Mediation of supersymmetry breaking in a class of string theory models

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Mediation of supersymmetry breaking in a class of string theory models

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Abstract: A consistent theory of supersymmetry breaking must have a hidden sector, an observable sector, and must be embedded in a locally supersymmetric theory which arises from string theory. For phenomenological reasons it must also transmit supersymmetry from the hidden to the visible sector in a dominantly flavor neutral manner. Also any such theory of supersymmetry breaking has to take into account the problem of quadratic divergences which arise once the theory is embedded in supergravity. A class of possible models that arise from GKP-KKLT type IIB string compactifications, incorporating all this while being consistent with gauge unification, with just the bare minimum of necessary supergravity/string theory moduli fields coupled to the minimally supersymmetric standard model, is presented. Such models give reasonable values for the soft masses, the μ and $B\mu$ terms and the gaugino masses. Assuming that an actual detailed realization exists, it is very likely that they are the simplest such possibility.

Keywords: Strings and branes phenomenology

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1 Introduction

Much of the discussion of supersymmetry breaking in the Minimal Supersymmetric Standard Model (MSSM) (and its generalizations) has been in the context of global supersymmetry. However it is well-known that in order to have a zero (or nearly zero) cosmological constant it is necessary to incorporate supergravity (SUGRA) effects. This is usually done by introducing a constant into the superpotential. The supergravity potential, unlike the globally supersymmetric one, is not positive definite and one can in principle use the constant to tune the cosmological constant (CC) to zero. But once it is admitted that a consistent theory needs to bring in supergravity effects, one needs to account for the potential effects of quadratic sensitivity to high scale physics of the low energy supersymmetry breaking parameters. Furthermore one has to consider effects of the additional fields (moduli) that are neutral under the standard model group but are essential ingredients in any consistent supergravity such as string theory.[1]

The moduli that occur in any string theory construction need to be stabilized, and in the recent literature there has been much discussion of how this may be done, particularly in the context of type IIB string theory.[2] In general one can find minima for the moduli sector potential which break supersymmetry. In fact generic minima would be expected to

[1]For reviews of the MSSM, SUSY breaking mechanisms and phenomenological SUGRA, see for example [36–38].

[2]For reviews see [39, 40].

have supersymmetry breaking at the natural scale of the theory - namely the string scale. Nevertheless it is possible to find non-generic points in this landscape which have a low or intermediate scale of supersymmetry breaking. If supersymmetry is to be relevant for phenomenology, the starting point of any string theory construction would have to be one of these points. In the following we will work with type IIB theory since in many respects this is the best understood, and we assume that such points exist with the MSSM living on some brane configuration. However we expect that similar arguments can be made in other string theory contexts, and we suspect that the phenomenology (assuming the relevant existence theorems can be established) is not likely to be very different since our arguments rest on some general features of the string theory input such as the tendency to have a 'no-scale' starting point at the classical level.

We work with $\kappa^2 = 8\pi G_N = M_P^{-2} = 1$. A general supergravity theory has a real analytic Kaehler potential $K = K(\Phi, \bar{\Phi})$ and a holomorphic superpotential $W(\Phi)$ where $\Phi = \{\Phi^A\}$, $A = 1, \ldots, N$, is the set of chiral scalar fields of the theory. The metric on field space is $K_{A\bar{B}} = \partial_A \partial_{\bar{B}} K$ and the metric on the gauge group (which is in general a chiral function of the neutral chiral superfields) is $f_{ab}(\Phi)$.

The embedding of a supersymmetry breaking theory in supergravity brings in additional effects that are not usually considered in the literature. The coefficient of the term that is quadratic in the cutoff in the one loop effective potential, is proportional to $\mathrm{Str}M^2(\Phi) \equiv \sum_J (-1)^{2J}(2J+1)\mathrm{tr}M^2(\Phi)$, where Φ is the set of chiral (super) fields in the theory and M^2 is the field dependent mass squared matrix. In a globally supersymmetric theory (even if the supersymmetry is spontaneously broken) this supertrace is zero and one has no quadratic divergence in the quantum theory. However in a SUGRA theory whose supersymmetry is spontaneously broken this supertrace does not vanish. Instead we have[3]

$$\mathrm{Str}M^2(\Phi) = (N-1)m_{3/2}^2(\Phi) - F^A(\Phi)(R_{A\bar{B}} + F_{A\bar{B}})(\Phi, \bar{\Phi})\bar{F}^{\bar{B}}(\bar{\Phi}), \tag{1.1}$$

where

$$R_{A\bar{B}} = \partial_A \partial_{\bar{B}} \ln \det K_{C\bar{D}}, \; F_{A\bar{B}} = -\partial_A \partial_{\bar{B}} \ln \det \Re f_{ab}, \tag{1.2}$$

and F^A is the F-term of the chiral multiplet Φ^A and $m_{3/2}^2(\Phi) = e^K|W|^2$ is the field dependent gravitino mass.

Let the indices $I, J, \ldots$ denote fields of the visible (MSSM) sector and $i, j, \ldots$ denote (closed string) moduli fields. Typically they are expected to get vaccum expectation values (vevs) of the order of the Planck scale or larger. Let the total number of such chiral superfields in the visible sector be N_v. Note that this number is taken to include GUT fields if any. In gauge mediated supersymmetry breaking (GMSB) models (for a review see [1]) there is a (gauge neutral) hidden sector which breaks supersymmetry that is distinct from the moduli sector. Let us denote the fields of this intermediate sector by indices $r, s, \ldots$. In our string theory context they could be open string moduli. In addition such models have a messenger sector which couples to this hidden sector and is charged under the gauge group and since these are expected to have only negligible F-terms we will denote them by the same indices as the MSSM fields. The distinction between these sectors may be

[3]See for example [2, 19].

understood in terms of the general formula [2] for the (unnormalized) soft mass terms in the visible sector,

$$\Delta M^2_{I\bar{J}} = -R_{I\bar{J}k\bar{l}}F^k F^{\bar{l}} - R_{I\bar{J}r\bar{s}}F^r F^{\bar{s}} - R_{I\bar{J}K\bar{L}}F^K F^{\bar{L}} + \frac{1}{3}F_I F_{\bar{J}} + K_{I\bar{J}}m^2_{3/2}. \tag{1.3}$$

(Note for simplicity of exposition we have ignored mixed terms such as $R_{I\bar{J}k\bar{L}}F^k F^{\bar{L}}$ above). The right hand side of this equation is to be evaluated at the minimum of the scalar potential. The general formula for the F-term is

$$\bar{F}^{\bar{A}} = e^{K/2}K^{\bar{A}B}D_B W = e^{K/2}K^{\bar{A}B}(\partial_B W + K_B W). \tag{1.4}$$

The different mechanisms and mediations may be distinguished in terms of the two physical scales $M_P\,(\simeq 10^{18}GeV) \to 1$ and the weak scale $G_F^{-1/2}\,(\simeq 100GeV) \to 10^{-16}$. Now the F-terms of the visible sector fields F^I are at most of the order of the squared Higgs vacuum expectation value (vev) or the Higgs vev times the gravitino mass i.e. $\sim 10^{-30}$ or $10^{-15}m_{3/2}$ (whichever is larger). On the other hand tuning the cosmological constant to zero implies

$$F^A \bar{F}^{\bar{B}} K_{A\bar{B}} - 3m^2_{3/2} = 0. \tag{1.5}$$

This means that (given that the Kaehler metric is positive definite) $|F^A| \lesssim m_{3/2}$. Now clearly the fourth term of (1.3) is much smaller (by a factor of 10^{-30}) than $m^2_{3/2}$ and so can be ignored (unless $m_{3/2} < 10^{-30}\,(\sim 10^{-6}eV)$ which we assume is not the case.

The first term in (1.3) is the classical contribution of the moduli which typically take Planck scale expectation values. In string theory for instance, in order for the four dimensional low energy approximation to be valid, these moduli must typically take values which are somewhat larger than the Planck scale. The curvature is of order one or less on the Planck scale so that the contribution of this term to the squared soft mass is at most $O(m^2_{3/2})$. This is called the moduli mediated (MMSB) contribution.

The second term can come from some hidden sector field (open string modulus) which acquires an F-term as in GMSB. As argued earlier all F-terms are $\lesssim O(m_{3/2})$. Classically the corresponding moduli space curvature is at most order one (and typically in models it is either zero or highly suppressed) and so this contribution would not dominate over the classical contribution from the (closed string) moduli sector. However there are wave function renormalization effects which effectively enhance loop effects since the moduli space curvature goes like ϕ^{-1} where ϕ is the lowest component of some scalar field. Thus the contribution of this to the soft mass is effectively like $|\epsilon F^\phi/\phi|^2$ (where $\epsilon = g^2/16\pi^2$ with g a gauge coupling). This would be enhanced over the MMSB contribution if the potential is such that ϕ, the hidden sector SUSY breaking field in GMSB, has a vev which is significantly smaller than the Planck scale (which could be the case for open string moduli). One does not really need sequestering in this case for the GMSB contribution to dominate the modulus contribution - all that is needed is that $F^\phi \leq F \leq (10^{-8})^2$ and $\phi \lesssim 10^{-3}$. The former inequality follows from (1.5) and the fact that any classical contribution to the soft masses from MMSB will be of order $m_{3/2}$ which should therefore be taken to be $O(10^{-16}) \sim 10^2 GeV$, while the latter is required so that the $O(\alpha/4\pi)$

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suppression in the flavor conserving GMSB contribution is compensated. However this gives a GMSB contribution at the same level as the MMSB contribution. So unless MMSB already has suppressed FCNC (in which case there is no need for any GMSB mechanism) we need to suppress the gravitino mass well beyond the usual MMSB value. Thus in typical GMSB models one has $F^{\phi} \sim (10^9 GeV)^2$ or less and $\phi \sim 10^{13} GeV$ or less with a gravitino mass (which in effect would be the size of possibly FCNC violating MMSB contributions) of around a few GeV or less.

The third mechanism is usually called anomaly mediated supersymmetry breaking (AMSB) [3–5] and is supposedly associated with Weyl (or conformal) anomalies in supergravity. As discussed in [6] (based on the work of [7]) this actually consists of two different contributions. One of them arises from the Weyl anomaly of the theory. This effect will be present even in the absence of matter fields - for example in supergravity coupled to super-Yang-Mills fields. In addition there is a contribution that arises from the mechanism pointed out by Dine and Seiberg [7] which in fact has nothing to do with Weyl anomalies. This is like GMSB in that the contribution to the soft masses arises from a quantum effect, but instead of having an intermediate scale sector as in GMSB, it relies on the fact that the Higgs field acquires a non-zero vev in the physical vacuum. This in turn leads to an F-term for the Higgses of the form $F^H \sim m_{3/2} H$. Now given that the classical contribution to the curvature is $O(1)$ or less, this gives a negligible contribution to the soft masses. However there is a quantum contribution which gives a moduli space curvature of the form $R \sim \epsilon^2/H^2$, giving a squared soft mass of $O(\epsilon^2 m_{3/2}^2)$. But this is usually much smaller than the contribution from MMSB and so the latter must be suppressed, i.e. the classical moduli contribution to SUSY breaking must be sequestered [3], unless its FCNC effects are negligible.

In this paper we will discuss a class of models which contain the minimal inputs that are necessary to have soft supersymmetry breaking terms in the MSSM, are consistent with the suppression of flavor violating terms, and which can be embedded in a supergravity/string theoretic framework. In section 2 we discuss the hidden sector which is responsible for supersymmetry breaking. This is the closed string moduli sector of the theory. Obviously we cannot start with a generic point on the landscape of string solutions since this will not have the tiny cosmological constant that is observed. Also it will most probably have large (i.e. string scale or Kaluza-Klein scale) supersymmetry breaking, so that we would certainly not be led to the MSSM, which in this bottom up approach is our starting point. Thus we need to be restricted to those points in the landscape which have a nearly zero cosmological constant and low energy supersymmetry breaking. We focus on those models where this happens in the simplest possible way. We will consider a moduli sector breaking supersymmetry in such a way that it is not directly passed on to the visible sector at the classical level. At one loop level there are quadratic divergences (with a cutoff that we will identify with the GUT/KK scale). This requires the retuning of the cosmological constant and it gives a flavor diagonal contribution to the soft supersymmetry breaking parameters that is proportional to $\frac{\Lambda^2}{16\pi^2} m_{3/2}^2$. In general however there is a FCNC violating term which needs to be suppressed by fine tuning the fluxes in an appropriate manner. In effect this is a derivation of a quantum version of the mSUGRA model. In addition to this there is the mechanism identified in [6, 7] which replaces what is usually presented as anomaly mediated supersymmetry breaking (AMSB). Thus the basic claim of this paper is that the simplest model of supersymmetry breaking that is consistent with all constraints (both theoretical and phenomenological) and which is independent of ad hoc uplift terms, is a version of mSUGRA which comes from an high energy quantum effect, plus the low-energy quantum effect identified in [6, 7].

2 The model

The moduli sector is taken to come from type IIB compactified on a Calabi-Yau orientifold [8, 9] with the visible sector being on a set of D3 branes. While a detailed construction of such a model is not yet available it is very plausible that one exists. Indeed it is likely that our arguments here apply to a whole class of such models since only very generic properties of such a construction are used. For simplicity we consider a model with just one Kaehler modulus T but a large number $h_{21} \gtrsim O(10^2)$ of complex structure moduli z^α, but it should be clear from the discussion that an extension to compactifications with several Kaehler moduli is straightforward. Also to stabilize the T modulus we will need non-perturbative terms as in KKLT [9].

The MSSM sector will have (schematically) quark/lepton SU(2) doublet superfields denoted by Q^i/L^i and the corresponding singlet conjugate fields U^{ci} D^{ci}, E^{ci} with i being a family index. The Higgs fields are H_u, H_d. For the Kaehler potential we take

$$\begin{aligned} K &= -3\ln(T+\bar{T}-(H_u\bar{H}_u+H_d\bar{H}_d+z^Q_{I\bar{J}}Q^I\bar{Q}^{\bar{J}}+(x_{IJ}Q^IQ^J+h.c.)) \\ &\quad -\ln(S+\bar{S})-\ln k(z,\bar{z}) \qquad (2.1) \\ &= K_{\text{mod}}+Z(T)_{I\bar{J}}\Phi^I\Phi^{\bar{J}}+\frac{1}{2}(X_{IJ}\Phi^I\Phi^J+h.c.)+\cdots \qquad (2.2) \\ K_{\text{mod}} &= -3\ln(T+\bar{T})-\ln(S+\bar{S})-\ln k(z,\bar{z}),\, Z_{I\bar{J}}=\frac{3z_{I\bar{J}}}{T+\bar{T}},\, X_{IJ}=\frac{3x_{IJ}}{T+\bar{T}}. \qquad (2.3) \end{aligned}$$

In the above Φ^I, $I=1,\dots N_v$ denotes all the visible sector fields. Note that in this model the space of dilaton-axion S and the complex structure moduli $\bar{z}^\alpha$, and the space T,Φ^I are not direct product spaces and the metric is not a direct sum of the metrics on these two spaces since x_{IJ} is in general a function of the complex structure moduli, though it vanishes when $z^\alpha=0$ [10]. Also the form given in the first line of (2.1) is valid only to linear order in the z^α. While many authors (see for example [11–13]) use a direct sum form for the Kaehler potential in the presence of D3 branes, this is only true if the complex structure moduli and the dilaton are frozen at zero.

For the moduli superpotential we have

$$W_{\text{mod}}=W_{\text{flux}}(S,z)+\sum_n A_n(S,z)e^{-a_nT}, \qquad (2.4)$$

while for the MSSM superpotential we take

$$W_{\text{MSSM}}=\tilde{\mu}H_uH_d+y_{uij}Q^iH_uU^{cj}+y_{Dij}Q^iH_dD^{cj}+y_{Eij}L^iH_dE^{cj}. \qquad (2.5)$$

In the above S is the dilaton-axion superfield and $z=\{z^\alpha\}$, $(\alpha=1,\dots,h_{21})$ denotes the set of complex structure moduli and T is the Kaehler modulus of some Calabi-Yau orientifold (with $h_{11} = 1$) compactification of type IIB string theory. The first term in (2.4) comes from internal magnetic fluxes and the second is a series of non-perturbative (NP) terms coming from condensing gauge groups associated with D7-branes [8, 9]. Also the MSSM sector is located on a stack of D3 branes. For details of the dependence of this superpotential on the closed string moduli see [10, 14]. The model has a R-parity symmetry under $\Phi(\theta) \to \pm\Phi(-\theta)$ with the plus sign for the Higgses and minus sign for quark and lepton superfields. There is also a PQ symmetry (if the μ-term is set to zero) with charges

$$PQ: Q = L = U^c = D^c = L^c = -\frac{1}{2},\ H^u = H^d = 1. \tag{2.6}$$

and all moduli having zero charge. The moduli potential is

$$\begin{aligned} V_{\rm mod} = & \frac{1}{k(z,\bar{z})(S+\bar{S})(T+\bar{T})^2}\left\{\frac{1}{3}(T+\bar{T})|\partial_T W_{\rm mod}|^2 - 2\Re\partial_T W_{\rm mod}\bar{W}_{\rm mod}\right\} \\ & + |F^S|^2 K_{S\bar{S}} + F^z F^{\bar{z}} k_{z\bar{z}}. \end{aligned} \tag{2.7}$$

Now if one ignores quantum corrections, one would want to look for a local minimum of this potential with zero cosmological constant (CC) and SUSY breaking only in the T direction,[4] i.e.

$$V_{\rm mod}|_0 = 0,\ F|_0^S = F^z|_0 = 0,\ F^T|_0 \neq 0. \tag{2.8}$$

There is certainly no obstruction to finding such a minimum and with a sufficient number of complex structure moduli and non-perturbative terms it is reasonable to expect that such a SUSY breaking minimum exists. However the T modulus - the scalar partner of the Goldstino - has zero mass if the CC is fine-tuned exactly to zero. It should be stressed though that this does *not* imply that this modulus is not stabilized, since we have included the non-perturbative terms which are explicitly T dependent. In other words the equation $\partial_T V = 0$ will have a non-trivial solution because of the first term of (2.7).[5]

In fact however as we will see in the next section the quantum corrections would have required us to re-fine-tune the CC if we had started with a zero value for it. So anticipating this what we really have to do is to start at the classical level by fine-tuning the CC to be a small (actually negative) value - much smaller, in absolute value, than $m_{3/2}^2$. In this case there is certainly no obstruction to having positive non-zero (squared) masses for all the moduli. Also there will be additional contributions to the masses of the visible sparticles, from the quantum corrections. In fact the sort of minimum we will start with is like the

[4]One could of course look for more general minima where the SUSY breaking is shared amongst all the moduli. However in practice in all models discussed so far one usually looks for minima in which one starts from the no-scale potential where $F_T \neq 0$, but the other F-terms are zero and then expect the addition of the non-perturbative (NP) terms that fix the T-modulus not to change this situation very much, at least for large volume compactifications (see for example [15]). Of course in the original toy model of KKLT $F_T = 0$, but this was a result of ignoring the non-trivial effects of the other moduli, which really cannot be frozen in the presence of the NP terms. See [41] for a discussion of these issues.

[5]The necessary conditions for stability in SUGRA models coming from string theory were first discussed in [30] and were generalized by Gómez-Reino et al. [42–44].

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one analyzed in the large volume scenario of [15].[6] The only difference is that unlike in that case we have to fine-tune $W_{\rm flux}$ at the minimum (by adjusting the fluxes) to a very small value, in order to get an intermediate volume scenario with the volume of the internal manifold $\mathcal{V} \sim T^{3/2} \sim 10^3$. We need this to preserve a field theoretic description up to the unifcation scale (which will be identified with the cut-off $\Lambda \sim 10^{-2}$ in the quantum theory) while having a gravitino mass in the $10TeV$ range. Furthermore (as discussed in the next section) the quantum contribution to the CC is $O(N\frac{\Lambda^2}{16\pi^2}m_{3/2}^2)$ where N is the number of chiral multiplets in the theory. Thus we expect a broken supersymmetric minimum $F|_0^S = F^z|_0 = 0$, $F^T|_0 \neq 0$, with a small negative cosmological constant $-|V_0|$ such that

$$|V_0| \sim \frac{m_{3/2}^2}{\mathcal{V}} \sim O\left(N\frac{\Lambda^2}{16\pi^2}m_{3/2}^2\right) \ll O(m_{3/2}^2). \tag{2.9}$$

It should also be stressed here that our framework does not need any ad hoc uplift terms to get an acceptable value for the CC. This will come about as a result of the fine tuning of the classical SUGRA CC against the quantum effects that are discussed below.

The curvature component relevant to the soft mass calculation in this model is $R_{T\bar{T}I\bar{J}} = \frac{1}{3}K_{T\bar{T}}Z_{I\bar{J}} + O(H^2)$ so that using the standard expression for soft masses, given for example in [16, 17], we have

$$m_{I\bar{J}}^2 = m_{3/2}^2 Z_{I\bar{J}} - F^T F^{\bar{T}} R_{T\bar{T}I\bar{J}} \sim O\left(m_{3/2}^2 \frac{\Lambda^2}{16\pi^2}\right) \ll m_{3/2}^2. \tag{2.10}$$

Similarly both the $B\mu$ and the trilinear couplings - the A-terms - are also suppressed.[7] In the next two sections we will consider the quantum effects.

3 Quadratic divergence issues and mSUGRA

It is well known that quadratic divergences are absent in (spontaneously broken) global supersymmetry, but this is not really relevant for phenomenology for well-known reasons. Any mechanism of supersymmetry breaking (such as say dynamical SUSY breaking) is incomplete unless it is embedded within supergravity. Then one needs to confront the problem of quadratic divergences. In the following we will discuss how the cosmological constant and the soft supersymmetry breaking parameters get affected by these divergences.

3.1 The cosmological constant and the soft masses

To one-loop order but keeping only the $O(\Lambda^2)$ (where Λ is the cutoff) corrections we have the following [18–20]) formulae for the potential (at a minimum) and the (unnormalized)

[6] In the analysis of [15] an α' correction is also included in the Kaehler potential though it is not really essential for the demonstration of the existence of a minimum as such. We can of course include this but have avoided doing so explicitly for simplicity since it does not change the qualitative features that we are discussing in this work.

[7] If we had fine-tuned the CC exactly to zero at the classical level we would have got zero for the soft masses and the μ, $B\mu$ and A terms as in the no-scale model.

soft mass terms.

$$V|_0 = (F^m \bar{F}^{\bar{n}} K_{m\bar{n}} - 3m_{3/2}^2)\left(1 + \frac{(N-5)\Lambda^2}{16\pi^2}\right) + \frac{\Lambda^2}{16\pi^2}(m_{3/2}^2(N-1) - F^T \bar{F}^{\bar{T}} R_{T\bar{T}}), \quad (3.1)$$

$$m_{I\bar{J}}^2 = V|_0 Z_{I\bar{J}} + (m_{3/2}^2 Z_{I\bar{J}} - F^T F^{\bar{T}} R_{T\bar{T}I\bar{J}})\left(1 + \frac{(N-5)\Lambda^2}{16\pi^2}\right)$$
$$-\frac{\Lambda^2}{16\pi^2}[m_{3/2}^2 R_{I\bar{J}} + m_{3/2}(F^T D_T R_{I\bar{J}} + F^{\bar{T}} D_{\bar{T}} R_{I\bar{J}}) + F^T F^{\bar{T}}(D_T D_{\bar{T}} R_{I\bar{J}}$$
$$-R_{\bar{T}}^{\bar{T}} R_{T\bar{T}I\bar{J}} - R_T^T R_{T\bar{T}I\bar{J}} + R_I^K R_{T\bar{T}K\bar{J}})]. \quad (3.2)$$

Here N is the total number of chiral scalar superfields. In writing these expressions we have kept, in the one loop correction terms, the classical fine tuning values (2.8) of the F-terms.

In estimating these corrections we will take the cutoff to be

$$\Lambda \sim M_{\rm GUT} \sim M_{KK} \sim 10^{16} GeV \sim 10^{-2} M_P \to \frac{\Lambda^2}{16\pi^2} \sim 10^{-6} M_P^2. \quad (3.3)$$

The first question that needs to be addressed is the fine-tuning of the cosmological constant. With the classical fine tuning (2.8) and using $R_{T\bar{T}} \simeq \frac{1}{3}(N_v+2)K_{T\bar{T}}$ (where N_v is the number of visible sector fields) we would obtain at one-loop a CC of order $\frac{\Lambda^2}{16\pi^2} m_{3/2}^2 (N - N_v - 3) = 10^{-6} m_{3/2}^2 M_P^2 (h_{21} - 1)$. Since we need the number of complex structure moduli to be of $O(10^2)$ in order to be able to fine tune the classical CC, this one loop correction leads to a CC (assuming that the gravitino is at least of order the SUSY mass splittings $10^{2-3} GeV$) that is a factor $\sim 10^{86}$ too large! Thus as we discussed before we need to change the classical starting point which ignored the fact that there are quantum corrections.[8] In other words to cancel the CC to the leading order in $N\Lambda^2/16\pi^2$ we need to add corrections to (2.8) and (2.9) such that (with $M_P = 1$)

$$3m_{3/2}^2 - F^m F^{\bar{n}} K_{m\bar{n}} = \frac{\Lambda^2}{16\pi^2}(m_{3/2}^2(N - 1 - (2 + N_v)) = \frac{\Lambda^2}{16\pi^2} m_{3/2}^2 (h_{21} - 1). \quad (3.4)$$

Note that since the r.h.s. of this equation is positive the classical CC (the negative of the l.h.s.) would have to be negative. In this case there is no obstruction to getting a positive squared mass for the T modulus and generically it will be $O(m_{3/2}^2)$.

The actual minimum around which we work in calculating the soft masses will also change the values of the F-terms of the moduli from those given in (2.8) to the following (with $|F^i| \equiv \sqrt{K_{i\bar{i}} F^i F^{\bar{i}}}$):

$$|F^T| = \sqrt{3} m_{3/2} + O\left(h_{21} \frac{\Lambda^2}{16\pi^2} m_{3/2}\right), \; |F^S| \lesssim O\left(\frac{\Lambda}{4\pi} m_{3/2}\right), \; |F^z| \lesssim O\left(\frac{1}{\sqrt{h_{21}}} \frac{\Lambda}{4\pi} m_{3/2}\right). \quad (3.5)$$

We will thus assume that one can find such a minimum by adjusting fluxes and there certainly is no obstruction to doing so.

Now let us calculate the soft masses by including the quantum corrections. The first term in (3.2) has now been re-fine-tuned to zero. However the second term is no longer

[8] For a discussion of the consequences for string phenomenology of this refinetuning problem see [45].

zero and there is an additional contribution from the third term. To calculate these we need the curvatures derived from the Kaehler potential (2.1):

$$R_{T\bar{T}I\bar{J}} = \frac{K_{T\bar{T}}z_{I\bar{J}}}{T+\bar{T}}+O(\Phi^2),\ R_{I\bar{J}K\bar{L}}=\frac{3}{(T+\bar{T})^2}(z_{I\bar{J}}z_{K\bar{L}}+z_{I\bar{L}}z_{K\bar{J}}-z_{IK}z_{\bar{L}\bar{K}})+O(\Phi^2)$$
$$R_{I\bar{J}} = \frac{N_v+1}{T+\bar{T}}z_{I\bar{J}}+O(\Phi),\ D_TR_{I\bar{J}}=O(\Phi),\ D_TD_{\bar{T}}R_{I\bar{J}}=O(\Phi)$$
$$R_I{}^K R_{T\bar{T}K\bar{J}} = \frac{N_v+1}{(T+\bar{T})^3}z_{I\bar{J}}+O(\Phi^2).$$

So (given that the MSSM fields Φ have values that are highly suppressed $< O(10^{-16})$ relative to the Planck scale) we find from (3.2) on using (3.4) that the the largest quantum contribution to the soft mass squared is

$$m'^2_{I\bar{J}} \sim (h_{21}-2N_v)\frac{\Lambda^2}{48\pi^2}m^2_{3/2}Z_{I\bar{J}} \sim (h_{21}-2N_v)10^{-6}m^2_{3/2}Z_{I\bar{J}}, \tag{3.6}$$

and is positive provided that $h_{21} > 2N_v$. It is also flavor diagonal. In fact it is precisely of the form assumed by mSUGRA models of supersymmetry breaking and is easily obtained for generic Calabi-Yau orientifold compactifications.

The flavor conserving two loop quantum corrections coming from fluctuations of light fields that we will consider in the next section, are in fact of the same order provided that the number of complex structure moduli is $O(10^2)$. In fact since N_v is also of the same order, this is a necessary condition to get positive squared masses. Of course the tuning of the cosmological constant already requires the number of cycles in the compactification manifold to be at least of this order. Thus this contribution is $O(10^{-4}m^2_{3/2})$. However if this (3.6) had been flavor violating the model (even with the flavor conserving effect of the next section) would have been in danger of being ruled out since the flavor violating effects need to be down by a factor of around 10^{-3} compared to the flavor conserving one. Note that the flavor conservation of the soft masses calculated in this section is entirely due to the fact that the visible sector field space metric factorizes into a modulus dependent factor and a matrix in generation space. This in turn is a reflection of the fact that all visible fields are from a stack of D3 branes. This would not have been the case if the visible sector came partially from D3 branes and partially from (wrapped) D7 branes for instance. Such a general embedding would have resulted in a metric $Z_{I\bar{J}} = f(M,\bar{M})z_{I\bar{J}} + g(M,\bar{M})z'_{I\bar{J}}$ where M denotes the set of moduli and the dilaton and $z_{I\bar{J}}, z'_{I\bar{J}}$ are in general different matrices so that the curvature would not have been proportional to $Z_{I\bar{J}}$, and hence we would have had flavor changing terms at an unacceptable level.

Nevertheless it should be noted that the above calculation of curvatures are done in the linearized (in the complex structure moduli z^α) solution given in [10]. It is indeed possible that the complete solution will yield a contribution to (3.6) that is not proportional to the matrix $Z_{I\bar{J}}$ and so in general will lead to fine tuning. For instance in general $z_{I\bar{J}}$ would be a function of the complex structure moduli $z^\alpha, \bar{z}^\beta$ and so there would be a contribution to the Ricci tensor in the MSSM directions of the form $R_{I\bar{J}} \sim K^{\alpha,\bar{\beta}}\partial_\alpha\partial_{\bar{\beta}}Z_{I\bar{J}}$ which is not (in general) proportional to $Z_{I\bar{J}}$. This could in principle give, from the second line of (3.2) a contribution as large as the one in (3.6). In this case we need additional fine-tuning to 1 part in 10^3 to achieve the necessary suppression of FCNC and this can be done by appropriate choices of the fluxes which determine the complex structure moduli.

It should be noted that (given the suppression of classical soft terms in our model) (3.6) by itself would give soft mass terms at an acceptable level provided that the gravitino mass is a factor of 10^2 larger than the soft mass - i.e. we would need a gravitino with $m_{3/2} \gtrsim 10TeV$. This is typical of so-called AMSB scenarios where the classical terms are "sequestered" [3] as is the case with our classical starting point (2.1). The point of our discussion here is to show that the quadratic divergences that are inevitably present, give a contribution which is competitive with the 'AMSB' effects.

As for the A terms, adding the quadratically divergent one-loop effects gives

$$A_{IJK} = e^{K_m/2} \frac{W_m^*}{|W_m|} \left\{ F^i D_i y_{IJK} \left(1 + \frac{N-5}{16\pi^2} \Lambda^2 \right) - \frac{\Lambda^2}{16\pi^2} O(F^T) \right\} \tag{3.7}$$

where the sum in the first term in parentheses excludes the T modulus (recall that the classical contribution is suppressed since in the no-scale model it would be exactly zero while here it is $O(\Lambda^2/16\pi^2)$). The second term consists of terms that are proportional to y_{IJK}. As shown in [10] the first term is proportional to y_{IJK} and hence when (due to quantum effects in our case) the F^i are turned on, no significant flavor violating effects are generated.

3.2 Consistency issues

Let us now check what the cancellation of the one-loop contribution to the cosmological constant implies for the F-terms of the moduli. Using (2.4) we have (assuming for simplicity that there is only one NP term)

$$F^{\bar{T}} = e^{K/2} K^{\bar{T}T} D_T W = e^{K/2} K^{\bar{T}T} (-aAe^{-aT} + K_T W) \tag{3.8}$$

Note that in this formula as well as in the arguments in the rest of this subsection the values of the moduli are understood to be taken at the local (negative CC) minimum of section II.

The requirement that the one loop CC contribution to the CC is cancelled then yields

$$3m_{3/2}^2 - K_{T\bar{T}} F^T F^{\bar{T}} = \frac{2\sqrt{3}a\Re A e^{-aT}}{k^{1/2}(S+\bar{S})^{1/2}(T+\bar{T})^{1/2}} m_{3/2} + O(e^{-2aT}) \sim h_{21} \frac{\Lambda^2}{16\pi^2} m_{3/2}^2, \tag{3.9}$$

where in the last relation we have used (3.4). This gives us an estimate of how large the non-perturbative contribution (at the minimum) should be:

$$Ae^{-aT} \sim a^{-1}(T+\bar{T})^{1/2} h_{21} \frac{\Lambda^2}{16\pi^2} m_{3/2}. \tag{3.10}$$

Let us check now that this gives a reasonable value for a. First we need to estimate the value of T. Using the fact that the Kaluza-Klein mass M_{KK} in Planck units is $1/T$,[9] we have

$$\frac{1}{T} \sim M_{KK} \sim \Lambda \sim 10^{-2} \Longrightarrow T \lesssim O(10^2) \tag{3.11}$$

[9]For a discussion of the scales involved in both the unwarped and warped cases see [46]. Note that we are actually discussing a class of type IIB solutions where warping can be ignored. It is completely unclear to us how to use the SUGRA formalism when warping is significant.

Assuming $A \sim O(1)$ and $m_{3/2} \sim 10TeV \sim 10^{-14}M_P$ from (3.10) we estimate $a \gtrsim O(1/10)$ which is a reasonable value since it would correspond to condensing gauge groups[10] of rank $N \sim 10-100$.

Let us ask how big the F-component of the complex structure moduli and the dilaton can be. In the presence of both imaginary anti-self-dual (IASD) fluxes (in the terminology of GKP [8]) and non-perturbative terms we have

$$F^{\bar{\alpha}} = K^{\bar{\alpha}\beta}e^{K/2}(D_\beta W_{\text{flux}} + K_\beta A e^{-aT}) = K^{\bar{\alpha}\beta}e^{K/2}(I_\beta + K_\beta A e^{-aT}) \tag{3.12}$$
$$F^S = K^{\bar{S}S}e^{K/2}(D_S W_{\text{flux}} + K_S A e^{-aT}) = K^{\bar{S}S}e^{K/2}(I + K_S A e^{-aT}) \tag{3.13}$$

Here I_β is an (2,1) flux and I is a (3.0) flux. Now the classical solution (in the absence of NP terms) requires that these IASD fluxes are zero. In finding a minimum for the classical potential that includes the non-perturbative terms such that the one-loop CC is cancelled, it is clear that we should not turn on IASD fluxes, since these would generically give large positive terms in the potential and violate the last two relations in (3.5). In that case using the estimate (3.10) we have

$$F^{\bar{\alpha}} \sim \frac{K^{\bar{\alpha}\beta}K_\beta}{k^{1/2}(S+\bar{S})2a\Re T}h_{21}\frac{\Lambda^2}{16\pi^2}m_{3/2}, \tag{3.14}$$
$$F^{\bar{S}} \sim \frac{K^{\bar{S}S}K_S}{k^{1/2}(S+\bar{S})2a\Re T}h_{21}\frac{\Lambda^2}{16\pi^2}m_{3/2}. \tag{3.15}$$

Finally we observe that for consistency these values of the F-terms of these moduli implies that their masses are considerably lower than the string scale. This can be seen by imposing the constraint that the mass of the scalar partner of the Goldstino should be of the order of $m_{3/2}$. Defining the unit vector in the Goldstino direction $u^m \equiv F^m/\sqrt{K_{m\bar{n}}F^m F^{\bar{n}}}$ we have

$$u^T = \frac{F^T}{\sqrt{F^T F_T(1+\epsilon^2)}} \sim \frac{e^{i\phi_T}}{\sqrt{K_{T\bar{T}}(1+\epsilon^2)}}$$
$$u^\alpha = \frac{F^\alpha}{\sqrt{F^T F_T(1+\epsilon^2)}} \sim \frac{\epsilon^\alpha e^{i\phi_\alpha}}{\sqrt{K_{T\bar{T}}(1+\epsilon^2)}}$$

where $\epsilon^\alpha \equiv |F^\alpha|/\sqrt{F^T F_T}$ and $\epsilon^2 = \epsilon^\alpha \epsilon_\alpha$ and we take $\alpha = 0, 1, \dots h_{12}$ with the index $\alpha = 0$ identified with the dilaton S. So for the squared mass of the sGoldstino we have

$$u^m V_{m\bar{n}} u^{\bar{n}} = \frac{K^{T\bar{T}}V_{T\bar{T}}}{1+\epsilon^2} + \frac{\epsilon^\alpha V_{\alpha\bar{\beta}}\epsilon^{\bar{\beta}}}{1+\epsilon^2} \sim O(m_{3/2}^2)$$

This tells us that $\epsilon^\alpha \sim m_{3/2}/m_\alpha$ so that we have the result[11] cd

$$F^\alpha \sim \frac{m_{3/2}}{m_\alpha}m_{3/2}.$$

Comparing with (3.14), (3.15) we see that this implies $m_S \sim m_z \sim 10^4 m_{3/2}$.

[10] In the KKLT picture this would come from the gauge theory on seven-branes wrapping a 4-cycle in the internal manifold.

[11] This has also been obtained in [47] though the argument there is different from the above and appears to depend on global SUSY.

3.3 μ and $B\mu$ terms

The expression for the effective μ term (after integrating out the moduli) is given by (see for example [16, 20] and references therein)

$$\mu_{IJ} = e^{K_{\rm mod}/2}\tilde{\mu}_{IJ} + m_{3/2}X_{IJ} - \bar{F}^{\bar{A}}\partial_{\bar{A}}X_{IJ} + O\left(\frac{\Lambda^2}{16\pi^2}m_{3/2}\right). \tag{3.16}$$

In this expression the second and third term are of the order of the supersymmetry breaking but there is no reason for first term (which comes from the original superpotential) to be of the same order - generically it would be $O(1)$ in Planck units. That of course would be a disaster since in that case there would be no electroweak symmetry breaking. This is the well known μ problem of the MSSM.

In our string theory based model of gravity mediated SUSY breaking with the MSSM located on D3 branes however $\tilde{\mu} = 0$ and the effective μ term emerges from the well-known Giudice-Masiero effect [21]. As shown by Graña et al [10] $\mu_{IJ} = -\bar{F}^{\bar{\alpha}}\partial_{\bar{\alpha}}X_{IJ}$ so using (3.14) and the fact that the sum over α has h_{21} terms, we get

$$\mu \sim O\left(\frac{h_{21}^2}{aT}\frac{\Lambda^2}{16\pi^2}m_{3/2}\right) \sim 10^{-2}m_{3/2}, \tag{3.17}$$

where we have used the value $aT \sim O(10)$ (see line after (3.11)) and $h_{21} \sim 3\times 10^2$.

Also using the calculation of the $B\mu$ term given in [10] we have

$$B\mu_{IJ} = V_{\rm classical}|_0 X_{IJ} \sim O\left(h_{21}\frac{\Lambda^2}{16\pi^2}m_{3/2}^2\right) \sim 10^{-3}m_{3/2}^2, \tag{3.18}$$

so that

$$\frac{B\mu}{\mu} \sim \frac{aT}{h_{21}}m_{3/2} \sim 3\times 10^{-2}m_{3/2}. \tag{3.19}$$

3.4 Gaugino mass

Let us now consider the gaugino masses. The general formula for these is

$$m_a = \frac{1}{2}(\Re f_a)^{-1}F^A\partial_A f_a(\Phi), \tag{3.20}$$

and we will only get a contribution if the gauge coupling function f_a depends on a chiral multiplet that acquires a non-vanishing F-term. In our case since the gauge theory on the D3 branes is independent of the moduli of the internal manifold and so the gaugino mass is suppressed relative to the gravitino mass. In particular the quadratic divergence in the potential led us to shift the minimum resulting in the new F-term values (3.4). Since the gauge coupling function on D3 branes depends (in the Einstein frame) on the dilaton this gives a non-vanishing contribution (since $f \sim S \sim 1/g^2$)

$$\frac{m_a}{g_a^2} = \frac{1}{2}F^S\partial_S f_a(S) \sim O\left(h_{21}\frac{\Lambda^2}{16\pi^2}m_{3/2}\right) \lesssim 10^{-3}m_{3/2}. \tag{3.21}$$

3.5 mSUGRA parameters

As discussed above we need to choose the cutoff $\Lambda \sim 10^{-2}$ and we took $h_{21} \sim 3 \times 10^2$. Taking $m_{3/2} \sim 10^4 GeV$ we get reasonable soft parameters except that the gaugino masses are too small. But as we shall see in the next section there is an 'AMSB' contribution to the gaugino masses which is much larger than the mSUGRA contribution. Our mSUGRA parameters are,

$$\mu \sim 10^{-2} m_{3/2} \sim 100GeV,\ m_s \sim 2\times 10^{-2} m_{3/2} \sim 200GeV,\ \frac{B\mu}{\mu} \sim \frac{aT}{h_{21}} m_{3/2} \sim 300GeV. \tag{3.22}$$

Note that a somewhat larger value of the gravitino mass (say $\sim 30TeV$ as in typical 'AMSB' scenarios) would also give acceptable values provided we take h_{21} to be a little larger for example $h_{21} \sim 4-5\times 10^2$. The gaugino masses however would still be of $< O(100GeV)$ and hence if this is the only contribution the model would be ruled out on phenomenological grounds. However as we discuss in the next section there are additional contributions.

4 SUSY breaking and AMSB

In the previous section we showed how mSUGRA like SUSY breaking terms arise at the cutoff scale Λ, in a model which can be naturally embedded in a type IIB string theoretic setup. These boundary conditions of mSUGRA need to be evolved down to the electro-weak scale in order to evaluate the actual predictions of this set up for the MSSM parameters. This calculation is just the same as in the usual mSUGRA set up and we will not go over it.

However there is a new contribution to any such theory that needs to be considered. This is usually assumed to be due to conformal anomalies and is referred to as anomaly mediated supersymmetry breaking (AMSB) [3–5, 22]. The most detailed SUGRA based derivation of the gaugino mass is given in the last citation and (in our conventions) reads

$$\frac{m_a}{g_a^2} = \Re[F^i \partial_i f_a(\Phi)| - \frac{1}{8\pi^2}(b'_a m_{3/2} + c_a F^i \partial_i K_m + 2T_R F^i \partial_i \ln Z_r)], \tag{4.1}$$

where $c_a = T(G_a) - \sum_r T_a(r)$ and $b'_a = 3T(G_a) - \sum_r T_a(r)$ with $T(G_a), T_a(r)$, being the traces of a squared gauge group generator in the adjoint and a matter representation r respectively. The sum over representations go over all states which are effectively massless at the cut off scale. In our approximately no-scale model with the MSSM on a stack of D3 branes the contribution of the first term was given in (3.21). We also see that there is a cancellation amongst the terms in the paranthesis in the second term in l.h.s. of (4.1) so that we effectively get from this formula the same result as before, namely

$$\frac{m_a}{g_a^2} \sim O(10^{-3} m_{3/2}). \tag{4.2}$$

If correct this would mean that the gaugino masses are well below the experimental upper limit, even for gravitino masses that are as high as $100TeV$, which is the highest one can tolerate without seriously affecting the hierarchy. This would imply that type IIB string theory with the MSSM on D3 branes can only give a split supersymmetry type of scenario. Thus we would need $m_{3/2} \sim 10^3 TeV$, giving gaugino masses $m_a \sim 1TeV$, but soft masses as well as μ and $B\mu/\mu$ would then be $O(10TeV)$ and the Higgs squared mass would be fine tuned (at a level of 1 part in 10^4). However as we will argue below this conclusion is not warranted.

The point is that as shown in [6] the arguments in [3–5, 22] need to be revised. Let us briefly summarize this discussion. The most important point is that the the so-called Weyl (or conformal) compensator chiral superfield C is a (non-propagating) field, and the theory needs to have enough gauge freedom so that it can for instance be set equal to unity to get the standard formulation of SUGRA. The Weyl anomaly at one-loop effectively prevents this, and Kaplunovsky and Louis (KL) [23] showed by a careful and detailed calculation how this anomaly could be cancelled thereby restoring the gauge symmetry. Their discussion led to a corrected form for the gauge coupling function in superspace (at the cutoff scale Λ)

$$H_a(\Phi;\Lambda) = f_a(\Phi) - \frac{3b'_a}{4\pi^2}\ln C - \frac{T_a(r)}{4\pi^2}\ln(e^{-\frac{1}{3}K_m}Z_r)|_{\text{holomorphic}}. \tag{4.3}$$

Projecting the F-term of this gives us the formula

$$\frac{m_a}{g_a^2} = \Re[F^i\partial_i f_a(\Phi)| - \frac{b'_a}{8\pi^2}\frac{F^C}{C} - \frac{T_a(r)}{4\pi^2}F^i\partial_i(\ln(e^{-\frac{1}{3}K_m}Z_r)], \tag{4.4}$$

where C, F^C are the lowest and highest components of the Weyl compensator superfield. The question is what is the value of the second term on the r.h.s. . As shown by KL, in the Kaehler-Einstein frame (which is the correct 'physical' frame in which standard SUGRA low energy results should be derived) $\frac{F^C}{C} = \frac{1}{3}K_iF^i$. Putting this in (4.4) we have

$$\begin{aligned}\frac{m_a}{g_a^2} = \Re H_a(\Phi;\Lambda)|_F &= \Re[F^i\partial_i f_a(\Phi)| - \frac{c_a}{8\pi^2}F^i\partial_i K_m - \frac{T_a(r)}{4\pi^2}F^i\partial_i(\ln Z_r)],\\ &= O(10^{-3}m_{3/2}) - \frac{b'_a}{8\pi^2}m_{3/2} + O\left(\frac{\Lambda}{4\pi}m_{3/2}\right).\end{aligned} \tag{4.5}$$

In the last line the first term is from (3.21), and we have used the values for the Kahler potential and F-terms for our model from (2.3), (2.4), (3.14), (3.15). This is the correct contribution from AMSB and indeed it gives a value for the gaugino masses (with $m_{3/2} \sim 30TeV$) that is of the right order.

However as shown in [7] (DS) and elaborated on in [6] there is an additional contribution which has nothing to do with Weyl anomalies but is a quantum effect in the effective action. This arises as follows. The gauge coupling function at the high (GUT) scale Λ is given by the superfield (4.3). At a low scale μ the one-loop beta function formula gives

$$H_a(\Phi;\mu) = H_a(\Phi;\Lambda) - \frac{b'_a}{8\pi^2}\ln\frac{\Lambda}{\mu}. \tag{4.6}$$

Now assume that there is an intermediate threshold which is generated by a superfield X (and its F-term) acquiring non-zero values in the ground state of the theory. The appropriate replacement of (4.6) in the Wilsonian effective action at the low scale μ is

$$H_a(\Phi, X;\mu) = H_a(\Phi;\Lambda) - \frac{b_a}{8\pi^2}\ln\frac{X}{\mu} - \frac{b'_a}{8\pi^2}\ln\frac{\Lambda}{X} \tag{4.7}$$

where b_a is the beta-function coefficient below the scale set by the vev X_0 of X. It may be obtained by integrating the one-loop beta function above and below the scale set by X and then using holomorphy. In effect this is the usual argument that the superspace gauge coupling function is one-loop exact. This is basically the argument given in DS [7] except that there the scale μ was not introduced and X_0 was eventually taken to zero. However that would clearly introduce infra-red divergences and in any case we should be in the Higgs phase where X has a non-zero expectation value and as DS argued their discussion really applies only in the Higgs phase. It should also be noted that this formula is exact for the Wilsonian coupling function. A similar formula is given in [24] in the context of gauge mediated supersymmetry breaking and is used in [5] to argue for what is often called deflected anomaly mediation in the literature.[12] Indeed as pointed out in both [7, 24] the X dependence in this formula follows from holomorphy and the necessity of reproducing the correct chiral anomaly from states that have been integrated out to get the effective theory at scales below that set by X.

Taking the F-term of this and replacing the field and its F-term by their ground state values, we have for the gaugino mass at the scale $\mu \to X_0$

$$\frac{m_a}{g_a^2} = \Re H_a(\Phi;\Lambda)|_F - \frac{b_a - b'_a}{8\pi^2}\frac{F^X}{X_0}. \tag{4.8}$$[13]

In our model the only possible threshold below the GUT scale is the weak scale set by the Higgs field itself. Thus we should take the gauge neutral field X to be the gauge neutral combination of the two MSSM Higgs superfields, i.e. we may put $X^2 = H_u H_d$. In effect this was what was done by Dine and Seiberg in their toy model outlining this general idea in [7] and the above mentioned chiral anomaly is in the global symmetry under which both H_u and H_d rotate by the same phase and the quark and lepton fields rotate by half the opposite phase - a symmetry which is explicitly broken by the μ term. X will acquire a non-zero value X_0 in the physical Higgs vacuum of the theory. The value which goes into (4.8) is thus

$$\frac{F^X}{X_0} = \frac{1}{2}\left(\frac{F^u}{v_u} + \frac{F^d}{v_d}\right), \tag{4.9}$$

where we have set the vevs of the charged components of the Higgs fields to zero and $v_u(v_d)$ are the vevs of the neutral Higgses $H_u^0(H_d^0)$. Now the F-terms may be computed from (2.5) and (2.2). The relevant term in the superpotential is $W \sim \mu H_u H_d$ and in the Kaehler

[12] Formulae for gaugino masses, given in the literature on GMSB, which have corrections which effectively involve D^2X etc (where D is the supercovariant spinor derivative) would necessarily take us to a higher derivative theory and in any case are not relevant for the Wilsonian coupling function that we are concerned with here. In other words to the extent that we confine ourselves to a two derivative Wilsonian action, the formula (4.7) is exact.

[13] We note that the calculation is similar to that given in [24] the main difference being that the effective messenger scale is the TeV scale and there is no messenger sector. It should also be noted that the GMSB constraint $F^X/X_0^2 < 1$ which applies from the necessity of ensuring non-negative messenger squared masses does not apply here. First of all there are no messengers but more importantly the mass formulae are given by the supergravity expression (1.3) and not by global SUSY formulae which are used in the GMSB context where the gravitino mass contribution maybe ignored.

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potential it is $K \sim Z(H_u\bar{H}_u + H_d\bar{H}_d)$. Then we have (with the indices for $H_u, H_d \to u, d$)

$$F^{\bar{u}} = e^{K/2}K^{\bar{u}u}(\partial_u W + K_u W) = e^{K/2}K^{\bar{u}u}\tilde{\mu}H_d + m_{3/2}\bar{H}_u \simeq m_{3/2}v_u, \tag{4.10}$$
$$F^{\bar{d}} = e^{K/2}K^{\bar{d}d}(\partial_d W + K_d W) = e^{K/2}K^{\bar{d}d}\tilde{\mu}H_u + m_{3/2}\bar{H}_d \simeq m_{3/2}v_u. \tag{4.11}$$

As is usual in the MSSM we have chosen the vevs to be real (this may in fact be done without loss of generality) and we have ignored the 'mu' term contribution since it is suppressed in our class of models. Thus we have

$$\frac{F^X}{X_0} = m_{3/2}. \tag{4.12}$$

Using this and (3.18) in (4.8) we get

$$\frac{m_a}{g_a^2}(\mu) = -\frac{b_a}{8\pi^2}m_{3/2}. \tag{4.13}$$

The above discussion is in fact the usual treatment of RG evolution that is used in the presence of thresholds. In our case this threshold is at the soft mass scale which is in fact the same as the Higgs scale v. Above this scale all superparticles would contribute to the evolution, while below one might expect only the standard model particles to contribute. This for instance is the assumption made in extrapolating from the standard model to the GUT scale to get unification, by accounting for the superpartners of the standard model particles which give a similar threshold effect. Note that these masses are of the same order of magnitude as the squark/slepton masses. For completeness we quote the values obtained for each separate gaugino

$$m_3 = -7\frac{\alpha_3}{4\pi}m_{3/2} \tag{4.14}$$
$$m_2 = -\frac{19}{6}\frac{\alpha_2}{4\pi}m_{3/2} \tag{4.15}$$
$$m_1 = \frac{41}{10}\frac{\alpha_1}{4\pi}m_{3/2} \tag{4.16}$$

In the usual discussion of AMSB it is claimed that there is a contribution from Weyl anomalies to the soft masses as well [3, 5]. This argument is based on the following reasoning. One starts with the *assertion* that the wave function renormalization should undergo the following replacement

$$Z\left(\Phi, \bar{\Phi}; \ln\frac{\Lambda}{\mu}\right) \to Z\left(\Phi, \bar{\Phi}; \ln\frac{\Lambda|C|}{\mu}\right) \tag{4.17}$$

in supergravity. Then one has for the squared soft masses

$$m^2 = -\ln Z|_{\theta^2\bar{\theta}^2} = -\frac{1}{4}|F^C|^2\frac{d^2\ln Z}{d\ln\Lambda^2}. \tag{4.18}$$

However not only is there no justification for the replacement (4.17), making Z dependent on C would in fact violate the Weyl gauge invariance of supergravity which is essential for the consistency of the formalism. In fact as discussed above, the addition of a $\ln C$ term to the gauge coupling function by KL [23] was designed to restore the Weyl invariance of the theory. The replacement (4.17) on the other hand would result in breaking the Weyl invariance making C a propagating field, and therefore it is incorrect. Formula (4.18) is therefore invalid.

Nevertheless there is a contribution to the soft mass that comes from a quantum effect that has nothing to do with Weyl anomalies. This was pointed out in [7] and the mechanism is a consequence of the formula (4.7). In the Higgs branch of the theory The radiatively generated soft mass at a scale $\mu \to X_0$

$$m_\Phi^2(X_0) = 2\sum_a c^a(r)\left(\frac{\alpha^a_{X^0}}{4\pi}\right)^2 (b^a - b'^a)\frac{|F_X|^2}{|X_0|^2}.^{14} \tag{4.19}$$

Here the sum is taken over the three gauge group factors and $\alpha^a = g^{a2}/4\pi$. Also $c^a(r)$ is the quadratic Casimir of the gauge group representation r of the observable field (squark or slepton). Then using (4.12) we have the contribution to the soft masses,

$$m_\Phi^2(X_0) = 2\sum_a c_\Phi^a\left(\frac{\alpha^a_{X^0}}{4\pi}\right)^2 (b^a - b'^a) m_{3/2}^2. \tag{4.20}$$

As we argued earlier, above the scale X_0 all superparticles would contribute to the evolution, while below one might expect only the standard model particles to contribute. This gives $b_3 - b_3' = 4$, $b_2 - b_2' = 25/6$, $b_1 - b_1' = 5/2$. We also have the values (with Q, L standing for the quark, lepton doublets respectively) $c_Q^{3.2.1} = \frac{4}{3}, \frac{3}{4}, \frac{1}{60}$; $c_L^{2,1} = \frac{3}{4}, \frac{3}{20}$; $c_{\tilde{u}}^{3,1} = \frac{4}{3}, \frac{4}{15}$ and $c_{\tilde{d}}^{3.1} = \frac{4}{3}, \frac{1}{15}$. The formula then gives the following generation independent contribution to the masses of the squark and sleptons.

$$m_Q^2 = \left[\frac{32}{3}\left(\frac{\alpha_3}{4\pi}\right)^2 + \frac{25}{4}\left(\frac{\alpha_2}{4\pi}\right)^2 + \frac{1}{12}\left(\frac{\alpha_1}{4\pi}\right)^2\right] m_{3/2}^2, \tag{4.21}$$

$$m_{\tilde{u}}^2 = \left[\frac{32}{3}\left(\frac{\alpha_3}{4\pi}\right)^2 + \frac{4}{3}\left(\frac{\alpha_1}{4\pi}\right)^2\right] m_{3/2}^2, \tag{4.22}$$

$$m_{\tilde{d}}^2 = \left[\frac{32}{3}\left(\frac{\alpha_3}{4\pi}\right)^2 + \frac{1}{3}\left(\frac{\alpha_1}{4\pi}\right)^2\right] m_{3/2}^2, \tag{4.23}$$

$$m_L^2 = \left[\frac{25}{4}\left(\frac{\alpha_2}{4\pi}\right)^2 + \frac{3}{4}\left(\frac{\alpha_1}{4\pi}\right)^2\right] m_{3/2}^2, \tag{4.24}$$

$$m_{\tilde{e}}^2 = 3\left(\frac{\alpha_1}{4\pi}\right)^2 m_{3/2}^2. \tag{4.25}$$

Let us now compare with the contribution from the quadratic divergence effects with $h_{21} \sim 3 - 5 \times 10^2$. For the squark masses the contribution from (3.22), is somewhat smaller than the values in (4.21)–(4.23) but it is an order of magnitude larger than the contribution (4.24), (4.25) to the slepton masses.

Note that the equations (4.14) to (4.16) and (4.21) to (4.25) give masses in the $O(10^3 GeV)$ to $O(10^2 GeV)$ range, provided we choose, as in the case of the usual presentation of AMSB phenomenology, a large mass for the gravitino ($\sim 30TeV$). Unlike

[14] Again this is very similar to the argument given in [24] for the corresponding GMSB calculation with similar caveats as in the gaugino case.

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that case however we do not require an additional mechanism to get non-negative squared slepton masses even though here we actually have such a contribution, namely (3.6), which as we observed above is much larger than this 'AMSB' contribution. Note that these give masses in the $O(10^3 GeV)$ to $O(10^2 GeV)$ range with the above choice for the gravitino mass.

5 Summary of phenomenology and conclusions

In this paper we have studied in detail the SUSY breaking phenomenology of a classically 'no-scale' like or 'sequestered' model (in the sense that the classical soft masses are suppressed relative to the gravitino mass) based on type IIB string theory with the MSSM coming from open string fluctuations on a stack of D3 branes. While there does not yet exist a complete chiral theory of this sort (with all moduli stabilized) it is plausible that once various technical difficulties are overcome such a model can indeed be constructed. If that turns out to be the case then its qualitative phenomenology would be that discussed in this paper. Actually it should be clear from the arguments that we have made, that this kind of phenomenology is quite generic for theories which are of this type (i.e. with suppressed classically generated soft terms). Thus we expect that similar phenomenological results emerge from a large class of string theoretic SUGRA models such as for instance IIB models with the MSSM on D7 branes.

These models have features that are similar to those found in all three standard mechanisms of SUSY mediation.

- Origin of SUSY breaking in moduli and transmitted by gravity as in mSUGRA.
- Soft parameters are due to quantum effects as in AMSB and GMSB.
- Gaugino masses mainly from 'AMSB'.
- $m_{3/2} \gtrsim 10TeV$ as in AMSB.

Such models have just two parameters that can be adjusted - the gravitino mass and the (integer) h_{21} (or in more general models the sum $h_{21}+h_{11}$). The cutoff is almost completely fixed once we demand that it should be higher than the scale at which the gauge couplings appear to unify, but below the string scale. Since at this point unification is the only concrete (albeit rather tenuous) evidence for supersymmetric physics, we strongly believe that it should be taken as an input. Since the cutoff should be definitely less than the string/Planck scale this limits us to the range $10^{16}GeV < \Lambda < 10^{18}GeV$. Furthermore as the string scale is expected to be somewhat below the Planck scale (perhaps $\leq 10^{17}GeV$) we are actually restricted to $\Lambda \sim 10^{16}GeV = 10^{-2}M_P$, which is the value that we have used. Also as we have seen in section 3 the effective perturbative parameter is $h_{21}\Lambda^2/16\pi^2M_P^2 \sim h_{21}10^{-6}$ with the above value of Λ. As we argued above, with a gravitino mass of around $30TeV$, the number of cycles in the Calabi-Yau manifold should not be much larger than 10^2 since the μ-term has to be well below $1TeV$ (see equation (3.18)) . On the other hand we cannot lower the gravitino mass below about $10TeV$ since in that case the gaugino masses (for the SU(2) × U(1) group) would be too low. Thus we see that this class of models must have

$$h_{21} \gtrsim 3 - 5 \times 10^2,\ m_{3/2} \sim 10 - 30TeV. \tag{5.1}$$

The universal scalar masses and the μ and $B\mu$ terms (at the unification scale Λ) are

$$\mu \sim \frac{B\mu}{\mu} \sim m_s \sim 100 - 500GeV, \tag{5.2}$$

and the gaugino masses (which are naturally computed at the MSSM scale via the 'AMSB' calculation of [6, 7] (see equations (4.14), (4.15), (4.16)) are

$$m_1 \sim 30 - 80GeV,\ m_2 \sim 40 - 100GeV,\ m_3 \sim 400 - 1000GeV.$$

Finally we should stress that so far there is no concrete string theoretic construction of the MSSM (living on D3 branes in type IIB or in any other string theory set up) with all the moduli stabilized. In this paper we have assumed that the SUSY breaking phenomenology of non-chiral constructions that has been discussed in the literature [10, 25] will hold for chiral models as well. Of course as we discussed before it is possible that these models will have FCNC terms that are proportional to h_{21}, like the flavor conserving terms, and in this case one would need a fine tuning of one part in 10^3 to suppress them.

Even with such additional fine-tuning it seems that this class of models is still the minimal possible and least fine-tuned one that can be embedded in string theory. Suppose for example there is a mSUGRA model with all moduli stabilized which does not have FCNC at the classical level. One would still have FCNC terms, but since the classical contribution to the scalar squared mass is now $O(m_{3/2}^2)$ the quantum contribution will be down by a factor $h_{21}\frac{\Lambda^2}{16\pi^2} \sim 10^{-4}$ and can be ignored. However now the gravitino mass is low $\sim 100GeV$ so there is a fine-tuning factor (according to the work of Douglas and collaborators [26]) $O(\frac{m_{3/2}(low)}{m_{3/2}(high)})^6 = O(10^{-12})$ when the high value is taken to be $\sim 10TeV$ as in the model discussed here. As for GMSB models, one might expect that the same factor applies, however it has been argued that such models are on a different branch [27]. Be that as it may, GMSB requires an additional sector - the so-called messenger sector - compared to the class of models discussed here.

It is clearly important to find detailed constructions that incorporate the MSSM within a string theory context where all moduli can be stabilized.[15] Even though one may not be able to make precise predictions (since by changing the fluxes one can make changes to the masses and couplings) the physics that we have discussed above would then be a qualitative prediction of string theory for LHC physics. This is so in the sense that the class of models that we have here are the minimal possible in terms of fine-tuning, and having just the sectors (namely a visible MSSM or GUT sector and a moduli sector) that necessarily have to be present.

[15]For recent progress on such constructions see [48–50].

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A On the relation to mirage mediation and large volume scenarios

The phenomenological consequences of the class of models discussed in this paper have a certain superficial resemblence to the so-called mirage mediation models of [28] (see also [29] for a recent variant of this) where the classical (plus nonperturbative) contribution is of the same order as the 'AMSB' one. However these models rely on the KKLT toy model with an uplift term. There are several problems with taking this model seriously for phenomenological purposes. Firstly (as pointed out in [30]) the logic of deriving a four dimensional theory from a ten dimensional theory requires that one starts from a classical supersymmetric vacuum of the ten-dimensional theory which enables one to organize the fluctuations around that point in 4D supergravity multiplets, and will then necessarily give a 4D $\mathcal{N} = 1$ SUGRA theory. If one starts with supersymmetry broken at the string level (even if the scale is suppressed by warping) there is simply no way of deriving a four dimensional supergravity. In fact the potential one gets is a runaway one for the Kaehler moduli - implying decompactification. The potential that is usually used is based on the assumption that one can add a non-perturbative term to the superpotential before one adds the Dbar term - but this is an inversion of the logic since the string theory starting point did not have such a non-perturbative term to begin with (especially if they arise from low energy gauge theoretic effects) while the Dbar brane is added at the string theory level i.e. in the ten dimensional theory.

If we ignore this, the uplift term is an explicit breaking of the $\mathcal{N} = 1$ four-dimensional supersymmetry (although in 10 dimensions it is a spontaneaous breaking caused by Dbar branes) at an intermediate scale $\sqrt{m_{3/2}M_P} \sim 10^{11}GeV$. The Dbar brane is located far down a throat so that its effective tension and hence the supersymmetry breaking, is warped down from the string/Planck scale to the above scale by the warp factor $e^{A_{\min}} \sim \sqrt{m_{3/2}/M_P}$ (see [28] equation (19)). If the MSSM/GUT branes are in the bulk (as is effectively the case in [28]) as opposed to being in the infra-red end of a throat region, this appears to give quantum effects that result in terms $O(m_{3/2}M_P\Lambda^2/16\pi^2)$ (rather than $O(m_{3/2}^2\Lambda^2/16\pi^2)$ as in this paper) in the potential. This would seem to introduce large corrections to soft masses etc when $\Lambda \sim M_{\rm GUT}$. However the overlap of the wave function of the MSSM states which are located at the UV end of the throat with the SUSY breaking fields at the IR end of the throat is exponentially suppressed. This results in an effective mass splitting at the UV end $\Delta m^2 \sim e^{2A_{\min}}m_{3/2}M_P = m_{3/2}^2$. and hence the quadratically divergent quantum contribution is again of the same order as in this paper. Another possibility would be to have the MSSM brane far down a throat region so that the effective cut off Λ is also warped down to an intermediate scale $\Lambda_{\rm eff}^2 = \frac{m_{3/2}}{M_P}\Lambda^2$, so that the estimate of the quantum contribution to the potential is again $O(m_{3/2}M_P\Lambda_{\rm eff}^2/16\pi^2) = O(m_{3/2}^2\Lambda^2/16\pi^2)$. In this case also the corrections to soft masses etc. will not significantly change the classical contributions. However in such a situation it is not entirely clear how to obtain a GUT theory since the effective cut off is far below the GUT scale though in the corresponding 5-dimensional scenario a possible resolution has been offered in [31]. Also in the string theoretic case the derivation of the effective four dimensional supergravity in the presence of significant warping has not yet been entirely resolved (for the problems associated with this and progress towards a resolution of this question see [32, 33]). It should also be mentioned that the mirage mediation scenario is not necessarily tied to having an anti-brane at the IR end of the throat. This could in principle be replaced by a conventional SUSY breaking sector at the IR end of a warped throat (for some discussion of this see [34]).[16]

One might of course avoid large corrections by taking the cutoff Λ to be much smaller than the GUT scale and this appears to be the case in the large volume scenario (LVS) discussed in [15]. In that and in subsequent work based on it, it is shown that in the absence of fine tuning a large volume scenario emerges (from GKP-KKLT type constructions) where the string scale is an intermediate scale (around $10^{12}GeV$). In this case there is a broken SUSY minimum with negative CC (exactly as required for the classical construction of this paper) and the authors use the uplift term of KKLT to lift the minimum to a small positive value. However the phenomenology appears to be insensitive to the uplift term (see for example [35]). But as we've argued in this paper, at this point in time, supersymmetric grand unification is the main piece of evidence for supersymmetry and should be taken seriously in model building. This can be achieved in the LVS scenario (if one fine-tunes the flux generated superpotential as we've done in this work) but then the quadratically divergent corrections that we have discussed in this paper will become relevant. A detailed discussion of this as well as an extention of the phenomenology discussed in this paper to the case when the standard model is located on D7 branes, will appear in future work.

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