SUPERSYMMETRY

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This review is divided into two parts:

SUPERSYMMETRY, PART I (THEORY)

(by H.E. Haber)

I.1. Introduction: Supersymmetry (SUSY) is a generalization of the space-time symmetries of quantum field theory that transforms fermions into bosons and vice versa. It also provides a framework for the unification of particle physics and gravity [1–4], which is governed by the Planck scale, $M_{\rm P} \approx 10^{19}$ GeV (defined to be the energy scale where the gravitational interactions of elementary particles become comparable to the gauge interactions). In particular, it is possible that supersymmetry will ultimately explain the origin of the large hierarchy of energy scales from the W and Z masses to the Planck scale. The stability of this hierarchy in the presence of radiative corrections is not possible in the Standard Model, but can be maintained in supersymmetric theories.

If supersymmetry were an exact symmetry of nature, then particles and their superpartners (which differ in spin by half a unit) would be degenerate in mass. Since this is not observed in data, supersymmetry cannot be an exact symmetry and must be broken. Nevertheless, the stability of the hierarchy of scales mentioned above can still be maintained if the supersymmetry breaking is *soft* [5] and the corresponding supersymmetry-breaking mass terms are no larger than a few TeV. (In softly-broken supersymmetry, the theory behaves like an unbroken supersymmetric theory at energy scales much larger than the size of the supersymmetry-breaking masses.) The most interesting theories of this type are theories of "low-energy" (or "weak-scale") supersymmetry, where the effective scale of supersymmetry breaking is tied to the scale of electroweak symmetry breaking [6–8]. The latter is characterized by the Standard Model Higgs vacuum expectation value, $v = 246$ GeV.

At present, there are no unambiguous experimental results that require the existence of low-energy supersymmetry. However, one tantalizing clue may be the observed unification of the three gauge couplings at an energy scale close to the Planck scale. The unification of gauge couplings does not occur in the Standard Model, but is achievable in the minimal supersymmetric extension of the Standard Model, and provides an additional motivation for seriously considering the low-energy supersymmetric framework [9]. If experimentation at future colliders uncovers evidence for supersymmetry, this would have a profound effect on the study of TeV-scale physics, and the development of a more fundamental theory of mass and symmetry-breaking phenomena in particle physics.

I.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 [3,10]. In addition, the MSSM contains two hypercharge $Y = \pm 1$ Higgs doublets, 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) [11,12]. All renormalizable supersymmetric interactions consistent with (global) B–L conservation (B =baryon number and L =lepton number) are included. Finally, the most general soft-supersymmetry-breaking terms are added [5]. To generate nonzero neutrino masses, extra structure is needed as discussed briefly in section I.8.

If supersymmetry is associated with the origin of the scale of electroweak interactions, then the mass parameters introduced by the soft-supersymmetry-breaking must be generally of order 1 TeV or below [13] (although models have been proposed in which some supersymmetric particle masses can be larger, in the range of 1–10 TeV [14]). Some lower bounds on these parameters exist due to the absence of supersymmetric-particle production at current accelerators [15]. Additional constraints arise from limits on the contributions of virtual supersymmetric particle exchange to a variety of Standard Model processes [16,17].

For example, the Standard Model global fit to precision electroweak data is quite good [18]. If all supersymmetric particle masses are significantly heavier than m_Z (in practice, masses greater than 300 GeV are sufficient [19]), then the effects of the supersymmetric particles decouple in loop-corrections to electroweak observables [20]. In this case, the Standard Model global fit to precision data and the corresponding MSSM fit yield similar results. On the other hand, regions of parameter space with light supersymmetric particle masses (just above the present day experimental limits) can in some cases generate significant one-loop corrections, resulting in a slight improvement or worsening of the overall global fit to the electroweak data depending on the choice of the MSSM parameters [21]. Thus, the precision electroweak data provide some constraints on the magnitude of the soft-supersymmetry-breaking terms.

There are a number of other low-energy measurements that are especially sensitive to the effects of new physics through virtual loops. For example, the virtual exchange of supersymmetric particles can contribute to the muon anomalous magnetic moment, $a_\mu \equiv \frac{1}{2}(g-2)_\mu$, and to the inclusive decay rate for $b \to s\gamma$. The most recent theoretical analysis of $(g-2)_\mu$ finds only a small deviation (less than two standard deviations) of the theoretical prediction from the experimentally observed value [22]. The theoretical prediction for $\Gamma(b \to s\gamma)$ agrees quite well (within the error bars) to the experimental observation [23]. In both cases, supersymmetric corrections could have generated an observable shift from the Standard Model prediction in some regions of the MSSM parameter space [23–25]. The absence of a significant deviation places interesting constraints on the low-energy supersymmetry parameters.

As a consequence of B–L invariance, the MSSM possesses a multiplicative R-parity invariance, where $R = (-1)^{3(B-L)+2S}$ for a particle of spin S [26]. Note that this implies that all the ordinary Standard Model particles have even R parity, whereas the corresponding supersymmetric partners have odd R parity. The conservation of R parity in scattering and decay processes has a crucial impact on supersymmetric phenomenology. For example, 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 into lighter states. However, R-parity invariance also implies that the lightest supersymmetric particle (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 [27]. (There are some model circumstances in which a colored gluino LSP is allowed [28], but we do not consider this possibility further here.) Consequently, the LSP in an R-parity-conserving theory is weakly interacting with ordinary matter, *i.e.*, it behaves like a stable heavy neutrino and will escape collider 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. Moreover, the LSP is a prime candidate for "cold dark matter" [29], a potentially important component of the non-baryonic dark matter that is required in many models of cosmology and galaxy formation [30]. Further aspects of dark matter can be found in Ref. [31].

In the MSSM, supersymmetry breaking is accomplished by including the most general renormalizable soft-supersymmetry-breaking terms consistent with the SU(3)×SU(2)×U(1) gauge symmetry and R-parity invariance. These terms parameterize our ignorance of the fundamental mechanism of supersymmetry breaking. If supersymmetry breaking occurs spontaneously, then a massless Goldstone fermion called the *goldstino* ($\widetilde{G}$) must exist. The goldstino would then be the LSP and could play an important role in supersymmetric phenomenology [32]. However, the goldstino is a physical degree of freedom only in models of spontaneously broken global supersymmetry. If the supersymmetry is a local symmetry, then the theory must incorporate gravity; the resulting theory is called supergravity. In models of spontaneously broken supergravity, the goldstino is "absorbed" by the *gravitino* ($\widetilde{g}_{3/2}$), the spin-3/2 partner of the graviton [33]. By this super-Higgs mechanism, the goldstino is removed from the physical spectrum and the gravitino acquires a mass ($m_{3/2}$).

It is very difficult (perhaps impossible) to construct a realistic model of spontaneously-broken low-energy supersymmetry where the supersymmetry breaking arises solely as a consequence of the interactions of the particles of the MSSM. A more viable scheme posits a theory consisting of at least two distinct sectors: a "hidden" sector consisting of particles that are completely neutral with respect to the Standard Model gauge group, and a "visible" sector consisting of the particles of the MSSM. There are no renormalizable tree-level interactions between particles of the visible and hidden sectors. Supersymmetry breaking is assumed to occur in the hidden sector, and to then be transmitted to the MSSM by some mechanism. Two theoretical scenarios have been examined in detail: gravity-mediated and gauge-mediated supersymmetry breaking.

Supergravity models provide a natural mechanism for transmitting the supersymmetry breaking of the hidden sector to the particle spectrum of the MSSM. In models of *gravity-mediated* supersymmetry breaking, gravity is the messenger of supersymmetry breaking [34,35]. More precisely, supersymmetry breaking is mediated by effects of gravitational strength (suppressed by an inverse power of the Planck mass). In this scenario, the gravitino mass is of order the electroweak-symmetry-breaking scale, while its couplings are roughly gravitational in strength [1,36]. Such a gravitino would play no role in supersymmetric phenomenology at colliders.

In *gauge-mediated* supersymmetry breaking, supersymmetry 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 [37,38]. The direct coupling of the messengers to the hidden sector generates a supersymmetry breaking spectrum in the messenger sector. Finally, supersymmetry 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 supersymmetry breaking. However, in models of gauge-mediated supersymmetry 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.

During the last few years, new approaches to supersymmetry breaking have been proposed, based on theories in which the number of space dimensions is greater than three. This is not a new idea-consistent superstring theories are formulated in ten spacetime dimensions, and the associated M-theory is based in eleven spacetime dimensions [39]. Nevertheless, in all approaches considered above, the string scale and the inverse size of the extra dimensions are assumed to be at or near the Planck scale, below which an effective four spacetime dimensional broken supersymmetric field theory emerges. More recently, a number of supersymmetry-breaking mechanisms have been proposed that are inherently extra-dimensional. In some cases, the size of the extra dimensions can be significantly larger than $M_{\rm P}^{-1}$; in some cases of order $({\rm TeV})^{-1}$ or even larger [40,41]. For example, in one approach, the fields of the MSSM live on some brane (a lower-dimensional manifold existing in a higher dimensional spacetime), while the sector of the theory that breaks supersymmetry lives on a second separated brane. Two examples of this approach are anomaly-mediated supersymmetry breaking of Ref. [42] and gaugino-mediated supersymmetry breaking of Ref. [43]; in both cases supersymmetry-breaking is transmitted through fields that live in the bulk (the higher dimensional space between the two branes). This setup has some features in common with both gravity-mediated and gauge-mediated supersymmetry breaking (*e.g.*, a hidden and visible sector and messengers). In another approach, one starts with a higher dimensional theory, which is compactified to four spacetime dimensions. In this approach, supersymmetry is broken by boundary conditions on the compactified space that distinguish between fermions and bosons [44](the so-called Scherk-Schwarz mechanism [45]). The phenomenology of such models can be strikingly different from the usual MSSM [46]. These approaches clearly deserve further investigation, although they will not be discussed further here.

I.3. Parameters of the MSSM: The parameters of the MSSM are conveniently described by considering separately the supersymmetry-conserving sector and the supersymmetry-breaking sector. A careful discussion of the conventions used in defining the MSSM parameters can be found in Ref. [47]. For simplicity, consider the case of one generation of quarks, leptons, and their scalar superpartners. The parameters of the supersymmetry-conserving sector consist of: (i) gauge couplings: g_s, g, and g', corresponding to the Standard Model gauge group SU(3)×SU(2)×U(1) respectively; (ii) a supersymmetry-conserving Higgs mass parameter μ; and (iii) Higgs-fermion Yukawa coupling constants: λ_u, λ_d, and λ_e (corresponding to the coupling of one generation of quarks, leptons, and their superpartners to the Higgs bosons and higgsinos).

The supersymmetry-breaking sector contains the following set of parameters: (i) gaugino Majorana masses M_3, M_2, and M_1 associated with the SU(3), SU(2), and U(1) subgroups of the Standard Model; (ii) five scalar squared-mass parameters for the squarks and sleptons, $M^2_{\widetilde{Q}}$, $M^2_{\widetilde{U}}$, $M^2_{\widetilde{D}}$, $M^2_{\widetilde{L}}$, and $M^2_{\widetilde{E}}$ [corresponding to the five electroweak gauge multiplets, *i.e.*, superpartners of $(u, d)_L$, u^c_L, d^c_L, $(\nu,\ e^-)_L$, and e^c_L, where the superscript c indicates a charge-conjugated fermion]; (iii) Higgs-squark-squark and Higgs-slepton-slepton trilinear interaction terms, with coefficients A_u, A_d, and A_e (these are the so-called "A parameters"); and (iv) three scalar Higgs squared-mass parameters-two of which (m^2_1 and m^2_2) contribute to the diagonal Higgs squared-masses, given by $m^2_1 + |\mu|^2$ and $m^2_2 + |\mu|^2$, and a third which contributes to the off-diagonal Higgs squared-mass term, $m^2_{12} \equiv B\mu$ (which defines the "B-parameter"). These three squared-mass parameters can be re-expressed in terms of the two Higgs vacuum expectation values, v_d and v_u (also called v_1 and v_2, respectively, in the literature), and one physical Higgs mass. Here, v_d (v_u) is the vacuum expectation value of the Higgs field which couples exclusively to down-type (up-type) quarks and leptons. Note that $v^2_d + v^2_u = 4m^2_W/g^2 = (246\ \text{GeV})^2$ is fixed by the W mass and the gauge coupling, whereas the ratio

$$\tan\beta = v_u/v_d \tag{1}$$

is a free parameter of the model.

The total number of degrees of freedom of the MSSM is quite large, primarily due to the parameters of the soft-supersymmetry-breaking sector. In particular, in the case of three generations of quarks, leptons, and their superpartners, $M^2_{\widetilde{Q}}$, $M^2_{\widetilde{U}}$, $M^2_{\widetilde{D}}$, $M^2_{\widetilde{L}}$, and $M^2_{\widetilde{E}}$ are hermitian 3×3 matrices, and the A parameters are complex 3×3 matrices. In addition, M_1, M_2, M_3, B, and μ are in general complex. Finally, as in the Standard Model, the Higgs-fermion Yukawa couplings, λ_f

($f = u$, d, and e), are complex 3×3 matrices that are related to the quark and lepton mass matrices via: $M_f = \lambda_f v_f/\sqrt{2}$, where $v_e \equiv v_d$ (with v_u and v_d as defined above). However, not all these parameters are physical. Some of the MSSM parameters can be eliminated by expressing interaction eigenstates in terms of the mass eigenstates, with an appropriate redefinition of the MSSM fields to remove unphysical degrees of freedom. The analysis of Ref. [48] shows that the MSSM possesses 124 independent parameters. Of these, 18 parameters correspond to Standard Model parameters (including the QCD vacuum angle $\theta_{\rm QCD}$), one corresponds to a Higgs sector parameter (the analogue of the Standard Model Higgs mass), and 105 are genuinely new parameters of the model. The latter include: five real parameters and three CP-violating phases in the gaugino/higgsino sector, 21 squark and slepton masses, 36 real mixing angles to define the squark and slepton mass eigenstates, and 40 CP-violating phases that can appear in squark and slepton interactions. The most general R-parity-conserving minimal supersymmetric extension of the Standard Model (without additional theoretical assumptions) will be denoted henceforth as MSSM-124 [49].

I.4. The supersymmetric-particle sector: Consider the sector of supersymmetric particles (*sparticles*) in the MSSM. The supersymmetric partners of the gauge and Higgs bosons are fermions, whose names are obtained by appending "ino" at the end of the corresponding Standard Model particle name. The *gluino* is the color octet Majorana fermion partner of the gluon with mass $M_{\widetilde{g}} = |M_3|$. The supersymmetric partners of the electroweak gauge and Higgs bosons (the *gauginos* and *higgsinos*) can mix. As a result, the physical mass eigenstates are model-dependent linear combinations of these states, called *charginos* and *neutralinos*, which are obtained by diagonalizing the corresponding mass matrices. The chargino-mass matrix depends on M_2, μ, $\tan\beta$, and m_W [50].

The corresponding chargino-mass eigenstates are denoted by $\widetilde{\chi}_1^+$ and $\widetilde{\chi}_2^+$, with masses

$$M^2_{\widetilde{\chi}_1^+,\widetilde{\chi}_2^+} = \tfrac{1}{2}\Big\{|\mu|^2 + |M_2|^2 + 2m_W^2 \mp \Big[\left(|\mu|^2 + |M_2|^2 + 2m_W^2\right)^2$$

$$- 4|\mu|^2|M_2|^2 - 4m_W^4 \sin^2 2\beta + 8m_W^2 \sin 2\beta \,\mathrm{Re}(\mu M_2)\Big]^{1/2}\Bigg\}, \quad (2)$$

where the states are ordered such that $M_{\widetilde{\chi}_1^+} \leq M_{\widetilde{\chi}_2^+}$. If CP-violating effects are neglected (in which case, M_2 and μ are real parameters), then one can choose a convention where $\tan\beta$ and M_2 are positive. (Note that the relative sign of M_2 and μ is meaningful. The sign of μ is convention-dependent; the reader is warned that both sign conventions appear in the literature.) The sign convention for μ implicit in Eq. (2) is used by the LEP collaborations [15] in their plots of exclusion contours in the M_2 vs. μ plane derived from the non-observation of $e^+e^- \to \widetilde{\chi}_1^+\widetilde{\chi}_1^-$.

The neutralino mass matrix depends on M_1, M_2, μ, $\tan\beta$, m_Z, and the weak mixing angle θ_W [50]. The corresponding neutralino eigenstates are usually denoted by $\widetilde{\chi}_i^0$ ($i = 1, \ldots 4$), according to the convention that $M_{\widetilde{\chi}_1^0} \leq M_{\widetilde{\chi}_2^0} \leq M_{\widetilde{\chi}_3^0} \leq M_{\widetilde{\chi}_4^0}$. If a chargino or neutralino eigenstate approximates a particular gaugino or higgsino state, it is convenient to employ the corresponding nomenclature. Specifically, if M_1 and M_2 are small compared to m_Z and $|\mu|$, then the lightest neutralino $\widetilde{\chi}_1^0$ would be nearly a pure *photino*, $\widetilde{\gamma}$, the supersymmetric partner of the photon. If M_1 and m_Z are small compared to M_2 and $|\mu|$, then the lightest neutralino would be nearly a pure *bino*, $\widetilde{B}$, the supersymmetric partner of the weak hypercharge gauge boson. If M_2 and m_Z are small compared to M_1 and $|\mu|$, then the lightest chargino pair and neutralino would constitute a triplet of roughly mass-degenerate pure *winos*, $\widetilde{W}^{\pm}$, and $\widetilde{W}_3^0$, the supersymmetric partners of the weak SU(2) gauge bosons. Finally, if $|\mu|$ and m_Z are small compared to M_1 and M_2, then the lightest neutralino would be nearly a pure *higgsino*. Each of the above cases leads to a strikingly different phenomenology.

The supersymmetric partners of the quarks and leptons are spin-zero bosons: the *squarks*, charged *sleptons*, and *sneutrinos*. For simplicity, only the one-generation case is illustrated below (using first-generation notation). For a given fermion f, there are two supersymmetric partners, $\widetilde{f}_L$ and $\widetilde{f}_R$, which are scalar partners of the corresponding left- and right-handed fermion. (There is no $\widetilde{\nu}_R$ in the MSSM.) However, in general, $\widetilde{f}_L$ and $\widetilde{f}_R$ are not mass-eigenstates, since there is $\widetilde{f}_L$-$\widetilde{f}_R$ mixing which is proportional in strength to the corresponding element of the scalar squared-mass matrix [51]

$$M_{LR}^2 = \begin{cases} m_d(A_d - \mu\tan\beta), & \text{for "down"-type } f \\ m_u(A_u - \mu\cot\beta), & \text{for "up"-type } f, \end{cases} \quad (3)$$

where m_d (m_u) is the mass of the appropriate "down" ("up") type quark or lepton. The signs of the A parameters are also convention-dependent; see Ref. [47]. Due to the appearance of the *fermion* mass in Eq. (3), one expects M_{LR} to be small compared to the diagonal squark and slepton masses, with the possible exception of the top-squark, since m_t is large, and the bottom-squark and tau-slepton if $\tan\beta \gg 1$.

The (diagonal) L- and R-type squark and slepton squared-masses are given by

$$M_{\widetilde{f}_L}^2 = M_{\widetilde{F}}^2 + m_f^2 + (T_{3f} - e_f\sin^2\theta_W)m_Z^2\cos 2\beta\,,$$

$$M_{\widetilde{f}_R}^2 = M_{\widetilde{R}}^2 + m_f^2 + e_f\sin^2\theta_W m_Z^2\cos 2\beta\,, \quad (4)$$

where $M_{\widetilde{F}}^2 \equiv M_{\widetilde{Q}}^2$ [$M_{\widetilde{L}}^2$] for $\widetilde{u}_L$ and $\widetilde{d}_L$ [$\widetilde{\nu}_L$ and $\widetilde{e}_L$], and $M_{\widetilde{R}}^2 \equiv M_{\widetilde{U}}^2$, $M_{\widetilde{D}}^2$, and $M_{\widetilde{E}}^2$ for $\widetilde{u}_R$, $\widetilde{d}_R$, and $\widetilde{e}_R$, respectively. In addition, $e_f = \frac{2}{3}$, $-\frac{1}{3}$, 0, -1 for $f = u$, d, ν, and e, respectively, $T_{3f} = \frac{1}{2}$ [$-\frac{1}{2}$] for up-type [down-type] squarks and sleptons, and m_f is the corresponding quark or lepton mass. Squark and slepton mass eigenstates, generically called $\widetilde{f}_1$ and $\widetilde{f}_2$ (these are linear combinations of $\widetilde{f}_L$ and $\widetilde{f}_R$), are obtained by diagonalizing the corresponding 2×2 squared-mass matrices.

In the case of three generations, the general analysis is more complicated. The scalar squared-masses [$M_{\widetilde{F}}^2$ and $M_{\widetilde{R}}^2$ in Eq. (4)], the fermion masses m_f, and the A parameters are now 3×3 matrices as noted in Section I.3. Thus, to obtain the squark and slepton mass eigenstates, one must diagonalize 6×6 mass matrices. As a result, intergenerational mixing is possible, although there are some constraints from the nonobservation of FCNC's [16,17]. In practice, because off-diagonal scalar mixing is appreciable only for the third generation, this additional complication can usually be neglected.

It should be noted that all mass formulae quoted in this section are tree-level results. One-loop corrections will modify all these results, and eventually must be included in any precision study of supersymmetric phenomenology [52].

I.5. The Higgs sector of the MSSM: Next, consider the MSSM Higgs sector [11,12,53]. Despite the large number of potential CP-violating phases among the MSSM-124 parameters, the tree-level MSSM Higgs sector is automatically CP-conserving. That is, unphysical phases can be absorbed into the definition of the Higgs fields such that $\tan\beta$ is a real parameter (conventionally chosen to be positive). Moreover, the physical neutral Higgs scalars are CP eigenstates. The model contains five physical Higgs particles: a charged Higgs boson pair ($H^{\pm}$), two CP-even neutral Higgs bosons (denoted by h^0 and H^0 where $m_h \leq m_H$), and one CP-odd neutral Higgs boson (A^0).

The properties of the Higgs sector are determined by the Higgs potential, which is made up of quadratic terms [whose squared-mass coefficients were mentioned above Eq. (1)] and quartic interaction terms whose coefficients are dimensionless couplings. The quartic interaction terms are manifestly supersymmetric at tree-level (and are modified by supersymmetry-breaking effects only at the loop level). In general, the quartic couplings arise from two sources: (i) the supersymmetric generalization of the scalar potential (the so-called "F-terms"), and (ii) interaction terms related by supersymmetry to the coupling of the scalar fields and the gauge fields, whose coefficients are proportional to the corresponding gauge couplings (the so-called "D-terms"). In the MSSM, F-term contributions to the quartic couplings are absent (although such terms may be present in extensions of the MSSM, *e.g.*, models with Higgs singlets). As a result, the strengths of the MSSM quartic Higgs interactions are fixed in terms of the gauge couplings. Due to the resulting constraint on the form of the two-Higgs-doublet scalar potential, all the tree-level MSSM Higgs-sector parameters depend only on two quantities: $\tan\beta$ [defined in Eq. (1)] and one Higgs mass (usually taken to be m_A). From these two quantities, one can predict the values of the remaining Higgs boson masses, an angle α (which measures the component of the original $Y = \pm 1$ Higgs doublet states in the physical CP-even neutral scalars), and the Higgs boson self-couplings.

When one-loop radiative corrections are incorporated, additional parameters of the supersymmetric model enter via virtual loops. The impact of these corrections can be significant [54]. For example, at tree-level, MSSM-124 predicts $m_h \leq m_Z|\cos 2\beta| \leq m_Z$ [11,12]. If this prediction were unmodified, it would be in conflict with the MSSM Higgs mass bounds obtained at LEP [55]. However, when radiative corrections are included, the light Higgs-mass upper bound may be significantly increased. The qualitative behavior of the radiative corrections can be most easily seen in the large top-squark mass limit, where in addition, both the splitting of the two diagonal entries [Eq. (4)] and the two off-diagonal entries [Eq. (3)] of the top-squark squared-mass matrix are small in comparison to the average of the two top-squark squared-masses, $M_{\rm S}^2 \equiv \frac{1}{2}(M_{\widetilde t_1}^2 + M_{\widetilde t_2}^2)$. In this case (assuming $m_A > m_Z$), the upper bound on the lightest CP-even Higgs mass at one-loop is approximately given by

$$m_h^2 \lesssim m_Z^2 + \frac{3g^2 m_t^4}{8\pi^2 m_W^2}\left\{\ln\left(M_{\rm S}^2/m_t^2\right) + \frac{X_t^2}{M_{\rm S}^2}\left(1 - \frac{X_t^2}{12M_{\rm S}^2}\right)\right\}, \quad (5)$$

where $X_t \equiv A_t - \mu\cot\beta$ is the top-squark mixing factor [see Eq. (3)]. A more complete treatment of the radiative corrections [56] shows that Eq. (5) somewhat overestimates the true upper bound of m_h. These more refined computations, which incorporate renormalization group improvement and the leading two-loop contributions, yield $m_h \lesssim 130$ GeV (with an accuracy of a few GeV) for $m_t = 175$ GeV and $M_S \lesssim 2$ TeV [56].

In addition, one-loop radiative corrections can introduce CP-violating effects in the Higgs sector, which depend on some of the CP-violating phases among the MSSM-124 parameters [57]. Although these effects are more model-dependent, they can have a non-trivial impact on the Higgs searches at future colliders.

I.6. Reducing the MSSM parameter freedom: In Sections I.4 and I.5 we surveyed the parameters that comprise the MSSM-124. However in its most general form, the MSSM-124 is not a phenomenologically-viable theory over most of its parameter space. This conclusion follows from the observation that a generic point in the MSSM-124 parameter space exhibits: (i) no conservation of the separate lepton numbers L_e, L_μ, and L_τ; (ii) unsuppressed FCNC's; and (iii) new sources of CP violation that are inconsistent with the experimental bounds. For example, the MSSM contains many new sources of CP violation [58]. In particular, some combination of the complex phases of the gaugino-mass parameters, the A parameters, and μ must be less than of order 10^{-2}–10^{-3} (for a supersymmetry-breaking scale of 100 GeV) to avoid generating electric dipole moments for the neutron, electron, and atoms in conflict with observed data [59,60]. As a result of the phenomenological deficiencies listed above, 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.

The MSSM-124 is also theoretically incomplete since it provides no explanation for the origin of the supersymmetry-breaking parameters (and in particular, why these parameters should conform to the exceptional points of the parameter space mentioned above). Moreover, there is no understanding of the choice of parameters that leads to the breaking of the electroweak symmetry. What is needed ultimately is a fundamental theory of supersymmetry breaking, which would provide a rationale for some set of soft-supersymmetry breaking terms that would be consistent with the phenomenological constraints referred to above. Presumably, the number of independent parameters characterizing such a theory would be considerably less than 124.

In the absence of a fundamental theory of supersymmetry breaking, there are two general approaches for reducing the parameter freedom of MSSM-124. In the low-energy 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 $M^2_{\widetilde{Q}}$, $M^2_{\widetilde{U}}$, $M^2_{\widetilde{D}}$, $M^2_{\widetilde{L}}$, $M^2_{\widetilde{E}}$, and the matrix A parameters are generation-independent (horizontal universality [7,48,61]). 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 [62]). In either case, L_e, L_μ, and L_τ are separately conserved, while tree-level FCNC's are automatically absent. In both cases, the number of free parameters characterizing the MSSM is substantially less than 124. Both scenarios are phenomenologically viable, although there is no strong theoretical basis for either scenario.

In the high-energy approach, one treats the parameters of the MSSM as running parameters and imposes a particular structure on the soft-supersymmetry-breaking terms at a common high-energy scale (such as the Planck scale, $M_{\rm P}$). Using the renormalization group equations, one can then derive the low-energy MSSM parameters. The initial conditions (at the appropriate high-energy scale) for the renormalization group equations depend on the mechanism by which supersymmetry breaking is communicated to the effective low energy theory. Examples of this scenario are provided by models of gravity-mediated and gauge-mediated supersymmetry breaking (see Section I.2). One bonus of such an approach is that one of the diagonal Higgs squared-mass parameters is typically driven negative by renormalization group evolution. Thus, electroweak symmetry breaking is generated radiatively, and the resulting electroweak symmetry-breaking scale is intimately tied to the scale of low-energy supersymmetry breaking.

One prediction of the high-energy approach that arises in most grand unified supergravity models and gauge-mediated supersymmetry-breaking models is the unification of the (tree-level) gaugino mass parameters at some high-energy scale $M_{\rm X}$, *i.e.*,

$$M_1(M_X) = M_2(M_X) = M_3(M_X) = m_{1/2}\,. \qquad (6)$$

Consequently, the effective low-energy gaugino mass parameters (at the electroweak scale) are related:

$$M_3 = (g_s^2/g^2)M_2\ , \qquad M_1 = (5g'^2/3g^2)M_2 \simeq 0.5M_2\ . \qquad (7)$$

In this case, the chargino and neutralino masses and mixing angles depend only on three unknown parameters: the gluino mass, μ, and $\tan\beta$. If in addition $|\mu| \gg M_1$, m_Z, then the lightest neutralino is nearly a pure bino, an assumption often made in supersymmetric particle searches at colliders.

In a certain class of supergravity models, tree-level masses for the gauginos are absent. The gaugino mass parameters arise at one-loop and do not satisfy Eq. (7). In this case, 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 [42]. This approach is called *anomaly-mediated* supersymmetry breaking. Eq. (7) is then replaced (in the one-loop approximation) by:

$$M_i \simeq \frac{b_i g_i^2}{16\pi^2} m_{3/2}\,, \tag{8}$$

where $m_{3/2}$ is the gravitino mass (assumed to be of order 1 TeV), and b_i are the coefficients of the MSSM gauge beta-functions corresponding to the corresponding U(1), SU(2) and SU(3) gauge groups: $(b_1, b_2, b_3) = (\frac{33}{5}, 1, -3)$. Eq. (8) yields $M_1 \simeq 2.8 M_2$ and $M_3 \simeq -8.3 M_2$, which implies that over most of the MSSM parameter space the lightest chargino pair and neutralino make up a nearly mass-degenerate triplet of winos. (For example, if $|\mu| \gg m_Z$, then Eq. (8) implies that $M_{\widetilde{\chi}_1^\pm} \simeq M_{\widetilde{\chi}_1^0} \simeq M_2$ [63].) The corresponding supersymmetric phenomenology differs significantly from the standard phenomenology based on Eq. (7), and is explored in detail in Ref. [64]. Anomaly-mediated supersymmetry breaking also generates (approximate) flavor-diagonal squark and slepton mass matrices. However, this yields negative squared-mass contributions for the sleptons in the MSSM. This fatal flaw may be possible to cure in approaches beyond the minimal supersymmetric model [65]. Alternatively, one may conclude that anomaly-mediation is not the sole source of supersymmetry-breaking in the slepton sector.

I.7. The constrained MSSMs: mSUGRA, GMSB, and SGUTs: One way to guarantee the absence of significant FCNC's mediated by virtual supersymmetric-particle exchange is to posit that the diagonal soft-supersymmetry-breaking scalar squared-masses are universal at some energy scale. In models of gauge-mediated supersymmetry breaking, scalar squared-masses are expected to be flavor-independent since gauge forces are flavor-blind. In the *minimal* supergravity (mSUGRA) framework [1–3], the soft-supersymmetry-breaking parameters at the Planck scale take a particularly simple form in which the scalar squared-masses and the A parameters are flavor-diagonal and universal [34]:

$$
\begin{aligned}
& M^2_{\widetilde{Q}}(M_{\rm P}) = M^2_{\widetilde{U}}(M_{\rm P}) = M^2_{\widetilde{D}}(M_{\rm P}) = m_0^2 \mathbf{1}\,, \\
& M^2_{\widetilde{L}}(M_{\rm P}) = M^2_{\widetilde{E}}(M_{\rm P}) = m_0^2 \mathbf{1}\,, \\
& m_1^2(M_{\rm P}) = m_2^2(M_{\rm P}) = m_0^2\,, \\
& A_U(M_{\rm P}) = A_D(M_{\rm P}) = A_L(M_{\rm P}) = A_0 \mathbf{1}\,,
\end{aligned} \tag{9}
$$

where $\mathbf{1}$ is a 3×3 identity matrix in generation space. Renormalization group evolution is then used to derive the values of the supersymmetric parameters at the low-energy (electroweak) scale. For example, to compute squark and slepton masses, one must use the *low-energy* values for $M^2_{\widetilde{F}}$ and $M^2_{\widetilde{R}}$ in Eq. (4). Through the renormalization group running with boundary conditions specified in Eq. (7) and Eq. (9), one can show that the low-energy values of $M^2_{\widetilde{F}}$ and $M^2_{\widetilde{R}}$ depend primarily on m_0^2 and $m^2_{1/2}$. A number of useful approximate analytic expressions for superpartner masses in terms of the mSUGRA parameters can be found in Ref. [66].

Clearly, in the mSUGRA approach, the MSSM-124 parameter freedom has been significantly reduced. For example, typical mSUGRA models give low-energy values for the scalar mass parameters that satisfy $M_{\widetilde{L}} \approx M_{\widetilde{E}} < M_{\widetilde{Q}} \approx M_{\widetilde{U}} \approx M_{\widetilde{D}}$, with the squark mass parameters somewhere between a factor of 1–3 larger than the slepton mass parameters (*e.g.*, see Ref. [66]). More precisely, the low-energy values of the squark mass parameters of the first two generations are roughly degenerate, while $M_{\widetilde{Q}_3}$ and $M_{\widetilde{U}_3}$ are typically reduced by a factor of 1–3 from the values of the first and second generation squark mass parameters, because of renormalization effects due to the heavy top quark mass.

As a result, one typically finds that four flavors of squarks (with two squark eigenstates per flavor) and $\widetilde{b}_R$ are nearly mass-degenerate. The $\widetilde{b}_L$ mass and the diagonal $\widetilde{t}_L$ and $\widetilde{t}_R$ masses are reduced compared to the common squark mass of the first two generations. (If $\tan\beta \gg 1$, then the pattern of third generation squark masses is somewhat altered; *e.g.*, see Ref. [67].) In addition, there are six flavors of nearly mass-degenerate sleptons (with two slepton eigenstates per flavor for the charged sleptons and one per flavor for the sneutrinos); the sleptons are expected to be somewhat lighter than the mass-degenerate squarks. Finally, third generation squark masses and tau-slepton masses are sensitive to the strength of the respective $\widetilde{f}_L$–$\widetilde{f}_R$ mixing, as discussed below Eq. (3).

Due to the implicit $m_{1/2}$ dependence in the low-energy values of $M^2_{\widetilde{Q}}$, $M^2_{\widetilde{U}}$, and $M^2_{\widetilde{D}}$, there is a tendency for the gluino in mSUGRA models to be lighter than the first- and second-generation squarks. Moreover, the LSP is typically the lightest neutralino, $\widetilde{\chi}^0_1$, which is dominated by its bino component. However, there are some regions of mSUGRA parameter space where the above conclusions do not hold. For example, one can reject those mSUGRA parameter regimes in which the LSP is a chargino. In general, if one imposes the constraints of supersymmetric particle searches and those of cosmology (say, by requiring the LSP to be a suitable dark matter candidate), one obtains significant restrictions to the mSUGRA parameter space. A recent compilation of benchmark mSUGRA points consistent with present data from particle physics and cosmology can be found in Ref. [68].

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. As previously noted, 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 (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. However, the mSUGRA approach is probably too simplistic. Theoretical considerations suggest that the universality of Planck-scale soft-supersymmetry-breaking parameters is not generic [69]. In particular, it is easy to write down effective operators at the Planck scale that do not respect flavor universality, and it is difficult to find a theoretical principle that would forbid them.

In contrast, in gauge-mediated supersymmetry breaking, universality of the fundamental soft-supersymmetry-breaking squark and slepton squared-mass parameters is guaranteed because the supersymmetry-breaking is communicated to the sector of MSSM fields via gauge interactions. In the minimal gauge-mediated supersymmetry-breaking (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 quite model-dependent, and lies somewhat outside the ansatz of gauge-mediated supersymmetry breaking. 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 supersymmetry-breaking dynamics can be very complex, and no complete model of gauge-mediated supersymmetry yet exists that is both simple and compelling.

It was noted in Section I.2 that the gravitino is the LSP in GMSB models. Thus, in such models, the next-to-lightest supersymmetric particle (NLSP) 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 superpartner plus a gravitino (*e.g.*, $\widetilde{\chi}_1^0 \to \gamma\widetilde{g}_{3/2}$, $\widetilde{\chi}_1^0 \to Z\widetilde{g}_{3/2}$ or $\widetilde{\tau}_R^\pm \to \tau^\pm\widetilde{g}_{3/2}$), with lifetimes and branching ratios that depend on the model parameters.

Different choices for the identity of the NLSP and its decay rate lead to a variety of distinctive supersymmetric phenomenologies [38,70]. For example, a long-lived $\widetilde{\chi}_1^0$-NLSP that decays outside collider detectors leads to supersymmetric decay chains with missing energy in association with leptons and/or hadronic jets (this case is indistinguishable from the canonical phenomenology of the $\widetilde{\chi}_1^0$-LSP). On the other hand, if $\widetilde{\chi}_1^0 \to \gamma\widetilde{g}_{3/2}$ is the dominant decay mode, and the decay occurs inside the detector, then nearly *all* supersymmetric particle decay chains would contain a photon. In contrast, the case of a $\widetilde{\tau}_R^\pm$-NLSP would lead either to a new long-lived charged particle (*i.e.*, the $\widetilde{\tau}_R^\pm$) or to supersymmetric particle decay chains with τ leptons.

Finally, grand unification [71] can impose additional constraints on the MSSM parameters. Perhaps one of the most compelling hints for low-energy supersymmetry is the unification of SU(3)×SU(2)×U(1) gauge couplings predicted by models of supersymmetric grand unified theories (SGUTs) [7,9,72](with the supersymmetry-breaking scale of order 1 TeV or below). Gauge coupling unification, which takes place at an energy scale of order 10^{16} GeV, is quite robust [73]. For example, successful unification depends weakly on the details of the theory at the unification scale. In particular, given the low-energy values of the electroweak couplings $g(m_Z)$ and $g'(m_Z)$, one can predict $\alpha_s(m_Z)$ by using the MSSM renormalization group equations to extrapolate to higher energies, and by imposing the unification condition on the three gauge couplings at some high-energy scale, $M_{\rm X}$. This procedure, which fixes $M_{\rm X}$, can be successful (*i.e.*, three running couplings will meet at a single point) only for a unique value of $\alpha_s(m_Z)$. The extrapolation depends somewhat on the low-energy supersymmetric spectrum (so-called low-energy "threshold effects"), and on the SGUT spectrum (high-energy threshold effects), which can somewhat alter the evolution of couplings. Ref. [74] summarizes the comparison of present data with the expectations of SGUTs, and shows that the measured value of $\alpha_s(m_Z)$ is in good agreement with the predictions of supersymmetric grand unification for a reasonable choice of supersymmetric threshold corrections.

Additional SGUT predictions arise through the unification of the Higgs-fermion Yukawa couplings (λ_f). There is some evidence that $\lambda_b = \lambda_\tau$ leads to good low-energy phenomenology [75], and an intriguing possibility that $\lambda_b = \lambda_\tau = \lambda_t$ may be phenomenologically viable [67,76] in the parameter regime where $\tan\beta \simeq m_t/m_b$. Finally, grand unification imposes constraints on the soft-supersymmetry-breaking parameters. For example, gaugino-mass unification leads to the relations given by Eq. (7). Diagonal squark and slepton soft-supersymmetry-breaking scalar masses may also be unified, which is analogous to the unification of Higgs-fermion Yukawa couplings.

In the absence of a fundamental theory of supersymmetry breaking, further progress will require a detailed knowledge of the supersymmetric-particle spectrum in order to determine the nature of the high-energy parameters. Of course, any of the theoretical assumptions described in this section could be wrong and must eventually be tested experimentally.

I.8. Beyond the MSSM: Non-minimal models of low-energy supersymmetry can also be constructed. One approach is to add new structure beyond the Standard Model at the TeV scale or below. The supersymmetric extension of such a theory would be a non-minimal extension of the MSSM. Possible new structures include: (i) the supersymmetric generalization of the see-saw model of neutrino masses [77,78]; (ii) an enlarged electroweak gauge group beyond SU(2)×U(1) [79]; (iii) the addition of new, possibly exotic, matter multiplets [*e.g.*, a vector-like color triplet with electric charge $\frac{1}{3}e$; such states sometimes occur as low-energy remnants in E_6 grand unification models]; and/or (iv) the addition of low-energy SU(3)×SU(2)×U(1) singlets [80]. A possible theoretical motivation for such new structure arises from the study of phenomenologically viable string theory ground states [81].

A second approach is to retain the minimal particle content of the MSSM but remove the assumption of R-parity invariance. The most general R-parity-violating (RPV) theory involving the MSSM spectrum introduces many new parameters to both the supersymmetry-conserving and the supersymmetry-breaking sectors. Each new interaction term violates either B

or L conservation. For example, consider new scalar-fermion Yukawa couplings derived from the following interactions:

$$(\lambda_L)_{pmn}\widehat{L}_p\widehat{L}_m\widehat{E}^c_n+(\lambda'_L)_{pmn}\widehat{L}_p\widehat{Q}_m\widehat{D}^c_n+(\lambda_B)_{pmn}\widehat{U}^c_p\widehat{D}^c_m\widehat{D}^c_n\,, \quad (10)$$

where p, m, and n are generation indices, and gauge group indices are suppressed. In the notation above, $\widehat{Q}$, $\widehat{U}^c$, $\widehat{D}^c$, $\widehat{L}$, and $\widehat{E}^c$ respectively represent $(u,d)_L$, u^c_L, d^c_L, $(\nu,\, e^-)_L$, and e^c_L and the corresponding superpartners. The Yukawa interactions are obtained from Eq. (10) by taking all possible combinations involving two fermions and one scalar superpartner. Note that the term in Eq. (10) proportional to λ_B violates B, while the other two terms violate L. Even if all the terms of Eq. (10) are absent, there is one more possible supersymmetric source of R-parity violation. In the notation of Eq. (10), one can add a term of the form $(\mu_L)_p\widehat{H}_u\widehat{L}_p$, where $\widehat{H}_u$ represents the $Y=1$ Higgs doublet and its higgsino superpartner. This term is the RPV generalization of the supersymmetry-conserving Higgs mass parameter μ of the MSSM, in which the $Y=-1$ Higgs/higgsino super-multiplet $\widehat{H}_d$ is replaced by the lepton/slepton super-multiplet $\widehat{L}_p$. The RPV-parameters $(\mu_L)_p$ also violate L.

Phenomenological constraints on various low-energy B- and L-violating processes can be used to derive limits on each of the coefficients $(\lambda_L)_{pmn}$, $(\lambda'_L)_{pmn}$, and $(\lambda_B)_{pmn}$ taken one at a time [82]. If more than one coefficient is simultaneously non-zero, then the limits are, in general, more complicated. All possible RPV terms cannot be simultaneously present and unsuppressed; otherwise the proton decay rate would be many orders of magnitude larger than the present experimental bound. One way to avoid proton decay is to impose B or L invariance (either one alone would suffice). Otherwise, one must accept the requirement that certain RPV coefficients must be extremely suppressed.

One particularly interesting class of RPV models is one in which B is conserved, but L is violated. It is possible to enforce baryon number conservation, while allowing for lepton number violating interactions by imposing a discrete baryon $\mathbf{Z_3}$ symmetry on the low-energy theory [83], in place of the standard $\mathbf{Z_2}$ R parity. In these models, supersymmetric phenomenology exhibits features that are quite distinct from that of the MSSM. The LSP is no longer stable, which implies that not all supersymmetric decay chains must yield missing-energy events at colliders. Both $\Delta L = 1$ and $\Delta L = 2$ phenomena are allowed (if L is violated), leading to neutrino masses and mixing [84], neutrinoless double-beta decay [85], sneutrino-antisneutrino mixing [78,86,87], and s-channel resonant production of the sneutrino in e^+e^- collisions [88]. Since the distinction between the Higgs and matter super-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. Note that if $\lambda'_L \neq 0$, then squarks can behave as leptoquarks since the following processes are allowed: $e^+\overline{u}_m \to \overline{\widetilde{d}}_n \to e^+\overline{u}_m$, $\overline{\nu}\overline{d}_m$, and $e^+d_m \to \widetilde{u}_n \to e^+d_m$. (As above, m and n are generation labels, so that $d_2 = s$, $d_3 = b$, *etc.*)

Of course, R-parity-violation can also enter via the soft-supersymmetry-breaking terms, leading to an explosion of unknown parameters (well beyond the 124 of the MSSM in the most general case). As in the MSSM, one can consider constrained versions of RPV supersymmetry, in which simplified assumptions are made about the structure of the supersymmetry breaking terms at some high energy scale. Moreover, one can make additional assumptions regarding the RPV parameters. For example, in the bilinear RPV model [89], the trilinear RPV terms of Eq. (10) (and the corresponding supersymmetry-breaking "A-parameters") are absent, and the only source of R-parity violation arises from $(\mu_L)_p$ and L-violating soft-supersymmetry-breaking "B-parameters" and squared-mass terms.

With the overwhelming evidence for neutrino masses and mixing [90], it is clear that any viable supersymmetric model of fundamental particles must incorporate some form of L violation in the low-energy theory [91]. The supersymmetric generalization of the see-saw mechanism and RPV supersymmetry provide two possible frameworks for non-zero neutrino masses.

For example, Ref. [92] demonstrates how one can fit both the solar and atmospheric neutrino data in the bilinear RPV supersymmetric model. In addition, experimental and theoretical constraints from collider physics also places some non-trivial restrictions on general R-parity-violating alternatives to the MSSM (see Refs. [82] and [93] for further details).

References

1. H.P. Nilles, Phys. Reports **110**, 1 (1984).
2. P. Nath, R. Arnowitt, and A.H. Chamseddine, *Applied* $N = 1$ *Supergravity* (World Scientific, Singapore, 1984).
3. S.P. Martin, in *Perspectives on Supersymmetry*, ed. G.L. Kane (World Scientific, Singapore, 1998) pp. 1–98.
4. S. Weinberg, *The Quantum Theory of Fields, Volume III: Supersymmetry* (Cambridge University Press, Cambridge, UK, 2000).
5. L. Girardello and M. Grisaru, Nucl. Phys. **B194**, 65 (1982);
   L.J. Hall and L. Randall, Phys. Rev. Lett. **65**, 2939 (1990);
   I. Jack and D.R.T. Jones, Phys. Lett. **B457**, 101 (1999).
6. E. Witten, Nucl. Phys. **B188**, 513 (1981).
7. S. Dimopoulos and H. Georgi, Nucl. Phys. **B193**, 150 (1981).
8. L. Susskind, Phys. Reports **104**, 181 (1984);
   N. Sakai, Z. Phys. **C11**, 153 (1981);
   R.K. Kaul, Phys. Lett. **109B**, 19 (1982).
9. For a review, see R.N. Mohapatra, in *Particle Physics 1999*, ICTP Summer School in Particle Physics, Trieste, Italy, 21 June—9 July, 1999, eds. G. Senjanovic and A.Yu. Smirnov (World Scientific, Singapore, 2000) pp. 336–394.
10. H.E. Haber and G.L. Kane, Phys. Reports **117**, 75 (1985).
11. K. Inoue *et al.*, Prog. Theor. Phys. **68**, 927 (1982) [erratum: **70,** 330 (1983)]; **71,** 413 (1984);
    R. Flores and M. Sher, Ann. Phys. (NY) **148**, 95 (1983).
12. J.F. Gunion and H.E. Haber, Nucl. Phys. **B272**, 1 (1986) [erratum: **B402**, 567 (1993)].
13. See, *e.g.*, R. Barbieri and G.F. Giudice, Nucl. Phys. **B305**, 63 (1988);
    G.W. Anderson and D.J. Castano, Phys. Lett. **B347**, 300 (1995); Phys. Rev. **D52**, 1693 (1995); Phys. Rev. **D53**, 2403 (1996);

J.L. Feng, K.T. Matchev, and T. Moroi, Phys. Rev. **D61**, 075005 (2000).

14. S. Dimopoulos and G.F. Giudice, Phys. Lett. **B357**, 573 (1995);
A. Pomarol and D. Tommasini, Nucl. Phys. **B466**, 3 (1996);
A.G. Cohen, D.B. Kaplan, and A.E. Nelson, Phys. Lett. **B388**, 588 (1996);
J.L. Feng, K.T. Matchev, and T. Moroi, Phys. Rev. Lett. **84**, 2322 (2000).
15. M. Schmitt, "Supersymmetry Part II (Experiment)," immediately following, in the printed version of the *Review of Particle Physics* (see also the Particle Listings immediately following).
16. See, *e.g.*, F. Gabbiani *et al.*, Nucl. Phys. **B477**, 321 (1996).
17. For a recent review and references to the original literature, see: A. Masiero and O. Vives, New Journal of Physics **4**, 4.1 (2002).
18. J. Erler and P. Langacker, "Electroweak Model and Constraints on New Physics," in the section on Reviews, Tables, and Plots in this *Review*.
19. P.H. Chankowski and S. Pokorski, in *Perspectives on Supersymmetry*, ed. G.L. Kane (World Scientific, Singapore, 1998) pp. 402–422.
20. A. Dobado, M.J. Herrero, and S. Penaranda, Eur. Phys. J. **C7**, 313 (1999); **C12**, 673 (2000); **C17**, 487 (2000).
21. G. Altarelli *et al.*, JHEP **0106**, 018 (2001);
W. de Boer and C. Sander, in the Proceedings of the 10th International Conference on Supersymmetry and Unification of Fundamental Interactions (SUSY-02), DESY Hamburg, Germany, 17–23 June 2002, edited by P. Nath and P.M. Zerwas, (DESY publications, Hamburg, Germany) pp. 1121–1126;
S. Heinemeyer and G. Weiglein, `hep-ph/0307177`, to appear in the Proceedings of the workshop "Electroweak precision data and the Higgs mass,"DESY Zeuthen, February 2003.
22. M. Davier *et al.*, preprint LAL 03-50 (2003) [`hep-ph/0308213`].
23. For a recent review and references to the literature, see T. Hurth, `hep-ph/0212304`, to be published in Rev. Mod. Phys. .
24. See, *e.g.*, M. Ciuchini *et al.*, Phys. Rev. **D67**, 075016 (2003).

25. U. Chattopadhyay and P. Nath, Phys. Rev. **D66**, 093001 (2002);
S.P. Martin and J.D. Wells, Phys. Rev. **D67**, 015002 (2003).

26. P. Fayet, Phys. Lett. **69B**, 489 (1977);
G. Farrar and P. Fayet, Phys. Lett. **76B**, 575 (1978).

27. J. Ellis *et al.*, Nucl. Phys. **B238**, 453 (1984).

28. S. Raby, Phys. Lett. **B422**, 158 (1998);
S. Raby and K. Tobe, Nucl. Phys. **B539**, 3 (1999);
A. Mafi and S. Raby, Phys. Rev. **D62**, 035003 (2000).

29. G. Jungman, M. Kamionkowski, and K. Griest, Phys. Reports **267**, 195 (1996).

30. A.R. Liddle and D.H. Lyth, Phys. Reports **213**, 1 (1993).

31. N.J.C. Spooner and M. Srednicki, in the section on "Dark Matter" in the *Review of Particle Physics*.

32. P. Fayet, Phys. Lett. **84B**, 421 (1979); Phys. Lett. **86B**, 272 (1979).

33. S. Deser and B. Zumino, Phys. Rev. Lett. **38**, 1433 (1977).

34. L.J. Hall, J. Lykken, and S. Weinberg, Phys. Rev. **D27**, 2359 (1983).

35. S.K. Soni and H.A. Weldon, Phys. Lett. **126B**, 215 (1983);
Y. Kawamura, H. Murayama, and M. Yamaguchi, Phys. Rev. **D51**, 1337 (1995).

36. A.B. Lahanas and D.V. Nanopoulos, Phys. Reports **145**, 1 (1987).

37. M. Dine and A.E. Nelson, Phys. Rev. **D48**, 1277 (1993);
M. Dine, A.E. Nelson, and Y. Shirman, Phys. Rev. **D51**, 1362 (1995);
M. Dine *et al.*, Phys. Rev. **D53**, 2658 (1996).

38. G.F. Giudice, and R. Rattazzi, Phys. Reports **322**, 419 (1999).

39. J. Polchinski, *String Theory, Volumes I and II* (Cambridge University Press, Cambridge, UK, 1998).

40. For a review of recent developments in models and the phenomenology of large extra dimensions, see J. Hewett and J. March-Russell, in the section on "Extra Dimensions" in the *Review of Particle Physics*.

41. These ideas are reviewed in: V.A. Rubakov, Sov. Phys. Usp. **44**, 871 (2001);
Y.A. Kubyshin, Lectures given at the 11th International School on Particles and Cosmology, Karbardino-Balkaria, Russia, 18–24 April 2001, `hep-ph/0111027`.

42. L. Randall and R. Sundrum, Nucl. Phys. **B557**, 79 (1999).

43. Z. Chacko, M.A. Luty, and E. Ponton, JHEP **0007**, 036 (2000);
D.E. Kaplan, G.D. Kribs, and M. Schmaltz, Phys. Rev. **D62**, 035010 (2000);
Z. Chacko *et al.*, JHEP **0001**, 003 (2000).

44. M. Quiros, to appear in the Proceedings of the 2002 Theoretical Advanced Study Institute (TASI-02), Boulder, CO, 3–28 June 2002 [hep-ph/0302189].

45. J. Scherk and J.H. Schwarz, Phys. Lett. **82B**, 60 (1979); Nucl. Phys. **B153**, 61 (1979).

46. See, *e.g.*, R. Barbieri, L.J. Hall, and Y. Nomura, Phys. Rev. **D66**, 045025 (2002); Nucl. Phys. **B624**, 63 (2002).

47. H.E. Haber, in *Recent Directions in Particle Theory, Proceedings of the 1992 Theoretical Advanced Study Institute in Particle Physics*, eds. J. Harvey and J. Polchinski (World Scientific, Singapore, 1993) pp. 589–686.

48. S. Dimopoulos and D. Sutter, Nucl. Phys. **B452**, 496 (1995);
D.W. Sutter, Stanford Ph. D. thesis, hep-ph/9704390.

49. H.E. Haber, Nucl. Phys. B (Proc. Suppl.) **62A-C**, 469 (1998).

50. Explicit forms for the chargino and neutralino mass matrices can be found in Appendix A of Ref. [12]; see also Ref. [47].

51. J. Ellis and S. Rudaz, Phys. Lett. **128B**, 248 (1983).

52. D.M. Pierce *et al.*, Nucl. Phys. **B491**, 3 (1997).

53. J.F. Gunion *et al.*, *The Higgs Hunter's Guide* (Perseus Publishing, Cambridge, MA, 1990);
M. Carena and H.E. Haber, Prog. in Part. Nucl. Phys. **50**, 63 (2003).

54. H.E. Haber and R. Hempfling, Phys. Rev. Lett. **66**, 1815 (1991);
Y. Okada, M. Yamaguchi, and T. Yanagida, Prog. Theor. Phys. **85**, 1 (1991);
J. Ellis, G. Ridolfi, and F. Zwirner, Phys. Lett. **B257**, 83 (1991).

55. ALEPH, DELPHI, L3 and OPAL Collaborations [LEP Higgs Working Group for Higgs boson searches], Phys. Lett. **B565**, 61 (2003).

56. See, *e.g.*, G. Degrassi *et al.*, Eur. Phys. J. **C28**, 133 (2003) and references contained therein.

57. A. Pilaftsis and C.E.M. Wagner, Nucl. Phys. **B553**, 3 (1999);

D.A. Demir, Phys. Rev. **D60**, 055006 (1999);
S.Y. Choi, M. Drees, and J.S. Lee, Phys. Lett. **B481**, 57 (2000);
M. Carena *et al.*, Nucl. Phys. **B586**, 92 (2000); Phys. Lett. **B495**, 155 (2000); Nucl. Phys. **B625**, 345 (2002).

58. S. Khalil, Int. J. Mod. Phys. **A18**, 1697 (2003).
59. W. Fischler, S. Paban, and S. Thomas, Phys. Lett. **B289**, 373 (1992);
S.M. Barr, Int. J. Mod. Phys. **A8**, 209 (1993);
T. Ibrahim and P. Nath, Phys. Rev. **D58**, 111301 (1998) [erratum: **D60**, 099902 (1999)];
M. Brhlik, G.J. Good, and G.L. Kane, Phys. Rev. **D59**, 115004 (1999).
60. A. Masiero and L. Silvestrini, in *Perspectives on Supersymmetry*, ed. G.L. Kane (World Scientific, Singapore, 1998) pp. 423–441.
61. H. Georgi, Phys. Lett. **169B**, 231 (1986);
L.J. Hall, V.A. Kostelecky, and S. Raby, Nucl. Phys. **B267**, 415 (1986).
62. Y. Nir and N. Seiberg, Phys. Lett. **B309**, 337 (1993);
S. Dimopoulos, G.F. Giudice, and N. Tetradis, Nucl. Phys. **B454**, 59 (1995);
G.F. Giudice *et al.*, JHEP **12**, 027 (1998);
J.L. Feng and T. Moroi, Phys. Rev. **D61**, 095004 (2000).
63. J.F. Gunion and H.E. Haber, Phys. Rev. **D37**, 2515 (1988).
64. J.L. Feng *et al.*, Phys. Rev. Lett. **83**, 1731 (1999);
T. Gherghetta, G.F. Giudice, and J.D. Wells, Nucl. Phys. **B559**, 27 (1999);
J.F. Gunion and S. Mrenna, Phys. Rev. **D62**, 015002 (2000).
65. See, *e.g.*, B. Murakami and J.D. Wells, Phys. Rev. **D68**, 035006 (2003) and references contained therein.
66. M. Drees and S.P. Martin, in *Electroweak Symmetry Breaking and New Physics at the TeV Scale*, eds. T. Barklow *et al.* (World Scientific, Singapore, 1996) pp. 146–215.
67. M. Carena *et al.*, Nucl. Phys. **B426**, 269 (1994).
68. M. Battaglia *et al.*, CERN-TH/2003-138 [hep-ph/0306219].
69. L.E. Ibáñez and D. Lüst, Nucl. Phys. **B382**, 305 (1992);
B. de Carlos, J.A. Casas, and C. Munoz, Phys. Lett. **B299**, 234 (1993);
V. Kaplunovsky and J. Louis, Phys. Lett. **B306**, 269 (1993);

A. Brignole, L.E. Ibáñez, and C. Munoz, Nucl. Phys. **B422**, 125 (1994) [erratum: **B436**, 747 (1995)].

70. For a review and guide to the literature, see J.F. Gunion and H.E. Haber, in *Perspectives on Supersymmetry*, ed. G.L. Kane (World Scientific, Singapore, 1998) pp. 235–255.

71. S. Raby , in the section on "Grand Unified Theories" in the *Review of Particle Physics*.

72. M.B. Einhorn and D.R.T. Jones, Nucl. Phys. **B196**, 475 (1982);
W.J. Marciano and G. Senjanovic, Phys. Rev. **D25**, 3092 (1982).

73. D.M. Ghilencea and G.G. Ross, Nucl. Phys. **B606**, 101 (2001).

74. S. Pokorski, Acta Phys. Polon. **B30**, 1759 (1999);
For a review, see N. Polonsky, *Supersymmetry: Structure and phenomena. Extensions of the standard model*, Lect. Notes Phys. **M68**, 1 (2001).

75. H. Arason *et al.*, Phys. Rev. Lett. **67**, 2933 (1991);
Phys. Rev. **D46**, 3945 (1992);
V. Barger, M.S. Berger, and P. Ohmann, Phys. Rev. **D47**, 1093 (1993);
M. Carena, S. Pokorski, and C.E.M. Wagner, Nucl. Phys. **B406**, 59 (1993);
P. Langacker and N. Polonsky, Phys. Rev. **D49**, 1454 (1994).

76. M. Olechowski and S. Pokorski, Phys. Lett. **B214**, 393 (1988);
B. Ananthanarayan, G. Lazarides, and Q. Shafi, Phys. Rev. **D44**, 1613 (1991);
S. Dimopoulos, L.J. Hall, and S. Raby, Phys. Rev. Lett. **68**, 1984 (1992);
L.J. Hall, R. Rattazzi, and U. Sarid, Phys. Rev. **D50**, 7048 (1994);
R. Rattazzi and U. Sarid, Phys. Rev. **D53**, 1553 (1996).

77. J. Hisano *et al.*, Phys. Lett. **B357**, 579 (1995);
J. Hisano *et al.*, Phys. Rev. **D53**, 2442 (1996);
J. Ellis *et al.*, Phys. Rev. **D66**, 115013 (2002).

78. Y. Grossman and H.E. Haber, Phys. Rev. Lett. **78**, 3438 (1997).

79. J.L. Hewett and T.G. Rizzo, Phys. Reports **183**, 193 (1989).

80. See, *e.g.*, U. Ellwanger, M. Rausch de Traubenberg, and C.A. Savoy, Nucl. Phys. **B492**, 21 (1997);

U. Ellwanger and C. Hugonie, Eur. Phys. J. **C25**, 297 (2002) and references contained therein.

81. K.R. Dienes, Phys. Reports **287**, 447 (1997).
82. H. Dreiner, in *Perspectives on Supersymmetry*, ed. G.L. Kane (World Scientific, Singapore, 1998) pp. 462–479.
83. L.E. Ibáñez and G.G. Ross, Nucl. Phys. **B368**, 3 (1992);
L.E. Ibáñez, Nucl. Phys. **B398**, 301 (1993).
84. For a review, see J.C. Romao, Nucl. Phys. Proc. Suppl. **81**, 231 (2000).
85. R.N. Mohapatra, Phys. Rev. **D34**, 3457 (1986);
K.S. Babu and R.N. Mohapatra, Phys. Rev. Lett. **75**, 2276 (1995);
M. Hirsch, H.V. Klapdor-Kleingrothaus, and S.G. Kovalenko, Phys. Rev. Lett. **75**, 17 (1995); Phys. Rev. **D53**, 1329 (1996).
86. M. Hirsch, H.V. Klapdor-Kleingrothaus, and S.G. Kovalenko, Phys. Lett. **B398**, 311 (1997).
87. Y. Grossman and H.E. Haber, Phys. Rev. **D59**, 093008 (1999).
88. S. Dimopoulos and L.J. Hall, Phys. Lett. **B207**, 210 (1988);
J. Kalinowski *et al.*, Phys. Lett. **B406**, 314 (1997);
J. Erler, J.L. Feng, and N. Polonsky, Phys. Rev. Lett. **78**, 3063 (1997).
89. A study of the phenomenology of bilinear R-parity-breaking supersymmetry and a guide to the literature can be found in D.A. Restrepo Quintero, `hep-ph/0111198`.
90. See the section on neutrinos in "Particle Listings" in the *Review of Particle Physics.*
91. For a recent review of neutrino masses in supersymmetry, see B. Mukhopadhyaya, `hep-ph/0301278`.
92. See, *e.g.*, M. Hirsch *et al.*, Phys. Rev. **D62**, 113008 (2000) [erratum: **D65**, 119901 (2002)];
M.A. Diaz, *et al.*, Phys. Rev. **D68**, 013009 (2003).
93. M. Bisset *et al.*, Phys. Rev. **D62**, 035001 (2000);
R. Barbier *et al.*, Report of the group on the R-parity violation, `hep-ph/9810232` (1998).