

Moduli stabilization and uplifting with dynamically generated F-terms

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ABSTRACT: We use the F-term dynamical supersymmetry breaking models with metastable vacua in order to uplift the vacuum energy in the KKLT moduli stabilization scenario. The main advantage compared to earlier proposals is the manifest supersymmetric treatment and the natural coexistence of a TeV gravitino mass with a zero cosmological constant. We argue that it is generically difficult to avoid anti de-Sitter supersymmetric minima, however the tunneling rate from the metastable vacuum with zero vacuum energy towards them can be very suppressed. We briefly comment on the properties of the induced soft terms in the observable sector.

KEYWORDS: Supersymmetry Breaking, Compactification and String Models, Supergravity Models.

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1. Introduction and conclusions

Chiral models of dynamical supersymmetry breaking with F-terms were constructed long time ago [1]. Explicit models with supersymmetry breaking ground state are generically relatively involved. More recently, Intriligator, Seiberg and Shih (ISS) proposed a simple, vector-like model with long-lived, metastable supersymmetry breaking vacua [2], whereas the ground state is supersymmetric.¹ On the other hand, in the last couple of years, convincing models of moduli stabilization in string theory were proposed, the prototype being the KKLT scenario [5], based on the orientifolds of IIB string theory flux compactifications [6]. One of the main problems of the KKLT scenario is the uplift of the vacuum energy to zero or positive values. The original proposal of using antibranes relies essentially on nonlinearly realized supersymmetry, whereas the latter attempts [7, 8] to uplift vacuum energy by D-terms, based on the suggestion in [9], lead generically to very heavy (close to the Planck mass) gravitino mass.²

Alternative uplifting using F-terms were already studied in [10–12]. As already stressed in [11], a generic F-type supersymmetry breaking with a supersymmetry breaking scale $\text{TeV} \ll \Lambda_{\text{SUSY}} \ll M_{\text{Pl}}$ can naturally produce the appropriate, intermediate energy scale, for an uplift with a gravitino mass in the TeV range. Dynamical supersymmetry breaking is certainly the best candidate to fulfill this criterion. Metastable vacua have by definition a positive contribution to the vacuum energy which could clearly realize the

¹See [3] for various extensions and string embedding of the ISS proposal and [4] for an earlier proposal.

²It would be very interesting to find explicit counter-examples to this claim.

uplifting required in the KKLT scenario. As we will see in this letter, dynamical supersymmetry breaking in metastable vacua of the ISS type does achieve the goal of uplifting the KKLT vacuum energy to zero, while keeping a TeV scale gravitino mass and therefore leading to low energy supersymmetry. We would like to emphasize, however, that the main ingredient in realizing the uplifting is not the metastable nature of the ISS model. Indeed, as we will briefly mention, other more traditional models [13] of dynamical supersymmetry breaking realize the uplifting in a qualitatively similar way. We argue by explicit examples in both cases that there are generically supersymmetric AdS minima generated by the supergravity interactions, with however Planckian vev's for some fields and therefore not fully trustable in the effective supergravity description. Even by considering seriously these AdS minima, we argue that tunneling from the Minkowski metastable vacuum to the AdS supersymmetric one can be very suppressed.

It would be very interesting to couple the Minimal Supersymmetric Standard Model to our present ISS-KKLT setup, to work out the low-energy phenomenology of the model and to compare it to the existing works [14] based on the original KKLT uplifting prescription relying on antibranes and nonlinearly realized supersymmetry.

The dynamically generated F-term uplifting method can also be combined with the moduli stabilization in type IIA strings [15]. Indeed, D-term uplifting is not available in type IIA strings with moduli stabilization, because of the strong constraints coming from gauge invariance [16]. There are no such constraints in our present setup, therefore there should be no fundamental obstacles in uplifting vacuum energy by non-supersymmetric metastable vacua in type IIA strings with all moduli stabilized.

The structure of this note is as follows. In section 2 we combine the KKLT model of moduli stabilization in type IIB strings with the ISS model of metastable supersymmetry breaking vacuum. We show that in this case the uplifting of the vacuum energy is naturally compatible with a TeV gravitino mass. We discuss supergravity corrections to the globally supersymmetric vacuum, the possibility of a new supersymmetric minimum induced by SUGRA interactions, the effects of gauging the color symmetry in the ISS model and the lifetime of the metastable vacuum. In section 3 we show that qualitatively similar results are obtained by replacing the ISS model with a more traditional model [13] of dynamical supersymmetry breaking. In section 4 we provide some general comments about the tree-level soft masses and under which conditions they could vanish. We then apply the general formulæ for the specific case of the model defined in section 2 and work out some tree-level soft terms, showing that generically tree-level soft masses are of the order of the gravitino mass, whereas gaugino masses can be suppressed in particular cases.

2. Metastable vacua and moduli stabilization

The model is defined by

$$\begin{aligned} W &= W_1(T) + W_2(\chi^i), \\ K &= -3 \ln(T + \bar{T}) + |\varphi|^2 + |\tilde{\varphi}|^2 + |\Phi|^2. \end{aligned} \tag{2.1}$$

In (2.1), χ^i denotes collectively the fields $\varphi_i^a, \tilde{\varphi}_a^{\bar{j}}, \Phi_j^i$ of the ISS model, where $i, \bar{j} = 1 \cdots N_f$ are flavor indices and $a, b = 1 \cdots N$ are color indices. Moreover, in (2.1)

$$\begin{aligned} W_1(T) &= W_0 + ae^{-bT}, \\ W_2(\chi^i) &= h\text{Tr}\tilde{\varphi}\Phi\varphi - h\mu^2\text{Tr}\Phi. \end{aligned} \tag{2.2}$$

Notice that the model is a straightforward combination of the ISS model of metastable supersymmetry breaking vacua with the KKLT model of moduli stabilization. As explained in [2], the sector $\varphi_i^a, \tilde{\varphi}_a^{\bar{j}}$ has a perturbative description in the free magnetic range $N_f > 3N$. The appropriate microscopic theory is an orientifold IIB/ Ω' , with the orientifold operation $\Omega' = \Omega(-1)^{F_L}I_6$, where $(-1)^{F_L}$ is the left spacetime fermion number and I_6 is the parity in the six internal coordinates. The theory contains D3 (O3) branes (orientifold planes) asked by the orientifold operation, with the D3 branes placed at singular points of the compact space in order to reduce supersymmetry to $\mathcal{N} = 1$. Typically there are also D7 (O7) branes (orientifold planes) if other orbifold operations are present. The constant W_0 is generated by 3-form closed string fluxes, as in [6], whereas the nonperturbative T -dependent superpotential could come from gaugino condensation on D7 branes [5] or D3 brane instantons. The gauge sector responsible for the nonperturbative ISS dynamics has a natural embedding on a stack of N D3 “color” branes, with a dynamical scale depending on the dilaton field S , which was already stabilized by three-form fluxes. The mesonic fields Φ are naturally interpreted as positions of a stack of N_f D7 “flavor” branes. This guarantees that their Kähler metric is independent at lowest order on the volume Kähler modulus T [17], as already assumed in (2.1). If the mesons had entered into the no-scale structure of the T -modulus in (2.2), as explained in [11] the vacuum of the theory would have had a marginally unstable direction. The quarks $\varphi, \tilde{\varphi}$ should come from open string in the D3-D7 sector. The Kähler potential of the (anti) quarks of the ISS model are less important for the whole construction. The reason is that they do not contribute to SUSY breaking in the rigid (global) limit and they play no role in the uplifting mechanism at the leading order in the expansion that we analyze in our paper. For this reason and in order to simplify some formulæ we take a canonical Kähler potential for them, too. Whereas a complete string/brane construction involves many other ingredients that we don’t address at all in the present paper, we believe that presenting our uplifting mechanism at the simplified supergravity level is a first step towards the explicit construction of a microscopic (string) theory. We do not attempt here a complete string construction underlying our effective theory, for recent progress see [3]. We point out nonetheless that global string constructions with finite internal space volume are needed in order to achieve this goal.

As transparent in (2.1), the KKLT and the ISS sectors are only coupled through gravitational interactions. In particular, as the ISS gauge group comes from D3 branes, the dynamical scale in the electric theory and therefore also the mass parameter μ in the magnetic theory superpotential (2.2) depend on the dilaton S , which we assume is already stabilized by NS-NS and RR three-form fluxes. We believe this decoupling is instrumental in getting the uplift of the vacuum energy. Another reason for forbidding a coupling of the dynamical supersymmetry breaking sector in the global supersymmetric limit to the

T -modulus is that it is unclear how to formulate the non-abelian Seiberg duality for field-dependent couplings.

At the global supersymmetry level and before gauging the color symmetry, the ISS model has a global symmetry $G = \text{SU}(N) \times \text{SU}(N_f) \times \text{SU}(N_f) \times \text{U}(1)_B \times \text{U}(1)' \times \text{U}(1)_R$, broken explicitly to $\text{SU}(N) \times \text{SU}(N_f) \times \text{U}(1)_B \times \text{U}(1)_R$ by the mass parameter μ . In the supergravity embedding (2.2), the R-symmetry $\text{U}(1)_R$ is explicitly broken. To start with, we consider the ungauged theory, in which the $\text{SU}(N)$ is part of the global symmetry group. At the global supersymmetry level, the metastable ISS vacuum is

$$\Phi_0 = 0 \quad , \quad \varphi_0 = \tilde{\varphi}_0^T = \begin{pmatrix} \mu I_N \\ 0 \end{pmatrix}, \quad (2.3)$$

where I_N is the $N \times N$ identity matrix and $\mu \ll \Lambda_m$, where $\Lambda_m \leq M_{\text{Pl}}$ denotes the mass scale associated with the Landau pole for the gauge coupling in the magnetic theory. The first question to address is the vacuum structure of the model. In order to answer this question, we start from the supergravity scalar potential

$$V = e^K \left[(K^{-1})^{i\bar{j}} D_i W D_{\bar{j}} \bar{W} - 3|W|^2 \right] + \frac{1}{2} (\text{Ref}_a) D_a^2, \quad (2.4)$$

where $\text{Ref}_a = 1/g_a^2$ define the gauge couplings. By using³ (2.1)–(2.2), we find

$$V = \frac{e^{\bar{\chi}_i \chi^i}}{(T + \bar{T})^3} \left\{ \frac{(T + \bar{T})^2}{3} \left| \partial_T W - \frac{3}{T + \bar{T}} W \right|^2 + \sum_i |\partial_i W + \bar{\chi}_i W|^2 - 3|W|^2 \right\}. \quad (2.5)$$

Since $\mu \ll M_{\text{Pl}}$, the vev's in the ISS model are well below the Planck scale. Then an illuminating way of rewriting the scalar potential (2.5) is to expand it in powers of the fields χ^i/M_{Pl} , in which case it reads⁴

$$V = \frac{1}{(T + \bar{T})^3} V_{\text{ISS}}(\chi^i, \bar{\chi}_i) + V_{\text{KKLT}}(T, \bar{T}) + \frac{\bar{\chi}_i \chi^i}{M_{\text{Pl}}^2} V_1(T, \bar{T}) + \frac{1}{M_{\text{Pl}}^3} [W_2(\chi^i) V_2(T, \bar{T}) + \chi^i \partial_i W_2 V_3(T, \bar{T}) + h.c.] + \dots, \quad (2.6)$$

where by comparing (2.6) with (2.5) we can check that $V_1 \sim m_{3/2}^2 M_{\text{Pl}}^2$, and $V_2, V_3 \sim m_{3/2} M_{\text{Pl}}^3$, where as usual $m_{3/2}^2 = |W|^2 \exp(K)$. Notice that the contribution to the vacuum energy from the ISS sector, in the global limit, is

$$\langle V_{\text{ISS}} \rangle = (N_f - N) h^2 \mu^4. \quad (2.7)$$

Since we are interested in small (TeV scale) gravitino mass, it is clear that the first two terms in the rhs of (2.6), V_{ISS} and V_{KKLT} , are the leading terms. Consequently, there should be a vacuum very close to an uplift KKLT vacuum $\langle T \rangle = T_0$ and the ISS vacuum

³The gauge D-term contributions do not exist in the ungauged case we are discussing in this section and will play essentially no role in the following sections.

⁴In most of the formulæ of this letter, $M_{\text{Pl}} = 1$. In some formulæ, however, we keep explicitly M_{Pl} .

$\langle \chi^i \rangle = \chi_0^i$. The KKLT uplift vacuum at the zeroth order T_0 is defined as the minimum of the zeroth order potential $\partial_{T_0} V_0 = 0$, obtained by inserting the ISS vacuum (2.3) into the supergravity scalar potential

$$V_0 = \frac{1}{(T + \bar{T})^3} \left[\frac{(T + \bar{T})^2}{3} |D_T W_1|^2 - 3|W_1|^2 + h^2(N_f - N)\mu^4 \right]. \quad (2.8)$$

In the limit $bT \gg 1$ and for zero cosmological constant, a good approximation for T_0 , considered to be real in what follows, is provided by

$$W_0 + \frac{ab(T_0 + \bar{T}_0)}{3} e^{-bT_0} = 0. \quad (2.9)$$

Notice that in this case T does contribute to supersymmetry breaking⁵

$$F^T \equiv e^{\frac{K}{2}} K^{T\bar{T}} \overline{D_T W} \simeq \frac{a}{(T_0 + \bar{T}_0)^{1/2}} e^{-bT_0}, \quad (2.10)$$

but by an amount suppressed by a factor of $1/b(T_0 + \bar{T}_0)$ compared to the naive expectation.

The cosmological constant at the lowest order is given by

$$\Lambda = V_{\text{KKLT}}(T_0, \bar{T}_0) + \frac{(N_f - N)h^2\mu^4}{(T_0 + \bar{T}_0)^3}, \quad (2.11)$$

which shows that the ISS sector plays the role of an uplifting sector of the KKLT model. In the zeroth order approximation and in the large volume limit $b(T_0 + \bar{T}_0) \gg 1$, we find that the condition of zero cosmological constant $\Lambda = 0$ implies roughly

$$3|W_0|^2 \sim h^2(N_f - N)\mu^4. \quad (2.12)$$

If we want to have a gravitino mass $m_{3/2} \sim W_0/(T_0 + \bar{T}_0)^{3/2}$ in the TeV range, we need small values of $\mu \sim 10^{-6} - 10^{-7}$. Since μ in the model [2] has a dynamical origin, this is natural. Moreover, the metastable vacuum of [2] has a significantly large lifetime exactly in this limit, more precisely when $\epsilon \equiv (\mu/\Lambda_m) \ll 1$. Therefore, a light (TeV range) gravitino mass is natural in our model and compatible with the uplift of the cosmological constant. We believe that this fact is an improvement over the D-term uplift models suggested in [9] and worked out in [8].

Notice that supergravity corrections give tree-level masses to the pseudo-moduli fields of the ISS model. As explained in more general terms in [2], these corrections are subleading with respect to masses arising from the one-loop Coleman-Weinberg effective potential in the global supersymmetric limit. This can be explicitly checked starting from the supergravity scalar potential (2.5) and expanding in small fluctuations around the vacuum (2.3) to the quadratic order.

⁵Notice that the leading order expression for W_0 in (2.9) is not enough for computing F^T , since the subleading terms neglected in (2.9) are needed as well. F^T can be computed directly, however, by keeping the leading terms in the eq. $\partial_T V = 0$.

2.1 The metastable vacuum and supergravity corrections

By coupling the T field to the ISS dynamical supersymmetry breaking system, we expect small deviations from the lowest order vacuum (2.3), (2.9). We expand

$$\chi^i = \chi_0^i + \delta\chi^i \quad , \quad T = T_0 + \delta T, \quad (2.13)$$

where χ_0^i are provided by (2.3), with $\delta\varphi \ll \varphi_0$ ($\delta\tilde{\varphi} \ll \tilde{\varphi}_0$) and $\delta T \ll T_0$. We now turn to the SUGRA corrections to the ISS metastable vacuum (2.13), by linearizing around the ISS-KKLT vacuum the field eqs,

$$\partial_\varphi V = \partial_{\tilde{\varphi}} V = \partial_\Phi V = \partial_T V = 0. \quad (2.14)$$

This can be done by starting from the expansion in the fields χ^i in (2.6), where

$$\begin{aligned} V_1 &= V_{\text{KKLT}} + \frac{|W_1|^2}{(T + \bar{T})^3}, \\ V_2 &= -\frac{1}{(T + \bar{T})^3} [(T + \bar{T})\overline{D_T W} + 3\overline{W_1}] \\ V_3 &= \frac{\overline{W_1}}{(T + \bar{T})^3}. \end{aligned} \quad (2.15)$$

Notice that in the zeroth order vacuum $V_1 \sim m_{3/2}^2 M_{\text{Pl}}^2$, and $V_2, V_3 \sim m_{3/2} M_{\text{Pl}}^3$, as well as $\partial_T V_1 \sim m_{3/2}^2 M_{\text{Pl}}^2$ and $\partial_T V_2, \partial_T V_3 \sim m_{3/2} M_{\text{Pl}}^3$. In order for the linearization to be well-defined, we need to include the Coleman-Weinberg one-loop quantum corrections to the scalar potential discussed in [2]. The reason is that at tree-level and in our zeroth order approximation, there are zero mass particles which, in addition to the Goldstone bosons of the broken symmetries, contain also pseudo-moduli which get their masses at one-loop. After including these corrections, we find at the leading order in the variations $\delta\chi^i, \delta T$ and for zero cosmological constant, that

$$\delta\chi^i \leq O(m_{3/2}) \quad , \quad \delta T \leq O\left(\frac{m_{3/2}}{M_{\text{Pl}}}\right). \quad (2.16)$$

Since in our framework $m_{3/2} \ll \mu$, the condition $\delta\varphi \ll \varphi_0$ is largely satisfied, showing that the expansion (2.13) is an excellent approximation. The precise values of the supergravity corrections (2.16) are not important for what follows. Notice that the small values for $\delta\varphi, \delta\Phi$ in (2.16) are in agreement with the arguments given in [2] stating that high energy microscopic effects in the magnetic theory should not affect significantly the metastable vacuum.

2.2 The SUGRA induced magnetic supersymmetric minimum

In the ISS model and in the case of ungauged $SU(N)$ symmetry, the ISS vacuum (2.3) is actually the true ground state. What happens in the supergravity embedding we are proposing here ? We will show that there is a new, AdS supersymmetric ground state generated by the SUGRA interactions. To find it, we search solutions of the type

$$\varphi = \begin{pmatrix} \varphi_1 \\ 0 \end{pmatrix}, \quad \tilde{\varphi}^T = \begin{pmatrix} \tilde{\varphi}_1 \\ 0 \end{pmatrix}, \quad \Phi = \begin{pmatrix} \Phi_1 & 0 \\ 0 & \Phi_2 \end{pmatrix}, \quad (2.17)$$

to the SUSY preserving equations

$$\begin{aligned}
 D_\varphi W = 0 &\rightarrow h\tilde{\varphi}_1\Phi_1 + \bar{\varphi}_1 W = 0, \\
 D_{\tilde{\varphi}} W = 0 &\rightarrow h\Phi_1\varphi_1 + \bar{\tilde{\varphi}}_1 W = 0, \\
 D_\Phi W = 0 &\rightarrow h\left(\tilde{\varphi}_1^i\varphi_{1,j} - \mu^2\delta_j^i\right) + (\bar{\Phi}_1)^i_j W = 0, \quad i, j = 1 \cdots N \\
 D_\Phi W = 0 &\rightarrow -h\mu^2\delta_m^n + (\bar{\Phi}_2)_m^n W = 0, \quad m, n = N+1 \cdots N_f, \\
 D_T W = 0 &\rightarrow abe^{-bT_m} + \frac{3}{T_m+T_m}W = 0.
 \end{aligned} \tag{2.18}$$

The eqs. (2.18) have the following solution:

$$\begin{aligned}
 \varphi_1 &= \mu_1 I_N, & \tilde{\varphi}_1 &= \mu_2 I_N, \quad \text{with } |\mu_1| = |\mu_2|, \\
 \Phi_1 &= (\mu_1\mu_2 - \mu^2)^{\frac{1}{2}} I_N, & \Phi_2 &= -\frac{\mu^2}{(\mu_1\mu_2 - \mu^2)^{\frac{1}{2}}} I_{N_f - N},
 \end{aligned} \tag{2.19}$$

$$\begin{aligned}
 abe^{-bT_m} - \frac{3h}{T_m+T_m}(\mu_1\mu_2 - \mu^2)^{\frac{1}{2}} &= 0, \\
 h^2(\mu_1\mu_2 - \mu^2) - |W|^2 &= 0.
 \end{aligned}$$

Since cosmological constant cancellation asks for $m_{3/2} \sim \langle W \rangle \sim h\mu^2$, where $m_{3/2}$ is the gravitino mass in the ISS-KKLT vacuum, for $\mu_i \sim \mu$ eq. (2.19) implies in particular $\Phi_2 \sim M_{\text{Pl}}$: the supersymmetric minimum (2.19) depends on the UV properties of the model and is not fully reliable in our effective field theory analysis. For $\mu_1\mu_2 \gg \mu^2$, all vev's are well below M_{Pl} , $\langle W \rangle \gg m_{3/2}M_{\text{Pl}}^2$ and the supersymmetric vacuum (2.19) would be within the validity of the effective supergravity. The second possibility is however incompatible with the condition (2.12) and for a TeV gravitino mass. Therefore we recover the conclusion that $\Phi_2 \sim M_{\text{Pl}}$.

Notice that the supersymmetric vacuum (2.19) survives the gauging of the $SU(N)$ symmetry. Indeed, the $SU(N)$ D-flatness conditions are satisfied, since $|\varphi_1|^2 = |\varphi_2|^2$ and $[\Phi, \Phi] = 0$ in (2.19).

2.3 Gauging the model: infrared description

In the ISS model, the $SU(N)$ symmetry is gauged and corresponds to the gauge group of the magnetic theory. In the electric description, the ISS model is the supersymmetric QCD with N_c colors and $N_c < N_f < 3N_c/2$ quark flavors Q, \tilde{Q} such that in the magnetic description with the gauge group $SU(N_f - N_c)$, the number of flavors is large $N_f > 3N$, where the magnetic theory is in the infrared-free phase. In this case the perturbative magnetic description, around the origin in field space, is reliable. The electric theory has a dynamical scale Λ and a mass term for the quarks $W = m_i^{\tilde{j}} Q^i \tilde{Q}_{\tilde{j}}$. There are N_c vacua described by

$$M_j^i = \left(\frac{1}{m}\right)_j^i (\det m)^{\frac{1}{N_c}} \Lambda^{\frac{3N_c - N_f}{N_c}}. \tag{2.20}$$

The perturbative treatment in the magnetic description translates into the constraint $m_a \ll \Lambda$, where a denotes here the number of light mass eigenvalues, which has to be equal or larger to $N_f + 1$ in order for the metastable vacua to exist. One of the open questions for the ISS model is a dynamical explanation for the constraint $m_a \ll \Lambda$. We believe

that a simple possibility is the following. At high energy there is an additional abelian “anomalous” symmetry $U(1)_X$, with mixed anomalies $U(1)_X SU(N_c)^2$ cancelled by the Green-Schwarz mechanism involving an axionic field a_X . This will render the gauge vector V_X massive and stabilize the complex modulus field containing the axion a_X . There will be an induced Fayet-Iliopoulos term, which in explicit string models is always cancelled by the vev of a scalar field $\langle N \rangle \ll M_{\text{Pl}}$. Mixed anomalies mean that the sum of the quark charges $X_Q + X_{\bar{Q}}$ is not zero and therefore the mass operator $m_i^{\bar{j}} Q^i \bar{Q}_{\bar{j}}$ is not gauge invariant. In generic models, the charge X_N is opposite compared to $X_Q + X_{\bar{Q}}$. We normalize $X_N = -1$ in what follows. Then the superpotential term $y_i^{\bar{j}} (N/M_{\text{Pl}})^{X_Q + X_{\bar{Q}}} Q^i \bar{Q}_{\bar{j}}$ is perturbatively allowed. Supersymmetry could be broken in the process [21], but it can also stay unbroken. In this last case, at energy scales well below the mass of the gauge boson A_X , the net effect of all this is to generate an effective mass term for the quarks of the electric theory $m \sim (\langle N \rangle / M_{\text{Pl}})^{X_Q + X_{\bar{Q}}}$. For large enough quark charges and/or small enough vev $\langle N \rangle$, the induced mass m can be very small.

Denoting by Λ_m the Landau pole of the magnetic theory, according to ISS, for arbitrary vev’s of Φ the quark flavors become massive and can be integrated out. By doing this and by coupling the resulting low-energy system to the KKLT model, we get a lagrangian described by

$$\begin{aligned}
 W &= W_0 + a e^{-bT} + N \left(\frac{h^{N_f} \det \Phi}{\Lambda_m^{N_f - 3N}} \right)^{1/N} - h \mu^2 \text{Tr} \Phi, \\
 K &= -3 \ln(T + \bar{T}) + \bar{\Phi} \Phi.
 \end{aligned}
 \tag{2.21}$$

Similarly to the global supersymmetry analysis of ISS [2], this action has $N_f - N$ supersymmetric vacua, which in the global limit are given by

$$\langle h \Phi \rangle = \Lambda_m \epsilon^{2N/(N_f - N)} I_{N_f} = \mu \frac{1}{\epsilon^{(N_f - 3N)/(N_f - N)}} I_{N_f},
 \tag{2.22}$$

where $\epsilon \equiv \mu / \Lambda_m$. The vacuum in the T -direction is simpler to describe by replacing the vev’s (2.22) in the superpotential (2.21). By doing this, we get an effective superpotential

$$W_{\text{eff}} = W_0 - \frac{(N_f - N) \mu^3}{\epsilon^{(N_f - 3N)/(N_f - N)}} + a e^{-bT}.
 \tag{2.23}$$

Since $W_0 < 0$ in the KKLT model, the effect of the supersymmetric Φ vev’s is to increase the absolute value of the (negative) constant in the superpotential. The values of the minimum for T and the corresponding negative cosmological constant are given approximately by

$$\begin{aligned}
 a b e^{-bT_s} + \frac{3}{T_s + \bar{T}_s} \left(W_0 - \frac{(N_f - N) \mu^3}{\epsilon^{(N_f - 3N)/(N_f - N)}} \right) &\simeq 0, \\
 V_0 &\simeq - \frac{3}{(T_s + \bar{T}_s)^3} \left| W_0 - \frac{(N_f - N) \mu^3}{\epsilon^{(N_f - 3N)/(N_f - N)}} \right|^2.
 \end{aligned}
 \tag{2.24}$$

The supersymmetric ISS vacuum is therefore AdS.

Notice that for $W_0 \gg \mu^3 / \epsilon^{(N_f - 3N)/(N_f - N)}$, we get $T_s \sim T_0$, with T_0 defined in (2.9), since in this case $W \simeq W_0$. If $W_0 \ll \mu^3 / \epsilon^{(N_f - 3N)/(N_f - N)}$, then $T_s < T_0$.

2.4 Lifetime of the metastable vacuum

The model we discussed in this paper has one metastable vacuum and two type of AdS supersymmetric minima. The metastable vacuum will tunnel to the supersymmetric AdS minimum (2.22)–(2.24). The purpose of this section is to provide a qualitative estimate of the lifetime of the metastable minimum, following [18, 19]. The bounce action is expected to come from the path in field space of minimum potential barrier between the metastable supersymmetry breaking vacuum and the supersymmetric vacua. Along this path, the bounce action cannot be computed analytically. For a triangular idealized approximation [19], the bounce action S_b is qualitatively

$$S_b \sim \frac{(\Delta\chi)^4}{\Delta V}, \tag{2.25}$$

where ΔV is the (minimum) barrier along the bounce and $\Delta\chi$ is the variation of the relevant field. For the tunneling between the metastable ISS vacuum (2.3) and the supersymmetric one (2.22) after gauging $SU(N)$, there are two cases. If $\mu \ll \epsilon^{(N_f-3N)/(N_f-N)} M_{\text{Pl}}$, we get

$$h\Delta\Phi \simeq \mu \frac{1}{\epsilon^{(N_f-3N)/(N_f-N)}} \quad , \quad \Delta V \sim \frac{3}{(T_s + \bar{T}_s)^3} |W_0|^2. \tag{2.26}$$

Then, by using the condition (2.12) for a zero vacuum energy in the metastable vacuum, we get

$$S_b \sim \frac{(T_s + \bar{T}_s)^3}{\epsilon^{4(N_f-3N)/(N_f-N)}} \gg 1, \tag{2.27}$$

which increases the lifetime of the metastable vacuum compared to the similar ISS analysis. The reason is that the energy difference between the metastable and the AdS supersymmetric minimum is decreased by the factor $1/(T_s + \bar{T}_s)^3$, resulting in an increase in the bounce action S_b . In the case where $\mu \gg \epsilon^{(N_f-3N)/(N_f-N)} M_{\text{Pl}}$, the vacuum energy of the supersymmetric vacuum (2.24) and consequently ΔV change. The bounce action in this case is

$$S_b \sim \frac{M_{\text{Pl}}^2}{\mu^2} \frac{(T_s + \bar{T}_s)^3}{\epsilon^{2(N_f-3N)/(N_f-N)}} \gg 1. \tag{2.28}$$

The metastable minimum could also tunnel to the supersymmetric minimum (2.19). Even by taking seriously the effective theory analysis in this case, we notice that the AdS supersymmetric minimum (2.19) is far away in the Φ field space from the ISS-KKLT metastable vacuum (2.3), (2.9). The tunneling probability to go to the AdS vacuum (2.19) is highly suppressed and irrelevant for all practical purposes.

3. Uplifting with supersymmetry breaking on the quantum moduli space

As mentioned in the introduction, the important ingredient from the F-term dynamical supersymmetry breaking sector is the intermediate scale for the resulting (positive) contribution to the vacuum energy and not the metastable nature of the vacuum. We discuss now a more conventional non-perturbative hidden sector which, in the global supersymmetry limit, has a non-supersymmetric ground state [13]. Since most of the analysis parallels that

already done for the ISS model, our discussion will be very brief. We consider a SQCD model with $N_c = N_f = 2$ colors and flavors. The effective action which puts together the KKLT moduli stabilization sector and the supersymmetry breaking sector is

$$\begin{aligned} W &= W_0 + ae^{-bT} + \lambda S^{ij} M_{ij} + X(PfM - \Lambda_2^4), \\ K &= -3 \ln(T + \bar{T}) + \text{Tr}\left(\frac{1}{\Lambda_2^2} |M|^2 + |S|^2\right), \end{aligned} \tag{3.1}$$

where $PfM = \epsilon^{ijkl} M_{ij} M_{kl}$, Λ_2 is the dynamical scale of the theory, $M_{ij} = Q_i^a Q_j^a$ are the mesons built up from the quarks Q_i^a with color indices $a = 1, 2$ and flavor indices $i, j = 1, 2, 3, 4$, whereas S^{ij} are gauge singlets. Both fields are antisymmetric in the flavor indices. In (3.1), X is a lagrange multiplier which enforces the eq. describing the quantum deformed moduli space $PfM = \Lambda_2^4$, whereas the factor of $(1/\Lambda_2^2)$ in the Kähler potential of the mesons is present since mesons have mass dimension two and have a dynamical origin. The supergravity scalar potential resulting from (3.1) is

$$\begin{aligned} V &= \frac{e^{\text{Tr}(|M|^2/\Lambda_2^2 + |S|^2)}}{(T + \bar{T})^3} \left\{ \frac{(T + \bar{T})^2}{3} \left| \partial_T W - \frac{3}{T + \bar{T}} W \right|^2 + \sum_{ij} |\lambda M_{ij} + \bar{S}_{ij} W|^2 \right. \\ &\quad \left. + \sum_{ij} \left| \lambda S^{ij} + 2X \epsilon^{ijkl} M_{kl} + \frac{\bar{M}^{ij}}{\Lambda_2^2} W \right|^2 + |PfM - \Lambda_2^4|^2 - 3|W|^2 \right\}. \end{aligned} \tag{3.2}$$

In the global limit, the strongly coupled sector breaks supersymmetry, since there is no solution to the supersymmetry eqs. $F^X = F^S = 0$. As explained in [13], the strongly coupled sector produces a contribution to the vacuum energy of order

$$V_0 \sim \lambda^2 \Lambda_2^4. \tag{3.3}$$

Even if at the global supersymmetric level the ground state breaks supersymmetry, similarly to the ISS model discussed in section 2.2, at the supergravity level we do find a supersymmetric AdS minimum. Indeed, by inserting the maximally SO(5) symmetric ansatz

$$\langle M \rangle = \begin{pmatrix} i\sigma_2 & 0 \\ 0 & i\sigma_2 \end{pmatrix} \Lambda_2^2, \quad , \quad \langle S \rangle = c \begin{pmatrix} i\sigma_2 & 0 \\ 0 & i\sigma_2 \end{pmatrix} \Lambda_2^2, \tag{3.4}$$

into the supersymmetry conditions $D_S W = D_M W = D_X W = D_T W = 0$, we find

$$\begin{aligned} \lambda + cW &= 0, \\ \lambda c + 2X + \frac{W}{\Lambda_2^2} &= 0, \\ abe^{-bT_0} + \frac{3}{T_0 + \bar{T}_0} \left(W_0 + ae^{-bT_0} + 4\lambda c \Lambda_2^4 \right) &= 0. \end{aligned} \tag{3.5}$$

If these conditions have a solution, the original supersymmetry breaking ground state becomes metastable. The condition for the uplifting of the vacuum energy in the metastable vacuum requires then $W_0 \sim \lambda \Lambda_2^2$. The last eq. in (3.6) leads then, for $bT_0 \gg 1$, to

$W \sim W_0$ in a first approximation, whereas T_0 is given again by (2.9). A TeV value for the gravitino mass asks therefore for $\Lambda_2^2 \sim m_{3/2} M_{\text{Pl}} \sim (10^{11} \text{ GeV})^2$. Combining the first two eqs. in (3.6), we then find $c \sim -\lambda/W_0$ and therefore $\langle S \rangle \sim M_{\text{Pl}}$. We find, analogously to section 2.2, Planckian values for the supersymmetric AdS vacuum, which means that the supersymmetry preserving vacuum is actually beyond the regime of validity of the effective lagrangian description. In contrast to section 2.2, however, the AdS vacuum energy itself is Planckian here $V_{\text{AdS}} \sim \lambda^2 M_{\text{Pl}}^4$.

By taking seriously this supersymmetric solution, the tunneling from the non-supersymmetric metastable vacuum proceed in the S -field direction in field space. Since $\Delta S \sim M_{\text{Pl}}$, whereas $\Delta V = |V_{\text{AdS}}| \sim \lambda^2 M_{\text{Pl}}^4$, we find for the bounce action $S_b \sim (1/\lambda^2)$. The tunneling probability $\exp(-S_b)$ is therefore suppressed only in the $\lambda \ll 1$ limit. This condition is the analog of the condition $m \ll \Lambda$ in the electric version of the ISS model, i.e. the quarks must have masses much smaller than the dynamical scale of the electric theory.

4. Soft terms and mass scales

4.1 General tree-level formulæ

The relevant couplings for our present discussion are the following terms in the Kähler potential and the superpotential arising from the perturbative expansion in the matter fields M^I

$$\begin{aligned} K &\rightarrow K + [(T + \bar{T})^{n_I} Z_{I\bar{J}} + \dots] M^I \bar{M}^{\bar{J}} + \dots \equiv K + K_{I\bar{J}} M^I \bar{M}^{\bar{J}}, \\ W &\rightarrow W + \frac{1}{6} W_{IJK} M^I M^J M^K, \end{aligned} \tag{4.1}$$

where \dots denote couplings to other (hidden-sector, messengers in gauge mediation models, etc) fields. In a manifestly supersymmetric approach, with both F and D-term contributions, the condition of zero cosmological constant is

$$K_{\alpha\bar{\beta}} F^\alpha F^{\bar{\beta}} + \sum_a (g_a^2/2) D_a^2 = 3m_{3/2}^2 M_{\text{Pl}}^2, \tag{4.2}$$

where $\alpha, \bar{\beta}$ refer to fields contributing to supersymmetry breaking and a is an index for anomalous U(1) gauge factors. Then the most general formulæ for soft terms of matter fields⁶ M^I ($F^I = 0$), are given by [7] (see also [22] for the heterotic strings case)

$$\begin{aligned} m_{I\bar{J}}^2 &= m_{3/2}^2 K_{I\bar{J}} - F^\alpha F^{\bar{\beta}} R_{\alpha\bar{\beta}I\bar{J}} - \sum_a g_a^2 D_a \left(\frac{1}{2} K_{I\bar{J}} - \partial_I \partial_{\bar{J}} \right) D_a, \\ A_{IJK} &= m_{3/2}^2 (3\nabla_I \nabla_J G_K + G^\alpha \nabla_I \nabla_J \nabla_K G_\alpha) - g_a^2 D_a \left(\frac{D_a}{2} \nabla_i \nabla_j G_k - \nabla_i \nabla_j \nabla_k D_a \right), \\ M_{1/2}^a &= \frac{1}{2} (Ref_a)^{-1} m_{3/2} G^\alpha \partial_\alpha f_a, \end{aligned} \tag{4.3}$$

⁶We don't write the analytic bilinear soft terms, since their discussion depends on the origin of the corresponding (μ -like) term in the superpotential.

where $G = K + \ln |W|^2$, $G_\alpha = \partial_\alpha G$, $\nabla_I G_J = G_{IJ} - \Gamma_{IJ}^K G_K$, etc., where $R_{\alpha\bar{\beta}I\bar{J}} = \partial_\alpha \partial_{\bar{\beta}} K_{I\bar{J}} - \Gamma_{\alpha I}^M K_{M\bar{N}} \Gamma_{\bar{\beta} \bar{J}}^{\bar{N}}$ is the Riemann tensor of the Kähler manifold and $\Gamma_{\alpha I}^M = K^{M\bar{N}} \partial_\alpha K_{\bar{N}I}$ are the Christoffel symbols. Moreover,

$$D_a = X_I^a M^I \partial_I K - \frac{\eta_a^\alpha}{2} \partial_\alpha K. \quad (4.4)$$

In (4.4), X_I^a denote $U(1)_a$ charges of charged fields M^I , and η_a^α are defined by the nonlinear gauge transformations of the moduli fields under (super-)gauge fields transformations

$$V_a \rightarrow V_a + \Lambda_a + \bar{\Lambda}_a \quad , \quad T_\alpha \rightarrow T_\alpha + \eta_a^\alpha \Lambda_a. \quad (4.5)$$

By using (4.4), we can also write the scalar masses in (4.3) as

$$m_{I\bar{J}}^2 = m_{3/2}^2 K_{I\bar{J}} - F^\alpha F^{\bar{\beta}} R_{\alpha\bar{\beta}I\bar{J}} - \sum_a g_a^2 D_a \left(\frac{1}{2} D_a - X_I^a - v_l X_l^a \partial_l + \frac{\eta_a^\alpha}{2} \partial_\alpha \right) K_{I\bar{J}}, \quad (4.6)$$

where v_l are vev's of charged scalar fields z^l with charge X_l^a . An interesting question is: in which simple cases do the tree-level contributions of order $m_{3/2}$ in (4.6) cancel each other? This question is particularly relevant in order to identify (classes of) models in which loop contributions and in particular the anomaly-mediated contributions [24] are important.

From a 4d point of view, we are aware of three simple cases:

- i) The well-known case of no-scale models [23], with $K_{T\bar{T}} |F^T|^2 = 3m_{3/2}^2 M_{\text{Pl}}^2$, $D_a = 0$, with matter fields having modular weights $n_I = -1$ in (4.1), when $|F^T|^2 R_{T\bar{T}I\bar{J}} = m_{3/2}^2 K_{I\bar{J}}$. This generalizes easily to the case of several Kähler moduli T_α . Starting from the effective lagrangian

$$K = - \sum_\alpha p_\alpha \ln(T_\alpha + \bar{T}_\alpha) + \prod_\alpha (T_\alpha + \bar{T}_\alpha)^{n_\alpha} |M^I|^2 + \dots, \quad (4.7)$$

the no-scale structure is defined by the condition that the superpotential W is *independent* of T_α and the (semi-)positivity of the scalar potential. Zero cosmological constant then implies

$$K^\alpha K_\alpha \equiv K_{\alpha\bar{\beta}} K^\alpha K^{\bar{\beta}} = 3 \quad \rightarrow \quad \sum_\alpha p_\alpha = 3. \quad (4.8)$$

The condition of having tree-level zero soft scalar masses and A-terms for matter fields M^I is then

$$\sum_\alpha n_\alpha^\alpha = -1. \quad (4.9)$$

- ii) When the following conditions are simultaneously satisfied:

- D-term contributions are much larger⁷ than the F-terms and cancel the cosmological constant $\sum_a (g_a^2/2) D_a^2 \simeq 3m_{3/2}^2$.

⁷We should keep in mind, however, that in supergravity with $\langle W \rangle \neq 0$, there is no pure D-breaking. This case assumes therefore $D_a \gg F^\alpha$, but F-terms have to exist.

- there are no (large) vev's of charged scalar fields $v_l = 0$.
- the matter fields are neutral under the U(1)'s symmetries and come from the D3 brane sector (or, more generally $n_I = -1$).

Indeed, in this case by using the Kähler potential

$$K = -3 \ln(T + \bar{T}) + (T + \bar{T})^{-1} |M^I|^2 + \dots, \tag{4.10}$$

then it can be easily checked that the D-term contributions precisely cancel the other terms in the soft terms in (4.3). The generalization of this D-dominated supersymmetry breaking case to the case of several moduli T_α is more involved and will not be discussed here.

- iii) A simple way to obtain tree-level zero soft masses is by geometric sequestering [24], i.e separating in the internal space the source of supersymmetry breaking from the matter fields. From a 4d viewpoint, the vanishing of the tree-level soft terms appear as non-trivial cancellations in the general formula (4.3). However this cancellation is protected from quantum corrections by the geometric separation of the source of supersymmetry breaking. A typical example, obtained by assuming that moduli fields (in particular the modulus T) were stabilized in a supersymmetric way, is that of a matter field M and a hidden sector field ϕ_h , which is the only source of supersymmetry breaking and of cancellation of the cosmological constant $G_h G^h = 3$. The 4d supergravity action is

$$\begin{aligned} K &= -3 \ln\left(1 - \frac{|M|^2}{3} - \frac{|\phi_h|^2}{3}\right), \\ W &= W_v(M) + W_h(\phi_h). \end{aligned} \tag{4.11}$$

It is also possible that a matter-like field C with couplings to the observable matter saturates the vacuum energy $K_{C\bar{C}} |F^C|^2 = 3m_{3/2}^2 M_{\text{Pl}}^2$ and by fine-tuning provides the cancellation of the tree-level soft scalar mass, see e.g. [12]. When neither of these cases occur, other manifestly supersymmetric uplifting mechanism are expected to lead to soft scalar masses of the order of the gravitino mass $m_{I\bar{J}}^2 \sim m_{3/2}^2$.

4.2 Soft terms with dynamical F-term uplifting

A particularly important question is the magnitude of the soft terms in the visible sector in the present setup. In order to answer this question, we first estimate the contribution to supersymmetry breaking from the various fields. By using the results of section 2, we find in the leading order

$$\begin{aligned} \overline{F^\varphi} &\equiv e^{K/2} K^{\varphi\bar{\varphi}} D_{\bar{\varphi}} W \simeq e^{K/2} K^{\varphi\bar{\varphi}} (\bar{\varphi}_0 W + \delta\Phi \partial_\Phi \partial_{\bar{\varphi}} W_2) \simeq 0, \\ \overline{F^{\bar{\varphi}}} &\simeq 0, \\ \overline{F^\Phi} &= e^{K/2} \begin{pmatrix} 0 & 0 \\ 0 & -h\mu^2 I_{N_f - N} \end{pmatrix}, \\ F^T &\simeq \frac{a}{(T_0 + \bar{T}_0)^{1/2}} e^{-bT_0} \simeq -\frac{3}{b} m_{3/2}. \end{aligned} \tag{4.12}$$

Notice that the main contribution to supersymmetry breaking comes from the magnetic mesonic fields Φ , which are the main responsible for the uplift of the vacuum energy

$$\text{Tr}(|F^\Phi|^2) \simeq 3m_{3/2}^2. \tag{4.13}$$

The transmission of supersymmetry breaking in the observable sector depends on the couplings of the observable fields M^I to the SUSY breaking fields Φ, T . The relevant couplings for our present discussion are the following terms in the Kähler metric of the matter fields M^I

$$K_{I\bar{J}} = (T + \bar{T})^{n_I} Z_{I\bar{J}} + \text{Tr}(|\Phi|^2) Z'_{I\bar{J}}, \tag{4.14}$$

where the form of the Φ coupling in the Kähler metric is dictated by the diagonal $SU(N_f)$ flavor symmetry left unbroken by the mass parameter μ in the ISS lagrangian. The Yukawa couplings W_{IJK} could also depend on T and Φ .

Then from (4.3) with no D-term contributions $D_a = 0$, we find that the F^T contribution is subleading by a factor $1/b^2(T + \bar{T})^2$ with respect to the other contributions. This has the nice feature that the flavor-dependent F^T contribution to scalar soft masses are subleading. The result for the (canonically normalized scalars) soft masses, at the leading order, is then given by

$$\begin{aligned} m_{I\bar{J}}^2 &= m_{3/2}^2 \delta_{I\bar{J}} + \frac{h^2(N_f - N)\mu^4}{(T + \bar{T})^3} (K^{-1} Z')_{I\bar{J}} \\ &\simeq m_{3/2}^2 (\delta_{I\bar{J}} + 3(K^{-1} Z')_{I\bar{J}}). \end{aligned} \tag{4.15}$$

If the coupling to the mesonic fields Φ is small, i.e the coefficients $Z'_{I\bar{J}}$ are suppressed, soft scalar masses in the observable (MSSM) sector are universal and are similar with the ones obtained in the “dilaton-dominated” scenario in the past. It would be very interesting to find physical reasons of why $Z'_{I\bar{J}}$ are small. The geometrical sequestering cannot be invoked in this case since the matter fields M and the mesons Φ do not fit into the structure (4.11). If the coeff. $Z'_{I\bar{J}}$ are of order one, the two terms in (4.15) are of the same order and the flavor problem of gravity mediation is back.

A similar conclusion holds for the other possible source of flavor violation, the A-terms. If the couplings of the mesons to the matter fields are small, we get at the leading order, for the canonically normalized scalars

$$A_{IJL} \simeq 3m_{3/2} w_{IJL}, \tag{4.16}$$

where w_{IJL} are the low-energy Yukawa couplings for the matter fields, related to the corresponding SUGRA couplings $W_{IJL} = \nabla_I \nabla_J \nabla_L W$ by

$$w_{IJL} = e^{K/2} (K^{-1/2})'_I (K^{-1/2})'_J (K^{-1/2})'_L W_{I'J'L'}. \tag{4.17}$$

Since A-terms are proportional to the Yukawa couplings, there are no flavor violations in this case.

Gaugino masses in the observable sector are determined by the gauge kinetic functions which in our case have generically the form

$$f_a = f_a^{(0)} + \alpha_a T + \beta_a (\text{Tr}\Phi), \tag{4.18}$$

where $f_a^{(0)}$ are provided by other moduli fields, stabilized in a supersymmetric manner. The form of coupling to the mesons in (4.18) is fixed by the diagonal $SU(N_f)$ flavor symmetry left unbroken by the mass parameter μ , whereas α_a are numbers of order one.⁸ The gaugino masses

$$M_a = \alpha_a F^T + \beta_a (\text{Tr} F^\Phi) \tag{4.19}$$

are of the order of the gravitino mass if β_a are of order one, whereas they are suppressed by the factor $1/b(T + \bar{T})$ if β_a are small. In this second case, the anomaly-mediated contributions [24, 25] are comparable to the tree-level ones. To conclude, we do not find a suppression of all the soft terms in the observable sector with respect to the gravitino mass. Therefore our results point towards a gravity-mediation type of supersymmetry breaking in the hidden sector, which in the case of small couplings of matter to hidden sector mesons are very similar to the dilaton-domination scenario and are therefore flavor blind at tree-level.⁹

We would like to briefly compare these results to the ones obtained in [14] by using the original KKLT uplifting mechanism with $\overline{D3}$ antibranes.¹⁰ By using a nonlinear supergravity approach, [14] found a (moderate) hierarchy $m_{3/2} \sim 4\pi^2 m_{\text{soft}}$. Let us try to understand better the difference with our results. As we discussed in the previous section, there are three ways of suppressing the tree-level soft masses for matter fields. The first is no-scale type models. The KKLT-type models are not of this type, since the F^T contribution is small. The second case is the dominant D-term breaking. This is probably the manifestly supersymmetric case which should correspond in the low energy limit to the analysis done in [14]. Knowing that pure D-term supersymmetry breaking does not exist, it could be difficult to realize a model along these lines. It is however very interesting to investigate this possibility in more detail.

We believe that a more detailed phenomenological analysis of the possible manifestly supersymmetric uplifting mechanisms deserves further investigation.

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⁸In a type IIB orientifold embedding, this happens if the observable sector lives on D7 branes.

⁹For other ways of getting flavor universality in compactifications with stabilized moduli, see e.g. [27].

¹⁰See also [26] for a model with a phenomenology similar to the one in [14].

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