

## SUPERSYMMETRY, PART I (THEORY)

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- I.1. Introduction
- I.2. Structure of the MSSM
  - I.2.1. R-parity and the lightest supersymmetric particle
  - I.2.2. The goldstino and gravitino
  - I.2.3. Hidden sectors and the structure of supersymmetry-breaking
  - I.2.4. Supersymmetry and extra dimensions
  - I.2.5. Split-supersymmetry
- I.3. Parameters of the MSSM
  - I.3.1. The supersymmetry-conserving parameters
  - I.3.2. The supersymmetry-breaking parameters
  - I.3.3. MSSM-124
- I.4. The supersymmetric-particle spectrum
  - I.4.1. The charginos and neutralinos
  - I.4.2. The squarks, sleptons and sneutrinos
- I.5. The supersymmetric Higgs sector
  - I.5.1. The tree-level Higgs sector
  - I.5.2. The radiatively-corrected Higgs sector
- I.6. Restricting the MSSM parameter freedom
  - I.6.1. Gaugino mass unification
  - I.6.2. The constrained MSSM: mSUGRA, CMSSM, ...
  - I.6.3. Gauge mediated supersymmetry breaking
  - I.6.4. The phenomenological MSSM
- I.7. Experimental data confronts the MSSM
  - I.7.1. Naturalness constraints and the little hierarchy
  - I.7.2. Constraints from virtual exchange of supersymmetric particles
- I.8. Massive neutrinos in low-energy supersymmetry
  - I.8.1. The supersymmetric seesaw
  - I.8.2. R-parity-violating supersymmetry
- I.9. Extensions beyond the MSSM

**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 [1]. The existence of such a non-trivial extension of the Poincaré symmetry of ordinary quantum field theory was initially surprising, and its form is highly constrained by theoretical principles [2]. Supersymmetry also provides a framework for the unification of particle physics and gravity [3–6] at the Planck energy scale,  $M_{\text{P}} \approx 10^{19}$  GeV, where the gravitational interactions become comparable in magnitude to the gauge interactions. Moreover, supersymmetry can provide an explanation of the large hierarchy between the energy scale that characterizes electroweak symmetry breaking (of order 100 GeV) and the Planck scale [7–10]. The stability of this large *gauge hierarchy* with respect to radiative quantum corrections is not possible to maintain in the Standard Model without an *unnatural* fine-tuning of the parameters of the fundamental theory at the Planck scale. In contrast, in a supersymmetric extension of the Standard Model, it is possible to maintain the gauge hierarchy with no fine-tuning of parameters, and provide a natural framework for elementary scalar fields.

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 superpartners have not (yet) been observed, supersymmetry must be a broken symmetry. Nevertheless, the stability of the gauge hierarchy can still be maintained if the supersymmetry breaking is *soft* [11,12], and the corresponding supersymmetry-breaking mass parameters are no larger than a few TeV. Whether this is still plausible in light of recent supersymmetry searches at the LHC [13] will be discussed in Section I.7.

In particular, soft-supersymmetry-breaking terms of the Lagrangian involve combinations of fields with total mass dimension of three or less, with some restrictions on the dimension-three terms as elucidated in Ref. 11. The impact of the soft terms becomes negligible at energy scales much larger than the size of the supersymmetry-breaking masses. Thus, a theory of

*weak-scale supersymmetry*, where the effective scale of supersymmetry breaking is tied to the scale of electroweak symmetry breaking, provides a *natural* framework for the origin and the stability of the gauge hierarchy [7–10].

The Standard Model cannot be the correct theory of fundamental particles and their interactions (applicable at all energy scales). However, no unambiguous experimental results currently exist that imply that the Standard Model breaks down at the TeV scale. The expectations of new physics beyond the Standard Model at the TeV scale are based primarily on three theoretical arguments. First, a natural explanation of the gauge hierarchy demands new physics at the TeV scale [10]. Second, the unification of the three Standard Model gauge couplings at a very high energy close to the Planck scale is possible if new physics beyond the Standard Model (which modifies the running of the gauge couplings above the electroweak scale) is present. The minimal supersymmetric extension of the Standard Model (MSSM), where superpartner masses lie below a few TeV, provides an example of successful gauge coupling unification [14]. Third, the existence of dark matter, which makes up approximately one quarter of the energy density of the universe, cannot be explained within the Standard Model of particle physics [15]. Remarkably, a stable weakly-interacting massive particle (WIMP) whose mass and interaction rate are governed by new physics associated with the TeV-scale can be consistent with the observed density of dark matter (this is the so-called *WIMP miracle*, which is reviewed in Ref. 16). The lightest supersymmetric particle is a promising (although not the unique) candidate for the dark matter [17–21]. Further aspects of dark matter can be found in Ref. 22.

Another phenomenon not explained by the Standard Model is the origin of the matter–antimatter asymmetry of the universe [23]. Models of baryogenesis must satisfy the three Sakharov conditions [24]:  $C$  and  $CP$  violation, baryon number violation and a departure from thermal equilibrium. For example, the matter–antimatter asymmetry in the early universe can be generated at the electroweak phase transition if the transition is sufficiently first-order [25]. These conditions are

not satisfied in the Standard Model, since the  $CP$  violation is too small and the phase transition is not strongly first-order [25]. In contrast, it is possible to satisfy these conditions in supersymmetric extensions of the Standard Model, where new sources of CP-violation exist and supersymmetric loops provide corrections to the temperature-dependent effective potential that can render the transition sufficiently first-order. The MSSM parameter space in which electroweak baryogenesis occurs is strongly constrained by LHC data [26]. However, extended supersymmetric models provide new opportunities for successful electroweak baryogenesis [27].

***I.2. Structure of the MSSM:*** The minimal supersymmetric extension of the Standard Model consists of the fields of the two-Higgs-doublet extension of the Standard Model and the corresponding supersymmetric partners [28,29]. A particle and its superpartner together form a supermultiplet. The corresponding field content of the supermultiplets of the MSSM and their gauge quantum numbers are shown in Table 1. The electric charge  $Q = T_3 + \frac{1}{2}Y$  is determined in terms of the third component of the weak isospin ( $T_3$ ) and the U(1) weak hypercharge ( $Y$ ).

The gauge supermultiplets consist of the gluons and their *gluino* fermionic superpartners and the  $SU(2) \times U(1)$  gauge bosons and their *gaugino* fermionic superpartners. The matter supermultiplets consist of three generations of left-handed and right-handed quarks and leptons, their scalar superpartners (*squarks* and *sleptons*), and the corresponding antiparticles. The Higgs supermultiplets consist of two complex Higgs doublets, their *higgsino* fermionic superpartners, and the corresponding antiparticles. The enlarged Higgs sector of the MSSM constitutes the minimal structure needed to guarantee the cancellation of anomalies from the introduction of the higgsino superpartners. Moreover, without a second Higgs doublet, one cannot generate mass for both “up”-type and “down”-type quarks (and charged leptons) in a way consistent with the underlying supersymmetry [30–32].

The fields of a given supermultiplet are components of a *superfield*. Vector superfields contain the gauge boson fields and

**Table 1:** The fields of the MSSM and their  $SU(3) \times SU(2) \times U(1)$  quantum numbers are listed. For simplicity, only one generation of quarks and leptons is exhibited. For each lepton, quark, and Higgs super-multiplet, there is a corresponding anti-particle multiplet of charge-conjugated fermions and their associated scalar partners.

Field Content of the MSSM						
Super-multiplets	Super-field	Bosonic fields	Fermionic partners	SU(3)	SU(2)	U(1)
gluon/gluino	$\widehat{V}_8$	$g$	$\widetilde{g}$	8	1	0
gauge/	$\widehat{V}$	$W^\pm, W^0$	$\widetilde{W}^\pm, \widetilde{W}^0$	1	3	0
gaugino	$\widehat{V}'$	$B$	$\widetilde{B}$	1	1	0
slepton/	$\widehat{L}$	$(\widetilde{\nu}_L, \widetilde{e}_L^-)$	$(\nu, e^-)_L$	1	2	-1
lepton	$\widehat{E}^c$	$\widetilde{e}_R^-$	$e_R^-$	1	1	-2
squark/	$\widehat{Q}$	$(\widetilde{u}_L, \widetilde{d}_L)$	$(u, d)_L$	3	2	1/3
quark	$\widehat{U}^c$	$\widetilde{u}_R$	$u_R$	3	1	4/3
	$\widehat{D}^c$	$\widetilde{d}_R$	$d_R$	3	1	-2/3
Higgs/	$\widehat{H}_d$	$(H_d^0, H_d^-)$	$(\widetilde{H}_d^0, \widetilde{H}_d^-)$	1	2	-1
higgsino	$\widehat{H}_u$	$(H_u^+, H_u^0)$	$(\widetilde{H}_u^+, \widetilde{H}_u^0)$	1	2	1

their gaugino partners. Chiral superfields contain the spin-0 and spin-1/2 fields of the matter or Higgs supermultiplets. A general supersymmetric Lagrangian is determined by three functions of the chiral superfields [6]: the superpotential, the Kähler potential, and the gauge kinetic function (which can be appropriately generalized to accommodate higher derivative terms [33]). Minimal forms for the Kähler potential and gauge kinetic function, which generate canonical kinetic energy terms for all the fields, are required for *renormalizable* globally supersymmetric theories. A renormalizable superpotential, which is at most cubic in the chiral superfields, yields supersymmetric Yukawa couplings and mass terms. A combination of gauge invariance and supersymmetry produces couplings of gaugino fields to matter (or Higgs) fields and their corresponding superpartners. The (renormalizable) MSSM Lagrangian is then constructed by including all possible supersymmetric interaction terms (of dimension four or less) that satisfy  $SU(3) \times SU(2) \times U(1)$  gauge invariance and  $B-L$  conservation (where  $B$  = baryon number

and  $L =$  lepton number). Finally, the most general soft-supersymmetry-breaking terms consistent with these symmetries are added [11,12,34].

Although the MSSM is the focus of much of this review, there is some motivation for considering non-minimal supersymmetric extensions of the Standard Model. For example, extra structure is needed to generate non-zero neutrino masses as discussed in Section I.8. In addition, in order to address some theoretical issues and tensions associated with the MSSM, it has been fruitful to introduce one additional singlet Higgs superfield. The resulting next-to-minimal supersymmetric extension of the Standard Model (NMSSM) [35] is considered further in Sections I.5, I.7 and I.9. Finally, one is always free to add additional fields to the Standard Model along with the corresponding superpartners. However, only certain choices for the new fields (*e.g.*, the addition of complete  $SU(5)$  multiplets) will preserve the successful gauge coupling unification. Some examples will be briefly mentioned in Section I.9.

***I.2.1. R-parity and the lightest supersymmetric particle:*** 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  $S$  [36]. 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, any initial state in a scattering experiment will involve ordinary (R-even) particles. Consequently, it follows that supersymmetric particles must be produced in pairs. In general, these particles are highly unstable and decay into lighter states. Moreover, 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 [19]. 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, as noted in Section I.1 and reviewed in Refs. [20,21], the stability of the LSP in R-parity-conserving supersymmetry makes it a promising candidate for dark matter.

**I.2.2. The goldstino and gravitino:** In the MSSM, supersymmetry breaking is accomplished by including the most general renormalizable soft-supersymmetry-breaking terms consistent with the  $SU(3)\times SU(2)\times U(1)$  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* ( $\tilde{G}_{1/2}$ ) must exist. The goldstino would then be the LSP, and could play an important role in supersymmetric phenomenology [37].

However, the goldstino degrees of freedom are physical only in models of spontaneously-broken global supersymmetry. If supersymmetry is a local symmetry, then the theory must incorporate gravity; the resulting theory is called supergravity [38]. In models of spontaneously-broken supergravity, the goldstino is “absorbed” by the *gravitino* ( $\tilde{G}$ ) [sometimes called  $\tilde{g}_{3/2}$  in the older literature], the spin-3/2 superpartner of the graviton [39]. By this super-Higgs mechanism, the goldstino is removed from the physical spectrum and the gravitino acquires a mass ( $m_{3/2}$ ). In processes with center-of-mass energy  $E \gg m_{3/2}$ , the goldstino–gravitino equivalence theorem [40] states that the interactions of the helicity  $\pm\frac{1}{2}$  gravitino (whose properties approximate those of the goldstino) dominate those of the helicity  $\pm\frac{3}{2}$  gravitino. The interactions of gravitinos with other light fields can be described by a low-energy effective Lagrangian that is determined by fundamental principles [41].

**I.2.3. Hidden sectors and the structure of supersymmetry breaking:** 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. An alternative 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 [34]. There are no renormalizable tree-level interactions between particles of the visible and hidden sectors. Supersymmetry breaking is assumed to originate in the hidden sector, and its effects are transmitted to the MSSM by some mechanism (often involving the mediation by particles that comprise an additional *messenger* sector). Two theoretical scenarios that exhibit this structure are 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 [42–46]. More precisely, supersymmetry breaking is mediated by effects of gravitational strength (suppressed by inverse powers 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 [3,47]. Such a gravitino typically plays no role in supersymmetric phenomenology at colliders (except perhaps indirectly in the case where the gravitino is the LSP [48]).

Under certain theoretical assumptions on the structure of the Kähler potential (the so-called sequestered form introduced in Ref. 49), supersymmetry breaking is due entirely to the super-conformal (super-Weyl) anomaly, which is common to all supergravity models [49]. In this case, gaugino masses are radiatively generated at one-loop and squark and slepton masses are flavor-diagonal. This approach is called *anomaly-mediated* supersymmetry breaking (AMSB). Indeed, anomaly mediation is more generic than originally conceived, and provides a ubiquitous source of supersymmetry breaking [50]. However in the simplest formulation of AMSB as applied to the MSSM, the squared-masses of the sleptons are negative. This is the so-called tachyonic slepton problem. It may be possible to cure



this fatal flaw in approaches beyond the minimal supersymmetric model [51]. Alternatively, one can assert that anomaly mediation is not the sole source of supersymmetry breaking, at least in the slepton sector.

In *gauge mediated* supersymmetry breaking (GMSB), gauge forces transmit the supersymmetry breaking to the MSSM. A typical structure of such models involves a hidden sector where supersymmetry is broken, a messenger sector consisting of particles (messengers) with nontrivial  $SU(3) \times SU(2) \times U(1)$  quantum numbers, and the visible sector consisting of the fields of the MSSM [52–54]. The direct coupling of the messengers to the hidden sector generates a supersymmetry-breaking spectrum in the messenger sector. Supersymmetry breaking is then transmitted to the MSSM via the virtual exchange of the messengers. In models of *direct gauge mediation*, there is no hidden sector. In particular, the sector in which the supersymmetry breaking originates includes fields that carry nontrivial Standard Model quantum numbers, which allows for the direct transmission of supersymmetry breaking to the MSSM [55]. In this case, no separate messenger sector distinct from the fundamental supersymmetry-breaking sector is required.

The gravitino mass in models of gauge mediated supersymmetry breaking is typically in the eV range (although in some cases it can be as large as a few GeV), which implies that  $\tilde{G}$  is the LSP [17]. In particular, the gravitino is a potential dark matter candidate (for a review and guide to the literature, see Ref. 21). Big bang nucleosynthesis also provides some interesting constraints on the gravitino and the properties of the next-to-lightest supersymmetric particle that decays into the gravitino LSP [56]. The couplings of the helicity  $\pm\frac{1}{2}$  components of  $\tilde{G}$  to the particles of the MSSM (which approximate those of the goldstino as previously noted in Section I.2.2) are significantly stronger than gravitational strength and amenable to experimental collider analyses.

The concept of a hidden sector is more general than supersymmetry. *Hidden valley* models [57] posit the existence of a hidden sector of new particles and interactions that are very weakly coupled to particles of the Standard Model. The

impact of a hidden valley on supersymmetric phenomenology at colliders can be significant if the LSP lies in the hidden sector [58].

***1.2.4. Supersymmetry and extra dimensions:***

Approaches to supersymmetry breaking have also been developed in the context of theories in which the number of space dimensions is greater than three. In particular, a number of supersymmetry-breaking mechanisms have been proposed that are inherently extra-dimensional [59]. The size of the extra dimensions can be significantly larger than  $M_{\text{P}}^{-1}$ ; in some cases of order  $(\text{TeV})^{-1}$  or even larger [60,61].

For example, in one approach the fields of the MSSM live on some brane (a lower-dimensional manifold embedded in a higher-dimensional spacetime), while the sector of the theory that breaks supersymmetry lives on a second spatially-separated brane. Two examples of this approach are anomaly-mediated supersymmetry breaking [49] and gaugino-mediated supersymmetry breaking [62]. 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).

Alternatively, one can consider a higher-dimensional theory that 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. This is the so-called Scherk-Schwarz mechanism [63]. The phenomenology of such models can be strikingly different from that of the usual MSSM [64].

***1.2.5. Split-supersymmetry:*** If supersymmetry is not connected with the origin of the electroweak scale, it may still be possible that some remnant of the superparticle spectrum survives down to the TeV-scale or below. This is the idea of *split-supersymmetry* [65], in which scalar superpartners of the quarks and leptons are significantly heavier (perhaps by many orders of magnitude) than 1 TeV, whereas the fermionic superpartners of the gauge and Higgs bosons have masses on the

order of 1 TeV or below (protected by some chiral symmetry). With the exception of a single light neutral scalar whose properties are practically indistinguishable from those of the Standard Model Higgs boson, all other Higgs bosons are also assumed to be very heavy.

The recent observation of a Higgs boson at a mass of 126 GeV [66,67] implies that there is an upper bound in the MSSM on the mass scale that characterizes the top squarks in the range of 10 to  $10^6$  TeV [69,70]. It is not difficult to formulate models in which gaugino masses are one-loop suppressed relative to the masses of the squarks and sleptons. The higgsino mass scale may or may not be likewise suppressed depending on the details of the model [71]. Recently, a number of authors have considered a supersymmetric spectrum with the squarks, slepton and the heavy Higgs masses around  $10^3$  TeV, whereas gaugino masses are around 1 TeV [70,71]. In this scenario, only the lighter gaugino states are kinematically accessible at the LHC. The one light Higgs scalar of the model (identified with the observed scalar state at 126 GeV) is indistinguishable from the Standard Model Higgs boson.

The supersymmetry breaking required to produce such a split-supersymmetry spectrum would destabilize the gauge hierarchy, and thus yield no natural explanation for the scale of electroweak symmetry breaking. Nevertheless, models of split-supersymmetry can account for the dark matter (which is assumed to be the LSP gaugino or higgsino) and gauge coupling unification, thereby preserving two of the good features of weak-scale supersymmetry.

***1.3. Parameters of the MSSM:*** The parameters of the MSSM are conveniently described by considering separately the supersymmetry-conserving and the supersymmetry-breaking sectors. A careful discussion of the conventions used here in defining the tree-level MSSM parameters can be found in Ref. 72. For simplicity, consider first the case of one generation of quarks, leptons, and their scalar superpartners.

***1.3.1. The supersymmetric-conserving parameters:***

The parameters of the supersymmetry-conserving sector consist of: (i) gauge couplings,  $g_s$ ,  $g$ , and  $g'$ , corresponding to the

Standard Model gauge group  $SU(3) \times SU(2) \times U(1)$  respectively; (ii) a supersymmetry-conserving higgsino mass parameter  $\mu$ ; and (iii) Higgs-fermion Yukawa coupling constants,  $\lambda_u$ ,  $\lambda_d$ , and  $\lambda_e$ , corresponding to the coupling of one generation of left- and right-handed quarks and leptons, and their superpartners to the Higgs bosons and higgsinos. Because there is no right-handed neutrino (and its superpartner) in the MSSM as defined here, a Yukawa coupling  $\lambda_\nu$  is not included here. The complex  $\mu$  parameter and Yukawa couplings enter via the most general renormalizable R-parity-conserving superpotential,

$$W = \lambda_d \widehat{H}_d \widehat{Q} \widehat{D}^c - \lambda_u \widehat{H}_u \widehat{Q} \widehat{U}^c + \lambda_e \widehat{H}_d \widehat{L} \widehat{E}^c + \mu \widehat{H}_u \widehat{H}_d, \quad (1)$$

where the superfields are defined in Table 1 and the gauge group indices are suppressed. The reader is warned that in the literature,  $\mu$  is sometimes defined with the opposite sign to the one given in Eq. (1).

***1.3.2. The supersymmetry-breaking parameters:***

The supersymmetry-breaking sector contains the following sets of parameters: (i) three complex gaugino Majorana mass parameters,  $M_3$ ,  $M_2$ , and  $M_1$ , associated with the  $SU(3)$ ,  $SU(2)$ , and  $U(1)$  subgroups of the Standard Model; (ii) five diagonal scalar squared-mass parameters for the squarks and sleptons,  $M_{\widehat{Q}}^2$ ,  $M_{\widehat{U}}^2$ ,  $M_{\widehat{D}}^2$ ,  $M_{\widehat{L}}^2$ , and  $M_{\widehat{E}}^2$  [corresponding to the five electroweak gauge multiplets, *i.e.*, superpartners of  $(u, d)_L$ ,  $u_L^c$ ,  $d_L^c$ ,  $(\nu, e^-)_L$ , and  $e_L^c$ , where the superscript  $c$  indicates a charge-conjugated fermion]; and (iii) three Higgs-squark-squark and Higgs-slepton-slepton trilinear interaction terms, with complex coefficients  $\lambda_u A_U$ ,  $\lambda_d A_D$ , and  $\lambda_e A_E$  (which define the so-called “ $A$ -parameters”). It is traditional to factor out the Yukawa couplings in the definition of the  $A$ -parameters (originally motivated by a simple class of gravity mediated supersymmetry-breaking models [3,5]). If the  $A$ -parameters defined as above are parametrically of the same order (or smaller) relative to other supersymmetry-breaking mass parameters, then only the third generation  $A$ -parameters are phenomenologically relevant. The reader is warned that the convention for the overall sign of the  $A$ -parameters varies in the literature.

Finally, we have (iv) three scalar squared-mass parameters: two of which ( $m_1^2$  and  $m_2^2$ ) are real parameters that contribute to the diagonal Higgs squared-masses, given by  $m_1^2 + |\mu|^2$  and  $m_2^2 + |\mu|^2$ , and a third that contributes to the off-diagonal Higgs squared-mass term,  $m_{12}^2 \equiv B\mu$  (which defines the complex “ $B$ -parameter”). The breaking of the electroweak symmetry  $SU(2) \times U(1)$  to  $U(1)_{EM}$  is only possible after introducing the supersymmetry-breaking Higgs squared-mass parameters. Minimizing the resulting tree-level Higgs scalar potential, 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 the CP-odd Higgs mass  $m_A$  (cf. Section I.5). Here,  $v_d$  [ $v_u$ ] is the vacuum expectation value of the neutral component of the Higgs field  $H_d$  [ $H_u$ ] that couples exclusively to down-type (up-type) quarks and leptons.

Note that  $v_d^2 + v_u^2 = 2m_W^2/g^2 \simeq (174 \text{ GeV})^2$  is fixed by the  $W$  mass and the  $SU(2)$  gauge coupling, whereas the ratio

$$\tan \beta = v_u/v_d \tag{2}$$

is a free parameter. It is convenient to choose the phases of the Higgs fields such that  $m_{12}^2$  is real and non-negative. In this case, we can adopt a convention where  $0 \leq \beta \leq \pi/2$ . The tree-level conditions for the scalar potential minimum relate the diagonal and off-diagonal Higgs squared-masses in terms of  $m_Z^2 = \frac{1}{4}(g^2 + g'^2)(v_d^2 + v_u^2)$ , the angle  $\beta$  and the CP-odd Higgs mass  $m_A$ :

$$\sin 2\beta = \frac{2m_{12}^2}{m_1^2 + m_2^2 + 2|\mu|^2} = \frac{2m_{12}^2}{m_A^2}, \tag{3}$$

$$\frac{1}{2}m_Z^2 = -|\mu|^2 + \frac{m_1^2 - m_2^2 \tan^2 \beta}{\tan^2 \beta - 1}. \tag{4}$$

One must also guard against the existence of charge and/or color breaking global minima due to non-zero vacuum expectation values for the squark and charged slepton fields. This possibility can be avoided if the  $A$ -parameters are not unduly large [43,73].

Note that supersymmetry-breaking mass terms for the fermionic superpartners of scalar fields and non-holomorphic

trilinear scalar interactions (*i.e.*, interactions that mix scalar fields and their complex conjugates) have not been included above in the soft-supersymmetry-breaking sector. These terms can potentially destabilize the gauge hierarchy [11] in models with gauge-singlet superfields. The latter are not present in the MSSM; hence as noted in Ref. 12, these so-called non-standard soft-supersymmetry-breaking terms are benign. However, the coefficients of these terms (which have dimensions of mass) are expected to be significantly suppressed compared to the TeV-scale in a fundamental theory of supersymmetry-breaking [74]. Consequently, we follow the usual approach and omit these terms from further consideration.

**I.3.3. MSSM-124:** The total number of independent physical parameters that define the MSSM (in its most general form) is quite large, primarily due to the soft-supersymmetry-breaking sector. In particular, in the case of three generations of quarks, leptons, and their superpartners,  $M_{\tilde{Q}}^2$ ,  $M_{\tilde{U}}^2$ ,  $M_{\tilde{D}}^2$ ,  $M_{\tilde{L}}^2$ , and  $M_{\tilde{E}}^2$  are hermitian  $3 \times 3$  matrices, and  $A_U$ ,  $A_D$ , and  $A_E$  are complex  $3 \times 3$  matrices. In addition,  $M_1$ ,  $M_2$ ,  $M_3$ ,  $B$ , and  $\mu$  are in general complex parameters. Finally, as in the Standard Model, the Higgs-fermion Yukawa couplings,  $\lambda_f$  ( $f = u, d$ , and  $e$ ), are complex  $3 \times 3$  matrices that are related to the quark and lepton mass matrices via:  $M_f = \lambda_f v_f$ , where  $v_e \equiv v_d$  [with  $v_u$  and  $v_d$  as defined above Eq. (2)].

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. 75 shows that the MSSM possesses 124 independent parameters. Of these, 18 correspond to Standard Model parameters (including the QCD vacuum angle  $\theta_{\text{QCD}}$ ), 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 [76].

**I.4. The supersymmetric-particle spectrum:** The supersymmetric particles (*sparticles*) differ in spin by half a unit from their Standard Model partners. The supersymmetric partners of the gauge and Higgs bosons are fermions, whose names are obtained by appending “ino” to the end of the corresponding Standard Model particle name. The gluino is the color-octet Majorana fermion partner of the gluon with mass  $M_{\tilde{g}} = |M_3|$ . The supersymmetric partners of the electroweak gauge and Higgs bosons (the gauginos and higgsinos) can mix due to  $SU(2) \times U(1)$  breaking effects. As a result, the physical states of definite mass are model-dependent linear combinations of the charged and neutral gauginos and higgsinos, called *charginos* and *neutralinos*, respectively. Like the gluino, the neutralinos are also Majorana fermions, which provide for some distinctive phenomenological signatures [77,78]. The supersymmetric partners of the quarks and leptons are spin-zero bosons: the *squarks*, charged *sleptons*, and *sneutrinos*, respectively. A complete set of Feynman rules for the sparticles of the MSSM can be found in Ref. 79. The MSSM Feynman rules also are implicitly contained in a number of Feynman diagram and amplitude generation software packages (see *e.g.*, Refs. [80–82]).

It should be noted that all mass formulae quoted in this section are tree-level results. Radiative loop corrections will modify these results and must be included in any precision study of supersymmetric phenomenology [83]. Beyond tree level, the definition of the supersymmetric parameters becomes convention-dependent. For example, one can define physical couplings or running couplings, which differ beyond the tree level. This provides a challenge to any effort that attempts to extract supersymmetric parameters from data. The Supersymmetry Les Houches Accord (SLHA) [84] has been adopted, which establishes a set of conventions for specifying generic file structures for supersymmetric model specifications and input parameters, supersymmetric mass and coupling spectra, and

decay tables. These provide a universal interface between spectrum calculation programs, decay packages, and high energy physics event generators. Ultimately, these efforts will facilitate the reconstruction of the fundamental supersymmetric theory (and its breaking mechanism) from high-precision studies of supersymmetric phenomena at future colliders.

**I.4.1. The charginos and neutralinos:** The mixing of the charged gauginos ( $\widetilde{W}^\pm$ ) and charged higgsinos ( $H_u^+$  and  $H_d^-$ ) is described (at tree-level) by a  $2 \times 2$  complex mass matrix [85–87]:

$$M_C \equiv \begin{pmatrix} M_2 & \frac{1}{\sqrt{2}}g v_u \\ \frac{1}{\sqrt{2}}g v_d & \mu \end{pmatrix}. \quad (5)$$

To determine the physical chargino states and their masses, one must perform a singular value decomposition [88,89] of the complex matrix  $M_C$ :

$$U^* M_C V^{-1} = \text{diag}(M_{\widetilde{\chi}_1^+}, M_{\widetilde{\chi}_2^+}), \quad (6)$$

where  $U$  and  $V$  are unitary matrices, and the right-hand side of Eq. (6) is the diagonal matrix of (non-negative) chargino masses. The physical chargino states are denoted by  $\widetilde{\chi}_1^\pm$  and  $\widetilde{\chi}_2^\pm$ . These are linear combinations of the charged gaugino and higgsino states determined by the matrix elements of  $U$  and  $V$  [85–87]. The chargino masses correspond to the *singular values* [88] of  $M_C$ , *i.e.*, the positive square roots of the eigenvalues of  $M_C^\dagger M_C$ :

$$M_{\widetilde{\chi}_1^+, \widetilde{\chi}_2^+}^2 = \frac{1}{2} \left\{ |\mu|^2 + |M_2|^2 + 2m_W^2 \mp \left[ (|\mu|^2 + |M_2|^2 + 2m_W^2)^2 - 4|\mu|^2|M_2|^2 - 4m_W^4 \sin^2 2\beta + 8m_W^2 \sin 2\beta \text{Re}(\mu M_2) \right]^{1/2} \right\}, \quad (7)$$

where the states are ordered such that  $M_{\widetilde{\chi}_1^+} \leq M_{\widetilde{\chi}_2^+}$ . Note that the relative phase of  $M_2$  and  $\mu$  is meaningful.

The mixing of the neutral gauginos ( $\widetilde{B}$  and  $\widetilde{W}^0$ ) and neutral higgsinos ( $\widetilde{H}_d^0$  and  $\widetilde{H}_u^0$ ) is described (at tree-level) by a  $4 \times 4$  complex symmetric mass matrix [85,86,90,91]:

$$M_N \equiv \begin{pmatrix} M_1 & 0 & -\frac{1}{2}g' v_d & \frac{1}{2}g' v_u \\ 0 & M_2 & \frac{1}{2}g v_d & -\frac{1}{2}g v_u \\ -\frac{1}{2}g' v_d & \frac{1}{2}g v_d & 0 & -\mu \\ \frac{1}{2}g' v_u & -\frac{1}{2}g v_u & -\mu & 0 \end{pmatrix}. \quad (8)$$



To determine the physical neutralino states and their masses, one must perform a Takagi-diagonalization [88,89,92,93] of the complex symmetric matrix  $M_N$ :

$$W^T M_N W = \text{diag}(M_{\tilde{\chi}_1^0}, M_{\tilde{\chi}_2^0}, M_{\tilde{\chi}_3^0}, M_{\tilde{\chi}_4^0}), \quad (9)$$

where  $W$  is a unitary matrix and the right-hand side of Eq. (9) is the diagonal matrix of (non-negative) neutralino masses. The physical neutralino states are denoted by  $\tilde{\chi}_i^0$  ( $i = 1, \dots, 4$ ), where the states are ordered such that  $M_{\tilde{\chi}_1^0} \leq M_{\tilde{\chi}_2^0} \leq M_{\tilde{\chi}_3^0} \leq M_{\tilde{\chi}_4^0}$ . The  $\tilde{\chi}_i^0$  are the linear combinations of the neutral gaugino and higgsino states determined by the matrix elements of  $W$  (which is denoted by  $N^{-1}$  in Ref. 85). The neutralino masses correspond to the singular values of  $M_N$ , *i.e.*, the positive square roots of the eigenvalues of  $M_N^\dagger M_N$ . Exact formulae for these masses can be found in Refs. [90] and [94]. A numerical algorithm for determining the mixing matrix  $W$  has been given by Ref. 95.

If a chargino or neutralino state 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  $\tilde{\chi}_1^0$  would be nearly a pure *photino*,  $\tilde{\gamma}$ , 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 *bin*o,  $\tilde{B}$ , 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*,  $\tilde{W}^\pm$ , and  $\tilde{W}_3^0$ , 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 chargino pair and neutralino would be nearly pure *higgsino* states, the supersymmetric partners of the Higgs bosons. Each of the above cases leads to a strikingly different phenomenology.

In the NMSSM, an additional Higgs singlet superfield is added to the MSSM. This superfield comprises two real Higgs scalar degrees of freedom and an associated neutral higgsino degree of freedom. Consequently, there are five neutralino mass

eigenstates that are obtained by a Takagi-diagonalization of the  $5 \times 5$  neutralino mass matrix. In many cases, the fifth neutralino state is dominated by its  $SU(2) \times U(1)$  singlet component, and thus is very weakly coupled to the Standard Model particles and their superpartners.

**I.4.2. The squarks, sleptons and sneutrinos:** For a given fermion  $f$ , there are two supersymmetric partners,  $\tilde{f}_L$  and  $\tilde{f}_R$ , which are scalar partners of the corresponding left-handed and right-handed fermion. (There is no  $\tilde{\nu}_R$  in the MSSM.) However, in general  $\tilde{f}_L$ – $\tilde{f}_R$  mixing is possible, in which case  $\tilde{f}_L$  and  $\tilde{f}_R$  are not mass eigenstates. For three generations of squarks, one must diagonalize  $6 \times 6$  matrices corresponding to the basis  $(\tilde{q}_{iL}, \tilde{q}_{iR})$ , where  $i = 1, 2, 3$  are the generation labels. For simplicity, only the one-generation case is illustrated in detail below. (The effects of second and third generation squark mixing can be significant and is treated in Ref. 96.)

Using the notation of the third family, the one-generation tree-level squark squared-mass matrix is given by [97]

$$\mathcal{M}^2 = \begin{pmatrix} M_{\tilde{Q}}^2 + m_q^2 + L_q & m_q X_q^* \\ m_q X_q & M_{\tilde{R}}^2 + m_q^2 + R_q \end{pmatrix}, \quad (10)$$

where

$$X_q \equiv A_q - \mu^* (\cot \beta)^{2T_{3q}}, \quad (11)$$

and  $T_{3q} = \frac{1}{2} [-\frac{1}{2}]$  for  $q = t$  [ $b$ ]. The diagonal squared-masses are governed by soft-supersymmetry-breaking squared-masses  $M_{\tilde{Q}}^2$  and  $M_{\tilde{R}}^2 \equiv M_{\tilde{U}}^2 [M_{\tilde{D}}^2]$  for  $q = t$  [ $b$ ], the corresponding quark masses  $m_t$  [ $m_b$ ], and electroweak correction terms:

$$L_q \equiv (T_{3q} - e_q \sin^2 \theta_W) m_Z^2 \cos 2\beta, \quad R_q \equiv e_q \sin^2 \theta_W m_Z^2 \cos 2\beta, \quad (12)$$

where  $e_q = \frac{2}{3} [-\frac{1}{3}]$  for  $q = t$  [ $b$ ]. The off-diagonal squark squared-masses are proportional to the corresponding quark masses and depend on  $\tan \beta$ , the soft-supersymmetry-breaking  $A$ -parameters and the higgsino mass parameter  $\mu$ . Assuming that the  $A$ -parameters are parametrically of the same order (or smaller) relative to other supersymmetry-breaking mass parameters, it then follows that  $\tilde{q}_L$ – $\tilde{q}_R$  mixing effects are small,

with the possible exception of the third generation, where mixing can be enhanced by factors of  $m_t$  and  $m_b \tan \beta$ .

In the case of third generation  $\tilde{q}_L$ - $\tilde{q}_R$  mixing, the mass eigenstates (usually denoted by  $\tilde{q}_1$  and  $\tilde{q}_2$ , with  $m_{\tilde{q}_1} < m_{\tilde{q}_2}$ ) are determined by diagonalizing the  $2 \times 2$  matrix  $\mathcal{M}^2$  given by Eq. (10). The corresponding squared-masses and mixing angle are given by [97]:

$$m_{\tilde{q}_{1,2}}^2 = \frac{1}{2} \left[ \text{Tr } \mathcal{M}^2 \mp \sqrt{(\text{Tr } \mathcal{M}^2)^2 - 4 \det \mathcal{M}^2} \right],$$

$$\sin 2\theta_{\tilde{q}} = \frac{2m_q |X_q|}{m_{\tilde{q}_2}^2 - m_{\tilde{q}_1}^2}. \quad (13)$$

The one-generation results above also apply to the charged sleptons, with the obvious substitutions:  $q \rightarrow \tau$  with  $T_{3\tau} = -\frac{1}{2}$  and  $e_\tau = -1$ , and the replacement of the supersymmetry-breaking parameters:  $M_{\tilde{Q}}^2 \rightarrow M_{\tilde{L}}^2$ ,  $M_{\tilde{D}}^2 \rightarrow M_{\tilde{E}}^2$ , and  $A_q \rightarrow A_\tau$ . For the neutral sleptons,  $\tilde{\nu}_R$  does not exist in the MSSM, so  $\tilde{\nu}_L$  is a mass eigenstate.

In the case of three generations, the supersymmetry-breaking scalar-squared masses [ $M_{\tilde{Q}}^2$ ,  $M_{\tilde{U}}^2$ ,  $M_{\tilde{D}}^2$ ,  $M_{\tilde{L}}^2$ , and  $M_{\tilde{E}}^2$ ] and the  $A$ -parameters [ $A_U$ ,  $A_D$ , and  $A_E$ ] are now  $3 \times 3$  matrices as noted in Section I.3.3. The diagonalization of the  $6 \times 6$  squark mass matrices yields  $\tilde{f}_{iL}$ - $\tilde{f}_{jR}$  mixing (for  $i \neq j$ ). In practice, since the  $\tilde{f}_L$ - $\tilde{f}_R$  mixing is appreciable only for the third generation, this additional complication can often be neglected (although see Ref. 96 for examples in which the mixing between the second and third generation squarks is relevant).

**I.5. The supersymmetric Higgs sector:** Consider first the MSSM Higgs sector [31,32,98]. Despite the large number of potential  $CP$ -violating phases among the MSSM-124 parameters, the tree-level MSSM Higgs sector is automatically  $CP$ -conserving. This follows from the fact that the only potentially complex parameter ( $m_{12}^2$ ) of the MSSM Higgs potential can be chosen real and positive by rephasing the Higgs fields, in which case  $\tan \beta$  is a real positive parameter. Consequently, the physical neutral Higgs scalars are  $CP$ -eigenstates. The MSSM Higgs sector contains five physical spin-zero particles: a charged Higgs

boson pair ( $H^\pm$ ), two  $CP$ -even neutral Higgs bosons (denoted by  $h^0$  and  $H^0$  where  $m_h < m_H$ ), and one  $CP$ -odd neutral Higgs boson ( $A^0$ ). The recent discovery of a Standard Model-like Higgs boson at the LHC with a mass of 126 GeV [66,67] strongly suggests that this state should be identified with  $h^0$ , although the possibility that the 126 GeV state should be identified with  $H^0$  cannot be completely ruled out [99].

In the NMSSM [35], the scalar component of the singlet Higgs superfield adds two additional neutral states to the Higgs sector. In this model, the tree-level Higgs sector can exhibit explicit  $CP$ -violation. If  $CP$  is conserved, then the two extra neutral scalar states are  $CP$ -even and  $CP$ -odd, respectively. These states can potentially mix with the neutral Higgs states of the MSSM. If scalar states exist that are dominantly singlet, then they are weakly coupled to Standard Model gauge bosons and fermions through their small mixing with the MSSM Higgs scalars. Consequently, it is possible that one (or both) of the singlet-dominated states is considerably lighter than the Higgs boson that was observed at the LHC.

***1.5.1 The Tree-level Higgs sector:*** The properties of the Higgs sector are determined by the Higgs potential, which is made up of quadratic terms [whose squared-mass coefficients were specified above Eq. (2)] and quartic interaction terms governed by dimensionless couplings. The quartic interaction terms are manifestly supersymmetric at tree level (although these are modified by supersymmetry-breaking effects 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. 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. (2)] 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  $\alpha$  (which measures the mixture of the original  $Y = \pm 1$  Higgs doublet states in the physical  $CP$ -even neutral scalars), and the Higgs boson self-couplings. Moreover, the tree-level mass of the lighter  $CP$ -even Higgs boson is bounded,  $m_h \leq m_Z |\cos 2\beta| \leq m_Z$  [31,32]. This bound can be substantially modified when radiative corrections are included, as discussed in Section I.5.2.

In the NMSSM, the superpotential contains a trilinear term that couples the two  $Y = \pm 1$  Higgs doublet superfields and the singlet Higgs superfield. The coefficient of this term is denoted by  $\lambda$ . Consequently, the tree-level bound for the mass of the lightest  $CP$ -even MSSM Higgs boson is modified [100],

$$m_h^2 \leq m_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta, \quad (14)$$

where  $v \equiv (v_u^2 + v_d^2)^{1/2} = 174$  GeV. If one demands that  $\lambda$  should stay finite below the Planck scale, then  $\lambda$  is constrained to lie below about 0.7. However, in light of the observed Higgs mass of 126 GeV, there is some phenomenological motivation for considering larger values of  $\lambda$  [101].

**I.5.2 The radiatively-corrected Higgs sector:** When radiative corrections are incorporated, additional parameters of the supersymmetric model enter via virtual supersymmetric particles that can appear in loops. The impact of these corrections can be significant [102]. The qualitative behavior of these 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 and the off-diagonal entries of the top-squark squared-mass matrix [Eq. (10)] are small in comparison to the geometric mean of the two top-squark squared-masses,  $M_S^2 \equiv M_{t_1} M_{t_2}$ . In this case (assuming  $m_A > m_Z$ ), the predicted upper bound for  $m_h$  is approximately given by

$$m_h^2 \lesssim m_Z^2 \cos^2 2\beta + \frac{3g^2 m_t^4}{8\pi^2 m_W^2} \left[ \ln \left( \frac{M_S^2}{m_t^2} \right) + \frac{X_t^2}{M_S^2} \left( 1 - \frac{X_t^2}{12M_S^2} \right) \right], \quad (15)$$

where  $X_t \equiv A_t - \mu \cot \beta$  is proportional to the off-diagonal entry of the top-squark squared-mass matrix (which for simplicity is

taken here to be real). The Higgs mass upper limit is saturated when  $\tan\beta$  is large (*i.e.*,  $\cos^2 2\beta \sim 1$ ) and  $X_t = \sqrt{6} M_S$ . This value of  $X_t$  defines the so-called *maximal mixing* scenario.

A more complete treatment of the radiative corrections [103] shows that Eq. (15) 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  $m_h \lesssim 135$  GeV in the large  $\tan\beta$  regime (with an accuracy of a few GeV) for  $m_t = 175$  GeV and  $M_S \lesssim 2$  TeV [103].

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 [104]. Although these effects are more model-dependent, they can have a non-trivial impact on the Higgs searches at LHC and at future colliders. A summary of the MSSM Higgs mass limits based on searches at LEP, Tevatron and LHC and the corresponding references to the literature can be found in Ref. 67.

In the NMSSM, the dominant radiative correction to Eq. (14) is the same as the one given in Eq. (15). However, in contrast to the MSSM, one does not need as large a boost from the radiative corrections to achieve a Higgs mass of 126 GeV. For example, the observed Higgs mass is easily achieved in the NMSSM parameter regime where  $\tan\beta \sim 2$  and  $\lambda \sim 0.7$ .

***I.6. Restricting 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  $L_e$ ,  $L_\mu$ , and  $L_\tau$ ; (ii) unsuppressed flavor-changing neutral currents (FCNCs); 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 [68]. In particular, in weak-scale supersymmetry

some combinations of the complex phases of the gaugino-mass parameters, the  $A$ -parameters, and  $\mu$  must be less than on the order of  $10^{-2}$ – $10^{-3}$  to avoid generating electric dipole moments for the neutron, electron, and atoms in conflict with observed data [105–107]. The non-observation of FCNCs [108–110] places additional strong constraints on the off-diagonal matrix elements of the squark and slepton soft-supersymmetry-breaking squared-masses and  $A$ -parameters (see Section I.3.3). 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” regions of the full parameter space.

The MSSM-124 is also theoretically incomplete as it provides no explanation for the origin of the supersymmetry-breaking parameters (and in particular, why these parameters should conform to the exceptional regions 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 a set of soft-supersymmetry-breaking terms that is consistent with all phenomenological constraints.

The successful unification of the  $SU(3) \times SU(2) \times U(1)$  gauge couplings in supersymmetric grand unified theories [8,65,111,112] suggests that the high-energy structure of the theory may be considerably simpler than its low-energy realization. In a top-down approach, the dynamics that governs the more fundamental theory at high energies is used to derive the effective broken-supersymmetric theory at the TeV scale. A suitable choice for the high energy dynamics is one that yields a TeV-scale theory that satisfies all relevant phenomenological constraints.

In this Section, we examine a number of theoretical frameworks that potentially yield phenomenologically viable regions of the MSSM-124 parameter space. The resulting supersymmetric particle spectrum is then a function of a relatively small number of input parameters. This is accomplished by imposing a simple structure on the soft-supersymmetry-breaking terms at a common high-energy scale  $M_X$  (typically chosen to be

the Planck scale,  $M_{\text{P}}$ , the grand unification scale,  $M_{\text{GUT}}$ , or the messenger scale,  $M_{\text{mess}}$ ). Using the renormalization group equations, one can then derive the low-energy MSSM parameters relevant for collider physics. 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, anomaly mediated and gauge mediated supersymmetry breaking, to be discussed in more detail below. In some of these approaches, one of the diagonal Higgs squared-mass parameters is driven negative by renormalization group evolution [113]. In such models, electroweak symmetry breaking is generated radiatively, and the resulting electroweak symmetry-breaking scale is intimately tied to the scale of low-energy supersymmetry breaking.

***1.6.1. Gaugino mass unification:***

One prediction that arises in many 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  $M_X$ :

$$M_1(M_X) = M_2(M_X) = M_3(M_X) = m_{1/2}. \quad (16)$$

Due to renormalization group running, the effective low-energy gaugino mass parameters (at the electroweak scale) are related:

$$M_3 = (g_s^2/g^2)M_2 \simeq 3.5M_2, \quad M_1 = (5g'^2/3g^2)M_2 \simeq 0.5M_2. \quad (17)$$

In this case, the chargino and neutralino masses and mixing angles depend only on three unknown parameters: the gluino mass,  $\mu$ , and  $\tan\beta$ . It then follows that the lightest neutralino must be heavier than 46 GeV due to the non-observation of charginos at LEP [114]. If in addition  $|\mu| \gg |M_1| \gtrsim m_Z$ , then the lightest neutralino is nearly a pure bino, an assumption often made in supersymmetric particle searches at colliders.

Although Eqs. (16) and (17) are often assumed in many phenomenological studies, a truly model-independent approach would take the gaugino mass parameters,  $M_i$ , to be independent



parameters to be determined by experiment. Indeed, an approximately massless neutralino *cannot* be ruled out at present by a model-independent analysis [115].

It is possible that the tree-level masses for the gauginos are absent. In this case, the gaugino mass parameters arise at one-loop and do not satisfy Eq. (17). For example, the gaugino masses in AMSB models arise entirely from a model-independent contribution derived from the super-conformal anomaly [49,116]. In this case, Eq. (17) is replaced (in the one-loop approximation) by:

$$M_i \simeq \frac{b_i g_i^2}{16\pi^2} m_{3/2}, \quad (18)$$

where  $m_{3/2}$  is the gravitino mass and 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,  $(b_1, b_2, b_3) = (\frac{33}{5}, 1, -3)$ . Eq. (18) yields  $M_1 \simeq 2.8M_2$  and  $M_3 \simeq -8.3M_2$ , which implies that the lightest chargino pair and neutralino comprise a nearly mass-degenerate triplet of winos,  $\widetilde{W}^\pm, \widetilde{W}^0$  (cf. Table 1), over most of the MSSM parameter space. For example, if  $|\mu| \gg m_Z$ , then Eq. (18) implies that  $M_{\widetilde{\chi}_1^\pm} \simeq M_{\widetilde{\chi}_1^0} \simeq M_2$  [117]. The corresponding supersymmetric phenomenology differs significantly from the standard phenomenology based on Eq. (17) [118,119].

Finally, it should be noted that the unification of gaugino masses (and scalar masses) can be accidental. In particular, the energy scale where unification takes place may not be directly related to any physical scale. One version of this phenomenon has been called *mirage unification* and can occur in certain theories of fundamental supersymmetry breaking [120].

### ***1.6.2. The constrained MSSM: mSUGRA, CMSSM, . . .***

In the *minimal* supergravity (mSUGRA) framework [3–5,42–44], a form of the Kähler potential is employed that yields minimal kinetic energy terms for the MSSM fields [46]. As a result, the soft-supersymmetry-breaking parameters at the high-energy scale  $M_X$  take a particularly simple form in which the scalar

squared-masses and the  $A$ -parameters are flavor-diagonal and universal [44]:

$$\begin{aligned}
 M_Q^2(M_X) &= M_U^2(M_X) = M_D^2(M_X) = m_0^2 \mathbf{1}, \\
 M_L^2(M_X) &= M_E^2(M_X) = m_0^2 \mathbf{1}, \\
 m_1^2(M_X) &= m_2^2(M_X) = m_0^2, \\
 A_U(M_X) &= A_D(M_X) = A_E(M_X) = A_0 \mathbf{1},
 \end{aligned}
 \tag{19}$$

where  $\mathbf{1}$  is a  $3 \times 3$  identity matrix in generation space. As in the Standard Model, this approach exhibits minimal flavor violation [121,122], whose unique source is the nontrivial flavor structure of the Higgs-fermion Yukawa couplings. The gaugino masses are also unified according to Eq. (16).

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 masses, one must use the *low-energy* values for  $M_Q^2$ ,  $M_U^2$ , and  $M_D^2$  in Eq. (10). Through the renormalization group running with boundary conditions specified in Eqs. (17) and (19), one can show that the low-energy values of  $M_Q^2$ ,  $M_U^2$ , and  $M_D^2$  depend primarily on  $m_0^2$  and  $m_{1/2}^2$ . A number of useful approximate analytic expressions for superpartner masses in terms of the mSUGRA parameters can be found in Ref. 123.

In the mSUGRA approach, four flavors of squarks (with two squark eigenstates per flavor) are nearly mass-degenerate. If  $\tan \beta$  is not very large,  $\tilde{b}_R$  is also approximately degenerate in mass with the first two generations of squarks. The  $\tilde{b}_L$  mass and the diagonal  $\tilde{t}_L$  and  $\tilde{t}_R$  masses are typically reduced relative to the common squark mass of the first two generations. 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. As noted below Eq. (10), third-generation squark masses and tau-slepton masses are sensitive to the strength of the respective  $\tilde{f}_L$ - $\tilde{f}_R$  mixing. The LSP is typically the lightest neutralino,  $\tilde{\chi}_1^0$ , which is dominated by its bino component. Regions of the

mSUGRA parameter space in which the LSP is electrically charged do exist but are not phenomenologically viable [19].

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  $m_0$ ,  $m_{1/2}$ ,  $A_0$ , the Planck-scale values for  $\mu$  and  $B$ -parameters (denoted by  $\mu_0$  and  $B_0$ ), and the gravitino mass  $m_{3/2}$ . Without additional model assumptions,  $m_{3/2}$  is independent of the parameters that govern the mass spectrum of the superpartners of the Standard Model [44]. In principle,  $A_0$ ,  $B_0$ ,  $\mu_0$ , and  $m_{3/2}$  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  $\mu_0$  and  $B_0$  in favor of  $m_Z$  and  $\tan\beta$  [the sign of  $\mu_0$ , denoted  $\text{sgn}(\mu_0)$  below, is not fixed in this process]. In this case, the MSSM spectrum and its interaction strengths are determined by five parameters:

$$m_0, A_0, m_{1/2}, \tan\beta, \text{ and } \text{sgn}(\mu_0), \quad (20)$$

and an independent gravitino mass  $m_{3/2}$  (in addition to the 18 parameters of the Standard Model). This framework is conventionally called the *constrained minimal supersymmetric extension of the Standard Model* (CMSSM).

In the early literature, additional conditions were obtained by assuming a simplified form for the hidden sector that provides the fundamental source of supersymmetry breaking. Two additional relations emerged among the mSUGRA parameters [42,46]:  $B_0 = A_0 - m_0$  and  $m_{3/2} = m_0$ . These relations characterize a theory that was called minimal supergravity when first proposed. In the more recent literature, it has been more common to omit these extra conditions in defining the mSUGRA model (in which case the mSUGRA model and the CMSSM are synonymous). The authors of Ref. 124 advocate

restoring the original nomenclature in which the mSUGRA model is defined with the extra conditions as originally proposed. Additional mSUGRA variations can be considered where different relations among the CMSSM parameters are imposed.

One can also relax the universality of scalar masses by decoupling the squared-masses of the Higgs bosons and the squarks/sleptons. This leads to the non-universal Higgs mass models (NUHMs), thereby adding one or two new parameters to the CMSSM depending on whether the diagonal Higgs scalar squared-mass parameters ( $m_1^2$  and  $m_2^2$ ) are set equal (NUHM1) or taken to be independent (NUHM2) at the high energy scale  $M_X^2$ . Clearly, this modification preserves the minimal flavor violation of the mSUGRA approach. Nevertheless, the mSUGRA approach and its NUHM generalizations are probably too simplistic. Theoretical considerations suggest that the universality of Planck-scale soft-supersymmetry-breaking parameters is not generic [125]. In particular, effective operators at the Planck scale exist that do not respect flavor universality, and it is difficult to find a theoretical principle that would forbid them.

In the framework of supergravity, if anomaly mediation is the sole source of supersymmetry breaking, then the gaugino mass parameters, diagonal scalar squared-mass parameters, and the supersymmetry-breaking trilinear scalar interaction terms (proportional to  $\lambda_f A_F$ ) are determined in terms of the beta functions of the gauge and Yukawa couplings and the anomalous dimensions of the squark and slepton fields [49,116,119]. As noted in Section I.2.3, this approach yields tachyonic sleptons in the MSSM unless additional sources of supersymmetry breaking are present. In the *minimal* AMSB (mAMSB) scenario, a universal squared-mass parameter,  $m_0^2$ , is added to the AMSB expressions for the diagonal scalar squared-masses [119]. Thus, the mAMSB spectrum and its interaction strengths are determined by four parameters,  $m_0^2$ ,  $m_{3/2}$ ,  $\tan \beta$  and  $\text{sgn}(\mu_0)$ .

The mAMSB scenario appears to be ruled out based on the observed value of the Higgs boson mass, assuming an upper limit on  $M_G$  of a few TeV, since the mAMSB constraint on  $A_F$  implies that the maximal mixing scenario cannot be achieved [cf. Eq. (15)]. Indeed, under the stated assumptions,

the mAMSB Higgs mass upper bound lies below the observed Higgs mass value [126]. Thus within the AMSB scenario, either an additional supersymmetry-breaking contribution to  $\lambda_f A_F$  and/or new ingredients beyond the MSSM are required.

**I.6.3. Gauge mediated supersymmetry breaking:** In contrast to models of gravity mediated supersymmetry breaking, the universality of the fundamental soft-supersymmetry-breaking squark and slepton squared-mass parameters is guaranteed in gauge mediated supersymmetry breaking (GMSB) because the supersymmetry breaking is communicated to the sector of MSSM fields via gauge interactions [53,54]. In GMSB models, the mass scale of the messenger sector (or its equivalent) is sufficiently below the Planck scale such that the additional supersymmetry-breaking effects mediated by supergravity can be neglected.

In the minimal GMSB approach, there is one effective mass scale,  $\Lambda$ , 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  $\mu$  and  $B$ -parameters is quite model-dependent, and lies somewhat outside the ansatz of gauge mediated supersymmetry breaking.

The simplest GMSB models appear to be ruled out based on the observed value of the Higgs boson mass. Due to suppressed  $A$  parameters, it is difficult to boost the contributions of the radiative corrections in Eq. (15) to obtain a Higgs mass as large as 126 GeV. However, this conflict can be alleviated in more complicated GMSB models [127]. To analyze these generalized GMSB models, it has been especially fruitful to develop model-independent techniques that encompass all known GMSB models [128]. These techniques are well-suited for a comprehensive analysis [129] of the phenomenological profile of gauge mediated supersymmetry breaking.

The gravitino is the LSP in GMSB models, as noted in Section I.2.3. As a result, the next-to-lightest supersymmetric particle (NLSP) now plays a crucial role in the phenomenology of supersymmetric particle production and decays. Note that

unlike the LSP, the NLSP can be charged. In GMSB models, the most likely candidates for the NLSP are  $\tilde{\chi}_1^0$  and  $\tilde{\tau}_R^\pm$ . The NLSP will decay into its superpartner plus a gravitino (*e.g.*,  $\tilde{\chi}_1^0 \rightarrow \gamma\tilde{G}$ ,  $\tilde{\chi}_1^0 \rightarrow Z\tilde{G}$ ,  $\tilde{\chi}_1^0 \rightarrow h^0\tilde{G}$  or  $\tilde{\tau}_R^\pm \rightarrow \tau^\pm\tilde{G}$ ), with lifetimes and branching ratios that depend on the model parameters. 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 a supersymmetric decay chain [130]. For example, in the slepton co-NLSP case, all three right-handed sleptons are close enough in mass and thus can each play the role of the NLSP.

Different choices for the identity of the NLSP and its decay rate lead to a variety of distinctive supersymmetric phenomenologies [54,131]. For example, a long-lived  $\tilde{\chi}_1^0$ -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 standard phenomenology of the  $\tilde{\chi}_1^0$ -LSP). On the other hand, if  $\tilde{\chi}_1^0 \rightarrow \gamma\tilde{G}$  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, in the case of a  $\tilde{\tau}_R^\pm$ -NLSP, the  $\tilde{\tau}_R^\pm$  would either be long-lived or would decay inside the detector into a  $\tau$ -lepton plus missing energy.

***1.6.4. The phenomenological MSSM:*** Of course, any of the theoretical assumptions described in this Section could be wrong and must eventually be tested experimentally. To facilitate the exploration of MSSM phenomena in a more model-independent way while respecting the constraints noted at the beginning of this Section, the phenomenological MSSM (pMSSM) has been introduced [132].

The pMSSM is governed by 19 independent real parameters beyond the Standard Model, which include the three gaugino mass parameters  $M_1$ ,  $M_2$  and  $M_3$ , the Higgs sector parameters  $m_A$  and  $\tan\beta$ , the Higgsino mass parameter  $\mu$ , five squark and slepton squared-mass parameters for the degenerate first and second generations ( $M_Q^2$ ,  $M_U^2$ ,  $M_D^2$ ,  $M_L^2$  and  $M_E^2$ ), the five corresponding squark and slepton squared-mass parameters for the third generation, and three third-generation  $A$ -parameters ( $A_t$ ,

$A_b$  and  $A_\tau$ ). As previously noted, the first and second generation  $A$ -parameters can be neglected as their phenomenological consequences are negligible.

***1.7. Experimental data confronts the MSSM:***

At present, there is no evidence for weak-scale supersymmetry from the data analyzed by the LHC experiments. Recent LHC data has been especially effective in ruling out the existence of colored supersymmetric particles (primarily the gluino and the first generation of squarks) with masses below about 1 TeV [13,133]. The precise mass limits are model dependent. For example, higher mass colored superpartners have been ruled out in the context of the CMSSM. In more generic frameworks of the MSSM, regions of parameter space can be identified in which lighter squarks and gluinos below 1 TeV cannot be definitely ruled out. Details of these constraints can be found in Ref. 13. Additional constraints arise from limits on the contributions of virtual supersymmetric particle exchange to a variety of Standard Model processes [108–110].

In light of the negative results in the search for supersymmetry, one must confront the tension that exists between the theoretical expectations for the magnitude of the supersymmetry-breaking parameters and the non-observation of supersymmetric phenomena.

***1.7.1 Naturalness constraints and the little hierarchy:***

In Section I, weak-scale supersymmetry was motivated as a natural solution to the hierarchy problem, which could provide an understanding of the origin of the electroweak symmetry-breaking scale without a significant fine-tuning of the fundamental parameters that govern the MSSM. In this context, the soft-supersymmetry-breaking masses must be generally of the order of 1 TeV or below [134]. This requirement is most easily seen in the determination of  $m_Z$  by the scalar potential minimum condition. In light of Eq. (4), to avoid the fine-tuning of MSSM parameters, the soft-supersymmetry-breaking squared-masses  $m_1^2$  and  $m_2^2$  and the higgsino squared-mass  $|\mu|^2$  should all be roughly of  $\mathcal{O}(m_Z^2)$ . Many authors have proposed quantitative measures of fine-tuning [134–136]. One of the simplest

measures is the one given by Barbieri and Giudice [134],

$$\Delta_i \equiv \left| \frac{\partial \ln m_Z^2}{\partial \ln p_i} \right|, \quad \Delta \equiv \max \Delta_i, \quad (21)$$

where the  $p_i$  are the MSSM parameters at the high-energy scale  $M_X$ , which are set by the fundamental supersymmetry-breaking dynamics. The theory is more fine-tuned as  $\Delta$  becomes larger.

One can apply the fine-tuning measure to any explicit model of supersymmetry breaking. For example, in the approaches discussed in Section I.6, the  $p_i$  are parameters of the model at the energy scale  $M_X$  where the soft-supersymmetry-breaking operators are generated by the dynamics of supersymmetry breaking. Renormalization group evolution then determines the values of the parameters appearing in Eq. (4) at the electroweak scale. In this way,  $\Delta$  is sensitive to all the supersymmetry-breaking parameters of the model (see e.g. Ref. 137).

As anticipated, there is a tension between the present experimental lower limits on the masses of colored supersymmetric particles [138] and the expectation that supersymmetry-breaking is associated with the electroweak symmetry-breaking scale. Moreover, this tension is exacerbated by the observed value of the Higgs mass ( $m_h \simeq 126$  GeV), which is not far from the the MSSM upper bound ( $m_h \lesssim 135$  GeV) [which depends on the top-squark mass and mixing as noted in Section I.5.2]. If  $M_{\text{SUSY}}$  characterizes the scale of supersymmetric particle masses, then one would crudely expect  $\Delta \sim M_{\text{SUSY}}^2/m_Z^2$ . For example, if  $M_{\text{SUSY}} \sim 1$  TeV then there must be at least a  $\Delta^{-1} \sim 1\%$  fine-tuning of the MSSM parameters to achieve the observed value of  $m_Z$ . This separation of the electroweak symmetry-breaking and supersymmetry-breaking scales is an example of the *little hierarchy problem* [140,141].

However, one must be very cautious when drawing conclusions about the viability of weak-scale supersymmetry to explain the origin of electroweak symmetry breaking [142]. First, one must decide the largest tolerable value of  $\Delta$  within the framework of weak-scale supersymmetry (should it be  $\Delta \sim 10$ ? 100? 1000?). Second, the fine-tuning parameter  $\Delta$  depends quite sensitively on the structure of the supersymmetry-breaking



dynamics, such as the value of  $M_X$  and relations among supersymmetry-breaking parameters in the fundamental high energy theory [143]. For example, in so-called focus point supersymmetry models [144], all squark masses can be as heavy as 5 TeV *without* significant fine-tuning. This can be attributed to a focusing behavior of the renormalization group evolution where certain relations hold among the high-energy values of the scalar squared-mass supersymmetry-breaking parameters.

Among the colored superpartners, the third generation squarks generically have the most significant impact on the naturalness constraints [145], while their masses are the least constrained by the LHC data. Hence, in the absence of any relation between third generation squarks and those of the first two generations, the naturalness constraints due to present LHC data can be considerably weaker than those obtained in the CMSSM. Indeed, models with first and second generation squark masses in the multi-TeV range do not generically require significant fine tuning. Such models have the added benefit that undesirable FCNCs mediated by squark exchange are naturally suppressed [146]. Other MSSM mass spectra that are compatible with moderate fine tuning have been considered in Refs. [133,143,147].

The lower bounds on squark and gluino masses may not be as large as suggested by the experimental analyses based on the CMSSM or simplified models [148]. For example, mass bounds for the gluino and the first and second generation squarks based on the CMSSM can often be evaded in alternative or extended MSSM models, *e.g.*, compressed supersymmetry [149] and stealth supersymmetry [150]. Moreover, experimental limits on the masses for the third generation squarks (which enter the fine-tuning considerations more directly) and color-neutral supersymmetric particles are less constrained than the masses of other colored supersymmetric states.

Finally, one can also consider extensions of the MSSM in which the degree of fine-tuning is relaxed. For example, it has already been noted in Section I.5.2 that it is possible to accommodate the observed Higgs mass more easily in the NMSSM due to contributions to  $m_h^2$  proportional to the parameter  $\lambda$ .

This means that we do not have to rely on a contribution to the Higgs mass from the radiative corrections to boost the Higgs mass sufficiently above its tree-level bound. In turn, this allows for smaller top squark masses, which are more consistent with the demands of naturalness. The reduction of the fine-tuning in various NMSSM models was initially advocated in Ref. 151, and more recently has been exhibited in Refs. [101,152]. Naturalness can also be relaxed in extended supersymmetric models with vector-like quarks [153] and in gauge extensions of the MSSM [154].

Thus, it is premature to conclude, after the first few years of LHC operation at just above half the design energy and with less than a quarter of the design luminosity, that weak-scale supersymmetry is on the verge of exclusion.

### *1.7.2 Constraints from virtual exchange of supersymmetric particles*

There are a number of low-energy measurements that are sensitive to the effects of new physics through indirect searches via supersymmetric loop effects. For example, the virtual exchange of supersymmetric particles can contribute to the muon anomalous magnetic moment,  $a_\mu \equiv \frac{1}{2}(g - 2)_\mu$ , as reviewed in Ref. 155. The Standard Model prediction for  $a_\mu$  exhibits a  $3.6\sigma$  deviation from the experimentally observed value [156]. This discrepancy is difficult to accommodate in the constrained supersymmetry models of Section I.6.2 and I.6.3 given the present sparticle mass bounds. Nevertheless, there are regions of the more general pMSSM parameter space that are consistent with the observed value of  $a_\mu$  [157].

The rare inclusive decay  $b \rightarrow s\gamma$  also provides a sensitive probe to the virtual effects of new physics beyond the Standard Model. Recent experimental measurements of  $B \rightarrow X_s + \gamma$  [159] are in very good agreement with the theoretical Standard Model predictions of Ref. 160. Since supersymmetric loop corrections can contribute an observable shift from the Standard Model predictions, the absence of any significant deviations places useful constraints on the MSSM parameter space [161].

The rare decay  $B_s \rightarrow \mu^+ \mu^-$  is especially sensitive to supersymmetric loop effects, with some loop contributions scaling as

$\tan^6 \beta$  when  $\tan \beta \gg 1$  [162]. The recent observation of this rare decay mode by the LHCb Collab. [163] is consistent with the predicted Standard Model rate. The absence of a *significant* deviation in these and other  $B$ -physics observables from their Standard Model predictions also places useful constraints on the MSSM parameter space [110,139,164].

Finally, we note that the constraints from precision electroweak observables [165] are easily accommodated in models of weak-scale supersymmetry [166]. Thus, robust regions of the MSSM parameter space, compatible with the results of direct and indirect searches for supersymmetry, remain unconstrained.

***1.8. Massive neutrinos in low-energy supersymmetry:***

In the minimal Standard Model and its supersymmetric extension, there are no right-handed neutrinos, and Majorana mass terms for the left-handed neutrinos are absent. However, given the overwhelming evidence for neutrino masses and mixing [167,168], any viable model of fundamental particles must provide a mechanism for generating neutrino masses [169]. In extended supersymmetric models, various mechanisms exist for producing massive neutrinos [170]. Although one can devise models for generating massive Dirac neutrinos [171], the most common approaches for incorporating neutrino masses are based on  $L$ -violating supersymmetric extensions of the MSSM, which generate massive Majorana neutrinos. Two classes of  $L$ -violating supersymmetric models will now be considered.

***1.8.1. The supersymmetric seesaw:*** Neutrino masses can be incorporated into the Standard Model by introducing  $SU(3) \times SU(2) \times U(1)$  singlet right-handed neutrinos ( $\nu_R$ ) and super-heavy Majorana masses (typically near the grand unification mass scale) for the  $\nu_R$ . In addition, one must also include a standard Yukawa couplings between the lepton doublets, the Higgs doublet, and the  $\nu_R$ . The Higgs vacuum expectation value then induces an off-diagonal  $\nu_L$ - $\nu_R$  masses on the order of the electroweak scale. Diagonalizing the neutrino mass matrix (in the three-generation model) yields three superheavy neutrino states, and three very light neutrino states that are identified with the light neutrinos observed in nature. This is the seesaw mechanism [172].

It is straightforward to construct a supersymmetric generalization of the seesaw model of neutrino masses [173,174]. In the seesaw-extended Standard Model, lepton number is broken due to the presence of  $\Delta L = 2$  terms in the Lagrangian (which include the Majorana mass terms for the superheavy neutrinos). Consequently, the seesaw-extended MSSM conserves R-parity. The supersymmetric analogue of the Majorana neutrino mass term in the sneutrino sector leads to sneutrino–antisneutrino mixing phenomena [174,175].

***1.8.2. R-parity-violating supersymmetry:*** In order to incorporate massive neutrinos in supersymmetric models while retaining the minimal particle content of the MSSM, one must relax the assumption of R-parity invariance. The most general R-parity-violating (RPV) model involving the MSSM spectrum introduces many new parameters to both the supersymmetry-conserving and the supersymmetry-breaking sectors [176]. Each new interaction term violates either  $B$  or  $L$  conservation. For example, starting from the MSSM superpotential previously given in Eq. (1) [suitably generalized to three generations of quarks, leptons and their superpartners], consider the effect of adding the following new terms:

$$(\lambda_L)_{pmn} \widehat{L}_p \widehat{L}_m \widehat{E}_n^c + (\lambda'_L)_{pmn} \widehat{L}_p \widehat{Q}_m \widehat{D}_n^c + (\lambda_B)_{pmn} \widehat{U}_p^c \widehat{D}_m^c \widehat{D}_n^c, \quad (22)$$

where  $p$ ,  $m$ , and  $n$  are generation indices, and gauge group indices are suppressed. Eq. (22) yields new scalar-fermion Yukawa couplings consisting of all possible combinations involving two Standard Model fermions and one scalar superpartner.

Note that the term in Eq. (22) proportional to  $\lambda_B$  violates  $B$ , while the other two terms violate  $L$ . Even if all the terms of Eq. (22) are absent, there is one more possible supersymmetric source of R-parity violation. Namely, one can add a term to the superpotential,  $(\mu_L)_p \widehat{H}_u \widehat{L}_p$ , where  $\widehat{H}_u$  is the  $Y = 1$  Higgs doublet superfield given in Table 1. This term is the RPV generalization of the higgsino mass parameter  $\mu$  of the MSSM, in which the  $Y = -1$  Higgs/higgsino supermultiplet  $\widehat{H}_d$  is replaced by the slepton/lepton supermultiplet  $\widehat{L}_p$ . The RPV-parameters  $(\mu_L)_p$  also violate  $L$ .

Phenomenological constraints derived from data on various low-energy  $B$ - and  $L$ -violating processes can be used to establish limits on each of the coefficients  $(\lambda_L)_{pmn}$ ,  $(\lambda'_L)_{pmn}$ , and  $(\lambda_B)_{pmn}$  taken one at a time [176,177]. If more than one coefficient is simultaneously non-zero, then the limits are in general more complicated [178]. 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 (and the stability of the proton), while allowing for lepton-number-violating interactions by imposing a discrete  $\mathbf{Z}_3$  baryon *triality* symmetry on the low-energy theory [179], in place of the standard  $\mathbf{Z}_2$  R-parity. Since the distinction between the Higgs and matter supermultiplets is lost in RPV models where  $L$  is violated, the mixing of sleptons and Higgs bosons, the mixing of neutrinos and neutralinos, and the mixing of charged leptons and charginos are now possible, leading to more complicated mass matrices and mass eigenstates than in the MSSM. Recent attempts to fit neutrino masses and mixing in this framework can be found in Ref. 180.

Alternatively, one can consider imposing a lepton parity such that all lepton superfields are odd [181,182]. In this case, only the  $B$ -violating term in Eq. (22) survives, and  $L$  is conserved. Models of this type have been considered in Ref. 183. Since  $L$  is conserved, the mixing of the lepton and Higgs superfields is forbidden. However, one expects that lepton parity cannot be exact due to quantum gravity effects. Remarkably, the standard  $\mathbf{Z}_2$  R-parity and the  $\mathbf{Z}_3$  baryon triality are stable with respect to quantum gravity effects, as they can be identified as residual discrete symmetries that arise from broken non-anomalous gauge symmetries [181].

The supersymmetric phenomenology of the RPV models exhibits features that are distinct from that of the MSSM [176].

The LSP is no longer stable, which implies that not all supersymmetric decay chains must yield missing-energy events at colliders. Indeed, the sparticle mass bounds obtained in searches for R-parity-conserving supersymmetry can be considerably relaxed in certain RPV models due to the absence of large missing transverse energy signatures [184]. This can alleviate some of the tension with naturalness discussed in Section I.7.1.

Nevertheless, the loss of the missing-energy signature is often compensated by other striking signals (which depend on which R-parity-violating parameters are dominant). For example, supersymmetric particles in RPV models can be singly produced (in contrast to R-parity-conserving models where supersymmetric particles must be produced in pairs). The phenomenology of pair-produced supersymmetric particles is also modified in RPV models due to new decay chains not present in R-parity-conserving supersymmetry models [176].

In RPV models with lepton number violation (these include low-energy supersymmetry models with baryon triality mentioned above), both  $\Delta L = 1$  and  $\Delta L = 2$  phenomena are allowed, leading to neutrino masses and mixing [185], neutrinoless double-beta decay [186], sneutrino-antisneutrino mixing [187], and resonant  $s$ -channel production of sneutrinos in  $e^+e^-$  collisions [188] and charged sleptons in  $p\bar{p}$  and  $pp$  collisions [189].

***I.9. Extensions beyond the MSSM:*** Extensions of the MSSM have been proposed to solve a variety of theoretical problems. One such problem involves the  $\mu$  parameter of the MSSM. Although  $\mu$  is a supersymmetric-*preserving* parameter, it must be of order the effective supersymmetry-breaking scale of the MSSM to yield a consistent supersymmetric phenomenology [190]. Any natural solution to the so-called  $\mu$ -problem must incorporate a symmetry that enforces  $\mu = 0$  and a small symmetry-breaking parameter that generates a value of  $\mu$  that is not parametrically larger than the effective supersymmetry-breaking scale [191]. A number of proposed mechanisms in the literature (*e.g.*, see Refs. [190–193]) provide concrete examples of a natural solution to the  $\mu$ -problem of the MSSM.

In extensions of the MSSM, new compelling solutions to the  $\mu$ -problem are possible. For example, one can replace  $\mu$  by the

vacuum expectation value of a new  $SU(3) \times SU(2) \times U(1)$  singlet scalar field. This is the NMSSM, which yields phenomena that were briefly discussed in Sections I.4, I.5 and I.7. The NMSSM superpotential consists only of trilinear terms whose coefficients are dimensionless. There are some advantages to extending the NMSSM further to the USSM [93] by adding a new broken  $U(1)$  gauge symmetry [194], under which the singlet field is charged.

Alternatively, one can consider a generalized version of the NMSSM (called the GNMSSM in Ref. 152), where all possible renormalizable terms in the superpotential are allowed, which yields new supersymmetric mass terms (analogous to the  $\mu$  term of the MSSM). Although the GNMSSM does not solve the  $\mu$ -problem, it does exhibit regions of parameter space in which the degree of fine-tuning is relaxed, as discussed in Section I.7.1.

It is also possible to add higher dimensional Higgs multiplets, such as Higgs triplet superfields [195], provided a custodial-symmetric model (in which the  $\rho$ -parameter of precision electroweak physics is close to 1 [165]) can be formulated. Such models can provide a rich phenomenology of new signals for future LHC studies.

All supersymmetric models discussed in this review possess self-conjugate fermions—the Majorana gluinos and neutralinos. However, it is possible to add additional chiral superfields in the adjoint representation. The spin-1/2 components of these new superfields can pair up with the gauginos to form Dirac gauginos [196,197]. Such states appear in models of so-called supersoft supersymmetry breaking [198], in some generalized GMSB models [199] and in R-symmetric supersymmetry [27,200]. Such approaches often lead to improved naturalness and/or significantly relaxed flavor constraints. The implications of models of Dirac gauginos on the observed Higgs boson mass and its properties is addressed in Ref. 201.

For completeness, we briefly note other MSSM extensions considered in the literature. These include an enlarged electroweak gauge group beyond  $SU(2) \times U(1)$  [202]; and/or the addition of new (possibly exotic) matter supermultiplets such as vector-like fermions and their superpartners [153,203].

## References

1. The early history of supersymmetry and a guide to the original literature can be found in *The Supersymmetric World—The Beginnings of the Theory*, edited by G. Kane and M. Shifman (World Scientific, Singapore, 2000).
2. R. Haag, J.T. Lopuszanski, and M. Sohnius, Nucl. Phys. **B88**, 257 (1975);  
S.R. Coleman and J. Mandula, Phys. Rev. **159** (1967) 1251.
3. H.P. Nilles, Phys. Reports **110**, 1 (1984).
4. P. Nath, R. Arnowitt, and A.H. Chamseddine, *Applied  $N = 1$  Supergravity* (World Scientific, Singapore, 1984).
5. S.P. Martin, in *Perspectives on Supersymmetry II*, edited by G.L. Kane (World Scientific, Singapore, 2010) pp. 1–153; see <http://zippy.physics.niu.edu/primer.html> for the latest version and errata.
6. S. Weinberg, *The Quantum Theory of Fields, Volume III: Supersymmetry* (Cambridge University Press, Cambridge, UK, 2000);  
P. Binétruy, *Supersymmetry : Theory, Experiment, and Cosmology* (Oxford University Press, Oxford, UK, 2006).
7. E. Witten, Nucl. Phys. **B188**, 513 (1981).
8. S. Dimopoulos and H. Georgi, Nucl. Phys. **B193**, 150 (1981).
9. N. Sakai, Z. Phys. **C11**, 153 (1981);  
R.K. Kaul, Phys. Lett. **109B**, 19 (1982);  
R.K. Kaul and M. Parthasarathi, Nucl. Phys. **B199**, 36 (1982).
10. L. Susskind, Phys. Reports **104**, 181 (1984).
11. L. Girardello and M. Grisaru, Nucl. Phys. **B194**, 65 (1982).
12. L.J. Hall and L. Randall, Phys. Rev. Lett. **65**, 2939 (1990);  
I. Jack and D.R.T. Jones, Phys. Lett. **B457**, 101 (1999).
13. O. Buchmüller and P. de Jong, “Supersymmetry Part II (Experiment),” in the 2013 web edition of the *Review of Particle Physics* at <http://pdg.lbl.gov>. See also there *Particle Listings: Other Searches—Supersymmetric Particles* .
14. For a review, see N. Polonsky, *Supersymmetry: Structure and phenomena. Extensions of the standard model*, Lect. Notes Phys. **M68**, 1 (2001).



15. G. Bertone, D. Hooper, and J. Silk, *Phys. Reports* **405**, 279 (2005).
16. D. Hooper, “TASI 2008 Lectures on Dark Matter,” in *The Dawn of the LHC Era, Proceedings of the 2008 Theoretical and Advanced Study Institute in Elementary Particle Physics*, Boulder, Colorado, 2–27 June 2008, edited by Tao Han (World Scientific, Singapore, 2009).
17. H. Pagels and J.R. Primack, *Phys. Rev. Lett.* **48**, 223 (1982).
18. H. Goldberg, *Phys. Rev. Lett.* **50**, 1419 (1983) [erratum: **103**, 099905 (2009)].
19. J. Ellis *et al.*, *Nucl. Phys.* **B238**, 453 (1984).
20. G. Jungman, M. Kamionkowski, and K. Griest, *Phys. Reports* **267**, 195 (1996);  
K. Griest and M. Kamionkowski, *Phys. Reports* **333**, 167 (2000).
21. F.D. Steffen, *Eur. Phys. J.* **C59**, 557 (2009).
22. M. Drees and G. Gerbier, “Dark Matter,” in the 2013 web edition of the *Review of Particle Physics* at [pdg.lbl.gov](http://pdg.lbl.gov).
23. For a review, see J.M. Cline, in *Particle Physics and Cosmology: the Fabric of Spacetime*, edited by F. Bernardeau, C. Grojean and J. Dalibard, Proceedings of Les Houches Summer School **86**, 53–116 (2007).
24. A.D. Sakharov, *JETP Lett.* **5**, 24 (1967); *Sov. Phys. JETP* **49**, 594 (1979).
25. Reviews of electroweak baryogenesis can be found in A.G. Cohen, D. Kaplan, and A. Nelson, *Ann. Rev. Nucl. Part. Sci.* **43**, 27 (1993);  
D.E. Morrissey and M.J. Ramsey-Musolf, *New J. Phys.* **14**, 125003 (2012).
26. T. Cohen, D.E. Morrissey, and A. Pierce, *Phys. Rev.* **D86**, 013009 (2012);  
D. Curtin, P. Jaiswal, and P. Meade, *JHEP* **1208**, 005 (2012);  
M. Carena *et al.*, *JHEP* **1302**, 001 (2013).
27. U. Sarkar and R. Adhikari, *Phys. Rev.* **D55**, 3836 (1997);  
R. Fok *et al.*, *Phys. Rev.* **D87**, 055018 (2013).
28. H.E. Haber and G.L. Kane, *Phys. Reports* **117**, 75 (1985).
29. M. Drees, R. Godbole, and P. Roy, *Theory and Phenomenology of Sparticles* (World Scientific, Singapore, 2005);  
H. Baer and X. Tata, *Weak Scale Supersymmetry: from*

- Superfields to Scattering Events* (Cambridge University Press, Cambridge, UK, 2006);  
 I.J.R. Aitchison, *Supersymmetry in Particle Physics: an elementary introduction* (Cambridge University Press, Cambridge, UK, 2007).
30. P. Fayet, Nucl. Phys. **B90**, 104 (1975).
  31. K. Inoue *et al.*, Prog. Theor. Phys. **67**, 1889 (1982);  
 R. Flores and M. Sher, Ann. Phys. (NY) **148**, 95 (1983).
  32. J.F. Gunion and H.E. Haber, Nucl. Phys. **B272**, 1 (1986)  
 [erratum: **B402**, 567 (1993)].
  33. I. Buchbinder, S. Kuzenko, and J. Yarevskaya, Nucl. Phys. **B411**, 665 (1994);  
 I. Antoniadis, E. Dudas, and D.M. Ghilencea, JHEP **0803**, 045 (2008).
  34. For an overview of the theory and models of the soft-supersymmetry-breaking Lagrangian, see D.J.H. Chung *et al.*, Phys. Reports **407**, 1 (2005).
  35. J. Ellis *et al.*, Phys. Rev. **D39**, 844 (1989);  
 U. Ellwanger and C. Hugonie, Eur. J. Phys. **C25**, 297 (2002);  
 U. Ellwanger, C. Hugonie, and A.M. Teixeira, Phys. Reports **496**, 1 (2010);  
 M. Maniatis, Int. J. Mod. Phys. A **25**, 3505 (2010).
  36. P. Fayet, Phys. Lett. **69B**, 489 (1977);  
 G. Farrar and P. Fayet, Phys. Lett. **76B**, 575 (1978).
  37. P. Fayet, Phys. Lett. **84B**, 421 (1979); Phys. Lett. **86B**, 272 (1979).
  38. D.Z. Freedman and A. Van Proeyen, *Supergravity* (Cambridge University Press, Cambridge, UK, 2012).
  39. S. Deser and B. Zumino, Phys. Rev. Lett. **38**, 1433 (1977);  
 E. Cremmer *et al.*, Phys. Lett. **79B**, 231 (1978).
  40. R. Casalbuoni *et al.*, Phys. Lett. **B215**, 313 (1988); Phys. Rev. **D39**, 2281 (1989);  
 A.L. Maroto and J.R. Pelaez, Phys. Rev. **D62**, 023518 (2000).
  41. Z. Komargodski and N. Seiberg, JHEP **0909**, 066 (2009);  
 I. Antoniadis *et al.*, Theor. Math. Phys. **170**, 26 (2012).
  42. A.H. Chamseddine, R. Arnowitt, and P. Nath, Phys. Rev. Lett. **49**, 970 (1982);  
 R. Barbieri, S. Ferrara, and C.A. Savoy, Phys. Lett. **119B**, 343 (1982);  
 L. Ibáñez, Nucl. Phys. **B218**, 514 (1982);  
 H.-P. Nilles, M. Srednicki, and D. Wyler, Phys. Lett.

- 120B**, 346 (1983); **124B**, 337 (1983);  
 E. Cremmer, P. Fayet, and L. Girardello, Phys. Lett. **122B**, 41 (1983);  
 N. Ohta, Prog. Theor. Phys. **70**, 542 (1983).
43. L. Alvarez-Gaumé, J. Polchinski, and M.B. Wise, Nucl. Phys. **B221**, 495 (1983).
44. L.J. Hall, J. Lykken, and S. Weinberg, Phys. Rev. **D27**, 2359 (1983).
45. S.K. Soni and H.A. Weldon, Phys. Lett. **126B**, 215 (1983);  
 Y. Kawamura, H. Murayama, and M. Yamaguchi, Phys. Rev. **D51**, 1337 (1995).
46. See, *e.g.*, A. Brignole, L.E. Ibáñez, and C. Muñoz, in *Perspectives on Supersymmetry II*, edited by G.L. Kane (World Scientific, Singapore, 2010) pp. 244–268.
47. A.B. Lahanas and D.V. Nanopoulos, Phys. Reports **145**, 1 (1987).
48. J.L. Feng, A. Rajaraman, and F. Takayama, Phys. Rev. Lett. **91**, 011302 (2003); Phys. Rev. **D68**, 063504 (2003); Gen. Rel. Grav. **36**, 2575 (2004).
49. L. Randall and R. Sundrum, Nucl. Phys. **B557**, 79 (1999).
50. F. D’Eramo, J. Thaler, and Z. Thomas, JHEP **1206**, 151 (2012); JHEP **1309**, 125 (2013);  
 S.P. de Alwis, Phys. Rev. **D77**, 105020 (2008); JHEP **1301**, 006 (2013).
51. See *e.g.*, I. Jack, D.R.T. Jones, and R. Wild, Phys. Lett. **B535**, 193 (2002);  
 B. Murakami and J.D. Wells, Phys. Rev. **D68**, 035006 (2003);  
 R. Kitano, G.D. Kribs, and H. Murayama, Phys. Rev. **D70**, 035001 (2004);  
 R. Hodgson *et al.*, Nucl. Phys. **B728**, 192 (2005);  
 D.R.T. Jones and G.G. Ross, Phys. Lett. **B642**, 540 (2006).
52. M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189**, 575 (1981);  
 S. Dimopoulos and S. Raby, Nucl. Phys. **B192**, 353 (1982); **B219**, 479 (1983);  
 M. Dine and W. Fischler, Phys. Lett. **110B**, 227 (1982);  
 C. Nappi and B. Ovrut, Phys. Lett. **113B**, 175 (1982);  
 L. Alvarez-Gaumé, M. Claudson, and M. Wise, Nucl. Phys. **B207**, 96 (1982).

53. 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).
54. G.F. Giudice and R. Rattazzi, Phys. Reports **322**, 419  
(1999).
55. E. Poppitz and S.P. Trivedi, Phys. Rev. **D55**, 5508  
(1997);  
H. Murayama, Phys. Rev. Lett. **79**, 18 (1997);  
M.A. Luty and J. Terning, Phys. Rev. **D57**, 6799 (1998);  
K. Agashe, Phys. Lett. **B435**, 83 (1998);  
N. Arkani-Hamed, J. March-Russell, and H. Murayama,  
Nucl. Phys. **B509**, 3 (1998);  
C. Csaki, Y. Shirman, and J. Terning, JHEP **0705**, 099  
(2007);  
M. Ibe and R. Kitano, Phys. Rev. **D77**, 075003 (2008).
56. M. Kawasaki *et al.*, Phys. Rev. **D78**, 065011 (2008).
57. M.J. Strassler and K. M. Zurek, Phys. Lett. **B651**, 374  
(2007);  
T. Han *et al.*, JHEP **0807**, 008 (2008).
58. M.J. Strassler, arXiv:hep-ph/0607160;  
K.M. Zurek, Phys. Rev. **D79**, 115002 (2009).
59. See *e.g.*, M. Quiros, in *Particle Physics and Cosmology:  
The Quest for Physics Beyond the Standard Model(s),  
Proceedings of the 2002 Theoretical Advanced Study Insti-  
tute in Elementary Particle Physics (TASI 2002)*, edited  
by H.E. Haber and A.E. Nelson (World Scientific, Singa-  
pore, 2004) pp. 549–601;  
C. Csaki, in *ibid.*, pp. 605–698.
60. See, *e.g.*, J. Parsons and A. Pomarol, “Extra Dimen-  
sions,” in the 2013 web edition of the *Review of Particle  
Physics* at <http://pdg.lbl.gov>.
61. See *e.g.*, V.A. Rubakov, Phys. Usp. **44**, 871 (2001);  
J. Hewett and M. Spiropulu, Ann. Rev. Nucl. Part. Sci.  
**52**, 397 (2002).
62. 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).
63. J. Scherk and J.H. Schwarz, Phys. Lett. **82B**, 60 (1979);  
Nucl. Phys. **B153**, 61 (1979).
64. See, *e.g.*, R. Barbieri, L.J. Hall, and Y. Nomura, Phys.  
Rev. **D66**, 045025 (2002); Nucl. Phys. **B624**, 63 (2002).

65. N. Arkani-Hamed and S. Dimopoulos, JHEP **0506**, 073 (2005);  
G.F. Giudice and A. Romanino, Nucl. Phys. **B699**, 65 (2004) [erratum: **B706**, 487 (2005)].
66. G. Aad *et al.* [ATLAS Collab.], Phys. Lett. **B716**, 1 (2012);  
S. Chatrchyan *et al.* [CMS Collab.], Phys. Lett. **B716**, 30 (2012).
67. M. Carena, *et al.*, “The Higgs Boson  $H^0$ ,” in the 2013 web edition of the *Review of Particle Physics* at [pdg.lbl.gov](http://pdg.lbl.gov).
68. S. Khalil, Int. J. Mod. Phys. **A18**, 1697 (2003).
69. G.F. Giudice and A. Strumia, Nucl. Phys. **B858**, 63 (2012).
70. A. Arvanitaki *et al.*, JHEP **1302**, 126 (2013);  
N. Arkani-Hamed *et al.*, [arXiv:1212.6971](https://arxiv.org/abs/1212.6971) [hep-ph].
71. L.J. Hall and Y. Nomura, JHEP **1201**, 082 (2012);  
M. Ibe and T.T. Yanagida, Phys. Lett. **B709**, 374 (2012).
72. H.E. Haber, in *Recent Directions in Particle Theory, Proceedings of the 1992 Theoretical Advanced Study Institute in Particle Physics*, edited by J. Harvey and J. Polchinski (World Scientific, Singapore, 1993) pp. 589–686.
73. J.M. Frere, D.R.T. Jones, and S. Raby, Nucl. Phys. **B222**, 11 (1983);  
J.P. Derendinger and C.A. Savoy, Nucl. Phys. **B237**, 307 (1984);  
J.F. Gunion, H.E. Haber, and M. Sher, Nucl. Phys. **B306**, 1 (1988);  
J.A. Casas, A. Lleyda, and C. Munoz, Nucl. Phys. **B471**, 3 (1996).
74. S.P. Martin, Phys. Rev. **D61**, 035004 (2000).
75. S. Dimopoulos and D. Sutter, Nucl. Phys. **B452**, 496 (1995);  
D.W. Sutter, Stanford Ph. D. thesis, [hep-ph/9704390](https://arxiv.org/abs/hep-ph/9704390).
76. H.E. Haber, Nucl. Phys. B (Proc. Suppl.) **62A-C**, 469 (1998).
77. R.M. Barnett, J.F. Gunion, and H.E. Haber, Phys. Lett. **B315**, 349 (1993);  
H. Baer, X. Tata, and J. Woodside, Phys. Rev. **D41**, 906 (1990).
78. S.M. Bilenky, E.Kh. Khristova, and N.P. Nedelcheva, Phys. Lett. **B161**, 397 (1985); Bulg. J. Phys. **13**, 283 (1986);  
G. Moortgat-Pick and H. Fraas, Eur. Phys. J. **C25**, 189 (2002).

79. J. Rosiek, Phys. Rev. **D41**, 3464 (1990) [erratum: hep-ph/9511250]. The most recent corrected version of this manuscript can be found on the author’s webpage, [www.fuw.edu.pl/~rosiek/physics/prd41.html](http://www.fuw.edu.pl/~rosiek/physics/prd41.html).
80. J. Alwall *et al.*, JHEP **0709**, 028 (2007). The MadGraph homepage is located at [madgraph.hep.uiuc.edu/](http://madgraph.hep.uiuc.edu/).
81. T. Hahn, Comput. Phys. Commun. **140**, 418 (2001); T. Hahn and C. Schappacher, Comput. Phys. Commun. **143**, 54 (2002). The FeynArts homepage is located at [www.feynarts.de/](http://www.feynarts.de/).
82. A. Pukhov *et al.*, INP MSU report 98-41/542 (arXiv: hep-ph/9908288); E. Boos *et al.* [CompHEP Collab.], Nucl. Instrum. Meth. A534, 250 (2004). The CompHEP homepage is located at <http://comphep.sinp.msu.ru>.
83. D.M. Pierce *et al.*, Nucl. Phys. **B491**, 3 (1997).
84. P. Skands *et al.*, JHEP **07** 036 (2004); B.C. Allanach *et al.*, Comput. Phys. Commun. **180**, 8 (2009). The Supersymmetry Les Houches Accord homepage is located at [home.fnal.gov/~skands/slha/](http://home.fnal.gov/~skands/slha/).
85. For further details, see *e.g.*, Appendix C of Ref. 28 and Appendix A of Ref. 32.
86. J.L. Kneur and G. Moultaka, Phys. Rev. **D59**, 015005 (1999).
87. S.Y. Choi *et al.*, Eur. Phys. J. **C14**, 535 (2000).
88. R.A. Horn and C.R. Johnson, *Matrix Analysis*, 2nd Edition (Cambridge University Press, Cambridge, UK, 2003).
89. H.K. Dreiner, H.E. Haber, and S.P. Martin, Phys. Reports **494**, 1 (2010).
90. S.Y. Choi *et al.*, Eur. Phys. J. **C22**, 563 (2001); **C23**, 769 (2002).
91. G.J. Gounaris, C. Le Mouel, and P.I. Porfyriadis, Phys. Rev. **D65**, 035002 (2002); G.J. Gounaris and C. Le Mouel, Phys. Rev. **D66**, 055007 (2002).
92. T. Takagi, Japan J. Math. **1**, 83 (1925).
93. S.Y. Choi *et al.*, Nucl. Phys. **B778**, 85 (2007).
94. M.M. El Kheishen, A.A. Aboshousha, and A.A. Shafik, Phys. Rev. **D45**, 4345 (1992); M. Guchait, Z. Phys. **C57**, 157 (1993) [erratum: **C61**, 178 (1994)].
95. T. Hahn, preprint MPP-2006-85, [physics/0607103](http://physics/0607103).

96. K. Hikasa and M. Kobayashi, Phys. Rev. **D36**, 724 (1987);  
 F. Gabbiani and A. Masiero, Nucl. Phys. **B322**, 235 (1989);  
 Ph. Brax and C.A. Savoy, Nucl. Phys. **B447**, 227 (1995).
97. J. Ellis and S. Rudaz, Phys. Lett. **128B**, 248 (1983);  
 F. Browning, D. Chang, and W.Y. Keung, Phys. Rev. **D64**, 015010 (2001);  
 A. Bartl *et al.*, Phys. Lett. **B573**, 153 (2003); Phys. Rev. **D70**, 035003 (2004).
98. J.F. Gunion *et al.*, *The Higgs Hunter's Guide* (Westview Press, Boulder, CO, 2000);  
 M. Carena and H.E. Haber, Prog. Part. Nucl. Phys. **50**, 63 (2003);  
 A. Djouadi, Phys. Reports **459**, 1 (2008).
99. P. Bechtle *et al.*, Eur. Phys. J. **C73**, 2354 (2013);  
 M. Carena *et al.*, Eur. Phys. J. **C73**, 2552 (2013).
100. H.E. Haber and M. Sher, Phys. Rev. **D35**, 2206 (1987).
101. L.J. Hall, D. Pinner, and J.T. Ruderman, JHEP **1204**, 131 (2012).
102. 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).
103. See, *e.g.*, G. Degrassi *et al.*, Eur. Phys. J. **C28**, 133 (2003);  
 S.P. Martin, Phys. Rev. **D75**, 055005 (2007);  
 P. Kant *et al.*, JHEP **1008**, 104 (2010).
104. 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);  
 M. Frank *et al.*, JHEP **0702**, 047 (2007);  
 S. Heinemeyer *et al.*, Phys. Lett. **B652**, 300 (2007).
105. 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);  
V.D. Barger *et al.*, Phys. Rev. **D64**, 056007 (2001);  
S. Abel, S. Khalil, and O. Lebedev, Nucl. Phys. **B606**, 151 (2001);  
K.A. Olive *et al.*, Phys. Rev. **D72**, 075001 (2005);  
G.F. Giudice and A. Romanino, Phys. Lett. **B634**, 307 (2006).
106. A. Masiero and L. Silvestrini, in *Perspectives on Supersymmetry*, edited by G.L. Kane (World Scientific, Singapore, 1998) pp. 423–441.
107. M. Pospelov and A. Ritz, Annals Phys. **318**, 119 (2005).
108. See, *e.g.*, F. Gabbiani *et al.*, Nucl. Phys. **B477**, 321 (1996);  
A. Masiero, and O. Vives, New J. Phys. **4**, 1 (2002).
109. For a review and references to the original literature, see: M.J. Ramsey-Musolf and S. Su, Phys. Reports **456**, 1 (2008).
110. M. Carena, A. Menon and C.E.M. Wagner, Phys. Rev. **D79**, 075025 (2009);  
S. Jager, Eur. Phys. J. **C59**, 497 (2009);  
W. Altmannshofer *et al.*, Nucl. Phys. **B830**, 17 (2010).
111. M.B. Einhorn and D.R.T. Jones, Nucl. Phys. **B196**, 475 (1982).
112. For a review, see R.N. Mohapatra, in *Particle Physics 1999*, ICTP Summer School in Particle Physics, Trieste, Italy, edited by G. Senjanovic and A.Yu. Smirnov (World Scientific, Singapore, 2000) pp. 336–394;  
W.J. Marciano and G. Senjanovic, Phys. Rev. **D25**, 3092 (1982).
113. L.E. Ibáñez and G.G. Ross, Phys. Lett. **B110**, 215 (1982).
114. J. Abdullah *et al.* [DELPHI Collab.], Eur. Phys. J. **C31**, 421 (2004).
115. H.K. Dreiner *et al.*, Eur. Phys. J. **C62**, 547 (2009).
116. G.F. Giudice *et al.*, JHEP **9812**, 027 (1998);  
A. Pomarol and R. Rattazzi, JHEP **9905**, 013 (1999);  
D.W. Jung and J.Y. Lee, JHEP **0903**, 123 (2009).
117. J.F. Gunion and H.E. Haber, Phys. Rev. **D37**, 2515 (1988);  
S.Y. Choi, M. Drees, and B. Gaissmaier, Phys. Rev. **D70**, 014010 (2004).
118. J.L. Feng *et al.*, Phys. Rev. Lett. **83**, 1731 (1999);  
J.F. Gunion and S. Mrenna, Phys. Rev. **D62**, 015002 (2000).



- 119. T. Gherghetta, G.F. Giudice, and J.D. Wells, Nucl. Phys. **B559**, 27 (1999).
- 120. M. Endo, M. Yamaguchi, and K. Yoshioka, Phys. Rev. **D72**, 015004 (2005);  
K. Choi, K.S. Jeong, and K.-I. Okumura, JHEP **0509**, 039 (2005);  
O. Loaiza-Brito *et al.*, AIP Conf. Proc. **805**, 198 (2006).
- 121. See *e.g.*, G. D’Ambrosio *et al.*, Nucl. Phys. **B465**, 155 (2002).
- 122. For a review of minimal flavor violation in supersymmetric theories, see C. Smith, Acta Phys. Polon. Supp. **3**, 53 (2010).
- 123. M. Drees and S.P. Martin, in *Electroweak Symmetry Breaking and New Physics at the TeV Scale*, edited by T. Barklow *et al.* (World Scientific, Singapore, 1996) pp. 146–215.
- 124. J.R. Ellis *et al.*, Phys. Lett. **B573**, 162 (2003); Phys. Rev. **D70**, 055005 (2004).
- 125. L.E. Ibáñez and D. Lüst, Nucl. Phys. **B382**, 305 (1992);  
B. de Carlos, J.A. Casas, and C. Muñoz, Phys. Lett. **B299**, 234 (1993);  
V. Kaplunovsky and J. Louis, Phys. Lett. **B306**, 269 (1993);  
A. Brignole, L.E. Ibáñez, and C. Muñoz, Nucl. Phys. **B422**, 125 (1994) [erratum: **B436**, 747 (1995)].
- 126. A. Arbey *et al.*, Phys. Rev. **D87**, 115020 (2013).
- 127. P. Draper *et al.*, Phys. Rev. **D85**, 095007 (2012).
- 128. P. Meade, N. Seiberg, and D. Shih, Prog. Theor. Phys. Suppl. **177**, 143 (2009);  
M. Buican *et al.*, JHEP **0903**, 016 (2009).
- 129. A. Rajaraman *et al.*, Phys. Lett. **B678**, 367 (2009);  
L.M. Carpenter *et al.*, Phys. Rev. **D79**, 035002 (2009).
- 130. S. Ambrosanio, G.D. Kribs, and S.P. Martin, Nucl. Phys. **B516**, 55 (1998).
- 131. For a review and guide to the literature, see J.F. Gunion and H.E. Haber, in *Perspectives on Supersymmetry II*, edited by G.L. Kane (World Scientific, Singapore, 2010) pp. 420–445.
- 132. A. Djouadi, J.L. Kneur, and G. Moultaka, Comput. Phys. Commun. **176**, 426–455 (2007);  
C.F. Berger *et al.*, JHEP **0902**, 023 (2009).
- 133. H. Baer *et al.*, JHEP **1205**, 109 (2012).

- 134. R. Barbieri and G.F. Giudice, Nucl. Phys. **B305**, 63 (1988).
- 135. 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);  
P. Athron and D.J. Miller, Phys. Rev. **D76**, 075010 (2007);  
M.E. Cabrera, J.A. Casas, and R.R. de Austri, JHEP **0903**, 075 (2009).
- 136. D.M. Ghilencea and G.G. Ross, Nucl. Phys. **B868**, 65 (2013).
- 137. G.L. Kane and S.F. King, Phys. Lett. **B451**, 113 (1999);  
M. Bastero-Gil, G.L. Kane, and S.F. King, Phys. Lett. **B474**, 103 (2000);  
J.A. Casas, J.R. Espinosa, and I. Hidalgo, JHEP **0401**, 008 (2004);  
J. Abe, T. Kobayashi, and Y. Omura, Phys. Rev. **D76**, 015002 (2007);  
R. Essig and J.-F. Fortin, JHEP **0804**, 073 (2008).
- 138. P. Bechtle *et al.*, JHEP **1206**, 098 (2012).
- 139. O. Buchmüller *et al.*, Eur. Phys. J. **C72**, 2243 (2012).
- 140. R. Barbieri and A. Strumia, hep-ph/0007265.
- 141. L. Giusti, A. Romanino, and A. Strumia, Nucl. Phys. **B550**, 3 (1999);  
H.C. Cheng and I. Low, JHEP **0309**, 051 (2003); **0408**, 061 (2004).
- 142. H. Baer, V. Barger, and D. Mickelson, Phys. Rev. **D88**, 095013 (2013).
- 143. H. Baer *et al.*, Phys. Rev. **D87**, 035017 (2013); 115028 (2013);  
J.L. Feng, Ann. Rev. Nucl. Part. Sci. **63**, 351 (2013).
- 144. J. Feng, K. Matchev, and T. Moroi, Phys. Rev. Lett. **84**, 2322 (2000); Phys. Rev. **D61**, 075005 (2000);  
J. Feng and F. Wilczek, Phys. Lett. **B631**, 170 (2005);  
D. Horton and G.G. Ross, Nucl. Phys. **B830**, 221 (2010).
- 145. M. Drees, Phys. Rev. **D33**, 1468 (1986);  
S. Dimopoulos and G.F. Giudice, Phys. Lett. **B357**, 573 (1995);  
A. Pomarol and D. Tommasini, Nucl. Phys. **B466**, 3 (1996).
- 146. M. Dine, A. Kagan, and S. Samuel, Phys. Lett. **B243**, 250 (1990);

- A.G. Cohen, D.B. Kaplan, and A.E. Nelson, Phys. Lett. **B388**, 588 (1996).
147. K.L. Chan, U. Chattopadhyay, and P. Nath, Phys. Rev. **D58**, 096004 (1998);  
R. Kitano and Y. Nomura, Phys. Rev. **D73**, 095004 (2006);  
M. Perelstein and C. Spethmann, JHEP **0704**, 070 (2007);  
H. Abe, T. Kobayashi, and Y. Omura, Phys. Rev. **D76**, 015002 (2007);  
D. Horton and G.G. Ross, Nucl. Phys. **B830**, 221 (2010);  
H. Baer *et al.*, JHEP **1010**, 018 (2010);  
M. Asano *et al.*, JHEP **1012**, 019 (2010);  
H. Baer, V. Barger, and P. Huang, JHEP **1111**, 031 (2011);  
C. Brust *et al.*, JHEP **1203**, 103 (2012);  
M. Papucci, J.T. Ruderman, and A. Weiler, JHEP **1209**, 035 (2012);  
H.K. Dreiner, M. Kramer, and J. Tattersall, Europhys. Lett. **99**, 61001 (2012).
148. N. Arkani-Hamed *et al.*, hep-ph/0703088;  
J. Alwall *et al.*, Phys. Rev. D **79**, 015005 (2009);  
J. Alwall, P. Schuster, and N. Toro, Phys. Rev. **D79**, 075020 (2009);  
D.S.M. Alves, E. Izaguirre, and J.G. Wacker, Phys. Lett. **B702**, 64 (2011); JHEP **1110**, 012 (2011);  
D. Alves *et al.*, J. Phys. G **39**, 105005 (2012).
149. S.P. Martin, Phys. Rev. **D75**, 115005 (2007); Phys. Rev. **D78**, 055019 (2009);  
T.J. LeCompte and S.P. Martin, Phys. Rev. **D85**, 035023 (2012).
150. J. Fan, M. Reece, and J.T. Ruderman, JHEP **1111**, 012 (2011).
151. R. Dermisek and J.F. Gunion, Phys. Rev. Lett. **95**, 041801 (2005); Phys. Rev. **D75**, 095019 (2007); Phys. Rev. **D76**, 095006 (2007).
152. G.G. Ross and K. Schmidt-Hoberg, Nucl. Phys. **B862**, 710 (2012); JHEP **1208**, 074 (2012);  
A. Kaminska, G.G. Ross, and K. Schmidt-Hoberg, JHEP **1311**, 209 (2013).
153. S.P. Martin and J.D. Wells, Phys. Rev. **D86**, 035017 (2012).
154. B. Bellazzini *et al.*, Phys. Rev. **D79**, 095003 (2009).
155. D. Stockinger, J. Phys. **G34**, R45 (2007).

156. J.P. Miller *et al.*, *Ann. Rev. Nucl. Part. Sci.* **62**, 237 (2012);  
E. de Rafael, *Nucl. Phys. Proc. Suppl.* **234**, 193 (2013).
157. M. Ibe, T. T. Yanagida, and N. Yokozaki, *JHEP* **1308**, 067 (2013).
158. M. Benayoun *et al.*, *Eur. Phys. J.* **C72**, 1848 (2012).
159. A. Limosani *et al.* [Belle Collab.], *Phys. Rev. Lett.* **103**, 241801 (2009);  
J.P. Lees *et al.* [BaBar Collab.], *Phys. Rev. Lett.* **109**, 191801 (2012); *Phys. Rev.* **D86**, 112008 (2012).
160. M. Misiak *et al.*, *Phys. Rev. Lett.* **98**, 022002 (2007);  
T. Becher and M. Neubert, *Phys. Rev. Lett.* **98**, 022003 (2007).
161. See, *e.g.*, M. Ciuchini *et al.*, *Phys. Rev.* **D67**, 075016 (2003);  
T. Hurth, *Rev. Mod. Phys.* **75**, 1159 (2003);  
F. Mahmoudi, *JHEP* **0712**, 026 (2007);  
K.A. Olive and L. Velasco-Sevilla, *JHEP* **0805**, 052 (2008).
162. S.R. Choudhury and N. Gaur, *Phys. Lett.* **B451**, 86 (1999);  
K.S. Babu and C.F. Kolda, *Phys. Rev. Lett.* **84**, 228 (2000);  
G. Isidori and A. Retico, *JHEP* **0111**, 001 (2001); *JHEP* **0209**, 063 (2002).
163. R. Aaij *et al.* [LHCb Collab.], *Phys. Rev. Lett.* **110**, 021801 (2013); **111**, 101805 (2013).
164. F. Mahmoudi, S. Neshatpour, and J. Orloff, *JHEP* **1208**, 092 (2012);  
A. Arbey *et al.*, *Phys. Rev.* **D87**, 035026 (2013).
165. J. Erler and A. Freitas, “Electroweak Model and Constraints on New Physics,” in this *Review*.
166. J.R. Ellis *et al.*, *JHEP* **0708**, 083 (2007);  
S. Heinemeyer *et al.*, *JHEP* **0808**, 087 (2008);  
G.-C. Cho *et al.*, *JHEP* **1111**, 068 (2011).
167. For the current status of neutrino masses and mixing, see:  
M.C. Gonzalez-Garcia *et al.*, *JHEP* **1212**, 123 (2012);  
D.V. Forero, M. Tortola, and J.W.F. Valle, *Phys. Rev.* **D86**, 073012 (2012);  
G.L. Fogli *et al.*, *Phys. Rev.* **D86**, 013012 (2012).
168. See the section on neutrinos in “Particle Listings—Leptons” in the 2013 web edition of the *Review of Particle Physics* at <http://pdg.lbl.gov>.
169. K. Zuber, *Phys. Reports* **305**, 295 (1998).

170. For a review of neutrino masses in supersymmetry, see *e.g.*, B. Mukhopadhyaya, *Proc. Indian National Science Academy* **A70**, 239 (2004);  
M. Hirsch and J.W.F. Valle, *New J. Phys.* **6**, 76 (2004).
171. F. Borzumati and Y. Nomura, *Phys. Rev.* **D64**, 053005 (2001).
172. P. Minkowski, *Phys. Lett.* **67B**, 421 (1977);  
M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by D. Freedman and P. van Nieuwenhuizen (North Holland, Amsterdam, 1979) p. 315;  
T. Yanagida, *Prog. Theor. Phys.* **64**, 1103 (1980);  
R. Mohapatra and G. Senjanovic, *Phys. Rev. Lett.* **44**, 912 (1980); *Phys. Rev.* **D23**, 165 (1981).
173. J. Hisano *et al.*, *Phys. Lett.* **B357**, 579 (1995);  
J. Hisano *et al.*, *Phys. Rev.* **D53**, 2442 (1996);  
J.A. Casas and A. Ibarra, *Nucl. Phys.* **B618**, 171 (2001);  
J. Ellis *et al.*, *Phys. Rev.* **D66**, 115013 (2002);  
A. Masiero, S.K. Vempati, and O. Vives, *New J. Phys.* **6**, 202 (2004);  
E. Arganda *et al.*, *Phys. Rev.* **D71**, 035011 (2005);  
F.R. Joaquim and A. Rossi, *Phys. Rev. Lett.* **97**, 181801 (2006);  
J.R. Ellis and O. Lebedev, *Phys. Lett.* **B653**, 411 (2007).
174. Y. Grossman and H.E. Haber, *Phys. Rev. Lett.* **78**, 3438 (1997);  
A. Dedes, H.E. Haber, and J. Rosiek, *JHEP* **0711**, 059 (2007).
175. M. Hirsch, H.V. Klapdor-Kleingrothaus, and S.G. Kovalenko, *Phys. Lett.* **B398**, 311 (1997);  
L.J. Hall, T. Moroi, and H. Murayama, *Phys. Lett.* **B424**, 305 (1998);  
K. Choi, K. Hwang, and W.Y. Song, *Phys. Rev. Lett.* **88**, 141801 (2002);  
T. Honkavaara, K. Huitu, and S. Roy, *Phys. Rev.* **D73**, 055011 (2006).
176. M. Chemtob, *Prog. Part. Nucl. Phys.* **54**, 71 (2005);  
R. Barbier *et al.*, *Phys. Reports* **420**, 1 (2005).
177. H. Dreiner, in *Perspectives on Supersymmetry II*, edited by G.L. Kane (World Scientific, Singapore, 2010) pp. 565–583.
178. B.C. Allanach, A. Dedes, and H.K. Dreiner, *Phys. Rev.* **D60**, 075014 (1999).
179. L.E. Ibáñez and G.G. Ross, *Nucl. Phys.* **B368**, 3 (1992);  
L.E. Ibáñez, *Nucl. Phys.* **B398**, 301 (1993).

180. A. Dedes, S. Rimmer, and J. Rosiek, JHEP **0608**, 005 (2006);  
B.C. Allanach and C.H. Kom, JHEP **0804**, 081 (2008);  
H.K. Dreiner *et al.*, Phys. Rev. **D84**, 113005 (2011).
181. L.E. Ibáñez and G.G. Ross, Nucl. Phys. **B368**, 3 (1992).
182. H.K. Dreiner, C. Luhn, and M. Thormeier Phys. Rev. **D73**, 075007 (2006).
183. K. Tamvakis, Phys. Lett. **B382**, 251 (1996);  
G. Eyal and Y. Nir, JHEP, **06**, 024 (1999);  
A. Florex *et al.*, Phys. Rev. **D87**, 095010 (2013).
184. B.C. Allanach and B. Gripaios, JHEP **1205**, 062 (2012);  
M. Asano, K. Rolbiecki, and K. Sakurai, JHEP **1301**, 128 (2013).
185. See *e.g.*, J.C. Romao, Nucl. Phys. Proc. Suppl. **81**, 231 (2000);  
Y. Grossman and S. Rakshit, Phys. Rev. **D69**, 093002 (2004).
186. 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).
187. Y. Grossman and H.E. Haber, Phys. Rev. **D59**, 093008 (1999).
188. 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).
189. H.K. Dreiner, P. Richardson, and M.H. Seymour, Phys. Rev. **D63**, 055008 (2001).
190. J.E. Kim and H.P. Nilles, Phys. Lett. **B138**, 150 (1984).
191. J.E. Kim and H.P. Nilles, Mod. Phys. Lett. **A9**, 3575 (1994).
192. G.F. Giudice and A. Masiero, Phys. Lett. **B206**, 480 (1988).
193. J.A. Casas and C. Munoz, Phys. Lett. **B306**, 288 (1993).
194. M. Cvetič *et al.*, Phys. Rev. **D56**, 2861 (1997) [erratum: **D58**, 119905 (1998)].
195. A. Delgado, G. Nardini, and M. Quiros, Phys. Rev. **D86**, 115010 (2012).
196. P. Fayet, Phys. Lett. **78B**, 417 (1978).

197. For a recent review, see K. Benakli, Fortsch. Phys. **59**, 1079 (2011).
198. P.J. Fox, A.E. Nelson, and N. Weiner, JHEP **0208**, 035 (2002).
199. K. Benakli and M.D. Goodsell, Nucl. Phys. **B816**, 185 (2009); **B840**, 1 (2010).
200. G.D. Kribs, E. Poppitz, and N. Weiner, Phys. Rev. **D78**, 055010 (2008).
201. K. Benakli, M.D. Goodsell, and F. Staub, JHEP **1306**, 073 (2013).
202. See *e.g.*, J.L. Hewett and T.G. Rizzo, Phys. Reports **183**, 193 (1989).
203. S.F. King, S. Moretti, and R. Nevzorov, Phys. Lett. **B634**, 278 (2006); Phys. Rev. **D73**, 035009 (2006).