \(m_{\tilde{\chi}_0^3,4} \sim m_{\tilde{\chi}_1^\pm} > m_{\tilde{\chi}_2^\pm} > m_{\tilde{\chi}_1^0}\) implies that most signatures will result from chain decays — e.g.

\[
\tilde{g} \to q\bar{q}\tilde{\chi}_1^\pm \to q\bar{q} + \left\{ \begin{array}{l} \ell^\pm \nu\tilde{\chi}_1^0 \\ q\bar{q}\tilde{\chi}_1^0 \end{array} \right\}.
\]

(It is important to note that \(\tilde{g}\) decay leads with equal probability to either \(\ell^+\) or \(\ell^-\).) Since masses and mass differences are both substantial, events will be characterized by fairly energetic jets and leptons and by large \(E_T\) from the very weakly-interacting, stable final \(\tilde{\chi}_1^0\)'s.

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♦ Tevatron and LHC

Extensive Monte Carlo studies have determined the region of m\text{SUGRA} parameter space for which direct discovery of sparticles will be possible. Cascade decays lead to events with jets, missing energy, and various numbers of leptons. The Tev33 option at Fermilab will cover [24] the most natural [21] portion of parameter space. The maximum reach at the LHC is in the \(1\ell + \text{jets} + E_T\) channel; one will be able to discover squarks and gluinos with masses up to several TeV [25]. Some particularly important types of events are the following.

- **pp \to g \to \text{jets} + E_T** and \(\ell^\pm \ell^\pm + \text{jets} + E_T\), the latter being the like-sign dilepton signal [26]. The mass difference \(m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0}\) can be determined from jet spectra end points [26–28], while \(m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0}\) can be determined from \(\ell\) spectra end points in the like-sign channel [26–28] — an absolute scale for \(m_{\tilde{\chi}_1^\pm}\) can be estimated (\(\pm 15\%\)) by separating the like-sign events into two hemispheres corresponding to the two \(\tilde{g}\)'s [26], by a similar separation in the jets+\(E_T\) channel [25], or variations thereof [27,28].

- **pp \to \tilde{\chi}_1^+ \tilde{\chi}_1^- \to (\ell^\pm \tilde{\chi}_1^0)(Z^*\tilde{\chi}_1^0)\), which yields a trilepton + \(E_T\) final state when \(Z^* \to \ell^+ \ell^-\); \(m_{\tilde{\chi}_1^0} - m_{\tilde{\chi}_1^0}\) is easily determined if enough events are available [29].

- **pp \to \ell \ell \to 2\ell + E_T\), detectable at the LHC for \(m_\ell \lesssim 300\) GeV [25], i.e. the region of parameter space favored by mixed dark matter cosmology [30].

- Squarks will be pair produced and, for \(m_0 \gg m_{1/2}\), would lead to \(\tilde{g}\tilde{g}\) events with two extra jets emerging from the primary \(\tilde{q} \to \tilde{g}\tilde{g}\) decays.

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♦ NLC

Important discovery modes include the following [31].

- **e^+e^- \to \tilde{\chi}_1^+ \tilde{\chi}_1^- \to (q\bar{q}\tilde{\chi}_1^0)\) or \(\ell\nu\tilde{\chi}_1^0\) + \(q\bar{q}\tilde{\chi}_1^0\) or \(\ell\nu\tilde{\chi}_1^0\); \(m_{\tilde{\chi}_1^\pm}\) and \(m_{\tilde{\chi}_1^0}\) will be well-measured using spectra end points and beam energy constraints.

- **e^+e^- \to \ell \ell \to (\ell\tilde{\chi}_1^0)(\ell\tilde{\chi}_1^-), (\nu\tilde{\chi}_1^-)(\bar{\nu}\tilde{\chi}_1^0)\), \ldots; masses will be well-measured.
The \( \tilde{q} \)’s and (if \( m_0 \) is big) \( \tilde{\ell} \)’s can be too heavy for pair production at \( \sqrt{s} = 500 \text{ GeV} \). If \( m_0 \sim 1 - 1.5 \text{ TeV} \) (the upper limit allowed by naturalness), then \( \sqrt{s} \gtrsim 2 - 3 \text{ TeV} \) is required. It could be that such energies will be more easily achieved in \( \mu^+\mu^- \) collisions than in \( e^+e^- \) collisions.

For certain choices of the mSUGRA parameters, it can happen that the phenomenology is much more ‘peculiar’ than the canonical situation described above. One such choice of boundary conditions is that termed the ‘Standard’ Snowmass 96 Point [28]:

\[
m_0 = 200 \text{ GeV}, \quad m_{1/2} = 100 \text{ GeV}, \quad A_0 = 0, \quad \tan\beta = 2, \quad \mu < 0.
\]

These boundary conditions predict that the gluino mass is approximately the same as the average squark mass; only the \( \tilde{t}_1 \) and \( \tilde{b}_L \) are lighter than the \( \tilde{g} \). Consequently, \( \tilde{g} \rightarrow \tilde{b}_L b \) almost 100% of the time. Also, masses are not very large: \( e.g. \, m_\tilde{g} = 285 \text{ GeV}, \, m_{\tilde{b}_L} = 266 \text{ GeV} \). At the LHC, there will be millions of spectacular events with \( \tilde{g}\tilde{g} \rightarrow b\tilde{b}_L \tilde{b}_L \rightarrow \ldots \). At the NLC, all \( \tilde{q}\tilde{q} \) thresholds are \( > \sqrt{s} = 500 \text{ GeV} \), but \( \tilde{\ell}\tilde{\ell} \) and \( \tilde{\chi}^\pm_1 \tilde{\chi}^\mp_1, \tilde{\chi}^0_1, \tilde{\chi}^0_{1,2} \) pair processes are all kinematically allowed.

There are particular choices for mSUGRA boundary conditions that have strong theoretical motivation in the context of strings and/or supergravity. These include:

- **No-Scale** [32]: \( m_{1/2} \neq 0, \, m_0 = A_0 = 0; \)
- **Dilaton or Dilaton-Like** [33]: \( m_{1/2} = -A_0 = \sqrt{3}m_0 \)

Here the dilaton is the string ‘modulus’ field associated with coupling constant strength in string models. To many, the dilaton-like boundary conditions appear to be particularly worthy of being taken seriously. To deviate significantly from dilaton-like boundary conditions in Calabi-Yau and Orbifold models requires going to an extreme in which the dilaton has nothing, or almost nothing, to do with supersymmetry breaking. The dilaton and no-scale models have many common features. Most importantly, the small (or zero) \( m_0 \) compared to \( m_{1/2} \) implies that sleptons are light, which in turn leads to excellent LEP2, TeV33, LHC, and NLC leptonic signals for supersymmetry [34], the only uncertainty being the overall rates as determined by the overall mass scale, \( m_{1/2} \). (\( \tan\beta \) and sign(\( \mu \)) are the only other free parameters in these models.)

A quite different set of mSUGRA boundary conditions is that motivated by the assumption that there is hidden-sector dynamical SUSY breaking without gauge singlets. In this case it is natural for all dimension-3 SUSY-breaking operators in the low-energy theory to be very small, \( i.e. \, m_{1/2}, A_0 \sim 0 \) (and possibly \( B_0 \sim 0 \) as well) [35]. Such boundary conditions, in combination with existing experimental constraints, imply that the gluino would be lighter than the lightest neutralino [36], and would tend to emerge as part of a relatively long-lived gluon-gluino bound state (denoted \( R^0 \)) with mass \( \sim 1.4 \text{ GeV} \).
(according to lattice calculations). Remarkably, this scenario cannot yet be absolutely excluded by accelerator experiments [37]. In addition, it is possible for the photino in such a model to have the necessary properties to be a dark matter candidate [38]. Ongoing experiments and analyses will be able to exclude this scenario in the near future. For example, for such boundary conditions the lightest chargino must have $\tilde{\chi}^\pm_1 < m_W$ and should be discovered at LEP2 once substantial luminosity is accumulated at $\sqrt{s} \sim 190$ GeV, despite the non-canonical nature of predicted signals. Searches for hadrons containing gluinos could also provide strong constraints [39].

2 Beyond mSUGRA: Non-universality

Despite the very attractive nature of the mSUGRA boundary conditions, it is certainly possible to find motivation for many possible sources of non-universality. The two general classes of non-universality are gaugino mass non-universality and scalar mass non-universality. Ref. [40] provides a useful review and references. Due to lack of space I discuss only (and very briefly) gaugino mass non-universality. Models characterized by such non-universality include:

- The O-II orbifold string model in which all matter fields lie in the $n = -1$ untwisted sector and SUSY-breaking is dominated by the overall size modulus (not the dilaton).

- $F$-term supersymmetry breaking with $F \neq \text{SU}(5)$ singlet, leading to $\mathcal{L} \sim \frac{\langle F \rangle_{ab}}{M_{\text{Planck}}} \lambda_{a} \lambda_{b}$, where $\lambda_{a,b}$ ($a,b = 1, 2, 3$) are the gaugino fields. If $F$ is an SU(5) singlet then $\langle F \rangle_{ab} \propto c \delta_{ab}$ and we get standard universality, $M_1 = M_2 = M_3$ (at $M_W$). But, more generally $F$ can belong to a non-singlet SU(5) representation: $F \in (24 \times 24)_{\text{symmetric}} = 1 \oplus 24 \oplus 75 \oplus 200$. which implies that $\langle F \rangle_{ab} = c_{a} \delta_{ab}$, with $c_{a}$ depending on the representation (an arbitrary superposition of representations is also possible).

The $M_{3,2,1}$ for these different cases are compared in Table 1.

I have room for only a few remarks regarding how the phenomenology changes as a function of boundary condition scenario. I focus on the most extreme changes. See Ref. [40] for more details and references.

- The general relations that $m_{\tilde{\chi}_1^0} \sim \min(M_1, M_2)$ and $m_{\tilde{\chi}_1^\pm} \sim M_2$ imply that $m_{\tilde{\chi}_1^\pm} \simeq m_{\tilde{\chi}_1^0}$ (both are winos) in the 200 and O-II scenarios, where $M_2 \lesssim M_1$. Some important consequences of this result are [41]:

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4) Readers not familiar with string terminology can simply take these specifications as names.
TABLE 1. $M_u$ at $M_U$ and $m_Z$ for the four $F_\Phi$ irr. reps. and in the $\delta_{GS} \sim -4$ O-II model. From Ref. [40].

| $F$ | $M_u$ | $m_Z$ |
|-----|-------|-------|
|     | $M_3$ | $M_2$ | $M_1$ | $M_3$ | $M_2$ | $M_1$ |
| 1   | 1     | 1     | 1     | $\sim 7$ | $\sim 2$ | $\sim 1$ |
| 24  | 2     | -3    | -1    | $\sim 14$ | $\sim -6$ | $\sim -1$ |
| 75  | 1     | 3     | -5    | $\sim 7$ | $\sim 6$ | $\sim -5$ |
| 200 | 1     | 2     | 10    | $\sim 7$ | $\sim 4$ | $\sim 10$ |
| O-II | $\delta_{GS} = -4$ | 1     | $53\over 1$ | $\sim 6$ | $\sim 10$ | $\sim 53\over 1$ |

- At the NLC, $e^+ e^- \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ would be very hard to see since the (invisible) $\tilde{\chi}_1^0$ would take all of the $\tilde{\chi}_1^\pm$ energy in the $\tilde{\chi}_1^\pm \rightarrow \ell^\pm \nu \tilde{\chi}_1^0$, $q \bar{q} \tilde{\chi}_1^0$ decays. One must employ $e^+ e^- \rightarrow \gamma \tilde{\chi}_1^+ \tilde{\chi}_1^-$.  

- At the LHC, the like-sign dilepton signal (coming again from $\tilde{\chi}_1^\pm \rightarrow \ell^\pm \nu \tilde{\chi}_1^0$ decays) would be very weak. In the O-II model, $m_{\tilde{g}} \sim m_{\tilde{\chi}_1^\pm} \simeq m_{\tilde{\chi}_1^0}$ means that the jets from $\tilde{g} \rightarrow q \bar{q} \tilde{\chi}_1^0$ decay would also be soft; the standard $E_T$ signature would be much weaker than normal and the maximum $m_{\tilde{g}}$ for which SUSY could be discovered would be smaller than in the usual case.

- Dark matter phenomenology would be substantially altered. In particular, the very close degeneracy $m_{\tilde{\chi}_1^\pm} \simeq m_{\tilde{\chi}_1^0}$ implies very similar $\tilde{\chi}_1^\pm, \tilde{\chi}_1^0$ densities (due to very similar Boltzmann factors) at freeze-out which in turn leads to $\tilde{\chi}_1^0, \tilde{\chi}_1^\pm$ ‘co-annihilation’ that greatly reduces relic dark matter ($\tilde{\chi}_1^0$) density.

- Generally speaking, if $m_{\tilde{g}}$ is not large then distinguishing between the five scenarios would be quite easy. For example, if we keep $m_0$ and $m_{\tilde{g}}$ the same as at the Snowmass 96 point described earlier, there would be many millions of $\tilde{g} \tilde{g}$ pair events at the LHC, and the different scenarios would lead to the following very different event characteristics [40]:

- 1: $E_T$, $\ell^+ \ell^-$, 4$b$’s;  
- 24: $E_T$, no $\ell^+ \ell^-$, 8$b$’s, with 2 pairs having mass $m_{h^0}$;  
- 75: traditional cascade chain decay signals;  
- 200: $E_T$, 4$b$’s, no leptons;  
- O-II: $\tilde{g} \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^0$ + soft and $\tilde{\chi}_1^- \rightarrow \tilde{\chi}_1^0$ + very soft; soft jet cuts required to observe.
3 R-parity violating models

I next consider the possibility that R-parity is violated. R-parity violation can come in two different forms:

- **Hard**: that is explicit terms in the superpotential of the form \( W_R = \lambda_{ijk} (\hat{L}_i \hat{L}_j \hat{E}_k + \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k + \lambda''_{ijk} \hat{U}_i \hat{D}_j \hat{D}_k) \).

- **Soft**: \( W_R = \mu_i \hat{L}_i \hat{H}_1 \)

I make a few brief remarks regarding the former. (Phenomenology for the latter is reviewed in Ref. [42].) First, \( \lambda, \lambda' \neq 0 \) implies lepton-number violation, while \( \lambda'' \neq 0 \) implies baryon-number violation. Both cannot be present; if they were, the proton lifetime would be short. Otherwise, current constraints are not terribly strong [43]: \( \lambda' \)'s < \( \sim 0.1 \) for superparticle masses in the > 100 GeV range. Phenomenologically, the crucial point is that unless the \( \lambda' \)'s are very, very small, the LSP \( \tilde{\chi}_1^0 \) will decay inside the detector. Thus, the signals for supersymmetry will no longer involve missing energy associated with the \( \tilde{\chi}_1^0 \). Impacts on phenomenology are substantial.

At the LHC [44]:

- If \( \lambda'' \neq 0 \), then \( \tilde{\chi}_1^0 \to 3j \). The large jet backgrounds imply that we would need to rely on the like-sign dilepton signal. For universal boundary conditions, this signal turns out to be sufficient for supersymmetry discovery out to \( m_\tilde{g} \) values somewhat above 1 TeV. However, if the leptons are very soft, as occurs if \( m_{\tilde{\chi}_1^\pm} \sim m_{\tilde{\chi}_1^0} \) (as in the 200 and O-II models) then the discovery reach would be much reduced — the combination of \( R \) and \( M_2 \lesssim M_1 \) would be a bad scenario for the LHC.

- If \( \lambda \neq 0 \), \( \tilde{\chi}_1^0 \to \mu^\pm e^\mp \nu_e, e^\pm e^\mp \nu_e \), and there would be many very distinctive multi-lepton signals.

- If \( \lambda' \neq 0 \), \( \tilde{\chi}_1^0 \to \ell 2j \) and again there would be distinctive multi-lepton signals.

At the NLC:

- Even \( e^+ e^- \to \tilde{\chi}_1^0 \tilde{\chi}_1^0 \to (3j)(3j), (2\ell \nu)(2\ell \nu), (\ell 2j)(\ell 2j) \) could yield an observable SUSY signal.

At HERA:

- Squark production via R-parity violating couplings could be an explanation for the HERA anomaly at high \( x \) and large \( Q^2 \) [45]. For example, if \( \lambda'_{113} \sim 0.04 - 0.1 \), a lepto-quark–like signal from \( e^+ d_R \to \tilde{t}_L \to e^+ d \) would be detected if \( m_{\tilde{t}_L} \lesssim 220 \) GeV. (We focus on the top squark since it can easily be the lightest squark in supergravity models.)
Finally, if R-parity is violated then supersymmetry will no longer provide a source for dark matter, the LSP no longer being stable. One would have to turn to neutrinos (which, however, only provide a source of hot dark matter, whereas some cold dark matter also seems to be needed).

4 Event-Motivated Models

A number of specific supersymmetry breaking parameter choices have been proposed in order to explain particular anomalous events seen at the Tevatron and at LEP. Here, I only have room to list some of these and make a few remarks. In all cases the $M_U$-scale boundary conditions would be required to be non-universal (assuming the usual desert between low energy and $M_U$).

CDF $ee\gamma\gamma$ event = $\tilde{e}\tilde{e}$, $\tilde{\chi}^0_1\tilde{\chi}^0_1$, ... production?

Let us focus on the $\tilde{e}\tilde{e}$ explanation. One proposed explanation of this event [46] as $\tilde{e}\tilde{e}$ production requires that the one-loop decay $\tilde{\chi}^0_2 \to \gamma\tilde{\chi}^0_1$ dominate over all tree level decays and that $BF(\tilde{e} \to e\tilde{\chi}^0_2) \gg BF(\tilde{e} \to e\tilde{\chi}^0_1)$. If both are true, then the process

$$\tilde{e}\tilde{e} \to (e\tilde{\chi}^0_2)(e\tilde{\chi}^0_2) \to (e\gamma\tilde{\chi}^0_1)(e\gamma\tilde{\chi}^0_1)$$

would explain the observed event provided masses are also appropriate. For the above decays to be dominant requires $\tilde{\chi}^0_2 = \tilde{\gamma}$ and $\tilde{\chi}^0_1 = \text{higgsino}$, which in turn is only the case if $M_2 \sim M_1 \tan\beta \sim 1 \ |\mu| < M_{1,2}$. This, in combination with the observed kinematics of the event and a cross section consistent with its having been detected, implies masses for all the neutralinos and charginos that are mostly in the $50 - 150$ GeV range. For such masses there should be a large number of other equally distinctive events in the $L \sim 100$ pb$^{-1}$ of Tevatron data currently being analyzed and clear signals should also emerge at LEP when run at $\sqrt{s} \sim 190$ GeV. At the Tevatron, $L \sim 100$ pb$^{-1}$ implies $N(2\ell + X + E_T) \geq 30$, $N(\gamma\gamma + X + E_T) \geq 2$, $N(\ell\gamma + X + E_T) \geq 15$, $N(2\ell\gamma + X + E_T) \geq 4$, $N(2\ell\gamma + X + E_T) \geq 2$, $N(3\ell + X + E_T) \geq 2$. At LEP190 with $L = 500$ pb$^{-1}$, one finds $N(2\ell + X + E_T) \geq 50$, $N(2\gamma + X + E_T) \geq 3$. In the above, $X =$additional leptons, photons, jets. Finally, I note that the small $\mu$ value needed implies that we must have soft mass non-universality in the form $m_{H_0}^2 \neq m_{\tilde{q},\tilde{\ell}}^2$.

CDF and D0 'di-lepton top events' that don't look like top events

There are two events in the CDF di-lepton top sample and one event in the D0 sample that have very low probability, in terms of their kinematics, to be top events. It has been proposed [47] that these could be from $\tilde{q} \to q\tilde{\chi} \to q(\nu\tilde{\ell})$ with $\tilde{\ell} \to \ell\tilde{\chi}^0_1$, where $\tilde{\chi} = \tilde{W}_3, \tilde{W}^\pm$. For this explanation to work in terms of
kinematics and cross section requires: \( m_\tilde{g} \simeq 330, m_\tilde{q} \simeq 310, m_\tilde{\tau} \simeq 220, m_\tilde{\nu}_\tau \sim 220, m_\tilde{e}_R \simeq 130, \mu \sim -400, M_1 \simeq 50, \) and \( M_2 \simeq 260. \) Note that these values are inconsistent with the previous \( ee\gamma\gamma \) event explanation. An abundance of other signals should be hidden in the full Tevatron data set for such parameter choices.

**The Four-Jet ‘Signal’ at ALEPH**

The 4-jet signal at ALEPH is well-known. It has \( M + M' \sim 106 \text{ GeV} \) and \( M - M' \sim 10 \text{ GeV} \) (i.e. \( M \sim 58, M' \sim 48 \text{ GeV} \)), significant ‘rapidity-weighted’ jet charge and no b jets. One interpretation [48] is as \( \bar{\nu}_e \nu_e \) production with \( \bar{e}_L \bar{e}_R \to 2j \) via \( H \) coupling \( W_R \equiv \chi'_{ijk}(\hat{L}\hat{Q})_L D^k_{R1}. \) \( M_1 \lesssim 100 \text{ GeV} \) is required for the needed \( \sigma \sim 1 - 2 \text{ pb} \) cross section level, and is also helpful in suppressing \( \bar{e}_L \bar{e}_L, \bar{e}_R \bar{e}_R \) production. To suppress (unobserved) \( \bar{e}_L \nu_e \) production we need the largest \( m_\tilde{\nu}_e \) possible, which occurs for \( \tan \beta \sim 1. \) For \( \bar{e}_L \to \nu_e \) to dominate over the more standard \( \bar{e}_L \to ee\bar{e}_R \) decays via \( \chi_1^0 \) exchange requires \( \lambda_{1ijk} \gtrsim \text{few} \times 10^{-4} \), which is entirely consistent with known bounds. In this picture, the \( \bar{e}_R \) must decay by mixing, \( \bar{e}_R \to \bar{e}_L \to \nu_1 d_k \) rather than via \( \bar{e}_R \to ee\bar{e}_R \) from virtual \( \chi_1^0, \hat{L} \) exchange; the required mixing angle, \( \sin \phi \gtrsim 10^{-4} \), is easily accommodated. One finds that a displaced vertex might be observable. The absence of \( \bar{\mu} \) and \( \bar{\tau} \) pair signals requires that these states have substantially heavier masses than the \( \bar{e} \), as would only be possible if slepton masses are non-universal. Finally, lots of signals would be expected at LEP2: for \( M_1 = 100 \text{ GeV}, m_\tilde{\nu}_e = 58 \text{ GeV} \), and \( \sqrt{s} = 186 \text{ GeV} \),

\[
\sigma(\bar{e}_L \bar{e}_R, \bar{e}_L \bar{e}_L, \bar{e}_R \bar{e}_R, \bar{\nu}_e L \bar{\nu}_e L, \chi_1^0 \chi_1^0) = 1.33, 0.23, 0.2, 0.29, 0.62 \text{ pb}
\]

5 Models with gauge-mediated supersymmetry breaking

Underlying all our previous discussions has been the implicit assumption that SUSY is almost certainly an exact local symmetry (as in strings) which, if not accidental, should be spontaneously broken in a ‘hidden sector’. The SUSY-breaking is then fed to the ordinary superfields in the form of an effective Lagrangian, \( \mathcal{L}_{\text{eff}} \). The mass expansion parameter for \( \mathcal{L}_{\text{eff}} \) is the mass scale of the sector responsible for the communication between the hidden sector and the ordinary superfields. For gravity-mediated communication (as appropriate in SUGRA/superstring theories), this mass scale is most naturally \( \sim M_{\text{Planck}} \) and one arrives at the \( \mathcal{L}_{\text{eff}} \sim (F) \frac{\lambda_a \lambda_a}{M_{\text{Planck}}} \) expression given earlier. For gaugino masses of order \( m_W \), this requires \( \langle F \rangle \equiv m_{\text{SUSY}}^2 m_W M_{\text{Planck}} \sim (10^{11} \text{ GeV})^2 \).

The gravity-mediated scenario is certainly very attractive but has its doubters, primarily based on the fact that keeping FCNC phenomena at a sufficiently small level is not guaranteed. Indeed, although gravity being ‘flavor blind’ suggests universal soft scalar masses, terms that violate universality
are certainly possible in the Kahler potential and (since not forbidden by symmetry) are generated at some level by radiative corrections, even if not present at tree-level. Further, a certain amount of FCNC and lepton-number violation can arise during evolution from $M_S \rightarrow M_U$, although, as noted earlier, $M_S \sim M_U$ is also possible non-perturbatively. A final point of concern is that EWSB via the RGE’s is not guaranteed for all possible boundary condition choices. All these problems can be resolved if the scale of supersymmetry breaking, $m_{\text{SUSY}}$, is much lower.

The specific and popular models that incorporate this idea are the gauge-mediated supersymmetry breaking (GMSB) models. In these models, $M_{\text{Planck}}$ above is replaced by $M$, the mass scale of a new messenger sector, and $m_{\text{SUSY}}^2 \sim m_W M$ can be much lower — $m_{\text{SUSY}} \sim 10 - 100$ TeV is often discussed. The GMSB models have rather few parameter, at least in current incarnations. They lead to dramatic signals, but the signals could well be somewhat different than at first anticipated because of issues related to the mass scale of supersymmetry breaking. I briefly outline the basics [49,4].

In the standard GMSB models there are three sectors.

I: First, there is the SUSY-breaking sector, sometimes called the ‘secluded’ sector, containing hidden particles that interact via strong gauge interactions which cause supersymmetry breaking characterized by a scale $\sqrt{F}$. The SUSY-breaking is then fed via two-loops in the strong interaction into

II: the ‘messenger’ sector, which contains a SU(3) $\times$ SU(2) $\times$ U(1) (but not necessarily SU(5)) singlet superfield $\hat{S}$ with non-zero vacuum expectation values for both its scalar component and its $F$-term component ($\langle S \rangle \neq 0$ and $\langle F_S \rangle \neq 0$). The two-loop communication between the two sectors implies that $\langle F_S \rangle \sim (\alpha_m^2/16\pi^2)F$, where $\alpha_m$ characterizes the strength of the gauge interactions that are responsible for the two-loop communication between the SUSY-breaking sector and the messenger sector. In addition to $\hat{S}$, the messenger sector must contain some messengers that transform under SU(3), SU(2) and U(1). In order to maintain actual gauge-coupling unification, these messengers should form a complete anomaly-free GUT representation (e.g. $5 + \overline{5}$ with messengers $\hat{q}, \hat{\overline{q}}$ and $\hat{\ell}, \hat{\overline{\ell}}$ superfield triplets and doublets in the SU(5) case). These messenger superfields must communicate with $\hat{S}$; a typical superpotential might be $W = \lambda_1 \hat{S}\hat{q}\hat{\overline{q}} + \lambda_2 \hat{S}\hat{\ell}\hat{\overline{\ell}}$. The mass scales of the component boson ($b$) and fermion ($f$) fields of these messenger superfields are given very roughly by

$$M \equiv m_{\text{mess}} \sim \lambda \langle S \rangle; \quad m_{\text{mess}}^2 \sim \left( \begin{array}{c} \lambda^2 \langle S \rangle^2 \\ \lambda^2 \langle F_S \rangle^2 \\ \lambda^2 \langle S \rangle^2 \end{array} \right)$$

(where $\lambda$ is the typical $\lambda_{1,2}$), implying that $\lambda \langle S \rangle^2 > \langle F_S \rangle$ is required to avoid a negative mass-squared eigenvalue. Ratios of $m_b^2$ eigenvalues $< 30$ (no fine tuning) suggests $\lambda \langle S \rangle^2/\langle F_S \rangle \gtrsim 1.05$. 
III: These messengers communicate SUSY-breaking to the normal superfield sector according to a scale characterized by \( \Lambda \equiv \langle F_S \rangle / \langle S \rangle \) (with \( M/\Lambda = \lambda \langle S \rangle^2 / \langle F_S \rangle > 1 \) required, as noted above). After integrating out the heavy messengers, and assuming that \( \tilde{S} \) is an SU(5) singlet, the masses of the particles important to low-energy phenomenology are as follows.

- The gauginos acquire masses at one-loop given by
  \[
  M_i(M) = k_i N_{5,10} g \left( \frac{\Lambda}{M} \right) \frac{\alpha_i(M)}{4\pi} \Lambda ,
  \]
  where \( k_2 = k_3 = 1, k_1 = 5/3 \), and \( N_{5,10} \) is the number of \( 5 + \bar{5} \) plus three times the number of \( 10 + \bar{10} \) messenger representations. \( N_{5,10} \leq 4 \) is required to avoid Landau poles. The gaugino mass ratios are the same as found for universal boundary conditions at \( M_U \) in mSUGRA models.

- The squarks/sleptons acquire masses-squared at two-loops: \( m^2_i(M) = \)
  \[
  2\Lambda^2 N_{5,10} f \left( \frac{\Lambda}{M} \right) \left[ c_3 \left( \frac{\alpha_3(M)}{4\pi} \right)^2 + c_2 \left( \frac{\alpha_2(M)}{4\pi} \right)^2 + \frac{5}{3} \left( \frac{Y}{2} \right)^2 \left( \frac{\alpha_1(M)}{4\pi} \right)^2 \right] ,
  \]
  with \( c_3 = 4/3 \) (triplets) \( c_2 = 3/4 \) (weak doublets), \( Y/2 = Q - T_3 \). Degeneracy among families is broken only by effects of order quark or lepton Yukawa couplings, implying no FCNC problems in the first two generations.

In the above, \( g(\Lambda/M), f(\Lambda/M) \to 1 \) for \( M/\Lambda \) not too near 1. For \( \Lambda = M \), \( g(1) = 1.4, f(1) = 0.7 \).

- Spontaneous SUSY-breaking leads to a goldstone fermion, the goldstino, \( G \). (In local SUSY, \( G \) is the longitudinal component of the gravitino.) \( G \) acquires mass determined by \( F \), the goldstino decay constant, where \( F \) is the largest \( F \)-term vev (here, the \( F \) of the secluded, true supersymmetry breaking sector):
  \[
  m_G = \frac{F}{3 M_{\text{Planck}}} \sim 2.5 \left( \frac{\sqrt{F}}{100 \text{ TeV}} \right)^2 \text{eV} ;
  \]
  As sketched earlier, \( \sqrt{F} \sim 10^8 \) TeV is appropriate in the usual mSUGRA/string-motivated models; the \( G \) is then fairly massive and, since it is also very weakly interacting, it is irrelevant to low-energy physics. In GMSB models, the much smaller \( \sqrt{F} \sim 100 - 1000 \) TeV values envisioned imply that the \( G \) will be the lightest supersymmetric particle. In this case, a very crucial constraint is that \( m_G \) be \( \lesssim 1 \) keV to avoid overclosing the universe; i.e. \( \sqrt{F} \lesssim 2000 \) TeV is required for consistency with cosmology.
Some useful observations affecting the phenomenology of a GMSB model are the following.

- FCNC etc. problems are solved since the gauge interactions are flavor-blind and, thus, so are the soft terms.

- As already noted, if $\hat{S}$ is an SU(5) singlet, then we obtain the usual mSUGRA prediction: $M_3 : M_2 : M_1 = 7 : 2 : 1$ at scale $\sim m_Z$.

- Eq. (4) implies
  \[
  \Lambda = \frac{\langle F_S \rangle}{\langle S \rangle} \sim \frac{80 \text{ TeV}}{N_{5,10}} \left( \frac{M_1}{100 \text{ GeV}} \right).
  \tag{7}
  \]

- Eqs. (4) and (5) imply (for $g = f = 1$)
  \[
  m_{\tilde{q}} : m_{\tilde{\tau}_L} : m_{\tilde{\tau}_R} : M_1 = 11.6 : 2.5 : 1.1 : \sqrt{N_{5,10}},
  \tag{8}
  \]
  which in turns implies that the lightest of the standard sparticles (referred to as the next-to-lightest supersymmetric particle, denoted NLSP — the goldstino being the LSP) is the $\tilde{B}$ for $N_{5,10} = 1$ or the $\tilde{\ell}_R$ for $N_{5,10} \geq 2$.

- $\mu$ and $B$ do not arise from the GMSB ansatz and are ‘model-dependent’, but EWSB driven by negative $m^2_{H^2}$ turns out to be completely automatic.

- The couplings of the $G$ are fixed by a supersymmetric Goldberger-Treiman relation:
  \[
  \mathcal{L} = \frac{1}{F} j^{\mu\alpha} \partial_\mu G_\alpha + h.c.
  \]
  where $j$ is the supercurrent connecting a SM particle to its superpartner. For $\sqrt{F}$ values in the range of interest, this coupling is very weak, implying that all the superparticles other than the NLSP undergo chain decay down to the NLSP. The NLSP finally decays to the $G$: e.g. $\tilde{B} \rightarrow \gamma G$ (and $ZG$ if $m_{\tilde{B}} > m_Z$) or $\tilde{\ell}_R \rightarrow \ell G$. The $c\tau$ for NLSP decay depends on $\sqrt{F}$; e.g. for $N_{5,10} = 1$

  \[
  (c\tau)_{\chi_0^0 = \tilde{B} \rightarrow \gamma G} \sim 130 \left( \frac{100 \text{ GeV}}{M_{\tilde{B}}} \right)^5 \left( \frac{\sqrt{F}}{100 \text{ TeV}} \right)^4 \mu\text{m}.
  \tag{9}
  \]

If $\sqrt{F} \sim 2000$ TeV (the upper limit from cosmology), then $c\tau \sim 21\mu$m for $M_{\tilde{B}} = 100$ GeV; $\sqrt{F} \sim 100$ TeV implies a short but vertexable decay length.
If the final decay of the NLSP occurs rapidly ($\sqrt{F} \sim 100$ TeV), then there are many highly observable signatures for a GMSB model [50]. For example, GMSB with the $\tilde{\chi}_0^0$ being the NLSP could explain the Tevatron $ee\gamma\gamma$ event as $\tilde{e}\tilde{e} \rightarrow e\tilde{\chi}_1^0 e\tilde{\chi}_1^0 \rightarrow e\gamma\gamma GG$ (the $G$'s yielding $E/\tau$). However, I will now argue that $\sqrt{F}$ is most naturally very near the $\sqrt{F} \sim 2000$ TeV upper bound, in which case rather few of the NLSP’s decay inside the detector. In fact, in their current incarnation the GMSB models have a significant problem of scale related to the value of $f \equiv F/\langle F_s \rangle$. As noted earlier, the communication between the true supersymmetry breaking sector and the messenger sector occurs at two-loops, implying

$$f \sim \left( \frac{g_m^2}{16\pi^2} \right)^{-2} \sim 2.5 \times 10^4 / g_m^4,$$

where $g_m$ refers to the gauge group responsible for SUSY-breaking, whereas Eqs. (3) and (7) imply

$$f = \frac{F}{\langle F_s \rangle} = \left( \frac{\sqrt{F}}{2000 \text{ TeV}} \right)^2 \frac{625\lambda N_{5,10}^2}{(M/\Lambda)} \left( \frac{100 \text{ GeV}}{M_1} \right)^2.$$  \hspace{2cm} (10)

The problem is that a value as large as $f \sim 2.5 \times 10^4 / g_m^4$ is generally rather inconsistent with basic phenomenological constraints if (as hoped) $g_m \lesssim 1$.

- To illustrate, consider $\lambda \sim 1, g_m = 1$ (i.e. $f = 2.5 \times 10^4$) and $M/\Lambda \sim 1$. (Recall $M/\Lambda > 1$ is required; the choice $M/\Lambda \sim 1$ minimizes the scale problem.) If $\sqrt{F} \sim 2000$ TeV, i.e. as large as possible consistent with $m_G \lesssim 1$ keV, then Eq. (10) implies: $M_1 \sim 16$ GeV (63 GeV) for $N_{5,10} = 1$ ($N_{5,10} = 4$). The former is experimentally ruled out. The latter might be acceptable, but for $N_{5,10} = 4$ the NLSP is the $\tilde{\ell}_R$ with $m_{\tilde{\ell}_R} \sim 35$ GeV, see Eq. (8), which is ruled out by $Z$ data. Although the inconsistency in this latter case is not very bad and could be resolved by modest increases in $\lambda$ and/or $g_m$, or corrections to our simple two-loop estimate for $f$, the $\tilde{\ell}_R$ NLSP would appear as a heavily ionizing track of substantial length in the detector when $\sqrt{F}$ is as large as assumed. Most probably such events would have been observed at the Tevatron in the $m_{\tilde{\ell}_R} \lesssim 100$ GeV range roughly consistent with $f \sim 2.5 \times 10^4$.

- A value of $\sqrt{F} \sim 100$ TeV, as taken in many phenomenological discussions and studies [50], is highly inconsistent with the two-loop expectations for $f$. (This has also been noted in Ref. [51], where it was used to motivate attempts to construct models in which the supersymmetry breaking sector and the messenger sector are one and the same.) For example, taking $\sqrt{F} = 100$ TeV, $M_1 = 100$ GeV, $M/\Lambda \sim 1, \lambda \sim 1$ and $N_{5,10} = 1$ ($N_{5,10} = 4$) in Eq. (10) results in $f = 1.56$ ($f = 25$).

- If an acceptable model with 1-loop communication between the SUSY-breaking sector and the messenger sector can be constructed, we would predict $f \sim 16\pi^2 / g_m^2 \sim 160 / g_m^2$. For $g_m = 1$ (i.e. $f = 16\pi^2$),
$M_1 = 100 \text{ GeV, } M/A \sim 1, \lambda \sim 1$ and $N_{5,10} = 1 (N_{5,10} = 4)$, Eq. (10) yields $\sqrt{F} \sim 1000 \text{ TeV} \left( \sqrt{F} \sim 250 \text{ TeV} \right)$. The associated NLSP lifetimes, see e.g. Eq. (9), would lead to easily detected vertices in events where sparticles are produced. Analysis of existing Tevatron data and forthcoming LEP2 data should readily uncover supersymmetric signals.

In short, for models in which the SUSY-breaking sector communicates at two (or even one) loops with the messenger sector, it seems to be inconsistent to take $\sqrt{F} \lesssim 100 \text{ TeV}$. This implies that the simple phenomenology in which the NLSP decays ($\tilde{\chi}_1^0 \to \gamma G$ or $\tilde{\ell}_R \to \ell G$) almost immediately is not relevant, and the GMSB explanation of the Tevatron $ee\gamma\gamma$ event would be very unlikely. The large $\sqrt{F}$ values required for even a modicum of consistency imply that experimentalists should be looking for events with vertices a substantial distance from the interaction point, quite possibly in association with heavily ionizing tracks. In the most natural case, where $\sqrt{F} \sim 1000 - 2000 \text{ TeV}$ and $m_{\tilde{\chi}_1^0}$ and $m_{\tilde{\ell}_R}$ are below 100 GeV, $c\tau$ — see e.g. Eq. (9) — is typically many tens to several hundreds of meters and only a fraction $\sim 2R/[(\gamma)c\tau]$ of the events will have at least one vertex inside a detector of radius $R$.

- If the NLSP is the $\tilde{\ell}_R$, the bulk of SUSY events will have several heavily ionizing tracks in the detector, a small fraction of which will suddenly terminate with the emission of a $\ell$. Since there is surely no background to such events, this possibility should be excludable using existing Tevatron data (given the significant production rates associated with the low sparticle masses required for scale consistency).

- If the NLSP is the $\tilde{B}$, the bulk of events will not have a vertex inside the detector and will appear as typical $E_T$ supersymmetry signal events; high rates would be required to uncover such events. However, the small, but significant (given the low sparticle masses), number of very distinctive events in which a photon (or $Z$) suddenly emerges in the middle of the electromagnetic or hadronic calorimeter should have a small background, in which case only a few events would be required for discovery. It would seem that this GMSB scenario might also be ruled out or confirmed with proper analysis of Tevatron, if not LEP2, data.

If the above kinds of events, consistent with the most natural $\sqrt{F} \sim 1000 - 2000 \text{ TeV}$ values are found, it will be useful to construct ‘far-out’ additions to the CDF and D0 detectors designed to reveal more of the decay vertices.

The cosmological implications/consistency of GMSB constitute an important issue [52,53], but one that is far too complex to elaborate on significantly here. I only note that if $\sqrt{F} \sim 100 \text{ TeV}$, as disfavored by the scale consistency discussed above, then $m_G \sim 2.5 \text{ eV}$ implies that the goldstino cannot make a significant contribution to the dark matter of the universe; however, the messenger sector might contain an appropriate dark matter candidate [52].
In a model with $\sqrt{F} \sim 1000$ TeV, as preferred by the scale arguments, the $G$ gives rise to a cosmologically significant abundance of warm dark matter; unfortunately, it would be invisible in halo detection experiments [52].

CONCLUSIONS

If supersymmetry is discovered, it will be a dream-come-true for both theorists and experimentalists. For theorists, it would a a triumph of aesthetic principles, naturalness, etc. For experimentalists, it would be a gold mine of experimental signals and analyses. As has always been the case in the past, the next step in theory will require experimental guidance and input. The many phenomenological manifestations and parameters of supersymmetry imply that many years of experimental work will be required before it will be possible to determine the precise nature of supersymmetry breaking and the associated boundary conditions. Our ultimate dream is that, armed with this information, we will be able to construct the ‘final’ theory.

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