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Constraints on Hidden Sector Gaugino Condensation

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When considering the subject of effective field theories from strings, the notion of "phenomenological viability" has in the past been a very loose standard. Indeed some of the well-known problems facing such low-energy theories seemed quite intractable, depressing the prospects of ever being able to refer to a meaningful superstring phenomenology. The problems to which we refer include the need to generate a hierarchy between the supersymmetry-breaking scale and the Planck scale, the cosmological dangers of moduli fields with Planck-suppressed interactions, the desire for a weakly-coupled effective quantum field theory, and most significantly the need to stabilize the dilaton [1].

Recently [2, 3, 4], however, it was shown that by incorporating postulated nonperturbative string-theoretical effects in a modular invariant low-energy field theory the above problems can be addressed in a simple manner with tuning required only in the vanishing of the cosmological constant. Having passed these initial tests it now becomes possible to ask for a slightly higher standard in "viability."

The philosophy behind this study is to probe this class of models in a series of phenomenological arenas to uncover relations between the dynamics of the hidden sector and the nature of our observable world. After a review in Section 1 of the class of models previously developed in [2, 3, 4, 5], we investigate in Section 2 the initial challenge of setting the supersymmetry-breaking scale that all effective field theories from strings must confront. This is largely a reiteration of results discussed in [4]. In Section 3 we turn to the pattern of soft supersymmetry-breaking parameters and look for the implications of current mass bounds arising from searches at LEP and the Tevatron. Finally, Section 4 considers the question of gauge coupling unification in the context of string theory.

1 Model

1.1 The Effective Lagrangian

The following is a condensation of material more fully presented in [3, 5] and aims to bring together the key points necessary for the subsequent discussion of phenomenological consequences. In those references, as here, the Kähler $U(1)$ superspace formalism of [6] is used throughout.

Supersymmetry breaking is implemented via condensation of gauginos charged under the hidden sector gauge group $\mathcal{G} = \prod_a \mathcal{G}_a$, which is taken to be a subgroup of E_8. For each gaugino condensate a vector superfield V_a is introduced and the gaugino condensate superfields $U_a \simeq \mathrm{Tr}(\mathcal{W}^\alpha \mathcal{W}_\alpha)_a$ are then identified as the (anti-)chiral projections of the vector superfields:

$$U_a = -\left(\mathcal{D}_{\dot\alpha}\mathcal{D}^{\dot\alpha} - 8R\right) V_a, \quad \overline{U}_a = -\left(\mathcal{D}^\alpha \mathcal{D}_\alpha - 8\overline{R}\right) V_a. \tag{1}$$

The dilaton field (in the linear multiplet formalism [7] used here) is the lowest component of the vector superfield $V = \sum_a V_a$: $\ell = V|_{\theta=\bar\theta=0}$. Note that none of the individual lowest components $V_a|_{\theta=\bar\theta=0}$ will appear in the effective theory component Lagrangian.

In the class of orbifold compactifications we will be considering there are three untwisted moduli chiral superfields T^I and matter chiral superfields Φ^A with Kähler potential

$$K = k(V) + \sum_I g^I + \sum_A e^{\sum_I q_I^A g^I} \left|\Phi^A\right|^2 + \mathcal{O}\left(\Phi^4\right), \quad g^I = -\ln\left(T^I + \overline{T}^I\right), \tag{2}$$

where the q_I^A are the modular weights of the fields Φ^A. The relevant part of the complete effective Lagrangian is then

$$\mathcal{L}_{\mathrm{eff}} = \mathcal{L}_{\mathrm{KE}} + \mathcal{L}_{\mathrm{VY}} + \mathcal{L}_{\mathrm{pot}} + \sum_a \mathcal{L}_{\mathrm{a}} + \mathcal{L}_{\mathrm{GS}}, \tag{3}$$

where

$$\mathcal{L}_{\mathrm{KE}} = \int \mathrm{d}^4\theta E\left[-2 + f(V)\right], \quad k(V) = \ln V + g(V), \tag{4}$$

is the Lagrangian density for the gravitational sector coupled to the vector multiplet and gives the kinetic energy terms for the dilaton, chiral multiplets, gravity superfields and tree-level Yang-Mills terms. Here the functions $f(V)$ and $g(V)$ represent nonperturbative corrections to the Kähler potential arising from string effects. The two functions f and g are related by the requirement that the Einstein term in (4) have canonical normalization:

$$V\frac{\mathrm{d}g(V)}{\mathrm{d}V} = -V\frac{\mathrm{d}f(V)}{\mathrm{d}V} + f(V), \tag{5}$$

and obey the weak-coupling boundary conditions: $f(0) = g(0) = 0$. In the presence of these nonperturbative effects the relationship between the dilaton and the effective field theory gauge coupling becomes $g^2/2 = \ell/(1 + f(\ell))$.

The second term in (3) is a generalization of the original Veneziano-Yankielowicz superpotential term [8],

$$\mathcal{L}_{\mathrm{VY}} = \frac{1}{8}\sum_a \int \mathrm{d}^4\theta \frac{E}{R} U_a \left[b'_a \ln\left(e^{-K/2}U_a\right) + \sum_\alpha b_a^\alpha \ln\left[(\Pi^\alpha)^{p_\alpha}\right]\right] + \mathrm{h.c.}, \tag{6}$$

which involves the gauge condensates U_a as well as possible gauge-invariant matter condensates described by chiral superfields $\Pi^\alpha \simeq \prod_A \left(\Phi^A\right)^{n_\alpha^A}$ [9]. Neither the gaugino nor the matter condensate superfields are taken to be propagating [10]. The coeffecients b'_a, b_a^α and p_α are determined by demanding the correct transformation properties of the expression in (6) under chiral and conformal transformations [3, 11] and yield the following relations:

$$b'_a = \frac{1}{8\pi^2}\left(C_a - \sum_A C_a^A\right), \quad \sum_{\alpha,A} b_a^\alpha n_\alpha^A p_\alpha = \sum_A \frac{C_a^A}{4\pi^2}, \quad p_\alpha \sum_A n_\alpha^A = 3 \quad \forall\alpha. \tag{7}$$

The final condition amounts to choosing the value of p_α so that the effective operator $(\Pi^\alpha)^{p_\alpha}$ has mass dimension three. In (7) the quantities C_a and C_a^A are the quadratic Casimir operators for the adjoint and matter representations, respectively. Given the above relations it is also convenient to define the combination

$$b_a \equiv b'_a + \sum_\alpha b_a^\alpha = \frac{1}{8\pi^2}\left(C_a - \frac{1}{3}\sum_A C_a^A\right) \tag{8}$$

which is proportional to the one-loop beta-function coefficient for the condensing gauge group $\mathcal{G}_a$.

The third term in (3) is a superpotential term for the matter condensates consistent with the symmetries of the underlying theory

$$\mathcal{L}_{\text{pot}} = \frac{1}{2}\int \mathrm{d}^4\theta \frac{E}{R} e^{K/2} W\left[(\Pi^\alpha)^{p_\alpha}, T^I\right] + \text{h.c.}. \tag{9}$$

We will adopt the same set of simplifying assumptions taken up in [3], namely that for fixed α, $b_a^\alpha \neq 0$ for only one value of a. Then $u_a = 0$ unless $W_\alpha \neq 0$ for every value of α for which $b_a^\alpha \neq 0$. We next assume that there are no unconfined matter fields charged under the hidden sector gauge group and ignore possible dimension-two matter condensates involving vector-like pairs of matter fields. This allows a simple factorization of the superpotential of the form

$$W\left[(\Pi^{p_\alpha}), T\right] = \sum_\alpha W_\alpha(T)\,(\Pi^\alpha)^{p_\alpha}, \tag{10}$$

where the functions W_α are given by

$$W_\alpha(T) = c_\alpha \prod_I \left[\eta\left(T^I\right)\right]^{2\left(p_\alpha q_I^\alpha - 1\right)}. \tag{11}$$

Here $q_I^\alpha = \sum_A n_\alpha^A q_I^A$ and the Yukawa coefficients c_α, while *a priori* unknown variables, are taken to be of $\mathcal{O}(1)$. The function $\eta(T^I)$ is the Dedekind function and its presence in (11) ensures the modular invariance of this term in the Lagrangian.

The remaining terms in (3) include the quantum corrections from light field loops to the unconfined Yang-Mills couplings and the Green-Schwarz (GS) counterterm introduced to ensure modular invariance.[1] The latter is given by the expression

$$\mathcal{L}_{\text{GS}} = \int \mathrm{d}^4\theta E V V_{\text{GS}}, \tag{12}$$

$$V_{\text{GS}} = b\sum_I g^I + \sum_A p_A e^{\sum_I q_I^A g^I}\left|\Phi^A\right|^2 + \mathcal{O}\left(\left|\Phi^A\right|^4\right), \tag{13}$$

[1]Not included in this paper are string loop corrections ($\mathcal{L}_{\text{th}}$) [13] which vanish for orbifold compactifications with no $N = 2$ supersymmetry sector [14].

where $b \equiv C_{E_8}/8\pi^2 \approx 0.38$ is proportional to the beta-function coefficient for the group E_8 and the coefficients p_A are as yet undetermined.

As for the operators $\mathcal{L}_a$ in (3), their rather involved form in curved superspace was worked out in [5] and will not be repeated here. Their importance for this work lies in their contributions to the supersymmetry-breaking gaugino masses at the condensation scale arising from the superconformal anomaly – a contribution that was recently emphasized by a number of authors [12]. We will return to these in Section 3.1.

1.2 Condensation and Dilaton Stabilization

The Lagrangian in (3) can be expanded into component form using the standard techniques of the Kähler superspace formalism of supergravity [6]. In reference [3] the bosonic part of the Lagrangian relevant to dilaton stabilization and gaugino condensation was presented and the equations of motion for the nonpropagating fields were solved. In particular, the equations of motion for the auxiliary fields of the condensates U^a give

$$\rho_a{}^2 = e^{-2\frac{b'_a}{b_a}} e^K e^{-\frac{(1+f)}{b_a \ell}} e^{-\frac{b}{b_a}\sum_I g^I} \prod_I \left|\eta\left(t^I\right)\right|^{\frac{4(b-b_a)}{b_a}} \prod_\alpha \left|b_a^\alpha/4c_\alpha\right|^{-2\frac{b_a^\alpha}{b_a}}, \qquad (14)$$

where $t_I \equiv T_I|_{\theta=\bar\theta=0}$ and $u_a = U_a|_{\theta=\bar\theta=0} \equiv \rho_a e^{i\omega_a}$.

Upon substituting for the gauge coupling via the relation $g^2/2 = \ell/(1+f(\ell))$ we recognize the expected one-instanton form for gaugino condensation. Expression (14) encodes more information, however, than simply the one-loop running of the gauge coupling. In [11] the loop corrections to the gauge coupling constants were computed using a manifestly supersymmetric Pauli-Villars regularization. The (moduli independent) corrections were identified with the renormalization group invariant [15]

$$\delta_a = \frac{1}{g_a^2(\mu)} - \frac{3b_a}{2}\ln\mu^2 + \frac{2C_a}{16\pi^2}\ln g_a^2(\mu) + \frac{2}{16\pi^2}\sum_a C_a^A \ln Z_a^A(\mu). \qquad (15)$$

Using the above expression it is possible to solve for the scale at which the $1/g^2(\mu)$ term becomes negligible relative to the $\ln g^2(\mu)$ term – effectively

looking for the "all loop" Landau pole for the coupling constant. This scale is related to the string scale by the relation

$$\mu_L{}^2 \sim \mu_{\rm str}{}^2 e^{-\frac{2}{3b_a g_a^2(\mu)}} \prod_A \left[Z_a^A\left(\mu_{\rm str}\right) / Z_a^A\left(\mu_L\right)\right]^{\frac{C_a^A}{12\pi^2 b_a}} . \tag{16}$$

Now comparing the effective Lagrangian given in Section 1.1 with the field theory loop calculation given in [11] shows that the two agree provided we identify the wave function renormalization coefficients Z_a^A with the quantity $|4W_\alpha/b_a^\alpha|^2$. This is precisely what is needed to produce the final product in the condensate expression given in (14), indicating that the condensation scale represents the scale at which the coupling becomes strong as would be computed using the so-called "exact" beta-function.

Note that this final factor introduces the unknown Yukawa coefficients c_α into the scale of supersymmetry breaking. Such dependence of the gaugino condensate on the parameters of the superpotential is not unexpected, and has in fact been demonstrated in the case of supersymmetric QCD as well as certain models of supersymmetric Yang-Mills theories coupled to chiral matter [16]. This last Yukawa-related factor has the virtue of allowing two different hidden sector configurations which result in the same beta-function to condense at widely different scales.

In order to go further and make quantitative statements about the scale of gaugino condensation (and hence supersymmetry breaking) it is necessary to specify some form for the nonperturbative effects represented by the functions f and g. The parameterization adopted in [4] was originally motivated by Shenker [17] and was of the form $\exp\left(-1/g_{\rm str}\right)$ where $g_{\rm str}$ is the string coupling constant. A consensus seems to be forming [18] around this characterization for string nonperturbative effects and the function $f\left(V\right)$ in (4) will be taken to be of the form

$$f\left(V\right) = \left[A_0 + A_1/\sqrt{V}\right] e^{-B/\sqrt{V}}, \tag{17}$$

which was shown [4] to allow dilaton stabilization at weak to moderate string coupling with parameters that are all of $\mathcal{O}\left(1\right)$. The benefits of invoking

string-inspired nonperturbative effects of the form of (17) have recently been explored by others in the literature [19].

The scalar potential for the moduli t_I is minimzed at the self-dual points $\langle t_I \rangle = 1$ or $\langle t_I \rangle = \exp(i\pi/6)$, where the corresponding F-components F_I of the chiral superfields T^I vanish. At these points the dilaton potential is given by

$$V(\ell) = \frac{1}{16\ell^2}\left(1 + \ell\frac{\mathrm{d}g}{\mathrm{d}\ell}\right)\left|\sum_a (1 + b_a\ell)\, u_a\right|^2 - \frac{3}{16}\left|\sum_a b_a u_a\right|^2 . \qquad (18)$$

As an example, the potential (18) can be minimized with vanishing cosmological constant and $\alpha_{\rm str} = 0.04$ for $A_0 = 3.25, A_1 = -1.70$ and $B = 0.4$ in expression (17).

2 Phenomenological Implications

2.1 Scale of Supersymmetry Breaking

With the adoption of (17) the scale of gaugino condensation can be obtained once the following are specified: (**1**) the condensing subgroup(s) of the original hidden sector gauge group E_8, (**2**) the representations of the matter fields charged under the condensing subgroup(s), (**3**) the Yukawa coefficients in the superpotential for the hidden sector matter fields and (**4**) the value of the string coupling constant at the compactification scale, which in turn determines the coefficients in (17) necessary to minimize the scalar potential (18).

A great deal of simplification in the above parameter space can be obtained by making the ansatz that all of the matter in the hidden sector which transforms under a given subgroup $\mathcal{G}_a$ is of the same representation, such as the fundamental representation. Then the sum of the coefficients b_a^α over the number of condensate fields labeled by α ($\alpha = 1, \ldots, N_c$), can be replaced by

one effective variable

$$\sum_{\alpha} b_a^{\alpha} \longrightarrow (b_a^{\alpha})_{\text{eff}} \quad (b_a^{\alpha})_{\text{eff}} = N_c b_a^{\text{rep}}. \tag{19}$$

In the above equation b_a^{rep} is proportional to the quadratic Casimir operator for the matter fields in the common representation and the number of condensates, N_c, can range from zero to a maximum value determined by the condition that the gauge group presumed to be condensing must remain asymptotically free. The redefinition in (19) essentially takes the coefficients b_a^{α}, which we are free to choose in our effective Lagrangian up to the conditions given in (7), and assigns the same value to each condensate.

The variable b_a^{α} can then be eliminated in (14) in favor of $(b_a^{\alpha})_{\text{eff}}$ provided the simultaneous redefinition $c_\alpha \longrightarrow (c_\alpha)_{\text{eff}}$ is made so as to keep the product in (14) invariant:

$$\left\langle \rho_+^2 \right\rangle \sim \left| \frac{(b_a^{\alpha})_{\text{eff}}}{4\,(c_\alpha)_{\text{eff}}} \right|^{-2(b_a^{\alpha})_{\text{eff}}/b_a}. \tag{20}$$

With the assumption of universal representations for the matter fields, this implies

$$(c_\alpha)_{\text{eff}} \equiv N_c \left(\prod_{\alpha=1}^{N_c} c_\alpha \right)^{\frac{1}{N_c}} \tag{21}$$

which we assume to be an $\mathcal{O}(1)$ number, if not slightly smaller.

From a determination of the condensate value ρ using (14) the supersymmetry-breaking scale can be found by solving for the gravitino mass, given by

$$M_{3/2} = \frac{1}{3} \langle |M| \rangle = \frac{1}{4} \left\langle \left| \sum_a b_a u_a \right| \right\rangle. \tag{22}$$

In [3] it was shown that in the case of multiple gaugino condensates the scale of supersymmetry breaking was governed by the condensate with the largest one-loop beta-function coefficient. Hence in the following it is sufficient to consider the case with just one condensate with beta-function coefficient denoted b_+:

$$M_{3/2} = \frac{1}{4} b_+ \langle |u_+| \rangle. \tag{23}$$

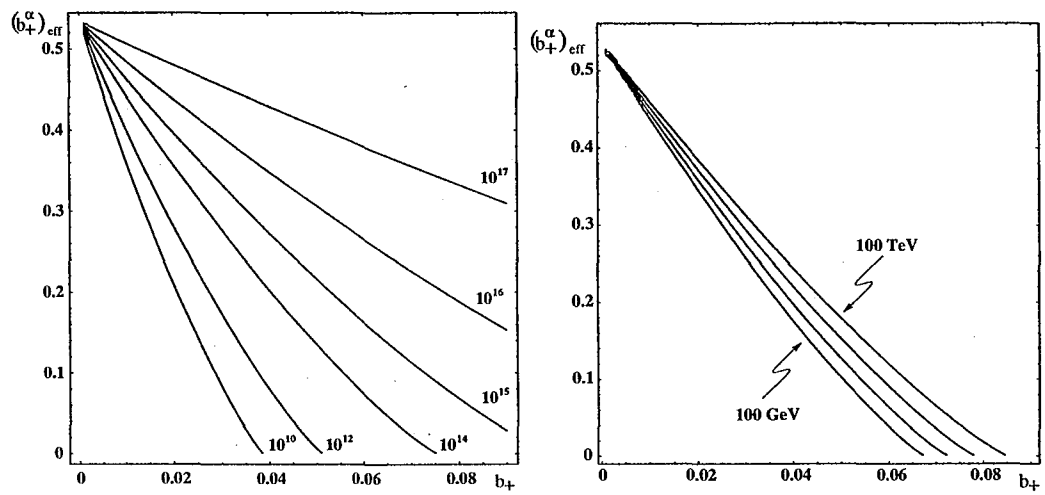

Figure 1: **Condensation Scale and Gravitino Mass.** Contours give the scale of gaugino condensation in GeV in the left panel and gravitino masses of 10^2 through 10^5 GeV in the right panel for $(c_\alpha)_{\rm eff} = 3$.

As an illustration of this point, the gravitino mass for the case of pure supersymmetric Yang-Mills $SU(5)$ condensation (no hidden sector matter fields) would be 4000 GeV. The addition of an additional condensation of pure supersymmetric Yang-Mills $SU(4)$ gauginos would only add an additional 0.004 GeV to the mass.

Now for given values of $(c_\alpha)_{\rm eff}$ and $g_{\rm str}$ the condensation scale

$$\Lambda_{\rm cond} = (M_{\rm Planck}) \left\langle \rho_+^2 \right\rangle^{1/6} \tag{24}$$

and gravitino mass (23) can be plotted in the $\left\{b_+, \left(b_+^\alpha\right)_{\rm eff}\right\}$ plane. The sharp variation of the condensate value with the parameters of the theory, as anticipated by the functional form in (14), is apparent in the contour plot of Figure 1.

The dependence of the gravitino mass on the group theory parameters is even more profound. Figure 1 gives contours for the gravitino mass between 100 GeV and 100 TeV. Clearly, the region of parameter space for which

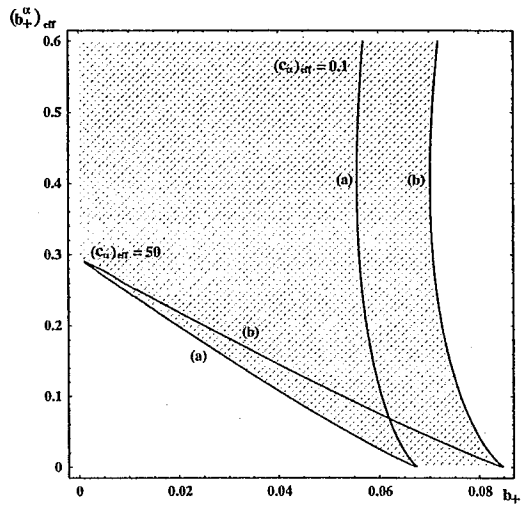

Figure 2: **Gravitino Mass Regions as a Function of Yukawa Parameter**. Gravitino mass contours for (a) 100 GeV and (b) 10 TeV are shown for $(c_\alpha)_{\text{eff}} = 50$ and $(c_\alpha)_{\text{eff}} = 0.1$ with $\alpha_{\text{str}} = 0.04$. The region between the two sets of curves can be considered roughly the region of phenomenological viability.

a phenomenologically preferred value of the supersymmetry-breaking scale occurs is a rather limited slice of the entire space available.

The variation of the gravitino mass as a function of the Yukawa parameters c_α is shown in Figure 2. On the horizontal axis there are no matter condensates ($b_a^\alpha = 0,\ \forall\alpha$) so there is no dependence on the variable $(c_\alpha)_{\text{eff}}$. For values of $(c_\alpha)_{\text{eff}} < 0.1$ the contours of gravitino mass in the TeV region lie beyond the limiting value of $b_+ \simeq 0.09$ and are thus in a region of parameter space which is inaccessible to a model in which the unified coupling at the string scale is $\alpha_{\text{str}} = 0.04$ or larger. For very large values of the effective Yukawa parameter the gravitino mass contours approach an asymptotic value very close to the case shown here for $(c_\alpha)_{\text{eff}} = 50$. We might therefore consider the shaded region between the two sets of contours as roughly the maximal region of viable parameter space for a given value of the unified coupling at the string scale.

2.2 Implications for the Hidden Sector

Having examined some of the universal constraints placed on any string-derived model proposing to describe low energy physics in Section 2.1 it is natural to ask whether the region of phenomenological viability (roughly the shaded area in Figure 2) can be used to constrain the matter content of the hidden sector.

Upon orbifold compactification the E_8 gauge group of the hidden sector is presumed to break to some subgroup(s) of E_8 and the set of all such possible breakings has been computed for Z_N orbifolds [20]. Under the working assumption that only the subgroup with the largest beta-function coefficient enters into the low-energy phenomenology, there are then a finite number of possible groups to consider:

$$\begin{cases} E_7, E_6 \\ SO(16), SO(14), SO(12), SO(10), SO(8) \\ SU(9), SU(8), SU(7), SU(6), SU(5), SU(4), SU(3) \end{cases} \tag{25}$$

For each of the above gauge groups equations (7) and (8) define a line in the $\left\{b_+, \left(b_+^\alpha\right)_{\text{eff}}\right\}$ plane. These lines will all be parallel to one another with horizontal intercepts at the beta-function coefficient for a pure Yang-Mills theory. The vertical intercept will then indicate the amount of matter required to prevent the group from being asymptotically free, thereby eliminating it as a candidate source for the supersymmetry breaking described in Section 2.1.

In Figure 3 we have overlaid these gauge lines on a plot similar to Figure 2. We restrict the Yukawa couplings of the hidden sector to the more reasonable range of $1 \leq (c_\alpha)_{\text{eff}} \leq 10$ and give three different values of the string coupling at the string scale. The choice of string coupling constant is made when specifying the boundary conditions for solving the dilaton scalar potential, as described in Section 1.2. Changing this boundary condition will affect the scale of gaugino condensation through equation (14), altering the

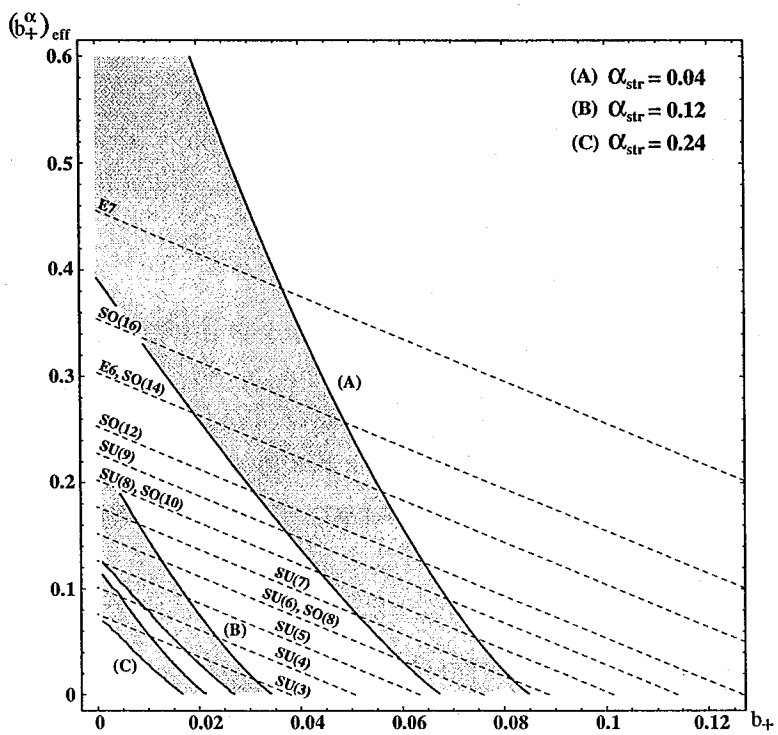

Figure 3: **Constraints on the Hidden Sector.** The shaded regions give three different "viable" regions depending on the value of the unified coupling strength at the string scale. The upper limit in each case represents a 10 TeV gravitino mass contour with $(c_\alpha)_{\text{eff}} = 1$, while the lower bound represents a 100 GeV gravitino mass contour with $(c_\alpha)_{\text{eff}} = 10$.

supersymmetry-breaking scale for a fixed point in the $\left\{b_+, \left(b_+^\alpha\right)_{\text{eff}}\right\}$ plane. Demanding larger values of g_{str} will result in the shifting of the contours of fixed gravitino mass towards the origin, as in Figure 3. Such large values of α_{str} have recently been invoked as part of a mechanism for stabilizing the dilaton and/or as a consequence of reconciling the apparent scale of gauge unification in the minimal supersymmetric standard model (MSSM) with the scale predicted by string theory [21]. We will return to such issues in Section 4.

A typical matter configuration would be represented in Figure 3 by a point

on one of the gauge group lines. As each field adds a discrete amount to $(b_a^\alpha)_{\rm eff}$ and the fields must come in gauge-invariant multiples, the set of all such possible hidden sector configurations is necessarily a finite one.[2] The number of possible configurations consistent with a given choice of $\{\alpha_{\rm str}, (c_\alpha)_{\rm eff}\}$ and supersymmetry-breaking scale $M_{3/2}$ is quite restricted. For example, Figure 3 immediately rules out hidden sector gauge groups smaller than SU(6) for weak coupling at the string scale ($g^2_{\rm str} \simeq 0.5$). Furthermore, even moderately larger values of the string coupling at unification become increasingly difficult to obtain as it is necessary to postulate a hidden sector with very small gauge group and particular combinations of matter to force the beta-function coefficient to small values.

3 Constraints from the Low-Energy Spectrum

3.1 Soft Supersymmetry-Breaking Terms

Simply requiring that the scale of supersymmetry breaking be in a reasonable range of energy values (*i.e.* within an order of magnitude of 1 TeV) can put significant constraints on the dynamics of the hidden sector. Requiring further that the *pattern* of supersymmetry breaking be consistent with observed electroweak symmetry breaking and direct experimental bounds on superpartner masses can restrict the parameter space even more.

The pattern of supersymmetry breaking is determined by the appearance of soft scalar masses, gaugino masses and trilinear couplings at the condensation scale. The gaugino masses in the one-condensate approximation, including the contribution from the quantum effects of light fields arising at

[2]For example, one cannot obtain values of b_+ arbitrarily close to zero in practical model building.

one loop from the superconformal anomaly, are given by [5]

$$m_{\lambda_a}|_{\mu=\Lambda_{\rm cond}} = -\frac{g_a^2(\mu)}{2}\left[\frac{3b_+(1+b'_a\ell)}{1+b_+\ell} - 3b_a + \sum_A \frac{C_a^A p_A(1+b_+\ell)}{4\pi^2 b_+(1+p_A\ell)}\right] M_{3/2}. \tag{26}$$

The incorporation of scalar masses and trilinear terms in the scalar potential for observable sector matter fields Φ^A depends on the form of the Kähler potential and the nature of the couplings of observable sector matter fields to the Green-Schwarz counterterm. Adopting the Kähler potential assumed in (2) and the counterterm of (13), the scalar masses are given in the one-condensate approximation by

$$m_A^2 = \frac{1}{16}\left\langle \rho_+^2 \frac{(p_A - b_+)^2}{(1+p_A\ell)^2}\right\rangle, \tag{27}$$

and the trilinear "A-terms" in the scalar potential are given by

$$V_A(\phi) = Ae^{K/2}W(\phi) \tag{28}$$

with

$$A = \frac{3}{4}\langle\bar{u}_+\rangle\left\langle\frac{(p_A - b_+)}{(1+p_A\ell)} + \frac{b_+}{(1+b_+\ell)}\right\rangle. \tag{29}$$

As noted in [4], the fact that (27) and (29) are independent of the modular weights q_I^A of the individual observable sector fields is the result of the vanishing of the auxiliary fields F^I in the vacuum. This is a manifestation of the so-called "dilaton dominated" scenario of supersymmetry breaking [22] for which flavor-changing neutral currents might be naturally suppressed. For this to indeed occur, however, it is also necessary to make the assumption that the couplings p_A are the same for the first and second generations of matter.

To analyze the low-energy particle spectrum it is necessary to choose a value of p_A for each generation of matter fields. If the Green-Schwarz term (13) is independent of the Φ^A so that $p_A = 0$, then from (27) $m_A = M_{3/2}$. We will call such a generation "light." On the other hand, it was postulated

in [4] that the Green-Schwarz term may well depend only on the combination $T^I + \bar{T}^I - \sum_A \left|\Phi_I^A\right|^2$, where Φ_I^A represents untwisted matter fields. Then for these multiplets $p_A = b$ and the scalar masses for these fields are in general an order of magnitude greater than the gravitino mass. We will call these generations "heavy."

The scalar masses (27) and A-terms (29) given above do not include the contributions proportional to the matter field wave-function renormalization coefficients arising from the superconformal anomaly (the analog to the gaugino mass terms studied in [5] and included in (26)). A systematic treatment of these contributions to the soft-breaking terms is currently underway [23], but their general size is comparable to the gaugino masses. In the following it has been checked that varying the initial soft terms by arbitrary amounts of this size has a negligible impact on the conclusions we report here.

Before giving the results of a numerical analysis using the renormalization group equations (RGEs) with the boundary conditions determined by equations (26), (27) and (29), it is worthwhile looking at what patterns of symmetry breaking are expected for various choices of the parameter p_A in the context of the MSSM. For any generation with non-negligible Yukawa couplings a good indicator that the stable minimum of the scalar potential will yield correct electroweak symmetry breaking is the relation

$$|A_t|^2 \leq 3\left(m_Q^2 + m_U^2 + m_{H_u}^2\right). \tag{30}$$

When this bound is not satisfied it is typical to develop minima away from the electroweak symmetry breaking point in a direction in which one of the scalar masses carrying electric or color charge becomes negative. For any heavy matter generation with a non-negligible coupling to a heavy Higgs field ($p_A = b$) equation (29) yields $A \approx 3m_A$ and so (30) is already nearly saturated at the condensation scale.

Another key factor in preventing dangerous color and charge-breaking minima is the ratio of scalar masses to gaugino masses and the degree of splitting between any light and heavy matter generations. In this model, both

of the hierarchies, $m_A^{\text{light}}/m_\lambda$ and $m_A^{\text{heavy}}/m_A^{\text{light}}$, will turn out to be $\mathcal{O}(10)$. This pattern of soft supersymmetry-breaking masses has been shown [24] to lie on the boundary of the region in MSSM parameter space for which light squark masses tend to be driven negative by two-loop effects arising from the heavier squarks. All of the above considerations suggest that compactification scenarios in which the observable sector matter fields couple universally to the Green-Schwarz counterterm with $p_A = b$ may have trouble reproducing the correct pattern of low-energy symmetry breaking.

3.2 RGE Viability Analysis Within the MSSM

To determine what region of parameter space in the $\left\{b_+, \left(b_+^\alpha\right)_{\text{eff}}\right\}$ plane is consistent with current experimental data it is necessary to run the soft supersymmetry-breaking parameters of equations (26), (27) and (29) from the condensation scale to the electroweak scale using the renormalization group equations. For this purpose we take the MSSM superpotential and matter content for the observable sector, keeping only the top, bottom and tau Yukawa couplings. In order to capture the significant two loop contributions to gaugino masses and scalar masses these parameters are run at two loops, while the other parameters are evolved using the one-loop RGEs. The equations used are in the $\overline{DR}$ scheme and can be found in [25]. The RGE analysis was performed on four different scenarios:

- **Scenario A**: All three generations light.
- **Scenario B**: Third generation light, first and second generations heavy.
- **Scenario C**: All three generations heavy.
- **Scenario D**: All matter heavy except for the two Higgs doublets which remain light ($p_A = 0$).

To protect against unwanted flavor changing neutral currents we have chosen the Green-Schwarz coefficients p_A to be universal throughout each matter

generation. While our scalars will turn out to be heavy enough that small deviations from universality (such as those arising from the superconformal anomaly discussed above) will not be problematic, the large hierarchies controlled by the values of the p_A would be untenable. The Higgs fields will be taken to couple to the Green-Schwarz counterterm identically to the third generation of matter, as we keep only the third generation Yukawa couplings in the MSSM superpotential. In Scenario D we relax this assumption.

In the boundary values of (26), (27) and (29) the values of $(c_\alpha)_{\text{eff}}$ and $(b_a^\alpha)_{\text{eff}}$ appear only indirectly through the determination of the value of the condensate $\left\langle \rho_+^2 \right\rangle$. It is thus convenient to cast all soft supersymmetry-breaking parameters in terms of the values of b_+ and $M_{3/2}$ using equation (23). While the gravitino mass itself is not strictly independent of b_+, it is clear from Figure 2 that we are guaranteed of finding a reasonable set of values for $\left\{(c_\alpha)_{\text{eff}}, \left(b_+^\alpha\right)_{\text{eff}}\right\}$ consistent with the choice of b_+ and $M_{3/2}$ provided we scan only over values $b_+ \leq 0.1$ for weak string coupling. This transformation of variables allows the slice of parameter space represented by the contours of Figure 3 to be recast as a two-dimensional plane for a given value of $\tan\beta$ and $\text{sgn}(\mu)$. The condensation scale (the scale at which the RG-running begins) is also a function of the gravitino mass in this framework, found by inverting equation (23).

Having chosen a set of input parameters $\left\{b_+, M_{3/2}, \tan\beta, \text{sgn}(\mu)\right\}$ for a particular scenario, the model parameters are run from the condensation scale Λ_{cond} given by (24) to the electroweak scale $\Lambda_{\text{EW}} = M_Z$, decoupling the scalar particles at a scale approximated by $\Lambda_{\text{scalar}} = m_A$. While treating all superpartners with a common scale sacrifices precision for expediency, the results presented below are meant to be a first survey of the phenomenology of this class of models.

At the electroweak scale the one-loop corrected effective potential $V_{1-\text{loop}} = V_{\text{tree}} + \Delta V_{\text{rad}}$ is computed and the effective mu-term $\bar{\mu}$ is calculated

$$\bar{\mu}^2 = \frac{\left(m_{H_d}^2 + \delta m_{H_d}^2\right) - \left(m_{H_u}^2 + \delta m_{H_u}^2\right)\tan\beta}{\tan^2\beta - 1} - \frac{1}{2}M_Z^2. \tag{31}$$

In equation (31) the quantities δm_{H_u} and δm_{H_d} are the second derivatives of the radiative corrections ΔV_{rad} with respect to the up-type and down-type Higgs scalar fields, respectively. These corrections include the effects of all third-generation particles. If the right hand side of equation (31) is positive then there exists some initial value of μ at the condensation scale which results in correct electroweak symmetry breaking with $M_Z = 91.187$ GeV [26].[3]

A set of input parameters will then be considered viable if at the electroweak scale the one-loop corrected mu-term $\bar{\mu}^2$ is positive, the Higgs potential is bounded from below, all matter fields have positive scalar mass-squareds and the spectrum of physical masses for the superpartners and Higgs fields satisfy the selection criteria given in Table 1.[4]

The first condition to be imposed on the scenarios considered here is correct electroweak symmetry breaking, defined by (31), with no additional scalar masses negative. This criterion alone rules out Scenario C, with all three generations coupling universally to the GS-counterterm and having large scalar masses. For the opposite case of no coupling to the GS-counterterm (Scenario A) the allowed region is displayed in Figure 4. In this scenario electroweak symmetry breaking requires $1.65 < \tan\beta < 4.5$, the lower bound being the value for which the top quark Yukawa coupling de-

[3]Note that we do not try to specify the origin of this mu-term (nor its associated "B-term") and merely leave them as free parameters in the theory – ultimately determined by the requirement that the Z-boson receive the correct mass.

[4]Though the inclusive branching ratio for $b \to s\gamma$ decays was not used as a criterion, an *a posteriori* check of the region of the parameter space where this class of models wants to live – namely relatively low $\tan\beta$ and gaugino masses with high scalar masses – indicates no reason to fear a conflict with the bounds from CLEO except possibly in the case $\text{sgn}(\mu) = -1$ for Scenario D [27].

Table 1: Superpartner and Higgs mass constraints imposed [28].

Gluino Mass	$m_{\tilde{g}}$	$>$	175 GeV
Lightest Neutralino Mass	$m_{\tilde{N}_1}$	$>$	15 GeV
Lightest Chargino Mass	$m_{\tilde{\chi}_1^\pm}$	$>$	70 GeV
Squark Masses	$m_{\tilde{q}}$	$>$	175 GeV
Slepton Masses	$m_{\tilde{l}}$	$>$	50 GeV
Light Higgs Mass	m_h	$>$	80 GeV
Pseudoscalar Higgs Mass	m_A	$>$	65 GeV

velops a Landau pole below the condensation scale. This restricted region of the $\tan\beta$ parameter space is a result of the large hierarchy between gaugino masses and scalar masses in these models and has been observed in more general studies of the MSSM parameter space [29].

Scenario B with its split generations can exist only for $0.08 \leq b_+ \leq 0.09$, where the hierarchy between the generations is small enough to prevent the two-loop effects of the heavy generations from driving the right-handed top squark to negative mass-squared values. Furthermore, proper electroweak symmetry breaking in this model requires the value of $\tan\beta$ to be in the uncomfortably narrow range $1.65 \leq \tan\beta \leq 1.75$, making this pattern of Green-Schwarz couplings phenomenologically unattractive.

As for Scenario D, the large third generation masses give an additional downward pressure on the Higgs mass-squareds in the running of the RGEs, allowing for a much wider allowed range of $\tan\beta$. In fact, electroweak symmetry is radiatively broken in the entire range of parameter space. However, as the value of b_+ is raised past the critical range $b_+ \simeq 0.08$, the scalar mass boundary values at the condensation scale start to become light enough that the right-handed stop is again driven to negative mass-squared values. This is shown in Figure 4 where the region between the upper and lower curves is excluded. While this region expands rapidly as the beta-function coeffi-

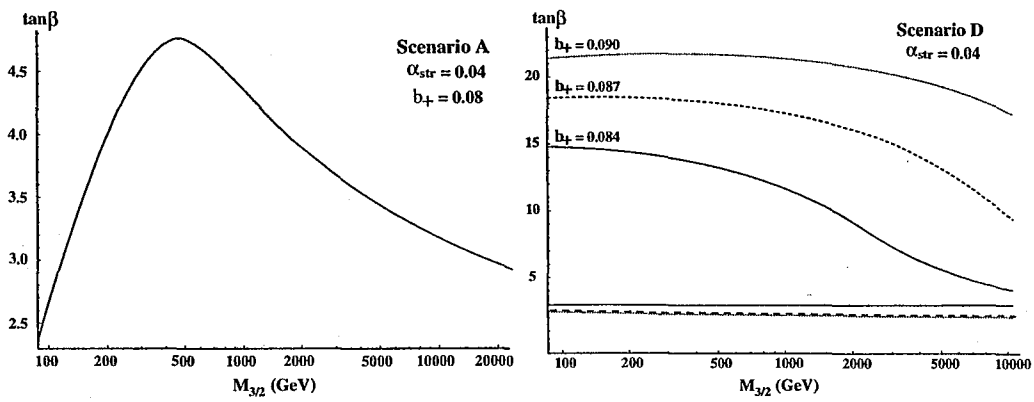

Figure 4: **Region with Correct Symmetry Breaking for Scenarios A and D.** The left panel gives the maximum value of $\tan\beta$ consistent with electroweak symmetry breaking and positive squark masses displayed as a function of the gravitino mass. The plot is shown with $b_+ = 0.08$ but the values are extremely insensitive to the choice of this parameter. The right panel shows three pairs of curves for $b_+ = 0.084, 0.087, 0.090$. For values of $\tan\beta$ between the curves the heavy scalar contribution at two loops to the running of $m^2_{U_3}$ drives its value negative.

cient is increased, the values of the beta-function coefficient consistent with $\alpha_{\rm str} \geq 0.04$ are nearly saturated when this effect arises.

The direct experimental constraints are most binding for the gaugino sector as they are by far the lightest superpartners in this class of models. Typical bounds reported from collider experiments are derived in the context of universal gaugino masses with a relatively large mass difference between the lightest chargino and the lightest neutralino. For most choices of parameters in the models studied here this is a valid assumption, but when the condensing group beta-function coefficient b_+ becomes relatively small (i.e. similar in size to the MSSM hypercharge value of $b_{U(1)} = 0.028$) the pieces of the gaugino mass arising from the superconformal anomaly (26) can become equal in magnitude to the universal term. Here there is a level crossing in the neutral gaugino sector. The lightest supersymmetric parti-

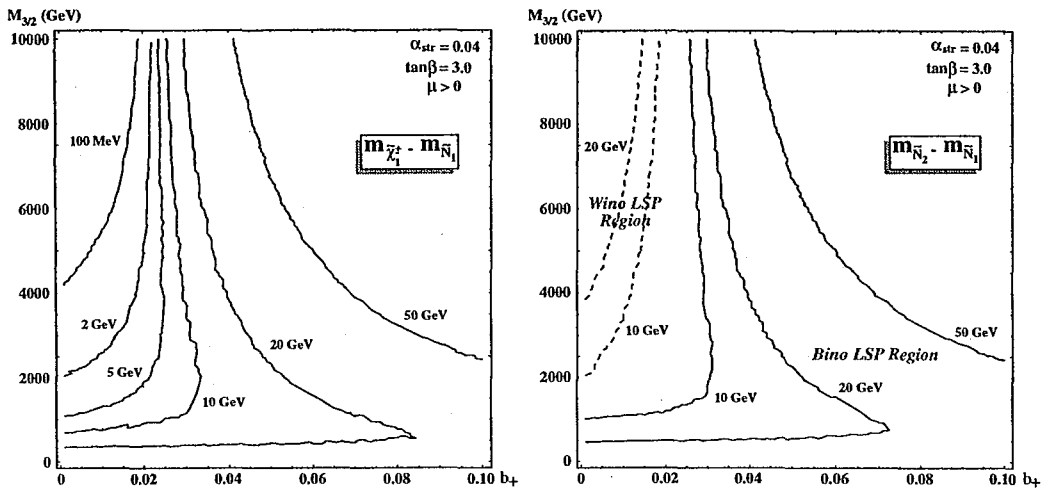

Figure 5: **The Physical Gaugino Sector in Scenario A.** The left panel gives the mass difference between the $\chi_1^{\pm}$ and the N_1 in GeV. Typical search algorithms at colliders assume a mass difference at least as large as 2 GeV. The right panel gives the difference in mass between the two lightest neutralinos N_2 and N_1. Note that a level crossing occurs and there exists a region in which the W_0 becomes the LSP, as is to be expected when the anomaly contribution to gaugino masses dominates.

cle (LSP) becomes predominately wino-like and the mass difference between the lightest chargino and lightest neutralino becomes negligible. This effect is displayed in Figure 5. The experimental constraints as normally quoted from LEP and the Tevatron cannot be applied in the region where the mass difference between the lightest neutralino and chargino falls below about 2 GeV. The phenomenology of such a gaugino sector has been studied recently in [30]. Note that when any scalar fields couple to the GS-counterterm (as in Scenario D) there is a large additional, universal contribution to the gaugino masses at the condensation scale in (26). This eliminates any region with a non-standard gaugino sector in these cases.

Figure 6 gives the binding constraints from Table 1 for Scenario A with $\tan\beta = 3$ and positive μ (the most restrictive case). The most critical

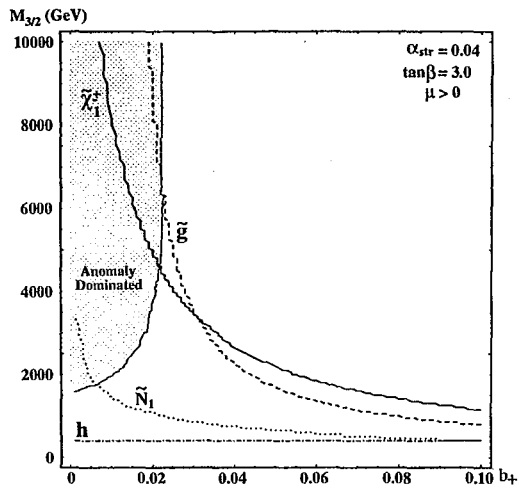

Figure 6: **Constraints from Table 1 for Scenario A.** Exclusion curves for lightest chargino (solid), gluino (dashed), lightest neutralino (dotted) and lightest Higgs mass (dashed-dotted) for weak coupling at the string scale. The region below the curves fails to meet the corresponding constraint from Table 1. The upper left corner represents the region where the difference in mass between the $\chi_1^{\pm}$ and the N_1 falls below 2 GeV and is thus not subject to the same observational constraints as standard minimal supergravity models.

constraints are for the lightest chargino and gluino.[5] The effect of varying $\tan\beta$ on these bounds is negligible over the range $1.65 < \tan\beta < 4.5$, as its effect is solely in the variation in the Yukawa couplings appearing at two loops in the gaugino mass evolution. The region for which the anomaly-induced contributions to the gaugino masses make the normal experimental constraints inoperative is represented by the shaded region in the upper left of the figure. In general, the light gaugino masses at the condensation scale require a large gravitino mass (and hence, a large set of soft scalar masses

[5]The gluino mass determination takes into account the difference between the running mass (M_3) and the physical gluino mass [31]. This difference is neglected for the other mass parameters.

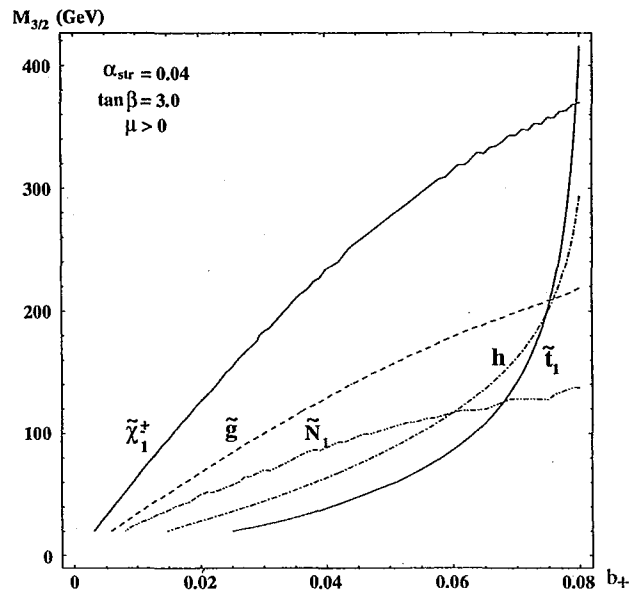

Figure 7: **Constraints from Table 1 for Scenario D**. Exclusion curves for lightest chargino (thick solid), gluino (dashed), lightest neutralino (dotted), lightest Higgs (dashed-dotted) and lightest stop (thin solid) mass. Curves are for weak coupling at the string scale. The region below the curves fails to meet the corresponding constraint from Table 1.

since $m_A = M_{3/2}$ in this scenario) in order to evade the observational bounds coming from LEP and the Tevatron. While current theoretical prejudice would disfavor such large soft scalar masses, this pattern of soft parameters may not necessarily be a sign of excessive fine-tuning [32]. Nevertheless, we refrain from making any statements about the "naturalness" of this class of models as we have not specified any mechanism for generating the mu-term.

Figure 7 gives the binding constraints from Table 1 for Scenario D with $\tan\beta = 3$ and positive μ. Note the change of scale in both axes for these plots relative to those of Scenario A. As in Figure 6, varying $\tan\beta$ over the range $1.65 < \tan\beta < 40$ has a negligible effect on the gaugino constraint contours and only a very small effect on the contours of constant stop mass. Here the gaugino masses start at much larger values so a lower supersymmetry-

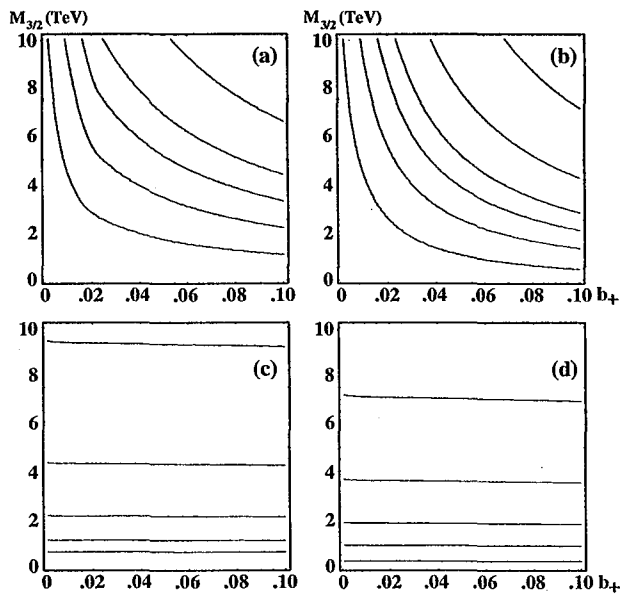

Figure 8: **Mass Contours for Scenario A.** *Panel A:* Contours for the lightest neutralino mass of 40, 80, 120, 160 and 240 GeV. *Panel B:* Contours of lightest chargino mass of 40, 80, 120, 160, 240 and 400 GeV. *Panel C:* Contours of lightest Higgs mass of 90, 100, 110, 120 and 130 GeV. *Panel D:* Contours of lightest stop mass of 200, 500, 1000, 2000 and 4000 GeV. All contours increase from the bottom to the top of each panel.

breaking scale is sufficient to evade the bounds from LEP and the Tevatron. Though the gravitino mass can now be much smaller, recall that the scalars in this scenario have masses at the condensation scale roughly an order of magnitude larger than the gravitino. Thus the typical size of scalar masses at the electroweak scale continues to be about 1 TeV for the first two generations and a few hundred GeV for the third generation scalars. As opposed to the case where all the matter fields of the observable sector decouple from the GS-counterterm, here smaller values of the condensing group beta-function coefficient *enhance* the gaugino masses via the last term in (26).

We end this section by giving mass contours for the lightest Higgs, chargino, neutralino and top-squark for $\tan\beta = 3$ and positive μ for Scenarios A and D

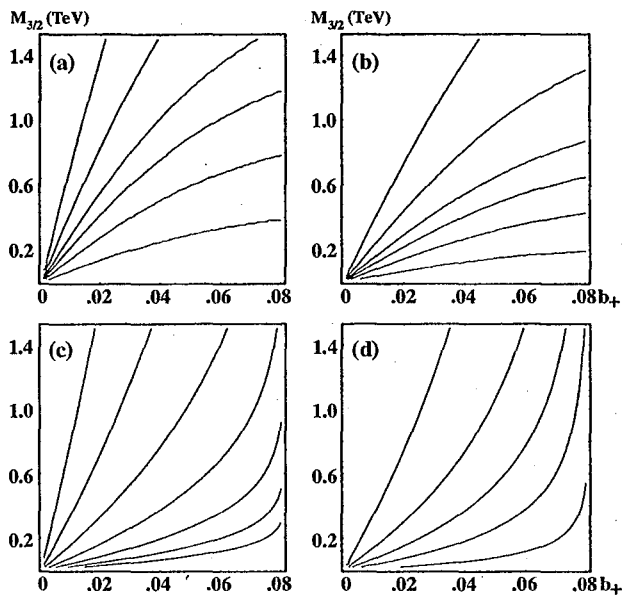

Figure 9: **Mass Contours for Scenario D.** *Panel A:* Contours for the lightest neutralino mass of 40, 80, 120, 160, 240 and 400 GeV. *Panel B:* Contours of lightest chargino mass of 40, 80, 120, 160, 240 and 400 GeV. *Panel C:* Contours of lightest Higgs mass of 80, 90, 100, 110, 120, 130 and 140 GeV. *Panel D:* Contours of lightest stop mass of 200, 500, 1000, 2000 and 4000 GeV. All contours increase from the bottom to the top of each panel.

in Figures 8 and 9, respectively.

4 Gauge Coupling Unification

In Section 2.2 the possibility of larger values of the unified coupling constant $g^2_{\rm str}$ at the string scale was considered in a very general way. It is well known [33] that the apparent unification of coupling constants at a scale $\Lambda_{\rm MSSM} \approx 2 \times 10^{16}$ GeV, assuming only the MSSM field content, is at odds

with the string prediction that unification must occur at a scale given by

$$M_{\text{string}}^2 = \lambda g_{\text{string}}^2 M_{\text{Planck}}^2 \tag{32}$$

where λ represents the (scheme-dependent) one-loop correction from heavy string modes. In [3] this factor was computed for the $\overline{MS}$ scheme and it is given by

$$\lambda = \frac{1}{2}\left(f+1\right)e^{g-1}. \tag{33}$$

For the vacua considered in this work this parameter is typically $\lambda \sim 0.19$.

Even after taking into account one-loop string corrections there is still an order of magnitude discrepancy between the scale of unification predicted by string theory and the apparent scale of unification as extrapolated from low energy measurements under the MSSM framework. One possible solution to the problem is the inclusion of additional matter fields in incomplete multiplets of SU(5) at some intermediate scale which will alter the running of the coupling constants, causing them to converge at some value higher than Λ_{MSSM} [34]. These solutions tend to involve slightly larger values of the coupling constant at the string scale than that of the MSSM ($\alpha_{\text{MSSM}}^{-1} \approx 24.7$).

In the model in question here, the intermediate scale (Λ_{cond}) at which this additional matter might appear is not independent of the scale of the superpartner spectrum ($\Lambda_{\text{SUSY}} \sim M_{3/2}$), but the two are in fact related by equation (23). Thus if we assume this additional matter has a typical mass of the condensation scale, each point in the $\left\{b_+, M_{3/2}\right\}$ plane can be tested for potential compatibility with string unification given a certain set of additional matter fields. We will not specify the origin of these fields (though such incomplete multiplets are not uncommon in string theory compactifications), but merely posit their existence with masses on the order of the condensation scale.

Our procedure for carrying out this investigation is similar to that used in the literature by a number of authors [35]. The standard model coupling constants α_3, α_2 and α_1 are determined from $\alpha_{\text{EM}}\left(M_Z\right) = 1/127.9$,

$\alpha_3(M_Z) = 0.119$ and $\sin^2\theta_{\rm EW}(M_Z) = .23124$ and these $\overline{MS}$ values are converted to the $\overline{DR}$ scheme. As we will not be concerned with performing a precision survey, these coupling constants are run at one loop from their values at the electroweak scale using only the standard model field content up to the scale $\Lambda = M_{3/2}$. At this scale the entire supersymmetric spectrum is added to the equations until the scale $\Lambda = \Lambda_{\rm cond}$ is reached. Here incomplete multiplets of SU(5) are added and the couplings are run to the scale at which the SU(2) and U(1) fine structure constants coincide. This scale will be defined as the string scale.

We now require $\alpha_3 = \alpha_2 = \alpha_1$ at this scale and invert equation (32) to find the implied Planck scale. Consistency requires that this value be the reduced Planck mass of 2.4×10^{18} GeV *and* that the QCD gauge coupling, when the renormalization group equations are solved in the reverse direction, give a value for α_3 at $\Lambda = M_Z$ within two standard deviations of the measured value.[6]

The results of the analysis for a typical choice of extra matter fields are shown in Figure 10, where a pair of vector like $\left(Q, \bar{Q}\right)$ and two pairs of vector-like $\left(D, \bar{D}\right)$'s are introduced at the condensation scale with quantum numbers identical to their MSSM counterparts. The two sigma window about the current best-fit value of α_3 can indeed accomodate a consistent Planck mass while allowing for perturbative unification of gauge couplings. From this base configuration additional **5**s and **10**s of SU(5) can be added at will to increase the value of the unified coupling at the string scale, but the contours of constant implied Planck mass shown in Figure 10 will not move significantly. While these combinations of matter fields have been known to allow for gauge coupling unification for some time [34], the relationships (23) and (32)

[6]It is worth remarking that even the celebrated supersymmetric SU(5) unification of couplings fails to predict the strong coupling at the electroweak scale at the level of two sigma and calls for a rather large value of $\alpha_3(M_Z)$[35]. This is usually taken as an indication of the size of model-dependent threshold corrections. We therefore demand no more from the models considered here.

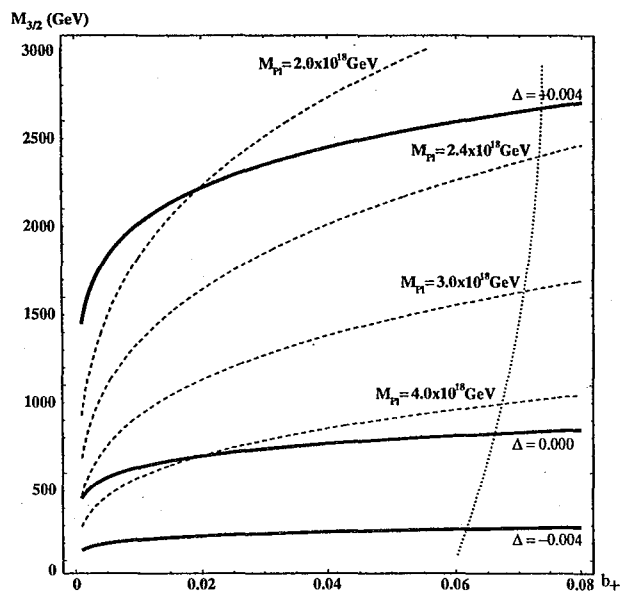

Figure 10: **Gauge Coupling Unification.** Results of adding one pair of $(Q, \bar{Q})$ and two pairs of $(D, \bar{D})$ at the condensation scale. Contours of constant implied Planck mass are overlaid on the region for which $\Delta = \alpha_3^{\rm RGE} - \alpha_3^{\rm obs}$ is within the two-sigma experimental limit of $\delta\alpha = \pm 0.004$. The dotted line represents the maximum value of b_+ consistent with $M_{3/2} \leq 10$ TeV and the RGE determined string coupling. The values of α_{str} here range from 0.044 at the $\Delta = +0.004$ contour to 0.050 at the $\Delta = -0.004$ contour.

between the various scales involved makes this a nontrivial accomplishment for this class of models.

Conclusion

The preceeding pages should be cause for guarded optimism with regard to string phenomenology. The initial challenge of dilaton stabilization has been met without resorting to strong coupling in the effective field theory nor requiring delicate cancellations. Reasonable values of the supersymmetry-breaking scale can be achieved over a fairly large region of the parameter

space, but a given combination of coupling strength at the string scale and hidden sector matter content will single out a tantalizingly small slice of this space. These successful combinations do not destroy the potential solutions to the coupling constant unification problem by the introduction of additional matter at the condensation scale. Tighter restriction on the hidden sector will require more precise knowledge of the size of Yukawa couplings in the corresponding superpotential.

Requiring a vacuum configuration which gives rise to successful electroweak symmetry breaking seems to demand that either the Green-Schwarz counterterm be independent of the matter fields or that all matter fields couple in a universal way but that the Higgs fields are distinct. The pattern of soft supersymmetry-breaking parameters in the former case pushes the theory towards large gravitino masses and very low values of $\tan\beta$. The low gaugino masses relative to scalar masses favors larger beta-function coefficients for the condensing group of the hidden sector, while smaller values may result in phenomenology in the gaugino sector similar to that of the "anomaly dominated" scenarios.

In the latter case a proper vacuum configuration and weak coupling at the string scale leave the value of $\tan\beta$ free to take its entire range of possible values. Larger beta-function coefficients for the condensing group allow a promising region with relatively light scalar partners of the third-generation matter fields and light gauginos.

A more realistic model may alter these results to some degree and uncertainty remains in the general size and nature of the Yukawa couplings of the hidden sector of these theories. Nevertheless this survey suggests that eventual measurement of the size and pattern of supersymmetry breaking in our observable world may well point towards a very limited choice of hidden sector configurations (and hence string theory compactifications) compatible with low energy phenomena.

Acknowledgements

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