Classical and quantum SUSY breaking effects in IIB local models

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ABSTRACT: We discuss the calculation of soft supersymmery breaking terms in type IIB string theoretic models in the Large Volume Scenario (LVS). The suppression of FCNC gives a lower bound on the size of the compactification volume. This leads to soft terms which are strongly suppressed relative to the gravitino mass so that the dominant contribution to the gaugino massses comes from the Weyl anomaly. The other soft terms are essentially generated by the renormalization group running from the string scale to the TeV scale.

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1 Introduction

Theories of supersymmetry (SUSY) breaking and transmission from a hidden to a visible sector has been the subject of much discussion over the last two to three decades. Much of this discussion has had little to do with string theory and often it has been conducted purely within a global SUSY framework. However a theory of supersymmetry breaking must necessarily be embedded within a ($\mathcal{N} = 1$) supergravity (SUGRA) which is derived from string theory. The following is a summary of the arguments leading to this assertion.

1. Adding a set of explicit soft SUSY breaking terms to a global theory (like the Minimally Supersymmetric Standard Model (MSSM)) leads to far too much arbitrariness - it does not give us a theory.

2. Spontaneous SUSY breaking in Global SUSY leads to a cosmological constant (CC) at the SUSY breaking scale which cannot be fine-tuned to zero.

3. A theory of SUSY breaking is therefore necessarily a SUGRA with a scalar potential which has a minimum that breaks SUSY spontaneously.

4. A SUGRA needs to be embedded in string theory in order to have a quantum mechanically consistent and complete theory.

The main problem in relating string theory to phenomenology is that the starting point of the theory is in ten dimensions. While there are only five weakly coupled string theories (which are related to each other through various dualities) the number of four dimensional 'compactifications' is extremely large. At the time of the second string revolution of the mid eighties, it was hoped that in spite of the existence of many compactifications, the number of theories with stabilized moduli (the fields governing the size and shape of the compact manifold) is small if not just one. However it was realized through the work of many authors (for a review see [1, 2]) culminating in that of [3] (GKP) and [4] (KKLT) that the number of such four dimensional models is extremely large. Thus at the current stage any discussion of the phenomenological consequences of string theory must proceed by first imposing a set of experimental inputs (in addition to requiring a compactification to four dimensional $\mathcal{N} = 1$ supergravity). These are:

- CC is tiny $\sim O((10^{-3}eV)^4)$

- No light scalars with gravitational strength coupling

- SUSY partner masses $\gtrsim O(100GeV)$

- Lightest Higgs $> 114GeV$

- Flavor changing neutral currents (FCNC) suppressed

- No large CP violating phases

The first of these is achieved by ensuring that there is a sufficiently large number of flux configurations such that there would be many solutions that realize this value. For typical Calabi-Yau compactifications of IIB string theory this certainly is the case. The second is achieved by a combination of fluxes and non-perturbative (NP) effects. The remaining four constraints are dependent on the particular mechanism of supersymmetry breaking and transmission.

In this paper we will discuss a theory of supersymmetry breaking that emerges from the so-called Large Volume Scenario (LVS) of type IIB string compactifications [5]. In addition to fluxes which stabilize the dilaton and complex structure moduli, as in the original KKLT model [4] non-perturbative (NP) effects may be used to stabilize the Kaehler moduli. Although there appears to have been some controversy about the latter in the literature the issue seems to have been settled - at least when the cycles in question are not wrapped by branes that support a gauge theory with chiral fermions. (For a recent comprehensive discussion and for references to earlier work see [6]).

However for a four-cycle which is wrapped by D7 branes carrying a chiral gauge theory (such as the MSSM) the situation appears to be different. In this case it has been argued in [7] that the chirality (of the MSSM) precludes the stabilization of the four cycle which they wrap by NP effects. The argument depends on the observation that in a D-brane construction of chiral theory, there would be an anomalous U(1) gauge group, which in effect requires the presence of charged matter field factors in the NP superpotential contribution that depends on the relevant four cycle volume. It has been argued in [7] (see also [8, 9])

that such matter fields must have zero vacuum values, so that effectively this contribution would be absent. This means also that this cycle would shrink below the string scale (this follows from examining the associated D-term potential). Effectively the situation becomes similar to that of having a D3 brane at a singularity (for more discussion see below). It is not clear to the author that this argument has been rigorously established, however it appears that requiring a reasonable phenomenology (in particular that MSSM fields should not acquire vaccum values at the scales at which the moduli are stabilized) seems to justify such a scenario. In any case we will take the attitude that such an outcome yields an interesting set up whose phenomenology is worth investigating.

After discussing the basic physical inputs and reviewing the pertinent results of [8] and [9], we go on to discuss the classical soft terms that arise from this class of theories (the details of the calculations are given in appendix A). In particular we will find that the classical soft mass squared is positive definite. This and the classical CC are both highly suppressed by a power of the (large) volume. Furthermore we show that the suppression of FCNC effects imply that there is a lower bound on the volume. If we ignore Weyl anomaly effects, then comparison of the classical FCNC effects with the flavor diagonal classical masses leads to a large $\mathcal{V} \gtrsim 10^{12}$ value.

However Weyl anomaly effects (usually called AMSB) changes the phenomenology of this class of theories. As shown in appendix B the Weyl anomaly gives an additional set of terms (calculated by Kaplunovsky and Louis (KL) [10]) to the gauge coupling function, leading to a contribution to the gaugino mass that is much larger than the classical one. This in turn drives the scalar masses by the mechanism of gaugino mediation [11, 12]. However the lower bound on the CYO volume implies a tension between having TeV scale soft masses (to address the hierarchy problem) and making the sGoldstino heavy enough to avoid the cosmological modulus problem.

2 Generalities

We follow the notation and discussion of [8] and [9]. We also set $M_P \equiv (8\pi G_N)^{-1/2} = 2.4 \times 10^{18} GeV = 1$.

The superpotential, Kaehler potential and gauge kinetic function for the theory under discussion are,

$$W = \hat{W}(\Phi) + \mu(\Phi)H_1 H_2 + \frac{1}{6}Y_{\alpha\beta\gamma}(\Phi)C^\alpha C^\beta C^\gamma + \ldots, \qquad (2.1)$$

$$K = \hat{K}(\Phi, \bar{\Phi}) + \tilde{K}_{\alpha\bar{\beta}}(\Phi, \bar{\Phi})C^\alpha C^{\bar{\beta}} + [Z(\Phi, \bar{\Phi})H_1 H_2 + h.c.] + \ldots \qquad (2.2)$$

$$f_a = f_a(\Phi). \qquad (2.3)$$

Here $\Phi = \{\Phi^A\}$ and C^α are chiral superfields (including the two Higgs doublets $H_{1,2}$) that correspond to the moduli and MSSM/GUT fields respectively. Also

$$\hat{K} = -2\ln\left(\mathcal{V} + \frac{\xi}{2}\left(\frac{(S+\bar{S})}{2}\right)\right) - \ln\left(i\int \Omega \wedge \bar{\Omega}(U,\bar{U})\right) - \ln(S+\bar{S}), \qquad (2.4)$$

$$\hat{W} = \int G_3 \wedge \Omega + \sum_i A_i e^{-a_i T^i}. \qquad (2.5)$$

Here $\mathcal{V}$ is the volume (in Einstein frame) of the internal manifold and the $\xi = -(\chi\zeta(3)/2(2\pi)^3)$ term is a correction term that is higher order in the α' expansion. For typical Calabi-Yau manifolds $\xi \sim O(1)$. S is the axio-dilaton, $U = \{U^m\}$ represents the set of $(m = 1,\ldots,h_{21})$ complex structure moduli and T^i $(i = 1,\ldots,h_{11})$ are the (complexified) Kaehler moduli. The type of Calabi-Yau manifolds that we consider are of the 'Swiss cheese' type. In the simplest such manifold consistent with our requirements the volume may be written as[1]

$$\mathcal{V} = \tau_b^{3/2} - \tau_s^{3/2} - \tau_a^{3/2}. \tag{2.6}$$

In the above the tau's are Kaehler moduli which control the volume of the four cycles with τ_b effectively determining the overall size of the CY. While in explicit calculations in the the rest of the paper, we will use (2.6) for the sake of simplicity it should be clear from the discussion that the results would hold even in a more general CY manifold of this type.

It turns out that in order to realize the Large Volume Scenario (LVS) with the MSSM located either on a $D3$ brane at a singularity or a seven-brane wrapping a four cycle on a CY orientifold, the total number of 2/4 cycles $h_{11} \geq 3$. At least one of these cycles (τ_s in the above) must either be wrapped by a Euclidean D3 instanton or by a stack of seven branes with a condensing gauge group that will generate non-perturbative effects. As pointed out in [7] the MSSM cannot be located on this cycle - hence the need for (at least) one more cycle (τ_a). It has been argued in [8, 9] that this cycle shrinks to zero unless stabilized at the string scale by non-perturbative string effects. In any case the F/D term associated with the corresponding Kaehler modulus is zero. Let us briefly review this argument.

The potential for the moduli is (assuming that the minimum would be at large $\mathcal{V}$ and expanding in it)

$$V = V_F + V_D. \tag{2.7}$$

$$V_F = \frac{4}{3} g_s (a|A|)^2 \frac{\sqrt{\tau_s} e^{-2a\tau_s}}{\mathcal{V}} - 2g_s a|AW_0| \frac{\tau_s e^{-a\tau_s}}{\mathcal{V}} + \frac{3}{8} \frac{\xi|W_0|^2}{g_s^{1/2}\mathcal{V}^3} + \ldots, \tag{2.8}$$

$$V_D = \frac{f}{2} D^2, \ D = f^{-1} k^i K_i. \tag{2.9}$$

In the above we've used the absence of a non-perturbative superpotential for the modulus of the MSSM cycle and the fact that the effective axionic partner of τ_s is stabilized at an odd multiple of π (giving the sign flip of the second term of V_F). The phases of A and W_0 (the flux superpotential) can then be set to zero without loss of generality [5]. The D-term comes from the anomalous U(1) (with Killing vector field k) living on the MSSM cycle under which the standard model fields and the modulus τ_a are charged and we have set the matter fields which are charged under this U(1) to zero following the arguments of [7] [8, 9]. The U(1) gauge coupling function is linear in τ_a and S. Also one has $K_a \sim \tau_a^\alpha/\mathcal{V}$, $\alpha > 0$,

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[1]In order to simplify the notation we have rescaled the moduli by replacing those used in [9] by $\tau_i \to \eta_i^{-1}\tau_i$. Correspondingly we have also replaced $a_i \to \eta_i a_i$.

and $D \propto 1/\mathcal{V}$. The F-term potential is minimized at

$$e^{-a\tau_s} \simeq \frac{3}{4}\frac{W_0}{aA\mathcal{V}}\sqrt{\tau_s}\left(1 - \frac{3}{4a\tau_s}\right), \tag{2.10}$$

$$\tau_s^{3/2} \simeq \frac{\hat{\xi}}{2}(1 + \frac{1}{2a\tau_s}), \tag{2.11}$$

where we've written $\hat{\xi} = (\frac{S+\bar{S}}{2})^{3/2}\xi$. Note that extremizing with respect to τ_s gives us an exponentially large volume and the three displayed terms in V_F are all of order $\mathcal{V}^{-3}$. This would mean that that at the classical (negative) minimum found in [5], the contribution to the F-term potential from the dilaton and complex-structure moduli[2] are zero. Also $V_D = 0$ since it is positive definite and of order $1/\mathcal{V}^2$. This would mean that, at least based on classical considerations, $\tau_a \to 0$, so that the standard model cycle shrinks to zero or at least shrinks below the string scale. Of course one might need to include all α' corrections in such a case, but even so the important point here is that both the D term and hence also the F-term of the MSSM cycle modulus τ_a become negligible at this classical minimum [9]. We also note for future use that (2.10) implies that

$$a\tau_s = |\ln m_{3/2}| + O(1), \tag{2.12}$$

and for $m_{3/2} \sim 10 - 100 TeV$ this is a number of $O(10)$.

The minimum found in [5] is at a negative value of the (classical) cosmological constant

$$V_0 = -\frac{3\hat{\xi}}{16a\tau_s}\frac{m_{3/2}^2}{\mathcal{V}}. \tag{2.13}$$

Note that here and in the rest of the paper we will be using the formula for the gravitino mass

$$m_{3/2}^2 = e^K|W|^2 \sim \frac{|W|^2}{\mathcal{V}^2}. \tag{2.14}$$

3 Classical soft terms

Let us first compute the classical soft mass term in this model. The general expression for the squared soft mass in SUGRA is [13][14]

$$m_{\alpha\bar{\beta}}^2 = V_{\text{class}}|_0\tilde{K}_{\alpha\bar{\beta}} + m_{3/2}^2\tilde{K}_{\alpha\bar{\beta}} - F^A F^{\bar{B}}R_{A\bar{B}\alpha\bar{\beta}} \tag{3.1}$$

$$= V_{\text{class}}|_0\tilde{K}_{\alpha\bar{\beta}} + m_{3/2}^2\tilde{K}_{\alpha\bar{\beta}} - F^b F^{\bar{b}}R_{b\bar{b}\alpha\bar{\beta}}$$

$$-2ReF^b F^{\bar{s}}R_{b\bar{s}\alpha\bar{\beta}} - F^s F^{\bar{s}}R_{s\bar{s}\alpha\bar{\beta}}. \tag{3.2}$$

In the second equality we've used the fact that in the classical vacuum before uplifting the only source of SUSY breaking are the F-terms of T^b and T^s. Also we will set $\tau_a \to 0$ (and hence also set the corresponding F-term to zero) following the arguments of [8, 9]. To

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[2]At this point we ignore uplifting issues.

calculate (3.2) explicitly we need these F-terms as well as the matter metric $K_{\alpha\bar{\beta}}$. We find (see appendix A for details)

$$F^b = -\tau^b \left(2 + \frac{3}{8} \frac{\hat{\xi}}{a\tau^s} \frac{1}{\mathcal{V}} + O\left(\frac{1}{(a\tau^s)^2 \mathcal{V}} \right) \right) m_{3/2}, \tag{3.3}$$

$$F^s = -\frac{3}{2} \frac{\tau^s}{a\tau^s} m_{3/2} (1 + O(\mathcal{V}^{-1})). \tag{3.4}$$

For the MSSM on D3 branes the matter metric can be calculated (see appendix A) from the formulae for the Kaehler coordinates given in [15] (assuming that the formulae given in that reference remain valid for D3 branes at a singularity). Since we expect D7 branes on a collapsed cycle to act like D3 branes, these formulae should be valid in that case too. We have

$$K_{\alpha\bar{\beta}} = \frac{c}{\mathcal{V} + \hat{\xi}/2} (\sqrt{\tau^b} \omega^b_{\alpha\bar{\beta}} - \sqrt{\tau^s} \omega^s_{\alpha\bar{\beta}}), \tag{3.5}$$

where $\omega^{b(s)}$ is the harmonic $(1,1)$ form associated with the big(small) modulus evaluated at the position of the D3 brane or the collapsed cycle wrapped by the D7 brane, and c is an $O(1)$ constant.

Let us pause here for a moment to discuss the validity of (3.5), beyond the context of smooth CY orientifolds within which it was derived (see appendix A) following the formula for the embedding of D3 branes given in [15]. The basic argument in appendix A depended essentially on the field redefinition that is necessary to obtain the correct holomorphic coordinates on moduli space, given in the above reference. This is schematically of the form (see equation (3.13) of [15])

$$T^i + T^{\bar{i}} = 2\tau^i + 2\mu l^2 i \omega^i_{\alpha\bar{\beta}} C^\alpha C^{\bar{\beta}} + (CUC + \bar{C}\bar{U}\bar{C}) \text{term} + \dots \tag{3.6}$$

This formula was obtained by comparing the effective action of IIB string theory with D3 branes located at some (generic smooth) point on the CY orientifold with the standard supergravity formula for the kinetic terms of the effective action which are written in terms of a Kaehler potential. Now unfortunately a similar derivation has not been given when the branes are located at a singularity. In so far as one still expects a supergravity description of the low energy physics, the question is to what extent a formula such as (3.6) remains valid. For us the essential feature of this formula is the dependence on the matter fields to quadratic order. In particular for the above calculation of the dependence of the metric on the Kaehler moduli, what is relevant is the second term. So our assumption here is that this structure, i.e. the proportionality to the 1,1 form ω^i (of the cycles which are stabilized) is preserved beyond the original calculation at a generic smooth point. It is hard to imagine that (apart from the possible modification of the coefficient) that this structure can be qualitatively modified when the D3 brane location is at some singularity. Indeed what we are using for our calculations are just very generic features of the formula for the metric (3.5). The numerical value of the normalized diagonal mass term (see equation (3.10) will certainly not change since it depends only on the part of the curvture on moduli space that is proportional to the matter metric $K_{\alpha\bar{\beta}}$ and is independent on

details of the dependence on the two ω's . Furthermore the relative numerical coefficients in (3.5) (provided they are not changed by more than $O(1)$ numbers) is not going to affect the qualitative features pertaining to the FCNC issue discussed below either. In fact what is relevant for the latter is the different dependence on the two moduli that is a feature of the two terms in the metric, and this is unlikely to change for the metric extracted from D3 branes at a singularity or D7 branes wrapping a collapsing cycle.

The dependence on ω^s tends to give FCNC effects and we'll postpone that discussion to the next section. Dropping that term the metric is of the form $K_{\alpha\bar\beta} = f(\tau^b, \tau^s)\omega^b_{\alpha\bar\beta}$ and the Riemann tensor can be calculated from the formula $R_{i\bar j\alpha\bar\beta} = (\partial_i\partial_{\bar j}\ln f)K_{\alpha\bar\beta}$. This gives (after using (2.11) and neglecting $O(1/(a\tau^s)^2\mathcal{V})$ terms[3])

$$R_{b\bar b\alpha\bar\beta} = \frac{1}{4(\tau^b)^2}(1 + \frac{15}{16}\frac{\hat\xi}{a\tau^s\mathcal{V}})K_{\alpha\bar\beta}, \tag{3.7}$$

$$R_{b\bar s\alpha\bar\beta} = -\frac{9}{16}\frac{(\tau^s)^{1/2}}{(\tau^b)^{5/2}}K_{\alpha\bar\beta}, \tag{3.8}$$

$$R_{s\bar s\alpha\bar\beta} = \frac{3}{16}\frac{(\tau^s)^{-1/2}}{(\tau^b)^{3/2}}K_{\alpha\bar\beta}. \tag{3.9}$$

Using the above and equations (3.3)(3.4) in (3.2) we find[4] to leading order in $1/a\tau_s\mathcal{V}$ (see appendix A for details),

$$m^2_{\alpha\bar\beta} = V_0 K_{\alpha\bar\beta} + \frac{3}{8}\frac{\hat\xi}{a\tau^s}\frac{m^2_{3/2}}{\mathcal{V}}K_{\alpha\bar\beta} = +\frac{3}{16}\frac{\hat\xi}{a\tau^s}\frac{m^2_{3/2}}{\mathcal{V}}K_{\alpha\bar\beta}. \tag{3.10}$$

Note that the second term differs in sign from the result quoted in [9]. The reason is that we have here included the contribution of the F^s terms (and also the sub leading corrections to $R_{b\bar b\alpha\bar\beta}$). Thus the classical squared masses (even after including the contribution of the negative CC) is actually positive.

3.1 Uplift issues

So far we have worked with the LVS minimum, which breaks supersymmetry, but has a negative cosmological constant. It was argued in [5] that at this minimum one expects the positive definite contributions to the potential coming from the dilaton and the complex-structure moduli to be actually zero since they scaled as $\mathcal{V}^{-2}$ whereas the negative minimum scaled as $\mathcal{V}^{-3}$. However this minimum needs to be uplifted and one way that could happen is if the dilaton (S) and the complex structure moduli (U^m) acquired F-terms. Generically

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[3]Typically $\hat\xi$ is a number of $O(1)$ and $a\tau^s$ is a number around 30 so these corrections can be safely ignored. For instance even if $\hat\xi \sim 5$ the ratio is expected to be of around 15-20% and that is the order of corrections that we would expect to these formulae. This is certainly does not affect the qualitative features of the calculations of the classical soft masses.

[4]To leading order in the matter fields $K_{\alpha\bar\beta}$ is the same as $\tilde K_{\alpha\bar\beta}$ so we will not distinguish between them in the rest of the paper.

one would expect at the new uplifted minimum (given that $V_0 \sim -m_{3/2}^2/(\mathcal{V} \ln m_{3/2})$)

$$F^S \bar{F}^{\bar{S}} K_{S\bar{S}} \lesssim \frac{m_{3/2}^2}{\ln m_{3/2} \mathcal{V}}, \tag{3.11}$$

$$F^m \bar{F}^{\bar{n}} K_{m\bar{n}} \lesssim \frac{m_{3/2}^2}{\ln m_{3/2} \mathcal{V}}. \tag{3.12}$$

Since much of the rest of the discussion focuses on the large modulus we will often replace $T^b \to T$ and $\tau^b \to \tau$. We have used above the argument that the seven brane when wrapping a shrinking 4-cycle should behave like a D3 brane. For the same reason we expect the leading classical contribution to the gauge coupling function for both the D3 and D7 cases to be

$$f_a = S + \kappa T^a, \tag{3.13}$$

(with $\kappa = 0$ in the D3 case). Since as argued in [8, 9] the F-term of the modulus T^a corresponding to the shrinking cycle is vanishingly small, we see that without the uplift contribution from the dilaton, there would be no (classical) contribution to the gaugino mass in both the D3 and the D7 brane cases.

After uplift (assuming that F^S, F^U take generic values consistent with (3.11)(3.12)) we have the following expressions for the gaugino mass M, the scalar mass m, the A term, the effective μ term[5] and B terms (see for example [13] for definitions and general formulae and appendix A):

$$M_a = \frac{F^i \partial_i f_a}{2 f_a} = \frac{F^S}{2S} \lesssim O\left(\frac{m_{3/2}}{\sqrt{\ln m_{3/2} \mathcal{V}}}\right), \tag{3.14}$$

$$m_{\alpha\bar{\beta}}^2 = (m_{3/2}^2 K_{\alpha\bar{\beta}} - F^i \bar{F}^{\bar{j}} R_{i\bar{j}\alpha\bar{\beta}}) = \left(O\left(\frac{m_{3/2}^2}{\ln m_{3/2} \mathcal{V}}\right)\right) K_{\alpha\bar{\beta}} + \ldots, \tag{3.15}$$

$$A_{\alpha\beta\gamma} = e^{K/2} F^i D_i y_{\alpha\beta\gamma} \lesssim O\left(\frac{m_{3/2}}{\mathcal{V}}\right) y_{\alpha\beta\gamma}, \tag{3.16}$$

$$\mu \sim B\mu/\mu \lesssim O\left(\sqrt{h_{21}} \frac{m_{3/2}}{\sqrt{\ln m_{3/2} \mathcal{V}}}\right). \tag{3.17}$$

Note that in the last equation h_{21} is the number of complex structure moduli.

3.2 Classical FCNC effects

In computing the soft mass using (3.5) we ignored the second term inside the parenthesis. Let us now compute its contribution (for details see appendix A). To do so we must evaluate the Riemann tensor. The leading contribution comes from the sectional curvature

$$\begin{aligned} R_{T\bar{T}\alpha\bar{\beta}} &= \partial_T \partial_{\bar{T}} K_{\alpha\bar{\beta}} - K^{\gamma\bar{\delta}} \partial_T K_{\alpha\bar{\delta}} \partial_{\bar{T}} K_{\gamma\bar{\beta}} + O(C) \\ &= \frac{1}{3} K_{T\bar{T}} c [\frac{\omega_b}{\tau_b} - \frac{7}{4} \frac{\omega_s}{\tau_b} \sqrt{\frac{\tau_s}{\tau_b}}]_{\alpha\bar{\beta}} + O(C). \end{aligned} \tag{3.18}$$

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[5]Assuming that the supersymmetric one is zero as is the case for the D3 branes [15].

The problem is that this is not proportional to $K_{\alpha\bar{\beta}}$ - see (3.5). The ω^s dependence in (3.18) gives (from the third term on the r.h.s. of (3.2)) an additional term to the expression in (3.10), i.e.

$$m^2_{\alpha\bar{\beta}} = m^2_{3/2}K_{\alpha\bar{\beta}} - F^T F^{\bar{T}} R_{T\bar{T}\alpha\bar{\beta}} + \ldots$$

$$= \frac{3}{16}\hat{\xi}\frac{m^2_{3/2}}{\ln m_{3/2}\mathcal{V}}K_{\alpha\bar{\beta}} + m^2_{3/2}\frac{3}{4}\sqrt{\frac{\tau_s}{\tau_b}}K'_{\alpha\bar{\beta}}, \tag{3.19}$$

$$\equiv (m^2\delta^\gamma_\alpha + \Delta m^{2\gamma}_\alpha)K_{\gamma\bar{\beta}}. \tag{3.20}$$

Here $K'_{\alpha\bar{\beta}} \equiv c\omega^s_{\alpha\bar{\beta}}/\tau_b$ and is not proportional to $K_{\alpha\bar{\beta}}$ so $\Delta m^{2\gamma}_\alpha = m^2_{3/2}\frac{3}{4}\sqrt{\frac{\tau_s}{\tau_b}}(K'K^{-1})^\gamma_\alpha$ is not a diagonal matrix. Also $m^2 = \frac{3}{16}\hat{\xi}\frac{m^2_{3/2}}{\ln m_{3/2}\mathcal{V}}$. If the two harmonic one-one forms ω^b and ω^s evaluated at the position of the D3 branes (or the collapsed cycle wrapped by the D7 branes) are of the same order of magnitude then the dominant contribution is a flavor violating one.

Clearly this would be a phenomenological disaster since at least for the first two generations the flavor violating non-diagonal contributions to the squared soft masses should be suppressed relative to the flavor conserving ones. The relevant bound may be expressed by the following relation (see for example [16])

$$\frac{\Delta m^2}{m^2} \lesssim 10^{-3}\frac{m}{500GeV}. \tag{3.21}$$

So if we want soft masses $m \lesssim 1TeV$ in order to address the hierarchy problem the FCNC effect must be suppressed by at least a factor 10^{-3}. Then the following alternatives may be pursued.

- To be consistent with (3.21) we need $\omega^s \lesssim 10^{-3}\frac{1}{\ln m_{3/2}\tau^b}\omega^b$ at the MSSM point. We can get this as follows. The small cycle may be regarded as a blow up (by a P^2) of a singularity. Then (at least in a non-compact Calabi-Yau) it has been shown in [17] that the corresponding harmonic 1,1 form falls off as R^{-6} where R is the distance to the singularity.[6] It is quite plausible that this behavior applies to the small cycle of a compact Calabi-Yau. Then indeed the desired suppression can be obtained if the distance R is identified with the location of the $D3$ brane (or the MSSM collapsed cycle) and is of the order of $\mathcal{V}^{1/6} \sim \tau^{1/4}$. The FCNC suppression is then obtained provided $\mathcal{V} \gtrsim 10^{12}$. This would imply a string scale that is well below the Planck scale since $M_{\text{string}} \sim M_P/\sqrt{\mathcal{V}} \lesssim 10^{12}GeV$. Any hope of getting a GUT scenario within LVS is then eliminated, but we would still have a viable intermediate scale phenomenology. This of course is consistent with the usual LVS scenario as discussed in [18] and references therein. Nevertheless it should be stressed that this conclusion holds only if we ignore quantum - in particular Weyl anomaly and gaugino mediation - contributions to scalar masses (see next section).

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[6]We thank Joe Conlon for suggesting this and bringing reference [17] to our attention.

- The alternative within LVS is to have a very heavy soft mass scale $\gtrsim 10^3 TeV$. In this case the SUSY solution to the hierarchy problem is much more fine-tuned than in the previous one. Nevertheless it is interesting to note that the constraint on the volume is now much weaker and reads $\mathcal{V} \gtrsim 10^3$. Note that in this case one might hope that the gaugino masses are still at the TeV scale (this could have been the case according to (3.14) if the dilaton F-term is not responsible for the uplift and the gaugino mass arises from loop corrections to the gauge coupling function). This scenario is essentially that of split supersymmetry. However again the Weyl anomaly and gaugino mediation effects will modify this.

- The third possibility is to consider compactifications with just one Kaehler modulus. In this case an LVS solution is not possible. But one could by including the α' corrections and race track terms find an intermediate volume ($\mathcal{V} \sim 10^{3-4}$) solution. Of course in this case W_0 the flux superpotential would have to be fine tuned to extremely low values in order to get TeV scale soft masses. This scenario has been discussed in [19].

In the following we will pursue only the first of these alternatives.

4   Quantum Effects

4.1   String loop effects

String loop contributions to the classical contributions considered in the previous section can be estimated. This can be done either from an effective field theory calculation as in [19, 20] or from arguments based on calculations in toy models in string theory [9]. They agree if certain cancellations take place. This is essential if the original LVS minimum is not to be destabilized. In this case these do not give a significant correction to the classical soft terms discussed above.

4.2   Weyl anomaly and gaugino masses

Significant corrections to the soft terms can arise from Weyl anomaly contributions [10, 21–25]. These are independent of the size of the compactification once the value of the gravitino mass is chosen. In particular the gaugino masses are given by the expressions (see appendix B for a discussion)

$$M_a = -b_a \left(\frac{\alpha_a}{4\pi}\right) m_{3/2}. \tag{4.1}$$

Here $a = 1, 2, 3$ index the three standard model gauge groups - respectively U(1), $SU(2)$, $SU(3)$, with couplings $\alpha_a = g_a^2/4\pi$. These expressions when evaluated at the UV scale (assumed to be at or close to the unification scale so that $\alpha_a \sim \alpha_{\rm GUT} \sim 1/25$) give

$$M_1 = \frac{33}{5}\frac{\alpha_{\rm GUT}}{4\pi}m_{3/2}, \ M_2 = \frac{\alpha_{\rm GUT}}{4\pi}m_{3/2}, \ M_3 = -3\frac{\alpha_{\rm GUT}}{4\pi}m_{3/2}. \tag{4.2}$$

These should be treated as initial values for the RG evolution down to the MSSM scale - and numerically have values $O(10^{-2}-10^{-3})m_{3/2}$. The important point here is that these values

are larger than the classical contribution which gave gaugino masses $\lesssim O(m_{3/2}/\ln m_{3/2}\sqrt{\mathcal{V}})$ (see (3.14)), unless $\mathcal{V} \lesssim 10^2$ which is far too small a volume for an LVS scenario because of the FCNC problem.

4.3 Gaugino mediation

Now using the above values as boundary conditions for the RG evolution down to the MSSM scale scalar masses are generated by the gaugino mediation mechanism [11, 12]. As shown in appendix B we get,

$$m_1^2 \sim m_2^2 \sim 10^{-6}m_{3/2}^2, \, m_3^2 \sim 10^{-4}m_{3/2}^2. \tag{4.3}$$

This is to be compared to the (diagonal) classical contribution (3.19) $m^2 \sim m_{3/2}^2 1/(\ln m_{3/2}\mathcal{V})$ which would dominate over (4.3) only if $\mathcal{V} \lesssim 10^3$. But in that case the flavor non-diagonal contribution (see (3.19)) would be far too large (even assuming the suppression by a factor $\mathcal{V}$ of the 1,1 form ω^s that we argued for in the discussion after (3.19)).

Thus we require that (4.3) dominates the FCNC contribution in (3.19) which is $\Delta m^2 \sim K'K^{-1}m_{3/2}^2/\sqrt{\tau^b} \sim m_{3/2}^2/(\tau^b)^2$ (see (3.19)(3.21)) , by at least a factor of 10^3. i.e. we need to have

$$\frac{\Delta m^2}{m_3^2} \sim \frac{10^4}{(\tau^b)^2} \lesssim 10^{-3}. \tag{4.4}$$

This gives

$$\mathcal{V} \sim (\tau^b)^{3/2} \gtrsim 10^5.$$

This would yield an effective string scale of $M_{\text{string}} \lesssim 1/\sqrt{\mathcal{V}} \sim 10^{-2.5}M_P \sim 10^{15.5}GeV$ which may just accommodate a GUT scenario.

Be that as it may, to have a SUSY solution to the hierarchy problem in a GUT scenario clearly needs fine tuning of the flux superpotential. From (4.3) we see that to get TeV scale squark masses $m_3 \sim 1TeV$, the gravitino mass must be $m_{3/2} \sim \frac{|W|}{\mathcal{V}} \sim 10^2TeV$. For $W \sim O(1)$ this gives $\mathcal{V} \sim 10^{13}$ (well above our lower limit) and a string scale $M_{\text{string}} \sim M_P/\sqrt{\mathcal{V}} \sim 10^{12}$ which is certainly well below the GUT scale. To get a scale close to the GUT scale would need a highly fine-tuned $W \sim 10^{-8}$.

5 Conclusions

In this paper we have discussed type IIB compactifications on "Swiss Cheese" type manifolds with the MSSM either on a D3 brane(s) at a singularity or on a stack of D7 branes which wrap a 4-cycle. The conflict between chirality and the generation of a non-perturbative superpotential leads to the conclusion that the latter (MSSM) cycle actually collapses below the string scale (or to a singularity) and its modulus does not contribute to SUSY breaking, so that effectively one can ignore it. The following results were derived in this set up.

1. The (classical) diagonal soft mass is given by $m_{\text{soft}} = \sqrt{\frac{3\hat{\xi}}{16\ln m_{3/2}\mathcal{V}}} m_{3/2}$. However there is also a flavor violating contribution to the mass matrix which needs to be suppressed. Assuming that the 1,1 form associated with the small cycle falls off with distance, as in the non-compact case studied in [17], we get a lower bound $\mathcal{V} \gtrsim 10^{12}$ in order that there is sufficient suppression of FCNC effects relative to the classical soft mass.

2. However the Weyl anomaly contribution to the gaugino mass actually dominates the classical contribution. Furthermore this generates, through the mechanism of gaugino mediation, a quantum contribution at the MSSM scale to the scalar soft masses, that is also larger than the classical contribution. Hence the scenario of item 1. above gets modified. These LVS models appear to give a string theoretic construction of a sequestered situation as envisaged in [21],[7]. The suppression of FCNC effects now only gives the weaker constraint $\mathcal{V} \gtrsim 10^5$ (see discussion after (4.4)).

3. The gravitino mass in this scenario is $m_{3/2} \sim 10^2 TeV$ if we wish to have TeV scale SUSY breaking MSSM soft terms. The gravitino gives no cosmological problems but the sGoldstino (light modulus) mass would be $m_{\text{mod}} \sim m_{3/2}/\sqrt{\mathcal{V}} < 1TeV$ so this scenario appears to suffer from the cosmological modulus problem. Again the lower bound on $\mathcal{V}$ is compatible with this estimate of the gravitino mass only if W_0 is highly fine-tuned to values around 10^{-7}. If $W_0 \sim O(1)$ then $\mathcal{V} \sim 10^{13}$, we have an intermediate string scale, no possibility of Grand Unification, and a very light modulus $\sim 100MeV$! Furthermore there would be a serious μ problem since as we see from (3.17) μ is highly suppressed for large volumes.

4. The resolution of the cosmological modulus problem necessitates raising the gravitino mass to $m_{3/2} \sim 10^3 TeV$. However the Weyl anomaly generated gaugino mass is now $\sim 10TeV$ with gaugino mediated soft masses which are at least a few TeV. We may take the volume $\mathcal{V}$ close to its minimum possible value 10^3, consistent with the now somewhat less stringent FCNC constraint (3.21) (since the squark mass is higher). Then we have a somewhat less fine-tuned $W_0 \sim 10^{-6}$ and a light modulus $\sim 10TeV$. But of course addressing the hierarchy problem would require somewhat more fine-tuning. If the cosmological modulus problem is taken seriously then this should be considered the preferred solution within this class of models.

The LVS compactification of type IIB string theory thus gives us an appealing class of models of supersymmetry breaking and transmission. It is consistent with all theoretical

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[7]To truly establish sequestering one would of course need to demonstrate that the couplings of the bulk and brane fields are suppressed at the quantum level as well. This would require a calculation of string loop effects beyond that considered in the literature. For some conjectures regarding this see [8] and references therein. It should also be noted that these assumptions are in agreement with the effective field theory estimates given in [19]. Essentially the point is that the quantum corrections are also expected to be suppressed by large volume factors, so that it seems unlikely that these corrections could significantly change the classical results. However it would be nice if this could be supported by a detailed calculation.

constraints and satisfies phenomenological constraints (and in the case of 4 above cosmological ones as well) . It is predictive and is expected to be ultra-violet complete. As argued in the introduction a bottom up approach ignores the necessary embedding of the low energy (or intermediate scale) hidden sector dynamics in the larger framework of a string theory. The essential point here is that one cannot ignore the dynamics of the string theory moduli and focus on some additional sector (as is usually done in GMSB) since in SUGRA, both open string fields and closed string moduli are coupled together in a highly non-linear fashion. This class of models, with LVS compactifications, generically break supersymmetry and provide, within a string theory context, a very compelling and phenomenologically viable scenario in which to discuss MSSM SUSY breaking. The detailed phenomenology of these models will be discussed elsewhere [26].

Acknowledgments

A   Soft scalar mass calculation

We give here some details of the calculation of the F-terms of the two Kaehler moduli τ^b, τ^s. Useful formulae for calculating the metrics inverse metrics and Riemann tensors for Kaehler potentials of the form $K = -n \ln Y$ are given in appendix A of [27]. In our case $n = 2$ and $Y = \mathcal{V} + \frac{\hat{\xi}}{2}$ where $\mathcal{V} = (\tau^b)^{3/2} - (\tau^s)^{3/2}$ where $\tau^i = \frac{1}{2}(T^i + \bar{T^i})$ with T^i being the holomorphic Kaehler moduli. As usual $K_i \equiv \partial_{T^i} K$, etc. We find

$$K^{i\bar{j}} K_{\bar{j}} \sim -2\tau^i - \frac{3}{2}\hat{\xi}\frac{\tau^i}{\mathcal{V}}, \tag{A.1}$$

and

$$K^{s\bar{s}} = -2\left(\mathcal{V} + \frac{\hat{\xi}}{2}\right)\left(-\frac{4}{3}(\tau^s)^{1/2} + 4(\tau^s)^2 + O(\mathcal{V}^{-1})\right), \tag{A.2}$$

$$K^{b\bar{s}} = 4\tau^b\tau^s(1 + O(\mathcal{V}^{-1})). \tag{A.3}$$

Also note that the stabilization of the axion corresponding to the small modulus results in a sign flip (i.e. the axion takes a value that is an odd multiple of π after choosing without loss of generality the phases of W_0, A to be zero) so that effectively the superpotential is $W = W_0 + Ae^{-aT^s} = W_0 - Ae^{-a\tau^s}$ [5]. Then we find (note that $F^i \equiv e^{K/2} K^{i\bar{j}}(\partial_{\bar{j}}\bar{W} + K_{\bar{j}}\bar{W})$) using the solution (2.10) (2.11),

$$F^b = -\tau^b\left(2 + \frac{3}{2} \cdot \frac{\hat{\xi}}{4a\tau^s}\frac{1}{\mathcal{V}}\right)m_{3/2}, \tag{A.4}$$

$$F^s = -\frac{3\tau^s}{2a\tau^s}m_{3/2}(1 + O(\mathcal{V}^{-1})). \tag{A.5}$$

(F^b was also calculated in [9]).

Here we compute the matter metric for the matter located on a D3 brane which sits at a point in the internal Calabi-Yau space. We expect that a similar formula will be valid for matter on a stack of D7 branes wrapping a collapsing four cycle.

The holomorphic Kaehler moduli T^i are related to the moduli τ^i (in terms of which the volume is given by (2.6)) by (see equation (3.13) of [15])

$$T^i + T^{\bar{i}} = 2\tau^i + 2\mu l^2 i\omega^i_{\alpha\bar{\beta}}C^\alpha C^{\bar{\beta}} + (CUC + \bar{C}\bar{U}\bar{C})\text{term} + \ldots \tag{A.6}$$

to linear order in the complex structure moduli U. Here ω^i are the harmonic 1,1 forms on the CY orientifold evaluated at the position of the D3 brane, C^α are the matter fields on the D3 brane, l is the axionic partner of the dilaton ($S \equiv e^{-\phi} - il$), and μ is the tension of the D3 brane. Writing $\mathcal{V}_i \equiv \partial\mathcal{V}/\partial\tau^i$,

$$K_\alpha \equiv \frac{\partial K}{\partial C^\alpha}|_{T,U,S} = -2\frac{\mathcal{V}_i}{Y}\frac{\partial\tau^i}{\partial C^\alpha}|_{T,U,S}$$

Differentiating (A.6) with respect to C keeping the moduli T, U, S fixed we have

$$0 = 2\frac{\partial\tau^i}{\partial C^\alpha} + 2\mu(i\omega^i_{\alpha\bar{\beta}})C^{\bar{\beta}} + O(UC)$$

$$0 = 2\frac{\partial^2\tau^i}{\partial C^\alpha \partial C^{\bar{\beta}}} + 2\mu(i\omega^i_{\alpha\bar{\beta}})$$

Hence we have

$$K_{\alpha\bar{\beta}} = 2\mu i\omega^i_{\alpha\bar{\beta}}\frac{\mathcal{V}_i}{Y} + O(C^2) = \frac{3\mu}{Y}(i\omega^b_{\alpha\bar{\beta}}\sqrt{\tau^b} - i\omega^s_{\alpha\bar{\beta}}\sqrt{\tau^s}) + O(C^2), \tag{A.7}$$

where in the last step we specialized to the Swiss cheese CY manifold (2.6) (with $\tau^a \to 0$). Note also for future reference that

$$Z = K_{H_1 H_2} \sim \frac{\mathcal{V}_i}{Y}\frac{\partial^2\tau^i}{\partial H_1 \partial H_2}. \tag{A.8}$$

To compute the soft masses (and other soft terms) we need the sectional curvatures.

First let us ignore the second term in parenthesis in (A.7). In this case the matter metric is conformal to the 1,1 harmonic form of the large modulus $K_{\alpha\bar{\beta}} = f(\tau^b, \tau^s)i\omega^b_{\alpha\bar{\beta}}$, $f \equiv 3\mu\sqrt{\tau^b}/Y$ and the relevant components of the Riemann tensor are easily computed using the formula $R_{i\bar{j}\alpha\bar{\beta}} = \partial_i\partial_{\bar{j}}\ln f K_{\alpha\bar{\beta}}$. This gives

$$R_{b\bar{b}\alpha\bar{\beta}} = \frac{1}{4(\tau^b)^2}\left(1 + \frac{15}{16}\frac{\hat{\xi}}{a\tau^s\mathcal{V}}\right)K_{\alpha\bar{\beta}}, \tag{A.9}$$

$$R_{b\bar{s}\alpha\bar{\beta}} = -\frac{9}{16}\frac{(\tau^s)^{1/2}}{(\tau^b)^{5/2}}K_{\alpha\bar{\beta}}, \tag{A.10}$$

$$R_{s\bar{s}\alpha\bar{\beta}} = \frac{3}{16}\frac{(\tau^s)^{-1/2}}{(\tau^b)^{3/2}}K_{\alpha\bar{\beta}}. \tag{A.11}$$

For future use we will recalculate $R_{b\bar{b}\alpha\bar{\beta}}$ keeping both terms in the parenthesis in (A.7) (but ignoring the $\hat{\xi}$ dependence for simplicity), and using the general formula for the Riemann tensor in Kaehler geometry

$$R_{b\bar{b}\alpha\bar{\beta}} = \partial_b \partial_{\bar{b}} K_{\alpha\bar{\beta}} - K^{\gamma\bar{\delta}} \partial_{\bar{b}} K_{\alpha\bar{\delta}} \partial_b K_{\gamma\bar{\beta}},$$

$$= \frac{3\mu}{4(\tau^b)^3} \left(i\omega^b_{\alpha\bar{\beta}} - \frac{7}{4} \sqrt{\frac{\tau^s}{\tau^b}} i\omega^s_{\alpha\bar{\beta}} \right), \tag{A.12}$$

$$= \frac{1}{3} K_{b\bar{b}} \left(K_{\alpha\bar{\beta}} - K'_{\alpha\bar{\beta}} \sqrt{\frac{\tau^s}{\tau^b}} \right). \tag{A.13}$$

Here we have defined

$$K'_{\alpha\bar{\beta}} \equiv \frac{9\mu}{4} \frac{i\omega^s_{\alpha\bar{\beta}}}{\tau^b},$$

to be compared with $K_{\alpha\bar{\beta}} \sim \frac{3\mu}{\tau^b} i\omega^b_{\alpha\bar{\beta}}$.

Let us use the above results to calculate soft masses. These are given by

$$m^2_{\alpha\bar{\beta}} = V_{\text{class}}|_0 \tilde{K}_{\alpha\bar{\beta}} + m^2_{3/2} \tilde{K}_{\alpha\bar{\beta}} - F^b F^{\bar{b}} R_{b\bar{b}\alpha\bar{\beta}}$$
$$-2 \text{Re} F^b F^{\bar{s}} R_{b\bar{s}\alpha\bar{\beta}} - F^s F^{\bar{s}} R_{s\bar{s}\alpha\bar{\beta}}.$$

First let us compute the flavor diagonal part - i.e. we will ignore the contribution to the matter metric from the harmonic form ω^s. We have

$$F^b \bar{F}^{\bar{b}} R_{b\bar{b}\alpha\bar{\beta}} = m^2_{3/2} \left(1 + \frac{21}{16} \frac{\hat{\xi}}{a\tau^s \mathcal{V}} \right) K_{\alpha\bar{\beta}}$$

$$2 F^b \bar{F}^{\bar{s}} R_{b\bar{s}\alpha\bar{\beta}} = -\frac{27}{8} \frac{\hat{\xi}}{2a\tau^s \mathcal{V}} m^2_{3/2} K_{\alpha\bar{\beta}}$$

$$F^s \bar{F}^{\bar{s}} R_{s\bar{s}\alpha\bar{\beta}} = \frac{27}{64} \frac{\hat{\xi}}{2(a\tau^s)^2 \mathcal{V}} m^2_{3/2} K_{\alpha\bar{\beta}}$$

So we have

$$m^2_{\alpha\bar{\beta}} = (V_{\text{class}}|_0 + \frac{3}{8} \frac{\hat{\xi}}{a\tau^s \mathcal{V}} m^2_{3/2}) \tilde{K}_{\alpha\bar{\beta}} + \text{flavor non} - \text{diagonal}$$

Now let us compute the flavor non-diagonal piece. The leading contribution comes from the extra contribution to $F^b \bar{F}^{\bar{b}} R_{b\bar{b}\alpha\bar{\beta}}$ coming from the term proportional to $K'_{\alpha\bar{\beta}}$ in the expression for the Riemann tensor (A.13). Using $F^b \bar{F}^{\bar{b}} K_{b\bar{b}} \sim 3m^2_{3/2}$ we find this to be

$$\Delta m^2_{\alpha\bar{\beta}} = \frac{3}{4} \sqrt{\frac{\tau^s}{\tau^b}} m^2_{3/2} K'_{\alpha\bar{\beta}}.$$

So unless $K'_{\alpha\bar{\beta}}$ is strongly suppressed relative to $\tilde{K}_{\alpha\bar{\beta}}$ (which means in effect the suppression at the position of the D3brane(s) of ω^s compared to ω^b) there would be a serious FCNC problem.

There is also an issue with the μ term (see (3.17)) that needs to be discussed. If the uplift comes mainly from giving F-terms to the complex structure moduli then from (3.12)

it follows that the $|F_m| \lesssim m_{3/2}/(\sqrt{h_{21}}\sqrt{\ln m_{3/2}\mathcal{V}})$ since in a basis in which the relevant metric is diagonal there are effectively h_{21} terms in the sum. Now the expression for the effective μ term has a contribution $F^m \partial_m Z$. From (A.8) we see that this gives the factor $\sqrt{h_{21}}$ in the upper bound (3.17).

B Gaugino masses

The physical gauge coupling (at the cutoff scale/GUT scale Λ) can be written in the following form [10][28][25]

$$H_i = f_i - \frac{3c_i}{8\pi^2}\tau - \sum_r \frac{T_i(r)}{4\pi^2}\tau_r - \frac{T(G_i)}{4\pi^2}\tau_i. \tag{B.1}$$

Here G_i is the gauge group and $c_i = T(G_i) - \sum_r T_i(r)$[8] where $T(G_i)$, $T_i(r)$ are respectively the trace of a squared generator in the adjoint and the matter representations of the group. The first term is the classical gauge coupling (say at the scale Λ). The second arises from the Weyl anomaly which arises when one does a chiral rotation from the supergravity frame to the Kaehler-Einstein frame. The third term comes from an anomaly associated with the field redefinition $C_\alpha \to e^{\tau z}C_\alpha$ needed to get canonical normalization of the MSSM fields and the last term arises from the redefinition of the gauge field pre-potential $V_i \to e^{(\tau_i + \bar{\tau}_i)/2}V_i$. The chiral fields τ, τ_r, τ_i are fixed by the relations

$$\tau + \bar{\tau} = \frac{1}{3}K|_{\text{harm}}, \tag{B.2}$$

$$\tau_r + \bar{\tau}_r = \ln \det \tilde{K}^{(r)}_{\alpha\bar{\beta}}, \tag{B.3}$$

$$\exp[-(\tau_i + \bar{\tau}_i)]|_{\text{harm}} = \frac{1}{2}(H_i + \bar{H}_i), \tag{B.4}$$

It should be emphasized that (as observed by Kaplunovsky and Louis) the first relation is precisely the one which accomplishes the transformations and field redefinitions that are needed to get to the Einstein-Kaehler frame. These are the same transformations that are done in for example Wess and Bagger [29] in component form to get to the final SUGRA action displayed in appendix G of that work. The last relation comes from supersymmetrizing the lowest component relation $\exp[-(\tau_i + \bar{\tau}_i)]|_0 = 1/g_{\text{phys}}^{(i)2} \equiv \Re H_i|_0$. To get the physical coupling function at a low scale μ ($\ll \Lambda$) we need to evaluate the right hand sides of (B.2)(B.3)(B.4) at the scale μ and make the replacement

$$f_i \to f_i - \frac{b_i}{16\pi^2}\ln\frac{\Lambda}{\mu}. \tag{B.5}$$

Here $b_i = 3T(G_i) - \sum_r T_i(r)$ and we used the fact that f_i is only renormalized at one-loop.

Projecting the lowest component of (B.1) (with the replacement(B.5)) and using the relations (B.2)(B.3)(B.4) gives us (the integrated form of) the NSVZ relation (with SUGRA

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[8]Note that this is the negative of the coefficient defined in [10].

and Weyl anomaly corrections)

$$\frac{1}{g_{\text{phys}}^{(i)2}} = \Re f_i - \frac{b_i}{16\pi^2}\ln\frac{\Lambda}{\mu} - \frac{c_i}{16\pi^2}K|_0$$
$$- \sum_r \frac{T_i(r)}{8\pi^2}\ln\det\tilde{K}_{\alpha\bar\beta}^{(r)}|_0 + \frac{T(G_i)}{8\pi^2}\ln\frac{1}{g_{\text{phys}}^{(i)2}}. \tag{B.6}$$

On the other hand projecting the F-term of (B.1) gives (after solving for $M_i/g_{\text{phys}}^{(i)2}$)

$$\frac{M_i}{g_{\text{phys}}^{(i)2}} = \frac{1}{2}(F^A\partial_A f_i - \frac{c_i}{8\pi^2}F^A K_A - \sum_r \frac{T_i(r)}{4\pi^2}F^A\partial_A\ln\det\tilde{K}_{\alpha\bar\beta}^{(r)})$$
$$\times(1 - \frac{T(G_i)}{8\pi^2}g_{\text{phys}}^{(i)2})^{-1}. \tag{B.7}$$

Let us evaluate this for the theory described in this paper. Since the gauge coupling function $f_a = S$ whose F-term is highly suppressed, the classical contribution can be ignored compared to the AMSB one coming from the last two terms in the first parenthesis. To one-loop (keeping only the large modulus ($T^b \equiv T$) contribution since the other contributions are highly suppressed) we have

$$M_i = -b_i\frac{g^{(i)2}}{16\pi^2}m_{3/2} = -b_i\frac{\alpha_i}{4\pi}m_{3/2} \tag{B.8}$$

To get this we used $F^T = -(T + \bar{T})m_{3/2}$, $K_T = -3/(T + \bar{T})$ and $\tilde{K}_{\alpha\bar\beta} = k_{\alpha\beta}/(T + \bar{T})$. Also note that there is no direct AMSB contribution to the scalar masses since the Weyl anomaly just affects the gauge coupling function - at least at the two derivative level[9].

Nevertheless there is a contribution to the scalar masses coming from 'gaugino mediation' (see [16] for a review). This comes about because of the contribution of gauginos to the running of the scalar masses. The relevant equation is (see section 11 of [16]) assuming GUT unification of coupling constants ($t \equiv \ln\mu/\Lambda$ and $c(r)$ is the quadratic Casimir in r)

$$\frac{dm_{\text{scalar}}^2}{dt} = -\frac{c(r)}{2\pi^2}g_i^2 M_i^2, \tag{B.9}$$

Integrating this using the beta function equations and the RG invariance of M_i/g_i^2,

$$m_{\text{scalar}}^2 = \frac{2c(r)}{b_i}[\frac{g_i^4(\mu)}{g_i^4(\Lambda)} - 1]M_i^2 \simeq 2c(r)\alpha_{\text{GUT}}\ln\frac{\Lambda}{\mu}M_i^2. \tag{B.10}$$

Note that in (B.9)we have only kept the dominant terms. Also we note that the squared scalar masses generated by this mechanism are always positive. Furthermore for μ at the TeV scale and Λ at the GUT scale and $\alpha_{\text{GUT}} \sim 1/25$, the scalar mass is of the order of the gaugino mass which in turn has a magnitude $m_3 \sim 10^{-2}m_{3/2}$, $m_2 \sim m_1 \sim 10^{-3}m_{3/2}$. Thus with $m_{3/2} \sim 100TeV$ we have TeV scale gluino and squark masses. These clearly dominate the classical masses that we obtained even in the most favorable case with $\mathcal{V} \sim 10^6$. These numbers would be somewhat different in the intermediate string scale case, but the above mechanism i.e. RG running from this intermediate scale, will still be the dominant mechanism for generating the scalar masses and the A-term. A detailed acount of the phenomenology in both cases will be presented in [26].

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[9]At the four derivative level of course there are curvature squared contributions.

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