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Impact of physical principles at very high-energy scales on the superparticle mass spectrum

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ABSTRACT: We survey a variety of proposals for new physics at high scales that serve to relate the multitude of soft supersymmetry breaking parameters of the MSSM. We focus on models where the new physics results in non-universal soft parameters, in sharp contrast with the usually assumed mSUGRA framework. These include (i) SU(5) and SO(10) grand unified (GUT) models, (ii) the MSSM plus a right-handed neutrino, (iii) models with effective supersymmetry, (iv) models with anomaly-mediated SUSY breaking and gaugino mediated SUSY breaking, (v) models with non-universal soft terms due to string dynamics, and (vi) models based on M-theory. We outline the physics behind these models, point out some distinctive features of the weak-scale sparticle spectrum, and allude to implications for collider experiments. To facilitate future studies, for each of these scenarios, we describe how collider events can be generated using the program ISAJET. Our hope is that detailed studies of a variety of alternatives will help point to the physics underlying SUSY breaking and how this is mediated to the observable sector, once sparticles are discovered and their properties measured.

KEYWORDS: Supersymmetric Effective Theories, Supersymmetric Standard Model, GUT.

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1. Introduction

The Minimal Supersymmetric Standard Model (MSSM) is a well-motivated extension of the Standard Model (SM) that includes broken supersymmetry (SUSY) at the weak scale [1]. To construct the MSSM, one postulates:

- the gauge group and the matter content of the SM, where the various fields are replaced by superfields:
$$\hat{Q}_i = \begin{pmatrix} \hat{u}_i \\ \hat{d}_i \end{pmatrix}, \qquad \hat{L}_i = \begin{pmatrix} \hat{\nu}_i \\ \hat{e}_i \end{pmatrix}, \quad \hat{U}^c_i, \quad \hat{D}^c_i, \quad \hat{E}^c_i,$$
where $i = 1, 2, 3$ corresponding to the various generations;

- an extended Higgs sector that includes two different SU(2) doublet Higgs superfields
$$\hat{H}_u(\mathbf{2}) = \begin{pmatrix} \hat{h}^+_u \\ \hat{h}^0_u \end{pmatrix} \qquad \text{and} \qquad \hat{H}_d(\bar{\mathbf{2}}) = \begin{pmatrix} \hat{h}^-_d \\ \hat{h}^0_d \end{pmatrix};$$
to allow superpotential Yukawa couplings (and hence, masses) for both up and down type fermions. The introduction of two doublets of Higgsinos is also just right to cancel the chiral anomaly from the gauginos.

- an R-parity conserving renormalizable superpotential,[1]
$$\hat{f} = \mu \hat{H}^a_u \hat{H}_{da} + f_u \epsilon_{ab} \hat{Q}^a \hat{H}^b_u \hat{U}^c + f_d \hat{Q}^a \hat{H}_{da} \hat{D}^c + f_e \hat{L}^a \hat{H}_{da} \hat{E}^c + \cdots,$$
where ϵ_{ab} is the completely antisymmetric SU(2) tensor with $\epsilon_{12} = 1$, and the ellipses refer to Yukawa couplings for the second and third generations;

- soft supersymmetry breaking (SSB) terms consistent with Lorentz invariance and SM gauge invariance,
$$\begin{aligned}\mathcal{L}_{\text{soft}} = & - \sum_r m_r^2 |\phi_r|^2 - \frac{1}{2} \sum_\lambda M_\lambda \bar{\lambda}_\alpha \lambda_\alpha + \left[B\mu \tilde{H}_d \tilde{H}_u + \text{h.c.} \right] + \\ & + \left[A_u f_u \epsilon \tilde{Q} \tilde{H}_u \tilde{u}^\dagger_R + A_d f_d \tilde{Q} \tilde{H}_d \tilde{d}^\dagger_R + A_e f_e \tilde{L} \tilde{H}_d \tilde{e}^\dagger_R + \cdots + \text{h.c.} \right],\end{aligned}$$
where contraction over the SU(2) indices is understood, and the ellipses again refer to terms of the second and third generation trilinear scalar couplings. In practice, because only third generation Yukawa couplings are sizeable, the A-parameters of just the third family are frequently relevant.

Although we have not shown this explicitly, the Yukawa couplings and the A-parameters are, in general, (complex) matrices in generation space. The resulting framework then requires ≥ 100 parameters beyond those of the SM [2], and hence is not very predictive. Since, the phenomenology that we consider is generally in-

[1] Our sign convention for the μ-term is defined by the chargino and neutralino mass matrices given in eqs. (33) and (34) of the review by X. Tata, ref. [1].

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sensitive to inter-generation mixing of quarks and squarks, we assume that these matrices are diagonal. Furthermore, since we do not discuss CP violating effects, we take the superpotential and soft SUSY breaking parameters to be real. Even so, a large number of additional parameters remains. Most of these occur in the SSB sector of the model, which simply reflects our ignorance of the mechanism of supersymmetry breaking. To gain predictivity, despite the lack of a compelling model of SUSY breaking, we must make additional simplifying assumptions about symmetries of interactions at energy scales not directly accessible to experiments, or postulate other physical principles that determine the origin of the soft SUSY breaking terms.

The most popular model in which to embed the MSSM is the minimal supergravity model (mSUGRA) [3]. In this model, supersymmetry is broken in a "hidden sector" which consists of fields which couple to the fields of the visible sector (the MSSM fields) only gravitationally. Within the framework of supergravity grand unification, the additional assumption that the vacuum expectation value (vev) of the gauge kinetic function does not break the unifying gauge symmetry leads to a common mass $m_{1/2}$ for all gauginos. In addition, it is usually assumed that there exists a common mass m_0 for all scalars and a common trilinear term A_0 for all soft SUSY breaking trilinear interactions. Universal soft SUSY breaking scalar masses are not, however, a consequence of the supergravity framework [4] but an additional assumption.

The (universal) soft parameters are assumed to be renormalized at some high scale $M_X \sim M_{\rm GUT} - M_{\rm Planck}$. These are assumed to have values comparable to the weak scale, $M_{\rm weak}$, resulting in an elegant solution to the fine-tuning problem associated with the Higgs sector. Motivated by the apparently successful gauge coupling unification in the MSSM, the scale $M_{\rm GUT} \simeq 2 \times 10^{16}$ GeV is usually adopted for the scale choice M_X. The resulting effective theory, valid at energy scales $E < M_{\rm GUT}$, is then just the MSSM with soft SUSY breaking terms that unify at $M_{\rm GUT}$. The soft SUSY breaking scalar and gaugino masses, the trilinear A terms and in addition a bilinear soft term B, the gauge and Yukawa couplings and the supersymmetric μ term are all then evolved from $M_{\rm GUT}$ to some scale $M \simeq M_{\rm weak}$ using renormalization group equations (RGE). The large top quark Yukawa coupling causes the squared mass of H_u to be driven to negative values, which signals the radiative breakdown of electroweak symmetry (REWSB); this then allows one to determine the value of μ^2 in terms of M_Z^2, possibly at the expense of some fine-tuning. Finally, it is customary to trade the parameter B for $\tan\beta$, the ratio of Higgs field vacuum expectation values. The resulting weak-scale spectrum of superpartners and their couplings can then be derived in terms of four continuous parameters plus one sign

$$m_0, \qquad m_{1/2}\,, \qquad A_0\,, \qquad \tan\beta \quad \text{and} \quad \text{sign}(\mu)\,, \tag{1.1}$$

in addition to the usual parameters of the standard model. This calculational procedure has been embedded into the event generator ISAJET [5] thereby allowing detailed predictions for the collider events within this framework.

The mSUGRA model, while highly predictive, rests upon a number of simplifying assumptions that are invalid in specific models of physics at energy scales $\sim M_{\rm GUT} - M_{\rm Planck}$. Thus, in the search for weak-scale supersymmetry, the mSUGRA model may give misleading guidance as to the possible event signatures expected at high-energy collider experiments. Indeed the literature is replete with models with non-universal soft SUSY breaking mass terms at the high scales. In this paper, we survey a variety of these models (as well as others that lead to universality) and comment on possible phenomenological implications, especially for high-energy collider experiments. For the most part, we restrict our attention to models which reduce to the R-parity conserving MSSM at scales $Q < M_{\rm GUT}$.

The event generator ISAJET (versions > 7.37) has recently been upgraded [5] to accomodate supersymmetric models with non-universal soft SUSY breaking masses at the GUT scale. To generate such models, the user must input the usual mSUGRA parameter set eq. (1.1), but may in addition select one or several of the following options:

$$\begin{aligned}
&NUSUG1 : M_1\,,\ M_2\,,\ M_3\,,\\
&NUSUG2 : A_t\,,\ A_b\,,\ A_\tau\,,\\
&NUSUG3 : m_{H_d}\,,\ m_{H_u}\,,\\
&NUSUG4 : m_{Q_1}\,,\ m_{D_1}\,,\ m_{U_1}\,,\ m_{L_1}\,,\ m_{E_1}\,,\\
&NUSUG5 : m_{Q_3}\,,\ m_{D_3}\,,\ m_{U_3}\,,\ m_{L_3}\,,\ m_{E_3}\,.
\end{aligned}$$

If one or more of the NUSUG$_i$ ($i = 1 \div 5$) inputs are selected, then the GUT scale universal soft breaking masses are overwritten and a weak-scale MSSM mass spectrum is generated. ISAJET then computes the corresponding branching fractions and sparticle cross sections, so that specific theoretical predictions for GUT scale SSB masses can be mapped onto explicit predictions for the high-energy collider events expected to arise from these models. In addition, the ISAJET keyword SSBCSC has been introduced in ISAJET versions ≥ 7.50. Using SSBCSC, the user may choose any scale between the weak scale and the Planck scale at which to impose the above SSB boundary conditions. We illustrate its use in section 11.3 where it is necessary to introduce boundary conditions at the string scale rather than at $M_{\rm GUT}$.

To facilitate the examination of these models by our experimental colleagues, we present here a survey of a number of well-motivated models which usually lead to non universality of SSB parameters. Our survey is far from exhaustive, but is meant to present a flavor of the range of possibilities available for such models. For each model, we

1. present a short description of the physics,

2. delineate the parameter space,

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3. indicate how, within the model framework, collider events may be generated using ISAJET, and

4. comment upon some of the general features of SUSY events expected at collider experiments.

The models selected include the following:

- SU(5) grand unified models with universal soft SUSY breaking masses at scales higher than $Q = M_{\mathrm{GUT}}$,
- SU(5) models where supersymmetry breaking occurs via non-singlet hidden sector superfields,
- the MSSM plus an intermediate-scale right-handed neutrino which leads to see-saw neutrino masses,
- models with extra D-term contributions to scalar masses that are generically present if the rank of the unifying gauge group exceeds 4,
- minimal and general SO(10) grand unified models with universal soft SUSY breaking masses at scales higher than $Q = M_{\mathrm{GUT}}$,
- grand unified models with group structure $G_{\mathrm{GUT}} \times G_H$, where G_H contains a hypercolor interaction used to solve the doublet-triplet splitting problem,
- effective supersymmetry models which lead to multi-TeV range scalar masses for the first two generations, but sub-TeV masses for third generation scalars and gauginos,
- anomaly-mediated SUSY breaking models (AMSB), where the hidden sector resides in different space-time dimensions from the visible sector,
- the minimal gaugino mediation model,
- 4-dimensional string models with Calabi-Yao or orbifold compactifications, and
- models inspired by M-theory with SUSY breaking by one or several moduli fields.

Space limitations preclude us from detailed discussions of these models. Here, we sketch the physics behind each model, and provide the reader with selected references where further details may be found. While much, but by no means all, of the material presented may be found in the literature, our hope is that the form in which we have presented it will facilitate, or even spur, closer examination of alternatives to the mSUGRA and gauge-mediated SUSY breaking models.

2. SU(5) grand unified model with the SSB universality scale higher than $M_{\rm GUT}$

As a working assumption, the scale at which all the SSB parameters are generated, is usually taken to be $M_{\rm GUT}$. If this scale is substantially higher than this (but smaller than the Planck scale), renormalization group (RG) evolution induces a non universality at the GUT scale. The effect can be significant if large representations are present. Here, we assume that supersymmetric SU(5) grand unification is valid at mass scales $Q > M_{\rm GUT} \simeq 2\times 10^{16}$ GeV, extending at most to the reduced Planck scale $M_P \simeq 2.4\times 10^{18}$ GeV. Below $Q = M_{\rm GUT}$, the SU(5) model breaks down to the MSSM with the usual $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry. This framework is well described in, for instance, the work of Polonsky and Pomarol [6].

In the SU(5) model, the $\hat{D}^c$ and $\hat{L}$ superfields are elements of a $\bar{\mathbf{5}}$ superfield $\hat{\phi}$, while the $\hat{Q}$, $\hat{U}^c$ and $\hat{E}^c$ superfields occur in the $\mathbf{10}$ representation $\hat{\psi}$. The Higgs sector is comprised of three super-multiplets: $\hat{\Sigma}(\mathbf{24})$ which is responsible for breaking SU(5), plus $\hat{\mathcal{H}}_1(\bar{\mathbf{5}})$ and $\hat{\mathcal{H}}_2(\mathbf{5})$ which contain the usual Higgs doublet superfields $\hat{H}_d$ and $\hat{H}_u$, respectively, which occur in the MSSM. The superpotential is given by,

$$\hat{f} = \mu_\Sigma tr\hat{\Sigma}^2 + \frac{1}{6}\lambda' tr\hat{\Sigma}^3 + \mu_H \hat{\mathcal{H}}_1\hat{\mathcal{H}}_2 + \lambda\hat{\mathcal{H}}_1\hat{\Sigma}\hat{\mathcal{H}}_2 + \frac{1}{4}f_t\epsilon_{ijklm}\hat{\psi}^{ij}\hat{\psi}^{kl}\hat{\mathcal{H}}_2^m + \sqrt{2}f_b\hat{\psi}^{ij}\hat{\phi}_i\hat{\mathcal{H}}_{1j}\,, \tag{2.1}$$

where a sum over families is understood. f_t and f_b are the top and bottom quark Yukawa couplings, λ and λ' are GUT Higgs sector self couplings, and μ_Σ and μ_H are superpotential Higgs mass terms.

Supersymmetry breaking is parametrized by the soft supersymmetry breaking terms:

$$\begin{aligned}\mathcal{L}_{\rm soft} = & -m_{\mathcal{H}_1}^2|\mathcal{H}_1|^2 - m_{\mathcal{H}_2}^2|\mathcal{H}_2|^2 - m_\Sigma^2 tr\{\Sigma^\dagger\Sigma\} - m_5^2|\phi|^2 - m_{10}^2 tr\{\psi^\dagger\psi\} - \\ & -\frac{1}{2}M_5\bar{\lambda}_\alpha\lambda_\alpha + \Big[B_\Sigma\mu_\Sigma tr\Sigma^2 + \frac{1}{6}A_{\lambda'}\lambda' tr\Sigma^3 + B_H\mu_H\mathcal{H}_1\mathcal{H}_2 + A_\lambda\lambda\mathcal{H}_1\Sigma\mathcal{H}_2 + \\ & +\frac{1}{4}A_t f_t\epsilon_{ijklm}\psi^{ij}\psi^{kl}\mathcal{H}_2^m + \sqrt{2}A_b f_b\psi^{ij}\phi_i\mathcal{H}_{1j} + \text{h.c.}\Big]\,.\end{aligned} \tag{2.2}$$

The various soft masses and gauge and Yukawa couplings evolve with energy according to the 15 renormalization group equations given in ref. [44, appendix A]. Here, we modify them to correspond with the sign conventions in ISAJET [5]:

$$\begin{aligned}\frac{dm_{10}^2}{dt} = & \frac{1}{8\pi^2}\Big[3f_t^2(m_{\mathcal{H}_2}^2 + 2m_{10}^2 + A_t^2) + \\ & + 2f_b^2(m_{\mathcal{H}_1}^2 + m_{10}^2 + m_5^2 + A_b^2) - \frac{72}{5}g_G^2M_5^2\Big]\,,\end{aligned} \tag{2.3}$$

$$\frac{dm_5^2}{dt} = \frac{1}{8\pi^2}\Big[4f_b^2(m_{\mathcal{H}_1}^2 + m_{10}^2 + m_5^2 + A_b^2) - \frac{48}{5}g_G^2M_5^2\Big]\,, \tag{2.4}$$

$$\frac{dm^2_{\mathcal{H}_1}}{dt} = \frac{1}{8\pi^2}\Bigg[4f_b^2(m^2_{\mathcal{H}_1}+m^2_{10}+m^2_5+A_b^2)+ + \frac{24}{5}\lambda^2(m^2_{\mathcal{H}_1}+m^2_{\mathcal{H}_2}+m^2_\Sigma+A^2_\lambda)-\frac{48}{5}g_G^2M_5^2\Bigg]\,, \tag{2.5}$$

$$\frac{dm^2_{\mathcal{H}_2}}{dt} = \frac{1}{8\pi^2}\Bigg[3f_t^2(m^2_{\mathcal{H}_2}+2m^2_{10}+A_t^2)+ + \frac{24}{5}\lambda^2(m^2_{\mathcal{H}_1}+m^2_{\mathcal{H}_2}+m^2_\Sigma+A^2_\lambda)-\frac{48}{5}g_G^2M_5^2\Bigg]\,, \tag{2.6}$$

$$\frac{dm^2_{\Sigma}}{dt} = \frac{1}{8\pi^2}\Bigg[\frac{21}{20}\lambda'^2(3m^2_\Sigma+A^2_{\lambda'})+ + \lambda^2(m^2_{\mathcal{H}_1}+m^2_{\mathcal{H}_2}+m^2_\Sigma+A^2_\lambda)-20g_G^2M_5^2\Bigg]\,, \tag{2.7}$$

$$\frac{dA_t}{dt} = \frac{1}{8\pi^2}\Bigg[9A_tf_t^2+4A_bf_b^2+\frac{24}{5}A_\lambda\lambda^2+\frac{96}{5}g_G^2M_5\Bigg]\,, \tag{2.8}$$

$$\frac{dA_b}{dt} = \frac{1}{8\pi^2}\Bigg[10A_bf_b^2+3A_tf_t^2+\frac{24}{5}A_\lambda\lambda^2+\frac{84}{5}g_G^2M_5\Bigg]\,, \tag{2.9}$$

$$\frac{dA_\lambda}{dt} = \frac{1}{8\pi^2}\Bigg[\frac{21}{20}A_{\lambda'}\lambda'^2+3A_tf_t^2+4A_bf_b^2+\frac{53}{5}A_\lambda\lambda^2+\frac{98}{5}g_G^2M_5\Bigg]\,, \tag{2.10}$$

$$\frac{dA_{\lambda'}}{dt} = \frac{1}{8\pi^2}\Bigg[\frac{63}{20}A_{\lambda'}\lambda'^2+3A_\lambda\lambda^2+30g_G^2M_5\Bigg]\,, \tag{2.11}$$

$$\frac{df_t}{dt} = \frac{f_t}{16\pi^2}\Bigg[9f_t^2+4f_b^2+\frac{24}{5}\lambda^2-\frac{96}{5}g_G^2\Bigg]\,, \tag{2.12}$$

$$\frac{df_b}{dt} = \frac{f_b}{16\pi^2}\Bigg[10f_b^2+3f_t^2+\frac{24}{5}\lambda^2-\frac{84}{5}g_G^2\Bigg]\,, \tag{2.13}$$

$$\frac{d\lambda}{dt} = \frac{\lambda}{16\pi^2}\Bigg[\frac{21}{20}\lambda'^2+3f_t^2+4f_b^2+\frac{53}{5}\lambda^2-\frac{98}{5}g_G^2\Bigg]\,, \tag{2.14}$$

$$\frac{d\lambda'}{dt} = \frac{\lambda'}{16\pi^2}\Bigg[\frac{63}{20}\lambda'^2+3\lambda^2-30g_G^2\Bigg]\,, \tag{2.15}$$

$$\frac{d\alpha_G}{dt} = -3\frac{\alpha_G^2}{2\pi}\,, \tag{2.16}$$

$$\frac{dM_5}{dt} = -3\frac{\alpha_GM_5}{2\pi}\,, \tag{2.17}$$

with $t=\log Q$.

To generate the weak-scale MSSM mass spectrum, one begins with the input parameters

$$\alpha_{\rm GUT}\,,\ f_t\,,\ f_b\,,\ \lambda\,,\ \lambda' \tag{2.18}$$

stipulated at $Q=M_{\rm GUT}$, where $f_b=f_\tau$ is obtained from the corresponding mSUGRA model. The first three of these can be extracted, for instance, from ISASUGRA, versions ≥ 7.44. The couplings $\lambda(M_{\rm GUT})$ and $\lambda'(M_{\rm GUT})$ are additional inputs, where $\lambda(M_{\rm GUT})\gtrsim 0.7$ [7] to make the triplet Higgsinos heavy enough to satisfy experimental bounds on the proton lifetime. The gauge and Yukawa couplings can be evolved via the RGEs to determine their values at $Q = M_P$. Assuming universality at M_P (this maximizes the effects of non universality at the GUT scale), one imposes

$$m_{10} = m_5 = m_{\mathcal{H}_1} = m_{\mathcal{H}_2} = m_\Sigma \equiv m_0 \,,$$
$$A_t = A_b = A_\lambda = A'_\lambda \equiv A_0 \,, \tag{2.19}$$

and evolves all the soft masses from M_P to $M_{\rm GUT}$. The MSSM soft breaking masses at $M_{\rm GUT}$ are specified via

$$m_Q^2 = m_U^2 = m_E^2 \equiv m_{10}^2 \,, \qquad m_D^2 = m_L^2 \equiv m_5^2 \,,$$
$$m_{H_1}^2 = m_{\mathcal{H}_1}^2 \,, \qquad m_{H_2}^2 = m_{\mathcal{H}_2}^2 \,, \tag{2.20}$$

which can serve as input to ISAJET [5] via the NUSUG$_i$ keywords. Since there is no splitting amongst the gaugino masses, the gaugino masses may be taken to be $M_1 = M_2 = M_3 \equiv m_{1/2}$ where $m_{1/2}$ is stipulated most conveniently at the GUT scale.

To obtain correct Yukawa unification, it is crucial to start with the correct weak-scale Yukawa couplings. To calculate the values of the Yukawa couplings at scale $Q = M_Z$, one begins with the pole masses $m_b = 4.9\,\mathrm{GeV}$ and $m_\tau = 1.784\,\mathrm{GeV}$. One may calculate the corresponding running masses in the $\overline{MS}$ scheme, and evolve m_b and m_τ up to M_Z using two-loop SM RGEs. At $Q = M_Z$, the SUSY loop corrections to m_b and m_τ must be included; ISAJET versions > 7.44 uses the approximate formulae of Pierce et al. [8]. A similar procedure is used to calculate the top quark Yukawa coupling at scale $Q = m_t$. SUSY particle mass spectra consistent with constraints from collider searches and with unified b and τ Yukawa couplings (to 5%) are then obtained (assuming universality of scalar masses at the scale M_P), but only for $\mu < 0$ and $30 \lesssim \tan\beta \lesssim 50$, where the allowed range is weakly sensitive to α_s.

Parameter	Scale	Value
m_0	M_P	150
$m_{1/2}$	$M_{\rm GUT}$	200
A_0	M_P	0
$\tan\beta$	$M_{\rm weak}$	35
μ	$M_{\rm weak}$	< 0
$g_{\rm GUT}$	$M_{\rm GUT}$	0.717
f_t	$M_{\rm GUT}$	0.534
$f_b = f_\tau$	$M_{\rm GUT}$	0.271
λ	$M_{\rm GUT}$	1
λ'	$M_{\rm GUT}$	0.1
$m_{10}^{1,2}$	$M_{\rm GUT}$	194.4
$m_5^{1,2}$	$M_{\rm GUT}$	180.8
m_{10}^3	$M_{\rm GUT}$	183.8
m_5^3	$M_{\rm GUT}$	1777.7
m_{H_d}	$M_{\rm GUT}$	107.8
m_{H_u}	$M_{\rm GUT}$	96.2
A_t	$M_{\rm GUT}$	-87.6
$A_b = A_\tau$	$M_{\rm GUT}$	-77.6

Table 1: Input and output parameters for an SU(5) case study.

To illustrate the extent of non universality due to SU(5) running of SSB masses between M_P and $M_{\rm GUT}$, we explicitly examine a typical case. The corresponding input parameters as well as the values of SSB parameters at $M_{\rm GUT}$ are listed in table 1. The GUT scale input parameters extracted from ISAJET for $\tan\beta = 35$ are $f_t = 0.534$ and $f_b = f_\tau = 0.271$ for the top, bottom and tau Yukawa couplings. We also adopt $\lambda = 1.0$ and $\lambda' = 0.1$ for the SU(5) Higgs couplings and $g_{\rm GUT} = 0.717$ for the unified SU(5) gauge coupling. At the Planck scale, we then take $m_0 = 150\,\mathrm{GeV}$ and $A_0 = 0\,\mathrm{GeV}$, parameters that are analogous to m_0 and A_0 at the GUT scale in the mSUGRA model. We take $m_{1/2}(M_{\rm GUT}) = 200\,\mathrm{GeV}$ for the universal gaugino masses.

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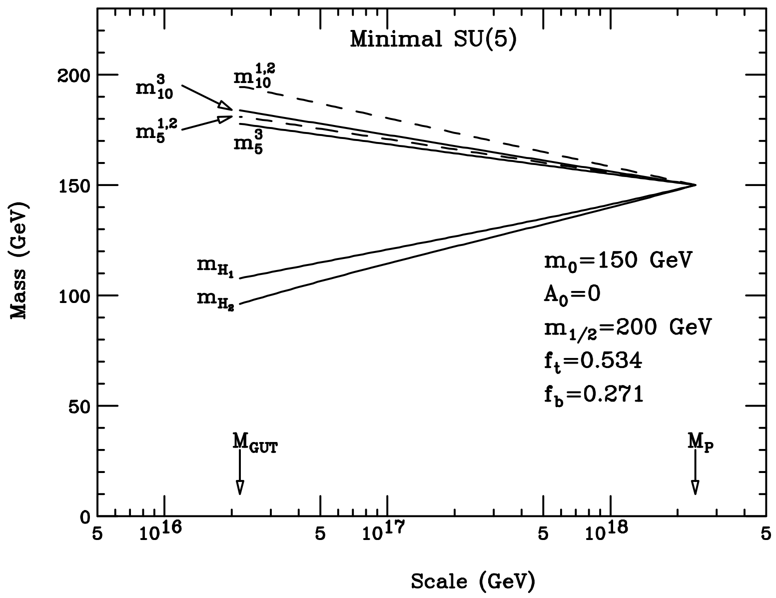

Figure 1: Running of the soft SUSY breaking masses between the Planck scale and the GUT scale in the minimal SU(5) model for $\tan\beta = 35$. At the GUT scale we have taken $\lambda = 1$ and $\lambda' = 0.1$ for the Higgs couplings, and $\alpha_{\rm GUT} = 0.041$ for the unified gauge coupling.

The evolution of SUSY mass parameters in the minimal SU(5) model between M_P and $M_{\rm GUT}$ is shown in figure 1, assuming universality at M_P. We see that the rather high value of λ induces a large splitting $m_5^2 \simeq m_{10}^2 > m_{\mathcal{H}_1}^2$, $m_{\mathcal{H}_2}^2$. Likewise, the large value of f_t is responsible for the splitting $m_{\mathcal{H}_1}^2 > m_{\mathcal{H}_2}^2$ at $M_{\rm GUT}$. The large t and b Yukawa couplings are also responsible for the split between third generation and the first two generation values of m_{10} and m_5. It is interesting to notice that reasonable values of the free parameters can give $\sim 100\%$ deviations from universality at $M_{\rm GUT}$. In the cases that we checked, it was typically the Higgs scalars that are split by the large amount from the other scalars, primarily because λ is large: for acceptable solutions, the corresponding non universality between matter scalar masses was typically $\sim 10 \div 20\%$. In table 2, we list the corresponding values of selected weak-scale sparticle masses for both the SU(5) case and mSUGRA. The shift in scalar masses in this case can be up to $\sim 20\%$, with the biggest shift occuring in the $\tilde{\ell}_R$ and $\tilde{\tau}_1$ masses.

Parameter	mSUGRA	SU(5)
$m_{\tilde{g}}$	512.0	515.0
$m_{\tilde{u}_L}$	468.0	484.0
$m_{\tilde{d}_R}$	454.7	463.2
$m_{\tilde{t}_1}$	335.6	337.8
$m_{\tilde{b}_1}$	375.4	375.2
$m_{\tilde{\ell}_L}$	212.8	235.4
$m_{\tilde{\ell}_R}$	174.5	213.8
$m_{\tilde{\tau}_1}$	124.3	151.1
$m_{\widetilde{W}_1}$	150.1	155.3
$m_{\widetilde{Z}_2}$	150.3	155.3
$m_{\widetilde{Z}_1}$	80.8	81.5
m_h	111.0	111.6
m_A	210.4	216.4
μ	−263.8	−304.3

Table 2: Weak-scale sparticle masses and parameters (GeV) for mSUGRA and for an SU(5) case study.

3. SU(5) models with non-universal gaugino masses

Since supergravity is not a renormalizable theory, in general we may expect a non-trivial gauge kinetic function, and hence the possiblity of non-vanishing gaugino masses if SUSY is broken. Expanding the gauge kinetic function as $f_{ab} = \delta_{ab} + \hat{\Phi}_{ab}/M_{\text{Planck}} + \cdots$, where the fields $\hat{\Phi}_{ab}$ transform as left-handed chiral superfields under supersymmetry transformations, and as the symmetric product of two adjoints under gauge symmetries, we parametrize the lowest order contribution to gaugino masses by,

$$\mathcal{L} \supset \int d^2\theta \hat{f}^a \hat{f}^b \frac{\hat{\Phi}_{ab}}{M_{\text{Planck}}} + \text{h.c.} \supset \frac{\langle F_\Phi \rangle_{ab}}{M_{\text{Planck}}} \lambda^a \lambda^b + \cdots , \tag{3.1}$$

where the λ^a are the gaugino fields, and F_Φ is the auxillary field component of $\hat{\Phi}$ that acquires a SUSY breaking vev .

If the fields F_Φ which break supersymmetry are gauge singlets, universal gaugino masses result. However, in principle, the chiral superfield which communicates supersymmetry breaking to the gaugino fields can lie in any representation in the symmetric product of two adjoints, and so can lead to gaugino mass terms[2] that (spontaneously) break the underlying gauge symmetry. We require, of course, that SM gauge symmetry is preserved. Non-universal gaugino masses have been previously considered by other authors [9, 10, 11, 12].

In the context of SU(5) grand unification, F_Φ belongs to an SU(5) irreducible representation which appears in the symmetric product of two adjoints:

$$(\mathbf{24}\times\mathbf{24})_{\text{symmetric}} = \mathbf{1} \oplus \mathbf{24} \oplus \mathbf{75} \oplus \mathbf{200} , \tag{3.2}$$

where only **1** yields universal masses. The relations amongst the various GUT scale gaugino masses have been worked out e.g. in ref. [12]. The relative GUT scale SU(3), SU(2) and U(1) gaugino masses M_3, M_2 and M_1 are listed in table 3 along with the approximate masses after RGE evolution to $Q \sim M_Z$. Here, motivated by the measured values of the gauge couplings at LEP, we assume that the vev of the SUSY-preserving scalar component of $\hat{\Phi}$ is neglible. Each of the three non-singlet models is as predictive as the canonical singlet case, and all are compatible with the unification of gauge couplings. These scenarios represent the predictive subset of the more general (and less predictive) case of an arbitrary superposition of these representations. The model parameters may be chosen to be,

$$m_0\,,\ M_3^0\,,\ A_0\,,\ \tan\beta \text{ and } \operatorname{sign}(\mu)\,, \tag{3.3}$$

where M_i^0 is the SU(i) gaugino mass at scale $Q = M_{\text{GUT}}$. M_2^0 and M_1^0 can then be calculated in terms of M_3^0 according to table 3. Sample spectra for each case are exhibited in table 4.

[2]The results of this section are not new, but in the interest of completeness we thought it fit to include a review of these models in this section.

F_Φ	$M_{\rm GUT}$ M_3	M_2	M_1	M_Z M_3	M_2	M_1
1	1	1	1	~ 6	~ 2	~ 1
24	2	-3	-1	~ 12	~ -6	~ -1
75	1	3	-5	~ 6	~ 6	~ -5
200	1	2	10	~ 6	~ 4	~ 10

Table 3: Relative gaugino masses at $M_{\rm GUT}$ and M_Z in the four possible F_Φ irreducible representations.

Parameter	$\Phi(\mathbf{1})$	$\Phi(\mathbf{24})$	$\Phi(\mathbf{75})$	$\Phi(\mathbf{200})$
$m_{\tilde g}$	394.9	397.7	409.3	404.0
$m_{\tilde u_L}$	356.8	372.0	457.9	406.2
$m_{\tilde t_1}$	243.7	283.8	255.2	295.0
$m_{\tilde b_1}$	328.8	342.0	356.5	340.2
$m_{\tilde\ell_L}$	154.4	191.1	360.1	372.7
$m_{\tilde\ell_R}$	123.2	112.3	310.4	589.8
$m_{\tilde\tau_1}$	120.9	111.6	309.8	372.1
$m_{\widetilde W_1}$	95.6	147.0	93.2	156.3
$m_{\widetilde Z_2}$	99.7	142.5	106.0	202.5
$m_{\widetilde Z_1}$	53.2	33.1	92.6	151.9
m_h	103.4	99.8	104.5	106.7
m_A	257.2	249.6	372.9	421.3
μ	215.2	173.0	104.7	197.2

Table 4: Weak-scale sparticle masses and parameters (GeV) for the four cases of singlet and non-singlet hidden sector vevs in SU(5). For each case, we take $(m_0,\ M_3^0,\ A_0) = (100,\ 150,\ 0\,{\rm GeV})$, with $\tan\beta = 5$ and $\mu > 0$.

The phenomenology of these models has recently been examined in ref. [13], and the SUSY reach presented for Fermilab Tevatron upgrade options for a variety of discovery channels. The results were found to be model-dependent. In particular, in the **24** model, a large splitting between weak-scale values of $m_{\widetilde Z_2}$, $m_{\widetilde W_1}$ and $m_{\widetilde Z_1}$ gave rise to large rates for events with isolated leptons, so that SUSY discovery should be easier in this case than in the mSUGRA model. A special feature of this model is the sizeable cross section for $(Z \to \ell\bar\ell) + jets + \not\!\!E_T$ events. Indeed, for certain ranges of model parameters, SUSY discovery seemed to be possible only via this channel. In contrast, for the **75** and **200** models, $m_{\widetilde Z_2}$, $m_{\widetilde W_1}$ and $m_{\widetilde Z_1}$ were all nearly degenerate, so that leptons arising from -ino decays were very soft and difficult to detect. Consequently, there was hardly any reach for SUSY in these models at the Tevatron via leptonic channels, and the best reach occurred typically in the $\not\!\!E_T + jets$ channels.

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4. The MSSM with a right-handed neutrino

Experimental evidence [14] strongly indicates the existence of neutrino oscillations, and almost certainly neutrino mass. The favoured interpretation is $\nu_\mu - \nu_\tau$ oscillations, with $\Delta m^2 \sim 10^{-2}$ eV^2 and near-maximal mixing. An attractive method for introducing neutrino mass into the MSSM is via the see-saw mechanism [15]. In this case, one can introduce an additional chiral superfield[3] ($\hat{N}^c$) which transforms as a gauge singlet (whose fermionic component is the left-handed anti-neutrino and scalar component is $\tilde{\nu}_R^\dagger$). A Majorana mass term for the right-handed neutrino is allowed and, because ν_R is a SM singlet, its mass may be large: $M_N \sim 10^{10} \div 10^{16}$ GeV. When electroweak symmetry is broken, a Dirac neutrino mass $m_D \sim m_\ell$ is also induced via the usual Higgs mechanism. The resulting neutrino mass matrix must be diagonalized, and one obtains a light physical neutrino mass $m_\nu \simeq m_D^2/M_N$ plus a dominantly singlet neutrino of mass $M \simeq M_N$.

The superpotential for the MSSM with a singlet neutrino superfield $\hat{N}^c$ (for just a single generation), is given by

$$\hat{f} = \hat{f}_{MSSM} + f_\nu \epsilon_{ij} \hat{L}^i \hat{H}_u^j \hat{N}^c + \frac{1}{2} M_N \hat{N}^c \hat{N}^c \,, \tag{4.1}$$

while the soft SUSY breaking terms now include

$$\mathcal{L} = \mathcal{L}_{MSSM} - m_{\tilde{\nu}_R}^2 |\tilde{\nu}_R|^2 + \left[A_\nu f_\nu \epsilon_{ij} \tilde{L}^i \tilde{H}_u^j \tilde{\nu}_R^\dagger + \frac{1}{2} B_\nu M_N \tilde{\nu}_R^2 + \text{h.c.} \right]. \tag{4.2}$$

The parameters A_ν, B_ν and $m_{\tilde{\nu}_R}$ are assumed to be comparable to the weak scale.

Many of the relevant RGEs have been presented in ref. [16]. Here we present the complete set needed for determining the sparticle spectrum at the weak scale. The one-loop RGEs for the gauge couplings and gaugino masses are unchanged from the MSSM case, since the $\hat{N}^c$ superfield is a gauge singlet. The Yukawa coupling RGEs are

$$\frac{df_t}{dt} = \frac{f_t}{16\pi^2} \left[6f_t^2 + f_b^2 + f_\nu^2 - \frac{16}{3} g_3^2 - 3g_2^2 - \frac{13}{15} g_1^2 \right], \tag{4.3}$$

$$\frac{df_b}{dt} = \frac{f_b}{16\pi^2} \left[f_t^2 + 6f_b^2 + f_\tau^2 - \frac{16}{3} g_3^2 - 3g_2^2 - \frac{7}{15} g_1^2 \right], \tag{4.4}$$

$$\frac{df_\tau}{dt} = \frac{f_\tau}{16\pi^2} \left[3f_b^2 + 4f_\tau^2 + f_\nu^2 - 3g_2^2 - \frac{9}{5} g_1^2 \right], \tag{4.5}$$

$$\frac{df_\nu}{dt} = \frac{f_\nu}{16\pi^2} \left[3f_t^2 + f_\tau^2 + 4f_\nu^2 - 3g_2^2 - \frac{3}{5} g_1^2 \right]. \tag{4.6}$$

[3] Our purpose here is to illustrate the effect of introducing singlet neutrino superfields on the SSB parameters and the SUSY spectrum. An explanation of the atmospheric neutrino data would, of course, require us to introduce more than one such superfield and also interactions that violate lepton flavour conservation, but as long as these have only small Yukawa couplings, their effect on the spectrum should be negligible.

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The RGEs for m_Q^2, m_U^2, m_D^2, m_E^2 and $m_{H_d}^2$ are all unchanged from the MSSM. However, for m_L^2, $m_{\tilde{\nu}_R}^2$ and $m_{H_u}^2$, we have

$$\frac{dm_L^2}{dt} = \frac{2}{16\pi^2}\left[-\frac{3}{5}g_1^2M_1^2 - 3g_2^2M_2^2 + f_\tau^2X_\tau + f_\nu^2X_n\right], \tag{4.7}$$

$$\frac{dm_{\tilde{\nu}_R}^2}{dt} = \frac{4}{16\pi^2}\left[f_\nu^2X_n\right], \tag{4.8}$$

$$\frac{dm_{H_u}^2}{dt} = \frac{2}{16\pi^2}\left[-\frac{3}{5}g_1^2M_1^2 - 3g_2^2M_2^2 + 3f_t^2X_t + f_\nu^2X_n\right], \tag{4.9}$$

where we have defined $X_n = m_L^2 + m_{\tilde{\nu}_R}^2 + m_{H_u}^2 + A_\nu^2$ and X_t and X_τ are given in ref. [17]. Finally, the RGEs for the A_i-parameters are given by

$$\frac{dA_t}{dt} = \frac{2}{16\pi^2}\left[\Sigma c_i g_i^2 M_i + 6f_t^2A_t + f_b^2A_b + f_\nu^2A_\nu\right], \tag{4.10}$$

$$\frac{dA_b}{dt} = \frac{2}{16\pi^2}\left[\Sigma c_i' g_i^2 M_i + 6f_b^2A_b + f_t^2A_t + f_\tau^2A_\tau\right], \tag{4.11}$$

$$\frac{dA_\tau}{dt} = \frac{2}{16\pi^2}\left[\Sigma c_i'' g_i^2 M_i + 3f_b^2A_b + 4f_\tau^2A_\tau + f_\nu^2A_\nu\right], \tag{4.12}$$

$$\frac{dA_\nu}{dt} = \frac{2}{16\pi^2}\left[\Sigma c_i''' g_i^2 M_i + 3f_t^2A_t + 4f_\nu^2A_\nu + f_\tau^2A_\tau\right], \tag{4.13}$$

where the c_i, c_i' and c_i'' are given in ref. [17], and $c_i''' = \{(3/5), 3, 0\}$. These RGEs apply for scales $Q > M_N$, while the MSSM RGEs are used below $Q = M_N$. Below the scale M_N the effective theory does not contain the right-handed neutrino or sneutrino, so that the running of the corresponding parameters is frozen at their values at this scale. The RGE for the parameter B_ν is irrelevant for our analysis.

This model has been explicitly included in ISAJET version $\geq$ 7.48, via the keyword SUGRHN, which allows, in addition to mSUGRA and/or NUSUG$_i$ inputs, the following:

$$m_{\nu_\tau}\,,\ M_N\,,\ m_{\tilde{\nu}_{\tau R}}\,,\ A_\nu\,, \tag{4.14}$$

where all masses are entered in GeV units. Then the neutrino Yukawa coupling is calculated, and the MSSM + RHN RGEs are used at scales $Q > M_N$, while MSSM RGEs are used below $Q = M_N$.

A sample spectrum of masses is shown in table 5, assuming $m_{\nu_\tau} = 10^{-9}$ GeV, $M_N = 10^{13}$ GeV, $m_{\tilde{\nu}_{\tau R}} = 200$ GeV and $A_\nu = 0$. The main effect is that the additional Yukawa coupling drives the third generation slepton masses to somewhat lower values than the massless neutrino case.

An upper limit on the parameter $\tan\beta$ occurs in mSUGRA for $\mu < 0$ due to a breakdown in the REWSB mechanism, where the H_u mass is not driven sufficiently negative by RG running. For the MSSM + RHN model, the additional Yukawa coupling f_ν aids somewhat in driving $m_{H_u}^2$ negative. It is natural to ask how much

Parameter	mSUGRA	MSSM + RHN
m_{ν_τ}	0	10^{-9}
M_N	—	10^{13}
$m_{\tilde{\nu}_{\tau R}}$	—	200
A_ν	—	0
$m_{\tilde{g}}$	511.5	511.5
$m_{\tilde{u}_L}$	485.1	485.1
$m_{\tilde{t}_1}$	343.1	344.2
$m_{\tilde{b}_1}$	386.6	386.4
$m_{\tilde{\ell}_L}$	250.4	250.4
$m_{\tilde{\ell}_R}$	218.9	218.9
$m_{\tilde{\tau}_1}$	144.2	140.5
$m_{\tilde{\tau}_2}$	257.1	252.6
$m_{\tilde{\nu}_\tau}$	220.1	211.6
$m_{\widetilde{W}_1}$	146.9	147.6
$m_{\widetilde{Z}_2}$	147.5	148.2
$m_{\widetilde{Z}_1}$	79.9	80.0
m_h	111.7	111.7
m_A	243.3	249.4
μ	263.2	267.6
A_τ	-66.6	-66.5
A_t	-383.3	-383.3

Table 5: Model parameters and weak-scale sparticle masses in GeV for mSUGRA and for the MSSM + right-handed neutrino model. For each case, we take $m_0 = 200\,\text{GeV}$, $m_{1/2} = 200\,\text{GeV}$, $A_0 = 0$, $\tan\beta = 40$ and $\mu > 0$.

the additional Yukawa coupling f_ν would help to increase the allowed range of $\tan\beta$ while still satisfying the REWSB constraint. As an example, we checked that for the case $m_0 = m_{1/2} = 200\,\text{GeV}$, $A_0 = 0$, and $\mu < 0$, for which $\tan\beta \leq 45$ in the mSUGRA framework, the inclusion of a right-handed neutrino with $m_N = 10^{13}$, (10^{10}) $((10^7))$ GeV, only increases this range to 45.3 (45.7) ((46)), assuming $f_\nu = f_t$ at the GUT scale.

5. Unifying gauge groups with rank ≥ 5: D-terms

In general, if the MSSM is embedded in a GUT gauge group with rank ≥ 5, and the GUT gauge group is spontaneously broken to a gauge group of lower rank, there are additional D-term contributions to scalar masses. The important thing is that these contributions [18] affect TeV scale physics even if the scale at which the GUT symmetry is broken is very large: since symmetry breaking is arranged to occur in a nearly D-flat direction, these D-term contributions to scalar masses are still of order the weak scale, even though the extra particles have masses $\sim M_{\rm GUT}$. The D-terms must be added to the various SUSY scalar mass squared parameters at the high scale at which the breaking occurs, so that these effectively lead to non-universal boundary conditions for scalar masses.

Kolda and Martin [19] have analysed these contributions for gauge groups which are subgroups of E_6, which encompasses a wide range of well-motivated GUT group choices. E_6 contains in addition to the SM $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge symmetry two additional $\mathrm{U}(1)$ symmetries labelled as $\mathrm{U}(1)_X$ and $\mathrm{U}(1)_S$. The D-term contributions to scalar masses can then be parametrized as,

$$\begin{aligned}
\Delta m_Q^2 &= \frac{1}{6}D_Y - \frac{1}{3}D_X - \frac{1}{3}D_S\,, \\
\Delta m_D^2 &= \frac{1}{3}D_Y + D_X - \frac{2}{3}D_S\,, \\
\Delta m_U^2 &= -\frac{2}{3}D_Y - \frac{1}{3}D_X - \frac{1}{3}D_S\,, \\
\Delta m_E^2 &= D_Y - \frac{1}{3}D_X - \frac{1}{3}D_S\,, \\
\Delta m_{H_d}^2 &= -\frac{1}{2}D_Y - \frac{2}{3}D_X + D_S\,, \\
\Delta m_{H_u}^2 &= \frac{1}{2}D_Y + \frac{2}{3}D_X + \frac{2}{3}D_S\,,
\end{aligned} \tag{5.1}$$

where D_Y is the usual D-term associated with weak hypercharge breaking. In light of our ignorance of the mechanism of gauge symmetry breaking, the contributions D_X and D_S can be treated as additional dimensionful parameters, that can range over positive as well as negative values.

5.1 Minimal $\mathrm{SO}(10)$ model with gauge symmetry breaking at $Q = M_{\rm GUT}$

A simple special case of the above arises if the GUT gauge group $\mathrm{SO}(10)$ is assumed to directly break to the SM gauge group at $Q = M_{\rm GUT}$ so the theory below this scale is the MSSM, possibly together with a right-handed neutrino and sneutrino. In this case, the three generations of matter superfields plus an additional SM gauge singlet right-handed neutrino superfield for each generation are each elements of the 16 dimensional spinor representation of $\mathrm{SO}(10)$, and so are taken to have a common mass m_{16} above $M_{\rm GUT}$. The Higgs superfields of the MSSM belong to a single 10 dimensional fundamental representation of $\mathrm{SO}(10)$, and acquire a mass m_{10}. At $Q = M_{\rm GUT}$, the gauge symmmetry breaking induces

$$D_X \neq 0\,; \qquad D_Y = D_S = 0 \tag{5.2}$$

so that at this scale the scalar masses are broken according to (5.1). The MSSM masses at $M_{\rm GUT}$ may then be written as

$$\begin{aligned} m_Q^2 = m_E^2 = m_U^2 &= m_{16}^2 + M_D^2 \\ m_D^2 = m_L^2 &= m_{16}^2 - 3M_D^2 \\ m_{H_{u,d}}^2 &= m_{10}^2 \mp 2M_D^2\,, \end{aligned} \tag{5.3}$$

where we have reparametrized $D_X = -3M_D^2$. If the right-handed neutrino mass is substantially below the GUT scale, the soft breaking sneutrino mass would evolve as in eq. (4.8); at the GUT scale it would then be given by,

$$m_{\tilde{\nu}_R}^2 = m_{16}^2 + 5M_D^2\,. \tag{5.4}$$

In *minimal* SO(10), the superpotential above $M_{\rm GUT}$ has the form,

$$\hat{f} = f\hat{\psi}\hat{\psi}\hat{\phi} + \cdots \tag{5.5}$$

with just a single Yukawa coupling per generation, where $\hat{\psi}$ and $\hat{\phi}$ are the **16** dimensional spinor and **10** dimensional Higgs superfields, respectively. We neglect possible inter-generational mixing and also assume that the right-handed neutrino has a mass $\sim M_{\rm GUT}$. The dots represent terms including for instance higher dimensional Higgs representations and interactions responsible for the breaking of SO(10). We assume here for simplicity that these couplings are suppressed relative to the usual Yukawa couplings.

In minimal SO(10), all the Yukawa couplings are unified above $M_{\rm GUT}$, which forces one into a region of very large $\tan\beta \sim 50$ which is actually excluded assuming universality of scalars if the constraint of radiative electroweak symmetry breaking is included. It has been suggested [20], and recently shown [21], that D-term contributions have the correct form to allow for Yukawa unified solutions to the SUSY particle mass spectrum consistent with radiative electroweak symmetry breaking.

The parameter space of this model can be taken as

$$m_{16}\,,\ m_{10}\,,\ M_D^2\,,\ m_{1/2}\,,\ A_0\,,\ {\rm sign}(\mu)\,, \tag{5.6}$$

where M_D^2 can be either positive or negative. Yukawa coupling unification forces $\tan\beta$ to be in the range $45 \div 52$ — for many purposes its exact value is irrelevant.

The parameter space of minimal SO(10) SUSY GUT models was explored in ref. [21]. It was found that, requiring Yukawa unification good to 5%, *no* solutions could be found for values of $\mu > 0$, while many solutions could be obtained for $\mu < 0$, but only for *positive* values of M_D^2. The D-term forces $m_{H_u} < m_{H_d}$ at $Q = M_{\rm GUT}$: this is necessary to drive $m_{H_u}^2$ negative before $m_{H_d}^2$, as is required for REWSB with $\tan\beta > 1$. Implications of this model for the dark matter relic density, $b \to s\gamma$ decay rate, and collider searches, are presented in ref. [22]

A sample spectrum from the mSUGRA model and a corresponding case in Yukawa-unified SO(10) are shown in table 6. The D-term splitting that ameliorates REWSB also leaves a distinct imprint on the masses of the matter scalars: the left-sleptons and right-down-type squarks have smaller GUT scale squared masses than their counterparts. This can be reflected in the weak-scale spectrum where left-sleptons can be lighter than right-sleptons, and the right bottom squark can be by far the lightest of all the squarks — perhaps, even within the kinematic reach of the Main Injector upgrade of the Tevatron, though its detection may be complicated. Note also the smaller absolute value of the μ parameter in the SO(10) case: this results in lighter charginos and neutralinos with substantial, or even dominant, higgsino components and a smaller $\widetilde{Z}_2 - \widetilde{Z}_1$ mass difference. Finally, we remark that for the case shown, the lighter $\tilde{\tau}$ is dominantly $\tilde{\tau}_L$.

It is well known [23] that SUSY models with $\mu < 0$ and large $\tan\beta$ yield a large rate for the decay $b \to s\gamma$. Indeed [22], this class of models is already severely constrained by experimental results on radiative b-decays. However, additional non universality between generations is possible in this framework, which could alter the gluino-loop contributions, and hence the final branching fraction for $b \to s\gamma$ decay.

Parameter	mSUGRA	SO(10)
m_{16}	1022.0	1022.0
m_{10}	1022.0	1315.0
M_D	0.0	329.8
$m_{1/2}$	232.0	232.0
A_0	−1350.0	−1350.0
$\tan\beta$	45.0	48.6
$m_{\tilde{g}}$	639.0	631.5
$m_{\tilde{u}_L}$	1130.6	1178.5
$m_{\tilde{d}_R}$	1121.7	970.1
$m_{\tilde{t}_1}$	553.0	512.3
$m_{\tilde{b}_1}$	657.1	187.1
$m_{\tilde{\ell}_L}$	1035.8	857.8
$m_{\tilde{\ell}_R}$	1026.8	1088.9
$m_{\tilde{\nu}_e}$	1032.8	854.1
$m_{\tilde{\tau}_1}$	725.4	623.6
$m_{\tilde{\nu}_\tau}$	897.8	619.5
$m_{\widetilde{W}_1}$	193.5	122.9
$m_{\widetilde{Z}_2}$	193.3	131.6
$m_{\widetilde{Z}_1}$	97.4	84.0
m_h	88.5	118.8
m_A	90.4	479.9
m_{H^+}	131.2	490.2
μ	−547.8	−150.5
$\langle\tilde{\tau}_1\|\tilde{\tau}_L\rangle$	0.14	0.99

Table 6: Model parameters and weak-scale sparticle masses for mSUGRA and Yukawa unified SO(10) with D-terms and GUT symmetry breaking at $M_{\rm GUT}$. Note the mSUGRA case has a somewhat smaller value of $\tan\beta$ than SO(10), since no mSUGRA solution could be obtained with $\tan\beta = 48.6$.

6. Mass splittings in SO(10) above $Q = M_{\rm GUT}$

6.1 Minimal SO(10)

As discussed above, the minimal SO(10) model contains three generations of matter superfields each in a **16** dimensional representation, and a *single* Higgs superfield in the **10** dimensional representation. The superpotential is as given in eq. (5.5) with f the common Yukawa coupling for the third generation. Other terms will also be present, including Yukawa couplings for the first two generations, as well as more complicated Higgs representations necessary for SO(10) breaking. We will assume the Yukawa couplings involving these fields are all small, so the dominant contribution to RGE running comes from just the superpotential (5.5). We also assume associated SO(10) soft SUSY breaking parameters: m_{16}, m_{10}, $m_{1/2}$ and A. Then the RGEs in the minimal SO(10) model are calculable. For the gauge coupling we have,

$$\frac{dg}{dt} = \frac{g^3}{16\pi^2}(S-24)\,, \tag{6.1}$$

where S is the sum of Dynkin indices of the various chiral superfields in the model. With the above minimal field content, $S = 7$. However, additional fields associated for instance with SO(10) breaking ought to be present, and will increase the value of S. The Yukawa coupling RGE is,

$$\frac{df}{dt} = \frac{1}{16\pi^2} f\left(14f^2 - \frac{63}{2}g^2\right). \tag{6.2}$$

For the gaugino mass we have the following RGE:

$$\frac{dm_{1/2}}{dt} = \frac{1}{16\pi^2} 2(S-24)g^2 m_{1/2}\,. \tag{6.3}$$

For the scalar masses we have:

$$\frac{dm_{16}^2}{dt} = \frac{1}{16\pi^2}\left[10f^2\left(2m_{16}^2 + m_{10}^2 + A^2\right) - 45g^2 m_{1/2}^2\right], \tag{6.4}$$

$$\frac{dm_{10}^2}{dt} = \frac{1}{16\pi^2}\left[8f^2\left(2m_{16}^2 + m_{10}^2 + A^2\right) - 36g^2 m_{1/2}^2\right]. \tag{6.5}$$

Finally, the RGE for the trilinear mass parameter is

$$\frac{dA}{dt} = \frac{1}{16\pi^2}(28f^2 A + 63g^2 m_{1/2})\,. \tag{6.6}$$

As an illustration, we adopt the minimal SO(10) case 5 spectra from ref. [21] for which Yukawa couplings unify at $M_{\rm GUT}$. The model parameters and mass spectrum is listed in the "$M_{\rm GUT}$ unification" column of table 7. We begin by using $f(M_{\rm GUT}) = 0.553$ and $g_{\rm GUT} = 0.706$ (as given by the minimal SO(10) model). We then evolve (using $S = 7$) from $M_{\rm GUT}$ to M_P to find the corresponding Planck scale gauge and Yukawa couplings. At M_P, we assume universality of the three generations with $m_{16} = 629.8\,\text{GeV}$, while $m_{10} = 836.2\,\text{GeV}$. At $M_{\rm GUT}$, we take $m_{1/2} = 348.8\,\text{GeV}$ and $A = -186.5$ GeV, with a D-term $M_D = 135.6\,\text{GeV}$. A Yukawa unified solution is obtained for $\tan\beta = 52.1$ and the corresponding spectrum is shown in the last column titled "M_P unification".

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Parameter	$M_{\rm GUT}$ unification	M_P unification
m_{Q_1}	644.2	677.2
m_{L_1}	584.4	621.0
m_{Q_3}	644.2	603.5
m_{L_3}	584.4	539.2
m_{H_d}	857.9	833.9
m_{H_u}	813.9	788.6
$m_{\tilde g}$	813.9	838.1
$m_{\tilde u_L}$	974.4	969.9
$m_{\tilde d_R}$	910.8	914.3
$m_{\tilde t_1}$	618.7	586.2
$m_{\tilde b_1}$	636.8	600.1
$m_{\tilde \ell_L}$	634.6	660.6
$m_{\tilde \ell_R}$	662.5	692.7
$m_{\tilde\nu_e}$	629.5	655.8
$m_{\tilde\tau_1}$	427.8	397.8
$m_{\tilde\nu_\tau}$	519.1	474.0
$m_{\widetilde W_1}$	106.3	130.0
$m_{\widetilde Z_2}$	126.1	153.9
$m_{\widetilde Z_1}$	87.5	105.2
m_h	93.7	104.7
m_A	93.9	105.2
m_{H^+}	137.1	144.9
μ	-113.9	-142.4

Table 7: Model parameters and weak-scale sparticle masses for Yukawa unified SO(10) with D-terms. The first model has universality of matter scalars at $M_{\rm GUT}$, while the second has universality at M_P. For both cases, we take $m_{16} = 629.8\,{\rm GeV}$ and $m_{10} = 836.2\,{\rm GeV}$. At $M_{\rm GUT}$, for both cases, we take $m_{1/2} = 348.8\,{\rm GeV}$, $A = -186.5\,{\rm GeV}$. We also take $\mu < 0$, $\tan\beta = 52.1$ and $M_D = 135.6$, where D-terms are imposed at $M_{\rm GUT}$.

In figure 2, we show by the solid lines the effect of running of SSB parameters between M_P and $M_{\rm GUT}$ for the minimal SO(10) model, for parameters as in table 7. The dashed lines show the corresponding situation for $S = 15$, i.e. *with one additional adjoint included*; in this case, the running of the gauge coupling between $M_{\rm GUT}$ and M_P (see eq. (6.1)) is somewhat slower. We see that the splitting δm_{16}^2 between the GUT scale mass parameters of the first (or second) and third generations is reduced, albeit by a small amount.[4]

[4]Since the right hand side of eq. (6.2) is more negative when $S = 15$ as compared to the $S = 7$ case, the corresponding f runs to smaller values in the former case. If we now consider the evolution of δm_{16}^2, for which the term depending on g drops out, we see that this difference runs the most for $S = 7$ for which f is largest. In this sense, the difference shown by the solid lines may be regarded as a bound.

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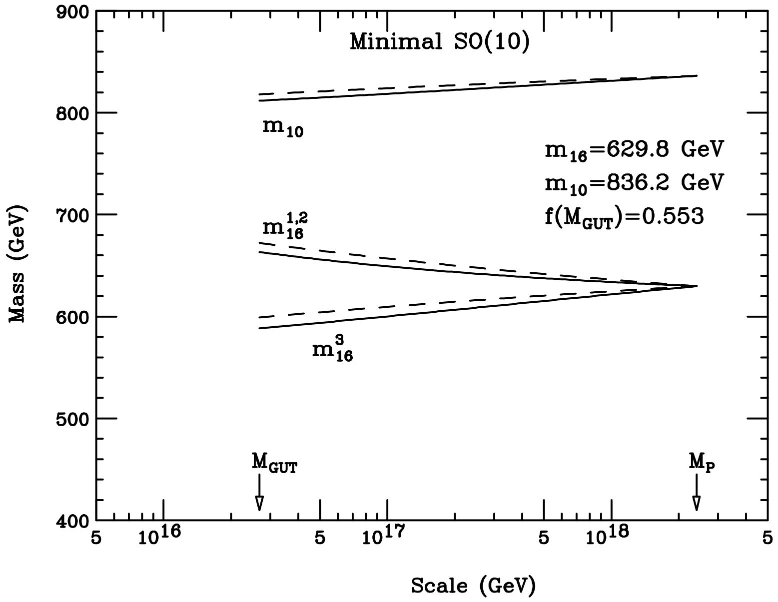

Figure 2: Running of the soft SUSY breaking masses between the Planck scale and the GUT scale in the *minimal* SO(10) model.

The effect of SO(10) running is that the first two generations of matter scalars run to higher masses, while the Higgs masses and third generation masses decrease somewhat. The corresponding weak-scale sparticle masses are listed in table 7, without and with the effect of Planck to GUT scale running. The main effect is a $\sim 27\%$ change in the mass difference between the (lightest) charged sleptons of the first and third generations.

6.2 General SO(10)

More generally, we may take the two MSSM Higgs doublets to live in different fundamental representations of SO(10): $\hat{H}_u \in \hat{H}_2$ and $\hat{H}_d \in \hat{H}_1$. Then the superpotential can be written as

$$\hat{f} = f_t \hat{\psi}\hat{\psi}\hat{H}_2 + f_b \hat{\psi}\hat{\psi}\hat{H}_1 \,, \tag{6.7}$$

so that there exist two Yukawa couplings above the GUT scale, and just $f_b = f_\tau$ unification, which can occur for a much wider range of $\tan\beta$ values [8], is required. In addition to the usual scalar masses, as in ref. [24], we include an off-diagonal mass term $m^2_{H_{12}}(H_1^\dagger H_2 + H_2^\dagger H_1)$. As in minimal SO(10), there should also be at least higher dimensional Higgs representations present responsible for SO(10) breaking, but again, we ignore these.

We give here the RGEs for the general SO(10) model, thereby completing the results of refs. [25, 24]. For the gauge coupling constant we have:

$$\frac{dg}{dt} = \frac{g^3}{16\pi^2}(S - 24) \,, \tag{6.8}$$

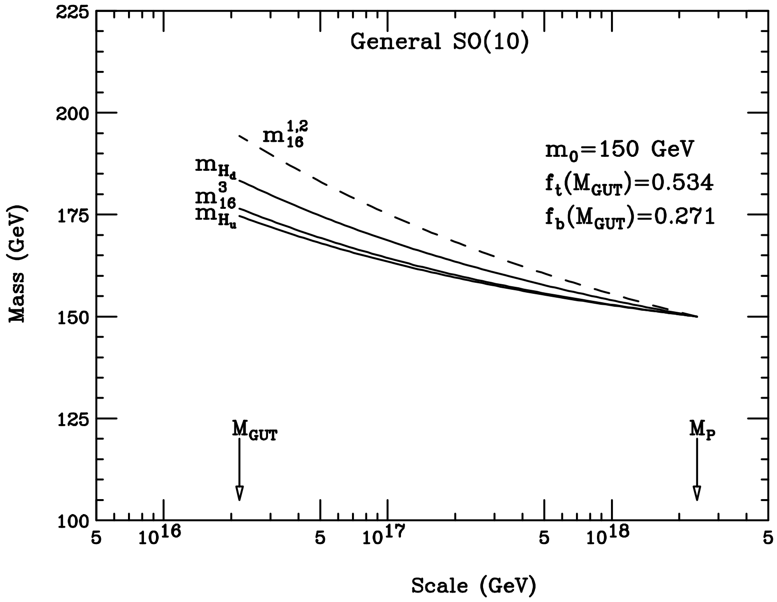

Figure 3: Running of the soft SUSY breaking masses between the Planck scale and the GUT scale in the general SO(10) model. The GUT scale Yukawa couplings here are the same as in the SU(5) case.

where S again is the sum of the Dynkin indices of the SO(10) fields. For just two 10 dimensional Higgs multiplets and 3 generations of matter, $S = 8$. For gaugino masses, we again have

$$\frac{dm_{1/2}}{dt} = \frac{1}{16\pi^2} 2(S-24) g^2 m_{1/2} \,. \tag{6.9}$$

The Yukawa coupling RGEs are:

$$\frac{df_t}{dt} = \frac{f_t}{16\pi^2}\left(14 f_t^2 + 14 f_b^2 - \frac{63}{2} g^2\right), \tag{6.10}$$

$$\frac{df_b}{dt} = \frac{f_b}{16\pi^2}\left(14 f_t^2 + 14 f_b^2 - \frac{63}{2} g^2\right). \tag{6.11}$$

The RGEs for the scalar masses are given by

$$\begin{aligned}\frac{dm_{16}^2}{dt} &= \frac{10}{16\pi^2}\Big[f_t^2(2m_{16}^2 + m_{H_2}^2) + f_b^2(2m_{16}^2 + m_{H_1}^2) + 2f_t f_b m_{H_{12}}^2 + \\ &\qquad + (A_t^2 f_t^2 + A_b^2 f_b^2) - \frac{9}{2} g^2 m_{1/2}^2\Big],\end{aligned} \tag{6.12}$$

$$\frac{dm_{H_1}^2}{dt} = \frac{8}{16\pi^2}\left[f_b^2(2m_{16}^2 + m_{H_1}^2) + f_t f_b m_{H_{12}}^2 + A_b^2 f_b^2 - \frac{9}{2} g^2 m_{1/2}^2\right], \tag{6.13}$$

$$\frac{dm_{H_2}^2}{dt} = \frac{8}{16\pi^2}\left[f_t^2(2m_{16}^2 + m_{H_2}^2) + f_t f_b m_{H_{12}}^2 + A_t^2 f_t^2 - \frac{9}{2} g^2 m_{1/2}^2\right], \tag{6.14}$$

$$\frac{dm_{H_{12}}^2}{dt} = \frac{4}{16\pi^2}\left[f_t f_b(4m_{16}^2 + m_{H_1}^2 + m_{H_2}^2 + 2A_t A_b) + (f_t^2 + f_b^2) m_{H_{12}}^2\right]. \tag{6.15}$$

Finally, the RGEs for the A-parameters are

$$\frac{dA_t}{dt} = \frac{1}{16\pi^2}(28f_t^2A_t + 20f_b^2A_b + 63g^2m_{1/2}) ,$$
$$\frac{dA_b}{dt} = \frac{1}{16\pi^2}(28f_b^2A_b + 20f_t^2A_t + 63g^2m_{1/2}) .$$

We show in figure 3 the running of SSB parameters in the general SO(10) model using GUT scale values of $g = 0.717$, $f_t = 0.534$ and $f_b = 0.271$, as in figure 1. Except for $m_{H_{12}}^2$ which is fixed to be zero at $Q = M_P$, the SSB parameters are also as in this figure. The main effect is again a significant splitting between first or second and third generation scalar masses at the GUT scale. Some splitting between m_{H_u} and m_{H_d} also occurs, with $m_{H_u}^2 < m_{H_d}^2$ as desired. The corresponding weak-scale sparticle masses are shown in table 8. The GUT scale SSB term splitting results in somewhat heavier scalars than in the mSUGRA case. For this example, because most of the weak-scale squark mass comes from the RG evolution, the effect is more pronounced for sleptons than squarks. In particular, this increase is just a few percent for squarks, but as much as 22% for sleptons.

parameter	mSUGRA	SO(10)
m_{16_1}	150.0	195.0
m_{16_3}	150.0	177.0
m_{H_d}	150.0	183.9
m_{H_u}	150.0	175.1
$m_{H_{12}}$	—	−46.6
$A_t(M_{\rm GUT})$	0.0	−117.8
$A_b(M_{\rm GUT})$	0.0	−120.0
$m_{\tilde g}$	512.0	492.9
$m_{\tilde u_L}$	468.0	468.1
$m_{\tilde d_R}$	454.7	457.2
$m_{\tilde t_1}$	335.6	319.1
$m_{\tilde b_1}$	375.4	360.1
$m_{\tilde\ell_L}$	212.8	241.3
$m_{\tilde\ell_R}$	174.5	213.2
$m_{\tilde\nu_e}$	197.2	227.7
$m_{\tilde\tau_1}$	124.3	139.6
$m_{\tilde\nu_\tau}$	189.3	201.4
$m_{\widetilde W_1}$	150.1	142.5
$m_{\widetilde Z_2}$	150.3	142.4
$m_{\widetilde Z_1}$	80.8	77.8
m_h	111.0	111.4
m_A	210.4	218.7
m_{H^+}	227.9	235.6
μ	−263.8	−285.2

Table 8: Model parameters and weak scale sparticle masses for mSUGRA model and general SO(10). The first model has universality of matter scalars at $M_{\rm GUT}$, while the second has universality at M_P. For both cases, we take $m_0 = 150$ GeV, $m_{1/2} = 200$ GeV, $A_0 = 0$, $\tan\beta = 35$ and $\mu < 0$, and list the values of the SSB masses at the GUT scale. This yields $f_t = 0.534$ and $f_b = f_\tau = 0.271$, with $g(M_{\rm GUT}) = 0.717$ as in table 1.

7. Supersymmetric missing partner models with hypercolor

In this variety of models, the gauge group is of the type $G_{\rm GUT} \times G_H$, where the first group is SU(5) or SO(10) and the second is related to a "hypercolor" interaction [26, 27]. While the weak SU(2) is completely contained in the first factor, colour SU(3) is not embedded in either of the factors. Although the gauge group is not simple, an approximate unification of the gauge coupling constants of the group $SU(3)_C \times SU(2) \times U(1)$ is achieved if the couplings of G_H are large enough. These models provide a solution to the doublet-triplet splitting problem by the missing partner mechanism. Since the MSSM gauginos do not belong to a single multiplet of a simple gauge group, their masses do not obey the usual unification condition [28], resulting in non universality of gaugino masses. However, if usual squarks and sleptons and the MSSM Higgs fields are singlets of G_H, universality of scalar masses is still possible, as for instance, in the $\mathrm{SU}(5)_{\mathrm{GUT}} \times \mathrm{SU}(3)_H \times \mathrm{U}(1)_H$ model of ref. [27], where the hypercharge U(1) is a combination of $\mathrm{U}(1)_H$ and a U(1) subgroup in the first factor.

In this case, the following relations among gauge couplings hold at the unification scale [28]:

$$\frac{1}{g_1^2} = \frac{1}{g_{\mathrm{GUT}}^2} + \frac{1}{15 g_{H1}^2}, \qquad \frac{1}{g_2^2} = \frac{1}{g_{\mathrm{GUT}}^2}, \qquad \frac{1}{g_3^2} = \frac{1}{g_{\mathrm{GUT}}^2} + \frac{1}{g_{H3}^2}, \tag{7.1}$$

where $\sqrt{3/5}g_1$, g_2, and g_3 are the gauge couplings of the $\mathrm{U}(1)_Y$, $\mathrm{SU}(2)_L$, and $\mathrm{SU}(3)_C$ SM groups, and g_{GUT}, g_{H3}, and g_{H1} are the $\mathrm{SU}(5)_{\mathrm{GUT}}$, $\mathrm{SU}(3)_H$, and $\mathrm{U}(1)_H$ unified groups, respectively. Clearly from eq. (7.1) we see that the unification of the gauge coupling constants from low-energy data is achieved if $g_{H1}^2 \gg g_{\mathrm{GUT}}^2$ and $g_{H3}^2 \gg g_{\mathrm{GUT}}^2$. In addition, considering that the prediction for α_s at the weak scale in SUSY GUT models (without threshold corrections) is higher than the world averaged experimental value, it was argued that the correction introduced by hypercolor moves the prediction for α_s in the correct direction. It was found that [28]:

$$\alpha_s(m_Z) \approx 0.130 - \frac{0.014}{\alpha_{H3}} - \frac{0.010}{15\alpha_{H1}}, \tag{7.2}$$

where $\alpha_i = g_i^2/4\pi$ and threshold corrections have been neglected. In order for α_s not to shift too much, we must have $\alpha_{H3} \gtrsim 0.6$ and $\alpha_{H1} \gtrsim 0.03$, though for α_{H1} as small as 0.03, $g_2^2 - g_1^2 = 0.18 g_1^2 g_2^2$.

Above the GUT scale there are three gauginos associated to the groups $\mathrm{SU}(5)_{\mathrm{GUT}}$, $\mathrm{SU}(3)_H$, and $\mathrm{U}(1)_H$ whose masses we denote $m_{1/2}$, M_{H3} and M_{H1}, respectively. Below the GUT scale we have the MSSM and the three MSSM gauginos are a linear combination of the former ones. The masses of the bino, wino, and gluino are then given by,[5]

$$\begin{aligned} M_1 &= g_1^2\left(\frac{m_{1/2}}{g_{\mathrm{GUT}}^2} + \frac{M_{H1}}{15 g_{H1}^2}\right), \\ M_2 &= m_{1/2}, \\ M_3 &= g_3^2\left(\frac{m_{1/2}}{g_{\mathrm{GUT}}^2} + \frac{M_{H3}}{g_{H3}^2}\right). \end{aligned} \tag{7.3}$$

The thing to note is that $M_{H1,3}/\alpha(H_{1,3})$ are renormalization group invariants (at one loop) so that $M_{H1,3}/g_{H1,3}^2$ need not be small even when $g_{H1,3}^2$ is large. The relative magnitude of the three masses $m_{1/2}$, M_{H3} and M_{H1} is unknown because it depends on the SUSY breaking mechanism. One might naively suppose that they are of the same order of magnitude; in this case, gaugino masses could be significantly

[5] A somewhat different model[26, 27] based on the group $\mathrm{SO}(10)_{\mathrm{GUT}} \times \mathrm{SO}(6)_H$ also has non-universal MSSM gaugino masses. However, since the hypercolor group is simple, there is one relation between them [29].

different at the GUT scale, though the magnitude of the non universality would be limited because, as noted above, the couplings g_{H1} and g_{H3} have to be considerably larger than $g_{\rm GUT}$. There is no reason, however, why M_{H1} and M_{H3} cannot be much larger than $m_{1/2}$. Indeed in scenarios with dilaton dominated SUSY breaking, we have [30, 31]

$$\frac{m_{1/2}}{g^2_{\rm GUT}} = \frac{M_{H1}}{g^2_{H1}} = \frac{M_{H3}}{g^2_{H3}}\,, \tag{7.4}$$

so that gaugino mass splittings of O(100%) are expected.

In figure 4a we plot non-universal gaugino masses of the MSSM as a function of a common hypercolor gaugino mass $M_{H1} = M_{H3} \equiv M_H$. We take $m_{1/2} = 200\,\mathrm{GeV}$, $g_{\rm GUT} = 0.716$, and two different choices for the hypercolor gauge couplings: $\alpha_{H1} = 0.1$ and $\alpha_{H3} = 0.7$ in solid lines, and $\alpha_{H1} = 0.5$ and $\alpha_{H3} = 0.8$ in dashed lines. As indicated in eq. (7.3) the wino mass M_2 is always equal to $m_{1/2} = 200\,\mathrm{GeV}$. The other two gaugino masses are larger (smaller) than M_2 if the hypercolor gaugino mass is larger (smaller) than $m_{1/2}$. The gluino mass deviates more from M_2 compared to the bino mass because of the factor 15 in eq. (7.3) and our choice of values for other parameters.

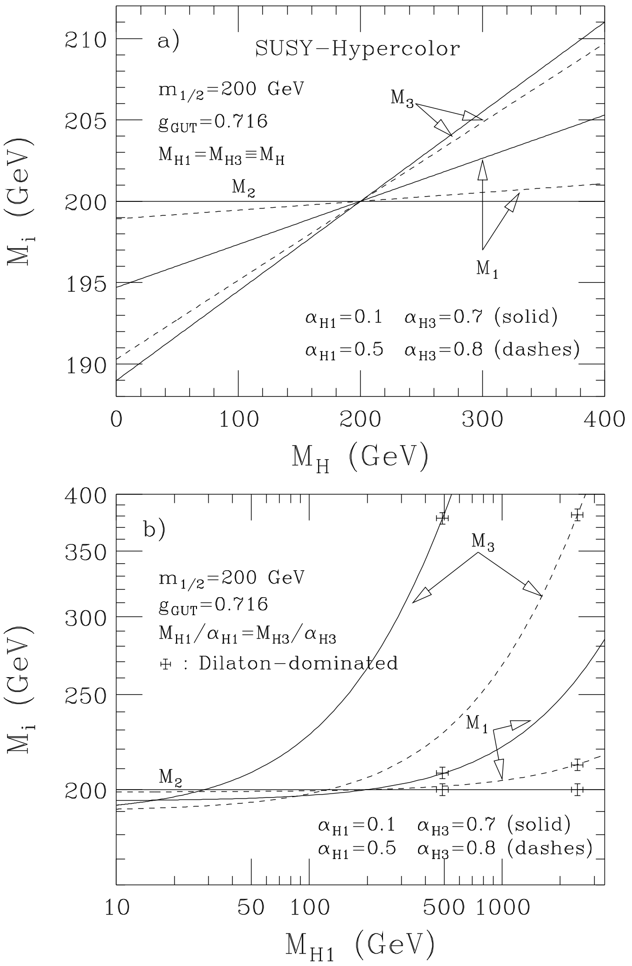

Figure 4: Non-universal gaugino masses M_i, $i = 1, 2, 3$, in a supersymmetric missing partner model with hypercolor, as a function of the common gaugino mass $M_{H1} = M_{H3} \equiv M_H$ in the hypercolor sector. Two values of the hypercolor gauge couplings are used.

The larger the hypercolor gauge couplings, the smaller the deviations from universality. In addition, if the common hypercolor mass is equal to $m_{1/2}$ there is no deviation from universality no matter the value of the hypercolor gauge couplings. In figure 4b, we show the same gaugino masses but assuming instead that $M_{H1}/g^2_{H1} = M_{H_3}/g^2_{H3}$. The three gauge couplings are chosen exactly as in frame a so that there is a large hierarchy between the masses of the gauginos of the three groups. The cross denotes the dilaton-dominated scenario for which point eq. (7.4) is satisfied. Indeed we see that very large non universality of gaugino masses may be possible.

In figure 5, we show several weak-scale sparticle masses versus the same parameter M_H as in figure 4 for parameter values corresponding to the solid curves in this figure. The two frames illustrate the results for the same choices of the gaugino masses as in figure 4. In frame a we see that the non-colored sparticle masses hardly vary at all versus M_H, while the gluino and squark masses can vary by up to 12%. This is presumably because the coloured sparticle masses run considerably more than those of uncoloured sparticles coupled with the fact that M_3 varies more with M_H than M_1 does, and M_2 does not change at all. The variation is, of course, much more dramatic in frame b. For very large values of M_{H1}, the coloured sparticles as well as the heavier chargino and neutralinos become very heavy, and may be in conflict with fine-tuning considerations. We also mention that although M_1 starts out larger than M_2 at the GUT scale (but not by a huge amount), M_1/M_2 is driven to a value close to 1/2 at the weak scale for acceptable values of M_{H1}: it would be interesting to examine whether precise measurements of masses and mixing angles could lead to observable deviations from expectations in mSUGRA or gauge-mediated SUSY breaking frameworks. In the same vein, we also mention that $m(\tilde{e}_R)$ also increases slowly from 132 GeV in mSUGRA to 134 GeV for the dilaton dominated scenario to 151 GeV for the extreme case with $M_{H1} = 3$ TeV, while $m(\tilde{e}_L)$ is roughly constant. This is because the RG evolution of $m(\tilde{e}_R)^2$ is due to hypercharge gauge interactions, and M_1 starts out bigger than M_2 (which is independent of M_{H1}).

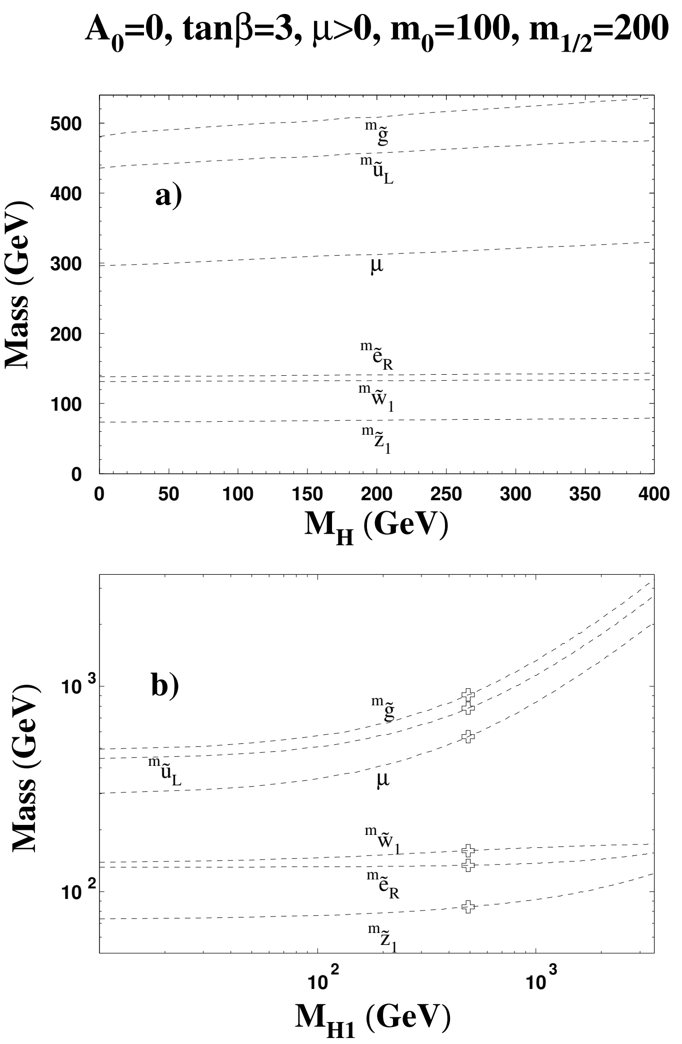

Figure 5: Various sparticle masses in the hypercolor model.

8. Models with effective supersymmetry

The SM exhibits accidental global symmetries which inhibit flavor-changing neutral currents (FCNC), lepton flavor violation (LFV), electric dipole moments (EDM) of electron and neutron, and proton decay, as opposed to the MSSM where degeneracy or alignment in the mass matrices has to be invoked. On the other hand, supersymmetry stabilizes the scalar masses under radiative corrections, contrary to the SM where it is hard to understand the hierarchy between the Higgs mass and the Planck scale. The models presented in this section [32, 33] aim to combine the good features of both the SM and the MSSM. There are two mass scales: gauginos, higgsinos, and third generation squarks are sufficiently light ($\lesssim$ 1 TeV) to naturally stabilize the Higgs mass and the electroweak scale, while the first two generations of squarks and sleptons (whose Yukawa couplings to Higgs are very small) are sufficiently heavy ($\widetilde{M} \sim 5$ to 20 TeV) to suppress FCNC, LFV, etc. This class of models, called Effective Supersymmetry, does not invoke degeneracy or alignment in the mass matrices.

In one of the realizations of Effective Supersymmetry [33], the first two generations of squarks and sleptons, together with the down-type Higgs, are composite, with constituents that carry a "superglue" charge, and have a mass $\sim \widetilde{M}$. Gauge superfields, third generation superfields and the up-type Higgs superfield are taken to be fundamental and neutral under superglue, with perturbative couplings to the constituents, so that their mass is suppressed relative to the mass of the composites. In this way, the spectrum is characterized as follows.

- Gaugino masses are light and can be non universal with masses given by $M_i = n_i(\alpha_i/4\pi)\widetilde{M}$, where n_i are numerical factors that can be as large as $\mathcal{O}(10)$.
- Left and right squark and slepton masses for the first two generations are of the order of $\widetilde{M}$.
- Left and right squark and slepton masses for the third generation are of the order of $(\lambda_3/4\pi)\widetilde{M}$; for $\lambda_3 \sim 1$, this is an order of magnitude smaller than $\widetilde{M}$.
- The down-Higgs mass satisfy $m_{H_d} \sim \widetilde{M}$. The up-Higgs mass on the other hand, is given by $m_{H_u} \sim (\lambda_H/4\pi)\widetilde{M}$, where λ_H is its perturbative coupling to the constituents. Therefore, there is only one Higgs in the low-energy theory and $\tan\beta \sim 4\pi/\lambda_H$ is large.
- The "μ-term" and the "$B\mu$-term" respectively satisfy $\mu \sim (\lambda_H/4\pi)\widetilde{M}$ and $B\mu \sim (\lambda_H/4\pi)\widetilde{M}^2$.

To obtain $m_{H_u} \sim 100\,\text{GeV}$, we require $\lambda_H/4\pi \sim 10^{-2}$, while $\lambda_3/4\pi \sim 10^{-1}$ ensures $m_{\tilde{t}} \lesssim 1$ TeV.

If the hierarchy of scalar masses is already present at the unification scale, then it has been shown that unless the stop mass squared at the unification scale is taken to be well above (1 TeV)^2, two-loop contributions to scalar renormalization group equations drive the top squark mass squared negative well before the weak scale, resulting in a breakdown of color symmetry [34]. Thus, this simple class of models seems to be ruled out by fine-tuning considerations. To account for this class of constraints, we have implemented the full set of two-loop MSSM RGEs in ISAJET versions ≥ 7.50.

Very recently, Hisano et al. [35] have identified scenarios in which first and second generation scalars can be much heavier than gauginos and scalars of the third generation, and for which the scalar masses are renormalization group invariant (so that the constraints of ref. [34] are not relevant) as long as gaugino masses are neglected in the RGEs. These constraints are also inapplicable in models in which the assumption of the scalar hierarchy is made for mass parameters at a scale $\sim 10 \div 50$ TeV, since then there are no large logs that drive m^2 to negative values. In this case, however, model-dependent *finite* contributions to δm^2 are no longer negligible, and need to be examined to discuss the viability of any particular model [36].

Yet another possibility has been considered in refs. [37, 24, 38]. These authors begin with *all* scalar masses initially at the multi-TeV scale at or above $M_{\rm GUT}$, and show that for certain choices of $M_{\rm GUT} - M_{\rm Planck}$ scale boundary conditions on the scalar masses and A-parameters-keeping gaugino masses at the weak scale — the third generation sfermion and Higgs masses are driven to weak-scale values, while scalars of the first two generations remain heavy. Such a scenario is particularly attractive in the context of minimal SO(10). In this case, with Yukawa coupling unification plus a singlet $\hat{N}^c$, particularly simple boundary conditions [38],

$$4m_{16}^2 = 2m_{10}^2 = A^2 \tag{8.1}$$

lead to sub-TeV scale third generation scalar masses, while first and second generation scalar masses can be as high as 20 TeV. If instead the boundary value of A is taken to be at the weak scale, the hierarchy generated [24] is somewhat smaller. Examples of sparticle mass spectra were not generated in ref. [24], where it was noted that this scenario shares the problem of obtaining correct radiative breaking of electroweak symmetry common to most high $\tan\beta$ scenarios: in examples shown in refs. [24, 38], the two Higgs SSB masses stay positive at all scales in their evolution to the weak scale, with $m_{H_u} > m_{H_d}$, contrary to what is needed for REWSB.

In a recent analysis [39] it has been shown that if the boundary conditions in eq. (8.1) are augmented by SO(10) D-terms, it is possible to obtain the desired inverted mass hierarchy amongst the squarks together with radiative electroweak symmetry breaking. This then yields a calculable model based on the gauge group SO(10) with (approximate) unification of Yukawa couplings. The analysis in ref. [39] took the right-handed neutrino mass to be fixed near $\sim 10^{13}$ GeV, and obtained "crunch" factor values S up to $\sim 5 \div 7$ for full SO(10) D-terms, and factors of S

up to 9 if splittings were applied only to the soft SUSY breaking Higgs masses. The crunch factor S is defined as

$$S = \frac{3(m_{u_L}^2 + m_{d_L}^2 + m_{u_R}^2 + m_{d_R}^2) + m_{\tilde{e}_L}^2 + m_{\tilde{e}_R}^2 + m_{\tilde{\nu}_e}^2}{3(m_{\tilde{t}_1}^2 + m_{\tilde{b}_1}^2 + m_{\tilde{t}_2}^2 + m_{\tilde{b}_2}^2) + m_{\tilde{\tau}_1}^2 + m_{\tilde{\tau}_2}^2 + m_{\tilde{\nu}_\tau}^2}\,.$$

These values are considerably below those quoted in ref. [38], where a more idealized case was considered.

Effective supersymmetry is not as mature a framework as mSUGRA or the gauge-mediated SUSY breaking. Except for the inverted hierarchy model of the previous paragraph, all the models discussed in this section suffer from incompleteness which preclude computations at as thorough a level. The scenario in ref. [33] involves new unknown strong dynamics at the 10 TeV scale. Models where the splitting between third generation scalars and those of the other generations has a dynamical origin [37, 24, 38] suffer from the fact that this dynamics does not break electroweak symmetry: the mass spectrum thus does not appear to be calculable unless deviations such as non universality are imposed. These considerations notwithstanding, collider events for generic effective SUSY models can be generated with ISAJET [5] by using the weak-scale $MSSMi$ keywords, with independent weak-scale SSB masses as inputs. One may enter multi-TeV scale first and second generation scalar masses, while using sub-TeV scale gaugino masses, third generation scalar masses and μ parameters. In the scenario of ref. [33], A-terms are $\mathcal{O}(100)$ GeV or smaller, while m_A is very large.

Sparticle mass spectra from the radiatively generated inverted mass hierarchy solution due to Bagger et al. are not possible without modifications that allow REWSB to occur. Two possibilities are the non-universalities due to SO(10) D-terms, or ad hoc Higgs sector splittings. These may be implemented in ISAJET using the NUSUG inputs along with the right-handed neutrino solution. In ISAJET, if a zero physical neutrino mass is entered, then the Yukawa couplings f_t and f_ν automatically unify. It remains to be seen whether the resulting inverted mass hierarchy is truly sufficient to solve problems due to FCNCs, LFVs and the EDM of the electron and neutron.

9. Anomaly-mediated SUSY breaking

In most models, soft SUSY breaking parameters of the low-energy effective theory are thought to receive contributions from gravitational or gauge interactions which are considered to be messengers of SUSY breaking in a hidden sector. It has recently been recognized [40, 41] that there is an additional contribution, that originates in the super-Weyl anomaly, which is always present when SUSY is broken. In models without SM gauge singlet superfields that can acquire a Planck scale vev , the usual supergravity contribution to gaugino masses is suppressed by an additional factor

$M_{\rm SUSY}/M_P$ relative to $m_{3/2} = M_{\rm SUSY}^2/M_P$, and the anomaly-mediated contribution can dominate. These contributions are determined in terms of the SUSY breaking scale by the corresponding β-functions.

$$M_i = \frac{\beta_g}{g} m_{3/2}\,, \tag{9.1}$$

where β_i is the one-loop β-function, defined by $\beta_{g_i} \equiv dg_i/d\ln\mu = -b_i g_i^3 + \cdots$. The gaugino masses are not universal, but given by the ratios of the respective β-functions.

In general, however, Kähler potential couplings between the observable sector and the hidden sector (Goldstino) field, which are generically not forbidden by a symmetry, result in large gravity contributions ($\sim m_{3/2}$) to scalar masses which would completely dominate the corresponding anomaly-mediated contributions. These gravity contributions can be strongly suppressed if the SUSY breaking and visible sectors reside on different branes, and are "sufficiently separated" in a higher dimensional space: in this case, the suppression is the result of geometry and not a symmetry, though then one has to wonder about the dynamics that results in such a geometry. The anomaly-mediated contribution is given by,

$$m_{\tilde{q}}^2 = -\frac{1}{4}\left\{\frac{d\gamma}{dg}\beta_g + \frac{d\gamma}{df}\beta_f\right\} m_{3/2}^2\,, \tag{9.2}$$

where β_g and β_f are the β-functions for gauge and Yukawa interactions, respectively, and $\gamma = \partial\ln Z/\partial\ln\mu$, with Z the wave function renormalization constant. Notice that this is comparable to the corresponding contribution to the gaugino masses. Furthermore, since Yukawa interactions are negligible for the first two generations, the anomaly-mediated contributions to scalar masses of the first two generations are essentially equal. Unfortunately, however [40], the anomaly contribution turns out to be negative for sleptons, necessitating additional sources for the squared masses of scalars. There are several proposals in the literature, but phenomenologically it suffices to add a universal contribution m_0^2 (which, of course, preserves the degeneracy between the first two generations of scalars) to eq. (9.2), and regard m_0 as an additonal parameter [42].

Finally, in the sign convention of ISAJET[6], the anomaly-mediated contribution to the trilinear SUSY breaking scalar coupling is given by,

$$A_f = +\frac{\beta_f}{f} m_{3/2}\,. \tag{9.3}$$

It is assumed that the ad hoc introduction of m_0^2 in eq. (9.2) does not affect the other relations.

[6]This is opposite to that used in ref. [42].

9.1 The minimal anomaly-mediated SUSY breaking model (AMSB)

In this framework, it is assumed that the anomaly-mediated SUSY breaking contributions to the soft-SUSY breaking contributions dominate, and further, that the introduction of the parameter m_0^2 is sufficent to circumvent the problem of negative squared masses for sleptons. The parameter space of the model consists of

$$m_0,\ m_{3/2}\,,\ \tan\beta \quad \text{and} \quad \mathrm{sign}(\mu)\,. \tag{9.4}$$

In this case, gaugino masses are given by

$$M_1 = \frac{33}{5}\frac{g_1^2}{16\pi^2}m_{3/2}\,, \tag{9.5}$$

$$M_2 = \frac{g_2^2}{16\pi^2}m_{3/2}\,, \tag{9.6}$$

$$M_3 = -3\frac{g_3^2}{16\pi^2}m_{3/2}\,. \tag{9.7}$$

Third generation scalar masses are given by

$$m_U^2 = \left(-\frac{88}{25}g_1^4 + 8g_3^4 + 2f_t\hat{\beta}_{f_t}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,, \tag{9.8}$$

$$m_D^2 = \left(-\frac{22}{25}g_1^4 + 8g_3^4 + 2f_b\hat{\beta}_{f_b}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,, \tag{9.9}$$

$$m_Q^2 = \left(-\frac{11}{50}g_1^4 - \frac{3}{2}g_2^4 + 8g_3^4 + f_t\hat{\beta}_{f_t} + f_b\hat{\beta}_{f_b}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,, \tag{9.10}$$

$$m_L^2 = \left(-\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + f_\tau\hat{\beta}_{f_\tau}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,, \tag{9.11}$$

$$m_E^2 = \left(-\frac{198}{25}g_1^4 + 2f_\tau\hat{\beta}_{f_\tau}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,, \tag{9.12}$$

$$m_{H_u}^2 = \left(-\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + 3f_t\hat{\beta}_{f_t}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,, \tag{9.13}$$

$$m_{H_d}^2 = \left(-\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + 3f_b\hat{\beta}_{f_b} + f_\tau\hat{\beta}_{f_\tau}\right)\frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2\,. \tag{9.14}$$

The A-parameters are given by

$$A_t = \frac{\hat{\beta}_{f_t}}{f_t}\frac{m_{3/2}}{16\pi^2}\,, \tag{9.15}$$

$$A_b = \frac{\hat{\beta}_{f_b}}{f_b}\frac{m_{3/2}}{16\pi^2}\,,\ \text{and} \tag{9.16}$$

$$A_\tau = \frac{\hat{\beta}_{f_\tau}}{f_\tau}\frac{m_{3/2}}{16\pi^2}\,. \tag{9.17}$$

In the above, we have

$$\hat{\beta}_{f_t} = 16\pi^2\beta_t = f_t\left(-\frac{13}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + 6f_t^2 + f_b^2\right), \tag{9.18}$$

$$\hat{\beta}_{f_b} = 16\pi^2\beta_b = f_b\left(-\frac{7}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + f_t^2 + 6f_b^2 + f_\tau^2\right), \tag{9.19}$$

$$\hat{\beta}_{f_\tau} = 16\pi^2\beta_\tau = f_\tau\left(-\frac{9}{5}g_1^2 - 3g_2^2 + 3f_b^2 + 4f_\tau^2\right). \tag{9.20}$$

The first two generations of squark and slepton masses are given by the corresponding formulae above with the Yukawa couplings set to zero. This model has been implemented in ISAJET versions ≥ 7.45, using the AMSB keyword, which allows input of the above parameter space set. In ISAJET, it is easiest to implement the above masses at scale $Q = M_{\rm GUT}$, and proceed with evolution to the weak scale. Then the B and μ^2 parameters are calculated in accord with the constraint from radiative electroweak symmetry breaking.

The most notable feature of this framework is the hierarchy of gaugino masses. The gluino is (as in mSUGRA) much heavier than the electroweak gauginos, but the novel feature is that $M_1/M_2 \sim 3.2$, so that the wino is by far the lightest supersymmetric particle (LSP). The wino LSP scenario has several implications for phenomenology, the most important of which is the near degeneracy of the chargino and the (wino-like) neutralino LSP. One-loop corrections [44, 43, 42, 45], which make the dominant contribution to the chargino-neutralino mass gap, have been included [5] in ISAJET v7.46 (in the gaugino limit). The phenomenology can be sensitive to this mass difference [43, 42].

In table 9, we show spectra generated from the minimal AMSB model for two values of m_0, with other parameters being the same. Note that the parameter $m_{3/2}$ should be selected typically above 25,000 GeV to avoid constraints from LEP experiments. From the spectra shown, we immediately see several well-known aspects of the AMSB spectrum. Most notably, we see that the $\widetilde{W}_1$ and $\widetilde{Z}_1$ are nearly degenerate in mass, so that in addition to the usual leptonic decay modes $\widetilde{W}_1 \to \widetilde{Z}_1\ell\nu$, the only other allowed (and in these cases dominant) decay of the chargino is $\widetilde{W}_1^\pm \to \widetilde{Z}_1\pi^\pm$. The chargino has a very small width, corresponding to a lifetime $\sim 1.5 \times 10^{-9}$ s, so that it would be expected to travel a significant fraction of a meter before decaying [42]. Secondly, the $\tilde{\ell}_L$ and $\tilde{\ell}_R$ are nearly mass degenerate. This degeneracy (which seems fortuitous) is much tighter than expected in the mSUGRA framework and certainly in the gauge-mediated SUSY breaking framework. Their mass scale is largely determined by the parameter m_0, and it is possible that for small enough m_0 slepton signals may be detectable at the next generation of e^+e^- colliders or even at the LHC. Another interesting feature (which may serve to distinguish the cases shown from mSUGRA) is that the $\tilde{\tau}_L - \tilde{\tau}_R$ mixing is near maximal. The prospects for measuring this have been discussed in ref. [46].

Parameter	AMSB (200)	AMSB (500)
m_0	200	500
$m_{3/2}$	35,000	35,000
$\tan\beta$	5	5
μ	> 0	> 0
$m_{\tilde{g}}$	816	825
$m_{\tilde{u}_L}$	754	872
$m_{\tilde{t}_1}$	512	588
$m_{\tilde{b}_1}$	666	758
$m_{\tilde{\ell}_L}$	156	484
$m_{\tilde{\ell}_R}$	154	483
$m_{\tilde{\tau}_1}$	132	476
$m_{\tilde{\tau}_2}$	173	489
$m_{\widetilde{W}_1}$	99.5	99.1
$m_{\widetilde{Z}_2}$	321	321
$m_{\widetilde{W}_1} - m_{\widetilde{Z}_1}$	0.171	0.172
m_h	114	113
m_A	654	806
μ	632	635
θ_τ	0.82	0.86
θ_b	0.11	0.074

Table 9: Model parameters and weak-scale sparticle masses in GeV for an anomaly-mediated SUSY breaking (AMSB) case study.

In the minimal AMSB framework, $m_{\widetilde{W}_1} - m_{\widetilde{Z}_1}$ is typically bigger than 160 MeV, so that $\widetilde{W}_1 \to \widetilde{Z}_1\pi$ is always allowed and the chargino typically decays within the detector [42]. The chargino would then manifest itself only as missing energy, unless the decay length is a few tens of cm, so that the chargino track can be established in the detector. The track would then seem to disappear [43] since the presence of the soft pion would be very difficult to detect. Some parameter regions with $m_{\widetilde{W}_1} - m_{\widetilde{Z}_1} < m_{\pi^\pm}$ may be possible; in this case, the chargino would mainly decay via $\widetilde{W}_1 \to \widetilde{Z}_1 e\nu$ and its decay length (depending on the mass difference) would be typically larger than several metres. It would then show up via a search for long-lived charged exotics.

There have been a number of alternative suggestions to cure the negative slepton mass squared problem [47]. Generally, these require the introduction of additional fields at energy scales higher than the weak scale. The mass spectrum in these scenarios differs from that of the minimal AMSB model sketched above, and characteristic features such as $m_{\widetilde{W}_1} \simeq m_{\widetilde{Z}_1}$ and $m_{\tilde{\ell}_L} \simeq m_{\tilde{\ell}_R}$ need not occur. These models are not hard wired into ISAJET, but can be generated using the NUSUG$_i$ inputs at a scale dictated by SSBCSC; in this case, the user must perform the calculation of the SSB masses of MSSM particles.

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10. Minimal gaugino mediation

Very recently, Schmaltz and Skiba [48] have proposed a model based on extra dimensions with branes, which is claimed to provide novel solutions to the SUSY flavour and CP problems. Within their framework, chiral supermultiplets of the observable sector reside on one brane whereas the SUSY breaking sector is confined to a different brane [40]. Gravity and gauge superfields propagate in the bulk, and hence, directly couple to fields on both the branes. As a result of their direct coupling to the SUSY breaking sector, gauginos acquire a mass. The scalar components of the chiral supermultiplets, however, can acquire a SUSY breaking mass only via their interactions with gauginos (or gravity) which feel the effects of SUSY breaking: as a result, these masses are suppressed relative to gaugino masses, and may be neglected in the first approximation. The same is true for the A- and B-parameters.

In the specific realization [48], to preserve the success of the unification of gauge couplings, it is assumed that there is grand unification (both SU(5) and SO(10) are discussed), and further, that the compactification scale M_c below which there are no Kaluza-Klein excitations, is larger than $M_{\rm GUT}$. Furthermore, since light-bulk fields have flavor-blind interactions by construction, it is argued that the scale $M_c \lesssim M_{\rm Planck}/10$ in order to sufficiently suppress flavour violating scalar couplings (due to heavy-bulk fields) that would be generically present. Based on the discussion in the previous paragraph, they take the boundary conditions for the soft SUSY breaking parameters of the MSSM to be, $m_0 = A_0 = B_0 = 0$ at the scale M_c, and argue that the spectrum is completely specified by the parameter set,

$$\mu\,, m_{1/2}\,, M_c\,, \tag{10.1}$$

where it is the grand unification assumption that leads to a universal gaugino mass above $Q = M_{\rm GUT}$. They refer to this as the Minimal Gaugino Mediation (MGM) model. The parameters $m_{1/2}$ and μ should be comparable, and are chosen to be $\sim M_{\rm weak}$. The REWSB constraints fix μ^2, while the requirement $B_0 = 0$ fixes $\tan\beta$. In ref. [48] it is shown that if $M_c \leq M_{\rm Planck}/10$ $\tan\beta$ lies between ~ 12 and ~ 18 (12–25) for the SU(5) (SO(10)) model with $\mathbf{5}+\bar{\mathbf{5}}$ ($\mathbf{16}+\bar{\mathbf{16}}$) Higgs supermultiplets in addition to the usual adjoint Higgs multiplet. The LSP may be the stau, the lightest neutralino or the gravitino. However, the latter has a weak-scale mass, and as in the mSUGRA framework, is irrelevant for collider phenomenology.

Our purpose here is to outline how to generate sample spectra in this framework using ISAJET [5], and examine some issues that have not been discussed in ref. [48]. For definiteness, we will choose the GUT group to be SU(5). This model is then a special case of our discussion in section 2, except that the SSB parameters now "unify" at the scale M_c rather than M_P (where they take on the special values). Our first observation is that the allowed range of $\tan\beta$ seems incompatible [8] with

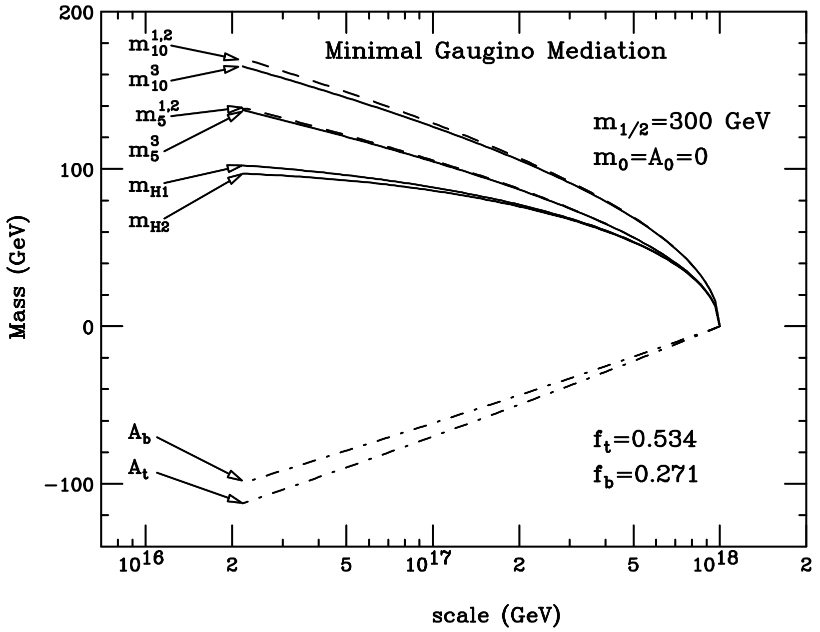

Figure 6: Renormalization group evolution of soft SUSY breaking SU(5) masses versus scale in the minimal gaugino mediation model. We take $\tan\beta = 35$ and $\mu < 0$ to achieve b-τ Yukawa coupling unification.

$\tan\beta \geq 30$ required for the unification of the τ and b Yukawa couplings.[7] For this reason, and also because the prediction for $\tan\beta$ could depend on how the μ problem might be solved, we will ignore the $B_0 = 0$ condition and treat $\tan\beta$ as a phenomenological parameter.[8] For our analysis, we modify the model parameters[9] to,

$$m_{1/2}\,, M_c\,, \tan\beta\,, \mathrm{sign}(\mu)\,. \tag{10.2}$$

As before, the user will have to obtain the values of the SSB parameters at $Q = M_{\mathrm{GUT}}$ using the RG equations of section 2, and input these into ISAJET for generating mass spectra and/or collider events as desired. As shown in table 10, we fix $m_{1/2}$ at the GUT scale and $\tan\beta$ at the weak scale.

In figure 6, we show the evolution of the various SSB parameters of the MSSM, starting with the MGM boundary conditions. Here, the unified gaugino mass is taken to be 300 GeV at $Q = M_{\mathrm{GUT}}$. The compactification scale is taken to be $M_c = 10^{18}$ GeV, and other parameters are fixed to be the same as in figure 1. We see that RG evolution results in GUT scale scalar masses and A-parameters that are

[7] Another possibility is the inclusion of a bilinear R-parity violating term in the tau sector. In this case, b-τ Yukawa unification can be achieved at smaller values of $\tan\beta$ [49].

[8] Moreover, if Higgs fields are also allowed to propogate in the bulk [50], we would expect $B_0 \sim m_{1/2} \sim m_{H_u} \sim m_{H_d}$.

[9] There are other coupling constants involving GUT scale physics, but we will see that these do not significantly change the spectrum.

substantial fractions of $m_{1/2}$; i.e. although we have no-scale [51] boundary conditions at the scale M_c, there are substantial deviations from these at $M_{\rm GUT}$. While the inter-generation splitting is small, the splittings between the **5** and the **10** dimensional matter multiplets, as well as between these and the Higgs multiplets is substantial.

In figure 7, we show the variation of several SSB masses at the scale $Q = M_{\rm GUT}$ with the unified gaugino mass $m_{1/2}$ for the same values of other parameters as in the previous figure. These masses then serve as inputs for ISAJET. We note that if $m_{1/2}$ is too small, the no-scale like boundary conditions lead to incorrect electroweak symmetry breaking or $m_{\tilde\tau_1} < m_{\widetilde Z_1}$. For instance, if $\tan\beta = 35$ (this allows unification of the b and τ Yukawa couplings) with other parameters as in figure 7, only values of $m_{1/2}$ larger than 275 GeV are phenomenologically acceptable.

In table 10 we show a sample spectrum for this model. We choose $m_{1/2} = 300\,{\rm GeV}$, $\tan\beta = 35$ and other parameters as in figure 7. The spectrum is not unlike that in the mSUGRA framework with small m_0 so that sleptons are relatively light and squarks are lighter than the gluino. The chargino and $\widetilde Z_2$ almost exclusively decay via $\widetilde W_1 \to \tilde\tau_1\nu_\tau$ and $\widetilde Z_2 \to \tilde\tau_1\tau$, respectively, so that cascade decays of gluinos and squarks will lead to multi-jet plus multi-tau events, with (soft) leptons as daughters of the tau. Except for h, this scenario is probably beyond the reach of the Tevatron, but it should be straightforward to study $\tilde\ell_R$ and $\tilde\tau_1$, and probably also detect $\widetilde W_1$ and $\tilde\nu$, at the NLC. At the LHC a variety of signals should be present.

Parameter	Scale	Value
m_0	M_c	0
A_0	M_c	0
$m_{1/2}$	$M_{\rm GUT}$	300
g_5	$M_{\rm GUT}$	0.717
f_t	$M_{\rm GUT}$	0.534
$f_b = f_\tau$	$M_{\rm GUT}$	0.271
λ	$M_{\rm GUT}$	1
λ'	$M_{\rm GUT}$	0.1
$\tan\beta$	$M_{\rm weak}$	35
μ	$M_{\rm weak}$	< 0
$m_{\tilde g}$	$M_{\rm weak}$	737.2
$m_{\tilde u_L}$	$M_{\rm weak}$	668.5
$m_{\tilde d_R}$	$M_{\rm weak}$	633.1
$m_{\tilde t_1}$	$M_{\rm weak}$	482.8
$m_{\tilde b_1}$	$M_{\rm weak}$	541.5
$m_{\tilde\ell_L}$	$M_{\rm weak}$	258.6
$m_{\tilde\ell_R}$	$M_{\rm weak}$	210.0
$m_{\tilde\tau_1}$	$M_{\rm weak}$	143.3
$m_{\widetilde W_1}$	$M_{\rm weak}$	240.2
$m_{\widetilde Z_2}$	$M_{\rm weak}$	240.0
$m_{\widetilde Z_1}$	$M_{\rm weak}$	124.8
m_h	$M_{\rm weak}$	115.6
m_A	$M_{\rm weak}$	311.2.4
μ	$M_{\rm weak}$	-411.5

Table 10: Input and output parameters for the Minimal Gaugino Mediation model case study described in the text.

We have also examined how the mass spectrum changes with variation of the superpotential couplings λ and λ'. These couplings cannot be too large in order that they remain perturbative up to M_c. For variation in this range, we found that m_{10}(GUT) and m_5(GUT) were insensitive to the choice of these couplings, while the GUT scale values of $m_{\mathcal{H}_1}$ and $m_{\mathcal{H}_2}$ as well as A_t and A_b vary by about 20% over the entire range of λ and λ' that we examined. The weak-scale spectrum and the μ value are, however, insensitive to the choice of these parameters; this is presumably because $m_{1/2}$ is significantly larger than the scalar masses at the GUT scale, so that RG evolution between the GUT

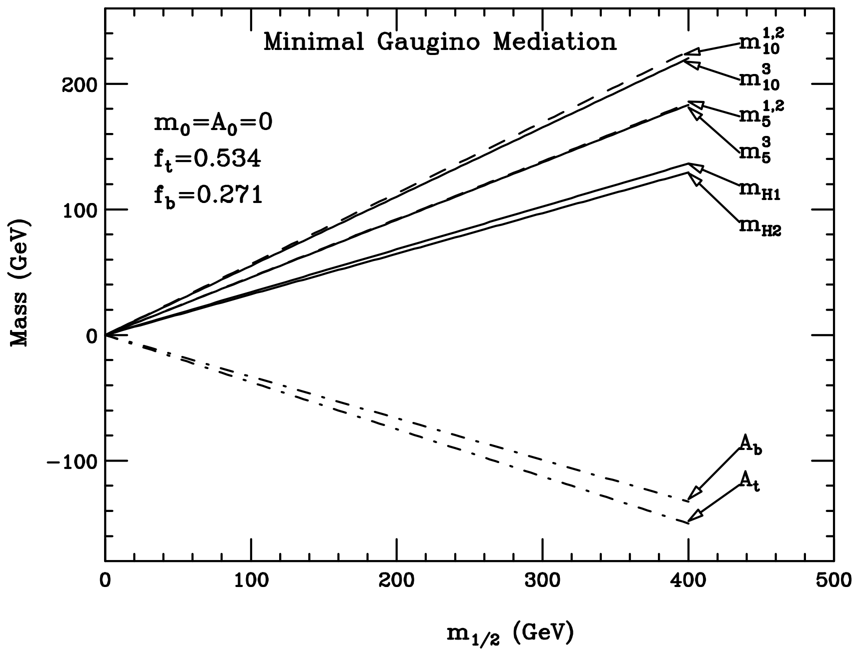

Figure 7: GUT scale values of SU(5) SSB masses in the minimal gaugino mediation model. We take $\tan\beta = 35$ and $\mu < 0$ to achieve b-τ Yukawa coupling unification. We take the compactification scale $M_c = 1 \times 10^{18}$ GeV. Models with $m_{1/2} < 275\,\text{GeV}$ lead to a breakdown in REWSB or a charged LSP.

and weak scales, rather than from m_0, makes the bulk of the contribution to scalar masses.

11. Models with non-universal soft terms due to 4-D superstring dynamics

Soft supersymmetry breaking terms obtained from $N = 1$ 4-dimensional superstrings, in general, exhibit non universality at the string scale [30, 31, 52, 53], a notable exception being when the dilaton is the dominant source of SUSY breaking. The soft supersymmetry breaking terms are determined by the Kähler potential K and the gauge kinetic functions f_a of the effective supergravity theory obtained from the string. The Kähler potential depends on the hidden sector fields, the dilaton S and the moduli T (there could be several), and the observable sector fields C_i, and it has the form,

$$K = -\log(S + S^*) + K_0(T, T^*) + \widetilde{K}_{ij}(T, T^*) C_i C^{*j} \,. \tag{11.1}$$

To avoid potential problems with FCNCs, we will assume that $\widetilde{K}_{ij} = \widetilde{K_i}\delta_{ij}$. In addition, the gauge kinetic function in any 4-dimensional superstring is given at tree level by

$$f_a = k_a S \,, \tag{11.2}$$

where k_a is the Kac-Moody level of the gauge factor G_a, with the entire group given by $G = \Pi_a G_a$. The Kac-Moody levels are usually taken $k_3 = k_2 = \frac{3}{5}k_1 = 1$. Beyond the tree level, f_a would in general also contain a dependence on the moduli fields.

Supersymmetry is broken when the auxiliary F-terms of the hidden sector fields acquire vacuum expectation values (vev). A convenient way to parametrize the vevs (in the case of one modulus) is as follows:

$$\begin{aligned} F^S &= \sqrt{3} C m_{3/2} K_{S\bar{S}}^{-1/2} \sin\theta e^{-i\gamma_S} \,, \\ F^T &= \sqrt{3} C m_{3/2} K_{T\bar{T}}^{-1/2} \cos\theta e^{-i\gamma_T} \,, \end{aligned} \tag{11.3}$$

where C is a constant defined by $C^2 = 1 + V_0/3m_{3/2}^2$, V_0 is the cosmological constant (the vev of the scalar potential), and $m_{3/2}$ is the gravitino mass. Here, $\sin\theta$ is the overlap between the goldstino and the fermionic component of the dilaton field. Therefore, $\sin\theta = 1$ in the limit where the SUSY breaking is completely due to the dilaton: i.e. $\langle F_S \rangle$ is the only relevant vev . The matrix $K_{n\bar{m}} \equiv \partial_n \partial_{\bar{m}} K$ is called the Kähler metric and γ_S and γ_T are possible complex phases.

The soft masses for scalar particles are determined by the Kähler potential in eq. (11.1) and are given by [31]

$$m_i^2 = 2m_{3/2}^2 (C^2 - 1) + m_{3/2}^2 C^2 (1 + N_i \cos^2\theta) \,, \tag{11.4}$$

with

$$N_i = \frac{-3(\log \widetilde{K}_i)_{T\bar{T}}}{(K_0)_{T\bar{T}}} \,.$$

We readily see that we can obtain non-universal scalar masses if $\cos\theta$ is different from zero. We mention that here we have for simplicity assumed that there is just one modulus field: multiple moduli are treated in ref. [52].

The gaugino masses are given by

$$M_a = \frac{1}{2}(\mathrm{Re} f_a)^{-1} F^m \partial_m f_a = \sqrt{3} C m_{3/2} (\mathrm{Re} f_a)^{-1} k_a \mathrm{Re} S e^{-i\gamma_S} \sin\theta \,, \tag{11.5}$$

where the gauge coupling constants are $\mathrm{Re} f_a = 1/g_a^2$. In the last equality, we have used the fact that (at tree level) the gauge kinetic function in eq. (11.2) depends only on the dilaton field S, so that the tree level gaugino masses are independent of the moduli sector. Model-dependent corrections to this may, however, be significant, particularly when dilaton contributions to SUSY breaking are small.

Expressions for A-parameters may also be found in ref. [31]. These depend on additional parameters, and generically also on the unknown phases γ_S and γ_T (as well as on additional direction cosines in the multi-moduli case). For the single modulus case, the form of A is given by,[10]

$$A_{ijk} = \sqrt{3} m_{3/2} C \left(e^{-i\gamma_S} \sin\theta + e^{-i\gamma_T} \omega_{ijk}(T, T^*) \cos\theta \right) , \tag{11.6}$$

[10]We have flipped the sign of A to conform to our convention where the soft trilinear term is written as $A_{ijk} f_{ijk} \tilde{C}_i \tilde{C}_j \tilde{C}_k$ in the lagrangean and not the scalar potential, with f_{ijk} being the corresponding superpotential coupling.

where ω_{ijk} depend on the Kähler and superpotentials. Fortunately, in many cases of interest, these model-dependent parameters either vanish or assume a simple form.

We should mention that these expressions for the soft-SUSY breaking masses and A-parameters are valid for these parameters renormalized at the string scale. As always, these have then to be evolved down to the weak scale for use in phenomenological analysis. We now consider some special cases to illustrate the forms of (string scale) non universality that may occur in this general framework.

11.1 Large-T limit of Calabi-Yau compactifications

Because of the complexity of the world-sheet instanton and sigma model contributions, the general form of the Kähler potential of generic Calabi-Yau $(2,2)$ compactifications is not known. The gauge group is $E_6 \times E_8$, with matter in the **27** dimensional representation of E_6. It is usual to analyze the large T (in practice $2-3 < |T| < 20-30$, large enough so that world sheet instanton contributions can be neglected, but not so large that string threshold corrections invalidate perturbation theory) limit of these theories. In this limit the Kähler potential takes a simple form [31]:

$$K = -\log(S+S^*) - 3\log(T+T^*) + \sum_i \frac{|C_i|^2}{T+T^*}\,, \tag{11.7}$$

and the gauge kinetic function is given by eq. (11.2) at tree level. In this case the gaugino mass is

$$m_{1/2} = \sqrt{3} C m_{3/2} \sin\theta e^{-i\gamma_S}\,, \tag{11.8}$$

while eq. (11.4) for the scalar masses reduces to,

$$m_0^2 = m_{3/2}^2 C^2 \sin^2\theta + 2m_{3/2}^2(C^2-1) \tag{11.9}$$

which simplifies even further if the cosmological constant vanishes ($C=1$). Notice that we find universality of soft scalar masses, even though we are not in the dilaton dominated SUSY breaking scenario.

In the $C=1$ case, we see that $|m_{1/2}| = \sqrt{3}m_0$, so that the gaugino mass always exceeds the scalar mass at the string scale. This relation obviously puts a significant constraint on SUSY phenomenology. Since this is a special case of the mSUGRA scenario whose phenomenological implications have been discussed at length in the literature, we will not mention this any further.

There are, however, arguments in the literature [54] that suggest that the observed cosmological constant (which is bounded to be smaller than $\sim (3\text{ meV})^4$) may not be directly connected to V_0; then, C could differ from unity, and the gaugino mass may (depending on the value of C and the goldstino angle θ) be even smaller than m_0, but for an appreciable effect, $C-1$ would have to deviate by many orders of magnitude[11] from the bound that would have resulted assuming V_0 was the observed cosmological constant.

Finally, in this limit, the parameters ω_{ijk} in eq. (11.6) vanish so that

$$A_{ijk} = \sqrt{3} m_{3/2} C e^{-i\gamma_S} \sin\theta \,.$$

In the single modulus large T case that we have been discussing, effects of the sigma-model loop contribution and the non-perturbative instanton contribution to the Kähler potential are known [55]. We still obtain universality of soft SUSY breaking parameters, with gaugino masses given by eq. (11.8) and scalar masses and the A-parameter (in the case $C = 1$) modified to,

$$m_0^2 = m_{3/2}^2 \left[1 - \cos^2\theta \left(1 - \Delta(T, T^*)\right)\right] \tag{11.10}$$

and

$$A = \sqrt{3} m_{3/2} \left[e^{-i\gamma_S} \sin\theta + \omega(T, T^*) e^{-i\gamma_T} \cos\theta\right] . \tag{11.11}$$

Here Δ and ω correspond to the sigma-model and instanton contributions (the latter are negligible): the numerical values of these are model dependent, but $\Delta \approx 0.4$ and $\omega = 0.17$ have been quoted [55] for a typical model. Notice that although these corrections do not lead to non universality, we lose the earlier prediction $m_{1/2} = \sqrt{3} m_0$: now, the soft scalar mass may even exceed the corresponding gaugino mass if $\cos^2\theta$ is sufficiently large.

11.2 General Calabi-Yau compactifications

There is no reason to believe that there is just a single modulus field T. In the multi-moduli case the parametrization of the vevs of the moduli in eq. (11.3) is modified to [52],

$$F^{T_i} = \sqrt{3} C m_{3/2} K_{T_i \bar{T}_i}^{-1/2} \cos\theta \Theta_i e^{-i\gamma_{T_i}} , \tag{11.12}$$

where we have assumed the Kähler metric to be diagonal to avoid any FCNC problems. Here Θ_i are direction cosines that parametrize the direction of the vev in moduli space. Indeed the more general case of an off-diagonal metric has also been examined in ref. [55] where a more general parametrization of the vevs of the moduli may be found. In this general case, the scalar masses are non-diagonal and the mass squared matrix assumes the form,

$$m_{ij}^2 = m_{3/2}^2 \left[\delta_{ij} - \cos^2\theta \left(\delta_{ij} - \Delta_{ij}(T_k, T_k^*)\right)\right] , \tag{11.13}$$

where Δ_{ij} depends on the moduli and on the direction of the vev in the moduli space. Notice that the model-dependent Δ_{ij} would be strongly constrained by experimental data on flavour mixing. We are, however, not aware of a realistic model in which such constraints may be analyzed. We also note that the presence of a (even diagonal) matrix Δ in eq. (11.13) would be a source of non universality of scalar masses.

[11] It should be appreciated that even $C = 1.1$ is an enormous value relative to the bound $C - 1 \lesssim 10^{-87}$ that we would get if we took V_0 to be related to the observed value of Λ.

11.3 Orbifold models with large threshold corrections

An example of such a model is the so-called O-I model discussed by Brignole et al. [31]. In orbifold compactifications the coefficient $\widetilde{K}_{ij}$ which determines the soft masses has the form $(T+T^*)^{n_i}$, where n_i is the modular weight of the matter field C_i. The Kähler potential is in this case:

$$K = -\log(S+S^*) - 3\log(T+T^*) + \sum_i |C_i|^2 (T+T^*)^{n_i} \,. \tag{11.14}$$

Gauge unification in good agreement with low-energy data is achieved by assigning the following modular weights for the massless fields: $n_Q = n_D = -1$, $n_U = -2$, $n_L = n_E = -3$, and $n_{H_d} + n_{H_u} = -5$ or -4, together with a large value for the modulus field, T $\approx$ 16, which then results in large threshold corrections. Under these conditions the gaugino masses are non-universal at the string scale:

$$\begin{aligned}
M_1 &= 1.18\sqrt{3}m_{3/2}\left[\sin\theta + 2.9\times 10^{-2}(B_1'/k_1)\cos\theta\right] ,\\
M_2 &= 1.06\sqrt{3}m_{3/2}\left[\sin\theta + 2.9\times 10^{-2}(B_2'/k_2)\cos\theta\right] ,\\
M_3 &= 1.00\sqrt{3}m_{3/2}\left[\sin\theta + 2.9\times 10^{-2}(B_3'/k_3)\cos\theta\right] ,
\end{aligned} \tag{11.15}$$

where $B_a' \equiv b_a' - k_a\delta_{GS}$ are given by $B_1' = -18 - k_1\delta_{GS}$, $B_2' = -8 - k_2\delta_{GS}$, and $B_3' = -6 - k_3\delta_{GS}$ if $n_{H_d} + n_{H_u} = -5$, and k_a as specified previously. Here, the parameter δ_{GS} is a model dependent negative integer and $m_{3/2}$ and θ are the gravitino mass and the goldstino angle as before. To obtain eqs. (11.15), it is assumed [31] that string threshold corrections lead to an apparent unification of the couplings at the "GUT scale" rather than at the string scale. Of course, since there is no GUT these couplings continue to evolve and diverge when evolved from the "GUT scale" to the one order of magnitude larger string scale. The coefficients in front of the gaugino mass formulae reflect just this difference in the gauge couplings at the string scale. In other words, if $\sin\theta = 1$, gaugino masses (while slightly different at the string scale) would be universal at $Q = M_{\rm GUT}$: non universality of GUT scale gaugino masses occurs only due to the loop correction proportional to $\cos\theta$ in eqs. (11.15). Finally, we note that if $n_{H_d} + n_{H_u} = -4$, the gaugino masses are obtained from eqs. (11.15) by modifying the coefficients B_i' to $B_1' = -17 - k_1\delta_{GS}$ and $B_2' = -7 - k_2\delta_{GS}$ while B_3' does not change.

The string-scale scalar masses and A-parameters depend on the modular weights, and (assuming zero cosmological constant) are given by,

$$\begin{aligned}
m_Q^2 = m_D^2 &= m_{3/2}^2\left[1 - (1-\delta_{GS}\times 10^{-3})^{-1}\cos^2\theta\right] ,\\
m_U^2 &= m_{3/2}^2\left[1 - 2(1-\delta_{GS}\times 10^{-3})^{-1}\cos^2\theta\right] ,\\
m_L^2 = m_E^2 &= m_{3/2}^2\left[1 - 3(1-\delta_{GS}\times 10^{-3})^{-1}\cos^2\theta\right] ,
\end{aligned} \tag{11.16}$$

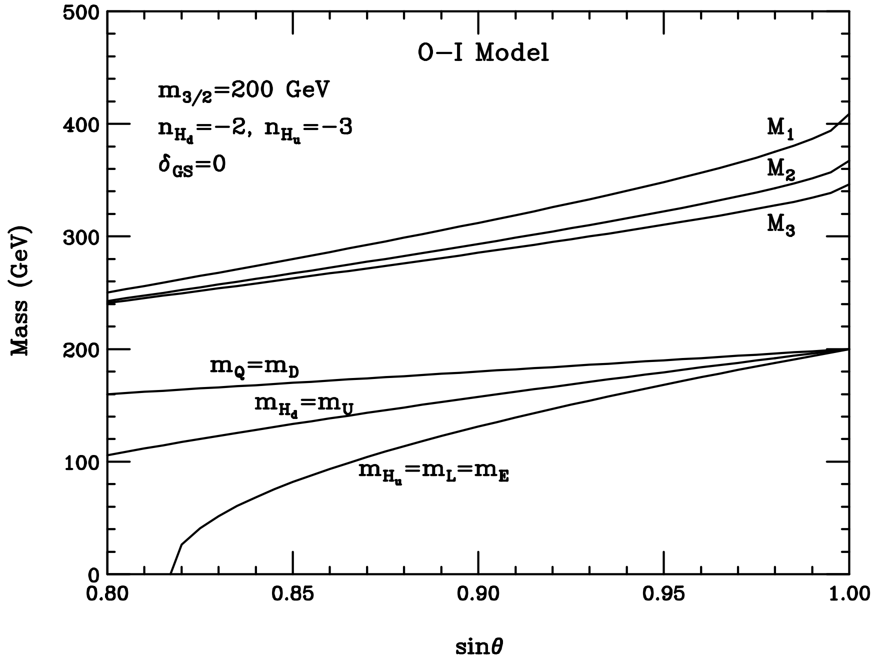

Figure 8: Soft SUSY breaking masses in the O-I superstring model, versus $\sin\theta$, for $m_{3/2} = 200\,\mathrm{GeV}$, $n_{H_D} + n_{H_u} = -5$ and $\delta_{GS} = 0$. The O-I model assumes $n_Q = n_D = -1$, $n_U = -2$, and $n_L = n_E = -3$.

and

$$A_{ijk} = \sqrt{3} m_{3/2} \sin\theta \pm m_{3/2} \cos\theta (1 - \delta_{GS} \times 10^{-3})^{-1/2} (3 + n_i + n_j + n_k)\,, \qquad (11.17)$$

where the terms with δ_{GS} come from radiative corrections, and the sign ambiguity reflects the possible relative phase between γ_S and γ_T (we take the A-parameters to be real). Note that if $\sin\theta = 1$, the scalar masses and A-parameters are universal at the string scale: RG evolution would then introduce a small non universality at M_{GUT}.

In figure 8 we plot different soft masses at the string scale as a function of $\sin\theta$ in the O-I model. There is a sign ambiguity since $\cos\theta$ could be negative. We have chosen $\cos\theta > 0$ and fixed $m_{3/2} = 200\,\mathrm{GeV}$, $\delta_{GS} = 0$, and, for the evaluation of gaugino masses, $n_{H_u} + n_{H_d} = -5$. We set the phases γ_S and γ_T to be zero. Scalar masses are universal in the dilaton dominated scenario and radiative corrections do not spoil this universality. On the contrary, gaugino masses are not universal at $\sin\theta = 1$, but as explained above, there is (approximate) universality at M_{GUT}. Values of $\cos\theta \gtrsim 1/\sqrt{3}$ ($\sin\theta \lesssim 0.8$) yield negative slepton soft squared masses and may be unacceptable;[12] hence the dilaton field is necessarily the most important

[12]It may be possible to have these squared masses negative at a high scale as long as they are positive near the weak scale.

Mass	O-I	O-II ($\sin\theta = 0.85$)	O-II ($\sin\theta = 0$)	O-II ($\sin\theta = 0.005$)
$m_{\tilde{g}}$	698	773	360	1693
$m_{\tilde{u}_L}$	633	701	4241	4436
$m_{\tilde{u}_R}$	602	677	4239	4409
$m_{\tilde{d}_R}$	607	673	4237	4400
$m_{\tilde{t}_1}$	442	523	2468	2637
$m_{\tilde{t}_2}$	651	708	3478	3672
$m_{\tilde{b}_1}$	583	650	3475	3669
$m_{\tilde{b}_2}$	608	674	4232	4394
$m_{\tilde{\ell}_L}$	223	292	4239	4289
$m_{\tilde{\ell}_R}$	132	223	4240	4272
$m_{\widetilde{W}_1}$	202	235	254	715
$m_{\widetilde{Z}_1}$	99	124	249	490
$m_{\widetilde{Z}_2}$	202	236	255	715
m_h	109	108	116	117
m_A	497	516	4506	4634
μ	422	416	1114	1379

Table 11: Model parameters and weak-scale sparticle masses in GeV for O-I and O-II models discussed in the text. For the O-I model, we take $\tan\beta = 4$, $n_{H_u} = -3$, $\sin\theta = 0.85$ and other parameters as in figure 8. For the O-II model, we fix $\tan\beta = 4$, $\delta_{GS} = -5$ and illustrate the results for the three cases discussed in the text. For both models we take $\cos\theta > 0$. We have fixed the string scale to be 4×10^{17} GeV.

source of SUSY breaking. Except close to the lowest acceptable values of θ, deviations from universality in the scalar sector are thus limited.

To facilitate simulation of such a scenario, we have introduced into ISAJET versions ≥ 7.50 the "SUSY Boundary Condition Scale" (SSBCSC keyword) option into ISAJET that allows the user to input a chosen scale $Q_{\max}$ up to which the MSSM is assumed to be valid. The values of SUSY breaking masses and A-parameters of the MSSM as given by any theory valid at the scale beyond $Q_{\max}$ would then be used as inputs to ISAJET, which would then evolve them down to the weak scale and generate SUSY events as usual. For the case at hand, $Q_{\max}$ would be the string scale, and the gaugino masses, scalar masses and A-parameters as given by eqs. (11.15)–(11.17), the boundary conditions for the RGE. We stress, however, that $Q_{\max}$ need not be larger than $M_{\rm GUT}$. For instance, in SO(10) models, $Q_{\max}$ would be the mass of the right-handed neutrino, or in E_6 models, the mass scale where the additional particles in the **27** dimensional representation and any extra Z' bosons all decouple, leaving the MSSM spectrum.

We give an example of the SUSY spectrum in the O-I scenario in table 11. In this example, we have fixed $\tan\beta = 4$, $\sin\theta = 0.85$ (with $\cos\theta > 0$) and have taken

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$n_{H_u} = -3$, with other parameters as in figure 8. Since the value of B depends on how μ is generated, we have treated $\tan\beta$ as a free parameter, and eliminated B in its favour, using the constraints given by radiative electroweak symmetry breaking. We fix the string scale to be 4×10^{17} GeV. Despite the fact that string scale slepton masses are considerably smaller than those of squarks (see figure 8), the spectrum is qualitatively very similar to that in the mSUGRA framework with $m_{\tilde{q}} \sim m_{\tilde{g}}$.

11.4 Orbifold models with small threshold corrections

In the O-I model, $\sin\theta$ was restricted to be large, so that the parameters of phenomenological interest were qualitatively similar to the mSUGRA scenario. To allow a wider range of $\sin\theta$ we consider a model where all the modular weights are -1. As noted in ref. [31] string threshold corrections cannot account for gauge coupling unification, which has then to be attributed to some different physics. Unlike the O-I model where a large value of $\mathrm{Re}\,T$ was needed to accommodate coupling constant unification, we will, following Brignole et al. [31] use $\mathrm{Re}\,T \approx 1.2$ and refer to this as the O-II model. As before, the gaugino masses are non-degenerate at the string scale (again, for $\sin\theta = 1$, these would be universal at "M_{GUT}") and given by:

$$\begin{aligned} M_1 &= 1.18\sqrt{3}m_{3/2}\left[\sin\theta + 4.6 \times 10^{-4}(B_1''/k_1)\cos\theta\right], \\ M_2 &= 1.06\sqrt{3}m_{3/2}\left[\sin\theta + 4.6 \times 10^{-4}(B_2''/k_2)\cos\theta\right], \\ M_3 &= 1.00\sqrt{3}m_{3/2}\left[\sin\theta + 4.6 \times 10^{-4}(B_3''/k_3)\cos\theta\right], \end{aligned} \tag{11.18}$$

with $B_1'' = 11 - k_1\delta_{GS}$, $B_2'' = 1 - k_2\delta_{GS}$, and $B_3'' = -3 - k_3\delta_{GS}$. On the other hand, the scalar masses ($V_0 = 0$) and A-parameters are all degenerate and equal to

$$m_Q^2 = m_D^2 = m_U^2 = m_L^2 = m_E^2 = m_{3/2}^2\left[1 - (1 - \delta_{GS} \times 10^{-3})^{-1}\cos^2\theta\right] \tag{11.19}$$

and

$$A_{ijk} = \sqrt{3}m_{3/2}\sin\theta\,, \tag{11.20}$$

at the string scale.

If $\sin\theta \sim 1$ the spectra should be the same as in the O-I model discussed previously. For smaller values of $\sin\theta$, the degeneracy in the string-scale scalar masses still remains. The most important difference between the two scenarios is that very small values of $\sin\theta$ are now permitted; i.e. the dilaton contribution need not necessarily dominate SUSY breaking. If $\sin\theta$ is very small so that the $\cos\theta$ terms are the dominant contributions to the gaugino mass, we see that (depending on the value of δ_{GS}) the GUT scale gluino mass may be much smaller than the corresponding electroweak gaugino masses. Indeed it is possible [56] to arrange scenarios where the gluino is the LSP[57]. The additional parameters also allow the possibility $M_1 \simeq M_2$ so that the lighter chargino and the two lighter neutralinos (and sometimes also the gluino) are all very degenerate. Such scenarios pose interesting experimental challenges [56].

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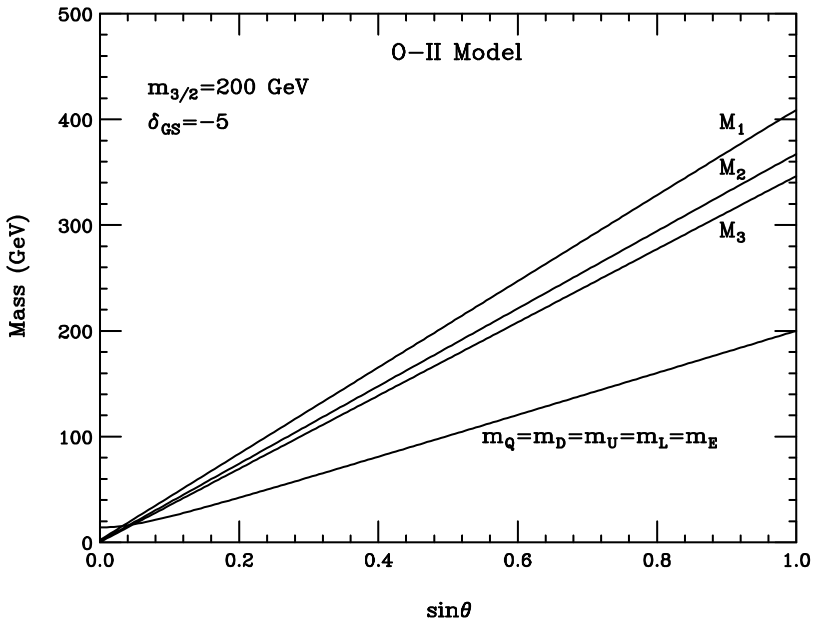

Figure 9: Soft SUSY breaking masses in the O-II superstring model, versus $\sin\theta$, for $m_{3/2} = 200\,\text{GeV}$ and $\delta_{GS} = -5$. The O-II model assumes $n_Q = n_D = n_U = n_L = n_E = -1$.

In figure 9 we illustrate the gaugino and scalar soft masses at the string scale as a function of $\sin\theta$ in the O-II model. Again, we take $m_{3/2} = 200\,\text{GeV}$, $\cos\theta > 0$, and ignore all phases. We choose $\delta_{GS} = -5$. The masses decrease as $\sin\theta$ decrease but they do not vanish at $\sin\theta = 0$ due to one-loop effects. Of course, for very small values of $\sin\theta$ phenomenological considerations require $m_{3/2}$ to be significantly larger. In the extreme case of moduli-dominated SUSY breaking, gaugino masses can be smaller than scalar masses, but generally speaking scalar masses are smaller than gaugino masses at the unification scale.

In the last three columns of table 11 we illustrate three examples of O-II model spectra. Again, we fix $\tan\beta = 4$, $\mu > 0$ and take $\delta_{GS} = -5$, to be in the region which can potentially yield [56] roughly equal masses for all the MSSM gauginos. First, we choose an O-II scenario with parameters close to those of the O-I model in the previous column: $m_{3/2} = 200\,\text{GeV}$ and $\sin\theta = 0.85$. This is the "typical" case for such a model. In this case, the $\sin\theta$ terms in eq. (11.18) completely dominate, and the resulting spectrum is again very similar to that in the mSUGRA framework (with $m_{\tilde{q}} \sim m_{\tilde{g}}$).

In the next column, we show a spectrum for the case $\sin\theta = 0$, the extreme case[13] of moduli-dominated SUSY breaking. Here, because of the small coefficient 4.6×10^{-4} in the expressions for gaugino masses, we have to choose $m_{3/2}$ to be

[13]It does not matter whether we take $\theta = 0$ or $\theta = \pi$ since the sign of the gaugino mass has no import for physics.

large. We fix $m_{3/2} = 60$ TeV. For this case, we have taken $m_t = 180\,\text{GeV}$, since we found that electroweak symmetry was not broken[14] for $m_t = 175\,\text{GeV}$. Since the (common) string-scale scalar mass is much bigger than the corresponding gaugino masses, the scalars are all roughly degenerate, and their spectrum is close to that of the corresponding mSUGRA spectrum with $m_{\tilde{q}} \sim m_{\tilde{\ell}} \gg m_{\tilde{g}}$ (i.e. $m_0 \gg m_{1/2}$). The gluinos, charginos and neutralino spectrum is quite different from that in the mSUGRA model: even though the lighter chargino and neutralinos are gaugino-like, $m_{\widetilde{W}_1} = m_{\widetilde{Z}_2} = m_{\widetilde{Z}_1} = 0.7 m_{\tilde{g}}$. This is because by choosing δ_{GS} we can adjust $M_1 : M_2 : M_3$ at the string scale. By a careful adjustment of parameters the gluino mass can even be brought closer to the chargino and neutralino masses. Experiments at the Fermilab Tevatron may be sensitive to this scenario.

To emphasize that the novel scenarios shown in ref. [56] obtain only for a very limited range of parameters, in the last column we show the spectrum for $\sin\theta = 0.005$ (with $\cos\theta > 0$), with all other parameters (including m_t) as for the $\sin\theta = 0$ case. We see that even for this tiny value of $\sin\theta$, the $\sin\theta$ terms in eq. (11.18) are comparable to (or even dominate) the $\cos\theta$ terms, and the spectrum is qualitatively different. While the sfermions are once again extremely heavy, the gluino, chargino and neutralino masses are now approximately as in the mSUGRA framework. Sparticle detection in this scenario would only be possible at the LHC. Our purpose in showing this (possibly unacceptably heavy) spectrum is only to emphasize the qualitative difference from the $\sin\theta = 0$ case. Of course, if $m_{3/2}$ is chosen to be 15 TeV, many more sparticles would be in the accessible range, but the spectrum would then be much like the canonical mSUGRA case with large m_0.

12. Models with non-universal soft terms due to M-theory dynamics

It was proposed that M-theory, i.e. , an 11-dimensional supergravity on a manifold where two E_8 gauge multiplets are restricted to the two 10-dimensional boundaries, is equivalent to the strong coupling limit of $E_8 \times E_8$ heterotic string theory [58]. It may be argued that M-theory is a better candidate than the weakly coupled string to explain low-energy physics and unification. After compactifying the 11-dimensional M-theory, a 4-dimensional effective theory emerges which can reconcile the reduced Planck scale $M_P \approx 2.4 \times 10^{18}\,\text{GeV}$, the grand unification scale $M_{\text{GUT}} \approx 3 \times 10^{16}\,\text{GeV}$, and α_{GUT}, in a way that the weakly coupled heterotic string theory

[14] The scalars start at a very large mass at the string scale, and the top Yukawa is not large enough to drive a Higgs mass squared eigenvalue negative at the scale $Q = \sqrt{m_{\tilde{t}_L} m_{\tilde{t}_R}}$ where the effective potential is evaluated in ISAJET [5]. We should mention that this is sensitive to the top mass radiative corrections that have been included [5] in ISAJET versions ≥ 7.48. These radiative corrections decrease the top Yukawa coupling by a few percent, and in this case, this is just sufficient to preclude electroweak symmetry breaking.

cannot. An interesting feature of the 4-dimensional effective SUGRA is that, in first approximation, the gauge kinetic function, the superpotential, and the Kähler potential do not change when moving from the weakly coupling heterotic string case to the M-theory case by changing the value of the modulus field.

12.1 One modulus case

Supersymmetry is broken when the auxiliary components of the dilaton field S and the modulus field T aquire non-zero vevs , as discussed in the last section. The low-energy effective supergravity theory [59] obtained from a specific Calabi-Yau compactification in M-theory is a Yang-Mills gauge theory with E_6 as the gauge group. The gauge kinetic function is given by,

$$f_{E_6} = S + \alpha T\,, \tag{12.1}$$

where α is an integer while the corresponding Kähler potential is

$$K = -\log(S+S^*) - 3\log(T+T^*) + \left[\frac{3}{T+T^*} + \frac{\alpha}{S+S^*}\right]\sum_i |C_i|^2\,, \tag{12.2}$$

where C_i again denote the observable fields. Adopting the same parametrization as in eq. (11.3) above, we find that with the Kähler potential of eq. (12.2) the soft SUSY breaking parameters are universal and given by,

$$m_{1/2} = \frac{\sqrt{3}Cm_{3/2}}{1+x}\left[\sin\theta e^{-i\gamma_S} + \frac{x}{\sqrt{3}}\cos\theta e^{-i\gamma_T}\right], \tag{12.3}$$

$$m_0^2 = m_{3/2}^2(3C^2-2) - \frac{3C^2 m_{3/2}^2}{(3+x)^2}\Big[x(6+x)\sin^2\theta + (3+2x)\cos^2\theta - -2\sqrt{3}\,x\sin\theta\cos\theta\cos(\gamma_S-\gamma_T)\Big], \tag{12.4}$$

and

$$A = \frac{\sqrt{3}Cm_{3/2}}{3+x}\left[(3-x)\sin\theta e^{-i\gamma_S} + \sqrt{3}x\cos\theta e^{-i\gamma_T}\right], \tag{12.5}$$

where

$$x \equiv \frac{\alpha(T+T^*)}{S+S^*}\,. \tag{12.6}$$

The range of x is $0 \le x \le 1$.

Note that in the weak coupling limit $x \to 0$, we recover from eqs. (12.3), (12.4) and (12.5) the gaugino and scalar masses as well as the A-parameter in the large T-limit of Calabi-Yau compactifications in eqs. (11.8) and (11.9), respectively.

In figure 10 we show the dependence on the goldstino angle of the universal gaugino and scalar masses, $m_{1/2}$ and m_0, respectively. We consider zero cosmological constant and three values of x. The strong coupling limit corresponds to $x = 1$ and for comparison $x = 0.5$ and $x = 0$ are also plotted.

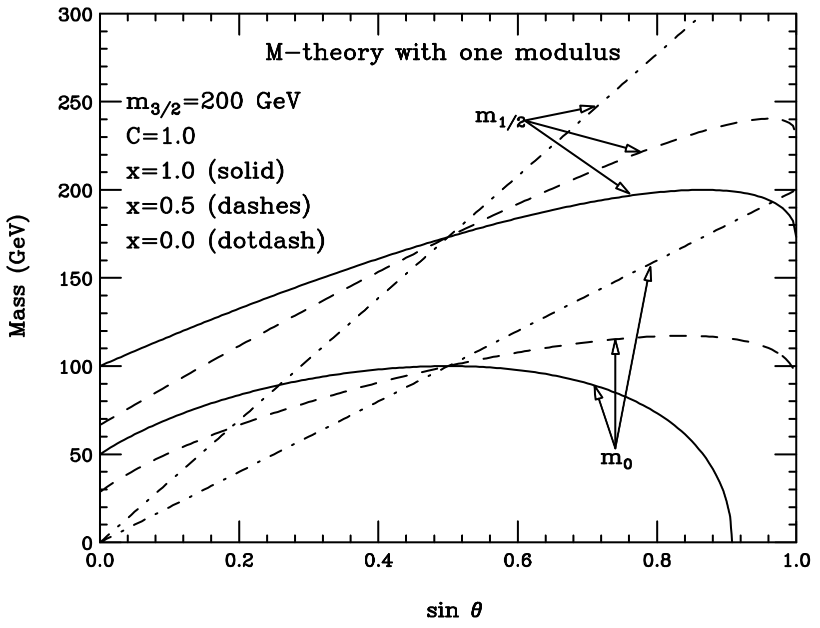

Figure 10: Universal gaugino ($m_{1/2}$) and scalar (m_0) masses as a function of $\sin\theta$ in M-theory with one modulus for three values of the parameter x and zero cosmological constant.

We remind the reader that the soft parameters obtained above are for an E_6 gauge theory. In order to obtain a realistic low-energy theory, we have to know how the symmetry group is reduced to the MSSM gauge group, which in turn will depend on the details of the theory at the high scale. It is possible that there may be additional TeV scale supermultiplets in the particle spectrum, or even extra gauge bosons [60]. Moreover, depending on how E_6 breaks to $\mathrm{SU}(3)\times\mathrm{SU}(2)\times\mathrm{U}(1)$, additional D-term contributions (see section 5) which break the universality of scalar masses may also be present.

12.2 Multi-moduli case

As before, the situation in the multi-moduli case can be more complicated. A toy example with three moduli fields and three observable fields has been considered in ref. [61]. The Kähler potential and gauge kinetic function of the effective theory is written as,

$$K = -\log(S+S^*) - \sum_{j=1}^{3}\log(T_j+T_j^*) + \left[2+\frac{2}{3}\sum_{j=1}^{3}\frac{\alpha_j(T_j+T_j^*)}{S+S^*}\right]\sum_{i=1}^{3}\frac{|C_i|^2}{T_i+T_i^*} \quad (12.7)$$

and

$$f_a = S + \sum_{i=1}^{3}\alpha_i T_i\,. \quad (12.8)$$

This then yields a universal mass for the gaugino and a universal A-parameter, but non-universal masses (and no mixing) for the scalars. While the gaugino and scalar masses as well as the A-parameter depend on the parameters and fields in the Kähler potential, the splitting δm^2 (between the scalars) appears to depend only on the orientation of the vevs of the auxiliary components of the moduli and on the goldstino angle θ. Since our focus is on sources of non universality in realistic scenarios that can potentially be of phenomenological interest, we merely note that multiple moduli could be a source of non universality of scalar masses, but do not exhibit results for this toy model here.

13. Concluding remarks

While weak-scale supersymmetry is a well-motivated idea, the physical principles that fix the multitude of SUSY breaking parameters are not known. Without any sparticle signals to provide clues, we do not have any guidance as to what these might be. The scale of this new physics may be as low as a few hundred TeV as in models with low-energy SUSY breaking mediated by gauge interactions, or as high as $M_{\rm GUT} - M_P$ as in frameworks where SUSY breaking is mainly mediated by gravity. Observable sparticle masses and mixing patterns, and via these weak-scale SUSY phenomenology, are determined by the physics behind SUSY breaking and how this is communicated to the observable sector. Turning this around, measurement of sparticle properties may provide clues about physics at energy scales that would be inaccessible to experiments in the foreseeable future.

Most early phenomenological analyses have been done within the framework of the mSUGRA model or the mSUGRA-motivated MSSM (where ad hoc relations between SSB parameters were assumed). In the last few years, phenomenological aspects of gauge-mediated SUSY breaking have also been examined in some detail. Both these models rest upon untested assumptions about physics at high energies. The good thing is that some of these assumptions will be directly testable if sparticles are discovered and their properties are measured [62]. Nevertheless, it seems worthwhile to look at other viable alternatives for physics at energy scales much beyond the weak scale, with a view to see if there are direct ramifications for sparticle signals in future experiments. A serious study of this would entail SUSY simulation at colliders in a wide variety of models with features different from the mSUGRA paradigm, which is characterized by universality of SSB parameters at a scale $Q \sim M_{\rm GUT}$.

Our study represents a first step in this direction. Here, we have surveyed a number of proposals for high-scale physics that lead to non universality of the SSB parameters in the MSSM, which we regard as the effective theory at a sufficently low mass scale. These range from relatively minor modifications of the mSUGRA SU(5) GUT model, where, e.g. unification of scalar masses and A-parameters is assumed to occur at M_P (so that RG evolution induces some non universality at $Q = M_{\rm GUT}$), to major modifications involving conceptually new ideas for high-scale physics (new hypercolour interactions, string physics) or the mediation of SUSY breaking to the

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observable sector (anomaly mediated SUSY breaking, gaugino mediated SUSY breaking). Other proposals that fall somewhat between these two extremes include models with larger unifying groups that naturally have additional non-universal contributions to scalar masses, or models where special boundary conditions on SSB parameters lead to unusual RG evolution and non-degeneracy of sparticle masses. For each of these scenarios, we have outlined the underlying physical ideas, delineated the parameter space in terms of which SUSY phenomenology might be analyzed, and discussed SUSY event generation using the simulation program ISAJET [5]. A variety of improvements to the ISAJET program have been made to allow event generation in the models discussed in this paper. These improvements are characterized by ISAJET keyword inputs, including NUSUG_i for non-universal masses, SUGRHN for models with a right-handed neutrino contribution, such as SO(10), SSBCSC for user choice as to when the MSSM becomes valid, and AMSB for anomaly-mediated SUSY breaking models. Where possible, we present sample spectra, and allude to the important phenomenological differences from the reference mSUGRA framework.

A detailed phenomenological analysis of each one of these scenarios is beyond the scope of the present work. Our hope though is that this study will facilitate and spur such analyses. Except for unusual cases where extreme degeneracies between sparticle masses result [56], we do not expect the reach of various future facilities (expressed in terms of physical sparticle masses) to qualitatively differ between the various scenarios. However, a careful examination of these will help us assess what we can hope to learn about high-scale physics if sparticles are discovered and their properties measured. Careful examination of physical implications of a variety of viable alternatives for the underlying theory will also help increase our understanding of the sort of analyses that might be needed to discriminate between these. In view of the potential pay-off, we believe that such studies will be very worthwhile.

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