Dynamically sequestered F-term uplifting in extra dimension

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ABSTRACT: We study moduli stabilization, the dynamical supersymmetry (SUSY) breaking, the uplifting of SUSY anti-de Sitter (AdS) vacuum and the sequestering of hidden sector in a five-dimensional supergravity model, where all modes of the visible sector and the hidden sector are originated from bulk fields. We clarify couplings between the visible and hidden sectors. The expressions for the visible sector soft SUSY breaking terms as well as the hidden sector potential are shown explicitly in our model. The sequestering is achieved dynamically by a wavefunction localization in extra dimension. We find that the tree-level soft scalar mass and the A-term can be suppressed at a SUSY breaking Minkowski minimum where the radius modulus is stabilized, while gaugino masses would be a mirage type.

KEYWORDS: Field Theories in Higher Dimensions, Supersymmetry Breaking,
Supersymmetric Effective Theories, Supergravity Models.

Contents

1. Introduction

Supersymmetric extensions of the standard model (SM) are promising candidates for the physics around a TeV scale. The supersymmetry (SUSY) can protect the electroweak scale $M_{EW} \sim 10^2$ GeV against radiative corrections of the order of the Planck scale $M_{Pl} \sim 10^{19}$ GeV, even with soft SUSY breaking if the scale of soft SUSY breaking terms are of $\mathcal{O}(M_{EW})$. That means, the unification of the gravity and the SM would be possible within such framework without violating the mass scale of the SM. Moreover, the minimal supersymmetric standard model (MSSM) suggests that the three gauge couplings in the SM are unified at the grand unification theory (GUT) scale, $M_{\mathrm{GUT}} \sim 10^{16}$ GeV. The lightest SUSY particle (LSP) can be a candidate for the cold dark matter, if anything like R-parity forbids LSP decays.

Because any SUSY particles have not been observed yet, the SUSY must be broken above the electroweak scale. The SUSY breaking generically introduces flavor violating masses, mixing and couplings (soft SUSY breaking terms), which are severely restricted by the flavor changing neutral current experiments. If the SUSY breaking effects are dominantly mediated from the SUSY breaking (hidden) sector to the visible (MSSM) sector by gauge interactions which are flavor blind, SUSY flavor violations can be suppressed. However, if the dominant SUSY breaking effects are mediated by gravitational interactions which are generically flavor dependent, there must exist certain mechanism to suppress flavor violations in the visible sector, which is sometimes called sequestering. The ultimate

situation for the sequestering is the case that the hidden sector is completely separated from the visible sector by the spatiality of extra dimensions [1]. In this case, the soft SUSY breaking terms in the visible sector are generated through the superconformal anomaly [1, 2]. However, due to the ultra-violet (UV) insensitive nature of anomaly mediation, the low energy behavior of MSSM soft terms is completely determined, resulting in tachyonic slepton masses.

On the other hand, if we consider superstring theory and its effective supergravity theory as a UV completion of the MSSM (or any SUSY SM), there generically exist moduli superfields in the four-dimensional (4D) effective theory. Vacuum expectation values (VEVs) of moduli correspond to sizes and shapes of extra dimensions and determine quantities in the 4D effective theory such as M_{Pl}, gauge and Yukawa couplings. Naively, the moduli are flat directions of the potential, but must be stabilized by some nontrivial effects such as fluxes and/or nonperturbative effects. The moduli stabilization is also relevant to soft SUSY breaking terms because moduli multiplets generically couple to the visible sector in the effective theory, and the auxiliary components of the moduli multiplets are also determined by the potential stabilizing the moduli themselves in supergravity.

It has been recognized that the realization of a SUSY breaking vacuum with an (almost) vanishing vacuum energy, which is required by the observations, is quite difficult in the conventional moduli stabilization schemes in supergravity theory. This is mainly because the moduli potential in the effective supergravity theory, in general, prefers SUSY preserving anti-de Sitter (AdS) vacua with a negative vacuum energy. Recently, a systematic way for realizing a SUSY breaking Minkowski minimum in a controllable manner was proposed, which we call the Kachru-Kallosh-Linde-Trivedi (KKLT) scenario [3]. In this scenario, we uplift the above mentioned SUSY AdS vacuum to a Minkowski minimum by a SUSY breaking vacuum energy generated in the so-called uplifting sector, which is assumed to be well sequestered from the light moduli as well as the visible sector. In the original KKLT scenario, the uplifting sector is formed by an anti D3-brane located at the tip of the warped throat in extra dimensions.

Because of the geometrical structure, the SUSY breaking anti-brane can be sequestered from the light moduli and the visible sector (on D7 branes). Soft SUSY breaking terms are calculated [4, 5] in the KKLT model. It was found that the tree level (light) modulus mediation is generically comparable to the anomaly mediation, resulting in the so-called mirage mediation [6] and leading to phenomenologically interesting aspects [7, 8]. Then, this model could be restricted from the viewpoint of flavor violation, if the light-modulus couplings to the visible sector are flavor dependent. The flavor dependence of the modulus couplings might be determined by the mechanism of generating Yukawa hierarchies for quarks and leptons. However, it is still a challenging issue to realize successful Yukawa structures within superstring models.

The anti-brane breaks SUSY explicitly in the effective supergravity theory. We can modify the KKLT scenario such that the uplifting energy is supplied by nonvanishing F-terms [9 – 12] and/or D-terms [13, 14] in the dynamical SUSY breaking (hidden) sector[1]

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[1]Note that we call the dominant SUSY breaking sector as the hidden sector, but the moduli sector is

such as the O'Raifeartaigh model [15] and the Intriligator-Seiberg-Shih (ISS) model [16].[2] The former is called as the F-term uplifting, which can realize a low energy SUSY models more easily than the latter. In this scenario, the soft terms are generically model dependent (see, e.g., ref. [17] and references therein), due to the possible direct couplings between the hidden and the visible sectors. If contact terms between the visible and hidden sectors are not suppressed, soft SUSY breaking scalar masses in the visible sector would be of $\mathcal{O}(m_{3/2})$. On the other hand, scalar masses would be suppressed in the case that the hidden sector is somehow sequestered from the visible sector. At any rate, these scalar masses depend on explicit forms of couplings between the hidden and visible sectors. Thus, it is important to clarify explicitly couplings between the visible and hidden sectors in a certain class of models and to study the moduli stabilization, the F-term uplifting scenario and soft SUSY breaking terms.

It would be difficult to study the above aspects explicitly in effective supergravity theories of superstring models on six extra dimensions with generic geometrical structure, although our naive motivation is originated from superstring theory. Obviously, a single extra dimensional model is the simplest in extra dimensional models, but it has a sufficiently rich structure, which would be useful to understand the above aspects in string-derived supergravity theories. Thus, in this paper we study five-dimensional (5D) supergravity on the orbifold S^1/Z_2 with a weakly warped factor. Indeed, the modulus stabilization and the uplifting of the vacuum energy in the 5D orbifold model have been studied in ref. [18]. Its setup is that the visible sector and the hidden sector are localized on different orbifold fixed points. That is, the SUSY breaking source and the uplifting sector are sequestered geometrically from the visible sector, and at the tree level, the sequestering in this setup is complete as it may be rather obvious from the geometrical viewpoint. In this paper, we study another setup that all modes are bulk fields. We also introduce kink-type masses for bulk modes in order to consider non-trivial wavefunction profiles of zero modes along the fifth dimension [19]. As a solution of the equation of motion, the zero mode wavefunction is quasi-localized exponentially at one of orbifold fixed points.[3] This setup includes the setup of ref. [18] as the limit that the hidden sector and the visible matter fields are localized sharply at different fixed points.[4] Similarly, as a certain limit our setup includes the setup that one of the visible and hidden sectors is localized at a orbifold fixed point and the other is originated from bulk modes. Because both the visible sector and the hidden sector are bulk modes in our setup, couplings between the hidden and visible sectors are constrained by the bulk SUSY and can be calculated up to brane-induced couplings. We clarify these couplings and study the moduli stabilization, the SUSY breaking and the F-term uplifting in our 5D supergravity model.

The following sections are organized as follows. In section 2, we construct our 5D su-

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not included in the hidden sector in our terminology.

[2]These models have the same behavior when heavy modes are integrated out.

[3]Such a wavefunction localization mechanism is useful to realize realistic Yukawa matrices for quarks and leptons. See, e.g., refs. [20, 21] and references therein.

[4]For the visible gauge sector, the brane-induced gauge kinetic function must be quite large compared with the bulk gauge kinetic function to realize the setup of ref. [18] from our model.

pergravity model. We also derive the 4D effective Lagrangian for our model. In section 3, we compute the Kähler potential and the superpotential for the radius modulus T and the hidden sector field X, which are relevant to the dynamical SUSY breaking and the uplifting. We also show the resulting soft SUSY breaking terms in terms of the auxiliary components F^T and F^X. Then, we discuss about the nature of SUSY breaking, F-term uplifting and sequestering in our model based on them. We summarize this paper in section 4. In appendix A, the basic structure of F-term uplifting is briefly reviewed as well as the original KKLT model.

2. Quasi-localized visible and hidden sectors

2.1 5D model

We consider a 5D supergravity compactified on an orbifold as an illustrating example of dynamical sequestering, SUSY breaking and F-term uplifting. For such a purpose, we start from the 5D off-shell (conformal) supergravity on an orbifold S^1/Z_2 [22] with the radius R. For simplicity and concreteness, we choose the simplest bulk supergravity characterized by a single Z_2-odd vector multiplet (graviphoton) and a single compensator hypermultiplet in addition to matter multiplets. Because of that, we have a single (radius) modulus T in the 4D effective theory, whose VEV is related to the orbifold radius as $\langle T \rangle = \pi R$, and the target manifold of the hypermultiplet is characterized by $USp(2, 2n_H)/USp(2) \times USp(2n_H)$, where n_H is the number of the matter hypermultiplets. The extensions to the case with more odd vector multiplets (i.e., more moduli) and/or with more compensators are straightforward.[5]

As for matter multiplets, we embed visible sector (MSSM) chiral multiplets Q^I into 5D hypermultiplets $\mathcal{H}^I = (\mathcal{Q}^I, \mathcal{Q}^I_c)$, where Q^I is the zero-mode of $\mathcal{Q}^I$, and the index I runs over all the quarks, leptons and Higgs fields. The hidden sector chiral multiplet X, which is responsible for the dynamical SUSY breaking and for the uplifting, is also assigned to a 5D hypermultiplet $\mathcal{H}^X = (\mathcal{X}, \mathcal{X}_c)$ where X is the zero-mode of $\mathcal{X}$. Zero modes of chiral multiplet partners $\mathcal{Q}^I_c$ and $\mathcal{X}_c$ in these hypermultiplets are projected out by orbifolding.

We gauge U(1) isometries of the compensator, visible (MSSM) and hidden hypermultiplets by the graviphoton with corresponding charges k, c_I and c_X, respectively [19, 23, 24]. Because the graviphoton has an odd Z_2-parity under the orbifold projection, the gauge coupling should change its sign across the orbifold fixed points located at $y = 0, \pi R$. This can be achieved by accompanying a periodic sign function $\epsilon(y)$ with the gauge coupling.[6] Then, these gauging procedures generate kink-type masses $\epsilon(y)\hat{k}$, $\epsilon(y)\hat{c}_I$ and $\epsilon(y)\hat{c}_X$ for the gravitino, visible and hidden hyperinos, respectively, as well as a bulk negative cosmological constant scaled by $\hat{k}$, where $\hat{k} = kM$, $\hat{c}_I = c_I M$, $\hat{c}_X = c_X M$ and M is the VEV of graviscalar (a scalar component of graviphoton multiplet). Thus, the equations of motion in the orbifold segment generate exponential profiles $e^{-\hat{k}y}$, $e^{-\hat{c}_I y}$ and $e^{-\hat{c}_X y}$ for the wave-

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[5]For example, the power 2/3 of the integrand in eq. (2.1) is different in the case with two compensators, from the viewpoint of the 4D effective theory.

[6]In supergravity, this can be achieved by the so-called four-form mechanism [25, 24].

functions [19] of the gravitino, visible and hidden hyperinos, respectively, in the orbifold slice of AdS$_5$ warped background geometry[7] $ds^2 = e^{-2\hat{k}y}dx^2 - dy^2$.

The bulk supergravity has an $N = 2$ SUSY (eight supercharges), and does not allow Yukawa interactions between hypermultiplets. To be phenomenologically viable, we introduce superpotential terms at the fixed point including the Yukawa interaction between quarks/leptons and the Higgs bosons. In addition, some superpotential terms for the hidden sector field X are necessary for triggering a dynamical SUSY breaking, generating a nonvanishing value of F^X, which would uplift the negative vacuum energy of the modulus sector. Actually, the fixed points respect only $N = 1$ SUSY (four supercharges) which survives under the orbifolding, where we can write any Kähler/superpotential terms and gauge kinetic functions if the other symmetries such as MSSM gauge symmetries do not forbid them.

2.2 4D effective theory

To study the nature of SUSY breaking, uplifting and sequestering, a 4D effective theory of our 5D model is desired. Especially, the information of contact interactions between Q^I and X in the Kähler potential is important in such 4D effective theory. For this end, we adopt the off-shell dimensional reduction method proposed by refs. [26, 27], which is based on an $N = 1$ superspace description of 5D conformal supergravity on an orbifold [28] and developed in subsequent works [29]. This method provides us a way for deriving the 4D off-shell effective action directly from the 5D off-shell supergravity action with generic boundary terms, respecting the $N = 1$ off-shell structure. The procedure is as follows. We start from the $N = 1$ off-shell description of 5D action. After some gauge transformation, we drop kinetic terms for Z_2-odd multiplets which are negligible at low energy. Then, these multiplets play a role of Lagrange multiplier and their equations of motion extract zero-modes from the Z_2-even multiplets.

After these steps, we find the 4D effective Lagrangian of our 5D model in the $N = 1$ superspace description,

$$\mathcal{L} = \int d^4\theta \, |C|^2 \Omega + \left\{ \int d^2\theta \, \left(f_a W^{a\alpha} W^a_\alpha + C^3 W \right) + \text{h.c.} \right\},$$

where C is the compensator chiral multiplet in the 4D $N = 1$ supergravity. The Kähler potential $K = -3\ln(-\Omega/3)$, the superpotential W and the gauge kinetic function f_a are given by

$$\Omega = \Omega^{(\text{bulk})} + \sum_{\vartheta=0,\pi} e^{-\frac{\vartheta}{\pi}k(T+\bar{T})}\Omega^{(\vartheta)} = -3e^{-K/3},$$

$$W = \sum_{\vartheta=0,\pi} e^{-3\frac{\vartheta}{\pi}kT}W^{(\vartheta)} + W^{(\text{np})},$$

$$f_a = f_a^{(\text{bulk})} + \sum_{\vartheta=0,\pi} f_a^{(\vartheta)},$$

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[7]The boundary tension terms which balance with the bulk cosmological constant is automatically supplied [24] by the four-form mechanism mentioned in the previous footnote.

and contributions from the bulk are found as

$$\Omega^{(\mathrm{bulk})} = -3 \int_0^{\mathrm{Re}\,T} dt\, e^{-2kt} \left(1 - e^{-2\left(c_X - \frac{3}{2}k\right)t}|X|^2 - \sum_I e^{-2\left(c_I - \frac{3}{2}k\right)t}|Q^I|^2 \right)^{\frac{2}{3}}, \quad (2.1)$$

$$f_a^{(\mathrm{bulk})} = k_a T.$$

The Yang-Mills gauge couplings to the hypermultiplets in $\Omega^{(\mathrm{bulk})}$ are omitted to simplify the expression, which are irrelevant to the following arguments. The nonperturbative effects such as gaugino condensations are encoded in the superpotential terms

$$W^{(\mathrm{np})} = \sum_{m=1}^{m_{\mathrm{np}}} B_m e^{-b_m} - \sum_{n=1}^{n_{\mathrm{np}}} A_n e^{-a_n T},$$

where the first constant and the second T-dependent term would come from the m_{np} boundary and the n_{np} bulk (zero-mode) gaugino condensations, respectively. Thus the natural orders of the constants are

$$b_m \sim a_n = \mathcal{O}(4\pi^2), \qquad B_m \sim A_n = \mathcal{O}(1).$$

The other quantities

$$\Omega^{(\vartheta)} = \Omega^{(\vartheta)}(X_\vartheta, \bar{X}_\vartheta, Q_\vartheta^I, \bar{Q}_\vartheta^I), \quad W^{(\vartheta)} = W^{(\vartheta)}(X_\vartheta, Q_\vartheta^I), \quad f_a^{(\vartheta)} = f_a^{(\vartheta)}(X_\vartheta),$$

originate from the Kähler potential, the superpotential and the gauge kinetic function, respectively, induced at the 4D fixed point $y = \vartheta R$ ($\vartheta = 0, \pi$). Within the framework of 5D orbifold supergravity, these can be arbitrary functions of the chiral multiplets X_ϑ and/or Q_ϑ^I originating from the bulk hypermultiplets (as well as the boundary own fields which are treated implicitly here and hereafter, if they are needed),

$$X_\vartheta = e^{-\frac{\vartheta}{\pi}\left(c_X - \frac{3}{2}k\right)T} X, \qquad Q_\vartheta^I = e^{-\frac{\vartheta}{\pi}\left(c_I - \frac{3}{2}k\right)T} Q^I,$$

for $\vartheta = 0$ and $\vartheta = \pi$, respectively. In order for the orbifold supergravity theory to be self consistent, we consider that all the boundary terms can be treated perturbatively, i.e., the constants in $\Omega^{(\vartheta)}$, $W^{(\vartheta)}$ and $f_a^{(\vartheta)}$ are assumed implicitly to be small compared with those originating from the bulk in this paper.

We consider the 4D effective theory around $Q^I = X = 0$, and then expand Ω in powers of Q^I and X as

$$\Omega = \hat{\Omega}_0(T, \bar{T}) + Y_{X\bar{X}}(T, \bar{T})|X|^2 + Y_{I\bar{J}}(T, \bar{T}, X, \bar{X})Q^I \bar{Q}^{\bar{J}} + \mathcal{O}(Q^4, X^4),$$

where

$$\hat{\Omega}_0(T, \bar{T}) = -3\alpha\, \frac{1 - \beta|e^{-kT}|^2}{2k}, \qquad Y_{X\bar{X}}(T, \bar{T}) = \alpha_X\, \frac{1 - \beta_X|e^{-(c_X - k/2)T}|^2}{c_X - k/2},$$

$$Y_{I\bar{I}}(T, \bar{T}, X, \bar{X}) = \alpha_I\, \frac{1 - \beta_I|e^{-(c_I - k/2)T}|^2}{c_I - k/2} + \frac{1}{3}\tilde{\alpha}_I\, \frac{1 - \tilde{\beta}_I|e^{-(c_I + c_X - 2k)T}|^2}{c_I + c_X - 2k}|X|^2, \quad (2.2)$$

and the coefficients are determined as

$$\alpha = 1 - \frac{2}{3}k\Omega^{(0)}|_0, \qquad\qquad \alpha\beta = 1 + \frac{2}{3}k\Omega^{(\pi)}|_0,$$

$$\alpha_X = 1 + \left(c_X - \frac{k}{2}\right)\Omega^{(0)}_{X\bar{X}}|_0, \qquad\qquad \alpha_X\beta_X = 1 - \left(c_X - \frac{k}{2}\right)\Omega^{(\pi)}_{X\bar{X}}|_0,$$

$$\alpha_I = 1 + \left(c_I - \frac{k}{2}\right)\Omega^{(0)}_{I\bar{I}}|_0, \qquad\qquad \alpha_I\beta_I = 1 - \left(c_I - \frac{k}{2}\right)\Omega^{(\pi)}_{I\bar{I}}|_0,$$

$$\tilde{\alpha}_I = 1 + 3(c_I + c_X - 2k)\Omega^{(0)}_{I\bar{I}X\bar{X}}|_0, \qquad \tilde{\alpha}_I\tilde{\beta}_I = 1 - 3(c_I + c_X - 2k)\Omega^{(\pi)}_{I\bar{I}X\bar{X}}|_0.$$

Here and hereafter, we use the notation that $F_{AB\cdots} = \partial_A\partial_B\cdots F$ and $F_{AB\cdots}|_0 = F_{AB\cdots}|_{Q^I=\bar{Q}^{\bar{I}}=X=\bar{X}=0}$ for a function $F = F(X, \bar{X}, Q^I, \bar{Q}^{\bar{I}})$ and indices $A, B, \ldots = (X, \bar{X}, I, \bar{I})$. In this paper, we assume that the boundary Kähler potential does not contain flavor mixing, i.e., $\Omega^{(\vartheta)}_{I\bar{J}}|_{Q^I=\bar{Q}^{\bar{I}}=0} = 0$ for $J \neq I$, for simplicity.[8] Note that there is no flavor mixing in the bulk Kähler potential. Then we obtain

$$Y_{I\bar{J}}(T, \bar{T}, X, \bar{X}) = 0 \qquad (J \neq I). \tag{2.3}$$

Moreover, when $\Omega^{(0),(\pi)}$ and their derivatives are sufficiently small, we have

$$\alpha = \alpha_X = \alpha_I = \tilde{\alpha}_I = \beta = \beta_X = \beta_I = \tilde{\beta}_I = 1.$$

The Kähler potential is calculated from Ω as

$$K = -3\ln\left(-\Omega/3\right) = \hat{K}(T, \bar{T}, X, \bar{X}) + \mathcal{O}(Q^2),$$

where

$$\hat{K}(T, \bar{T}, X, \bar{X}) = \hat{K}_0(T, \bar{T}) + Z_{X\bar{X}}(T, \bar{T})|X|^2 + \mathcal{O}(|X|^4), \tag{2.4}$$

and

$$\hat{K}_0(T, \bar{T}) = -3\ln\left(-\hat{\Omega}_0(T, \bar{T})/3\right),$$

$$Z_{X\bar{X}}(T, \bar{T}) = -3Y_{X\bar{X}}(T, \bar{T})/\hat{\Omega}_0(T, \bar{T}).$$

Similarly, we also expand the superpotential W and the gauge kinetic function f_a as

$$W = \hat{W}(T, X) + \frac{1}{6}\lambda_{IJK}(T, X)Q^I Q^J Q^K + \mathcal{O}(Q^4),$$

$$f_a = f_a^{(0)}|_0 + f_a^{(\pi)}|_0 + k_a T + \left(f_{a,X}^{(0)}|_0 + f_{a,X}^{(\pi)}|_0\, e^{-(c_X - 3k/2)T}\right)X + \mathcal{O}(X^2), \tag{2.5}$$

where

$$\hat{W}(T, X) = c - \sum_{n=0}^{n_{\mathrm{np}}} A_n e^{-a_n T} + \left(W_X^{(0)}|_0 + W_X^{(\pi)}|_0\, e^{-(c_X + 3k/2)T}\right)X + \mathcal{O}(X^2), \tag{2.6}$$

$$\lambda_{IJK}(T, X) = W_{IJK}^{(0)}|_0 + W_{IJK}^{(\pi)}|_0\, e^{-(c_I + c_J + c_K - 3k/2)T}$$
$$+ \left(W_{IJK,X}^{(0)}|_0 + W_{IJK,X}^{(\pi)}|_0\, e^{-(c_I + c_J + c_K + c_X - 3k)T}\right)X + \mathcal{O}(X^2). \tag{2.7}$$

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[8]We would study a more general case with $\Omega^{(\vartheta)}_{I\bar{J}}|_{Q^I=\bar{Q}^{\bar{I}}=0} \neq 0$ for $J \neq I$ in a separate work [30].

Note that the constant terms in $W^{(\vartheta)}$ are now encoded in the constant c and the zeroth component of A_n as

$$
c = W^{(0)}|_0 + \sum_{m=1}^{m_{\rm np}} B_m e^{-b_m},
$$

$$
A_n = \{W^{(\pi)}|_0, A_1, A_2, \ldots, A_{n_{\rm np}}\}, \qquad a_n = \{3k, a_1, a_2, \ldots, a_{n_{\rm np}}\},
$$

where $n = 0, 1, 2, \ldots, n_{\rm np}$.

In the above expressions, we find the Kähler/superpotential terms for the modulus T and the hidden sector field X, which carry the informations of dynamical SUSY breaking and the uplifting structures. We can also compute soft supersymmetry breaking terms from the above expressions, in terms of the F-component of the radius modulus T and the hidden sector field X. With the resulting soft terms, we can analyze the nature of sequestering in our 5D model.

3. Hidden sector potential and soft terms

In order to obtain a SUSY breaking Minkowski vacuum with (almost) vanishing vacuum energy, we consider the scenario of F-term uplifting, which would be realized by the scalar potential of the modulus and the hidden sector:

$$
V_F = \hat{K}_{i\bar{j}} F^i \bar{F}^{\bar{j}} - 3|m_{3/2}|^2,
$$

where $i, j, \ldots = (T, X)$,

$$
F^i = -e^{\hat{K}/2} \hat{K}^{\bar{j}i}(\bar{\hat{W}}_{\bar{j}} + \hat{K}_{\bar{j}}\bar{\hat{W}}), \qquad m_{3/2} = e^{\hat{K}/2}\hat{W},
$$

and $\hat{K}_{i\bar{j}}\hat{K}^{\bar{j}k} = \delta_i{}^k$, $\hat{K}^{\bar{i}j}\hat{K}_{j\bar{k}} = \delta^{\bar{i}}{}_{\bar{k}}$. By substituting $\hat{K}(T, \bar{T}, X, \bar{X})$ and $\hat{W}(T, X)$ shown in eqs. (2.4) and (2.6), respectively, we obtain the modulus and the hidden sector F-term scalar potential of our model. Note that the above potential is evaluated in the Einstein frame where $|C|^2 = e^{K/3}M_{Pl}^2$, and we measure all the mass scales in the unit $M_{Pl} = 1$ in the following. We also restrict ourselves to the region $X \ll 1$ where the expansion in powers of X is valid.

As for the tree-level soft SUSY breaking terms of the visible fields, in this paper we focus on[9] the gaugino masses M_a, the scalar masses $m_{I\bar{J}}^2$ and the A-terms A_{IJK}. These are defined as

$$
\mathcal{L}_{\rm soft} = -m_I^2 |Q^I|^2 - \frac{1}{2}\left(M_a \lambda^a \lambda^a + \frac{1}{6} y_{IJK} A_{IJK} Q^I Q^J Q^K + {\rm h.c.}\right),
$$

where *all the kinetic terms are canonically normalized*, and $y_{IJK} = Y_{I\bar{I}}^{-1/2} Y_{J\bar{J}}^{-1/2} Y_{K\bar{K}}^{-1/2} \lambda_{IJK}$ is the physical Yukawa coupling. Note that there is no flavor mixing in the soft scalar masses $m_{I\bar{J}}^2 = \delta_{IJ} m_I^2$ due to eq. (2.3). These soft terms are generated through the mediation by

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[9]The so-called μ-term and B-term would be discussed in a separate work [30].

the radius modulus T as well as the direct couplings to the SUSY breaking field X. Such effects are summarized in the following general formula [31, 5]:

$$M_a = F^i \partial_i \ln(\text{Re} f_a),$$
$$m_I^2 = -F^i \bar{F}^{\bar{j}} \partial_i \partial_{\bar{j}} \ln Y_{I\bar{I}},$$
$$A_{IJK} = F^i \partial_i \ln \left(Y_{I\bar{I}}^{-1} Y_{J\bar{J}}^{-1} Y_{K\bar{K}}^{-1} \lambda_{IJK} \right),$$

where $i, j, \ldots = (T, X)$. Here we assume that the total vacuum energy is vanishing at the minimum where the soft terms are evaluated, which would be realized by the F-term uplifting.

Substituting f_a, $Y_{I\bar{I}}$ and λ_{IJK} shown in eqs. (2.5), (2.2) and (2.7), respectively, we find the expressions for the above soft terms as

$$M_a = \frac{F^T + (2/k_a)(f_{a,X}^{(0)} + f_{a,X}^{(\pi)} e^{-(c_X - 3k/2)T}) F^X}{T + \bar{T} + (2/k_a)(f_a^{(0)} + f_a^{(\pi)})} + \mathcal{O}(X),$$

$$m_I^2 = \beta_I (c_I - k/2)^2 \frac{|e^{(c_I - k/2)T}|^2}{(|e^{(c_I - k/2)T}|^2 - \beta_I)^2} |F^T|^2$$
$$- \frac{\tilde{\alpha}_I}{3\alpha_I} \frac{c_I - k/2}{c_I + c_X - 2k} \frac{|e^{(c_I - k/2)T}|^2 - \tilde{\beta}_I |e^{-(c_X - 3k/2)T}|^2}{|e^{(c_I - k/2)T}|^2 - \beta_I} |F^X|^2 + \mathcal{O}(X),$$

$$A_{IJK} = -\left\{ \left(\frac{\beta_I (c_I - k/2)}{|e^{(c_I - k/2)T}|^2 - \beta_I} + (I \leftrightarrow J) + (I \leftrightarrow K) \right) \right.$$
$$\left. + \frac{c_I + c_J + c_K - 3k/2}{W_{IJK}^{(0)} e^{(c_I + c_J + c_K - 3k/2)T} + W_{IJK}^{(\pi)}} \right\} F^T$$
$$+ \frac{W_{IJK,X}^{(0)} + W_{IJK,X}^{(\pi)} e^{-(c_I + c_J + c_K + c_X - 3k)T}}{W_{IJK}^{(0)} + W_{IJK}^{(\pi)} e^{-(c_I + c_J + c_K - 3k/2)T}} F^X + \mathcal{O}(X), \quad (3.1)$$

where we omit the symbol $|_0$. In our 5D model, the visible sector fields Q_I and the hidden sector field X are quasi-localized with the wavefunctions $e^{-\hat{c}_I y}$ and $e^{-\hat{c}_X y}$, respectively, whose effects are encoded in the above expressions as exponential factors. In order to suppress the contributions from the direct coupling, i.e., F^X in the soft terms, it is favored that Q^I and X are localized away from each other.

Taking into account the warp factor $e^{-\hat{k}y}$ of the background geometry, we have basically two choices of kink mass parameters for such sequestering of the hidden sector,

$$\begin{cases} \text{(i)} \ c_I - \frac{k}{2} \equiv -\tilde{c}_I \ < \ 0 \ (^{\forall}I), \ c_X - \frac{k}{2} \equiv \tilde{c}_X \ > \ 0, \\ \text{(ii)} \ c_I - \frac{k}{2} \equiv \tilde{c}_I \ > \ 0 \quad (^{\forall}I), \ c_X - \frac{k}{2} \equiv -\tilde{c}_X \ < \ 0. \end{cases} \quad (3.2)$$

Without loss of generality, we can assume that

$$k > 0.$$

The opposite case $k < 0$ is achieved by exchanging the quantities originating from two fixed points $y = 0$ and $y = \pi R$ each other.

3.1 UV uplifting

First we consider the case (i) defined in eq. (3.2). In this case, the hidden sector field X is localized toward the $y = 0$ (UV) fixed point. The hidden-sector Kähler potential (2.4) and the superpotential (2.6) are determined by

$$\hat{K}_0(T, \bar{T}) = -3 \ln \left(\alpha \frac{1 - \beta |\epsilon_k(T)|^2}{2k} \right), \qquad Z_{X\bar{X}}(T, \bar{T}) = \frac{2k\alpha_X}{\alpha \tilde{c}_X} \frac{1 - \beta_X |\epsilon_X(T)|^2}{1 - \beta |\epsilon_k(T)|^2},$$

$$\hat{W}(T, X) = c - \sum_{n=0}^{n_{np}} A_n e^{-a_n T} + \left\{ W_X^{(0)}|_0 + W_X^{(\pi)}|_0 \epsilon_k^2(T) \epsilon_X(T) \right\} X + \mathcal{O}(X^2).$$

Here and hereafter we use epsilon parameters

$$\epsilon_I(T) = e^{-\tilde{c}_I T}, \quad \epsilon_X(T) = e^{-\tilde{c}_X T}, \quad \epsilon_k(T) = e^{-kT},$$

whose vacuum values can be exponentially suppressed. Especially, $\epsilon_k(\pi R)$ determines the scale at the infra-red (IR) boundary $y = \pi R$ and also the Kaluza-Klein (KK) resonance scale $M_{KK} = \mathcal{O}(\epsilon_k(\pi R)k)$. We compare the above Kähler potential and the superpotential[10] with those of the ISS-KKLT model (A.10) or the ISS-racetrack model (A.19) reviewed as basic models of F-term uplifting in appendix A. In the following, we assume an (approximate) R-symmetry in the hidden sector by assigning the R-charge 2 for X. Then the quadratic and higher powers of X in the hidden sector superpotential are forbidden (or suppressed[11]). This is a requirement for a dynamical SUSY breaking [33].

For example, for

$$A_0 = W^{(0)}|_0 = 0, \qquad n_{np} = 1, \qquad W_X^{(\pi)}|_0 = 0,$$

the above superpotential is in the same form as the ISS-KKLT model (A.10), with the identification

$$A_1 = A, \qquad a_1 = a, \qquad W_X^{(0)}|_0 = \mu^2.$$

Only the difference is that the above modulus Kähler potential $\hat{K}_0(T, \bar{T})$ carries $\epsilon_k(T)$ due to the warped background geometry. In the limit $k \to 0$, this is reduced to $\hat{K}_0(T, \bar{T}) = -3 \ln(T + \bar{T}) + \mathcal{O}(\Omega^{(0)}|_0, \Omega^{(\pi)}|_0)$, that is, to the effective modulus Kähler potential of the KKLT model (A.10) with some corrections from boundary constants. However, even with a finite value of k, as long as the following relation,

$$1 - a(\partial_T \hat{K}_0)^{-1}|_{T=T_0} = 1 - (3k)^{-1} a(1 - |\epsilon_k(T_0)|^{-2}) \ll A_1 c^{-1},$$

---

[10]We implicitly assume some heavy modes living at the fixed points in our model, which generate a SUSY breaking mass m_X for X at the one-loop level as in the ISS-KKLT or ISS-racetrack model shown in eq. (A.11) in appendix A.

[11]If there exist such R-symmetry breaking terms with higher powers of X, SUSY vacua would exist in our model. However, such SUSY points are far away from the SUSY breaking local minimum if the coefficients of the R-breaking terms are suppressed [32].

is satisfied (see eq. (A.3)) and also the T-X mixing is small as in eq. (A.12), the reference point,

$$T_0 \simeq a^{-1}\ln(Ac^{-1}), \qquad X_0 = \mathcal{O}((\mu/m_X)^2 c),$$

satisfying $\hat{W}_T + \hat{K}_T\hat{W} = 0$, $V_X = 0$ ($\hat{W}_X + \hat{K}_X\hat{W} \neq 0$), almost represents the SUSY breaking minimum (see eqs. (A.13) and (A.14)) up to certain small deviations δT and δX as in the ISS-KKLT model shown in appendix A. Therefore, if the parameters satisfy a weak warping condition

$$c \ll |\epsilon_k(T_0)|^2, \qquad \text{i.e.,} \qquad 2kT_0 \ll \ln c^{-1}, \tag{3.3}$$

as well as the low energy SUSY conditions

$$\ln c^{-1} = \mathcal{O}(4\pi^2), \qquad \mu^2 = \mathcal{O}(\tilde{c}),$$

(see eqs. (A.2) and (A.15), respectively), we can realize a kind of the ISS-KKLT model effectively in our 5D framework.

The vacuum energy at the SUSY breaking minimum is vanishing if we tune the corresponding parameters as in eq. (A.16), i.e.,

$$\tilde{c} \simeq \mu^2/\sqrt{3Z_{X\bar{X}}}|_{T=T_0} + \mathcal{O}(\mu^2),$$

where the SUSY AdS vacuum in the modulus T sector alone is uplifted to a SUSY breaking Minkowski minimum by a dynamically generated F-term in the hidden sector X. With negligible Kähler mixing $\hat{K}_{X\bar{T}}, \hat{K}_{XT} = \mathcal{O}(X) \ll 1$, the modulus mass m_T and the SUSY breaking order parameters F^T, F^X, F^C are expressed (see eq. (A.18)) as

$$m_T \simeq -e^{\hat{K}/2}\hat{K}^{T\bar{T}}\hat{W}_{TT}|_{T=T_0,X=X_0},$$
$$F^T \simeq -\frac{\sqrt{3}\hat{K}_T}{\hat{K}_{T\bar{T}}}\left(\sqrt{3} + \frac{\sqrt{\hat{K}^{X\bar{X}}}\hat{W}_{TX}}{\hat{K}_T\hat{W}}\right)\frac{|m_{3/2}|^2}{m_T}\Bigg|_{T=T_0,X=X_0},$$
$$F^X \simeq -\sqrt{\frac{3}{\hat{K}_{X\bar{X}}}}\, m_{3/2}\Bigg|_{T=T_0,X=X_0}, \qquad F^C \simeq Cm_{3/2}|_{T=T_0,X=X_0}.$$

Because we are considering a weak warping (3.3), m_T, F^T, F^C and F^X can be also estimated by those in the ISS-KKLT model, (A.5), (A.6) and (A.17), respectively. That is, the modulus T is heavier than the gravitino by a loop factor, and the tree-level modulus mediation is comparable to the anomaly mediation,

$$m_T = \mathcal{O}(4\pi^2 m_{3/2}), \quad \frac{F^T}{T+\bar{T}} = \mathcal{O}(m_{3/2}/4\pi^2), \quad F^X \sim \frac{F^C}{C} = \mathcal{O}(m_{3/2}). \tag{3.4}$$

On the other hand, for

$$c = 0, \qquad A_0 = W^{(0)}|_0 = 0, \qquad n_{np} = 2, \qquad W_X^{(0)}|_0 = 0,$$

the above superpotential has the same form as one in the ISS-racetrack model (A.19), if we tune the parameters as, e.g., $kT_0 \ll \tilde{c}_X T_0 = \mathcal{O}(4\pi^2)$. However, in this case, a heavier modulus mass can be obtained[12] without affecting the size of F^T and F^C, that is,

$$
m_T = \mathcal{O}(\epsilon_k^{-2}(4\pi^2)^2 m_{3/2}), \quad \frac{F^T}{T + \bar{T}} = \mathcal{O}(m_{3/2}/4\pi^2), \quad F^X \sim \frac{F^C}{C} = \mathcal{O}(m_{3/2}).
$$

Note that the modulus mass m_T is heavier by a factor ϵ_k^{-2} than the one in the ISS-racetrack model (A.8).

So far, we have confirmed that our 5D model can realize the ISS-KKLT-type moduli stabilization, SUSY breaking and F-term uplifting, under the assumption of weak warping (3.3). It might be possible to construct a different class of F-term uplifting model with a strong warping, e.g., $k = \mathcal{O}(4\pi^2)$ and then $\epsilon_k = \mathcal{O}(c)$. In this case, we have to be careful about the fact that the KK scale is quite low and the effects of non-zero modes can be enhanced at low energy. This case is beyond the scope of this paper, and we would study it elsewhere. (For the case with $k = a/3 = \mathcal{O}(4\pi^2)$, see ref. [35].)

Next, we study the nature of sequestering in the case (i), where Q^I are localized toward the $y = \pi R$ (IR) fixed point. The tree-level soft terms for visible fields are found as

$$
M_a = \frac{F^T + k_a^{-1}\left(f_{a,X}^{(0)} + \epsilon_k^{-1}\epsilon_X f_{a,X}^{(\pi)}\right)F^X}{T + \bar{T} + 2k_a^{-1}(f_a^{(0)} + f_a^{(\pi)})} + \mathcal{O}(X),
$$

$$
m_I^2 = \frac{\beta_I \tilde{c}_I^2}{\beta_I - \epsilon_I^2}\left(\frac{\epsilon_I^2}{\beta_I - \epsilon_I^2}|F^T|^2 - \frac{\tilde{\alpha}_I}{3\alpha_I\beta_I}\frac{\epsilon_I^2 - \tilde{\beta}_I \epsilon_k^{-2}\epsilon_X^2}{\tilde{c}_I(\tilde{c}_X - \tilde{c}_I - k)}|F^X|^2\right) + \mathcal{O}(X),
$$

$$
A_{IJK} = -\left\{\left(\frac{\beta_I \tilde{c}_I}{\beta_I - \epsilon_I^2} + (I \leftrightarrow J) + (I \leftrightarrow K)\right) - \frac{(\tilde{c}_I + \tilde{c}_J + \tilde{c}_K)W_{IJK}^{(\pi)}}{W_{IJK}^{(0)}\epsilon_I\epsilon_J\epsilon_K + W_{IJK}^{(\pi)}}\right\}F^T
$$

$$
+ \frac{W_{IJK,X}^{(0)}\epsilon_I\epsilon_J\epsilon_K + W_{IJK,X}^{(\pi)}\epsilon_k^{-1}\epsilon_X}{W_{IJK}^{(0)}\epsilon_I\epsilon_J\epsilon_K + W_{IJK}^{(\pi)}}F^X + \mathcal{O}(X), \tag{3.5}
$$

where we omit the symbol $|_0$.

We assume that the order parameters F^T, F^C and F^X are given by (3.4). For $f_{a,X}^{(0)} = \mathcal{O}(1)$, we find $M_a = \mathcal{O}(m_{3/2})$. If we assume that the hidden sector R-symmetry is preserved also in the visible sector gauge kinetic functions, i.e., $f_{a,X}^{(0)} = f_{a,X}^{(\pi)} = 0$, the F^X can not contribute to the gaugino mass. In this case, the gaugino mass is a mirage-type, where the tree-level modulus mediation and the anomaly mediation are comparable to each other.

The tree-level contributions to scalar masses m_I^2 are suppressed compared with the anomaly mediation if ϵ_I^2, $\epsilon_k^{-2}\epsilon_X^2 \ll 1/(4\pi^2)$. In this case, the sequestering is maximal. The limit $\tilde{c}_I \to \infty$, that is, $\epsilon_I \to 0$, corresponds to the complete localization of Q^I at the fixed point $y = \pi R$. Then, Q_I is a chiral multiplet living only at the fixed point $y = \pi R$ and its scalar mass m_I has no contribution due to F^T. Similarly, the limit $\tilde{c}_X \to \infty$ makes

---

[12]This enhancement of modulus mass would play a role to avoid the so-called moduli-induced gravitino/neutralino problem [34] in the KKLT-type scenario. Note that the modulus mass is already enhanced by a loop factor in the ISS-racetrack model (A.8) compared with the original KKLT model [12].

X live only at the fixed point $y = 0$, and the scalar mass m_I of Q_I has no contribution due to F^X. At any rate, the modulus mediation is always subdominant compared with the direct mediation unless all the epsilon parameters are of $\mathcal{O}(1)$. Note that the modulus mediation and the direct mediation typically give a positive and a negative contribution to scalar masses squared, respectively.

The contact term between X and Q^I would be induced by loop effects through the gravitational interaction even if X and Q^I are completely localized at opposite fixed points [36]. Such loop effects would lead to corrections to scalar masses squared Δm_I^2, which are proportional to $|F^X|^2$. However, such corrections are suppressed by the one-loop factor and the warp factor ϵ_k^2. (See also ref. [37].)[13] Thus, such corrections are negligible compared with the anomaly mediation in a weakly warped case, e.g. $\epsilon_k \sim M_{\mathrm{GUT}}/M_{Pl}$ and even in the case with $\epsilon_k \sim 1/(4\pi^2)$.

For ϵ_I^2 ($^\forall I$), $\epsilon_I \epsilon_J \epsilon_K$, $\epsilon_k^{-1}\epsilon_X \ll 1/(4\pi^2)$, the tree-level contributions to the A-term are suppressed compared with the anomaly mediation if $W_{IJK}^{(\pi)}|0 \neq 0$, and the maximal sequestering is achieved. However, if $W_{IJK}^{(\pi)}|_0 = 0$, the A-term becomes $A_{IJK} \simeq -(\tilde{c}_I + \tilde{c}_J + \tilde{c}_K)F^T$ [20, 21] which can be a mirage-type for $\tilde{c}_I$, $\tilde{c}_J$, $\tilde{c}_K = \mathcal{O}(1)$.

3.2 IR uplifting

Next we consider the case (ii) defined in eq. (3.2). In this case, the hidden sector field X is localized toward the $y = \pi R$ (IR) fixed point. For $k = 0$, i.e., a flat extra dimension, the case (ii) is physically equivalent to the case (i) under the exchange of two fixed points. The difference between the case (i) and the case (ii) is enhanced for large k. Thus, we only consider such a case,

$$\epsilon_k \lesssim \epsilon_X, \ \epsilon_I, \qquad (\text{i.e. } k \gtrsim \tilde{c}_X, \tilde{c}_I), \tag{3.6}$$

for $^\forall I$ in the following. In this case, the hidden sector (and the modulus) Kähler and the superpotential are obtained as

$$\hat{K}(T, \bar{T}, \tilde{X}, \bar{\tilde{X}}) = -3\ln\left(\alpha \frac{1 - \beta|\epsilon_k(T)|^2}{2k}\right) + \frac{2k\alpha_X}{\alpha\tilde{c}_X}\frac{\beta_X - |\epsilon_X(T)|^2}{1 - \beta|\epsilon_k(T)|^2}|\tilde{X}|^2 + \mathcal{O}(|\tilde{X}|^4),$$

$$\hat{W}(T, \tilde{X}) = c - \sum_{n=0}^{n_{\mathrm{np}}} A_n e^{-a_n T} + \left\{W_X^{(0)}|_0\epsilon_X(T) + W_X^{(\pi)}|_0\epsilon_k^2(T)\right\}\tilde{X} + \mathcal{O}(\tilde{X}^2),$$

where we redefined X as

$$\tilde{X} = \epsilon_X^{-1}(T)X.$$

The superpotential can be the ISS-KKLT type (A.10) or the ISS-racetrack type (A.19), for the case of $W_X^{(0)}|_0 = 0$ or $W_X^{(\pi)}|_0 = 0$ with the identification $\epsilon_X(T)W_X^{(0)}|_0 = \mu^2(T)$ or $\epsilon_k^2(T)W_X^{(\pi)}|_0 = \mu^2(T)$, respectively, where a sizable T-X mixing exists for $k \gtrsim \tilde{c}_X \gtrsim 4\pi^2$. As shown previously, the effect of warping in the Kähler potential is not relevant to the

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[13]Other sources of contact terms would be bulk vector multiplets [38, 14], which are (assumed to be) absent in our 5D model.

analysis in appendix A, only when a weak warping condition (3.3) is satisfied, and this restricts ϵ_X and ϵ_I as $c \ll \epsilon_X, \epsilon_I$ due to eq. (3.6). Then the condition corresponding to eq. (A.15) can not be satisfied, that is, the analysis in appendix A is not valid in this case, and the structure of F-term uplifting can be different from the conventional one due to eq. (3.6). We would study also this case elsewhere as well as the above mentioned strong warping case. At any rate, for the weak warping (3.3) *without* the condition (3.6), the physics should be almost the same as the case (i) as mentioned above. Then we would realize the ISS-KKLT model and the ISS-racetrack model effectively, also in the case (ii).

The tree-level soft terms for visible fields in the case (ii) are derived as

$$
M_a = \frac{F^T + 2k_a^{-1}\left(\epsilon_X f_{a,X}^{(0)} + \epsilon_k^{-1} f_{a,X}^{(\pi)}\right) F^{\tilde{X}}}{T + \bar{T} + 2k_a^{-1}(f_a^{(0)} + f_a^{(\pi)})} + \mathcal{O}(\tilde{X}),
$$

$$
m_I^2 = \frac{\beta_I \tilde{c}_I^2}{1 - \beta_I \epsilon_I^2}\left(\frac{\epsilon_I^2}{1 - \beta_I \epsilon_I^2}|F^T|^2 - \frac{\tilde{\alpha}_I}{3\alpha_I \beta_I}\frac{\epsilon_X^2 - \tilde{\beta}_I \epsilon_k^{-2}\epsilon_I^2}{\tilde{c}_I(\tilde{c}_I - \tilde{c}_X - k)}|F^{\tilde{X}}|^2\right) + \mathcal{O}(\tilde{X}),
$$

$$
A_{IJK} = -\left\{\left(\frac{\beta_I \tilde{c}_I \epsilon_I^2}{1 - \beta_I \epsilon_I^2} + (I \leftrightarrow J) + (I \leftrightarrow K)\right) + \frac{(\tilde{c}_I + \tilde{c}_J + \tilde{c}_K)W_{IJK}^{(\pi)}\epsilon_I\epsilon_J\epsilon_K}{W_{IJK}^{(0)} + W_{IJK}^{(\pi)}\epsilon_I\epsilon_J\epsilon_K}\right\} F^T
$$

$$
+ \frac{W_{IJK,X}^{(0)}\epsilon_X + W_{IJK,X}^{(\pi)}\epsilon_k^{-1}\epsilon_I\epsilon_J\epsilon_K}{W_{IJK}^{(0)} + W_{IJK}^{(\pi)}\epsilon_I\epsilon_J\epsilon_K} F^{\tilde{X}} + \mathcal{O}(\tilde{X}),
$$

where we again omit the symbol $|_0$. The difference from eq. (3.5) is just the position of ϵ_k^{-1} factors aside from the exchange of $y = 0$ and $y = \pi R$ fixed points with each other.

If we consider the case satisfying eq. (3.6) (although above we have not showed the corresponding vacuum in the hidden sector), the contribution from the direct coupling to the gaugino mass M_a is enhanced by ϵ_k^{-1} for $f_{a,X}^{(\pi)} \neq 0$, the scalar mass m_I is of $\mathcal{O}(F^{\tilde{X}})$ or larger, but the A-term can be suppressed if $W_{IJK}^{(0)} \neq 0$.

3.3 Patterns of soft terms

Here, let us summarize resultant patterns of soft SUSY breaking terms in our model. First recall the sizes of F^T, F^X and F^C, (3.4), that is, F^X is larger than F^T by a factor of $\mathcal{O}(4\pi^2)$ and the modulus mediation is comparable to the anomaly mediation. When $f_{a,X}^{(0)}$, $f_{a,X}^{(\pi)}e^{-(c_X-3k/2)T} = \mathcal{O}(1)$, and $e^{(c_I-k/2)T}$ or $e^{-(c_X-3k/2)T}$ is not suppressed in eq. (3.1), the F-term of X, F^X, is dominant in all of soft terms, which are obtained as

$$
M_a = \mathcal{O}(m_{3/2}), \qquad m_I^2 = \mathcal{O}(m_{3/2}^2), \qquad A_{IJK} = \mathcal{O}(m_{3/2}),
$$

that is, the visible sector is not sequestered from the dominant SUSY breaking source. Their explicit ratios to $m_{3/2}$ are model-dependent like generic spectrum due to the gravity mediation.

When $f_{a,X}^{(0),(\pi)}$ are sufficiently suppressed, the size of gaugino masses is estimated as $M_a = \mathcal{O}(m_{3/2}/(4\pi^2))$. On the other hand, unless $e^{(c_I-k/2)T}$ or $e^{-(c_X-3k/2)T}$ is not suppressed in eq. (3.1), the size of scalar masses is estimated as $m_I^2 = \mathcal{O}(m_{3/2}^2)$. However,

those are tachyonic in a natural parameter region like $\alpha_I \sim \tilde{\alpha}_I \sim \beta_I \sim \tilde{\beta}_I \sim 1$. Thus, the contribution of F^X to scalar masses must be sequestered except in a certain model like a negative value of $\tilde{\alpha}_I/\alpha_I$ and/or $\tilde{\beta}_I/\beta_I$. Note that when we suppress the contribution of F^X to scalar masses by requiring $e^{(c_I - k/2)T} \ll 1$ as well as $e^{-(c_X - 3k/2)T} \ll 1$, the contribution from F^T is always suppressed.[14] In this case, the modulus mediation contribution to the gaugino masses is obtained as

$$M_a = \frac{F^T}{T + \bar{T} + (2/k_a)(f_a^{(0)} + f_a^{(\pi)})} = \mathcal{O}(m_{3/2}/(4\pi^2)),$$

and the F^X contributions to the scalar masses in the case (i), i.e. the UV uplifting (3.5), are obtained as

$$m_I^2 = -\frac{\tilde{c}_I \tilde{\alpha}_I}{3\alpha_I \beta_I} \frac{\epsilon_I^2 - \tilde{\beta}_I \epsilon_k^{-1} \epsilon_X}{(\tilde{c}_X - \tilde{c}_I - k)} |F^X|^2.$$

When $\alpha_I = \tilde{\alpha}_I = \beta_I = \tilde{\beta}_I = 1$, this reduces to

$$m_I^2 = -\frac{\tilde{c}_I}{3} \frac{\epsilon_I^2 - \epsilon_k^{-2} \epsilon_X^2}{(\tilde{c}_X - \tilde{c}_I - k)} |F^X|^2.$$

This scalar mass must be suppressed as $|m_I^2| \leq O(M_a^2)$ as the above reason. Otherwise, scalar masses become tachyonic even at a low-energy scale when we include radiative corrections due to gaugino masses. Then, the anomaly mediation is comparable, that is, the mirage mediation, and magnitudes of soft masses are estimated as

$$M_a = \mathcal{O}(m_{3/2}/(4\pi^2)), \qquad m_I^2 = \mathcal{O}(m_{3/2}^2/(4\pi^2)^2).$$

In order to avoid the SUSY flavor problem, the universal parameters, $c_I = c_J$ for $I \neq J$, are generically favored, as far as the modulus and the direct mediations are concerned. However, if all the parameters c_I are chosen so that $\epsilon_I \ll 1/(4\pi^2)$ $(^\forall I)$ is satisfied, the anomaly mediation is dominant in scalar masses and A-terms. In that case, nonuniversal parameters, $c_I \neq c_J$ $(I \neq J)$, would not cause problematic SUSY flavor violations.

4. Summary and discussions

We studied concretely moduli stabilization, SUSY breaking, F-term uplifting and sequestering, altogether in a simple supergravity model with a single extra dimension. These issues are realized in a fully dynamical way by the use of wavefunction localization in extra dimension, allowing explicit calculations. Contributions due to the uplifting sector make soft scalar masses squared tachyonic and their absolute values are always larger than contributions due to the modulus F-term. We found that the sequestering in the soft scalar mass and the A-term can be achieved within the framework of F-term uplifting in our 5D

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[14]This is because our model has only a single extra dimension. If we would extend our scenario to models with more than one extra dimensions, we could obtain soft scalar masses, where the contribution from the dominant SUSY breaking F^X is sequestered but some moduli F-terms have significant contributions.

model. Because the radius modulus is stabilized by a KKLT-type potential, the gaugino mass is a mirage-type if the visible sector gauge kinetic function preserves the hidden sector R-symmetry which is responsible for the dynamical SUSY breaking, and then the tree-level modulus mediation is comparable to the anomaly mediation. It is notable that the TeV scale mirage mediation [8] can solve the so-called little hierarchy problem [39] within the MSSM, due to the gluino and wino mass unification at the TeV scale [40]. Note also that our 5D model might have a corresponding conformal field theory (CFT) description[15] due to the AdS/CFT correspondence [41].

We have only considered typical cases (3.2) from the viewpoint of sequestering, i.e., all the generations of quarks and leptons localize toward the common fixed point. However, as mentioned in the introduction, it is known that a certain class of realistic Yukawa matrices for quarks and leptons can be obtained by the wavefunction localization in 5D supergravity theory. In this case, either the light generation or the heavy generation would be forced to localize toward the same fixed points as the SUSY breaking field X. For such a generation, sequestering can not occur, and the squark/slepton might receive a soft scalar mass with a large magnitude from the direct coupling. We would study more detailed flavor structure of our model with such realistic Yukawa couplings in a separate work [30]. (For the case of radion domination or Scherk-Schwarz (SS) SUSY breaking [43],[16] see refs. [20, 21] and references therein.)

We also restricted to the case with a single modulus. If we consider more Z_2-odd vector multiplets in 5D, we have multiple moduli which can cause moduli mixing in the gauge kinetic function, and then in the nonperturbative superpotential. Such mixing effects could play important roles in the moduli stabilization [47, 14] after integrating out heavy moduli [4, 48]. It is also important to study models with more than one extra dimensions and it would be interesting to investigate whether the situation would change, e.g., whether contributions due to the uplifting sector make soft scalar masses squared tachyonic or not, and whether their absolute values are always larger than the modulus F-term contributions.

Our model is not directly related to a certain string model known until now. However, the result of this paper would help us to understand some basic features of the moduli stabilization, the SUSY breaking, the realization of Minkowski (de Sitter) vacuum and the sequestering in higher-dimensional supergravity models and superstring models.

Acknowledgments

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[15]See ref. [42] for a realization of ISS sector in this direction.

[16]Note that the SS twist is prohibited in our 5D supergravity model [44]. We can introduce such a twist in a different type of orbifold supergravity [45, 46].

A. Basic structure of F-term uplifting

In this appendix, we review the scenario of F-term uplifting [10 – 12] based on the KKLT-type SUSY AdS vacua.

A.1 KKLT model

First we briefly review the original KKLT model [3 – 5]. The F-term potential of the 4D N=1 supergravity theory is given by

$$V_F = K_{I\bar{J}} F^I \bar{F}^{\bar{J}} - 3|m_{3/2}|^2,$$

where F^I is the F-component of the I-th chiral multiplet and $m_{3/2}$ is the gravitino mass given respectively as $F^I = -e^{K/2} K^{I\bar{J}} (\bar{W}_{\bar{J}} - K_{\bar{J}} \bar{W})$ and $m_{3/2} = e^{K/2} W$.

Before introducing an uplifting sector, the KKLT model assumes the following Kähler potential and superpotential,

$$K = -3\ln(T + \bar{T}), \qquad W = c - Ae^{-aT}, \tag{A.1}$$

in the 4D effective supergravity theory, where T is a light modulus, the first constant term in W, c, originates from a flux, the second T-dependent term comes from a nonperturbative effect such as a gaugino condensation and then $a = \mathcal{O}(4\pi^2)$, $A = \mathcal{O}(1)$. We measure all the mass scales in the unit $M_{Pl} = 1$. In order to realize low scale SUSY breaking (TeV scale gravitino mass), we consider a tiny flux constant,

$$\ln c^{-1} = \mathcal{O}(4\pi^2). \tag{A.2}$$

Then, the SUSY stationary condition of the scalar potential, $F^T = 0$, is satisfied by

$$aT_0 = \ln(Ac^{-1}) + \ln(1 - aK_T^{-1}|_{T=T_0}) \simeq \ln(Ac^{-1}) = \mathcal{O}(4\pi^2). \tag{A.3}$$

This stationary solution corresponds to a SUSY AdS vacuum of the scalar potential with a negative vacuum energy $V_{\text{SUSY}} = -3|m_{3/2}|^2|_{T=T_0} = \mathcal{O}(c^2)$.

In the original KKLT model, this SUSY AdS minimum is uplifted to a Minkowski minimum by introducing an anti D3-brane. The anti-brane breaks $N = 1$ SUSY explicitly in the 4D effective supergravity, and generates an uplifting potential energy

$$U = \int d^4\theta\, |C|^4 \xi \theta^2 \bar{\theta}^2 = \xi e^{2K/3}, \tag{A.4}$$

where ξ is a constant. The total scalar potential is now given by $V = V_F + U$ and the previous minimum is shifted as $T = T_0 + \delta T$. We tune the constant $\xi = 3e^{-2K/3}|m_{3/2}|^2|_{T=T_0}$ so that $V = 0$ at the leading order in a $\delta T/T_0$ expansion. Then, we find the shift of T at this Minkowski minimum, $\delta T/T_0 \sim 1/(aT_0)^2 = \mathcal{O}(1/(4\pi^2)^2)$, and the modulus mass,

$$m_T \simeq aT_0\, m_{3/2} = \mathcal{O}(4\pi^2 m_{3/2}). \tag{A.5}$$

The SUSY breaking order parameters are estimated as

$$\frac{F^T}{T + \bar{T}} \simeq \frac{m_{3/2}}{aT_0} = \mathcal{O}(m_{3/2}/4\pi^2), \qquad \frac{F^C}{C} \simeq m_{3/2} = \mathcal{O}(c). \tag{A.6}$$

Here we find that the tree level modulus mediation is comparable to the one-loop anomaly mediation, $F^T \sim F^C/4\pi^2$, that is, the so-called mirage mediation [6].

If we consider a racetrack model,

$$W = Ce^{-cT} - Ae^{-aT}, \tag{A.7}$$

instead of the KKLT superpotential (A.1), the modulus mass m^T (F^T) is more enhanced (suppressed) as

$$m_T \simeq (aT_0)(cT_0)\, m_{3/2} \;=\; \mathcal{O}((4\pi^2)^2 m_{3/2}), \tag{A.8}$$

$$\frac{F^T}{T+\bar{T}} \simeq \frac{m_{3/2}}{(aT_0)(cT_0)} \;=\; \mathcal{O}(m_{3/2}/(4\pi^2)^2). \tag{A.9}$$

Thus the modulus mediation is negligible compared with the anomaly mediation in this case.

A.2 ISS-KKLT model

In the original KKLT model, the uplifting potential (A.4) is a kind of an explicit SUSY breaking term in the low energy effective theory. Instead, we can consider the case that the uplifting potential is supplied by the F-term of a dynamical SUSY breaking sector X which is included in the F-term potential V_F itself. When X is somehow sequestered from T, the picture that the AdS SUSY vacua existing in the T sector alone is uplifted by F^X generated by the X sector, would be valid. Then, we assume the following Kähler and superpotential,

$$K = -3\ln(T+\bar{T}) + Z_{X\bar{X}}(T,\bar{T})|X|^2, \qquad W \;=\; c - Ae^{-aT} + \mu^2(T)X. \tag{A.10}$$

The tadpole of X would appear as a low energy effective superpotential term in the dynamical SUSY breaking sector, such as the O'Raifeartaigh model [15] and the Intriligator-Seiberg-Shih (ISS) model [16], after integrating out heavy modes, and we call the model (A.10) the ISS-KKLT model. The effect of the heavy modes appears at low energy as a one-loop correction to the above Kähler potential,

$$\Delta K = -\Lambda^{-2} Z^{(1)}(T,\bar{T})|X|^4,$$

and then the correction to the scalar potential in this case is expressed as a SUSY breaking mass term of X,

$$\Delta V_F = m_X^2 |X|^2 + \mathcal{O}(|X|^4), \qquad m_X^2 \;=\; e^K (4\mu^4 Z^{(1)})/(Z_{X\bar{X}}\Lambda)^2, \tag{A.11}$$

where Λ is the mass scale of the heavy modes.[17]

If the T-X mixing is small,

$$|K_{T\bar{X}}| \ll K_{T\bar{T}},\; K_{X\bar{X}} = \mathcal{O}(1), \qquad |K_{TX}| \ll |K_{TT}|, \qquad |W_{TX}| \ll |W_{TT}|, \tag{A.12}$$

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[17]In Ref [12], the $Z_{X\bar{X}}$ (Z in the notation of ref. [12]) dependence of m_X^2 should be replaced by $Z_{X\bar{X}}^{-2}$, which is a typographical error.

the solution

$$T_0 \simeq a^{-1}\ln(Ac^{-1}), \qquad X_0 \simeq 2(\mu_0/m_X)^2\tilde{c}, \tag{A.13}$$

satisfying

$$W_T + K_T W = 0, \qquad V_X = 0 \quad (W_X + K_X W \neq 0), \tag{A.14}$$

would be a good *reference point* of the SUSY breaking minimum, where $Z_0 = Z_{X\bar{X}}(T_0, \bar{T}_0)$, $\mu_0 = \mu(T_0)$ and $\tilde{c} = c - Ae^{-aT_0}$. Here we have assumed

$$m_X^2 = \mathcal{O}(\mu_0^2/(4\pi^2)), \qquad \mu_0^2 = \mathcal{O}(\tilde{c}). \tag{A.15}$$

We expand the potential around this point by substituting $T = T_0 + \delta T$, $X = X_0 + \delta T$ and find $\delta T/T_0 \sim 1/(aT_0)^2 = \mathcal{O}(1/(4\pi^2)^2)$, $\delta X/X_0 \sim 1/(aT_0) = \mathcal{O}(1/(4\pi^2))$. The vacuum energy at this SUSY breaking minimum is vanishing, $V = 0$, if we tune the parameters as

$$\tilde{c} \simeq \mu_0^2/\sqrt{3Z_0} + \mathcal{O}(\mu_0^4), \tag{A.16}$$

which is consistent with the above assumption $\mu_0^2 = \mathcal{O}(c)$. The leading moduli mass m_T and F^T are the same as those in the original KKLT model, while we obtain

$$F^X \simeq \sqrt{3/Z_0}\, m_{3/2} = \mathcal{O}(m_{3/2}). \tag{A.17}$$

In general, the Kähler mixing at the reference point (A.13) is suppressed $K_{T\bar{X}}, K_{TX} \propto X_0 = \mathcal{O}(c)$ satisfying the first two conditions in eq. (A.12), in the ISS-type model where the VEV of X can be significantly small. In such a case, we find general expressions [12],

$$m_T \simeq -e^{K/2}K^{T\bar{T}}W_{TT}\Big|_{T=T_0,X=X_0},$$
$$F^T \simeq -\frac{\sqrt{3}K_T}{K_{T\bar{T}}}\left(\sqrt{3} + \frac{\sqrt{K^{X\bar{X}}}W_{TX}}{K_T W}\right)\frac{|m_{3/2}|^2}{m_T}\Bigg|_{T=T_0,X=X_0},$$
$$F^X \simeq -\sqrt{\frac{3}{K_{X\bar{X}}}}\, m_{3/2}\Bigg|_{T=T_0,X=X_0}, \qquad F^C \simeq Cm_{3/2}\big|_{T=T_0,X=X_0}. \tag{A.18}$$

This can be also adopted to the case with a sizable value of W_{TX} under the assumption that the reference point (A.13) is stable.

We can generalize the ISS-KKLT model to the case with, e.g., $\mu^2(T) = Be^{-bT}$ where the superpotential mixing W_{TX} is sizable [12]. In this case, the perturbation around the reference point (A.13) becomes unstable if the superpotential terms of the modulus sector is KKLT-type. We can stabilize the reference point by considering a racetrack-type modulus sector (A.7), i.e.,

$$W = Ce^{-cT} - Ae^{-aT} + Be^{-bT}X, \tag{A.19}$$

which we call the ISS-racetrack model. In this case, the deviations from the reference point (A.13), are estimated as $\delta T/T_0 \sim bT_0/(aT_0\, cT_0)^2 = \mathcal{O}(1/(4\pi^2)^3)$, $\delta X/X_0 \sim$

$(bT_0)^2/(aT_0 \, cT_0)^2 = \mathcal{O}(1/(4\pi^2)^2)$. Because of the racetrack structure, the modulus mass is enhanced as eq. (A.8), but unlike eq. (A.9), F^T is estimated by eq. (A.18) as

$$\frac{F^T}{T+\bar{T}} \simeq \frac{b}{(aT_0)(cT_0)} m_{3/2} \; = \; \mathcal{O}(m_{3/2}/4\pi^2),$$

due to the enhancement factor b from the T-X mixing, W_{TX} (i.e., the second term in the parenthesis of F^T in eq. (A.18)). Thus the tree-level modulus mediation is comparable to the anomaly mediation, in spite of a larger mass hierarchy between the modulus and the gravitino mass (A.8).

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