Low-Energy Supersymmetry and its Phenomenology

Howard E. Haber[a*]

[a]Santa Cruz Institute for Particle Physics
University of California, Santa Cruz, CA 95064, USA

The structure of low-energy supersymmetric models of fundamental particles and interactions is reviewed, with an emphasis on the minimal supersymmetric extension of the Standard Model (MSSM) and some of its variants. Various approaches to the supersymmetry-breaking mechanism are considered. The implications for the phenomenology of Higgs bosons and supersymmetric particles at future colliders are discussed.

1. INTRODUCTION

The Standard Model provides a remarkably successful description of the properties of the quarks, leptons and spin-1 gauge bosons at energy scales of $\mathcal{O}(100)$ GeV and below. However, the Standard Model is not the ultimate theory of the fundamental particles and their interactions. At an energy scale above the Planck scale, $M_{\rm P} \simeq 10^{19}$ GeV, quantum gravitational effects become significant and the Standard Model must be replaced by a more fundamental theory that incorporates gravity. It is also possible that the Standard Model breaks down at some energy scale, Λ, below the Planck scale. In this case, the Standard Model degrees of freedom are no longer adequate for describing the physics above Λ and new physics must enter. Thus, the Standard Model is not a *fundamental* theory; at best, it is an *effective field theory* [1]. At an energy scale below Λ, the Standard Model (with higher-dimension operators to parameterize the new physics at the scale Λ) provides an extremely good description of all observable phenomena.

In an effective field theory, all parameters of the low-energy theory (*i.e.* masses and couplings) are calculable in terms of parameters of a more fundamental, renormalizable theory that describes physics at the energy scale Λ. All low-energy couplings and fermion masses are logarithmically sensitive to Λ. In contrast, scalar squared-masses are *quadratically* sensitive to Λ. The Higgs mass

(at one-loop) has the following form:

$$m_h^2 = (m_h^2)_0 + \frac{cg^2}{16\pi^2}\Lambda^2\,, \qquad (1)$$

where $(m_h^2)_0$ is a parameter of the fundamental theory and c is a constant, presumably of $\mathcal{O}(1)$, which is calculable within the low-energy effective theory. The "natural" value for the scalar squared-mass is $g^2\Lambda^2/16\pi^2$. Thus, the expectation for Λ is

$$\Lambda \simeq \frac{4\pi m_h}{g} \sim \mathcal{O}(1\ {\rm TeV})\,. \qquad (2)$$

If Λ is significantly larger than 1 TeV (often called the hierarchy and naturalness problem in the literature [2]), then the only way for the Higgs mass to be of order the scale of electroweak symmetry breaking is to have an "unnatural" cancellation between the two terms of eq. (1). This seems highly unlikely given that the two terms of eq. (1) have completely different origins.

A viable theoretical framework that incorporates weakly-coupled Higgs bosons and satisfies the constraint of eq. (2) is that of "low-energy" or "weak-scale" supersymmetry [3–5]. In this framework, supersymmetry is used to relate fermion and boson masses and interaction strengths. Since fermion masses are only logarithmically sensitive to Λ, boson masses will exhibit the same logarithmic sensitivity if supersymmetry is exact. Since no supersymmetric partners of Standard Model particles have yet been found, supersymmetry cannot be an exact

symmetry of nature. Thus, Λ should be identified with the supersymmetry-breaking scale. The naturalness constraint of eq. (2) is still relevant, so in the framework of low-energy supersymmetry, the scale of supersymmetry breaking should not be much larger than about 1 TeV in order that the naturalness of scalar masses be preserved. The supersymmetric extension of the Standard Model would then replace the Standard Model as the effective field theory of the TeV scale. One advantage of the supersymmetric approach is that the effective low-energy supersymmetric theory *can* be valid all the way up to the Planck scale, while still being natural! The unification of the three gauge couplings at an energy scale close to the Planck scale, which does not occur in the Standard Model, is seen to occur in the minimal supersymmetric extension of the Standard Model, and provides an additional motivation for seriously considering the low-energy supersymmetric framework [6].

2. STRUCTURE OF THE MSSM

The minimal supersymmetric extension of the Standard Model (MSSM) consists of taking the Standard Model and adding the corresponding supersymmetric partners. (For a review of the structure of the MSSM, see *e.g.*, refs. [4,5].) In addition, the MSSM contains two hypercharge $Y = \pm 1$ Higgs doublets (H_u and H_d, respectively), which is the minimal structure for the Higgs sector of an anomaly-free supersymmetric extension of the Standard Model. The supersymmetric structure of the theory also requires (at least) two Higgs doublets to generate mass for both "up"-type and "down"-type quarks (and charged leptons) [7,8]. The supersymmetric spectrum can be described by a set of superfields that incorporates both the Standard Model fields and their superpartners as shown in Table 1. Note that Table 1 lists the interaction eigenstates. Particles with the same $\mathrm{SU}(3) \times \mathrm{U}(1)_{\mathrm{EM}}$ quantum numbers can mix. The physical (mass) eigenstates are determined from the full interaction Lagrangian of the theory. For example, the physical *charginos* $[\widetilde{\chi}_j^{\pm} \ (j = 1, 2)]$ are linear combinations of the charged winos $(\widetilde{W}^{\pm})$ and charged higgsi-

Table 1
The MSSM Particle Spectrum

Superfield	Boson Fields	Fermionic Partners
Gauge Multiplets		
$\widehat{G}$	g	$\widetilde{g}$
$\widehat{V}^a$	W^a	$\widetilde{W}^a$
$\widehat{V}'$	B	$\widetilde{B}$
Matter Multiplets		
$\widehat{L}$ leptons $\begin{cases} \widetilde{L}^j = (\widetilde{\nu}, \widetilde{e}^-)_L \\ \widetilde{E} = \widetilde{e}_R^+ \end{cases}$		$(\nu, e^-)_L$ e_L^c
$\widehat{Q}$ $\widehat{U}$ quarks $\begin{cases} \widetilde{Q}^j = (\widetilde{u}_L, \widetilde{d}_L) \\ \widetilde{U} = \widetilde{u}_R^* \\ \widetilde{D} = \widetilde{d}_R^* \end{cases}$		$(u, d)_L$ u_L^c d_L^c
$\widehat{H}_d$ $\widehat{H}_u$ Higgs $\begin{cases} H_d^i \\ H_u^i \end{cases}$		$(\widetilde{H}_d^0, \widetilde{H}_d^-)_L$ $(\widetilde{H}_u^+, \widetilde{H}_u^0)_L$

nos $(\widetilde{H}_u^+, \widetilde{H}_d^-)$, while the physical *neutralinos* $[\widetilde{\chi}_k^0 \ (k = 1, \ldots, 4)]$ are linear combinations of the neutral wino $(\widetilde{W}^3)$, bino $(\widetilde{B})$ and neutral higgsinos $(\widetilde{H}_u^0, \widetilde{H}_d^0)$.

The renormalizable supersymmetric interactions are fixed once the superpotential (a cubic polynomial of superfields) is chosen. In the MSSM, the most general $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ invariant cubic superpotential is given by $W \equiv W_{\mathrm{RPC}} + W_{\mathrm{RPV}}$, where

$$W_{\mathrm{RPC}} = (y_e)_{mn} \widehat{H}_d \widehat{L}_m \widehat{E}_n + (y_d)_{mn} \widehat{H}_d \widehat{Q}_m \widehat{D}_n$$
$$- (y_u)_{mn} \widehat{H}_u \widehat{Q}_m \widehat{U}_n - \mu \widehat{H}_d \widehat{H}_u, \quad (3)$$

and

$$W_{\mathrm{RPV}} = (\lambda_L)_{mnp} \widehat{L}_m \widehat{L}_n \widehat{E}_p + (\lambda_L')_{mnp} \widehat{L}_m \widehat{Q}_n \widehat{D}_p$$
$$+ (\lambda_B)_{mnp} \widehat{U}_m \widehat{D}_n \widehat{D}_p - (\mu_L)_n \widehat{L}_n \widehat{H}_u. \quad (4)$$

In eqs. (3) and (4), m, n and p label the generation number, and SU(2) doublet fields are contracted by the invariant antisymmetric tensor ϵ_{ij} (with $\epsilon_{12} = 1$). For example, the Yukawa interactions are obtained from the trilinear terms

of eq. (3) by taking all possible combinations involving two fermions and one scalar superpartner. The gauge multiplets couple to matter multiplets in a manner consistent with supersymmetry and the SU(3)×SU(2)×U(1) gauge symmetry.

In contrast to the Standard Model, the terms of dimension ≤ 4 in the supersymmetric Lagrangian do not automatically preserve baryon number (B) and lepton number (L). In order to be consistent with the (approximately) conserved baryon number and lepton number observed in nature, one must impose constraints on the terms in the superpotential [eqs. (3) and (4)]. In particular, the terms of W_{RPC} conserve B and L, whereas the terms of W_{RPV} violate either B or L as indicated by the subscript of the corresponding coefficients. If all terms in W_{RPV} were present, some of the coefficients would have to be extremely small in order to avoid proton decay at a catastrophic rate.

In the MSSM, one enforces an approximate lepton and baryon number invariance by imposing R-parity-invariance. Any particle of spin S, baryon number B and lepton number L then possess a multiplicatively-conserved R-parity given by $R = (-1)^{3(B-L)+2S}$. Then, all operators in the R-parity-conserving (RPC) model of dimension ≤ 4 exactly conserve B and L. In the RPC model, R-parity conservation is equivalent to imposing a $\mathbf{Z_2}$ matter parity on the superpotential according to the quantum numbers specified in Table 2, which results in $W_{\mathrm{RPV}} = 0$. Thus, the supersymmetric interactions of the MSSM depend on the following parameters:

- SU(3)×SU(2)×U(1) gauge couplings, g_3, $g_2 \equiv g$, and $g_1 \equiv \sqrt{\frac{5}{3}}\, g'$;
- (complex) Higgs-fermion Yukawa coupling matrices, $\mathbf{y_u}$, $\mathbf{y_d}$, and $\mathbf{y_e}$; and
- a supersymmetric Higgs mass parameter, μ.

Since the MSSM is a model of three generations of quarks, leptons and their superpartners, $\mathbf{y_u}$, $\mathbf{y_d}$, and $\mathbf{y_e}$ are complex 3×3 matrices. However, not all these degrees of freedom are physical, as emphasized below.

Note that in the (minimal version of the) Standard Model and in the MSSM, neutrinos are ex-

Table 2
Matter discrete symmetries

symmetry	$\widehat{Q}_n$	$\widehat{U}_n$	$\widehat{D}_n$	$\widehat{L}_n$	$\widehat{E}_n$	$\widehat{H}_U$	$\widehat{H}_D$
$\mathbf{Z_2}$	-1	-1	-1	-1	-1	$+1$	$+1$
$\mathbf{Z_3}$	ω	ω^{-1}	ω^{-1}	$+1$	$+1$	$+1$	$+1$

$\omega \equiv e^{2i\pi/3}$

actly massless if only terms of dimension ≤ 4 in the Lagrangian are considered. Majorana neutrino masses may be generated in either theory by introducing the appropriate dimension-5 gauge invariant $\Delta L = 2$ terms. In the MSSM, such terms can be realized via an RPC supersymmetric extension of the seesaw mechanism, which would require the addition of a SU(3)×SU(2)×U(1) singlet superfield [9].

One can also avoid B violation in an R-parity-violating (RPV) theory by imposing a discrete $\mathbf{Z_3}$ triality shown in Table 2. This is the unique choice for a (generation-independent) discrete symmetry with no discrete gauge anomalies in a model consisting only of the MSSM superfields [10]. In particular, the $\mathbf{Z_3}$ discrete symmetry requires that $\lambda_B = 0$ in eq. (4) [in fact, all operators of dimension ≤ 5 preserve B], while all other terms in the superpotential, including all possible L-violating terms, are permitted. Hence, one must check that the magnitude of the L-violating interactions is consistent with experimental bounds on L-violating processes [11]. One good feature of this model is that it provides a mechanism for non-zero neutrino masses without requiring the introduction of a new superfield beyond those already contained in the MSSM.

At this stage, one does not yet have a realistic model of fundamental particles and their interactions, since supersymmetry (SUSY) is unbroken. However, the fundamental origin of SUSY-breaking is not known at present. Without a fundamental theory of SUSY-breaking, the best we can do is to parameterize our ignorance and introduce the most general renormalizable soft SUSY-breaking terms [12] consistent with the SU(3)×SU(2)×U(1) gauge symmetry and any additional discrete matter symmetries that have been imposed. It is here where most of the

new supersymmetric model parameters reside.[1] If supersymmetry is relevant for explaining the scale of electroweak interactions, then the mass parameters introduced by the soft SUSY-breaking terms must be of order 1 TeV or below [14].

In the MSSM, the soft-SUSY-breaking terms consist of [2] (following the notation of ref. [5]):

$$
\begin{aligned}
V_{\text{soft}} =\ & m_{H_d}^2 |H_d|^2 + m_{H_u}^2 |H_u|^2 - (b H_d H_u + \text{h.c.}) \\
& + (m_Q^2)_{mn} \widetilde{Q}_m^{i*} \widetilde{Q}_n^i + (m_{\overline{u}}^2)_{mn} \widetilde{U}_m^* \widetilde{U}_n \\
& + (m_{\overline{d}}^2)_{mn} \widetilde{D}_m^* \widetilde{D}_n \\
& + (m_L^2)_{mn} \widetilde{L}_m^{i*} \widetilde{L}_n^i + (m_{\overline{e}}^2)_{mn} \widetilde{E}_m^* \widetilde{E}_n \\
& + \big[(a_e)_{mn} \widetilde{H}_d \widetilde{L}_m \widetilde{E}_n + (a_d)_{mn} \widetilde{H}_d \widetilde{Q}_m \widetilde{D}_n \\
& \qquad - (a_u)_{mn} \widetilde{H}_u \widetilde{Q}_m \widetilde{U}_n + \text{h.c.} \big] \\
& + \tfrac{1}{2} \big[M_3\, \widetilde{g}\,\widetilde{g} + M_2 \widetilde{W}^a \widetilde{W}^a + M_1 \widetilde{B}\widetilde{B} + \text{h.c.} \big] \ . \quad (5)
\end{aligned}
$$

In the RPV model, additional A and b parameters can be added, corresponding to the additional terms that appear in the superpotential.[3]

The Higgs scalar potential receives contributions from both the SUSY-conserving and the SUSY-breaking sector. The SU(2)×U(1) electroweak symmetry is spontaneously broken only when the terms of V_{soft} [eq. (5)] are included. The neutral Higgs fields acquire vacuum expectation values: $\left\langle H_d^0 \right\rangle = v_d/\sqrt{2}$ and $\left\langle H_u^0 \right\rangle = v_u/\sqrt{2}$, where $v^2 \equiv v_d^2 + v_u^2 = (246\ \text{GeV})^2$ is determined by the Z mass, and $\tan\beta \equiv v_u/v_d$ is a free parameter.

Not all of the parameters appearing in eqs. (3) and (5) represent independent degrees of freedom. By suitable redefinitions of the various fields of the model, one can remove all unphysical degrees of freedom and identify the correct

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[1]It was argued in ref. [13] that additional dimension-3 SUSY-breaking terms, which were *not* characterized as "soft" by ref. [12], should be considered in models such as the MSSM that contain no gauge-singlet superfields.

[2]Here, we omit the more general dimension-3 SUSY-breaking terms advocated in ref. [13]. In addition, we note that dimension-4 SUSY-breaking terms can also be generated by high-energy-scale physics, albeit with very small coefficients [15].

[3]In the literature, a different matrix A-parameter is often defined via the relation: $\mathbf{a_f} \equiv \mathbf{y_f A_f}$ [$\mathbf{f = u,\ d,\ e}$], where the $\mathbf{y_f}$ are the Higgs-fermion Yukawa coupling matrices, and a B-parameter is defined by $b \equiv \mu B$, where μ is the supersymmetric Higgs mass parameter.

Table 3
Parameter Count in the Standard Model and MSSM Extensions

model	real	phases	TOTAL
SM	17	2	19
MSSM	79	45	124
$(\text{MSSM})_{\text{B}}$	157	122	279
$(\text{MSSM})_{\text{RPV}}$	175	140	315

number of physical parameters of the model [16]. For example, the MSSM is characterized by 124 independent parameters, of which 18 correspond to Standard Model parameters (including θ_{QCD}), one corresponds to a Higgs sector parameter (the analogue of the Standard Model Higgs mass), and 105 are genuinely new parameters of the model. All together, among the various complex parameters of the model, there are 45 phases that cannot be removed; the remaining 79 real parameters consist of masses, real couplings and mixing angles. Thus, an appropriate name for the minimal supersymmetric extension of the Standard Model described above is MSSM-124 [17]. In Table 3, we compare the parameter count of the Standard Model with three possible versions of low-energy supersymmetry based on the fields of the MSSM. The RPC model is denoted in Table 3 simply by MSSM, whereas the RPV model with a $\mathbf{Z_3}$ triality that preserves B is denoted by $(\text{MSSM})_{\text{B}}$. Finally, $(\text{MSSM})_{\text{RPV}}$ denotes the most general RPV model in which all terms in W_{RPV} [eq. (4)] and the corresponding soft SUSY-breaking terms are allowed.

Even in the absence of a fundamental theory of SUSY-breaking, one is hard-pressed to regard MSSM-124 as a fundamental theory. For example, no fundamental explanation is provided for the origin of electroweak symmetry breaking. Moreover, MSSM-124 is not a phenomenologically viable theory over most of its parameter space. Among the phenomenological deficiencies are: (i) no conservation of the separate lepton numbers L_e, L_μ, and L_τ; (ii) unsuppressed flavor-changing neutral currents (FCNC's) [18];

and (iii) new sources of CP-violation that are inconsistent with the experimental bounds [19,20]. As a result, almost the entire MSSM-124 parameter space is ruled out! This theory is viable only at very special "exceptional" points of the full parameter space.

A truly Minimal SSM does not (yet) exist. In the present usage, the word "minimal" in MSSM refers to the minimal particle spectrum and the associated R-parity invariance. The MSSM particle content must be supplemented by assumptions about the origin of SUSY-breaking that lie outside the low-energy domain of the model. Moreover, a comprehensive map of the phenomenologically acceptable region of MSSM-124 parameter space does not yet exist. This presents a formidable challenge to supersymmetric particle searches that must impose some parameter constraints while trying to ensure that the search is as complete as possible. Ultimately, the goal of any supersymmetric particle search is to measure as many of the 124 parameters as is feasible and to determine any additional parameters that would characterize a possible departure from the minimal SSM structure.

3. REDUCING THE MSSM PARAMETER FREEDOM

There are two general approaches for reducing the parameter freedom of MSSM-124. In the low-energy (or "bottom-up") approach, an attempt is made to elucidate the nature of the exceptional points in the MSSM-124 parameter space that are phenomenologically viable. Consider the following two possible choices. First, one can assume that the squark and slepton squared-mass matrices and the matrix A-parameters, $\mathbf{A_f}$, are proportional to the 3×3 unit matrix (horizontal universality [21,22,16]). Alternatively, one can simply require that all the aforementioned matrices are flavor diagonal in a basis where the quark and lepton mass matrices are diagonal (flavor alignment [23]). In these approaches, the number of free parameters characterizing the MSSM is substantially less than 124. Moreover, L_e, L_μ and L_τ are separately conserved, while tree-level FCNC's are automatically absent. The resulting models

are phenomenologically viable, although there is no strong theoretical basis for either approach.

In the high-energy (or "top-down") approach, the MSSM-124 parameter freedom is reduced by imposing theoretical constraints on the structure of SUSY-breaking. It is very difficult (perhaps impossible) to construct a model of low-energy supersymmetry where the SUSY-breaking arises solely as a consequence of the interactions of the particles of the MSSM. A more viable scheme posits a theory in which the fundamental source of SUSY-breaking originates in a sector that is distinct from the fields that make up the MSSM. For lack of a better term, we will call this new sector the direct SUSY-breaking (DSB) sector.[4] The SUSY-breaking inherent in the DSB-sector is subsequently transmitted to the MSSM spectrum by some mechanism.

Integrating out the physics associated with the DSB sector, one obtains initial conditions for the MSSM parameters at some high energy scale. Renormalization group (RG) evolution then allows one to evolve down to the electroweak scale and derive the full MSSM particle spectrum. One bonus of this analysis is that one of the diagonal Higgs squared-mass parameters is typically driven negative by the RG-evolution. Thus, electroweak symmetry breaking is generated radiatively, and the resulting electroweak symmetry-breaking scale is intimately tied to the scale of low-energy SUSY-breaking.

Two theoretical scenarios have been examined in detail: gravity-mediated and gauge-mediated SUSY-breaking. In both these approaches, supersymmetry is spontaneously broken, in which case a massless Goldstone fermion, the *goldstino*, arises. Its coupling to a particle and its superpartner is fixed by the supersymmetric Goldberger-Treiman relation [25]

$$\mathcal{L}_{\text{int}} = -\frac{1}{F} j^{\mu\alpha} \partial_\mu \widetilde{G}_\alpha + \text{h.c.}, \qquad (6)$$

[4]In the literature, the DSB-sector usually refers to the sector of dynamical SUSY-breaking. Here, we shall interpret the word *dynamical* in its broadest sense. Dynamical SUSY-breaking can be non-perturbative in nature (*e.g.*, gaugino condensation [24]) or perturbative in nature. Examples of the latter include tree-level O'Raifeartaigh (or F-type) breaking and Fayet-Iliopoulos (or D-type) breaking. For additional details, see ref. [5].

where $j^{\mu\alpha}$ is the supercurrent, which depends bilinearly on all the fermion–boson superpartner pairs of the theory and $\widetilde{G}_\alpha$ is the spin-1/2 goldstino field (with spinor index α). In particular, $\sqrt{F}$ is the scale of direct SUSY-breaking which occurs in the DSB-sector (typically, $\sqrt{F} \gg m_Z$). When gravitational effects are included, the goldstino is "absorbed" by the *gravitino* ($\widetilde{g}_{3/2}$), the spin-3/2 partner of the graviton. By this super-Higgs mechanism [26], the goldstino is removed from the physical spectrum and the gravitino acquires a mass ($m_{3/2}$). In models where the gravitino mass is generated at tree-level (see, *e.g.*, ref. [27] for further discussion), one finds:

$$m_{3/2} = \frac{F}{\sqrt{3}M_\mathrm{P}}\,, \tag{7}$$

where M_P is the reduced Planck mass. The helicity $\pm\frac{1}{2}$ components of the gravitino behave approximately like the goldstino, whose couplings to particles and their superpartners are determined by eq. (6). In particular, the goldstino couplings are enhanced by a factor of M_P^2/F relative to couplings of gravitational strength. In contrast, the helicity $\pm\frac{3}{2}$ components of the gravitino always couple with gravitational strength to particles and their superpartners, and thus can be neglected in phenomenological studies.

In many models, the DSB-sector is comprised of fields that are completely neutral with respect to the Standard Model gauge group. In such cases, the DSB-sector is also called the "hidden sector." The fields of the MSSM are said to reside in the "visible sector," and the model is constructed such that no renormalizable tree-level interactions exist between fields of the visible and hidden sectors. A third sector, the so-called "messenger sector," is often employed in models to transmit the SUSY-breaking from the hidden sector to the visible sector. However, it is also possible to construct models in which the DSB-sector is not strictly hidden and contains fields that are charged with respect to the Standard Model gauge group.

3.1. Gravity-mediated SUSY-breaking

All particles feel the gravitational force. In particular, particles of the hidden sector and the visible sector can interact via the exchange of gravitons. Thus, supergravity (SUGRA) models provide a natural mechanism for transmitting the SUSY-breaking of the hidden sector to the particle spectrum of the MSSM. In models of gravity-mediated SUSY-breaking, Planck-scale physics is the messenger of SUSY-breaking [28,29].

In the *minimal* supergravity (mSUGRA) framework [3], the soft SUSY-breaking parameters at the Planck scale take a particularly simple form:

$$\begin{aligned}
\mathbf{m_Q^2}(M_\mathrm{P}) &= \mathbf{m_{\bar{u}}^2}(M_\mathrm{P}) = \mathbf{m_{\bar{d}}^2}(M_\mathrm{P}) \\
&= \mathbf{m_L^2}(M_\mathrm{P}) = \mathbf{m_{\bar{e}}^2}(M_\mathrm{P}) = m_0^2\mathbf{1}\,, \\
m_{H_u}^2(M_\mathrm{P}) &= m_{H_d}^2(M_\mathrm{P}) = m_0^2\,, \\
\mathbf{a_f}(M_\mathrm{P}) &= A_0\mathbf{y_f}(M_\mathrm{P})\,, \qquad \mathbf{f} = \mathbf{u, d, e}\,, \tag{8}
\end{aligned}$$

where $\mathbf{1}$ is the 3×3 identity matrix in generation space. Note that the last condition is equivalent to $\mathbf{A_f}(M_\mathrm{P}) = A_0\mathbf{1}$. Thus, the 3×3 matrices above respect horizontal universality at the Planck scale.

In addition, the gauge couplings and gaugino mass parameters are assumed to unify at some high energy scale, M_X. (The latter condition is automatic in models of supersymmetric grand unification, where M_X is the unification scale.) The unification relation

$$M_1(M_X) = M_2(M_X) = M_3(M_X) = m_{1/2} \tag{9}$$

implies that the low-energy gaugino mass parameters (approximately) satisfy:

$$M_3 = \frac{g_3^2}{g_2^2}M_2 \simeq 3.5M_2\,,$$

$$M_1 = \tfrac{5}{3}\tan^2\theta_W M_2 \simeq 0.5M_2\,. \tag{10}$$

One can count the number of independent parameters in the mSUGRA framework. In addition to 18 Standard Model parameters (excluding the Higgs mass), one must specify m_0, $m_{1/2}$, A_0, and Planck-scale values for μ and B-parameters (denoted by μ_0 and B_0). In principle, A_0, B_0 and μ_0 can be complex, although in the mSUGRA approach, these parameters are taken (arbitrarily) to be real. Renormalization group evolution is used to compute the low-energy values of the mSUGRA parameters, which then fixes all the

parameters of the low-energy MSSM. In particular, the two Higgs vacuum expectation values, v_u and v_d (or equivalently, m_Z and $\tan\beta$) can be expressed as a function of the Planck-scale supergravity parameters. The simplest procedure is to remove μ_0 and B_0 in favor of m_Z and $\tan\beta$ (the sign of μ_0 is not fixed in this process). In this case, the MSSM spectrum and its interaction strengths are determined by five parameters: m_0, A_0, $m_{1/2}$, $\tan\beta$, and the sign of μ_0, in addition to the 18 parameters of the Standard Model. The requirement of radiative electroweak symmetry-breaking imposes an additional constraint on the possible range of the mSUGRA parameters. In particular, one finds that $1 \lesssim \tan\beta \lesssim m_t/m_b$. In principle, one should also include the mass of the gravitino, $m_{3/2}$ (or equivalently, $\sqrt{F}$), in the list of independent parameters. In mSUGRA, one arranges the scale of the hidden sector SUSY-breaking such that $\sqrt{F} \sim 3 \times 10^{10}$ GeV. In this case, the gravitino mass is of order the electroweak symmetry-breaking scale [see eq. (7)], while its couplings to the MSSM fields are extremely weak. Such a gravitino would play no role in supersymmetric phenomenology at colliders.

Recently, attention has been given to a class of supergravity models in which eq. (10) does not hold. In models where no tree-level gaugino masses are generated, one finds a model-independent contribution to the gaugino mass whose origin can be traced to the super-conformal (super-Weyl) anomaly which is common to all supergravity models [30]. This approach has been called *anomaly-mediated* SUSY-breaking (AMSB). The gaugino mass parameters (in the one-loop approximation) are given by:

$$M_i \simeq \frac{b_i g_i^2}{16\pi^2} m_{3/2}\,, \qquad (11)$$

where the b_i are the coefficients of the MSSM gauge beta-functions corresponding to the corresponding U(1), SU(2) and SU(3) gauge groups: $(b_1, b_2, b_3) = (\frac{33}{5}, 1, -3)$. Anomaly-mediated SUSY-breaking also generates (approximate) flavor-diagonal squark and slepton mass matrices. However, in the MSSM this cannot be the sole source of SUSY-breaking in the slepton sector, since the latter yields negative squared-mass con-

tributions for the sleptons. A possible RPV solution, involving only the superfields of the MSSM, has been advocated in ref. [31]

3.2. Gauge-mediated SUSY-breaking

In *gauge-mediated* SUSY-breaking (GMSB), SUSY-breaking is transmitted to the MSSM via gauge forces. A typical structure of such models involves a hidden sector where supersymmetry is broken, a "messenger sector" consisting of particles (messengers) with SU(3)×SU(2)×U(1) quantum numbers, and the visible sector consisting of the fields of the MSSM [32,33]. The direct coupling of the messengers to the hidden sector generates a SUSY-breaking spectrum in the messenger sector. Finally, SUSY-breaking is transmitted to the MSSM via the virtual exchange of the messengers. If this approach is extended to incorporate gravitational phenomena, then supergravity effects will also contribute to SUSY-breaking. However, in models of gauge-mediated SUSY-breaking, one usually chooses the model parameters in such a way that the virtual exchange of the messengers dominates the effects of the direct gravitational interactions between the hidden and visible sectors. In this scenario, the gravitino mass is typically in the eV to keV range, and is therefore the LSP. The helicity $\pm\frac{1}{2}$ components of $\widetilde{g}_{3/2}$ behave approximately like the goldstino; its coupling to the particles of the MSSM is significantly stronger than a coupling of gravitational strength.

In the minimal GMSB approach, there is one effective mass scale, Λ, that determines all low-energy scalar and gaugino mass parameters through loop-effects (while the resulting A-parameters are suppressed). In order that the resulting superpartner masses be of order 1 TeV or less, one must have $\Lambda \sim 100$ TeV. The origin of the μ and B-parameters is model dependent and lies somewhat outside the GMSB ansatz. The simplest models of this type are even more restrictive than mSUGRA, with two fewer degrees of freedom. However, minimal GMSB is not a fully realized model. The sector of SUSY-breaking dynamics can be very complex, and no complete GMSB model yet exists that is both simple and compelling.

4. THE MSSM HIGGS SECTOR

4.1. The Higgs Sector in Low-Energy Supersymmetry

In the MSSM, the Higgs sector is a two-Higgs-doublet model with Higgs self-interactions constrained by supersymmetry [7,8,34,35]. Moreover, in spite of the large number of potential CP-violating phases among the MSSM-124 parameters, the tree-level MSSM Higgs sector is automatically CP-conserving. In particular, unphysical phases can be absorbed into the definition of the Higgs fields such that $\tan\beta$ is real and positive. As a result, the physical neutral Higgs scalars are CP-eigenstates. There are five physical Higgs particles in this model: a charged Higgs pair ($H^{\pm}$), two CP-even neutral Higgs bosons (denoted by h^0 and H^0 where $m_{h^0} \leq m_{H^0}$) and one CP-odd neutral Higgs boson (A^0). At tree level, $\tan\beta$ and one Higgs mass (usually chosen to be m_{A^0}) determine the tree-level Higgs-sector parameters. These include the other Higgs masses, an angle α [which measures the component of the original $Y = \pm 1$ Higgs doublet states in the physical CP-even neutral scalars], the Higgs boson self-couplings, and the Higgs boson couplings to particles of the Standard Model and their superpartners.

When one-loop radiative corrections are incorporated, the Higgs masses and couplings depend on additional parameters of the supersymmetric model that enter via virtual loops. One of the most striking effects of the radiative corrections to the MSSM Higgs sector is the modification of the upper bound of the light CP-even Higgs mass, as first noted in ref. [36]. When $\tan\beta \gg 1$ and $m_{A^0} \gg m_Z$, the *tree-level* prediction for m_{h^0} corresponds to its theoretical upper bound, $m_{h^0}^{\max} = m_Z$. Including radiative corrections, the theoretical upper bound is increased, primarily because of an incomplete cancellation of the top-quark and top-squark (stop) loops (these effects actually cancel in the exact supersymmetric limit). The relevant parameters that govern the stop sector are the average of the two stop squared-masses: $M_{\mathrm{SUSY}}^2 \equiv \frac{1}{2}(M_{\tilde{t}_1}^2 + M_{\tilde{t}_2}^2)$, and the off-diagonal element of the stop squared-mass matrix: $m_t X_t \equiv m_t(A_t - \mu\cot\beta)$. The qualitative behavior of the radiative corrections can be most easily seen in the large top squark mass limit, where in addition the splitting of the two diagonal entries and the off-diagonal entry of the stop squared-mass matrix are both small in comparison to M_{SUSY}^2. In this case, the upper bound on the lightest CP-even Higgs mass is approximately given by

$$m_{h^0}^2 \lesssim m_Z^2 + \frac{3g^2 m_t^4}{8\pi^2 m_W^2}\left[\ln\left(\frac{M_{\mathrm{SUSY}}^2}{m_t^2}\right)\right.$$
$$\left.+\frac{X_t^2}{M_{\mathrm{SUSY}}^2}\left(1-\frac{X_t^2}{12M_{\mathrm{SUSY}}^2}\right)\right]\,, (12)$$

More complete treatments of the radiative corrections include the effects of stop mixing, renormalization group improvement, and the leading two-loop contributions, and imply that eq. (12) somewhat overestimates the true upper bound of m_{h^0} (see ref. [37] for the most recent results). Nevertheless, eq. (12) correctly reflects some noteworthy features of the more precise result. First, the increase of the light CP-even Higgs mass bound beyond m_Z can be significant. This is a consequence of the m_t^4 enhancement of the one-loop radiative correction. Second, the dependence of the light Higgs mass on the stop mixing parameter X_t implies that (for a given value of M_{SUSY}) the upper bound of the light Higgs mass initially increases with X_t and reaches its *maximal* value at $X_t = \sqrt{6}M_{\mathrm{SUSY}}$. This point is referred to as the *maximal mixing* case (whereas $X_t = 0$ corresponds to the *minimal mixing* case).

Taking m_{A^0} large, fig. 1 illustrates that the maximal value of the lightest CP-even Higgs mass bound is realized at large $\tan\beta$ in the case of maximal mixing. Allowing for the uncertainty in the measured value of m_t and the uncertainty inherent in the theoretical analysis, one finds for $M_{\mathrm{SUSY}} \lesssim 2$ TeV that $m_{h^0} \lesssim m_{h^0}^{\max}$, where

$$m_{h^0}^{\max} \simeq 122 \text{ GeV}, \quad \text{minimal stop minimal},$$
$$m_{h^0}^{\max} \simeq 135 \text{ GeV}, \quad \text{maximal stop mixing}. \quad (13)$$

The h^0 mass bound in the MSSM quoted above does not in general apply to non-minimal supersymmetric extensions of the Standard Model. If additional Higgs singlet and/or triplet fields
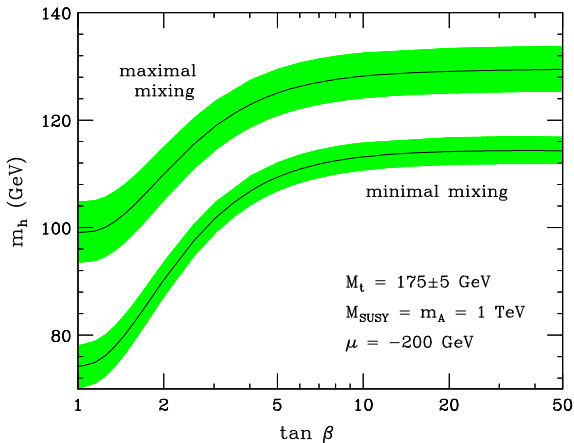
Figure 1. The radiatively corrected light CP-even Higgs mass is plotted as a function of $\tan\beta$, for the maximal mixing [upper band] and minimal mixing cases. The impact of the top quark mass is exhibited by the shaded bands; the central value corresponds to $M_t = 175$ GeV, while the upper [lower] edge of the bands correspond to increasing [decreasing] M_t by 5 GeV.

are introduced, then new Higgs self-coupling parameters appear, which are not significantly constrained by present data. For example, in the simplest non-minimal supersymmetric extension of the Standard Model (NMSSM), the addition of a Higgs singlet adds a new Higgs self-coupling parameter, λ [38]. The mass of the lightest neutral Higgs boson can be raised arbitrarily by increasing the value of λ (analogous to the behavior of the Higgs mass in the Standard Model!). Under the assumption that all couplings stay perturbative up to the Planck scale, one finds in essentially all cases that $m_{h^0} \lesssim 200$ GeV, independent of the details of the low-energy supersymmetric model [39].

4.2. MSSM Higgs searches at colliders

There is presently no definitive experimental evidence for the Higgs boson. Experimental limits on the charged and neutral Higgs masses have been obtained at LEP. For the charged Higgs boson, $m_{H^\pm} > 78.7$ GeV [40]. This is the most

model-independent bound (it is valid for more general non-supersymmetric two-Higgs doublet models) and assumes only that the H^+ decays dominantly into $\tau^+\nu_\tau$ and/or $c\bar{s}$. The LEP limits on the masses of h^0 and A^0 are obtained by searching simultaneously for $e^+e^- \to Z \to Zh^0$ and $e^+e^- \to Z \to h^0A^0$. At LEP1, the intermediate Z is real and the final state Z is virtual, while at LEP2, the intermediate Z is virtual and the final state Z is real. The ZZh^0 and Zh^0A^0 couplings that govern these two decay rates are proportional to $\sin(\beta-\alpha)$ and $\cos(\beta-\alpha)$, respectively. Thus, one can use the LEP data to obtain simultaneous limits on m_{h^0} and m_{A^0}, since the two tree-level masses determine α and β. However, radiative corrections can be significant, so the final limits depend on the choice of MSSM parameters that govern the radiative corrections (of which, the third generation squark parameters are the most important). The present LEP 95% CL lower limits are $m_{A^0} > 90.2$ GeV and $m_{h^0} > 89.5$ GeV [40].

The Higgs mass limits quoted above were based on the assumption of $M_{\rm SUSY} = 1$ TeV and maximal stop mixing.[5] Under these assumptions, the LEP MSSM Higgs search excludes the region of $0.53 < \tan\beta < 2.25$ At future colliders, the MSSM Higgs search will extend the excluded region in the m_{A^0}—$\tan\beta$ plane. Eventually, at least one Higgs boson (h^0) must be discovered, or else low-energy supersymmetry in its minimal form must be discarded. It is possible that additional Higgs bosons (H^0, A^0 and/or $H^\pm$) will also be discovered. But, unlike h^0 whose mass is bounded from above [eq. (13)], the masses of H^0, A^0 and $H^\pm$ are not so restricted (although heavy Higgs boson masses must be below a few TeV in order to preserve the naturalness of the theory).

In the region of MSSM Higgs parameter space where $m_{A^0} \gg m_Z$, one can show that $m_{H^0} \sim m_{H^\pm} \sim m_{A^0}$ [where the corresponding mass differences are of $\mathcal{O}(m_Z^2/m_{A^0})$]. This parameter range is called *the decoupling limit* [41], since it

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[5]Although this tends to be a conservative assumption (that is, other choices ensure that *more* of the m_{A^0}—$\tan\beta$ plane is covered), there are a number of other parameter regimes in which certain Higgs search strategies become more problematical.

corresponds to the case where the Higgs sector of the effective low-energy theory is equivalent to that of the one-Higgs-doublet Standard Model. In particular, the properties of h^0 will be nearly indistinguishable from those of the Standard Model Higgs boson, while the heavier Higgs bosons may be too heavy for discovery at the next generation of future colliders.

Let us consider briefly the prospects for discovering and making precision measurements of the MSSM Higgs sector at the upgraded Tevatron, LHC and the next e^+e^- linear collider (LC). The upgraded Tevatron begins running in 2001 at $\sqrt{s} = 2$ TeV, with an initial goal of reaching 2 fb^{-1} of integrated luminosity per year. It has been suggested that an ambitious program could achieve a total integrated luminosity of 15 fb^{-1} by the end of 2007. LHC expect to begin taking data in 2006 at $\sqrt{s} = 14$ TeV, with an initial goal of 10 fb^{-1} per year. Eventually, the integrated luminosity is expected to reach 100 fb^{-1} per year. Finally, the next generation e^+e^- LC is now currently under development. Initially, one expects the LC to operate at $\sqrt{s} = 500$ GeV, with an integrated luminosity of 50 fb^{-1} per year, although there have been suggestions that the luminosity could be improved by nearly an order of magnitude.

The discovery reach of the Standard Model Higgs boson was analyzed by the Tevatron Higgs Working Group in ref. [42]. The relevant production mechanism is $q_i\bar{q}_j \to Vh_{\mathrm{SM}}$, where $V = W$ or Z. In all cases, it was assumed that h_{SM} would be observed via $h_{\mathrm{SM}} \to b\bar{b}$. The most relevant final state signatures corresponded to events in which the vector boson decayed leptonically ($W \to \ell\nu$, $Z \to \ell^+\ell^-$ and $Z \to \nu\bar{\nu}$, where $\ell = e$ or μ), resulting in $\ell\nu b\bar{b}$, $\nu\bar{\nu}b\bar{b}$ and $\ell^+\ell^- b\bar{b}$ final states. This analysis can be reinterpreted in terms of the search for the CP-even Higgs boson of the MSSM. In the MSSM at large $\tan\beta$, the enhancement of the $A^0 b\bar{b}$ coupling (and a similar enhancement of either the $h^0 b\bar{b}$ or $H^0 b\bar{b}$ coupling), provides a new search channel: $q\bar{q}$, $gg \to b\bar{b}\phi$ (where ϕ is a neutral Higgs boson with enhanced couplings to $b\bar{b}$). Combining both sets of analyses, the Tevatron Higgs Working Group obtained the anticipated 5σ Higgs discovery contours for the maximal mix-

ing scenario as a function of total integrated luminosity per detector (combining both CDF and DØ data sets) shown in fig. 2 [42].
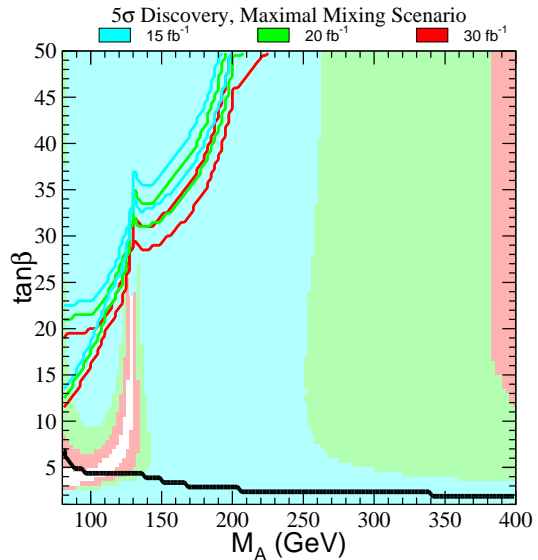

Figure 2. The anticipated Tevatron 5σ discovery region on the m_{A^0}—$\tan\beta$ plane, for the maximal mixing scenario and two different search channels: $q\bar{q} \to V\phi$ [$\phi = h^0$, H^0], $\phi \to b\bar{b}$ (shaded regions) and gg, $q\bar{q} \to b\bar{b}\phi$ [$\phi = h^0$, H^0, A^0], $\phi \to b\bar{b}$ (region in the upper left-hand corner bounded by the solid lines; the different lines correspond to CDF and DØ simulations). Different integrated luminosities are explicitly shown by the color coding (or shades of gray). The region below the solid black line near the bottom of the plot is excluded by the absence of observed $e^+e^- \to Z\phi$ events at LEP2. Taken from ref. [42].

Fig. 2 shows that a total integrated luminosity of about 20 fb^{-1} per experiment is necessary in order to assure a significant, although not exhaustive, coverage of the MSSM parameter space. If the anticipated 15 fb^{-1} integrated luminosity is achieved, the discovery reach will significantly extend beyond that of LEP. Nevertheless, the

MSSM Higgs boson could still evade capture at the Tevatron. We would then turn to the LHC to try to obtain a definitive Higgs boson discovery.

Over nearly[6] all the remaining region of the MSSM Higgs sector parameter space, at least one of the MSSM Higgs bosons will be detectable at the LHC, assuming that the machine runs at its design luminosity of 100 pb^{-1} per year, and under the assumption that the current detector design capabilities are achieved [44]. The LC would provide complete coverage of the MSSM Higgs sector parameter space (by extending the LEP-2 Higgs search) once the center-of-mass energy of the machine is above 300 GeV [45].

However, complete coverage of the MSSM parameter space only means that at least one Higgs boson of the MSSM can be detected. For example, in the decoupling limit (where $m_{A^0} \gtrsim 2m_Z$), h^0 will surely be discovered at either the Tevatron or the LHC (and can be easily observed at the LC). But, detection of the heavier non-minimal Higgs states H^0, A^0, and/or $H^\pm$ is not guaranteed. At the LHC, there is a region [44] in the m_{A^0}—$\tan\beta$ parameter space characterized (very roughly) by $3 \lesssim \tan\beta \lesssim 10$ and $m_{A^0} \gtrsim 300$ GeV such that only the h^0 will be detectable.[7] The LHC can make Higgs measurements (branching ratios and couplings) with some precision [46], and thus one can begin to check, if no other MSSM Higgs bosons are detected, whether the properties of h^0 deviate from those expected of the Standard Model Higgs boson ($h_{\rm SM}$).

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[6]At the present writing, not all possible MSSM parameter regimes have been studied with full detector simulations. Ref. [43] has pointed out that for certain special choices of MSSM Higgs sector parameters, it may be very difficult to achieve a 5σ discovery of *any* Higgs boson at the LHC. Moreover, it is important to note that in other regions of the Higgs sector parameter space, the LHC search strategies depend on the observation of small signals (above significant Standard Model backgrounds) in more than one channel. The present estimates of the statistical significance of the Higgs signal rely on theoretical determinations of both signal and background rates as well as simulations of detector performance.

[7]At large $\tan\beta$, the couplings of charged leptons to A^0 (and H^0 in the decoupling limit) are enhanced by a factor of $\tan\beta$. Thus, gluon-gluon fusion to A^0 or H^0 followed by Higgs decay to $\tau^+\tau^-$ and $\mu^+\mu^-$ should be observable at the LHC if $\tan\beta$ is large enough (the value of the latter depends on m_{A^0}).

At the LC, H^0A^0 and H^+H^- pair production are not kinematically allowed if $m_{A^0} \gtrsim \sqrt{s}/2$. Moreover, the production rate for h^0A^0 (although it may be kinematically allowed) is suppressed in the decoupling limit. Thus, the non-minimal Higgs states are not directly detectable at the LC for $m_{A^0} \gtrsim \sqrt{s}/2$. In the latter case, only the h^0 can be observed with properties that are nearly identical to that of $h_{\rm SM}$. More specifically, relative to the $h_{\rm SM}$ couplings to fermion pairs,

$$\frac{g_{h^0 f\bar{f}}}{g_{h_{\rm SM} f\bar{f}}} = 1 + \mathcal{O}\left(\frac{m_Z^2}{m_{A^0}^2}\right). \quad (14)$$

Thus, precision measurements of the h^0 branching ratios could reveal small discrepancies that would provide an indication of the value of m_{A^0} even if the heavy MSSM Higgs boson cannot be kinematically produced at the LC.

Battaglia and Desch [47] showed that for an LC with $\sqrt{s} = 350$ GeV and an ambitious integrated luminosity of 500 fb^{-1}, one could measure BR($h^0 \to b\bar{b}$) to within an accuracy of about $\pm 2.5\%$. To evaluate the significance of such a measurement, the percentage deviation of BR($h^0 \to b\bar{b}$) in the MSSM relative to that of $h_{\rm SM}$ was computed in ref. [48]. The result of the computation is a set of contours in the m_{A^0}—$\tan\beta$ plane, each one corresponding to a fixed percentage deviation of the BR. The results are sensitive to radiative corrections (which depend on the MSSM spectrum). One would expect the deviation to vanish at large values of m_{A^0} corresponding to the decoupling limit. However, at large $\tan\beta$, the approach to decoupling can be slowed due to $\tan\beta$-enhanced radiative corrections. For example, for $M_{\rm SUSY} = 1$ TeV and maximal stop mixing, ref. [48] finds that the contour corresponding to a 3% deviation in the BR starts at $m_{A^0} \simeq 600$ GeV for $\tan\beta = 3$ and slowly increases in m_{A^0} until it reaches $m_{A^0} \simeq 1$ TeV for $\tan\beta = 50$. This means that a precision measurement of BR($h^0 \to b\bar{b}$) can provide evidence for the non-minimal Higgs sector, with some sensitvity to large values of the heavy Higgs masses even if they lie above the center-of-mass energy of the LC.

5. THE LSP AND NLSP

The phenomenology of low-energy supersymmetry depends crucially on the properties of the lightest supersymmetric particle (LSP). In R-parity conserving low-energy supersymmetry, all Standard Model particles are R-even while their superpartners are R-odd. Thus, starting from an initial state involving ordinary (R-even) particles, it follows that supersymmetric particles must be produced in pairs. In general, these particles are highly unstable and decay quickly into lighter states. However, R-parity invariance also implies that the LSP is absolutely stable, and must eventually be produced at the end of a decay chain initiated by the decay of a heavy unstable supersymmetric particle.

In order to be consistent with cosmological constraints, a stable LSP is almost certainly electrically and color neutral [49]. Consequently, the LSP is weakly-interacting in ordinary matter, *i.e.* it behaves like a stable heavy neutrino and will escape detectors without being directly observed. Thus, the canonical signature for conventional R-parity-conserving supersymmetric theories is missing (transverse) energy, due to the escape of the LSP.

In mSUGRA models, the LSP is typically the lightest neutralino, $\widetilde{\chi}_1^0$, and it tends to be dominated by its $\widetilde{B}$ component (an LSP with significant higgsino components is possible only if $\tan\beta$ is large). However, there are some regions of mSUGRA parameter space where other possibilities for the LSP are realized. For example, there are regions of mSUGRA parameter space where the LSP is a chargino. These regions must be excluded since we reject the possibility of charged relic particles surviving the early universe [50]. In addition, one may impose cosmological constraints (such that the relic LSP's do not "overclose" the universe by contributing a mass density that is larger than the critical density of the universe [51]) to rule out additional regions of mSUGRA parameter space. The condition that the relic density of LSP's constitutes a significant part of the dark matter imposes even further restrictions on the mSUGRA parameter space [52].

In more general SUGRA models, the nature of the LSP need not be so constrained. One can envision a $\widetilde{\chi}_1^0$-LSP which has an arbitrary mixture of gaugino and higgsino components by relaxing gaugino mass unification. A nearly pure higgsino LSP is possible in the region where the (low-energy) gaugino Majorana mass parameters satisfy $M_1 \simeq M_2 \gtrsim \mu$ [53]. The sneutrino (in particular, the $\widetilde{\nu}_\tau$) can be a viable LSP, although it is unlikely to be a major component of the dark matter [54] unless the supersymmetric model incorporates some lepton number violation [55]. In AMSB models, eq. (11) yields $M_1 \simeq 2.8M_2$ and $M_3 \simeq -8.3M_2$, which implies that the lightest chargino pair and neutralino make up a nearly mass-degenerate triplet of winos; the corresponding LSP is approximately a pure $\widetilde{W}^3$.

In most GMSB models, the mass of the gravitino lies in the eV—keV regime. Thus, in this scenario the gravitino will be the LSP and the next-to-lightest supersymmetric particle (NLSP) also plays a crucial role in the phenomenology of supersymmetric particle production and decay. Note that unlike the LSP, the NLSP can be charged. In GMSB models, the most likely candidates for the NLSP are $\widetilde{\chi}_1^0$ and $\widetilde{\tau}_R^\pm$. The NLSP will decay into its Standard Model superpartner plus a gravitino [either $\widetilde{\chi}_1^0 \to N\widetilde{g}_{3/2}$ ($N = \gamma$, Z, or h^0) or $\widetilde{\tau}_R^\pm \to \tau^\pm \widetilde{g}_{3/2}$], with a lifetime that is quite sensitive to the model parameters. In a small range of parameter space, it is possible that several of the next-to-lightest supersymmetric particles are sufficiently degenerate in mass such that each one behaves as the NLSP.[8] In this case, these particles are called co-NLSP's [56]. Different choices for the identity of the NLSP and its decay rate lead to a variety of distinctive supersymmetric phenomenologies [57,58].

Since a light gravitino is stable in R-parity conserving GMSB models, it is also a candidate for dark matter [59]. Although very light gravitinos (eV masses) will not contribute significantly to the total mass density of the universe, the requirement that their relic density not overclose

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[8]For example, if $\widetilde{\tau}_R^\pm$ and $\widetilde{\chi}_1^0$ are nearly degenerate in mass, then neither $\widetilde{\tau}_R^\pm \to \tau^\pm \widetilde{\chi}_1^0$ nor $\widetilde{\chi}_1^0 \to \widetilde{\tau}_R^\pm \tau^\mp$ are kinematically allowed decays. In this case, $\widetilde{\tau}_R^\pm$ and $\widetilde{\chi}_1^0$ are co-NLSP's, and each decays dominantly into its Standard Model superpartner plus a gravitino.

the universe implies that $m_{3/2} \lesssim$ few keV in the usual early universe scenarios. Alternative scenarios do exist in which the gravitino of GMSB models can be somewhat heavier. In SUGRA-based models, the gravitino is typically not the LSP and thus is unstable, although its lifetime is quite long and is relevant only in considerations of early universe cosmology. The pertinent issues are nicely summarized in ref. [33].

More generally, in any R-parity-conserving supersymmetric model it is important to check that the relic density of LSP's does not overclose the universe. This can lead to useful constraints on the parameters of the model. Even if the LSP is a viable dark matter candidate, one should not necessarily constrain the parameters of the MSSM by requiring the LSP to be a major component of the dark matter. It may turn out that the main component of the dark matter has another source. Some examples are: the QCD axion, its supersymmetric partner (the axino [60]) or the lightest stable particle in the GMSB messenger sector [61].

6. CLASSES OF SUPERSYMMETRIC SIGNALS AT FUTURE COLLIDERS

The lack of knowledge of the origin and structure of the SUSY-breaking parameters implies that the predictions for low-energy supersymmetry and the consequent phenomenology depend on a plethora of unknown parameters. Many details of supersymmetric phenomenology are strongly dependent on the underlying assumptions of the model. Nevertheless, one can broadly classify supersymmetric signals at future colliders by considering the various theoretical approaches described in Section 3.

6.1. Missing energy signatures

In R-parity conserving low-energy supersymmetry, supersymmetric particles are produced in pairs. The subsequent decay of a heavy supersymmetric particle generally proceeds via a multistep decay chain [62–64], ending in the production of at least one supersymmetric particle that (in conventional models) is weakly interacting and escapes the collider detector. Thus, supersym-

metric particle production yields events that contain at least two escaping non-interacting particles, leading to a missing energy signature. At hadron colliders, it is only possible to detect missing transverse energy (E_T^{miss}), since the center-of-mass energy of the hard collision is not known on an event-by-event basis.

In conventional SUGRA-based models, the weakly-interacting LSP's that escape the collider detector (which yields large missing transverse energy) are accompanied by energetic jets and/or leptons. This is the "smoking-gun" signature of low-energy supersymmetry. In contrast, in AMSB models, the $\widetilde{\chi}_1^0$ is the LSP but the lightest neutralino and chargino are nearly degenerate in mass. If the mass difference is $\lesssim 100$ MeV, then $\widetilde{\chi}_1^+$ is long-lived and decays outside the detector [65,66]. In this case, some supersymmetric events would yield *no* missing energy and two semi-stable charged particles that pass through the detector.

In conventional GMSB models with a gravitino-LSP, all supersymmetric events contain at least two NLSP's, and the resulting signature depends on the NLSP properties. Four physically distinct possible scenarios emerge:

- The NLSP is electrically and color neutral and long-lived, and decays outside of the detector to its associated Standard Model partner and the gravitino.

- The NLSP is the sneutrino and decays invisibly into $\nu\widetilde{g}_{3/2}$ either inside or outside the detector.

In either of these two cases, the resulting missing-energy signal is then similar to that of the SUGRA-based models where $\widetilde{\chi}_1^0$ or $\widetilde{\nu}$ is the LSP.

- The NLSP is the $\widetilde{\chi}_1^0$ and decays inside the detector to $N\widetilde{g}_{3/2}$, where $N = \gamma$, Z or a neutral Higgs boson.

In this case, the gravitino-LSP behaves like the neutralino or sneutrino LSP of the SUGRA-based models. However, the missing energy events of the GMSB-based model are characterized by the associated production of (at least) two N's, one for each NLSP. Note that if $\widetilde{\chi}_1^0$ is lighter than

the Z and h^0 then $\mathrm{BR}(\widetilde{\chi}_1^0 \to \gamma \widetilde{g}_{3/2}) = 100\%$, and all supersymmetric production will result in missing energy events with at least two associated photons.

- The NLSP is a charged slepton (typically $\widetilde{\tau}_R$ in GMSB models if $m_{\widetilde{\tau}_R} < m_{\widetilde{\chi}_1^0}$), which decays to the corresponding lepton partner and gravitino.

If the decay is prompt, then one finds missing energy events with associated leptons. If the decay is not prompt, one observes a long-lived heavy semi-stable charged particle with *no* associated missing energy (prior to the decay of the NLSP).

There are also GMSB scenarios in which there are several nearly degenerate so-called co-NLSP's, any one of which can be produced at the penultimate step of the supersymmetric decay chain. The resulting supersymmetric signals would consist of events with two (or more) co-NLSP's, each one of which would decay according to one of the four scenarios delineated above. For additional details on the phenomenology of the co-NLSP's, see ref. [56].

In R-parity violating SUGRA-based models the LSP is unstable. If the RPV-couplings are sufficiently weak, then the LSP will decay outside the detector, and the standard missing energy signal applies. If the LSP decays inside the detector, the phenomenology of RPV models depends on the identity of the LSP and the branching ratio of possible final state decay products. If the latter includes a neutrino, then the corresponding RPV supersymmetric events would result in missing energy (through neutrino emission) in association with hadron jets and/or leptons. However, other decay chains are possible depending on the relative strengths of λ_L, λ'_L and λ_B [see eq. (4)]. Other possibilities include decays into charged leptons in association with jets (with no neutrinos), and decays into purely hadronic final states. Clearly, these latter events would contain little missing energy. If R-parity violation is present in GMSB models, the RPV decays of the NLSP can easily dominate over the NLSP decay to the gravitino. In this case, the phenomenology of the NLSP resembles that of the LSP of SUGRA-based RPV models [67].

6.2. Lepton (e, μ and τ) signatures

Once supersymmetric particles are produced at colliders, they do not necessarily decay to the LSP (or NLSP) in one step. The resulting decay chains can be complex, with a number of steps from the initial decay to the final state [63]. Along the way, decays can produce real or virtual W's, Z's, charginos, neutralinos and sleptons, which then can produce leptons in their subsequent decays. Thus, many models yield large numbers of supersymmetric events characterized by one or more leptons in association with missing energy, with or without hadronic jets.

One signature of particular note is events containing like-sign di-leptons [68]. The origin of such events is associated with the Majorana nature of the gaugino. For example, $\widetilde{g}\widetilde{g}$ production, followed by gluino decay via

$$\widetilde{g} \to q\bar{q}\widetilde{\chi}_1^\pm \to q\bar{q}\ell^\pm \nu \widetilde{\chi}_1^0 \tag{15}$$

can result in like-sign leptons since the $\widetilde{g}$ decay leads with equal probability to either ℓ^+ or ℓ^-. If the masses and mass differences are both substantial (which is typical in mSUGRA models, for example), like-sign di-lepton events will be characterized by fairly energetic jets and isolated leptons and by large E_T^{miss} from the LSP's. Other like-sign di-lepton signatures can arise in a similar way from the decay chains initiated by the heavier neutralinos.

Distinctive tri-lepton signals [69] can result from $\widetilde{\chi}_1^\pm \widetilde{\chi}_2^0 \to (\ell^\pm \nu \widetilde{\chi}_1^0)(\ell^+ \ell^- \widetilde{\chi}_1^0)$. Such events have little hadronic activity (apart from initial state radiation of jets off the annihilating quarks at hadron colliders). These events can have a variety of interesting characteristics depending on the fate of the final state neutralinos.

If the soft-SUSY-breaking slepton masses are flavor universal at the high energy scale M_X (as in mSUGRA models) and $\tan\beta \gg 1$, then the $\widetilde{\tau}_R$ will be significantly lighter than the other slepton states. As a result, supersymmetric decay chains involving (s)leptons will favor $\widetilde{\tau}_R$ production, leading to a predominance of events with multiple τ-leptons in the final state.

In GMSB models with a charged slepton NLSP, the decay $\widetilde{\ell} \to \ell \widetilde{g}_{3/2}$ (if prompt) yields at least two leptons for every supersymmetric event in as-

sociation with missing energy. In particular, in models with a $\widetilde{\tau}_R$ NLSP, supersymmetric events will characteristically contain at least two τ's.

In RPV models, decays of the LSP (in SUGRA models) or NLSP (in GMSB models) mediated by RPV-interactions proportional to λ_L and λ'_L will also yield supersymmetric events containing charged leptons. However, if the only significant RPV-interaction is the one proportional to λ'_L, then such events would contain little missing energy (in contrast to the GMSB signature described above).

6.3. b-quark signatures

The phenomenology of gluinos and squarks depends critically on their relative masses. If the gluino is heavier, it will decay dominantly into $q\widetilde{q}$,[9] while the squark can decay into quark plus chargino or neutralino. If the squark is heavier, it will decay dominantly into a quark plus gluino, while the gluino will decay into the three-body modes $q\bar{q}\widetilde{\chi}$ (where $\widetilde{\chi}$ can be either a neutralino or chargino, depending on the charge of the final state quarks). A number of special cases can arise when the possible mass splitting among squarks of different flavors is taken into account. For example, models of supersymmetric mass spectra have been considered where the third generation squarks are lighter than the squarks of the first two generations. If the gluino is lighter than the latter but heavier than the former, then the only open gluino two-body decay mode could be $b\widetilde{b}$. In such a case, all $\widetilde{g}\widetilde{g}$ events will result in at least four b-quarks in the final state (in associated with the usual missing energy signal, if appropriate). More generally, due to the flavor independence of the strong interactions, one expects three-body gluino decays into b-quarks in at least 20% of all gluino decays.[10] Additional b-quarks can arise from both top-quark and top-squark decays, and from neutral Higgs bosons produced somewhere in the chain decays [70]. Finally, at large $\tan\beta$,

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[9]In the following, the notation $q\widetilde{q}$ means either $q\bar{\widetilde{q}}$ or $\overline{q}\widetilde{q}$.
[10]Here we assume the approximate degeneracy of the first two generations of squarks, as suggested from the absence of FCNC decays. In many models, the b-squarks tend to be of similar mass or lighter than the squarks of the first two generations.

the enhanced Yukawa coupling to b-quarks can increase the rate of b-quark production in neutralino and chargino decays occurring at some step in the gluino chain decay. These observations suggest that many supersymmetric events at hadron colliders will be characterized by b-jets in association with missing energy [64,71].

6.4. Signatures involving photons

In mSUGRA models, most supersymmetric events do not contain isolated energetic photons. However, some areas of low-energy supersymmetric parameter space do exist in which final state photons can arise in the decay chains of supersymmetric particles. If one relaxes the condition of gaugino mass unification [eq. (9)], interesting alternative supersymmetric phenomenologies can arise. For example, if $M_1 \simeq M_2$, then the branching ratio for $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 \gamma$ can be significant [72]. In the model of ref. [58], the $\widetilde{\chi}_1^0$-LSP is dominantly higgsino, while $\widetilde{\chi}_2^0$ is dominantly gaugino. Thus, many supersymmetric decay chains end in the production of $\widetilde{\chi}_2^0$, which then decays to $\widetilde{\chi}_1^0 \gamma$. In this picture, the pair production of supersymmetric particles often yields two photons plus associated missing energy.

In GMSB models with a $\widetilde{\chi}_1^0$-NLSP, all supersymmetric decay chains would end up with the production of $\widetilde{\chi}_1^0$. Assuming that $\widetilde{\chi}_1^0$ decays inside the collider detector, one possible decay mode is $\widetilde{\chi}_1^0 \to \gamma\widetilde{g}_{3/2}$. In many models, the branching ratio for this radiative decay is significant (and could be as high as 100% if other possible two-body decay modes are not kinematically allowed). In the latter case, supersymmetric pair production would also yield events with two photons in association with large missing energy. The characteristics of these events differ in detail from those of the corresponding events expected in the model of ref. [58].

6.5. Kinks and long-lived heavy particles

In most SUGRA-based models, all supersymmetric particles in the decay chain decay promptly until the LSP is reached. The LSP is exactly stable and escapes the collider detector. However, exceptions are possible. In particular, if there is a supersymmetric particle that is

just barely heavier than the LSP, then its (three-body) decay rate to the LSP will be significantly suppressed and it could be long lived. For example, in AMSB models where $m_{\widetilde{\chi}_1^\pm} \simeq m_{\widetilde{\chi}_1^0}$, the $\widetilde{\chi}_1^\pm$ can be sufficiently long lived to yield a detectable vertex, or perhaps even exit the detector [65,66].

In GMSB models, the NLSP may be long-lived, depending on its mass and the scale of SUSY-breaking, $\sqrt{F}$. The couplings of the NLSP to the helicity $\pm\frac{1}{2}$ components of the gravitino are fixed by eq. (6). For $\sqrt{F} \sim 100$—10^4 TeV, this coupling is very weak, implying that all the supersymmetric particles other than the NLSP undergo chain decays down to the NLSP (the branching ratio for the direct decay to the gravitino is negligible). The NLSP is unstable and eventually decays to the gravitino. For example, in the case of the $\widetilde{\chi}_1^0$-NLSP (which is dominated by its $\widetilde{B}$ component), one can use eq. (6) to obtain $\Gamma(\widetilde{\chi}_1^0 \rightarrow \gamma\widetilde{g}_{3/2}) = m_{\widetilde{\chi}_1^0}^5 \cos^2\theta_W/16\pi F^2$. It then follows that

$$(c\tau)_{\widetilde{\chi}_1^0} \simeq 130 \left(\frac{100 \text{ GeV}}{m_{\widetilde{\chi}_1^0}}\right)^5 \left(\frac{\sqrt{F}}{100 \text{ TeV}}\right)^4 \mu\text{m} . (16)$$

For simplicity, assume that $\widetilde{\chi}_1^0 \rightarrow \gamma\widetilde{g}_{3/2}$ is the dominant NLSP decay mode. If $\sqrt{F} \sim 10^4$ TeV, then the decay length for the NLSP is $c\tau \sim 10$ km for $m_{\widetilde{\chi}_1^0} = 100$ GeV; while $\sqrt{F} \sim 100$ TeV implies a short but vertexable decay length. A similar result is obtained in the case of a charged NLSP. Thus, if $\sqrt{F}$ is sufficiently large, the charged NLSP will be semi-stable and may decay outside of the collider detector.

Finally, if R-parity violation is present, the decay rate of the LSP in SUGRA-based models (or the NLSP in R-parity-violating GMSB models) could be in the relevant range to yield visible secondary vertices.

6.6. Exotic supersymmetric signatures

If R-parity is not conserved, supersymmetric phenomenology exhibits many features that are quite distinct from those of the MSSM [67]. Both $\Delta L = 1$ and $\Delta L = 2$ phenomena are allowed (if L is violated), leading to neutrino masses and mixing [73], neutrinoless double beta decay [74], and sneutrino-antisneutrino mixing [75]. Further,

since the distinction between the Higgs and matter multiplets is lost, R-parity violation permits the mixing of sleptons and Higgs bosons, the mixing of neutrinos and neutralinos, and the mixing of charged leptons and charginos, leading to more complicated mass matrices and mass eigenstates than in the MSSM.

Some of the consequences for collider signatures have already been mentioned. Most important, the LSP in RPV SUGRA models is no longer stable, which implies that not all supersymmetric decay chains must yield missing-energy events at colliders. In particular, if $\widetilde{\chi}_1^0$ is the LSP, then its RPV decays contain visible particles:

$$\widetilde{\chi}_1^0 \rightarrow \underbrace{(jjj)}_{\lambda_B \neq 0}, \quad \underbrace{(\ell\ell^{(\prime)}\nu)}_{\lambda_L \neq 0}, \quad \underbrace{(\ell jj, \nu jj)}_{\lambda_L^\prime \neq 0}, \qquad (17)$$

where j stands for a hadronic jet (in this case arising from single quarks or anti-quarks), and the relevant RPV-coupling [in the notation of eq. (4)] is indicated below the corresponding channel. Thus, even $e^+e^- \rightarrow \widetilde{\chi}_1^0\widetilde{\chi}_1^0$ pair production and sneutrino decay via $\widetilde{\nu} \rightarrow \nu\widetilde{\chi}_1^0$ become visible (since the $\widetilde{\chi}_1^0$ now decays into visible channels). Likewise, a sneutrino LSP would also yield visible decay modes. For example, $\widetilde{\nu}$—h^0 mixing could lead to a substantial $\widetilde{\nu} \rightarrow b\bar{b}$ decay branching ratio. The relative strength of $\widetilde{\nu} \rightarrow q\bar{q}$, $\ell^+\ell^-$ also depends on the strength of the L-violating RPV couplings. In addition, $\widetilde{\chi}_1^0 \rightarrow \bar{\nu}\widetilde{\nu}$ (or its charge-conjugated state) would be visible once the sneutrino decayed.

A number of other phenomenological consequences are noteworthy. For $\lambda_L \neq 0$ sneutrino resonance production in e^+e^- [76] collisions becomes possible. For $\lambda_L^\prime \neq 0$, squarks can be regarded as leptoquarks [77] since the following processes are allowed: $e^+\overline{u}_m \rightarrow \bar{\widetilde{d}}_n \rightarrow e^+\overline{u}_m$, $\overline{\nu}\overline{d}_m$ and $e^+d_m \rightarrow \widetilde{u}_n \rightarrow e^+d_m$ (where m and n are generation labels). The same term responsible for the processes displayed above could also generate purely hadronic decays for sleptons and sneutrinos: $e.g.$, $\widetilde{\ell}_p^- \rightarrow \overline{u}_m d_n$ and $\widetilde{\nu}_p \rightarrow \overline{q}_m q_n$ ($q = u$ or d). If such decays were dominant, then the pair production of sleptons in e^+e^- events would lead to hadronic four-jet events with jet pairs of equal mass [78], a signature quite different from

the missing energy signals expected in the MSSM. Alternatively, $\lambda_L \neq 0$ could result in substantial branching fractions for $\widetilde{\ell} \to \ell\nu$ and $\widetilde{\nu} \to \ell^+\ell^-$ decays. Sneutrino pair production would then yield events containing four charged leptons with two lepton pairs of equal mass.

7. CONCLUSIONS

Low-energy supersymmetry remains the most elegant solution to the naturalness and hierarchy problems, while providing a possible link to Planck scale physics and the unification of particle physics and gravity. Nevertheless, the origin of the soft supersymmetry-breaking terms and the details of their structure remain a mystery. There are many theoretical ideas, but we still cannot be certain which region of the MSSM-124 parameter space (or some non-minimal extension thereof) is the one favored by nature. The key theoretical breakthroughs will surely require experimental guidance and input.

Thus, we must rely on experiments at future colliders to uncover evidence for low-energy supersymmetry. Canonical supersymmetric signatures at future colliders are well analyzed and understood. Much of the recent efforts have been directed at trying to develop strategies for precision measurements to establish the underlying supersymmetric structure of the interactions and to distinguish among models. However, we are far from understanding all possible facets of MSSM-124 parameter space (even restricted to those regions that are phenomenologically viable). For example, the phenomenological implications of the potentially new CP-violating phases that can arise in the MSSM and its consequences for collider physics have only recently begun to attract attention. Moreover, the variety of possible non-minimal models of low-energy supersymmetry presents additional challenges to experimenters who plan on searching for supersymmetry at future colliders.

If supersymmetry is discovered, it will provide a plethora of experimental signals and theoretical analyses. The variety of phenomenological manifestations and parameters of supersymmetry suggest that many years of experimental and theoretical work will be required before it will be possible to determine the precise nature of supersymmetry-breaking and the implications for a more fundamental theory of particle interactions.

ACKNOWLEDGMENTS

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