

Moduli stabilization with Fayet-Iliopoulos uplift

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ABSTRACT: In the recent years, phenomenological models of moduli stabilization were proposed, where the dynamics of the stabilization is essentially supersymmetric, whereas an O’Raifeartaigh supersymmetry breaking sector is responsible for the ”uplift” of the cosmological constant to zero. We investigate the case where the uplift is provided by a Fayet-Iliopoulos sector. We find that in this case the modulus contribution to supersymmetry breaking is larger than in the previous models. A first consequence of this class of constructions is for gauginos, which are heavier compared to previous models. In some of our explicit examples, due to a non-standard gauge-mediation type negative contribution to scalars masses, the whole superpartner spectrum can be efficiently compressed at low-energy. This provides an original phenomenology testable at the LHC, in particular sleptons are generically heavier than the squarks.

KEYWORDS: Supersymmetry Breaking, Supersymmetry Phenomenology, Supergravity Models, dS vacua in string theory.

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1. Introduction

Recently, Kachru et al. [1] proposed a strategy to stabilize the moduli in the context of type IIB string theory orientifold, following earlier work [2]. The KKLT set-up involves three steps to achieve a SUSY breaking Minkowski vacuum, while stabilizing all moduli. We will consider in this study a KKLT-like model where all the complex-structure moduli are fixed by the introduction of background fluxes for NS and RR forms. All steps except the last one (uplifting through the addition of one anti D3-brane, analyzed in detail in [3]) can be understood within the context of an effective supergravity. Whereas several attempts [4] tried to use the D-term to uplift the supersymmetric minima, it was shown that this can work only for a gravitino mass of the order of the GUT scale. It was however possible to obtain TeV gravitino mass by introducing corrections to the Kahler metric [5, 6]. Other works insisted on the possibility of using F-terms of matter fields in a decoupled sector to uplift the anti-de Sitter minima through metastable vacua [7–9].

In this note we describe a new way to obtain de Sitter space with a TeV gravitino mass by using a Fayet-Iliopoulos (FI) model [10] as uplift sector. The uplifting is realized through the appearance of a non-zero F -term induced by the vev 's of matter-fields charged under an anomalous $U(1)_X$. The F -term is directly induced through the D-term contribution in the minimization procedure. Moreover, the $U(1)_X$ invariance of the superpotential implies a natural coupling between the moduli fields and the matter charged fields under the $U(1)_X$, which changes substantially the pattern of soft breaking mass terms compared to KKLT. The framework can be naturally realized in orientifolds with internal magnetic fields and is simple enough to be able to address detailed phenomenological questions. One of the main advantages compared to previous uplifts [3, 8, 9] is a larger contribution of the modulus to supersymmetry breaking, which increases the tree-level gaugino masses. Moreover, due to the details of the model mostly related to anomaly cancelation, it is natural to introduce messenger-like fields which realize a very particular version of the gauge mediation proposed some time ago by Poppitz and Trivedi [11] (see also [12]), in which gauge mediation contributions to scalar masses are negative. This naturally leads us to a mixed gravity-gauge mediation scenario, where gauge contributions of a non-standard type [11] are generated at high scale and compete with the gravity contributions. The resulting soft spectrum at low-energy has new features compared to other supersymmetry breaking schemes, in particular the spectrum is compressed, i.e. gauginos and scalar masses have values closer to each other than in mSUGRA, gauge mediation or the mixed modulus-anomaly mediation. For related phenomenological analysis of string compactifications with stabilized moduli, see e.g. [13]. The plan of our paper is the following. In section 2 we review the various uplift mechanisms, insisting on the (non)decoupling of the sector realizing the cancelation of the cosmological constant. In section 3 we define our working model, based on a FI sector with an anomalous $U(1)_X$ gauge symmetry, and analyze its vacuum structure. Section 4 presents a microscopic realization in terms of string orientifold models with internal magnetic fields and invoking stringy and spacetime instantons effects in order to obtain the main couplings of our model. In section 5 we couple our supersymmetry breaking sector to MSSM and analyze the resulting superpartner spectrum at high and low-energy from the viewpoint of electroweak symmetry breaking. In section 6, by using anomaly cancelation arguments, we enlarge our model by adding messenger like fields, chiral with respect to the $U(1)_X$ symmetry. The messenger fields have a peculiar spectrum, in particular $Str M^2 > 0$ and will generate, via gauge mediation diagrams, non-standard gauge contributions [11], which will change the low-energy spectrum in an interesting way. We end with some brief summary of results and conclusions. The appendix contain a more detailed derivation of the crucial term coupling the SUSY breaking sector to the modulus sector.

2. Uplifting and decoupling

The philosophy advocated in [1] to stabilize moduli with zero cosmological constant was to separate the process into three steps:

- Add all possible fluxes in order to stabilize most of the (in type II, the dilaton and the complex) moduli fields.

- Add additional (nonperturbative in type IIB) effects in order to stabilize the remaining (Kahler moduli in type II) moduli. The corresponding dynamics is supersymmetric, generically generating a negative cosmological constant.
- Uplift the vacuum energy to zero by a source of supersymmetry breaking which perturbs only slightly the steps above.

The last step was realized originally in [1] by adding an anti D3 brane at the end of the throat in the internal manifold, while later on it was argued [7, 8] and explicitly shown [9] that this can be naturally realized with the help of a decoupled sector breaking dynamical supersymmetry in the rigid limit. In its manifestly supersymmetric realization, by denoting collectively T_α the moduli left unstabilized after the first above and by χ_i the fields responsible for dynamical supersymmetry breaking and the uplift of the vacuum energy

$$K^{i\bar{j}} D_i W \bar{D}_{\bar{j}} \bar{W} = 3 m_{3/2}^2 M_P^2, \tag{2.1}$$

the decoupling of the two sectors is symbolically described in an effective supergravity action by writing

$$\begin{aligned} W &= W_1(T_\alpha) + W_2(\chi_i), \\ K &= K_1(T_\alpha, \bar{T}_\alpha) + K_2(\chi_i, \bar{\chi}_i). \end{aligned} \tag{2.2}$$

The result of this decoupling is the generation of the scalar potential of the form

$$V \simeq V_{\text{SUSY}}(T_\alpha, \bar{T}_\alpha) + \frac{1}{(T_\alpha + \bar{T}_\alpha)^p} V_{\text{uplift}}(\chi_i, \bar{\chi}_i) + \frac{\chi_i \bar{\chi}_i}{M_P^2} V_1(T_\alpha, \bar{T}_\alpha) + \dots, \tag{2.3}$$

where the index p depends on details of the uplift sector and the term V_1 represents the first term in an expansion which mixes non-trivially, due to supergravity interactions, the modulus sector with the uplift sector. For the case of interest $\langle \chi_i \rangle / M_P \ll 1$, the decoupling is very efficient and has the main consequence of perturbing very little the supersymmetric modulus stabilization dynamics. This reflects itself in the very small contribution of the modulus to supersymmetry breaking, which was estimated in [9], for the case of one modulus, to be

$$K^{T\bar{T}} D_T W \bar{D}_{\bar{T}} \bar{W} \sim \frac{1}{(T + \bar{T})^2} K^{i\bar{j}} D_i W \bar{D}_{\bar{j}} \bar{W} \simeq \frac{3 m_{3/2}^2 M_P^2}{(T + \bar{T})^2}. \tag{2.4}$$

The small contribution of the modulus to supersymmetry breaking in this class of uplifting mechanism has the important outcome that generically the gauginos are much lighter than the gravitino [3, 9]. Consequently, in order to find accurate predictions, one-loop contributions, in particular the anomaly-mediated ones are needed, resulting in the so-called mirage unification of gaugino masses.

It was clear from the very beginning that, while such a decoupling renders the uplifting easy to realize, it is by no means mandatory for the stabilization with zero vacuum energy. It is indeed conceivable to contemplate the possibility of a sector breaking supersymmetry that, due to various reasons, in particular gauge invariance consistency constraints, has

a non-trivial coupling to the modulus (KKLT) sector. This non-decoupling was actually forced upon us by gauge invariance in the D-term attempts to uplift the vacuum energy [4], having as a result a very heavy gravitino mass. Notice that in the original version with anti D3 branes [1], a naive attempt to couple more strongly the two sectors by increasing V_{uplift} results actually in a run-away potential which destroys the minimum.

In the next sections we provide explicit examples where this non-decoupling is successfully realized.¹ Similarly to the D-term uplifting models, the non-decoupling is unavoidable due to gauge invariance constraints. In the present case, due to the use of a FI uplift sector, the novelty is the presence of a new supersymmetry breaking source which generates a positive vacuum energy, similar to the F-term uplifting models [8, 9]. As a result, compared to (2.4), we get a modulus contribution to SUSY breaking bigger than in [8, 9] by a factor $(4 + q)/3$, where q is a $U(1)_X$ charge that will be defined more precisely later on.

3. The model and its vacuum structure

Our model in its globally supersymmetric limit is a variant of the Fayet-Iliopoulos model of supersymmetry breaking. It has two charged fields Φ_{\pm} of $U(1)_X$ charges ± 1 , a constant term W_0 relevant, as usual, for the supergravity generalization, and a new term parametrized by a constant a which couples Φ_- to the modulus under consideration T . This last term is the main novelty and ensure the gauge invariance of the nonperturbative superpotential term. In the language of $\mathcal{N} = 1$ Supergravity (SUGRA), we consider the gauge invariant superpotential

$$W = W_0 + m \phi_+ \phi_- + a \phi_-^q e^{-bT}, \tag{3.1}$$

where W_0 is an effective parameter coming from having integrated out all complex structure moduli through the use of fluxes. It is worth noticing here that following what was demonstrated in the first reference of [3], we expect that the inclusion of the moduli we integrated out (which are not charged under $U(1)_X$) does not change the dynamic of the model. A Fayet-Iliopoulos (FI) term is generated in the 4D effective action of the form²

$$V_D = \frac{4\pi}{T + \bar{T}} D^2 = \frac{4\pi}{T + \bar{T}} \left(|\phi_+|^2 - |\phi_-|^2 + \frac{\xi^2}{T + \bar{T}} \right)^2. \tag{3.2}$$

The non-perturbative potential is either generated by Euclidean D3-branes or by gaugino condensation [15]. The charged field ϕ_- restores the gauge invariance of the nonperturbative modulus-dependent superpotential [16–18].

Indeed, the $U(1)_X$ gauge transformations act on various fields as

$$\begin{aligned} \delta V_X &= \Lambda_X + \bar{\Lambda}_X, & \delta \Phi_i &= -2q_i \Phi_i \Lambda_X, \\ \delta T &= \delta_{GS} \Lambda_X, \end{aligned} \tag{3.3}$$

¹Recently, another example of non-decoupled sectors was provided in the context of heterotic strings in [14].

²The exact form of the D-term depends in principle on the precise form of the Kahler metric. We take $K = |\phi_+|^2 + |\phi_-|^2 - 3 \ln(T + \bar{T})$ in what follows, but we will comment later on about other options in the analysis. The term ξ^2 can be interpreted as $\xi^2 = 3/2\delta_{GS}$ if the FI term arises from non-trivial fluxes for the gauge fields living on the $D7$ -branes.

where q_i are the charges of the fields Φ_i . Gauge invariance forces the Kahler potential for the modulus T to be of the form $K(T + \bar{T} - \delta_{GS}V_X)$. This leads in turn to the FI term

$$\xi_{FI} = -\frac{\delta_{GS}}{2}\partial_T K = \frac{3\delta_{GS}}{2}\frac{1}{T + \bar{T}} \quad (3.4)$$

and fixes $\xi^2 = 3\delta_{GS}/2$. From eq. (3.1), it is clear that imposing $q = (b/3)b \xi^2$ ensures the gauge invariance of the model. The numerical values we will be interested in what follows are

$$\xi \sim M_P \quad , \quad m \ll M_P \quad , \quad W_0 \ll M_P^3 \quad . \quad (3.5)$$

The first requirement in (3.5) is natural from the string theory viewpoint, whereas the third relation is needed in order to get $m_{3/2} \sim \text{TeV}$; the landscape picture of string theory could be invoked in order to achieve this [1]. The smallness of the mass term m in our model is then an outcome from a proper cancellation (uplift) of the cosmological constant. The most natural explanation for it, in our opinion, is in terms of stringy instanton effects recently discussed in the literature [19], which can provide values $m \sim \exp(-S_E)M_P$, where S_E is the area of the euclidian brane responsible for the mass term.

From eq. (3.1) we can deduce explicitly the F-part of the scalar potential given by

$$V_F = e^K \left(K^{i\bar{j}} D_i W \bar{D}_{\bar{j}} \bar{W} - 3 |W|^2 \right) \quad , \quad (3.6)$$

where $K^{i\bar{j}}$ is the inverse of the $K_{i\bar{j}} = \partial^2 K / \partial \Phi_i \partial \bar{\Phi}_{\bar{j}}$ metric and D_i is the Kahler covariant derivative: $D_i W = \partial_i W + (\partial_i K)W$. Using a conventional Kahler potential of the form $K = |\phi_+|^2 + |\phi_-|^2 - 3 \ln(T + \bar{T})$, we can rewrite eq. (3.6) as:³

$$V_F = \frac{1}{(T + \bar{T})^3} \left[\frac{(T + \bar{T})^2}{3} |W_T - \frac{3}{T + \bar{T}} W|^2 + |D_+ W|^2 + |D_- W|^2 - 3 |W|^2 \right] \quad . \quad (3.7)$$

The scalar potential is given explicitly as

$$\begin{aligned} V(\phi_+, \phi_-, T) = & \frac{1}{(2\text{Re}[T])^3} \left[\frac{(2\text{Re}[T])^2}{3} |ab\phi_-^q e^{-bT}|^2 \right. \\ & + 2\text{Re}[T] \left(ab\phi_-^q e^{-bT} \bar{W} + \bar{a}\bar{b}\bar{\phi}_-^q e^{-\bar{b}\bar{T}} W \right) \\ & \left. + |m\phi_+ + aq\phi_-^{q-1} e^{-bT} + \bar{\phi}_- W|^2 + |m\phi_- + \bar{\phi}_+ W|^2 \right] \\ & + \frac{4\pi}{2\text{Re}[T]} \left[|\phi_+|^2 - |\phi_-|^2 + \frac{\xi^2}{2\text{Re}[T]} \right]^2 \quad . \quad (3.8) \end{aligned}$$

³The Kahler metric of the charged fields Φ_{\pm} can be more complicated and can also depend on T . We checked explicitly that with the Kahler potential $K = -3 \ln(T + \bar{T} - |\Phi_+|^2 - |\Phi_-|^2)$ we obtain very similar results. The main changes are that, in going from the SUGRA basis to a low-energy one, the rescaling of the parameters and fields is different. Therefore keeping the same numerical values of the low-energy parameters, those at high-energy have to be slightly modified. Nonetheless, these new values are consistent with the modified approximations needed with such a changed Kahler metric, and the whole procedure remains valid.

The nonperturbative term has significant consequences both on the resolution of the equation of motion for ϕ_+ and ϕ_- , and on the uplifting mechanism. It is important to notice here that due to the intricate coupling between ϕ_- and T , in solving the equations of motions, we cannot neglect the supergravity corrections to $|F_+|^2$ and $|F_-|^2$ in the scalar potential. In what follows we define as usual

$$F^i = e^{K/2} K^{i\bar{j}} \overline{D_{\bar{j}}W} . \tag{3.9}$$

Asking for a zero cosmological constant at the minimum, we can find immediately a relation at first order between the gravitino mass and the parameters of the model by anticipating that the uplift is mainly induced by F_+ :

$$|F_+|^2 \simeq 3m_{3/2}^2 M_P^2 \quad \rightarrow \quad |m\phi_-| = \sqrt{3} |W_0| . \tag{3.10}$$

Solving now the equations $\partial V_{\text{eff}}/\partial\phi_+ = \partial V_{\text{eff}}/\partial\phi_- = 0$, using the approximations allowed by the choice of the parameters, and always fixing the cosmological constant to zero, we obtain at the first order⁴

$$D = |\phi_+|^2 - |\phi_-|^2 + \frac{\xi^2}{2\text{Re}[T]} = \frac{2\text{Re}[T]}{8\pi} \frac{m^2}{(2\text{Re}[T])^3} . \tag{3.11}$$

$$\phi_- = \sqrt{\frac{\xi^2}{2\text{Re}[T]}} = \sqrt{\frac{3q}{2b\text{Re}[T]}} , \tag{3.12}$$

$$\phi_+ = -\frac{3q}{2b\text{Re}[T]} \left[\frac{2aqe^{-bT}}{m} \left(\frac{3q}{2b\text{Re}[T]} \right)^{\frac{q-3}{2}} - \frac{1}{\sqrt{3}} \right] . \tag{3.13}$$

We can check our approximation by defining a parameter $\tilde{\epsilon}$ which will be fundamental in the calculation of the soft breaking term. $\tilde{\epsilon}$ measures the contribution of T to the uplift:

$$\tilde{\epsilon} = 2\text{Re}[T] \frac{abe^{-bT}\phi_-^q + 3W/(2\text{Re}[T])}{\sqrt{3}m\phi_-} = \frac{F_T}{F_+} . \tag{3.14}$$

Solving $\partial_T V(T, \phi_+, \phi_-)_{\phi_+, \phi_-} = 0$ with the reasonable hypothesis⁵

$$a e^{-bT} \ll W_0 \ll m \tag{3.15}$$

and $\phi_+ \ll \phi_-$ (hypothesis that we check a posteriori), we obtain at first order

$$\tilde{\epsilon} = \frac{4+q}{2b\text{Re}[T]} - \frac{2}{\sqrt{3}}\phi_+ , \tag{3.16}$$

which gives for a typical KKLT value $2b\text{Re}[T] = 60$, $\tilde{\epsilon}(q=1) \sim \frac{1}{12}$ and $\tilde{\epsilon}(q=2) \sim \frac{1}{6}$ which is bigger than the values obtained in sequestered F-term uplifting [9], where

$$\tilde{\epsilon}_{\text{F-uplift}} = \frac{3}{2b\text{Re}[T]} . \tag{3.17}$$

⁴For simplicity, we take all the parameters to be real and we choose the real positive solution for the vev of ϕ_- . The general case of complex parameters does not change significantly the results.

⁵If the conditions (3.15) are violated, it turns out not to be possible to realize the uplift of the cosmological constant with a TeV gravitino mass.

It turns out that the numerical solution of the equation of motion for T is very close to the supersymmetric minimum for T , whereas the numerical solutions for ϕ_+ , ϕ_- are very close to the analytical ones (3.12), (3.13); the deviation from it can be parameterized by expanding in a perturbative parameter of the theory (which is ae^{-bT}/m in our specific case), checking in the meantime the analytical consistency of the whole procedure. In fact, this procedure is remnant of the one used in the original KKLT paper where the authors noticed that the term induced by the anti D3-brane which is proportional to $1/Re[T]^2$, does uplift the potential without disturbing significantly the shape and the value of $Re[T]$ at the minimum. F-term uplifting exhibit similar features in the sense that it can be seen (see eq. (3.7)) as an uplift proportional to $|D_i W|^2/Re[T]$. However, it is important to point two main differences with KKLT models [1] and F-term uplifting ones [8, 9]. Indeed, firstly the F-term breaking parameters F_i are induced by the D-term, which imposes a non vanishing vev for ϕ_+ and especially ϕ_- at the minimum of the potential. Secondly, the gauge invariant term $ae^{-bT}\phi_-^q$ in W imposes more constraints on the parameter space, linking directly the F_T and F_+ in the minimization procedure. It turns out that F_T is more important in this case and participate more to the cosmological constant cancelation. One of the main consequences appears on the gaugino masses ($M_i \propto F_T/(2Re[T])$), which are heavier than in previous uplift schemes. One of the main difference with the models inspired by D-term uplifting is the possibility to achieve a TeV scale SUSY breaking.

At the first order, the value of T at the minimum respects the condition $F_T = 0$, i.e.

$$abe^{-bT}\phi_-^q = -\frac{3W}{2Re[T]} \simeq -\frac{3W_0}{2Re[T]} . \tag{3.18}$$

The mass of the gravitino is given by $W/(2Re[T])^{3/2}$. To illustrate the procedure, we apply the minimization condition to find a phenomenological viable point in the parameter space. We fix W_0 , b and q . ξ^2 is given by the gauge invariance constraint, $t = Re[T]$ (and $m_{3/2}$) are obtained by the minimization procedure, whereas m is fine-tuned to ensure a zero cosmological constant. For the numerical values provided in figure 1, we obtain:

$$\begin{aligned} m_{3/2} &= 3.3 \text{ TeV}, & \sqrt{D} &= 52.3 \text{ TeV}, & t &= 60. M_P , \\ \phi_+ &= -2.9 \cdot 10^{-2} M_P, & \phi_- &= 0.22 M_P . \end{aligned} \tag{3.19}$$

Concerning the contribution of various fields to the uplift, we obtain $F_T \sim F_-$ and $\tilde{\epsilon} = F_T/F_+ \sim 1/12$.

4. Microscopic definition of the model

The setup we are considering is very similar to the one proposed in [20, 18, 4], with slight modifications. We start with type IIB string propagating on a Calabi-Yau manifold, orientifolded with an involution $\Omega' = \Omega\sigma$, $\sigma^2 = 1$ which generate non-dynamical O7 and O3 orientifold planes. They ask for consistency the introduction of D7 and D3 branes. The non-trivial dynamics we will be concerned happen on the D7 branes. The relevant ingredients for our discussion are two stacks of D7⁽¹⁾ and D7^(N) branes , giving rise to

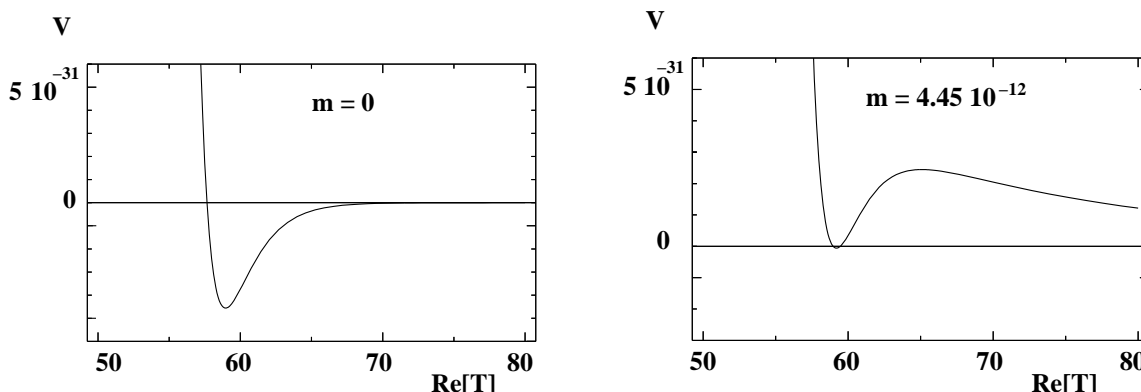


Figure 1: Scalar potential for $m = 0$ (left) and $m \neq 0$ (right) for $W_0 = -4.3 \times 10^{-13}$, $q = 1$ and $b = 0.5$. The other parameters are determined by gauge invariance conditions, minimization of the potential and zero cosmological constant : $m = 3.15 \times 10^{-12}$, $t = 60$. and $\xi^2 = 3$, which gives a gravitino mass of 3.3 TeV.

an $U(1)_X$ and an $U(N)$ gauge groups. The stack $D7^{(N)}$ wraps a four-cycle of volume V , which suitably combined with an axion obtained by wrapping the RR four-form over the four-cycle $a \sim \int C^{(4)}$, forms the complex Kahler modulus $T = V + ia$. The massless chiral open string spectrum for an arbitrary number of stacks of branes can be given a more geometrical interpretation by performing three T-dualities in a IIA setting with intersecting $D6^{(a)}$ branes [22]. In IIA orientifolds with $D6$ branes at angles, each stack $D^{(a)}$, containing M_a coincident branes, has a mirror $D^{(a')}$ with respect to the O6 planes. The chiral spectrum for type II orientifold Calabi-Yau compactification with intersecting branes contains chiral fermions in

sector	representation	multiplicity of states
$D^{(a)} - D^{(b)}$	(\bar{M}_a, M_b)	I_{ab}
$D^{(a')} - D^{(b)}$	(M_a, M_b)	$I_{a'b}$
$D^{(a')} - D^{(a)}$	$\frac{M_a(M_a - 1)}{2}$	$\frac{1}{2}(I_{a'a} + I_{Oa})$
$D^{(a')} - D^{(a)}$	$\frac{M_a(M_a + 1)}{2}$	$\frac{1}{2}(I_{a'a} - I_{Oa})$,

(4.1)

where I_{ab} is the intersection number between the stacks $D^{(a)}$ and $D^{(b)}$, $I_{a'b}$ is the intersection number between the images $D^{(a')}$ and $D^{(b)}$, whereas I_{Oa} is the intersection number between the stack $D^{(a)}$ and the O6 planes. In the original type II language, the intersection numbers are mapped into magnetic fluxes [21, 22].

For the two stack case discussed above and in the type IIA picture, we take the $U(1)_X$ brane to intersect along a six-dimensional subspace with the O-planes. This means that the spectrum of the states stretched between the $U(1)_X$ brane and its image is non-chiral and is described in four-dimensional language by fields ϕ_{\pm} of $U(1)_X$ charges ± 2 . We take the multiplicity of these states, which correspond to the symmetric representations

in (4.1), to be equal to one. If the second stack $U(N)$ does not intersect the O-planes, the symmetric/antisymmetric representations are absent and only the byfundamental chiral multiplets $Q = (N, 1)$ and $\tilde{Q} = (\bar{N}, 1)$, of multiplicity $N_f < N$, are charged under the non-abelian gauge group. In the type IIB language, the axion field coupling to the $U(N)$ gauge fields get charged under the $U(1)_X$ gauge field of the first stack if the two stacks intersect over a two-dimensional cycle on which the magnetic flux is non-trivial [23]. In this case we get the typical Stueckelberg couplings

$$\frac{1}{2}(\partial_\mu a + \delta_{GS} A_\mu)^2, \quad (4.2)$$

rendering the $U(1)_X$ gauge field massive. The supersymmetric description of this phenomenon is precisely the one described in eqs. (3.2)–(3.4).

If $N_f < N$, the non-abelian stack $U(N)$ will undergo gaugino condensation and generate a non-perturbative ADS type superpotential in terms of the "mesonic" fields $M = Q\tilde{Q}$.

$$W_{np} = (N - N_f) \left(\frac{e^{-2\pi T}}{\det M} \right)^{\frac{1}{N - N_f}}. \quad (4.3)$$

It was shown long ago in a similar heterotic context [16] and updated recently for orientifolds [18, 4, 23] that, once the Green-Schwarz anomaly cancellation conditions are imposed, the $U(1)_X$ charges of the mesons are precisely such that the gauge variation of T in (4.3) is compensated by that of the mesons. In addition to the D-term potential (3.2) and the nonperturbative term (4.3), the other terms in the superpotential defining our model are

$$W_1 = W_0 + \lambda \phi_- Q \tilde{Q} + m \phi_+ \phi_- . \quad (4.4)$$

The constant W_0 can be generated by closed string three-form fluxes [2, 1] which stabilize the dilaton and the complex structure moduli, whereas the second term in (4.4) is a disk-level perturbative open string coupling. The last term, which will turn out to be crucial for our purposes, deserves a special discussion. Unless the $U(1)_X$ stack and its image are parallel to each other in some internal subspace, the mass m cannot have a perturbative origin (like for example Wilson lines). In what follows we will advocate a non-perturbative origin $m \ll M_P$. There are two possibilities for generating an exponentially small mass term m . The first option is provided by stringy instanton effects [19]. The instantons under consideration can be $E(-1)$ instantons or $E3$ instantons wrapping cycles different than the one defining the Kahler modulus T under consideration. The resulting parameter m is then proportional to $m \sim \exp(-S_E) M_P$, where S_E is the instanton action. The other option uses a second sector undergoing spacetime nonperturbative dynamics. This could arise, for example, if the $U(1)_X$ brane is part of a bigger stack of branes $U(M) = U(1)_X \times SU(M)$, with the non-abelian part $SU(M)$ undergoing non-perturbative phenomena, for example gaugino condensation $\langle \lambda\lambda \rangle = \Lambda_M^3$. Then an open string perturbative coupling

$$\int d^2\theta \frac{W^{\alpha,M} W_{\alpha,M}}{M_P^2} \phi_+ \phi_- \rightarrow \frac{\Lambda_M^3}{M_P^2} \int d^2\theta \phi_+ \phi_- , \quad (4.5)$$

generates a hierarchically small mass parameter $m = (\Lambda_M^3/M_P^2)$. We can notice that the main assumption made here is that the gaugino condensation scale depends on other moduli, which were already stabilized by high-energy dynamics. Integrating these moduli out could also add higher-order terms in the superpotential and in the Kahler potential, but provided the vev's of $\phi_{\pm} \ll 1$ in Planck units, we expect these contributions not to destabilize our vacuum.

Whereas it is fair to say that constructing a complete, global model along these lines could be a difficult task, there is no conceptual obstruction to the implementation of the ingredients that we need in order to define completely our model in a semi-realistic compactification.

Finally, by invoking the stringy instanton effects described previously or, alternatively, by integrating out the quarks Q, \tilde{Q} of the hidden sector as described in detail in the appendix, we arrive at the generic form of the superpotential

$$W = W_0 + m\phi_+\phi_- + a\phi_-^q e^{-bT}, \tag{4.6}$$

that was defining our model analyzed in the previous section.

5. Soft-breaking terms

In what follows we investigate the effects of supersymmetry breaking in the observable sector, that we take for simplicity to be the Minimal Supersymmetric Standard Model (MSSM). Irrespective on which type of brane MSSM sit (D7 or D3 branes), if they contain magnetic fluxes the gauge kinetic functions contain a T-dependence

$$f_a = \frac{c_a}{4\pi}T + f_a^{(0)}, \tag{5.1}$$

where c_a are positive numbers⁶ and $f_a^{(0)}$ effective constants generated by the couplings of the MSSM branes to other, stabilized fluxes (e.g. the dilaton S). By denoting in what follows by i, j matter fields and by greek indices α any field contributing to SUSY breaking, a relevant quantity for computing the soft terms [24] is the coupling of the matter fields metric $K_{i\bar{j}}$ to the SUSY breaking fields. This can in turn be parameterized as

$$K_{i\bar{j}} = (T + \bar{T})^{n_i} \left[\delta_{i\bar{j}} + (T + \bar{T})^{m_{ij}} |\phi_+|^2 Z'_{i\bar{j}} + (T + \bar{T})^{p_{ij}} |\phi_-|^2 Z''_{i\bar{j}} + (T + \bar{T})^{l_{ij}} (\phi_+\phi_- Z'''_{i\bar{j}} + \text{h.c.}) + O(|\phi_i|^4) \right], \tag{5.2}$$

where $G = K + \log |W|^2$, $K_{i\bar{j}} = \partial_i \partial_{\bar{j}} K$, i and j representing the matter fields, not participating to the SUSY breaking mechanism ($G_i = 0$). The metric $K_{i\bar{j}}$ in (5.2) is written as an expansion in powers of the charged vev fields $\phi_{\pm}/M_P \ll 1$, up the quadratic order.

⁶In the rest of the paper we consider $c_a > 0$. This is easier to obtain in a string setup and also safer for phenomenological purposes, since for $c_a < 0$ there is a serious danger of destabilizing the vacuum.

5.1 Scalar masses

For the calculation of the scalar mass, we use the classical formulas at the linear order in the D-term [18, 25]

$$\tilde{m}_0^2|_{i\bar{j}} = m_{3/2}^2 \left[G_{i\bar{j}} - G^\alpha G^{\bar{\beta}} R_{i\bar{j}\alpha\bar{\beta}} \right] + \sum_a g_a^2 D_a \partial_i \partial_{\bar{j}} D_a, \quad (5.3)$$

with the standard definitions

$$R_{i\bar{j}\alpha\bar{\beta}} = \partial_i \partial_{\bar{j}} G_{\alpha\bar{\beta}} - \Gamma_{i\alpha}^m G_{m\bar{n}} \Gamma_{\bar{j}\bar{\beta}}^{\bar{n}} \quad , \quad \Gamma_{i\alpha}^m = G^{m\bar{k}} \partial_\alpha G_{i\bar{k}}. \quad (5.4)$$

For the uncharged⁷ scalar mass terms we obtain, after normalization of the kinetic terms:

$$(\tilde{m}_0^2)_{i\bar{j}} = m_{3/2}^2 \left[\delta_{i\bar{j}} + \frac{n_i}{(T + \bar{T})^2} |G^T|^2 \delta_{i\bar{j}} - |G^+|^2 (T + \bar{T})^{m_{ij} + \frac{n_i - n_j}{2}} Z'_{i\bar{j}} - |G^-|^2 (T + \bar{T})^{p_{ij} + \frac{n_i - n_j}{2}} Z''_{i\bar{j}} \right]. \quad (5.5)$$

Notice that the contribution to the scalar masses coming from the moduli, depending on the unknown moduli weights n_i , is suppressed compared to the universal first term. This comes actually from the uplift field Φ_+ , via the purely supergravity interactions, as in the mSUGRA case. The third term, coming from the main uplift field Φ_+ , is also negligibly small if $r_{ij} \equiv m_{ij} + (n_i - n_j)/2 \leq -1$, whereas it is comparable to the universal contribution for $r_{ij} = 0$ and dominant for $r_{ij} > 0$. Whereas this last case cannot arise in a string compactification, the case $r_{ij} = 0$ could and deserve a more detailed study from the viewpoint of possible flavor-dependent Φ_+ couplings. Since Φ_+ is a charged and therefore open-string/brane-localized field, whereas the modulus T is a closed/bulk field, the pattern of the flavor dependence of their respective couplings to MSSM fields is clearly different. In particular, whereas it is very difficult to suppress the mixed modulus-MSSM fields couplings (first term in the rhs of (5.2)) in the Kahler potential, this can be easily realized for the uplift open field (the second term in the rhs of (5.2)) Φ_+ .⁸ In this last case (or if the Φ_+ couplings are flavor-universal), the scalar masses (5.5) do not generate dangerous FCNC effects. In conclusion, under reasonable assumptions, the dependence of the soft masses on the unknown quantities $n_i, m_{ij}, Z'_{i\bar{j}}, Z''_{i\bar{j}}$ is weak and can be neglected in a first approximation. For the phenomenological analysis performed in the next section we analyze in detail the universal case, where the gravity-mediated contributions are dominated by the universal term $(\tilde{m}_0^2)_{i\bar{j}} \simeq m_{3/2}^2 \delta_{i\bar{j}}$.

5.2 Gaugino masses

The gaugino mass for a general gauge kinetic function f_a is given by [25]:

$$M_a = \frac{\partial_T f_a}{\text{Re}[f_a]} e^{K/2} K^{T\bar{T}} D_T W. \quad (5.6)$$

⁷We anticipate, for reasons that we discuss later on, that the MSSM fields are neutral with respect to the $U(1)_X$ symmetry.

⁸These comments also apply to a generic F-term uplift [8, 9].

With the hypothesis of a gauge kinetic function given in (5.1), we obtain

$$M_a = m_{3/2} \alpha_a \frac{(T + \bar{T})}{3} \frac{D_T W}{W} = m_{3/2} \alpha_a \frac{(T + \bar{T})}{3} G_T, \quad (5.7)$$

where

$$\alpha_a = \frac{c_a}{c_a + 4\pi f_a^{(0)}/T}. \quad (5.8)$$

For the phenomenological analysis performed in the next section we analyze in detail the unified case $\alpha_a \simeq 1$, which is easily realized for $4\pi f_a^{(0)} \ll c_a T$.

5.3 Trilinear couplings

The general formulas for the trilinear couplings including the D-term contribution can be found in [18, 25]

$$A_{KLM} = e^G [3 + G^\alpha \nabla_\alpha] \nabla_K \nabla_L \nabla_M G + \sum_a g_a^2 D_a \nabla_K \nabla_L \nabla_M D_a, \quad (5.9)$$

where $\nabla_i G = \partial_i G = G_i$, $\nabla_i G_j = G_{ij} - \Gamma_{ij}^k G_k$, etc. It is easy to show that the last contribution in (5.9) coming from the D-term is in our case negligible. Applying it to our special case, after normalization of the kinetic terms and for a typical superpotential for matter fields of the form $W_m = \frac{1}{6} W_{KLM} Q_K Q_L Q_M$, we get

$$\begin{aligned} A_{KLM} = m_{3/2} \left[3W_{KLM}^0 - \frac{G^T}{2(T + \bar{T})} (n_K + n_L + n_M - 3) W_{KLM}^0 \right. \\ \left. + G^T \partial_T W_{KLM}^0 - 3G^+ \bar{\phi}_+ \left((T + \bar{T})^{\frac{n_K - n_i}{2} + m_{K_i}} Z'_{K_i} W_{iLM}^0 \right)_{\text{symm.}} \right. \\ \left. - 3G^- \bar{\phi}_- \left((T + \bar{T})^{\frac{n_K - n_i}{2} + p_{K_i}} Z''_{K_i} W_{iLM}^0 \right)_{\text{symm.}} \right], \quad (5.10) \end{aligned}$$

where *symm.* denotes the symmetrized parts in the (KLM) indices and

$$W_{KLM}^0 = e^{\frac{K}{2}} (K^{-1/2})_K^{K'} (K^{-1/2})_L^{L'} (K^{-1/2})_M^{M'} W_{K'L'M'} = (T + \bar{T})^{-\frac{(3+n_K+n_L+n_M)}{2}} W_{KLM} \quad (5.11)$$

are the low-energy (for canonically normalized fields) Yukawa couplings.

Comments similar to the ones concerning the flavor-dependence of soft masses apply here. Analogously to the discussion concerning soft scalar masses, the dependence of trilinear A -couplings on the unknown quantities n_i, Z'_{ij}, Z''_{ij} can be neglected under reasonable assumptions. For the phenomenological analysis performed in the next section we analyze in detail the gravity-universal case $A_{KLM} = 3m_{3/2} W_{KLM}^0$.

5.4 μ and B_μ terms

The μ parameter and bilinear coupling arises in our model through a Giudice-Masiero mechanism [26]. We will suppose a Kahler metric of the form

$$K = K_0 + Z(T, \bar{T}) [H_1 H_2 + \text{h.c}] \quad (5.12)$$

where $Z(T, \bar{T})$ is a modular function ensuring the modular invariance of the term $Z(T, \bar{T})H_1H_2$. With this convention, we can deduce

$$\mu = m_{3/2} Z(T, \bar{T}) \quad , \quad B_\mu = m_{3/2}^2 [2Z(T, \bar{T}) + G^\alpha \nabla_\alpha Z(T, \bar{T})] + \sum_a g_a^2 D_a \nabla_{H_1} \nabla_{H_2} D_a \quad (5.13)$$

with $D_a \nabla_{H_1} \nabla_{H_2} D_a = -\frac{3}{2} \xi D \partial_T Z(T, \bar{T})$ in our case. The modular function $Z(T, \bar{T})$ allows a certain flexibility of μ and $B\mu$ terms with respect to the gravitino mass. We will use this flexibility in order to determine the appropriate parameters from the analysis of electroweak symmetry breaking in the next section.

5.5 Phenomenology

If we apply the previous soft term calculations to the numerical example of eq. (3.19) we obtain $\tilde{m}_0 = 3.3$ TeV and $M_a = 330$ GeV. In this case, we have a splitting (a factor 10) between scalar masses \tilde{m}_0 and gaugino masses M_a , smaller by a factor of two than in the classical KKLT case. This implies that the one loop contributions (AMSB) are less important here compared to the tree-level one. As we already mentioned, this comes from the fact that G_T participate more actively to the SUSY breaking and therefore its contribution to the gaugino masses is more important compared to the KKLT or classical F-term uplifting cases. We show the spectrum of some typical points in table 1. The absolute value of μ is determined by the minimisation condition of the Higgs potential (assuming CP conservation), but its sign is not fixed. Furthermore, instead of B it is more convenient to use the low energy parameter $\tan \beta = \langle H_2^0 \rangle / \langle H_1^0 \rangle$, which is a function of B and the other parameters. We first derive the soft breaking terms as boundary conditions valid at the GUT scale. For concrete analysis, we will use the standard value of the unification scale $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV. We then evolve these soft terms down to the electroweak scale (at two loops), and impose several theoretical and experimental constraints. The low energy mass spectrum is calculated using the Fortran package SUSPECT [37] and its routines were described in detail in ref. [37]. The evaluation of the $b \rightarrow s\gamma$ branching ratio, the anomalous moment of the muon and the relic neutralino density is carried out using the routines provided the program micrOMEGAS2.0 [38]. Minimizing the Higgs potential in the MSSM leads to the standard relation

$$\mu^2 = \frac{-m_{H_2}^2 \tan^2 \beta + m_{H_1}^2}{\tan^2 \beta - 1} - \frac{1}{2} M_Z^2, \quad (5.14)$$

This minimization condition is imposed at the scale $M_{\text{SUSY}} = \sqrt{\tilde{m}_{\tilde{t}_1} \tilde{m}_{\tilde{t}_2}}$. Eq. (5.14) can be approximated in most cases by

$$\mu^2 \approx -m_{H_2}^2 - \frac{1}{2} M_Z^2 \quad (5.15)$$

When the right hand side is negative, electroweak breaking cannot occur. The Higgs mass parameter $m_{H_2}^2$ is positive at the GUT scale, but decreases with decreasing scale down to M_{SUSY} , through the contributions it receives from RG running, $\frac{\partial m_{H_2}^2}{\partial \log \mu} \approx 6y_t^2(m_{H_2}^2 + m_{U_3}^2 + m_{Q_3}^2 + A_t^2)$. Typically, the value of $m_{H_2}^2$ at the scale M_{SUSY} depends mainly on the soft breaking terms $m_{U_3}^2$ and $m_{Q_3}^2$.

	A	B
W_0	$-7 \cdot 10^{-13}$	$-4.3 \cdot 10^{-13}$
m	$7.3 \cdot 10^{-12}$	$4.5 \cdot 10^{-12}$
a	1	1
b	0.3	0.5
q	1	1
$\tan \beta$	30	15
t	98.3	59.4
μ (GeV)	810	1070
$B\mu$ (GeV) ²	(400) ²	(870) ²
$m_{\chi_1^0}$	110	140
$m_{\chi_1^+}$	220	290
$m_{\tilde{g}}$	760	950
m_h	120	120
m_A	2220	3290
$m_{\tilde{t}_1}$	1380	1770
$m_{\tilde{t}_2}$	1920	2610
$m_{\tilde{c}_1}, m_{\tilde{u}_1}$	2580	3300
$m_{\tilde{b}_1}$	1910	2610
$m_{\tilde{b}_2}$	2310	3230
$m_{\tilde{s}_1}, m_{\tilde{d}_1}$	2580	3300
$m_{\tilde{\tau}_1}$	2290	3200
$m_{\tilde{\tau}_2}$	2420	3230
$m_{\tilde{\mu}_1}, m_{\tilde{e}_1}$	2550	3270

Table 1: Sample spectra. All superpartner masses are in GeV, whereas W_0 , m and t are given in Planck units. Both spectra give a relic density much above the WMAP constraint.

6. Anomalies and gauge messengers

6.1 Anomalies and messengers

Anomaly arguments that we discussed in section 3 and coupling of the MSSM gauge couplings to the T-modulus introduced in the previous section strongly suggest that there should be fields carrying Standard Model quantum numbers charged under the additional $U(1)_X$. Indeed, the couplings (5.1) generate, through the shift of T under $U(1)_X$ gauge transformation (3.3) mixed $U(1)_X - G_a^2$ anomalies, with $G_a = SU(3), SU(2)_L, U(1)_Y$ being a SM group factor. These anomalies imply some SM-charged fields have to carry *positive* (for positive coefficients c_a in (5.1)) $U(1)_X$ charges in order to cancel, via the 4d Green-Schwarz mechanism, the mixed anomalies. There are two generic possibilities that realize this, that we consider in turn:

- The SM quarks and leptons carry $U(1)_X$ charges. In this case, the squarks and the sleptons will acquire D-term soft masses $\tilde{m}_0^2 \sim D \sim 100$ TeV. If we wish to keep

	A+GMSB	B+GMSB
W_0	$-7 \cdot 10^{-13}$	$-4.3 \cdot 10^{-13}$
m	$7.3 \cdot 10^{-12}$	$4.5 \cdot 10^{-12}$
a	1	1
b	0.3	0.5
q	1	1
$\tan \beta$	30	15
t	97.3	59.4
λ	$1.7 \cdot 10^{-3}$	$1.1 \cdot 10^{-3}$
N_{Mess}	6	6
μ (GeV)	186	216
$B\mu$ (GeV) ²	(330) ²	(730) ²
$m_{\chi_1^0}$	120	150
$m_{\chi_1^+}$	160	200
$m_{\tilde{g}}$	850	1060
m_h	120	121
m_A	1740	2770
$m_{\tilde{t}_1}$	990	1220
$m_{\tilde{t}_2}$	1280	1710
$m_{\tilde{c}_1}, m_{\tilde{u}_1}$	1950	2420
$m_{\tilde{b}_1}$	1250	1700
$m_{\tilde{b}_2}$	1930	2690
$m_{\tilde{s}_1}, m_{\tilde{d}_1}$	1950	2420
$m_{\tilde{\tau}_1}$	2130	2870
$m_{\tilde{\tau}_2}$	2160	2960
$m_{\tilde{\mu}_1}, m_{\tilde{e}_1}$	2300	2910
Ωh^2	0.12	0.12

Table 2: Sample spectra including gauge mediation contribution. All superpartner masses are in GeV, whereas W_0 , m and t are in Planck units. The last line correspond to the relic abundance, within WMAP bounds in each case.

some light superpartners and to minimize the fine-tuning of the electroweak scale, one possibility would be to give a charge to the first two generations only [34]. The large hierarchy between the first two and the third generation of squarks can generate various problems, in particular the third generation could become tachyonic through the RGE running towards low-energy [35].

- All MSSM fields (quarks, leptons and Higgses) carry no $U(1)_X$ charges. In this case, there should be additional fields carrying both SM and $U(1)_X$ charges. In order to preserve perturbative gauge coupling unification and be able to give these states a large mass, we only consider complete vector-like $SU(5)$ multiplets, called generically M and \tilde{M} in what follows. Notice that these fields have precisely the features of the

so-called "messenger" fields in gauge-mediation scenarios [28].

These arguments strongly suggest therefore to introduce heavy messengers which can contribute significantly to the soft SUSY masses breaking terms. In our model, coupling the charged field ϕ_- to the messengers pushes naturally the messenger scale up to the GUT scale, still giving rise to important contributions to the scalar masses.

The superpotential is of the form:

$$W_{\text{mess}} = \lambda \phi_- M \tilde{M}, \tag{6.1}$$

where M and \tilde{M} represent the messenger fields of charges q and \tilde{q} respectively. Without loss of generality, we will take $q = \tilde{q} = +1/2$ thorough the rest of the analysis. Notice that the messenger fields, vector-like wrt to SM gauge interactions, are chiral wrt the anomalous $U(1)_X$ symmetry. In (6.1), λ is the low-energy coupling, related to the high-energy supergravity coupling by a formula similar to (5.11). Notice that for zero (or positive) modular weights for Φ_- , M , \tilde{M} , the low energy coupling λ is highly suppressed wrt the high-energy ones by inverse powers of $T + \bar{T}$.

In general, adding messengers to a supersymmetry breaking sector generates a new, supersymmetry preserving vacuum. This is because in order to generate gaugino masses we have to explicitly break R-symmetry, which in turns generically restores supersymmetry [36]. In our case, however, due to the presence of the $U(1)_X$ gauge symmetry, this does not happen; even in the presence of messenger fields, there is no supersymmetry preserving vacuum. This is an important difference compared to standard gauge mediation models of supersymmetry breaking.

Another very important outcome of the charged nature of messenger fields is a new D-term contribution to scalar messenger masses.

The scalar messenger mass matrix is

$$M_{\text{mess}}^2 = \begin{pmatrix} (\lambda\phi_-)^2 + \frac{1}{2}g_X^2 D & \lambda F_- \\ \lambda F_- & (\lambda\phi_-)^2 + \frac{1}{2}g_X^2 D \end{pmatrix} \tag{6.2}$$

Once diagonalized the messenger scalar mass matrix, the two eigenvalues are:

$$m_-^2 = \left[(\lambda\phi_-)^2 + \frac{1}{2}g_X^2 D \right] - \lambda F_- \quad m_+^2 = \left[(\lambda\phi_-)^2 + \frac{1}{2}g_X^2 D \right] + \lambda F_-, \tag{6.3}$$

whereas the fermion mass is given by:

$$m_f = \lambda \phi_- . \tag{6.4}$$

Notice that

$$(Str M^2)_{\text{mess.}} = 2 g_X^2 D \neq 0 . \tag{6.5}$$

By standard gauge-mediation type diagrams, gaugino masses are induced at one-loop, whereas scalar masses are induced at two-loops. Due to (6.5), the computation of the scalar masses is slightly different compared to the standard gauge-mediation models, as shown

by Poppitz and Trivedi [11]. In particular the result is not anymore UV finite, there is a logarithmically divergence term which will play a crucial role in what follows.

In the context of our model, the uplift relation (3.10) has very strong phenomenological implications. Indeed, if D-term contributions appear in the scalar soft mass terms of the visible sector, through the two loops-suppressed GMSB mechanism, it turns out that their magnitude is automatically of the same order as the gravity (i.e. $m_{3/2}$) contribution. This is clearly seen from our numerical example (3.19), in particular from the values of the D -term.

6.2 Soft masses

The exact calculation of the radiatively induced gaugino and scalar masses is performed in [11, 28]. For one messenger multiplet, we obtain for the gaugino mass

$$M_a^{\text{GMSB}} = \frac{g_a^2 m_f S_Q}{8\pi^2} \frac{y_- \log y_- - y_+ \log y_+ - y_- y_+ \log(y_-/y_+)}{(y_- - 1)(y_+ - 1)} \quad (6.6)$$

and for the scalar masses

$$(\tilde{m}_0^{\text{GMSB}})^2 = \sum_a \frac{g_a^4}{128\pi^4} m_f^2 C_a S_Q F(y_-, y_+, \Lambda_{\text{UV}}^2/m_f^2) \quad (6.7)$$

with $y_i = m_i^2/m_f^2$ and where g_a is the corresponding SM gauge coupling (unified at high scale), C_a is the Casimir in the MSSM scalar fields representations (normalized as $C_a(N) = (N^2 - 1)/(2N)$ for the fundamental representation of $SU(N)$ gauge group, while for $U(1)_Y$ it is simply Y^2) and S_Q the Dynkin index of the messenger representation (normalized to $1/2$ for a fundamental of $SU(N)$). The function F is given by

$$\begin{aligned} F(y_-, y_+, \Lambda_{\text{UV}}^2/m_f^2) &= -(2y_- + 2y_+ - 4) \log \frac{\Lambda_{\text{UV}}^2}{m_f^2} \\ &\quad + 2(2y_- + 2y_+ - 4) + (y_- + y_+) \log y_- \log y_+ \\ &\quad + G(y_-, y_+) + G(y_+, y_-), \end{aligned} \quad (6.8)$$

where

$$\begin{aligned} G(y_-, y_+) &= 2y_- \log y_- + (1 + y_-) \log^2 y_- - \frac{1}{2}(y_- + y_+) \log^2 y_- \\ &\quad + 2(1 - y_-) \text{Li}_2\left(1 - \frac{1}{y_-}\right) + 2(1 + y_-) \text{Li}_2(1 - y_-) \\ &\quad - y_- \text{Li}_2\left(1 - \frac{y_-}{y_+}\right). \end{aligned} \quad (6.9)$$

$\text{Li}_2(x)$ above refers to the dilogarithm function and is defined by $\text{Li}_2(x) = -\int_0^1 dz z^{-1} \log(1 - xz)$. After an expansion in the perturbative parameter $\epsilon = \lambda F_- / (\lambda \phi_-)^2$ the mass terms become

$$M_a^{\text{GMSB}} = S_Q \frac{m_0 g_a^2}{8\pi^2} \left(\frac{\phi_+}{\phi_-} \right), \quad (6.10)$$

where m_0 is the low-energy mass parameter of the FI model, equal to $m_0 = m/(T + \bar{T})^{3/2}$ for our model in section 3, and

$$\begin{aligned}
 (\tilde{m}_0^{\text{GMSB}})^2 &= \sum_a \frac{g_a^4}{128\pi^4} C_a S_Q \left[-2g_X^2 D \log \left(\frac{\Lambda_{\text{UV}}}{\lambda\phi_-} \right)^2 + 2g_X^2 D + G(y_-, y_+) + G(y_+, y_-) \right] \\
 &= \frac{m_0^2}{64\pi^4} \sum_a g_a^4 C_a S_Q \left[1 - \log \left(\frac{\Lambda_{\text{UV}}}{\lambda\phi_-} \right)^2 + \left(\frac{\phi_+}{\phi_-} \right)^2 \right], \tag{6.11}
 \end{aligned}$$

where in the last line we used (3.11).

One important feature of eq. (6.11) is the presence of the $\log(\Lambda_{\text{UV}}/\lambda\phi_-)$ term in the soft scalar masses. This logarithmic divergence arises typically in the presence of anomalous $U(1)_X$ that gives a non-vanishing supertrace (6.5) for the messengers superfields [11]. In low energy-GMSB, it usually limits the scale beyond which "new physics" occurs, because the scalars become tachyonic already for $\Lambda_{\text{UV}}/\lambda\phi_-$ around 50. In our specific case, the running is much shorter: from the FI scale (ϕ_-) to the Planck scale (a factor less than 10).

Some remarks are in order concerning the anomaly-mediation contribution to soft terms. For scalar masses, they are completely negligible compared to both gravity and the non-standard GMSB contributions (6.11). For gaugino masses, they are much smaller than the gravity contribution, whereas they are suppressed wrt to the standard GMSB contributions (6.10) only by the number of messenger fields $1/N_{\text{mess}}$. Since we consider relatively large values $N_{\text{mess}} = 6$ in our analysis, we can neglect also the anomaly contributions to gaugino masses in what follows.

6.3 Phenomenological effects

In the complete model, scalar and gaugino masses get contributions both from gravity and the gauge mediation diagrams

$$\begin{aligned}
 (\tilde{m}_0^2) &= (\tilde{m}_0^2)_{\text{grav.}} + N_{\text{Mess}} (\tilde{m}_0^{\text{GMSB}})^2, \\
 M_a &= (M_a)_{\text{grav.}} + N_{\text{Mess}} (M_a^{\text{GMSB}}). \tag{6.12}
 \end{aligned}$$

The negative contribution to the scalar masses \tilde{m}_0 induced by the ultraviolet divergence has strong consequences on the mass spectrum and the phenomenology of the model. It reduces significantly the masses in the left-handed squark sector (the more charged under the SM gauge group) and can have repercussion in the neutralino sector through M_{H_1} . In addition, decreasing the value of $m_{U_3}^2$ and $m_{Q_3}^2$ with gauge mediation naturally decreases the value of μ^2 through eq. (5.15).

We show in table (2) (table (1)) the spectrum with (without) the gauge mediation contributions, after including the RG evolution to low-energy. The scalar spectrum and nature of neutralino (through the μ parameter) are considerably altered. On the other hand, the positive GMSB contribution to gauginos compresses even more the supersymmetric spectrum, especially for a large number N_{Mess} of messenger fields. Notice that, whereas for traditional messenger masses (i.e. around 100 – 1000 TeV), the RG running up to the unification scale forces $N_{\text{Mess}} \leq 3$ in order to avoid strong coupling effects, in

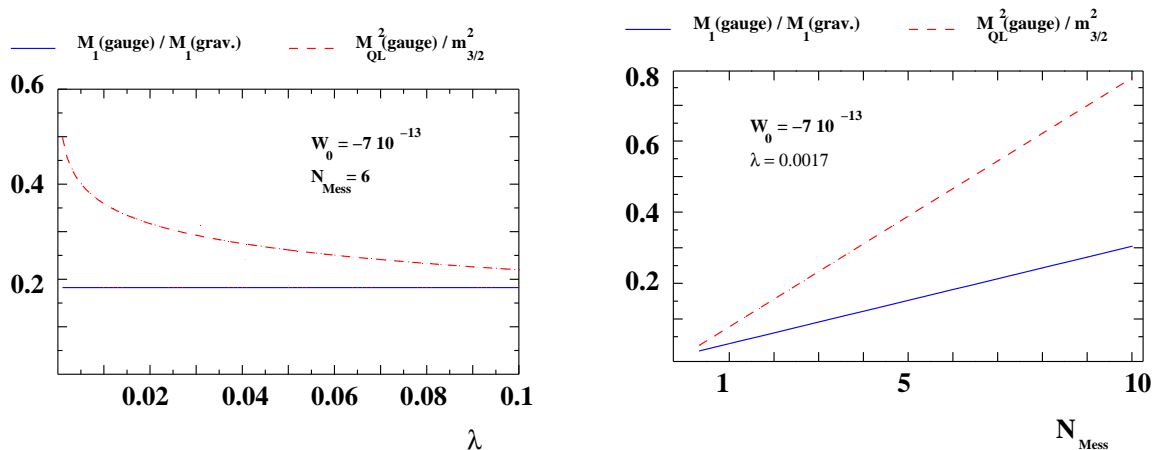


Figure 2: Ratio of the (gauge mediated)/(gravity mediated) for the gaugino masses (blue line) and the scalar masses (red dashed line) as function of λ (left) and the number of messenger (right).

our case since messengers have masses of order 10^{17} GeV, the number of messengers can be larger. By using this (and/or also the alternative possibility of enhancing the negative contribution to scalar masses by decreasing the coupling λ) we can obtain the efficiently compressed spectrum displayed in table 2.

Notice that, in contrast to other scenarios (see for example [27]) where the gauge/gravity relative contributions are completely fixed, in our case, due to the presence of the two charged fields Φ_{\pm} , the gauge and the gravity contributions to soft terms are governed by different parameters. It is instructive to see in figures (2) the dependence of the gauge contribution to soft terms as a function of the relevant parameters of the model (λ and N_{Mess}). For low values of λ , the gauge contribution to the scalar mass becomes important, and even of the same order of magnitude than the gravity contribution for $\lambda \sim 10^{-3}$. Indeed, smaller values of λ implies lighter messenger and thus a larger running between M_{mess} and Λ_{UV} . Gaugino masses are not affected by λ . The number of messenger acts directly on the scalar and gaugino masses, and the gauge contribution becomes relevant in both cases for $N_{\text{Mess}} \sim 6$.

7. Conclusions

In this work, we tried to combine the various ingredients that a microscopic string theory can provide in order to successfully stabilize moduli fields with a TeV gravitino mass. This is realized via a supergravity version of the Fayet-Iliopoulos (FI) model with an anomalous $U(1)_X$ gauge symmetry, in which a (non-linearly charged) modulus field T plays an instrumental role, whereas in turn the FI sector plays a crucial role in supersymmetry breaking and for getting a zero vacuum energy. Due to the non-decoupling between the modulus and the uplift sector, the contribution of the modulus T to supersymmetry breaking is

higher than in previous schemes. This increases the numerical values of the gaugino mass and renders less important the loop contributions to soft terms.

Due to the intricate definition of the model, anomaly arguments strongly suggests the presence in the spectrum of charged fields which have properties similar to messenger fields of gauge mediation of supersymmetry breaking. Due to the charged nature of our messengers, their superpotential couplings are R-symmetric. As a consequence, in contrast to standard gauge mediation scenarios, our messengers do not restore supersymmetry; there is no new supersymmetric vacuum state due to their presence and couplings. Our model has therefore a completely stable non-supersymmetric ground state, which is difficult to realize in more standard gauge-mediation scenarios. Whereas, for the standard reasons, our messengers are vector-like with respect to SM gauge interactions, due to anomaly cancellations they are however chiral with respect to the $U(1)_X$ interactions. Consequently, due to their coupling to the FI supersymmetry breaking sector, they have a particular spectrum, in particular $(StrM^2)_{\text{mess.}} \sim D \neq 0$. The resulting mixed gravity/gauge mediation scenario is therefore of non-standard type: the two-loop gauge mediation contributions to MSSM scalars are negative [11]

The mass scales in the problem are such that gravity and gauge mediation contributions to scalar soft masses are comparable and compete with each other, providing an original predictive spectrum and phenomenology. Indeed, squarks, which are the heaviest superpartners in most mediation scenarios, are here typically the lightest scalars since they get the biggest negative contributions from the non-standard gauge-mediation contribution (6.11). Taking into account the larger than usual gravity-induced gaugino masses and the additional positive contribution to them coming from gauge mediation diagrams, we end up with an original low-energy spectrum in which the whole superpartners spectrum is more compressed than in the usual mSUGRA, gauge or mixed modulus-anomaly mediation scenarios. The complete mixed model of section 6 has the Higgsinos as the LSP and a good relic abundance, compatible with the WMAP bounds. The peculiar details of the spectrum, like the universality of gaugino masses at the unification scale and the negative GMSB contribution to scalar masses, rendering squarks lighter than sleptons, could be tested at LHC and certainly deserve a more focused study.

Finally, whereas in the present paper we get a bigger modulus contribution to supersymmetry breaking than in previous O’Rafaartaigh type models, the main contribution still comes from the uplift sector and more precisely in our case from the Φ_+ field. It is a very interesting and open question to find explicit realizations, with complete moduli stabilization and zero vacuum energy, of the string-inspired supersymmetry breaking parametrizations [25], which assumes moduli/dilaton domination. To our knowledge, this does not seem to be realized in the current known models of moduli stabilization.

A. Dynamical origin for the superpotential

The aim of this appendix is to justify the coupling

$$W = \dots + ae^{-bT} \phi_-^q + \dots \tag{A.1}$$

that we considered in the superpotential of our model. We will show that this term has its origin in strong coupling regime effects for the non-abelian gauge group of the hidden sector, and that in particular it is induced by "integrating out" the mesonic fields which are the right degrees of freedom describing the theory in such regime.

As discussed in section 4, the microscopic description of the model implies $N_f < N$ chiral multiplets in the representations $Q = (N, 1)$ and $\tilde{Q} = (\bar{N}, 1)$ of the gauge group $U(N) \times U(1)_X$. Naturally, at a characteristic scale

$$\Lambda = M_P e^{-\frac{2\pi}{3N-N_f} T}, \tag{A.2}$$

the non-abelian gauge group will enter in a strong coupling regime, it will undergo a gaugino condensation and the model is properly described in terms of the mesonic field $M = Q\tilde{Q}$, of generic charge q' under the $U(1)_X$ gauge group (once normalized at -1 the charge of the field ϕ_-). A non-perturbative ADS potential is generated:

$$W_{np} = (N - N_f) \left(\frac{\Lambda^{3N-N_f}}{\det M} \right)^{\frac{1}{N-N_f}}. \tag{A.3}$$

Including the most general disk-level perturbative open string coupling, and the coupling between ϕ_- and ϕ_+ discussed in section 4, the gauge invariant superpotential reads⁹

$$W = (N - N_f) \left(\frac{\Lambda^{3N-N_f}}{\det M} \right)^{\frac{1}{N-N_f}} + \left(\frac{\phi_-}{M_P} \right)^{q'} \lambda_i^{\bar{j}} M_P M_j^i + m\phi_+\phi_-. \tag{A.4}$$

The auxiliary fields and the D-term can now be calculated [16]:

$$\begin{aligned} (\bar{F}_{M^\dagger})^{\bar{i}}_i &= 2 \left[- (M^{-1})^{\bar{j}}_i \left(\frac{\Lambda^{3N-N_f}}{\det M} \right)^{\frac{1}{N-N_f}} + \left(\frac{\phi_-}{M_P} \right)^{q'} \lambda_i^{\bar{j}} M_P \right] \left[(M^\dagger M)^{\frac{1}{2}} \right]^{\bar{i}}_j, \\ \bar{F}_{\phi_-} &= q' \left(\frac{\phi_-}{M_P} \right)^{q'-1} \text{Tr}(\lambda M) + m\phi_+, \\ \bar{F}_{\phi_+} &= m\phi_-, \\ D &= \left[q' \text{Tr} \left(M^\dagger M \right)^{1/2} - |\phi_-|^2 + |\phi_+|^2 + \frac{\xi^2}{T + \bar{T}} \right]. \end{aligned} \tag{A.5}$$

From a simple analysis of the equations of motions for the mesons, ϕ_- and ϕ_+ , it is possible to see that in the minimum, under the conditions

$$\Lambda^2 < m^2 \ll \frac{\xi^2}{T + \bar{T}} < M_P^2 \tag{A.6}$$

and in particular requiring¹⁰

$$q' N_f (\det \lambda)^{\frac{1}{N}} \left(\frac{M_P}{m} \right) \left(\frac{\langle \phi_- \rangle}{M_P} \right)^{q' \frac{N_f}{N} - 2} \left(\frac{\Lambda}{M_P} \right)^{\frac{3N-N_f}{N}} \ll 1, \tag{A.7}$$

⁹For the aim of this appendix, the term W_0 is completely irrelevant.

¹⁰This in order to assure that $|\langle \phi_+ \rangle| \ll |\langle \phi_- \rangle|$, in accord with what happens in the usual FI global model.

the contribution of the D-term is negligible, $\langle \phi_- \rangle^2 \sim \frac{\xi^2}{T+T}$ and the value of the F-terms for the mesons are very small. These F-terms satisfy the relation

$$\frac{\langle F_M \rangle}{\langle M \rangle} \sim \frac{\langle F_{\phi_-} \rangle}{\langle \phi_- \rangle} \sim q' N_f (\det \lambda)^{\frac{1}{N}} \left(\frac{\Lambda}{M_P} \right)^{\frac{3N-N_f}{N}} \left(\frac{\langle \phi \rangle}{M_P} \right)^{\frac{q' N_f}{N} - 2} M_P. \quad (\text{A.8})$$

Since in the potential the contributions of these F-terms are respectively proportional to $|\langle F_M \rangle|^2 / |\langle M \rangle|$ and $|\langle F_{\phi_-} \rangle|^2$, as long as the vev's for the mesons $|\langle M \rangle|$ is very small compared to $|\langle \phi_- \rangle|^2$, we are therefore allowed to integrate out the mesons M_j^i and the effective superpotential we obtain has the form

$$W^{\text{eff}} = N \left(\frac{\Lambda}{M_P} \right)^{\frac{3N-N_f}{N}} \left(\frac{\phi_-}{M_P} \right)^{\frac{q' N_f}{N}} (\det \lambda)^{\frac{1}{N}} M_P^3 + m \phi_+ \phi_- . \quad (\text{A.9})$$

Nonetheless, as one can show resolving in the first approximation the equation $\bar{F}_{M^\dagger} = 0$, the vev of the meson M_j^i is approximatively

$$\langle M_j^i \rangle \sim N_f (\lambda^{-1})_j^i (\det \lambda)^{\frac{1}{N}} \left(\frac{\langle \phi_- \rangle}{M_P} \right)^{q' \frac{N_f - N}{N}} \left(\frac{\Lambda}{M_P} \right)^{\frac{3N - N_f}{N}} M_P^2 \quad (\text{A.10})$$

and then, using (A.7) and $\lambda_i^{\bar{j}} \sim \delta_i^{\bar{j}}$, we have

$$\frac{|\langle M \rangle|}{|\langle \phi_- \rangle|^2} \ll \left(\frac{m}{M_P} \right) \left(\frac{\langle \phi_- \rangle}{M_P} \right)^{-q'}. \quad (\text{A.11})$$

Therefore, for m small enough (A.6) and for reasonable q' , this ratio is actually $\ll 1$ and the "integration out" is consistent. By re-writing the superpotential (A.9) by using the definitions

$$\begin{aligned} a &= N (\det \lambda)^{\frac{1}{N}}, \\ e^{-bT} &= \left(\frac{\Lambda}{M_P} \right)^{\frac{3N - N_f}{N}}, \\ q &= \frac{q' N_f}{N}, \end{aligned} \quad (\text{A.12})$$

we find exactly the form of the superpotential used in the equation (3.1) once the right powers of M_P restored.

As a numerical example, in order to check if the results obtained in the paper agree with a reasonable nonperturbative scale Λ , we can consider the special case $N = 2$, $N_f = 1$, $q' = 2$, $\lambda_i^{\bar{j}} \sim \delta_i^{\bar{j}}$, i.e. the case studied numerically in the section 5.5. With the choice of the parameters done in that example, we can evaluate

$$\Lambda^{5/2} \sim a e^{-bt} M_P^{5/2} \sim 10^{-13} M_P^{5/2}, \quad (\text{A.13})$$

which means that in this scenario we expect that the non-abelian gauge group $U(N)$ enters in a strong coupling regime at a sensible scale of order $\Lambda \sim 10^{14}$ GeV. Moreover, we can check if in this case the approximations done for the "integration out" step are good. Actually, since $m < 10^{-11} M_P$ and $\langle \phi_- \rangle \sim 10^{-1} M_P$, it is clear that (A.11) is verified and that therefore the whole procedure is consistent.

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