SU(5) x U(1): A String Paradigm of a TOE and its Experimental Consequences

JORGE L. LOPEZ\(^{(a),(b)}\), D. V. NANOPoulos\(^{(a),(b),(c)}\), and A. ZICHICHI\(^{(d)}\)

\(^{(a)}\)Center for Theoretical Physics, Department of Physics, Texas A&M University
College Station, TX 77843–4242, USA
\(^{(b)}\)Astroparticle Physics Group, Houston Advanced Research Center (HARC)
The Woodlands, TX 77381, USA
\(^{(c)}\)CERN, Theory Division, 1211 Geneva 23, Switzerland
\(^{(d)}\)CERN, 1211 Geneva 23, Switzerland

ABSTRACT

We present a string-inspired/derived supergravity model based on the flipped $SU(5) \times U(1)$ structure supplemented by a minimal set of additional matter representations such that unification occurs at the string scale ($\sim 10^{18}$ GeV). This model is complemented by two string supersymmetry breaking scenarios: the $SU(N,1)$ no-scale supergravity model and a dilaton-induced supersymmetry breaking scenario. Both imply universal soft supersymmetry breaking parameters: $m_0 = 0, A = 0$ and $m_0 = \frac{1}{\sqrt{3}} m_{1/2}, A = -m_{1/2}$ respectively. In either case the models depend on only three parameters: $m_t, \tan \beta$, and $m_{\tilde{g}}$. We present a comparative study of the sparticle and Higgs spectra of both models and conclude that even though both can be partially probed at the Tevatron, LEP II, and HERA, a larger fraction of the parameter space of the no-scale model is actually accessible. In both cases there is a more constrained version which allows to determine $\tan \beta$ in terms of $m_t, m_{\tilde{g}}$. In the strict no-scale case we find that the value of $m_t$ determines the sign of $\mu$ ($\mu > 0 : m_t \lesssim 135$ GeV, $\mu < 0 : m_t \gtrsim 140$ GeV) and whether the lightest Higgs boson mass is above or below 100 GeV. In the more constrained version of the dilaton scenario, $\tan \beta \approx 1.4 - 1.6$ and $m_t \lesssim 155$ GeV, $61$ GeV $\lesssim m_h \lesssim 91$ GeV follow. Thus, continuing Tevatron top-quark searches and LEP II Higgs searches could probe this restricted scenario completely.

\(^{1}\)To appear in the Proceedings of the INFN Eloisatron Project 26th Workshop “From Superstrings to Supergravity”, Erice, Italy, Dec. 5-12, 1992; D. Nanopoulos and A. Zichichi conference speakers.
1 Introduction

The ultimate unification of all particles and interactions has string theory as the best candidate. If this theory were completely understood, we would be able to show that string theory is either inconsistent with the low-energy world or supported by experimental data. Since our present knowledge of string theory is at best fragmented and certainly incomplete, it is important to consider models which incorporate as many stringy ingredients as possible. The number of such models is expected to be large, however, the basic ingredients that such “string models” should incorporate fall into few categories: (i) gauge group and matter representations which unify at a calculable model-dependent string unification scale; (ii) a hidden sector which becomes strongly interacting at an intermediate scale and triggers supersymmetry breaking with vanishing vacuum energy and hierarchically small soft supersymmetry breaking parameters; (iii) acceptable high-energy phenomenology, e.g., gauge symmetry breaking to the Standard Model (if needed), not-too-rapid proton decay, decoupling of intermediate-mass-scale unobserved matter states, etc.; (iv) radiative electroweak symmetry breaking; (v) acceptable low-energy phenomenology, e.g., reproduce the observed spectrum of quark and lepton masses and the quark mixing angles, sparticle and Higgs masses not in conflict with present experimental bounds, not-too-large neutralino cosmological relic density, etc.

All the above are to be understood as constraints on potentially realistic string models. Since some of the above constraints can be independently satisfied in specific models, the real power of a string model rests in the successful satisfaction of all these constraints within a single model.

String model-building is at a state of development where large numbers of models can be constructed using various techniques (so-called formulations) \[1\]. Such models provide a gauge group and associated set of matter representations, as well as all interactions in the superpotential, the Kähler potential, and the gauge kinetic function. The effective string supergravity can then be worked out and thus all the above constraints can in principle be enforced. In practice this approach has never been followed in its entirety: sophisticated model-building techniques exist which can produce models satisfying constraints (i), (iii), (iv) and part of (v); detailed studies of supersymmetry breaking triggered by gaugino condensation have been performed for generic hidden sectors; and extensive explorations of the soft-supersymmetry breaking parameter space satisfying constraints (iii), (iv), and (v) have been conducted.

In searching for good string model candidates, we are faced with two kinds of choices to be made: the choice of the gauge and matter content of the model, and the choice of the supersymmetry breaking mechanism. Fortunately, a string theory theorem provides significant enlightenment regarding the first choice: models whose gauge groups are constructed from level-one Kac-Moody algebras do not allow adjoint or higher representations in their spectra \[2\]. This implies that the traditional GUT groups (\(SU(5)\), \(SO(10)\), \(E_6\)) are excluded since the GUT symmetry would remain unbroken. Exceptions to this theorem exist if one uses the technically complicated higher-level Kac-Moody algebras \[3\], but these models are beset with
If one imposes the aesthetic constraint of unification of the Standard Model non-abelian gauge couplings, then flipped $SU(5) \times U(1)$ \cite{4, 5, 6, 7} emerges as the prime candidate, as we shortly discuss. String models without non-abelian unification, such as the standard-like models of Refs. \cite{8, 9} and the Pati-Salam–like model of Ref. \cite{10} possess nonetheless gauge coupling unification at the string scale, even though no larger structure is revealed past this scale. However, the degree of phenomenological success which some of these models enjoy, usually rests on some fortuitous set of vanishing couplings which are best understood in terms of remnants of higher symmetries.

Besides the very economic GUT symmetry breaking mechanism in flipped $SU(5)$ \cite{1, 3} – which allows it to be in principle derivable from superstring theory \cite{3} – perhaps one of the more interesting motivations for considering such a unified gauge group is the natural avoidance of potentially dangerous dimension-five proton decay operators \cite{7}. In Ref. \cite{11} we constructed a supergravity model based on this gauge group, which has the additional property of unifying at a scale $M_U = \mathcal{O}(10^{18})$ GeV, as expected to occur in string-derived versions of this model \cite{12}. As such, this model constitutes a blueprint for string model builders. In fact, in Ref. \cite{13} one such model was derived from string and served as inspiration for the field theory model in Ref. \cite{11}. The string unification scale should be contrasted with the naive unification scale, $M_U = \mathcal{O}(10^{16}$ GeV), obtained by running the Standard Model particles and their superpartners to very high energies. This apparent discrepancy of two orders of magnitude \cite{13} creates a gap which needs to be bridged somehow in string models. It has been shown \cite{13} that the simplest solution to this problem is the introduction in the spectrum of heavy vector-like particles with Standard Model quantum numbers. The minimal such choice \cite{14}, a quark doublet pair $Q, \bar{Q}$ and a $1/3$-charge quark singlet pair $D, \bar{D}$, fit snugly inside a $10, \overline{10}$ pair of flipped $SU(5)$ representations, beyond the usual $3 \cdot (10 + \overline{5} + 1)$ of matter and $10, \overline{10}$ of Higgs.

In this model, gauge symmetry breaking occurs due to vacuum expectation values (vevs) of the neutral components of the $10, \overline{10}$ Higgs representations, which develop along flat directions of the scalar potential. There are two known ways in which these vevs (and thus the symmetry breaking scale) could be determined:

(i) In the conventional way, radiative corrections to the scalar potential in the presence of soft supersymmetry breaking generate a global minimum of the potential for values of the vevs slightly below the scale where supersymmetry breaking effects are first felt in the observable sector \cite{13}. If the latter scale is the Planck scale (in a suitable normalization) then $M_U \sim M_{Pl}/\sqrt{8\pi} \sim 10^{18}$ GeV.

(ii) In string-derived models a pseudo $U_A(1)$ anomaly arises as a consequence of truncating the theory to just the massless degrees of freedom, and adds a contribution to its $D$-term, $D_A = \sum q^A \langle \phi_i \rangle^2 + \epsilon$, with $\epsilon = g^2 \text{Tr} U_A(1)/192\pi^2 \sim (10^{18}$ GeV$)^2$ \cite{17}. To avoid a huge breaking of supersymmetry we need to demand $D_A = 0$ and therefore the fields charged under $U_A(1)$ need to get suitable vevs. Among these one generally finds the symmetry breaking Higgs fields, and thus $M_U \sim 10^{18}$ GeV follows.

In general, both these mechanisms could produce somewhat lower values of $M_U$. However, $M_U > 10^{16}$ GeV is necessary to avoid too rapid proton decay due
to dimension-six operators \[18\]. In these more general cases the \(SU(5)\) and \(U(1)\) gauge couplings would not unify at \(M_U\) (only \(\alpha_2\) and \(\alpha_3\) would), although they would eventually “superunify” at the string scale \(M_{SU} \sim 10^{18}\) GeV. To simplify matters, below we consider the simplest possible case of \(M_U = M_{SU} \sim 10^{18}\) GeV. We also draw inspiration from string model-building and regard the Higgs mixing term \(\mu h\bar{h}\) as a result of an effective higher-order coupling \[19, 20, 21\], instead of as a result of a light singlet field getting a small vev (i.e., \(\lambda h\bar{h}\phi \rightarrow \lambda \langle \phi \rangle h\bar{h}\)) as originally considered \[5, 7\]. An additional contribution to \(\mu\) is also generically present in supergravity models \[22, 21, 23\].

The choice of supersymmetry breaking scenario is less clear. Below we show that the phenomenologically acceptable choices basically fall in two categories:

1. The no-scale ansatz \[24\], which ensures the vanishing of the (tree-level) cosmological constant even after supersymmetry breaking. This framework also arises in the low-energy limit of superstring theory \[23\]. In a theory which contains heavy fields, the minimal no-scale structure \(SU(1, 1)\) \[26\] is generalized to \(SU(N, 1)\) \[27\] which implies that the scalar fields do not feel the supersymmetry breaking effects. In practice this means that the universal scalar mass \((m_0)\) and the universal cubic scalar coupling \((A)\) are set to zero. The sole source of supersymmetry breaking is the universal gaugino mass \((m_{1/2})\), i.e.,

\[
m_0 = 0, \quad A = 0. \tag{1}
\]

2. The dilaton \(F\)-term scenario, which also leads to universal soft supersymmetry breaking parameters \[23\]

\[
m_0 = \frac{1}{\sqrt{3}} m_{1/2}, \quad A = -m_{1/2}. \tag{2}
\]

In either case, after enforcement of the above constraints, the low-energy theory can be described in terms of just three parameters: the top-quark mass \((m_t)\), the ratio of Higgs vacuum expectation values \((\tan \beta)\), and the gluino mass \((m_{\tilde{g}} \propto m_{1/2})\). Therefore, measurement of only two sparticle or Higgs masses would determine the remaining thirty. Moreover, if the hidden sector responsible for these patterns of soft supersymmetry breaking is specified, the gravitino mass will also be determined and the supersymmetry breaking sector of the theory will be completely fixed.

In sum, we see basically two unified string supergravity models emerging as good candidates for phenomenologically acceptable string models, both of which include a flipped \(SU(5)\) observable gauge group supplemented by matter representations in order to unify at the string scale \(M_U \sim 10^{18}\) GeV \[13, 16\], and supersymmetry breaking is parametrized by either of the scenarios in Eqs. \(1, 2\).

We should remark that a real string model will include a hidden sector in addition to the observable sector discussed in what follows. The model presented here tacitly assumes that such hidden sector is present and that it has suitable properties. For example, the superpotential in Eq. \(1\) below, in a string model will receive
contributions from cubic and higher-order terms, with the latter generating effective observable sector couplings once hidden sector matter condensates develop \[19\]. The hidden sector is also assumed to play a fundamental role in triggering supersymmetry breaking via \textit{e.g.}, gaugino condensation. This in turn would make possible the mechanism for gauge symmetry breaking discussed above. Probably the most important constraint on this sector of the theory is that it should yield one of the two supersymmetry breaking scenarios outlined above.

This paper is organized as follows. In Sec. 2 we present the string-inspired model with all the model-building details which determine in principle the masses of the new heavy vector-like particles. We also discuss the question of the possible reintroduction of dangerous dimension-five proton decay operators in this generalized model. We then impose the constraint of flipped \(SU(5)\) unification and string unification to occur at \(M_U = 10^{18}\) GeV to deduce the unknown masses. In Sec. 3 we discuss the various supersymmetry breaking scenarios. In Sec. 4 we consider the experimental predictions for all the sparticle and one-loop corrected Higgs boson masses in these models, and deduce several simple relations among the various sparticle masses. In Sec. 5 we repeat this analysis for special more constrained cases of the chosen supersymmetry breaking scenario. In Sec. 6 we discuss the prospects for experimental detection of these particles at Fermilab, LEPI,II, and HERA. Finally, in Sec. 7 we summarize our conclusions.

2 The Model: Gauge-Matter Structure and Properties

The model we consider is a generalization of that presented in Ref. [3], and contains the following flipped \(SU(5)\) fields:

1. three generations of quark and lepton fields \(F_i, \bar{f}_i, l^c_i, i = 1, 2, 3\);
2. two pairs of Higgs 10,\(\overline{10}\) representations \(H_i, \bar{H}_i, i = 1, 2\);
3. one pair of “electroweak” Higgs 5,\(\overline{5}\) representations \(h, \bar{h}\);
4. three singlet fields \(\phi_{1,2,3}\).

Under \(SU(3) \times SU(2)\) the various flipped \(SU(5)\) fields decompose as follows:

\[
F_i = \{Q_i, d^c_i, \nu^c_i\}, \quad \bar{f}_i = \{L_i, u_i^c\}, \quad l^c_i = e^c_i,
\]
\[
H_i = \{Q_H, d_H^c, \nu_{H_i}\}, \quad \bar{H}_i = \{Q_{\bar{H}_i}, d_{\bar{H}_i}, \nu_{\bar{H}_i}\},
\]
\[
h = \{H, D\}, \quad \bar{h} = \{\bar{H}, \bar{D}\}.
\]

The most general effective\(^2\) superpotential consistent with \(SU(5) \times U(1)\) symmetry is given by

\[
W = \lambda^y_{ij} F_i F_j h + \lambda^y_{ij} F_i \bar{f}_j \bar{h} + \lambda^y_{ij} \bar{f}_i l^c_j h + \mu h \bar{h} + \lambda^y_{ij} H_i H_j h + \lambda^y_{ij} \bar{H}_i \bar{H}_j \bar{h}
\]

\(^2\)To be understood in the string context as arising from cubic and higher order terms \[28, 13\].
Symmetry breaking is effected by non-zero vevs $\langle \nu_{H}^{c} \rangle = V_i$, $\langle \nu_{H}^{c} \rangle = \bar{V}_i$, such that $V_1^2 + V_2^2 = \bar{V}_1^2 + \bar{V}_2^2$.

### 2.1 Higgs doublet and triplet mass matrices

The Higgs doublet mass matrix receives contributions from $\mu H \bar{H}$ and $\lambda_{2}^{ij} H_i \bar{f}_j h \rightarrow \lambda_{2}^{ij} V_i L_j H$. The resulting matrix is

$$
\mathcal{M}_2 = \begin{pmatrix}
\frac{\mu}{H} & \lambda_{1}^{ij} V_i \\
\lambda_{2}^{ij} V_i & \lambda_{2}^{ij} V_i \\
\lambda_{2}^{ij} V_i & \lambda_{2}^{ij} V_i
\end{pmatrix}.
$$

To avoid fine-tunings of the $\lambda_{2}^{ij}$ couplings we must demand $\lambda_{2}^{ij} \equiv 0$, so that $H$ remains light.

The Higgs triplet matrix receives several contributions: $\mu H \bar{H} \rightarrow \mu D \bar{D}$; $\lambda_{2}^{ij} H_i \bar{f}_j h \rightarrow \lambda_{2}^{ij} V_i d^c_i D$; $\lambda_{2}^{ij} H_i \bar{H}_j h \rightarrow \lambda_{2}^{ij} \bar{V}_i d^c_i \bar{D}$; $\lambda_{2}^{ij} \bar{H} \bar{H} \rightarrow \lambda_{2}^{ij} \bar{V}_i d^c_i \bar{D}$; $\lambda_{2}^{ij} \bar{H} \bar{H} \rightarrow \lambda_{2}^{ij} \bar{V}_i d^c_i \bar{D}$. The resulting matrix is

$$
\mathcal{M}_3 = \begin{pmatrix}
\bar{D} & d^c_{H_1} & d^c_{H_2} & d^c_{i_1} & d^c_{i_2} & d^c_{i_3} \\
\lambda_{2}^{ij} V_i & \lambda_{1}^{ij} V_i & \lambda_{2}^{ij} V_i & \lambda_{i_1}^{ij} V_i & \lambda_{i_2}^{ij} V_i & \lambda_{i_3}^{ij} V_i
\end{pmatrix}.
$$

Clearly three linear combinations of $\{\bar{D}, d^c_{H_1,2}, d^c_{i1,2,3}\}$ will remain light. In fact, such a general situation will induce a mixing in the down-type Yukawa matrix $\lambda_{i}^{ij} F_i F_j h \rightarrow \lambda_{i}^{ij} Q_i d^c_j H$, since the $d^c_i$ will need to be re-expressed in terms of these mixed light eigenstates.\(^{[4]}\) This low-energy quark-mixing mechanism is an explicit realization of the general extra-vector-abeysance (EVA) mechanism of Ref. \([29]\). As a first approximation though, in what follows we will set $\lambda_{1}^{ij} = 0$, so that the light eigenstates are $d^c_{i1,2,3}$.

---

\(^{[3]}\)The zero entries in $\mathcal{M}_3$ result from the assumption $\langle \phi_k \rangle = 0$ in $\lambda_{6}^{ijk} F_i \bar{H}_j \phi_k$.

\(^{[4]}\)Note that this mixing is on top of any structure that $\lambda_{2}^{ij}$ may have, and is the only source of mixing in the typical string model-building case of a diagonal $\lambda_2$ matrix.
2.2 Neutrino see-saw matrix

The see-saw neutrino matrix receives contributions from: $\lambda^{ij}_{2} F_i \bar{f_j} h \rightarrow m^i_{\nu} \nu^c_j; \lambda^{ijk}_6 F_i \bar{h_j} \phi_k \rightarrow \lambda^{ijk}_6 \bar{\nu^c_j} \phi_k; \mu^{ij} \phi_i \phi_j$. The resulting matrix is:

$$M_{\nu} = \begin{pmatrix} \nu_j & \nu^c_j & \phi_j \\ m^i_{\nu} & 0 & 0 \\ \phi_i & \lambda^{ijk}_6 \bar{V}_k & \mu^{ij} \end{pmatrix}. \quad (9)$$

2.3 Numerical scenario

To simplify the discussion we will assume, besides $\lambda^{ij}_1 = \lambda^{ij}_2 \equiv 0$, that

$$\lambda^{ij}_4 = \delta^{ij} \lambda^{(i)}_4, \quad \lambda^{ij}_5 = \delta^{ij} \lambda^{(i)}_5, \quad \lambda^{ij}_6 = \delta^{ij} \delta^{ik} \lambda^{(i)}_6, \quad \mu^{ij} = \delta^{ij} \mu_i, \quad w^{ij} = \delta^{ij} w_i. \quad (10, 11)$$

These choices are likely to be realized in string versions of this model and will not alter our conclusions below. In this case the Higgs triplet mass matrix reduces to

$$M_3 = \begin{pmatrix} D & d^c_{H_1} & d^c_{H_2} \\ \mu & \lambda^{(1)}_4 V_1 & \lambda^{(2)}_4 V_2 \\ \lambda^{(1)}_5 V_1 & 0 & \mu \\ \lambda^{(2)}_5 V_1 & 0 & w_2 \end{pmatrix}. \quad (12)$$

Regarding the (3, 2) states, the scalars get either eaten by the $X, Y$ $SU(5)$ heavy gauge bosons or become heavy Higgs bosons, whereas the fermions interact with the $\tilde{X}, \tilde{Y}$ gauginos through the following mass matrix:

$$M_{(3,2)} = \begin{pmatrix} Q_{H_1} & Q_{H_2} & \tilde{Y} \\ Q_{H_1} & Q_{H_2} & 0 \\ \tilde{X} & 0 & g_5 V_1 \\ 0 & g_5 V_2 & 0 \end{pmatrix}. \quad (13)$$

The lightest eigenvalues of these two matrices (denoted generally by $d^c_H$ and $Q_H$ respectively) constitute the new relatively light particles in the spectrum, which are hereafter referred to as the “gap” particles since with suitable masses they bridge the gap between unification masses at $10^{16}$ GeV and $10^{18}$ GeV.

---

5We neglect a possible higher-order contribution which could produce a non-vanishing $\nu^c_i \nu^c_j$ entry.  
6In Ref. the discrete symmetry $H_1 \rightarrow -H_1$ was imposed so that these couplings automatically vanish when $H_2, \bar{H}_2$ are not present. This symmetry (generalized to $H_i \rightarrow -H_i$) is not needed here since it would imply $w^{ij} \equiv 0$, which is shown below to be disastrous for gauge coupling unification.
Guided by the phenomenological requirement on the gap particle masses, \( M_{Q_H} \gg M_{d_H} \) \cite{16}, we consider the following explicit numerical scenario

\[
\lambda_i^{(2)} = \lambda_5^{(2)} = 0, \quad V_1, \tilde{V}_1, V_2, \tilde{V}_2 \sim V \gg w_1 \gg w_2 \gg \mu, \tag{14}
\]

which would need to be reproduced in a viable string-derived model. From Eq. (12) we then get

\[
M_{d_H^2} = M_{d_H^2} = w_2, \quad \text{and all other mass eigenstates} \sim V. \quad \text{Furthermore,} \quad \mathcal{M}_{(3,2)} \quad \text{has a characteristic polynomial}
\]

\[
\lambda^3 - \lambda^2 (w_1 + w_2) - \lambda (2V^2 - w_1 w_2) + (w_1 + w_2)V^2 = 0,
\]

which has two roots of \( \mathcal{O}(V) \) and one root of \( \mathcal{O}(w_1) \). The latter corresponds to \( \sim (Q_H^1 - Q_H^2) \) and \( \sim (Q_H^2 - Q_H^3) \). In sum then, the gap particles have masses \( M_{Q_H} \sim w_1 \) and \( M_{d_H} \sim w_2 \), whereas all other heavy particles have masses \( \sim V \).

The see-saw matrix reduces to

\[
\mathcal{M}_\nu = \begin{pmatrix}
\nu_i & \nu_i^c & \phi_i \\
\nu_i^c & \mu_i & 0 \\
\phi_i & 0 & \lambda(i)\tilde{V}_i
\end{pmatrix}, \tag{15}
\]

for each generation. The physics of this see-saw matrix has been discussed in Ref. \cite{30} and more generally in Ref. \cite{31}, where it was shown to lead to an interesting amount of hot dark matter (\( \nu_\tau \)) and an MSW-effect (\( \nu_\mu, \nu_e \)) compatible with all solar neutrino data. Moreover, the out-of-equilibrium decays of the \( \nu^c \) “flipped neutrino” fields in the early Universe induce a lepton number asymmetry which is later processed into a baryon number asymmetry by non-perturbative electroweak processes \cite{32,31}. All these phenomena can occur in the same region of parameter space.

### 2.4 Proton decay

The dimension-six operators mediating proton decay in this model are highly suppressed due to the large mass of the \( X,Y \) gauge bosons (\( \sim M_U = 10^{18} \text{ GeV} \)). Higgsino mediated dimension-five operators exist and are naturally suppressed in the minimal model of Ref. \cite{5}. The reason for this is that the Higgs triplet mixing term \( \mu h\tilde{h} \to \mu D\tilde{D} \) is small (\( \mu \approx M_Z \)), whereas the Higgs triplet mass eigenstates obtained from Eq. (8) by just keeping the \( 2 \times 2 \) submatrix in the upper left-hand corner, are always very heavy (\( \sim V \)). The dimension-five mediated operators are then proportional to \( \mu/V^2 \) and thus the rate is suppressed by a factor or \( \mu/V^2 \ll 1 \) relative to the unsuppressed case found in the standard \( SU(5) \) model.

In the generalized model presented here, the Higgs triplet mixing term is still \( \mu D\tilde{D} \). However, the exchanged mass eigenstates are not necessarily all very heavy. In fact, above we have demanded the existence of a relatively light (\( \sim w_1 \)) Higgs triplet state (\( d_H^i \)). In this case the operators are proportional to \( \mu \alpha_i \tilde{\alpha}_i / M_i^2 \), where \( M_i \) is the mass of the \( i \)-th exchanged eigenstate and \( \alpha_i, \tilde{\alpha}_i \) are its \( D, \tilde{D} \) admixtures. In the scenario described above, the relatively light eigenstates (\( d_H^1, d_H^2 \)) contain no \( D, \tilde{D} \) admixtures, and the operator will again be \( \propto \mu/V^2 \).
Note however that if conditions (14) (or some analogous suitability requirement) are not satisfied, then diagonalization of $M_3$ in Eq. (12) may re-introduce a sizeable dimension-five mediated proton decay rate, depending on the value of the $\alpha_i$, $\bar{\alpha}_i$ coefficients. To be safe one should demand \[33, 34\]

$$\frac{\mu \alpha_i \bar{\alpha}_i}{M_i^2} \lesssim \frac{1}{10^{17} \text{GeV}},$$

(16)

For the higher values of $M_{d_H}$ in Table 1 (see below), this constraint can be satisfied for not necessarily small values of $\alpha_i$, $\bar{\alpha}_i$.

### 2.5 Gauge coupling unification

Since we have chosen $V \sim M_U = M_{SU} = 10^{18} \text{GeV}$, this means that the Standard Model gauge couplings should unify at the scale $M_U$. However, their running will be modified due to the presence of the gap particles. Note that the underlying flipped $SU(5)$ symmetry, even though not evident in this respect, is nevertheless essential in the above discussion. The masses $M_Q$ and $M_{d_H}$ can then be determined, as follows

$$\ln \frac{M_{QH}}{m_Z} = \pi \left( \frac{1}{2\alpha_e} - \frac{1}{3\alpha_3} - \frac{\sin^2 \theta_w - 0.0029}{\alpha_e} \right) - 2 \ln \frac{M_U}{m_Z} - 0.63, \quad (17)$$

$$\ln \frac{M_{d_H}}{m_Z} = \pi \left( \frac{1}{2\alpha_e} - \frac{7}{3\alpha_3} + \frac{\sin^2 \theta_w - 0.0029}{\alpha_e} \right) - 6 \ln \frac{M_U}{m_Z} - 1.47, \quad (18)$$

where $\alpha_e$, $\alpha_3$ and $\sin^2 \theta_w$ are all measured at $M_Z$. This is a one-loop determination (the constants account for the dominant two-loop corrections) which neglects all low- and high-energy threshold effects, but is quite adequate for our present purposes. As shown in Table 1 (and Eq. (18)) the $d_H$ mass depends most sensitively on $\alpha_3(M_Z) = 0.118 \pm 0.008 \ [35]$, whereas the $Q_H$ mass and the unified coupling are rather insensitive to it. The unification of the gauge couplings is shown in Fig. 1 (solid lines) for the central value of $\alpha_3(M_Z)$. This figure also shows the case of no gap particles (dotted lines), for which $M_U \approx 10^{16} \text{GeV}$.

### 3 The Model: Supersymmetry Breaking Scenarios

Supersymmetry breaking in string models can generally be triggered in a phenomenologically acceptable way by non-zero $F$-terms for: (a) any of the moduli fields of the string model ($\langle F_M \rangle$) \[36\], (b) the dilaton field ($\langle F_D \rangle$) \[23\], or (c) the hidden matter fields ($\langle F_H \rangle$) \[37\]. It has been recently noted \[23\] that much model-independent information can be obtained about the structure of the soft supersymmetry breaking

---

\[7\]Here we assume a common supersymmetric threshold at $M_Z$. In fact, the supersymmetric threshold and the $d_H$ mass are anticorrelated. See Ref. \[10\] for a discussion.
Table 1: The value of the gap particle masses and the unified coupling for \( \alpha_3(M_Z) = 0.118 \pm 0.008 \). We have taken \( M_U = 10^{18} \) GeV, \( \sin^2 \theta_w = 0.233 \), and \( \alpha_e^{-1} = 127.9 \).

| \( \alpha_3(M_Z) \) | \( M_{\delta} \) (GeV) | \( M_{Q_H} \) (GeV) | \( \alpha(M_U) \) |
|-----------------|-----------------|-----------------|------------|
| 0.110           | \( 4.9 \times 10^4 \) GeV | \( 2.2 \times 10^{12} \) GeV | 0.0565     |
| 0.118           | \( 4.5 \times 10^6 \) GeV | \( 4.1 \times 10^{12} \) GeV | 0.0555     |
| 0.126           | \( 2.3 \times 10^8 \) GeV | \( 7.3 \times 10^{12} \) GeV | 0.0547     |

Figure 1: The running of the gauge couplings in the flipped SU(5) model for \( \alpha_3(M_Z) = 0.118 \) (solid lines). The gap particle masses have been derived using the gauge coupling RGEs to achieve unification at \( M_U = 10^{18} \) GeV. The case with no gap particles (dotted lines) is also shown; here \( M_U \approx 10^{16} \) GeV.
parameters in generic string supergravity models if one neglects the third possibility \(\langle F_M \rangle = 0\) and assumes that either: (i) \(\langle F_M \rangle \gg \langle F_D \rangle\), or (ii) \(\langle F_D \rangle \gg \langle F_M \rangle\).

In case (i) the scalar masses are generally not universal, i.e., \(m_i = f_im_0\) where \(m_0\) is the gravitino mass and \(f_i\) are calculable constants, and therefore large flavor-changing-neutral-currents (FCNCs) are potentially dangerous \([38]\). The gaugino masses arise from the one-loop contribution to the gauge kinetic function and are thus suppressed \((m_{1/2} \sim (\alpha/4\pi)m_0) \[33, 10, 23\]\). The experimental constraints on the gaugino masses then force the squark and slepton masses into the TeV range \([40]\). It is interesting to note that this supersymmetry breaking scenario is not unlike that required for the minimal \(SU(5)\) supergravity model in order to have the dimension-five proton decay operators under control \([33, 34]\), which requires \(m_{1/2}/m_0 \lesssim 1/3\). This constraint entails potential cosmological troubles: the neutralino relic density is large and one needs to tune the parameters to have the neutralino mass be very near the Higgs and \(Z\) resonances \([34, 41, 42]\). Clearly, such cosmological constraints are going to be exacerbated in the case (i) scenario \((m_{1/2}/m_0 \ll 1)\) and will likely require real fine-tuning of the model parameters.

An important exception to case (i) occurs if \(f_i \equiv 0\) and all scalar masses at the unification scale vanish \((\langle F_M \rangle_{m_0=0} = 0)\), as is the case in unified no-scale supergravity models \([24]\). This special case automatically restores the much needed universality of scalar masses, and in the context of no-scale models also entails \(A = 0\), see Eq. (1). A special case of this scenario occurs when the bilinear soft-supersymmetry breaking mass parameter \(B(M_U)\) is also required to vanish. With the additional ingredient of a flipped \(SU(5)\) gauge group, all the above problems are naturally avoided \([11]\), and interesting predictions for direct \([43, 44, 45, 46]\) and indirect \([47, 48, 49]\) experimental detection follow.

If supersymmetry breaking is triggered by \(\langle F_D \rangle\) (case (ii)), one obtains universal soft-supersymmetry gaugino and scalar masses and trilinear interactions \([23]\) and the soft-supersymmetry breaking parameters in Eq. (2) result. As well, there is a special more constrained case where \(B(M_U) = 2m_0 = \frac{2}{\sqrt{3}}m_{1/2}\) is also required, if one demands that the \(\mu\) parameter receive contributions solely from supergravity \([23]\). With the complement of a flipped \(SU(5)\) structure, this model has also been seen to avoid all the difficulties of the generic \(\langle F_M \rangle\) scenario \([50]\). This supersymmetry breaking scenario has been studied recently also in the context of the minimal supersymmetric Standard Model (MSSM) in Ref. \([51]\).

Therefore, in what follows we restrict ourselves to the two supersymmetry breaking scenarios in Eqs. (1,2) and their special cases \((B(M_U) = 0\) and \(B(M_U) = 2m_0\), respectively).

### 4 Phenomenology: General Case

The procedure to extract the low-energy predictions of the models outlined above is rather standard by now (see e.g., Ref. [52]): (a) the bottom-quark and tau-lepton masses, together with the input values of \(m_t\) and \(\tan \beta\) are used to determine the
respective Yukawa couplings at the electroweak scale; (b) the gauge and Yukawa couplings are then run up to the unification scale $M_U = 10^{18}$ GeV taking into account the extra vector-like quark doublet ($\sim 10^{12}$ GeV) and singlet ($\sim 10^{6}$ GeV) introduced above \[16, 11\]; (c) at the unification scale the soft-supersymmetry breaking parameters are introduced (according to Eqs. (13)) and the scalar masses are then run down to the electroweak scale; (d) radiative electroweak symmetry breaking is enforced by minimizing the one-loop effective potential which depends on the whole mass spectrum, and the values of the Higgs mixing term $|\mu|$ and the bilinear soft-supersymmetry breaking parameter $B$ are determined from the minimization conditions; (e) all known phenomenological constraints on the sparticle and Higgs masses are applied (most importantly the LEP lower bounds on the chargino and Higgs masses), including the cosmological requirement of not-too-large neutralino relic density.

4.1 Mass ranges

We have scanned the parameter space for $m_t = 130, 150, 170$ GeV, $\tan \beta = 2 \rightarrow 50$ and $m_{1/2} = 50 \rightarrow 500$ GeV. Imposing the constraint $m_{\tilde{g}, \tilde{q}} < 1$ TeV we find

$$\langle F_M \rangle_{m_0=0} : \quad m_{1/2} < 475 \text{ GeV}, \quad \tan \beta \lesssim 32, \quad \langle F_D \rangle : \quad m_{1/2} < 465 \text{ GeV}, \quad \tan \beta \lesssim 46.$$  

(19) \hspace{1cm} (20)

These restrictions on $m_{1/2}$ cut off the growth of most of the sparticle and Higgs masses at $\approx 1$ TeV. However, the sleptons, the lightest Higgs, the two lightest neutralinos, and the lightest chargino are cut off at a much lower mass, as follows:

$$\langle F_M \rangle_{m_0=0} : \quad \begin{cases} m_{\tilde{e}_R} < 190 \text{ GeV}, & m_{\tilde{e}_L} < 305 \text{ GeV}, \quad m_{\tilde{\nu}} < 295 \text{ GeV} \\ m_{\tilde{\tau}_1} < 185 \text{ GeV}, & m_{\tilde{\tau}_2} < 315 \text{ GeV} \\ m_h < 125 \text{ GeV} \\ m_{\chi_1^0} < 145 \text{ GeV}, & m_{\chi_2^0} < 290 \text{ GeV}, \quad m_{\chi_1^+} < 290 \text{ GeV} \\ m_{\chi_R} < 325 \text{ GeV}, & m_{\tilde{\nu}} < 400 \text{ GeV} \\ m_{\tilde{\tau}_1} < 325 \text{ GeV}, & m_{\tilde{\tau}_2} < 400 \text{ GeV} \\ m_h < 125 \text{ GeV} \\ m_{\chi_1^0} < 145 \text{ GeV}, & m_{\chi_2^0} < 285 \text{ GeV}, \quad m_{\chi_1^+} < 285 \text{ GeV} \end{cases}$$

(21)

$$\langle F_D \rangle : \quad \begin{cases} m_{\tilde{g}} \gtrsim 245 \,(260) \text{ GeV} \\ m_{\tilde{q}} \gtrsim 240 \,(250) \text{ GeV} \end{cases} \quad \langle F_D \rangle : \quad \begin{cases} m_{\tilde{g}} \gtrsim 195 \,(235) \text{ GeV} \\ m_{\tilde{q}} \gtrsim 195 \,(235) \text{ GeV} \end{cases}$$

(22) \hspace{1cm} (23)

It is interesting to note that due to the various constraints on the model, the gluino and (average) squark masses are bounded from below,

$$\langle F_M \rangle_{m_0=0} : \quad \begin{cases} m_{\tilde{g}} \gtrsim 245 \,(260) \text{ GeV} \\ m_{\tilde{q}} \gtrsim 240 \,(250) \text{ GeV} \end{cases} \quad \langle F_D \rangle : \quad \begin{cases} m_{\tilde{g}} \gtrsim 195 \,(235) \text{ GeV} \\ m_{\tilde{q}} \gtrsim 195 \,(235) \text{ GeV} \end{cases}$$

(23)

for $\mu > 0(\mu < 0)$. Relaxing the above conditions on $m_{1/2}$ simply allows all sparticle masses to grow further proportional to $m_{\tilde{g}}$.

\*In this class of supergravity models the three sneutrinos ($\tilde{\nu}$) are degenerate in mass. Also, $m_{\tilde{\mu}_L} = m_{\tilde{e}_L}$ and $m_{\tilde{\mu}_R} = m_{\tilde{e}_R}$. 

11
Table 2: The value of the $c_i$ coefficients appearing in Eq. (25), the ratio $c_3 = m_\tilde{g}/m_{1/2}$, and the average squark coefficient $\bar{c}_\tilde{q}$, for $\alpha_3(M_Z) = 0.118 \pm 0.008$. Also shown are the $a_i, b_i$ coefficients for the central value of $\alpha_3(M_Z)$ and both supersymmetry breaking scenarios ($M : \langle F_M \rangle_{m_0=0}, D : \langle F_D \rangle$). The results apply as well to the second-generation squark and slepton masses.

| $i$ | $c_i \ (0.110)$ | $c_i \ (0.118)$ | $c_i \ (0.126)$ | $i$ | $a_i(M)$ | $b_i(M)$ | $a_i(D)$ | $b_i(D)$ |
|-----|-----------------|-----------------|-----------------|-----|----------|----------|----------|----------|
| $\tilde{\nu}, \tilde{e}_L$ | 0.406 | 0.409 | 0.413 | $\tilde{e}_L$ | 0.302 | +1.115 | 0.406 | +0.616 |
| $\tilde{e}_R$ | 0.153 | 0.153 | 0.153 | $\tilde{\nu}$ | 0.302 | −2.089 | 0.406 | −1.153 |
| $\tilde{u}_L, \tilde{d}_L$ | 3.98 | 4.41 | 4.97 | $\tilde{u}_L$ | 0.991 | −0.118 | 1.027 | −0.110 |
| $\tilde{u}_R$ | 3.68 | 4.11 | 4.66 | $\tilde{d}_L$ | 0.991 | +0.164 | 1.027 | +0.152 |
| $\tilde{d}_R$ | 3.63 | 4.06 | 4.61 | $\tilde{d}_R$ | 0.950 | −0.033 | 0.989 | −0.030 |
| $c_{\tilde{q}}$ | 1.95 | 2.12 | 2.30 | $c_{\tilde{q}}$ | 3.82 | 4.07 | 4.80 | $\bar{c}_{\tilde{q}}$ | 1.95 | 2.12 | 2.30 | $\bar{c}_{\tilde{q}}$ | 3.82 | 4.07 | 4.80 |

4.2 Mass relations

The neutralino and chargino masses show a correlation observed before in this class of models \[^3, \[^{11}\] namely

$$m_{\chi_1^0} \approx \frac{1}{2} m_{\chi_2^0}, \quad m_{\chi_1^\pm} \approx M_2 = (\alpha_2/\alpha_3)m_\tilde{g} \approx 0.28m_\tilde{g}. \quad \text{(24)}$$

This is because throughout the parameter space $|\mu|$ is generally much larger than $M_W$ and $|\mu| > M_2$. In practice we find $m_{\chi_2^0} \approx m_{\chi_1^\pm}$ to be satisfied quite accurately, whereas $m_{\chi_1^0} \approx \frac{1}{2} m_{\chi_2^0}$ is only qualitatively satisfied, although the agreement is better in the $\langle F_D \rangle$ case. In fact, these two mass relations are much more reliable than the one that links them to $m_\tilde{g}$. The heavier neutralino ($\chi_{3,4}^0$) and chargino ($\chi_{2}^\pm$) masses are determined by the value of $|\mu|$; they all approach this limit for large enough $|\mu|$. More precisely, $m_{\chi_2^0}$ approaches $|\mu|$ sooner than $m_{\chi_1^0}$ does. On the other hand, $m_{\chi_1^0}$ approaches $m_{\chi_2^0}$ rather quickly.

The first- and second-generation squark and slepton masses can be determined analytically

$$\bar{m}_i = \left[ m_{1/2}^2(c_i + \xi_0^2) - d_i \tan^2 \beta - 1 + M_W^2 \right]^{1/2} = a_i m_{\tilde{g}} \left[ 1 + b_i \left( \frac{150}{m_{\tilde{g}}} \right)^2 \tan^2 \beta - 1 \right]^{1/2}, \quad \text{(25)}$$

where $d_i = (T_{3i} - Q) \tan^2 \theta_w + T_{3i} (\text{e.g., } d_{\tilde{u}_L} = \frac{1}{2} - \frac{1}{6} \tan^2 \theta_w, d_{\tilde{e}_R} = -\tan^2 \theta_w)$, and $\xi_0 = m_0/m_{1/2} = 0, \frac{1}{\sqrt{3}}$. The coefficients $c_i$ can be calculated numerically in terms of the low-energy gauge couplings, and are given in Table 2 for $\alpha_3(M_Z) = 0.118 \pm 0.008$.

\[^9\]These are renormalized at the scale $M_Z$. In a more accurate treatment, the $c_i$ would be renormalized at the physical sparticle mass scale, leading to second order shifts on the sparticle masses.
In the table we also give $c_{\tilde{g}} = m_{\tilde{g}}/m_{\tilde{t}/2}$. Note that these values are smaller than what is obtained in the minimal $SU(5)$ supergravity model (where $c_{\tilde{g}} = 2.90$ for $\alpha_3(M_Z) = 0.118$) and therefore the numerical relations between the gluino mass and the neutralino masses are different in that model. In the table we also show the resulting values for $\alpha_1, \alpha_2$ for the central value of $\alpha_3(M_Z)$. Note that the apparently larger $\tan \beta$ dependence in the $\langle F_M \rangle_{m_0=0}$ case (i.e., $|b_1(M)| > |b_1(D)|$) is actually compensated by a larger minimum value of $\bar{m}_{\tilde{g}}$ in this case (see Eq. (23)).

The “average” squark mass, $m_{\tilde{g}} = \frac{1}{2}(m_{\tilde{u}_L} + m_{\tilde{u}_R} + m_{\tilde{d}_L} + m_{\tilde{d}_R} + m_{\tilde{e}_L} + m_{\tilde{e}_R} + m_{\tilde{\tau}_L} + m_{\tilde{\tau}_R}) = (m_{\tilde{g}}/c_{\tilde{g}})\sqrt{\bar{c}_{\tilde{g}} + \xi^2}$, with $\bar{c}_{\tilde{g}}$ given in Table 2, is determined to be

$$m_{\tilde{g}} = \begin{cases} (1.00, 0.95, 0.95)m_{\tilde{g}}, & \langle F_M \rangle_{m_0=0} \\ (1.05, 0.99, 0.98)m_{\tilde{g}}, & \langle F_D \rangle \end{cases}$$

for $\alpha_3(M_Z) = 0.110, 0.118, 0.126$ (the dependence on $\tan \beta$ is small). The squark splitting around the average is $\approx 2\%$.

These masses are plotted in Fig. 3. The thickness and straightness of the lines shows the small $\tan \beta$ dependence, except for $\tilde{\nu}$. The results do not depend on the sign of $\mu$, except to the extent that some points in parameter space are not allowed for both signs of $\mu$: the $\mu < 0$ lines start-off at larger mass values. Note that

$$\langle F_M \rangle_{m_0=0} : \begin{cases} m_{\tilde{e}_R} \approx 0.18m_{\tilde{g}} \\ m_{\tilde{e}_L} \approx 0.30m_{\tilde{g}} \\ m_{\tilde{e}_R}/m_{\tilde{e}_L} \approx 0.61 \end{cases} \quad \langle F_D \rangle : \begin{cases} m_{\tilde{e}_R} \approx 0.33m_{\tilde{g}} \\ m_{\tilde{e}_L} \approx 0.41m_{\tilde{g}} \\ m_{\tilde{e}_R}/m_{\tilde{e}_L} \approx 0.81 \end{cases}$$

The third generation squark and slepton masses cannot be determined analytically. In Fig. 3 we show $\tilde{t}_{1,2}, \tilde{b}_{1,2}, \tilde{\tau}_{1,2}$ for the choice $m_\tau = 150$ GeV. The variability on the $\tilde{\tau}_{1,2}$ and $\tilde{b}_{1,2}$ masses is due to the $\tan \beta$-dependence in the off-diagonal element of the corresponding 2 $\times$ 2 mass matrices ($\propto m_{\tau,b}(A_{\tau,b} + \mu \tan \beta)$). The off-diagonal element in the stop-squark mass matrix ($\propto m_t(A_t + \mu / \tan \beta)$) is rather insensitive to $\tan \beta$ but still effects a large $\tilde{t}_1 - \tilde{t}_2$ mass splitting because of the significant $A_t$ contribution. Note that both these effects are more pronounced for the $\langle F_D \rangle$ case since there $|A_{t,b,\tau}|$ are larger than in the $\langle F_M \rangle_{m_0=0}$ case. The lowest values of the $\tilde{t}_1$ mass go up with $m_\tau$ and can be as low as

$$m_{\tilde{t}_1} \gtrsim \begin{cases} 160, 170, 190 (155, 150, 170) \text{ GeV}; & \langle F_M \rangle_{m_0=0} \\ 88, 112, 150 (92, 106, 150) \text{ GeV}; & \langle F_D \rangle \end{cases}$$

for $m_\tau = 130, 150, 170$ GeV and $\mu > 0$ ($\mu < 0$).

The one-loop corrected lightest CP-even ($h$) and CP-odd ($A$) Higgs boson masses are shown in Fig. 4 as functions of $m_{\tilde{g}}$ for $m_\tau = 150$ GeV. Following the methods of Ref. [14] we have determined that the LEP lower bound on $m_h$ becomes $m_h \gtrsim 60$ GeV, as the figure shows. The largest value of $m_h$ depends on $m_t$; we find

$$m_h < \begin{cases} 106, 115, 125 \text{ GeV}; & \langle F_M \rangle_{m_0=0} \\ 107, 117, 125 \text{ GeV}; & \langle F_D \rangle \end{cases}$$
Figure 2: The first-generation squark and slepton masses as a function of the gluino mass, for both signs of $\mu$, $m_t = 150\,\text{GeV}$, and both supersymmetry breaking scenarios under consideration. The same values apply to the second generation. The thickness of the lines and their deviation from linearity are because of the small $\tan\beta$ dependence.
Figure 3: The $\tilde{\tau}_{1,2}$, $\tilde{b}_{1,2}$, and $\tilde{t}_{1,2}$ masses versus the gluino mass for both signs of $\mu$, $m_t = 150$ GeV, and both supersymmetry breaking scenarios. The variability in the $\tilde{\tau}_{1,2}$, $\tilde{b}_{1,2}$, and $\tilde{t}_{1,2}$ masses is because of the off-diagonal elements of the corresponding mass matrices.
for $m_t = 130, 150, 170$ GeV. It is interesting to note that the one-loop corrected values of $m_h$ for $\tan \beta = 2$ are quite dependent on the sign of $\mu$. This phenomenon can be traced back to the $\tilde{t}_1 - \tilde{t}_2$ mass splitting which enhances the dominant $\tilde{t}$ one-loop corrections to $m_h$ [54], an effect which is usually neglected in phenomenological analyses. The $\tilde{t}_{1,2}$ masses for $\tan \beta = 2$ are drawn closer together than the rest. The opposite effect occurs for $\mu < 0$ and therefore the one-loop correction is larger in this case. The sign-of-$\mu$ dependence appears in the off-diagonal entries in the $\tilde{t}$ mass matrix $\propto m_t(A_t + \mu/\tan \beta)$, with $A_t < 0$ in this case. Clearly only small $\tan \beta$ matters, and $\mu < 0$ enhances the splitting. The $A$-mass grows fairly linearly with $m_{\tilde{g}}$ with a $\tan \beta$-dependent slope which decreases for increasing $\tan \beta$, as shown in Fig. 4. Note that even though $m_A$ can be fairly light, we always get $m_A > m_h$, in agreement with a general theorem to this effect in supergravity theories [55]. This result also implies that the channel $e^+e^- \rightarrow hA$ at LEPI is not kinematically allowed in this model.

4.3 Neutralino relic density

The computation of the neutralino relic density (following the methods of Refs. [56, 57]) shows that $\Omega_{\chi h^2} \lesssim 0.25 (0.90)$ in the no-scale (dilaton) model. This implies that in these models the cosmologically interesting values $\Omega_{\chi} h^2 \lesssim 1$ occur quite naturally. These results are in good agreement with the observational upper bound on $\Omega_{\chi} h^2$ [58]. Moreover, fits to the COBE data and the small and large scale structure of the Universe suggest [59] a mixture of $\approx 70\%$ cold dark matter and $\approx 30\%$ hot dark matter together with $h_0 \approx 0.5$. The hot dark matter component in the form of massive tau neutrinos has already been shown to be compatible with the flipped $SU(5)$ model we consider here [30, 31], whereas the cold dark matter component implies $\Omega_{\chi} h^2 \approx 0.17$ which is reachable in these models.

5 Phenomenology: Special Cases

5.1 The strict no-scale case

We now impose the additional constraint $B(M_U) = 0$ to be added to Eq. (1), and obtain the so-called strict no-scale case. Since $B(M_Z)$ is determined by the radiative electroweak symmetry breaking conditions, this added constraint needs to be imposed in a rather indirect way. That is, for given $m_{\tilde{g}}$ and $m_t$ values, we scan the possible values of $\tan \beta$ looking for cases where $B(M_U) = 0$. The most striking result is that solutions exist only for $m_t \lesssim 135$ GeV if $\mu > 0$ and for $m_t \gtrsim 140$ GeV if $\mu < 0$. That is, the value of $m_t$ determines the sign of $\mu$. Furthermore, for $\mu < 0$ the value of $\tan \beta$ is determined uniquely as a function of $m_t$ and $m_{\tilde{g}}$, whereas for $\mu > 0$, $\tan \beta$ can be double-valued for some $m_t$ range which includes $m_t = 130$ GeV (but does not include $m_t = 100$ GeV). In Fig. 3 (top row) we plot the solutions found in this manner for
Figure 4: The one-loop corrected $h$ and $A$ Higgs masses versus the gluino mass for both signs of $\mu$, $m_t = 150\text{ GeV}$, and the two supersymmetry breaking scenarios. Representative values of $\tan\beta$ are indicated.
the indicated $m_t$ values.

All the mass relationships deduced in the previous section apply here as well. The tan $\beta$-spread that some of them have will be much reduced though. The most noticeable changes occur for the quantities which depend most sensitively on tan $\beta$. In Fig. 5 (bottom row) we plot the one-loop corrected lightest Higgs boson mass versus $m_{\tilde{g}}$. The result is that $m_h$ is basically determined by $m_t$; only a weak dependence on $m_{\tilde{g}}$ exists. Moreover, for $m_t \lesssim 135 \text{ GeV} \Leftrightarrow \mu > 0$, $m_h \lesssim 105 \text{ GeV}$; whereas for $m_t \gtrsim 140 \text{ GeV} \Leftrightarrow \mu < 0$, $m_h \gtrsim 100 \text{ GeV}$. Therefore, in the strict no-scale case, once the top-quark mass is measured, we will know the sign of $\mu$ and whether $m_h$ is above or below 100 GeV.

For $\mu > 0$, we just showed that the strict no-scale constraint requires $m_t \lesssim 135 \text{ GeV}$. This implies that $\mu$ cannot grow as large as it did previously in the general case. In fact, for $\mu > 0$, $\mu_{\text{max}} \approx 745 \text{ GeV}$ before and $\mu_{\text{max}} \approx 440 \text{ GeV}$ now. This smaller value of $\mu_{\text{max}}$ has the effect of cutting off the growth of the $\chi^0_{3,4}, \chi^\pm_2$ masses at $\approx \mu_{\text{max}} \approx 440 \text{ GeV}$ (c.f. $\approx 750 \text{ GeV}$) and of the heavy Higgs masses at $\approx 530 \text{ GeV}$ (c.f. $\approx 940 \text{ GeV}$).

### 5.2 The special dilaton scenario case

In our analysis above, the radiative electroweak breaking conditions were used to determine the magnitude of the Higgs mixing term $\mu$ at the electroweak scale. This quantity is ensured to remain light as long as the supersymmetry breaking parameters remain light. In a fundamental theory this parameter should be calculable and its value used to determine the Z-boson mass. From this point of view it is not clear that the natural value of $\mu$ should be light. In specific models on can obtain such values by invoking non-renormalizable interactions \[20, 21\]. Another contribution to this quantity is generically present in string supergravity models \[22, 21, 23\]. The general case with contributions from both sources has been effectively dealt with in the previous section. If one assumes that only supergravity-induced contributions to $\mu$ exist, then it can be shown that the $B$-parameter at the unification scale is also determined \[23\],

$$B(M_U) = 2m_0 = \frac{2}{\sqrt{3}}m_{1/2},$$

which is to be added to the set of relations in Eq. \(\frac{2}{\sqrt{3}}\). This new constraint effectively determines tan $\beta$ for given $m_t$ and $m_{\tilde{g}}$ values and makes this restricted version of the model highly predictive.

From the outset we note that only solutions with $\mu < 0$ exist. This is not a completely obvious result, but it can be partially understood as follows. In tree-level approximation, $m_A^2 > 0 \Rightarrow \mu B < 0$ at the electroweak scale. Since $B(M_U)$ is required to be positive and not small, $B(M_Z)$ will likely be positive also, thus forcing $\mu$ to be negative. A sufficiently small value of $B(M_U)$ and/or one-loop corrections to

---

\[10\] The values of tan $\beta$ shown in Fig. 5 (top row) differ somewhat from those shown previously in Ref. \[11\] (Fig. 8) because in that paper $m_b = 4.5 \text{ GeV}$ was used, whereas throughout the present paper $m_b = 4.9 \text{ GeV}$ is used instead.
Figure 5: The value of $\tan \beta$ versus $m_{\tilde{g}}$ in the strict no-scale case (where $B(M_U) = 0$) for the indicated values of $m_t$. Note that the sign of $\mu$ is determined by $m_t$ and that $\tan \beta$ can be double-valued for $\mu > 0$. Also shown is the one-loop corrected lightest Higgs boson mass. Note that if $\mu > 0$ (for $m_t < 135$ GeV) then $m_h < 105$ GeV; whereas if $\mu < 0$ (for $m_t > 140$ GeV) then $m_h > 100$ GeV.

$m_A^2$ could alter this result, although in practice this does not happen. A numerical iterative procedure allows us to determine the value of $\tan \beta$ which satisfies Eq. (30), from the calculated value of $B(M_Z)$. We find that

$$\tan \beta \approx 1.57 - 1.63, 1.37 - 1.45, 1.38 - 1.40 \quad \text{for } m_t = 130, 150, 155 \text{ GeV} \quad (31)$$

is required. Since $\tan \beta$ is so small ($m_h^{\text{tree}} \approx 28 - 41$ GeV), a significant one-loop correction to $m_h$ is required to increase it above its experimental lower bound of $\approx 60$ GeV [44]. This requires the largest possible top-quark masses and a not-too-small squark mass. However, perturbative unification imposes an upper bound on $m_t$ for a given $\tan \beta$ [60], which in this case implies

$$m_t \lesssim 155 \text{ GeV}, \quad (32)$$

which limits the magnitude of $m_h$

$$m_h \lesssim 74, 87, 91 \text{ GeV} \quad \text{for } m_t = 130, 150, 155 \text{ GeV}. \quad (33)$$

Lower values of $m_t$ are disfavored experimentally.
Table 3: The range of allowed sparticle and Higgs masses in the restricted dilaton scenario. The top-quark mass is restricted to be \( m_t < 155 \) GeV. All masses in GeV.

| \( m_t \) | \( \tilde{g} \) | \( \chi_1^0 \) | \( \chi_2^0, \chi_1^\pm \) | tan \( \beta \) | \( h \) | \( \tilde{t} \) | \( \tilde{q} \) | \( A, H, H^+ \) |
|---|---|---|---|---|---|---|---|---|
| 130 | 335 – 1000 | 38 – 140 | 75 – 270 | 1.57 – 1.63 | 61 – 74 | 110 – 400 | 335 – 1000 | > 400 |
| 150 | 260 – 1000 | 24 – 140 | 50 – 270 | 1.37 – 1.45 | 64 – 87 | 90 – 400 | 260 – 1000 | > 400 |
| 155 | 640 – 1000 | 90 – 140 | 170 – 270 | 1.38 – 1.40 | 84 – 91 | 210 – 400 | 640 – 1000 | > 970 |

In Table 3 we give the range of sparticle and Higgs masses that are allowed in this case. Clearly, continuing top-quark searches at the Tevatron and Higgs searches at LEPI,II should probe this restricted scenario completely.

6 Prospects for Experimental Detection

The sparticle and Higgs spectrum shown in Figs. 2,3,4,5, and Table 3 can be explored partially at present and near future collider facilities, as we discuss below for each supersymmetry breaking scenario considered above. First, we want to point out that there are two indirect experimental constraints which restrict these models in a more general way [47, 48]: (i) the recently experimentally determined range for the \( b \rightarrow s\gamma \) rare decay mode [61] 
\[
BR(b \rightarrow s\gamma) = (0.6 - 5.5) \times 10^{-4}
\]  
(at 95% CL); and (ii) the precise LEP electroweak measurements which constrain the \( \epsilon_{1,2,3} \) parameters [62]. The first constraint is particularly effective in removing acceptable points in parameter space since in these models very small values of \( BR(b \rightarrow s\gamma) \) are not uncommon [17, 63]. The second constraint basically imposes an upper bound on the top-quark mass of \( \approx 175 \) GeV. However, for \( 150 \) GeV < \( m_t < 175 \) GeV a progressively stricter upper bound on the chargino mass (and therefore on all sparticle and Higgs masses) is required, \( i.e., 50 \) GeV < \( m_{\chi_1^\pm} < 100 \) GeV, in order to keep \( \epsilon_1 \) below its current 90% CL upper limit. This implies that the choice \( m_t = 170 \) GeV above is rather constrained [64]. Setting aside these indirect constraints on the parameter space of these models, we now discuss the prospects for direct experimental detection.
6.1 Tevatron

(a) The search and eventual discovery of the top quark will narrow down the three-dimensional parameter space of these models considerably. Moreover, in the two special cases discussed in the previous section this measurement will be very important: (i) in the strict no-scale case it will determine the sign of $\mu$ ($\mu > 0$ if $m_t < \sim 135$ GeV; $\mu < 0$ if $m_t > \sim 140$ GeV) and whether the Higgs mass is above or below $\approx 100$ GeV, and (ii) it may rule out the restricted dilaton scenario if $m_t > 150$ GeV.

(b) The trilepton signal in $p\bar{p} \rightarrow \chi_2^0 \chi_1^\pm X$, where $\chi_2^0$ and $\chi_1^\pm$ both decay leptonically, is a clean test of supersymmetry [64] and in particular of this class of models [43]. The trilepton rates in the no-scale model have been given in Ref. [43]. In Fig. 6 we show these for the case $m_t = 130$ GeV. One can show that with $L = 100$ pb$^{-1}$ of integrated luminosity, chargino masses as high as $\approx 175$ GeV could be explored, although some regions of parameter space for lighter chargino masses would remain unexplored. We expect that somewhat weaker results will hold for the dilaton model, since the sparticle masses are heavier in that model, especially the sleptons which enhance the leptonic branching ratios when they are light enough.

(c) The relation $m_{\tilde{q}} \approx m_{\tilde{g}}$ for the $\tilde{u}_{L,R}, \tilde{d}_{L,R}$ squark masses should allow to probe the low end of the squark and gluino allowed mass ranges, although the outlook is more promising for the dilaton model since the allowed range starts off at lower values of $m_{\tilde{g}, \tilde{q}}$ (see Eq. (23)). An important point distinguishing the two models is that the average squark mass is slightly below (above) the gluino mass in the no-scale (dilaton) model, which should have an important bearing on the experimental signatures and rates [65]. In the dilaton case the $\tilde{t}_1$ mass can be below 100 GeV for sufficiently low $m_t$, and thus may be detectable. As the lower bound on $m_t$ rises, this signal becomes less accessible. The actual reach of the Tevatron for the above processes depends on its ultimate integrated luminosity.

6.2 LEPI,II

(a) In the class of models we consider, the lightest Higgs boson has couplings to gauge bosons and fermions which are close to those of the Standard Model (SM) Higgs boson, and therefore experimental lower bounds to the SM Higgs mass have been shown to apply slightly weakened to the supersymmetric Higgs [44]. Since the lower bound on the SM Higgs boson mass could still be pushed up several GeV at LEPI, the strict dilaton scenario (which requires $m_h \approx 61 - 91$ GeV) could be further constrained at LEPI and definitely tested at LEPII. At LEPII the SM Higgs mass could be explored up to roughly the beam energy minus 100 GeV [66]. This will allow exploration of the low tan $\beta$ values in both models,
Figure 6: The number of trilepton events at the Tevatron per 100 pb$^{-1}$ in the no-scale model for $m_t = 130$ GeV. Note that with 200 pb$^{-1}$ and 60% detection efficiency it should be possible to probe chargino masses as high as 175 GeV.

although the strict no-scale case will probably be out of reach (see Figs. [4][5]). The $e^+e^- \rightarrow hA$ channel will be open for large $\tan\beta$ and low $m_{\tilde{g}}$.

(b) Chargino masses below the kinematical limit ($m_{\chi^\pm_1} \lesssim 100$ GeV) should not be a problem to detect through the “mixed” mode with one chargino decaying leptonically and the other one hadronically [45], i.e., $e^+e^- \rightarrow \chi^+_1\chi^-_1$, $\chi^+_1 \rightarrow \chi^0q\bar{q}'$, $\chi^-_1 \rightarrow \chi^0l\bar{\nu}_l$. In Fig. 7 (top row) we show the corresponding event rates in the no-scale model. Note that $m_{\chi^\pm_1}$ can be as high as $\approx 290$ GeV in these models.

(c) Selectron, smuon, and stau pair production is partially accessible for both the no-scale and dilaton models, although more so in the no-scale case. In Fig. 6 (bottom row) we show the rates for the most promising (dilepton) mode in $e^+e^- \rightarrow \tilde{e}_R\tilde{e}_R$ production in the no-scale model.

6.3 HERA

The elastic and deep-inelastic contributions to $e^-p \rightarrow \tilde{e}_R\chi^0_1$ and $e^-p \rightarrow \tilde{\nu}\chi^-_1$ in the no-scale model should push the LEPI lower bounds on the lightest selectron, the lightest neutralino, and the sneutrino masses by $\approx 25$ GeV with $\mathcal{L} = 100$ pb$^{-1}$ [16]. In Fig. 8 we show the elastic plus deep-inelastic contributions to the total supersymmetric signal ($ep \rightarrow \text{susy} \rightarrow eX + \not{p}$) versus the lightest selectron mass ($m_{\tilde{e}_R}$) and the sneutrino mass ($m_{\tilde{\nu}}$) in the no-scale model. These figures show the “reach” of HERA in each of these variables. With $\mathcal{L} = 1000$ pb$^{-1}$ HERA should be competitive with LEPII as far as the no-scale model is concerned. In the dilaton scenario, because of the
Figure 7: The number of “mixed” events (1-lepton+2jets+$\bar{p}$) events per $\mathcal{L} = 100\text{ pb}^{-1}$ at LEPII versus the chargino mass in the no-scale model (top row). Also shown (bottom row) are the number of di-electron events per $\mathcal{L} = 100\text{ pb}^{-1}$ from selectron pair production versus the lightest selectron mass.

somewhat heavier sparticle masses, the effectiveness of HERA is reduced, although probably both channels may be accessible.

7 Conclusions

We have presented the simplest, string-derivable, supergravity model which has as gauge group flipped $SU(5)$ with supplementary matter representations to ensure uni-

fication at the string scale ($\sim 10^{18}\text{ GeV}$). This basic structure is complemented by two possible string supersymmetry breaking scenario: $SU(N,1)$ no-scale supergravity and dilaton-induced supersymmetry breaking. These two variants should be considered to be idealizations of what their string-derived incarnation should be. The specification of the hidden sector is crucial to the determination of the supersymmetry breaking scenario at work. A thorough exploration of the parameter spaces of the two models yields interesting results for experimental detection at present or near future collid-
ers. In this regard, the no-scale model is more within reach than the dilaton model, because of its generally lighter spectrum. In both supersymmetry breaking scenario considered, there ia a more constrained special case which allows \( \tan \beta \) to be determined in terms of \( m_t \) and \( m_{\tilde{g}} \). In the strict no-scale case we find a striking result: if \( \mu > 0, \ m_t \lesssim 135 \text{ GeV} \), whereas if \( \mu < 0, \ m_t \gtrsim 140 \text{ GeV} \). Therefore the value of \( m_t \) determines the sign of \( \mu \). Furthermore, we found that the value of \( m_t \) also determines whether the lightest Higgs boson is above or below 100 GeV. In the restricted dilaton case there is an upper bound on the top-quark mass \( (m_t \lesssim 155 \text{ GeV}) \) and the lightest Higgs boson mass \( (m_h \lesssim 91 \text{ GeV}) \). Thus, continuing Tevatron top-quark searches and LEPI,II Higgs searches could probe this restricted scenario completely. In Table 4 we give a summary of the general properties of these models and a comparison of their spectra.

We conclude that these well motivated string-inspired/derived models (especially their strict versions) could soon be probed experimentally. The various ingredients making up these models are likely to be present in actual fully string-derived
Table 4: Major features of the $SU(5) \times U(1)$ string-inspired/derived model and a comparison of the two supersymmetry breaking scenarios considered. (All masses in GeV).

| SU(5) × U(1) | SU(5) × U(1) |
|--------------|--------------|
| • Easily string-derivable, several known examples | • Parameters 3: $m_{1/2}$, $\tan \beta$, $m_t$ |
| • Symmetry breaking to Standard Model due to vevs of 10, $\overline{10}$ and tied to onset of supersymmetry breaking | • Universal soft-supersymmetry breaking automatic |
| • Natural doublet-triplet splitting mechanism | • Dark matter: $\Omega h_0^2 < 0.25$ |
| • Proton decay: $d = 5$ operators very small | • $m_{1/2} < 475$ GeV, $\tan \beta < 32$ |
| | • $m_\tilde{g} > 195$ GeV, $m_\tilde{q} > 195$ GeV |
| | • $m_\tilde{q} \approx 1.01 m_\tilde{g}$ |
| | • $m_{\tilde{t}_1} > 90$ GeV |
| | • $m_{\tilde{e}_R} \approx 0.30 m_\tilde{g}$ |
| | • $m_{\tilde{e}_L}/m_{\tilde{e}_R} \approx 0.61$ |
| | • $60$ GeV $< m_h < 125$ GeV |
| | • $2 m_{\chi_1^0} \approx m_{\chi_2^0} \approx m_{\chi_1^+} \approx 0.28 m_\tilde{g} \lesssim 290$ |
| | • $m_{\chi_3^0} \approx m_{\chi_4^0} \approx m_{\chi_3^\pm} \sim |\mu|$ |
| | • Spectrum easily accessible soon |
| | • Spectrum accessible soon |
| | | |
| | | | | |
| | | | | |
| | | | | |
| | | | | |
| | | | | |
models which yield the set of supersymmetry breaking parameters in Eqs. (1,2). The search for such models is imperative, although it may not be an easy task since in traditional gaugino condensation scenario Eqs. (1,2) are usually not reproduced (see however Refs. [67, 68]). Moreover, the requirement of vanishing vacuum energy may be difficult to fulfill, as a model with these properties and all the other ones outlined in Sec. 1 is yet to be found. This should not be taken as a discouragement since the harder it is to find the correct model, the more likely it is to be in some sense unique.

Acknowledgements

This work has been supported in part by DOE grant DE-FG05-91-ER-40633. The work of J.L. has been supported by an SSC Fellowship.
References

[1] See *e.g.*, *String theory in four dimensions*, ed. by M. Dine (North-Holland, Amsterdam, 1988); *Superstring construction*, ed. by A. N. Schellekens (North-Holland, Amsterdam, 1989).

[2] See *e.g.*, J. Ellis, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B 245 (1990) 375; A. Font, L. Ibáñez, and F. Quevedo, Nucl. Phys. B 345 (1990) 389.

[3] D. Lewellen, Nucl. Phys. B 337 (1990) 61; J. A. Schwarz, Phys. Rev. D 42 (1990) 1777.

[4] S. Barr, Phys. Lett. B 112 (1982) 219, Phys. Rev. D 40 (1989) 2457; J. Derendinger, J. Kim, and D. V. Nanopoulos, Phys. Lett. B 139 (1984) 170.

[5] I. Antoniadis, J. Ellis, J. Hagelin, and D. V. Nanopoulos, Phys. Lett. B 194 (1987) 231.

[6] I. Antoniadis, J. Ellis, J. Hagelin, and D. V. Nanopoulos, Phys. Lett. B 231 (1989) 65.

[7] J. Ellis, J. Hagelin, S. Kelley, and D. V. Nanopoulos, Nucl. Phys. B 311 (1988/89) 1.

[8] L. Ibáñez, H. Nilles, and F. Quevedo, Phys. Lett. B 187 (1987) 25; L. Ibáñez, J. Mas, H. Nilles, and F. Quevedo, Nucl. Phys. B 301 (1988) 157; A. Font, L. Ibáñez, F. Quevedo, and A. Sierra, Nucl. Phys. B 331 (1990) 421.

[9] A. Faraggi, D. V. Nanopoulos, and K. Yuan, Nucl. Phys. B 335 (1990) 347; A. Faraggi, Phys. Lett. B 278 (1992) 131, Phys. Lett. B 274 (1992) 47, Nucl. Phys. B 387 (1992) 239.

[10] I. Antoniadis, G. Leontaris, and J. Rizos, Phys. Lett. B 245 (1990) 161.

[11] J. L. Lopez, D. V. Nanopoulos, and A. Zichichi, Texas A & M University preprint CTP-TAMU-68/92, CERN-TH.6667/92, and CERN-PPE/92-188.

[12] I. Antoniadis, J. Ellis, R. Lacaze, and D. V. Nanopoulos, Phys. Lett. B 268 (1991) 188; S. Kalara, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B 269 (1991) 84.

[13] J. L. Lopez, D. V. Nanopoulos, and K. Yuan, Nucl. Phys. B 399 (1993) 654.

[14] J. Ellis, S. Kelley, and D. V. Nanopoulos, Phys. Lett. B 249 (1990) 441; F. Anselmo, L. Cifarelli, and A. Zichichi, Nuovo Cimento 105A (1992) 1335.

[15] I. Antoniadis, J. Ellis, S. Kelley, and D. V. Nanopoulos, Phys. Lett. B 272 (1991) 31; D. Bailin and A. Love, Phys. Lett. B 280 (1992) 26.
[16] S. Kelley, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B 278 (1992) 140; G. Leontaris, Phys. Lett. B 281 (1992) 54.

[17] For a review see e.g., J. L. Lopez and D. V. Nanopoulos, in Proceedings of the 15th Johns Hopkins Workshop on Current Problems in Particle Theory, August 1991, p. 277; ed. by G. Domokos and S. Kovesi-Domokos (World Scientific, Singapore 1992).

[18] J. Ellis, S. Kelley, and D. V. Nanopoulos, Phys. Lett. B 260 (1991) 131.

[19] J. L. Lopez and D. V. Nanopoulos, Phys. Lett. B 251 (1990) 73, Phys. Lett. B 256 (1991) 150, and Phys. Lett. B 268 (1991) 359.

[20] J. E. Kim and H. P. Nilles, Phys. Lett. B 138 (1984) 150 and Phys. Lett. B 263 (1991) 79; E. J. Chun, J. E. Kim, and H. P Nilles, Nucl. Phys. B 370 (1992) 105.

[21] J. Casas and C. Muñoz, Phys. Lett. B 306 (1993) 288.

[22] G. Giudice and A. Masiero, Phys. Lett. B 206 (1988) 480.

[23] V. Kaplunovsky and J. Louis, Phys. Lett. B 306 (1993) 269.

[24] For a review see A. B. Lahanas and D. V. Nanopoulos, Phys. Rep. 145 (1987) 1.

[25] E. Witten, Phys. Lett. B 155 (1985) 151.

[26] J. Ellis, C. Kounnas, and D. V. Nanopoulos, Nucl. Phys. B 241 (1984) 406.

[27] J. Ellis, C. Kounnas, and D. V. Nanopoulos, Nucl. Phys. B 247 (1984) 373.

[28] S. Kalara, J. Lopez, and D. V. Nanopoulos, Phys. Lett. B 245 (1990) 421, Nucl. Phys. B 353 (1991) 650.

[29] S. Kelley, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B 261 (1991) 424.

[30] J. Ellis, J. L. Lopez, and D. V. Nanopoulos, Phys. Lett. B 292 (1992) 189.

[31] J. Ellis, J. L. Lopez, D. V. Nanopoulos, and K. Olive, Phys. Lett. B 308 (1993) 70.

[32] J. Ellis, D. V. Nanopoulos, and K. Olive, Phys. Lett. B 300 (1993) 121.

[33] M. Matsumoto, J. Arafune, H. Tanaka, and K. Shiraishi, Phys. Rev. D 46 (1992) 3966; R. Arnowitt and P. Nath, Phys. Rev. Lett. 69 (1992) 725; P. Nath and R. Arnowitt, Phys. Lett. B 287 (1992) 89; J. Hisano, H. Murayama, and T. Yanagida, Phys. Rev. Lett. 69 (1992) 1014 and Tohoku University preprint TU–400 (July 1992).
[34] J. L. Lopez, D. V. Nanopoulos, and A. Zichichi, Phys. Lett. B 291 (1992) 255; J. L. Lopez, D. V. Nanopoulos, and H. Pois, Phys. Rev. D 47 (1993) 2468; J. L. Lopez, D. V. Nanopoulos, H. Pois, and A. Zichichi, Phys. Lett. B 299 (1993) 262.

[35] S. Bethke, in Proceedings of the XXVI International Conference on High Energy Physics, Dallas, August 1992, ed. by J. R. Sanford (AIP Conference Proceedings No. 272), p. 81.

[36] J. Derendinger, L. Ibáñez, and H. Nilles, Phys. Lett. B 155 (1985) 65; M. Dine, N. Seiberg, and E. Witten, Phys. Lett. B 156 (1985) 55; A. Font, L. Ibáñez, D. Lüst, and F. Quevedo, Phys. Lett. B 245 (1990) 401; H. Nilles and M. Olechowski, Phys. Lett. B 248 (1990) 268; P. Binetruy and M. Gaillard, Nucl. Phys. B 358 (1991) 121.

[37] I. Antoniadis, J. Ellis, A. Lahanas, and D. V. Nanopoulos, Phys. Lett. B 241 (1990) 24.

[38] J. Ellis and D. V. Nanopoulos, Phys. Lett. B 110 (1982) 44.

[39] L. Ibáñez and D. Lüst, Nucl. Phys. B 382 (1992) 305.

[40] B. de Carlos, J. Casas, and C. Muñoz, Nucl. Phys. B 399 (1993) 623 and Phys. Lett. B 299 (1993) 234.

[41] R. Arnowitt and P. Nath, Phys. Lett. B 299 (1993) 58 and 307 (1993) 403(E); P. Nath and R. Arnowitt, Phys. Rev. Lett. 70 (1993) 3696.

[42] J. L. Lopez, D. V. Nanopoulos, and K. Yuan, Texas A & M University preprint CTP-TAMU-14/93 (to appear in Phys. Rev. D).

[43] J. L. Lopez, D. V. Nanopoulos, X. Wang, and A. Zichichi, Texas A & M University preprint CTP-TAMU-76/92 and CERN-PPE/92-194 (to appear in Phys. Rev. D).

[44] J. L. Lopez, D. V. Nanopoulos, H. Pois, X. Wang, and A. Zichichi, Phys. Lett. B 306 (1993) 73.

[45] J. L. Lopez, D. V. Nanopoulos, H. Pois, X. Wang, and A. Zichichi, Texas A & M University preprint CTP-TAMU-89/92 and CERN-PPE/93-16.

[46] J. L. Lopez, D. V. Nanopoulos, X. Wang, and A. Zichichi, Texas A & M University preprint CTP-TAMU-15/93 and CERN-PPE/93-64.

[47] J. L. Lopez, D. V. Nanopoulos, and G. Park, Texas A & M University preprint CTP-TAMU-16/93 (to appear in Phys. Rev. D).
[48] J. L. Lopez, D. V. Nanopoulos, G. Park, H. Pois, and K. Yuan, Texas A & M University preprint CTP-TAMU-19/93.

[49] R. Gandhi, J. L. Lopez, D. V. Nanopoulos, K. Yuan, and A. Zichichi (in preparation).

[50] J. L. Lopez, D. V. Nanopoulos, and A. Zichichi, Texas A & M University preprint CTP-TAMU-31/93 and CERN-TH.6903/93.

[51] R. Barbieri, J. Louis, and M. Moretti, CERN-TH.6856/93.

[52] S. Kelley, J. L. Lopez, D. V. Nanopoulos, H. Pois, and K. Yuan, Nucl. Phys. B 398 (1993) 3.

[53] P. Nath and R. Arnowitt, Phys. Lett. B 289 (1992) 368.

[54] J. Ellis, G. Ridolfi, and F. Zwirner, Phys. Lett. B 262 (1991) 477.

[55] M. Drees and M. Nojiri, Phys. Rev. D 45 (1992) 2482.

[56] J. L. Lopez, D. V. Nanopoulos, K. Yuan, Nucl. Phys. B 370 (1992) 445.

[57] S. Kelley, J. L. Lopez, D. V. Nanopoulos, H. Pois, and K. Yuan, Phys. Rev. D 47 (1993) 2461.

[58] See e.g., E. Kolb and M. Turner, The Early Universe (Addison-Wesley, 1990).

[59] See e.g., R. Schaefer and Q. Shafi, Nature 359 (1992) 199; A. N. Taylor and M. Rowan-Robinson, Nature 359 (1992) 393.

[60] L. Durand and J. L. Lopez, Phys. Lett. B 217 (1989) 463, Phys. Rev. D 40 (1989) 207.

[61] E. Thorndike, talk given at the 1993 Meeting of the American Physical Society, Washington D. C., April 1993.

[62] M. Peskin and T. Takeuchi, Phys. Rev. Lett. 65 (1990) 964; W. Marciano and J. Rosner, Phys. Rev. Lett. 65 (1990) 2963; D. Kennedy and P. Langacker, Phys. Rev. Lett. 65 (1990) 2967; G. Altarelli and R. Barbieri, Phys. Lett. B 253 (1990) 161; G. Altarelli, R. Barbieri, and S. Jadach, Nucl. Phys. B 369 (1992) 3; R. Barbieri, M. Frigeni, and F. Caravaglios, Phys. Lett. B 279 (1992) 169; G. Altarelli, R. Barbieri, and F. Caravaglios, CERN-TH.6770/93.

[63] J. L. Lopez, D. V. Nanopoulos, G. Park, and A. Zichichi (in preparation).

[64] P. Nath and R. Arnowitt, Mod. Phys. Lett. A 2 (1987) 331; R. Barbieri, F. Caravaglios, M. Frigeni, and M. Mangano, Nucl. Phys. B 367 (1991) 28; H. Baer and X. Tata, Phys. Rev. D 47 (1993) 2739.
[65] See e.g., J. White, talk presented at the “Recent Advances in the Superworld” Workshop, The Woodlands, Texas, April, 1993; R. M. Barnett, J. Gunion, and H. Haber, LBL-34106; H. Baer, C. Kao, and X. Tata, FSU-HEP-930527.

[66] See e.g., J. Alcaraz, M. Fetcini, M. Pieri, and B. Zhou, CERN-PPE/93-28.

[67] See e.g., J. Casas, Z. Lalak, C. Muñoz, and G. Ross, Nucl. Phys. B 347 (1990) 243; A. de la Macorra and G. Ross, OUTP-92-14P.

[68] I. Antoniadis, J. Ellis, E. Floratos, D. V. Nanopoulos, and T. Tomaras, Phys. Lett. B 191 (1987) 96; M. Porrati and F. Zwirner, Nucl. Phys. B 326 (1989) 162 and references therein.