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Supersymmetry breaking in warped geometry

K. Choi[1,a], D.Y. Kim[1,b], I.-W. Kim[1,c], T. Kobayashi[2,d]

[1] Department of Physics, Korea Advanced Institute of Science and Technology, Daejeon 305-701, Korea
[2] Department of Physics, Kyoto University, Kyoto 606-8502, Japan

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Abstract. We examine the soft supersymmetry breaking parameters in supersymmetric theories on a slice of AdS_5 which generate the hierarchical Yukawa couplings by dynamically quasi-localizing the bulk matter fields in an extra dimension. Such models can be regarded as the AdS dual of the recently studied 4-dimensional models which contain a supersymmetric CFT to generate the hierarchical Yukawa couplings. In such models, if supersymmetry breaking is mediated by the bulk radion superfield and/or some brane chiral superfields, potentially dangerous flavor-violating soft parameters are suppressed with an appropriate correlation with the Yukawa coupling suppression, thereby avoiding the SUSY flavor problem in a natural manner. We present some models of radion-dominated supersymmetry breaking which yield a highly predictive form of soft parameters in this framework, and discuss the constraints from flavor-changing rare processes. Most of the discussions in this paper can be applied also to models with a flat extra dimension in which the Yukawa hierarchy is generated by quasi-localizing the bulk matter fields in the extra dimension.

1 Introduction

A warped extra dimension can provide a rationale for various hierarchical structures in particle physics. For instance, in 5-dimensional (5D) theory on a slice of AdS_5 with AdS curvature k and orbifold radius R, the small warp factor $e^{-\pi kR}$ can generate a huge hierarchy between the 4D Planck scale ($M_{Pl} \sim 10^{18}$ GeV) and the weak scale ($M_W \sim 10^2$ GeV) [1], and/or hierarchical quark and lepton masses [2, 3], and/or small neutrino masses [4]. The underlying dynamical reason for small M_W/M_{Pl} and small fermion masses is the quasi-localization of the gravity and fermion zero modes in the extra dimension.

In this paper, we wish to examine soft supersymmetry (SUSY) breaking parameters in models with a warped extra dimension in which the hierarchical Yukawa couplings are generated by quasi-localizing the bulk matter fields in the extra dimension [5]. To this end, we consider supersymmetric 5D theory on a slice of AdS_5 in which all quark and lepton fields arise from 5D bulk hypermultiplets [2]. The graviphoton of the model can be gauged appropriately to have non-zero bulk cosmological constant and hypermultiplet masses which would make the gravity zero mode quasi-localized at one boundary ($y = 0$) and the matter zero modes quasi-localized at any of the two boundaries ($y = 0$ or π). In this type of models, the small warp factor and the small Yukawa couplings have a common origin, i.e. the dynamical quasi-localization of the zero modes associated with the gauging of the graviphoton. As a result, the small Yukawa couplings are typically given by $y_{mn} \sim e^{-c\pi kR}$ where c is a constant of order unity, which would require that the radion is stabilized at a value giving the warp factor: $e^{-\pi kR} \approx 10^{-2}$–$10^{-5}$ and the corresponding Kaluza–Klein (KK) scale: $M_{KK} \approx k e^{-\pi kR} \approx 10^{16}$–$10^{13}$ GeV. This type of models can be regarded as the AdS dual of the recently studied 4D models [6] which contain a supersymmetric CFT sector generating the Yukawa hierarchy through the renormalization group evolution.

It has been noted that supersymmetric CFT dynamics can provide also a mechanism to suppress dangerous flavor-violating soft parameters [6]. This suggests that a similar suppression of flavor violation can take place in AdS_5 models also. As we will see, flavor-violating soft parameters in AdS_5 models are indeed suppressed under a reasonable assumption on SUSY breaking. Furthermore, if SUSY is broken by the F-term of the radion superfield [7–9], the model leads to a concrete prediction on soft parameters which can be tested by future experiments. In this paper, we first discuss general forms of the soft parameters in supersymmetric AdS_5 model and later present some models of radion-dominated SUSY breaking in AdS_5 which can pass all constraints on flavor violation without severe fine tuning of the parameters. Most of our results on the Yukawa couplings and soft parameters can be applied straightforwardly to the models with a *flat* extra dimension in which the bulk matter fields are quasi-localized in the extra dimension. For this, one can simply take the limit that

[a] e-mail: kchoi@hep.kaist.ac.kr
[b] e-mail: kimdoyoung@hep.kaist.ac.kr
[c] e-mail: iwkim@hep.kaist.ac.kr
[d] e-mail: kobayash@gauge.scphys.kyoto-u.ac.jp

the AdS curvature becomes zero, while the hypermultiplet masses remain non-zero.

2 Yukawa couplings and soft parameters in supersymmetric AdS$_5$ model

To proceed, let us consider a generic 5D gauge theory coupled to the minimal 5D supergravity (SUGRA) on S^1/Z_2. The action of the model is given by [2, 10, 11]

$$S = \int \mathrm{d}^5 x \sqrt{-G} \times \Bigg[\frac{1}{2} \left(\mathcal{R} + \bar{\Psi}^i_M \gamma^{MNP} D_N \Psi_{iP} - \frac{3}{2} C_{MN} C^{MN} + 12k^2 \right) + \frac{1}{g^2_{5a}} \Bigg(-\frac{1}{4} F^{aMN} F^a_{MN} + \frac{1}{2} D_M \phi^a D^M \phi^a + \frac{\mathrm{i}}{2} \bar{\lambda}^{ai} \gamma^M D_M \lambda^a_i \Bigg) + |D_M h^i_I|^2 + \mathrm{i}\bar{\Psi}_I \gamma^M D_M \Psi_I + \mathrm{i} c_I k \epsilon(y) \bar{\Psi}_I \Psi_I + \ldots \Bigg], \tag{1}$$

where $\mathcal{R}$ is the 5D Ricci scalar, Ψ^i_M $(i = 1, 2)$ are the symplectic Majorana gravitinos, $C_{MN} = \partial_M B_N - \partial_N B_M$ is the graviphoton field strength, and y is the 5th coordinate with a fundamental range $0 \le y \le \pi$. Here ϕ^a, A^a_M and λ^{ia} are 5D scalar, vector and symplectic Majorana spinors constituting a 5D vector multiplet, h^i_I and Ψ_I are a 5D scalar and Dirac spinor constituting the Ith hypermultiplet with kink mass $c_I k \epsilon(y)$. The AdS curvature k and the kink mass $c_I k$ are related with the gauging of the graviphoton as indicated by the following covariant derivatives [10, 12]:

$$\begin{aligned} D_M h^i_I &= \partial_M h^i_I - \mathrm{i} \left(\frac{3}{2} (\sigma_3)^i_j - c_I \delta^i_j \right) k \epsilon(y) B_M h^j_I + \ldots, \\ D_M \Psi_I &= \partial_M \Psi_I + \mathrm{i} c_I k \epsilon(y) B_M \Psi_I + \ldots, \\ D_M \lambda^{ai} &= \partial_M \lambda^{ai} - \mathrm{i} \frac{3}{2} (\sigma_3)^i_j k \epsilon(y) B_M \lambda^{aj} + \ldots, \end{aligned} \tag{2}$$

where the ellipses stand for the other gauge couplings. Note that we set the 5D Planck mass $M_5 = 1$ and all dimensionful parameters, e.g. the 5D gauge coupling g_{5a} and the AdS curvature k, are defined in this unit. Although not required within 5D SUGRA, it is not unreasonable to assume that the graviphoton charges (c_I) are *quantized* in an appropriate unit, i.e. c_I/c_J are rational numbers, which assumption we will adopt throughout this paper. We note that in the CFT view, the c_I are related to the anomalous dimensions γ_I of the chiral operators built from CFT-charged superfields [6]. At a superconformal fixed point, the γ_I are determined simply by group theory factors, so they are quantized.

With appropriate values of the brane cosmological constants at the orbifold fixed points $(y = 0, \pi)$, the ground state geometry of the action (1) is given by a slice of AdS$_5$ having the following form of 5D metric:

$$G_{MN} \mathrm{d}x^M \mathrm{d}x^N = \mathrm{e}^{-2kRy} \eta_{\mu\nu} \mathrm{d}x^\mu \mathrm{d}x^\nu + R^2 \mathrm{d}y^2. \tag{3}$$

It is convenient to write the 5D action (1) in $N = 1$ superspace [8, 13, 14]. For the 5D SUGRA multiplet, we keep only the radion superfield

$$T = \left(R + \mathrm{i} B_5, \frac{1}{2} (1 + \gamma_5) \Psi^{i=2}_5 \right)$$

and replace the other fields by their vacuum expectation values. For the 5D vector multiplets and hypermultiplets, one needs appropriate R- and B_5-dependent field redefinitions to construct the corresponding $N = 1$ superfields [8, 12]. After such a field redefinition, the relevant piece of the action is given by[1] [8, 13]

$$\begin{aligned} &\int \mathrm{d}^5 x \left[\int \mathrm{d}^4 \theta \, \frac{T + T^*}{2} \left(\hat{H}^*_I \hat{H}_I + \hat{H}^{c*}_I \hat{H}^c_I \right) + \left\{ \int \mathrm{d}^2 \theta \frac{1}{4 g^2_{5a}} T W^{a\alpha} W^a_\alpha + \mathrm{h.c.} \right\} \right] \\ &= \int \mathrm{d}^5 x \left[\int \mathrm{d}^4 \theta \frac{T + T^*}{2} \left(\mathrm{e}^{\left(\frac{1}{2} - c_I \right)(T + T^*) k |y|} H^*_I H_I + \mathrm{e}^{\left(\frac{1}{2} + c_I \right)(T + T^*) k |y|} H^{c*}_I H^c_I \right) + \left\{ \int \mathrm{d}^2 \theta \, \frac{1}{4 g^2_{5a}} T W^{a\alpha} W^a_\alpha + \mathrm{h.c.} \right\} \right], \end{aligned} \tag{4}$$

where W^a_α is the chiral spinor superfield for the vector superfield $\mathcal{V}^a$ containing (A^a_μ, λ^a) with $\lambda^a = \frac{1}{2}(1 - \gamma_5) \lambda^{a1}$, H_I and H^c_I are chiral superfields containing (h^1_I, ψ_I) and (h^{2*}_I, ψ^c_I), respectively, with $\psi_I = \frac{1}{2}(1 - \gamma_5) \Psi_I$, $\psi^c{}_I = \frac{1}{2}(1 + \gamma_5) \Psi_I$. Here we consider two superfields bases for hypermultiplets, $(\hat{H}_I, \hat{H}^c_I)$ and (H_I, H^c_I), which are related to each other by

$$\begin{aligned} \hat{H}_I &= \mathrm{e}^{\left(\frac{1}{2} - c_I \right) T k |y|} H_I, \\ \hat{H}^c_I &= \mathrm{e}^{\left(\frac{1}{2} + c_I \right) T k |y|} H^c_I. \end{aligned}$$

As the theory is orbifolded by $Z_2 : y \to -y$, all 5D fields should have a definite boundary condition under Z_2. The 5D SUGRA multiplet is assumed to have the standard boundary condition leaving the 4D $N = 1$ SUSY unbroken. To give a massless 4D gauge multiplet, the vector superfield $\mathcal{V}^a$ is required to be Z_2-even:

$$\mathcal{V}^a(-y) = \mathcal{V}^a(y). \tag{5}$$

On the other hand, the hypermultiplet can have any Z_2-boundary condition:

$$\begin{aligned} H_I(-y) &= z_I H_I(y), \\ H^c_I(-y) &= -z_I H_I(y), \end{aligned} \tag{6}$$

where $z_I = \pm 1$, and then a massless 4D chiral superfield Q_I originates from either H_I $(z_I = 1)$ or H^c_I $(z_I = -1)$.

In the superfield basis of (H_I, H^c_I), the Ith 4D superfield Q_I appears as the *y-independent* mode of (H_I, H^c_I).

[1] Note that our H_I and H^c_I differ from [8] which is related to ours by $H_I \to \mathrm{e}^{-\left(\frac{3}{2} - c_I \right) T k |y|} H_I$ and $H^c_I \to \mathrm{e}^{-\left(\frac{3}{2} + c_I \right) T k |y|} H^c_I$.

However, in another superfield basis of $(\hat{H}_I, \hat{H}_I^c)$ for which the 5D action (4) has the standard form of a Kähler metric, the wavefunction of Q_I appears to be *quasi-localized* near one of the orbifold fixed points:

$$\hat{H}_I \sim \mathrm{e}^{-q_I T k|y|} \quad (z_I = 1) \quad \text{or} \quad \hat{H}_I^c \sim \mathrm{e}^{-q_I T k|y|} \quad (z_I = -1), \tag{7}$$

where

$$q_I = z_I c_I - \frac{1}{2}.$$

So, Q_I is quasi-localized near $y = 0$ if $q_I > 0$, while it is quasi-localized at $y = \pi$ if $q_I < 0$. As is well known, such a quasi-localization of the matter zero modes can naturally generate the hierarchical Yukawa couplings. It is in fact possible to achieve the quasi-localization of the matter zero modes in a flat extra dimension [13,15]. However in a 5D SUGRA context, the quasi-localization of the matter zero modes has the same origin as the quasi-localization of the gravity zero mode, i.e. the gauging of the graviphoton as (2). In this sense, the warped extra dimension with localized 4D gravity can be considered as a more natural ground than the flat extra dimension for the dynamical localization of the matter zero modes. Furthermore, the quasi-localization of the matter zero modes in AdS_5 has an interpretation in terms of supersymmetric 4D CFT as the renormalization group evolution at a superconformal fixed point [6].

In addition to the bulk action (4), there can be brane actions at the fixed points $y = 0, \pi$. General covariance requires that the 4D metric in the brane action should be the 4D component of the 5D metric at the fixed point. Using the general covariance and also the R- and B_5-dependent field redefinitions which have been made to construct $N = 1$ superfields, one can easily find the T-dependence of the brane actions [8, 16, 17]. For instance, the brane actions which would be relevant for Yukawa couplings and soft parameters are given by[2]

$$\begin{aligned} S_{\text{brane}} = \int \mathrm{d}^5 x \Bigg[&\delta(y) \left\{ \int \mathrm{d}^4\theta L_{I\bar{J}}(Z, Z^*) \Phi_I \Phi_J^* \right. \\ &+ \left. \left(\int \mathrm{d}^2\theta \frac{1}{4} \omega_a(Z) W^{a\alpha} W^a_\alpha + \lambda_{IJK}(Z) \Phi_I \Phi_J \Phi_K + \text{h.c.} \right) \right\} \\ &+ \delta(y - \pi) \left\{ \int \mathrm{d}^4\theta \mathrm{e}^{-(q_I \pi k T + q_J \pi k T^*)} L'_{I\bar{J}}(Z', Z'^*) \Phi_I \Phi_J^* \right. \\ &+ \left(\int \mathrm{d}^2\theta \frac{1}{4} \omega'_a(Z') W^{a\alpha} W^a_\alpha \right. \\ &+ \left. \left. \mathrm{e}^{-(q_I + q_J + q_K)\pi k T} \lambda'_{IJK}(Z') \Phi_I \Phi_J \Phi_K + \text{h.c.} \right) \right\} \Bigg], \end{aligned} \tag{8}$$

where Z and Z' denote generic 4D chiral superfields living *only* on the brane at $y = 0$ and $y = \pi$, respectively, and

$$\Phi_I = H_I \quad (z_I = 1) \quad \text{or} \quad \Phi_I = H_I^c \quad (z_I = -1).$$

[2] The chiral anomaly of the R- and B_5-dependent field redefinition induces T-dependent pieces in ω_a and ω'_a [12]. But they are loop suppressed and not very relevant for the discussion in this paper.

Here $L_{I\bar{J}}$ ($L'_{I\bar{J}}$) are generic hermitian functions of Z and Z^* (Z' and Z'^*), and ω_a and λ_{IJK} (ω' and λ'_{IJK}) are generic holomorphic functions of Z (Z').

If there is no light gauge-singlet 5D field other than the minimal 5D SUGRA multiplet, SUSY breaking would be mediated by the radion superfield T, and/or the brane superfields Z, Z', and/or the 4D SUGRA multiplet which always participate in the mediation of SUSY breaking through the conformal anomaly [18]. The relative importance of anomaly mediation (compared to radion mediation) depends on the details of radion stabilization. Although anomaly mediation can give an important (or even dominant) contribution to soft parameters in some cases [19]; they are loop suppressed, so subleading, unless the F-component of the conformal compensator Ω is much bigger than the F-components of T, Z and Z'. In this paper, we will assume F^Ω is either comparable to or less than F^T, F^Z and $F^{Z'}$, and we thus focus on the soft SUSY breaking parameters induced by F^T, F^Z and $F^{Z'}$.

The 4D Yukawa couplings and soft parameters can be most easily studied by constructing the effective action of massless 4D superfields. As we have noted, in the superfield basis of $(\mathcal{V}^a, H_I, H_I^c)$, the *y-independent* constant modes correspond to the massless 4D superfields of the model. Let V^a denote the constant mode of $\mathcal{V}^a$, and Q_I be the constant mode of H_I ($z_I = 1$) or of H_I^c ($z_I = -1$). Here we will assume that all visible 4D gauge and matter fields are in $\{V^a, Q_I\}$, and we will examine their Yukawa couplings and soft SUSY breaking parameters when the SUSY breaking is mediated by T and/or Z, Z'. Those Yukawa couplings and soft parameters at the Kaluza–Klein scale $M_{\text{KK}} \approx k\mathrm{e}^{-\pi k R}$ can be evaluated from the 4D effective action which can be written as

$$\begin{aligned} &\left[\int \mathrm{d}^4\theta Y_{I\bar{J}} Q_I Q_J^* \right] \\ &+ \left[\int \mathrm{d}^2\theta \left(\frac{1}{4} f_a W^{a\alpha} W^a_\alpha + \tilde{y}_{IJK} Q_I Q_J Q_K \right) + \text{h.c.} \right], \end{aligned} \tag{9}$$

where $Y_{I\bar{J}}$ are hermitian wavefunction coefficients, f_a are holomorphic gauge kinetic functions, and $\tilde{y}_{IJK}$ are holomorphic Yukawa couplings. Using (4) and (8), we find

$$\begin{aligned} Y_{I\bar{J}} &= Y_I \delta_{IJ} + L_{I\bar{J}}(Z, Z^*) + \frac{L'_{I\bar{J}}(Z', Z'^*)}{\mathrm{e}^{(q_I \pi k T + q_J \pi k T^*)}}, \\ f_a &= \frac{2\pi}{g_{5a}^2} T + \omega_a(Z) + \omega'_a(Z'), \\ \tilde{y}_{IJK} &= \lambda_{IJK}(Z) + \frac{\lambda'_{IJK}(Z')}{\mathrm{e}^{(q_I + q_J + q_K)\pi k T}}, \end{aligned} \tag{10}$$

where

$$Y_I = \frac{1}{q_I k} (1 - \mathrm{e}^{-q_I \pi k (T + T^*)}). \tag{11}$$

Note that 5D SUSY in bulk enforces that the Yukawa couplings of Q_I originate entirely from the brane action (8).

It is straightforward to compute the soft parameters for generic forms of the brane wavefunction coefficients $L_{I\bar{J}}$ and $L'_{I\bar{J}}$ [20]. Here we will assume that $L_{I\bar{J}}$ and $L'_{I\bar{J}}$ are (approximately) invariant under the phase rotation of Z, Z', and also $\langle Z\rangle \ll 1$, $\langle Z'\rangle \ll 1$ in the unit with $M_5 = 1$. Then $L_{I\bar{J}}$ and $L'_{I\bar{J}}$ can be expanded as

$$L_{I\bar{J}}(Z, Z^*) = L^0_{I\bar{J}} - \kappa_{I\bar{J}} ZZ^*,$$
$$L'_{I\bar{J}}(Z', Z'^*) = L'^0_{I\bar{J}} - \kappa'_{I\bar{J}} Z'Z'^*, \tag{12}$$

where $L^0_{I\bar{J}}$, $L'^0_{I\bar{J}}$, $\kappa_{I\bar{J}}$ and $\kappa'_{I\bar{J}}$ are field-independent constants. The effects of the brane kinetic coefficients $L^0_{I\bar{J}}, L'^0_{I\bar{J}}$ in low-energy physics are suppressed by their ratios to the bulk kinetic coefficients:

$$\epsilon_{I\bar{J}} \equiv \frac{L^0_{I\bar{J}}}{\sqrt{Y_I Y_J}} \quad \text{or} \quad \frac{L'^0_{I\bar{J}} \mathrm{e}^{-(q_I+q_J)\pi kR}}{\sqrt{Y_I Y_J}}\,, \tag{13}$$

which correspond to the brane thickness divided by the width of the zero mode wavefunction. Obviously, $\epsilon_{I\bar{J}}$ becomes smaller for a larger orbifold length. It is in fact an attractive possibility that the 5D theory is strongly coupled as $g^2_{5a} \approx 24\pi^3/\Lambda$ and $1/M_5^3 \approx 24\pi^3/\Lambda^3$ at the cutoff scale Λ, for which the orbifold length has the maximum value [21]:

$$\Lambda\pi R \approx 24\pi^3\,.$$

Applying the dimensional analysis rule to the bulk and brane kinetic coefficients in the strong coupling limit [21], one easily finds

$$L^0_{I\bar{J}}, L'^0_{I\bar{J}} \approx \frac{3\pi}{4}\frac{1}{\Lambda}\,,$$

where the matter fields are normalized in such a way that their bulk coefficients are given by (11). If $|q_I|k\pi R \ll 1$, so the matter zero mode Q_I is equally spread over the extra dimension, the brane to bulk ratios are given by $\epsilon_{IJ} \approx 1/\Lambda R$. On the other hand, if $|q_I|k\pi R > 1$, so Q_I is quasi-localized with a wavefunction width $1/|q_I|k$, we have

$$\epsilon_{I\bar{J}} \approx \sqrt{q_I q_J} k L^0_{I\bar{J}} \;\text{ or }\; \sqrt{q_I q_J} k L'^0_{I\bar{J}} \approx \frac{3\pi}{4}\frac{q_I \pi kR}{\Lambda\pi R} \approx 10^{-2}\,, \tag{14}$$

where we choose $|q_I|\pi kR \approx 5$ as a typical value of the quasi-localization factors, which can be inferred from the Yukawa coupling structure given by (18). In the following, we will adopt this strong coupling estimate of the brane to bulk ratios, and ignore $L^0_{I\bar{J}}$ and $L'^0_{I\bar{J}}$, while keeping in mind that they can give corrections suppressed by the factor of $\mathcal{O}(10^{-2})$. However, we will keep the other brane coefficients $\kappa_{I\bar{J}}$ and $\kappa'_{I\bar{J}}$ in the calculation since they can be important for SUSY breaking if F^Z or $F^{Z'}$ is a major source of SUSY breaking.

Let y_{IJK}, M_a, $m^2_{I\bar{J}}$, and A_{IJK} denote the Yukawa couplings, gaugino masses, soft scalar masses, trilinear scalar couplings, respectively, for the *canonically normalized* matter superfields $Q_I = \phi^I + \theta\psi^I + \theta^2 F^I$ and gauginos λ^a which are renormalized at M_{KK}:

$$\frac{1}{2} y_{IJK}\phi_I\psi_J\psi_K - \frac{1}{2}M_a\lambda^a\lambda^a$$
$$-\frac{1}{2}m^2_{I\bar{J}}\phi^I\phi^{J*} - \frac{1}{6}A_{IJK}\phi^I\phi^J\phi^K + \text{h.c.} \tag{15}$$

One then finds from (10):

$$\begin{aligned}
y_{IJK} &= (Y_I Y_J Y_K)^{-1/2}\left(\lambda_{IJK} + \frac{\lambda'_{IJK}}{\mathrm{e}^{(q_I+q_J+q_K)\pi kT}}\right),\\
M_a &= \frac{F^T}{2R} + \frac{1}{2}g_a^2\left(\frac{\partial\omega_a}{\partial Z}F^Z + \frac{\partial\omega'_a}{\partial Z'}F^{Z'}\right),\\
m^2_{I\bar{J}} &= (Y_I Y_J)^{-1/2}\\
&\quad\times\left[\frac{\pi^2 q_I k\delta_{IJ}|F^T|^2}{\mathrm{e}^{q_I\pi k(T+T^*)}-1} + \kappa_{I\bar{J}}\left|F^Z\right|^2\right.\\
&\qquad\left. + \frac{\kappa'_{I\bar{J}}|F^{Z'}|^2}{\mathrm{e}^{(q_I\pi kT + q_J\pi kT^*)}}\right],\\
A_{IJK} &= -(Y_I Y_J Y_K)^{-1/2}\\
&\quad\times\left[F^T\frac{\partial}{\partial T}\ln\left(\frac{\lambda_{IJK}+\lambda'_{IJK}\mathrm{e}^{-(q_I+q_J+q_K)\pi kT}}{Y_I Y_J Y_K}\right)\right.\\
&\quad\times\left(\lambda_{IJK} + \frac{\lambda'_{IJK}}{\mathrm{e}^{(q_I+q_J+q_K)\pi kT}}\right) + F^Z\frac{\partial\lambda_{IJK}}{\partial Z}\\
&\quad\left. + \frac{F^{Z'}}{\mathrm{e}^{(q_I+q_J+q_K)\pi kT}}\frac{\partial\lambda'_{IJK}}{\partial Z'}\right],
\end{aligned}\tag{16}$$

where g_a^2 are 4D gauge couplings, and F^T, F^Z and $F^{Z'}$ denote the F-components of T, Z and Z', respectively.

To be more concrete, let us consider the case that both the Yukawa couplings and the brane SUSY breaking come entirely from λ_{IJK} and F^Z at $y = 0$. We further assume that the wavefunctions of 4D Higgs fields are localized at $y = 0$, which would be necessary to generate the top quark Yukawa coupling of order one. Let Q_H and Q_m denote the Higgs and quark/lepton superfields, respectively, and specify only the quark/lepton flavor indices (m, n) in the Yukawa couplings and trilinear A parameters. We then find

$$\begin{aligned}
y_{mn} &= \frac{\lambda_{mn}}{\sqrt{Y_H Y_m Y_n}},\\
M_a &= \frac{F^T}{2R} + \frac{1}{2}g_a^2\frac{\partial\omega_a}{\partial Z}F^Z,\\
m^2_{m\bar{n}} &= \frac{1}{\sqrt{Y_m Y_n}}\left[\frac{\pi^2 q_m k\delta_{mn}}{\mathrm{e}^{2\pi q_m R}-1}\left|F^T\right|^2 + \kappa_{m\bar{n}}|F^Z|^2\right],\\
A_{mn} &= -y_{mn}\left[F^T\frac{\partial}{\partial T}\ln\left(\frac{1}{Y_H Y_m Y_n}\right) + F^Z\frac{\partial}{\partial Z}\ln(\lambda_{mn})\right],
\end{aligned}\tag{17}$$

where λ_{mn} denote the holomorphic Yukawa couplings in the brane action (8) at $y = 0$, and

$$\begin{aligned}
\frac{1}{\sqrt{Y_H}} &= \left(\frac{1}{q_H k}\left(1-\mathrm{e}^{-q_H\pi k(T+T^*)}\right)\right)^{-1/2}\\
&\approx \sqrt{q_H k} \quad (q_H > 0),
\end{aligned}$$

$$\frac{1}{\sqrt{Y_m}} = \left(\frac{1}{q_m k}\left(1-\mathrm{e}^{-q_m \pi k(T+T^*)}\right)\right)^{-1/2}$$
$$\approx \begin{cases} \sqrt{q_m k} & (q_m > 0), \\ \sqrt{|q_m| k}\mathrm{e}^{-\pi |q_m| kR} & (q_m < 0). \end{cases}$$

As we have noted, Q_m with $q_m < 0$ is quasi-localized at $y = \pi$, while Q_m with $q_m > 0$ is quasi-localized at $y = 0$. As a result, in the case that both the Yukawa couplings and the brane SUSY breaking arise from $y = 0$, the Yukawa couplings, A parameters and squark/slepton masses (renormalized at M_{KK}) involving the light quark/lepton superfields Q_m with $q_m < 0$ are exponentially suppressed as

$$y_{mn} = \mathcal{O}(\mathrm{e}^{-(|q_m|+|q_n|)\pi kR}),$$
$$A_{mn} = \mathcal{O}(\mathrm{e}^{-(|q_m|+|q_n|)\pi kR} M_a),$$
$$m^2_{m\bar{n}} = \mathcal{O}(\mathrm{e}^{-(|q_m|+|q_n|)\pi kR} M_a^2), \quad (18)$$

while those involving Q_m with $q_m > 0$ are unsuppressed. On the other hand, there is no localization of V^a, so no suppression of M_a. Note that although the squark/slepton masses from F^Z are suppressed *only* for Q_m with $q_m < 0$, the squark/slepton masses from F^T are suppressed for *all* Q_m with $q_m \neq 0$. This can easily be understood by noting that the radion mediation has a correspondence with the Scherk–Schwarz SUSY breaking [9]. The Scherk–Schwarz SUSY breaking is due to a twist of the boundary condition, so its effects are (exponentially) suppressed for any quasi-localized mode independently of the position of localization. At any rate, the above results show that potentially dangerous flavor-violating soft parameters are suppressed with an appropriate correlation with the Yukawa coupling suppression, thereby ameliorating the SUSY flavor problem in a natural manner.

In fact, for each model in which both the Yukawa couplings and the brane SUSY breaking arise from $y = 0$, there exists a dual model in which the Yukawa couplings and the brane SUSY breaking arise from $y = \pi$. Once we change the sign of all q_I of the involved hypermultiplets, the dual model gives the same 4D Yukawa couplings and soft parameters. To see this, let us consider the case that both the Yukawa couplings and the brane SUSY breaking come entirely from λ'_{IJK} and $F^{Z'}$ at $y = \pi$. It is then straightforward to find

$$y_{mn} = \frac{\mathrm{e}^{-(q_H+q_m+q_n)\pi kR}\lambda'_{mn}}{\sqrt{Y_H Y_m Y_n}},$$
$$M_a = \frac{F^T}{2R} + \frac{1}{2} g_a^2 \frac{\partial \omega'_a}{\partial Z'} F^{Z'},$$
$$m^2_{m\bar{n}} = \frac{1}{\sqrt{Y_m Y_n}} \times \left[\frac{\pi^2 q_m k \delta_{mn}}{\mathrm{e}^{2\pi q_m R}-1}|F^T|^2 + \kappa'_{m\bar{n}} \mathrm{e}^{-(q_m+q_n)\pi kR}|F^{Z'}|^2\right],$$
$$A_{mn} = -y_{mn} \times \left[F^T \frac{\partial}{\partial T}\ln\left(\frac{\mathrm{e}^{-(q_H+q_m+q_n)\pi kT}}{Y_H Y_m Y_n}\right) + F^{Z'}\frac{\partial}{\partial Z'}\ln(\lambda'_{mn})\right]. \quad (19)$$

It is rather obvious that these results give (17) when (λ'_{mn}, $\kappa'_{m\bar{n}}$, Z', q_H, q_m) are replaced by (λ_{mn}, $\kappa_{m\bar{n}}$, Z, $-q_H$, $-q_m$). So in this case, we need to localize the Higgs fields at $y = \pi$ by assuming $q_H < 0$ in order to have a large top quark Yukawa coupling, and localize the light quark/lepton fields at $y = 0$ by assuming $q_m > 0$ in order to suppress the Yukawa couplings and flavor-violating soft parameters.

The results of (17) and (19) show that if the brane SUSY breaking takes place *only* at one of the orbifold fixed points from which the Yukawa couplings arise, the flavor-violating soft parameters of light squark/slepton generations which are induced by F^T and *one* of $F^Z, F^{Z'}$ are suppressed as (18). So, the AdS_5 models with quasi-localized bulk matter fields can avoid (or at least ameliorate) the SUSY flavor problem while generating the hierarchical Yukawa couplings[3]. Note that all of our results can be applied also to models with a *flat* extra dimension which generate the hierarchical Yukawa couplings by quasi-localizing the bulk matter fields. For this, one can simply take the limit $k \to 0$ while keeping $q_I k$ non-zero. In the next section, we will consider the radion-dominated SUSY breaking scenario in this framework, which provides a concrete prediction on soft parameters, and present some models which can pass the constraints on flavor violation without severe fine tuning of parameters.

The results (16), (17) and (19) are obtained from the 4D effective action (9) without including the possible threshold corrections due to massive KK modes. The (exponential) suppression of y_{mn} and A_{mn} for light generations is expected to be stable against KK threshold corrections as it is due to the quasi-localization of the zero modes in the extra dimension. In the 4D effective SUGRA point of view, the suppression of y_{mn} and A_{mn} is a consequence of an exponentially large Kähler metric, $Y_{m\bar{n}} \approx \mathrm{e}^{2\pi|q_m|kR}\delta_{mn}/|q_m|k$, for the case of (17), or of an exponentially small holomorphic Yukawa couplings, $\tilde{y}_{mn} = \mathrm{e}^{-(|q_m|+|q_n|)\pi kR}\lambda'_{mn}$, for the case of (19). These features also suggest that the suppressions of y_{mn} and A_{mn} are stable against KK threshold corrections.

On the other hand, $m^2_{m\bar{n}}$ generically get threshold corrections of order $M_a^2/8\pi^2$. However the geometric interpretation for the suppressed flavor-violating soft parameters suggets that the flavor-violating part of the KK threshold correction to $m^2_{m\bar{n}}$ is suppressed also by $\mathrm{e}^{-(|q_m|+|q_n|)\pi kR}$ as well as by the loop factor, which means that the KK threshold correction gives just a subleading piece of flavor violation in $m^2_{m\bar{n}}$. There can be additional corrections to the soft parameters which are induced by non-renormalizable SUGRA interactions in 4D effective action [23], but they are suppressed by $M^2_{\mathrm{KK}}/8\pi^2 M^2_{\mathrm{Pl}} \sim \mathrm{e}^{-2\pi kR}/8\pi^2$, so they are small enough. There are also the model-independent SUSY breaking effects of $\mathcal{O}(F^{\Omega}/8\pi^2)$ mediated by the 4D superconformal anomaly [18], where F^{Ω} is the F-component of the chiral compensator superfield Ω. If F^{Ω} is comparable to or less than $F^T, F^Z, F^{Z'}$, as we have assumed, one can

[3] Another way to ameliorate the SUSY flavor problem using an extra dimension has been suggested in [22].

ignore the anomaly-mediated contributions also, and then the leading radiative corrections to Yukawa couplings and soft parameters come from the standard renormalization group running down to M_W starting from the boundary values at $M_{\rm KK}$ given by (17) or (19).

We note that the idea of the AdS/CFT correspondence suggests a CFT framework which would reproduce the main features of our AdS models. Indeed, models involving a superconformal sector have been proposed to generate hierarchical Yukawa couplings as well as exponentially suppressed soft masses [6,19,24]. It is then easy to see the correspondence: $q_I \pi kR \to \gamma_I \ln(\Lambda/M_c)$, where γ_I is the anomalous dimension of Q_I driven by the coupling to the SC sector, and Λ and M_c are the cutoff scale and the decoupling scale of the superconformal sector, respectively. The AdS_5 models discussed here provide a perturbative framework to generate the hierarchical Yukawa couplings while suppressing the dangerous flavor-violating soft parameters.

3 Radion-dominated scenario

The results of (17) and (19) show that potentially dangerous flavor-violating soft parameters can be naturally suppressed in AdS_5 models. However, the resulting soft parameters still involve many adjustable free parameters, particularly in the contributions from the brane SUSY breaking. Obviously, the model becomes much more predictive in the radion-dominated scenario [7] which we will discuss in some detail in this section.

To proceed, let us consider the case that the Yukawa couplings come from the brane action at $y = 0$ and the Higgs zero modes are localized at $y = 0$. To be definite, we assume that the Kaluza–Klein scale is given by $M_{\rm KK} = k\mathrm{e}^{-\pi kR} \sim 10^{16}\,\mathrm{GeV}$; however, the results are not sensitive to the precise value of $M_{\rm KK}$. In the radion-dominated scenario for this case, the Yukawa couplings and soft parameters renormalized at $M_{\rm KK}$ are given by

$$\begin{aligned} y_{mn} &= \frac{\lambda_{mn}\ln(1/\epsilon)}{\pi R\sqrt{Y_H}}\sqrt{\frac{\phi_m\phi_n}{(1-\epsilon^{2\phi_m})(1-\epsilon^{2\phi_n})}}, \\ M_a &= \frac{F^T}{2R}, \\ A_{mn} &= 2y_{mn}\ln(1/\epsilon) \\ &\quad\times\left(\frac{\phi_m}{\epsilon^{-2\phi_m}-1}+\frac{\phi_n}{\epsilon^{-2\phi_n}-1}+\frac{\phi_H}{\epsilon^{-2\phi_H}-1}\right)\frac{F^T}{2R}, \\ m^2_{m\bar{n}} &= \delta_{mn}\left(2\ln(1/\epsilon)\frac{\phi_m}{\epsilon^{\phi_m}-\epsilon^{-\phi_m}}\left|\frac{F^T}{2R}\right|\right)^2, \qquad (20) \end{aligned}$$

where $\epsilon \approx 0.2$ denotes the Cabibbo angle, $Y_H \approx 1/q_H k$, and

$$\begin{aligned} \phi_m &= \frac{q_m \pi kR}{\ln(1/\epsilon)} = \left(z_m c_m - \frac{1}{2}\right)\frac{\pi kR}{\ln(1/\epsilon)}, \\ \phi_H &= \frac{q_H \pi kR}{\ln(1/\epsilon)} = \left(z_H c_H - \frac{1}{2}\right)\frac{\pi kR}{\ln(1/\epsilon)}, \end{aligned}$$

for the kink mass $c_I k$ and the Z_2-boundary condition factor z_I $(=\pm 1)$ for the 5D hypermultiplets (H_I, H_I^c) whose zero modes correspond to the quark/lepton superfields $(I = m)$ or the Higgs superfields $(I = H)$. As we have noted, c_I and thus ϕ_I are likely to be quantized. The observed quark and lepton masses indicate that ϕ_m are integers for $\epsilon \approx 0.2$, so that all Yukawa couplings are given by an integer power of ϵ in their order of magnitudes.

Let $\psi_m = \{q_i, u_i, d_i, \ell_i, e_i\}$ $(i = 1,2,3)$ denote the known three generations of the left-handed quark doublets (q_i), up-type antiquark singlets (u_i), down-type antiquark singlets (d_i), lepton doublets (ℓ_i), and anti-lepton singlets (e_i). The Yukawa couplings can be written as

$$\mathcal{L}_{\rm Yukawa} = y^u_{ij}H_2 q_i u_j + y^d_{ij}H_1 q_i d_j + y^\ell_{ij}H_1 \ell_i e_j, \quad (21)$$

and the squark/sleptons $\varphi_m = \{\tilde{q}_i, \tilde{u}_i, \tilde{d}_i, \tilde{\ell}_i, \tilde{e}_i\}$ have the soft SUSY breaking couplings:

$$\begin{aligned} \mathcal{L}_{\rm soft} = -\Big(&A^u_{ij}H_2\tilde{q}_i\tilde{u}_j + A^d_{ij}H_1\tilde{q}_i\tilde{d}_j + A^\ell_{ij}H_1\tilde{\ell}_i\tilde{e}_j \\ &+ m^{2(\tilde{q})}_{i\bar{j}}\tilde{q}_i\tilde{q}^*_j + m^{2(\tilde{u})}_{i\bar{j}}\tilde{u}_i\tilde{u}^*_j + m^{2(\tilde{d})}_{i\bar{j}}\tilde{d}_i\tilde{d}^*_j \\ &+ m^{2(\tilde{\ell})}_{i\bar{j}}\tilde{\ell}_i\tilde{\ell}^*_j + m^{2(\tilde{e})}_{i\bar{j}}\tilde{e}_i\tilde{e}^*_j\Big). \qquad (22) \end{aligned}$$

Let us first consider the quark/squark sector. There can be several different choices of $\phi_m \equiv \phi(Q_m)$ which would yield the observed quark masses and mixing angles [25]. Here we will consider one example which would belong to the class of models with minimal flavor violation:

$$\begin{aligned} &\phi(H_1) = \phi(H_2) = 2, \\ &\phi(q_i) = (-3,-2,2), \\ &\phi(u_i) = (-5,-2,2), \\ &\phi(d_i) = (-3,-3,-2), \qquad (23) \end{aligned}$$

for which $\tan\beta \sim \epsilon^2 m_t/m_b$ is not large. These values of $\phi(Q_m)$ give the following forms of the Yukawa coupling matrices:

$$\begin{aligned} y^u_{ij} &= \begin{pmatrix} \epsilon^8\lambda^u_{11} & \epsilon^5\lambda^u_{12} & \epsilon^3\lambda^u_{13} \\ \epsilon^7\lambda^u_{21} & \epsilon^4\lambda^u_{22} & \epsilon^2\lambda^u_{23} \\ \epsilon^5\lambda^u_{31} & \epsilon^2\lambda^u_{32} & \lambda^u_{33} \end{pmatrix}, \\ y^d_{ij} &= \begin{pmatrix} \epsilon^6\lambda^d_{11} & \epsilon^6\lambda^d_{12} & \epsilon^5\lambda^d_{13} \\ \epsilon^5\lambda^d_{21} & \epsilon^5\lambda^d_{22} & \epsilon^4\lambda^d_{23} \\ \epsilon^3\lambda^d_{31} & \epsilon^3\lambda^d_{32} & \epsilon^2\lambda^d_{33} \end{pmatrix}, \qquad (24) \end{aligned}$$

where λ^u_{ij} and λ^d_{ij} are the coefficients of order unity[4]. The soft parameters renormalized at $M_{\rm KK}$ are determined to be[5]

$$\frac{A^u_{ij}}{y^u_{ij}} = 2M_{1/2}\ln 5\begin{pmatrix} 8 & 5 & 3 \\ 7 & 4 & 2 \\ 5 & 2 & 6\epsilon^4 \end{pmatrix},$$

[4] Here λ^u_{ij} and λ^d_{ij} are redefined from λ_{mn} in (20) to include the coefficients of order one, depending on $\ln(1/\epsilon)/\pi R, \phi(q_i)$, $\phi(u_i)$, and $\phi(d_i)$.

[5] Here we have ignored small corrections suppressed by high powers of ϵ.

$$\frac{A^d_{ij}}{y^d_{ij}} = 2M_{1/2}\ln 5 \begin{pmatrix} 6 & 6 & 5 \\ 5 & 5 & 4 \\ 3 & 3 & 2 \end{pmatrix} ,$$

$$m^{2(\tilde{q})}_{i\bar{j}} = (2\ln 5)^2 |M_{1/2}|^2 \begin{pmatrix} 9\epsilon^6 & 0 & 0 \\ 0 & 4\epsilon^4 & 0 \\ 0 & 0 & 4\epsilon^4 \end{pmatrix}$$

$$\approx |M_{1/2}|^2 \begin{pmatrix} 6\times 10^{-3} & 0 & 0 \\ 0 & 6\times 10^{-2} & 0 \\ 0 & 0 & 6\times 10^{-2} \end{pmatrix} ,$$

$$m^{2(\tilde{u})}_{i\bar{j}} = (2\ln 5)^2 |M_{1/2}|^2 \begin{pmatrix} 25\epsilon^{10} & 0 & 0 \\ 0 & 4\epsilon^4 & 0 \\ 0 & 0 & 4\epsilon^4 \end{pmatrix}$$

$$\approx |M_{1/2}|^2 \begin{pmatrix} 3\times 10^{-5} & 0 & 0 \\ 0 & 6\times 10^{-2} & 0 \\ 0 & 0 & 6\times 10^{-2} \end{pmatrix} ,$$

$$m^{2(\tilde{d})}_{i\bar{j}} = (2\ln 5)^2 |M_{1/2}|^2 \begin{pmatrix} 9\epsilon^6 & 0 & \\ 0 & 9\epsilon^6 & 0 \\ 0 & 0 & 4\epsilon^4 \end{pmatrix} \quad (25)$$

$$\approx |M_{1/2}|^2 \begin{pmatrix} 6\times 10^{-3} & 0 & 0 \\ 0 & 6\times 10^{-3} & 0 \\ 0 & 0 & 6\times 10^{-2} \end{pmatrix} ,$$

where $M_{1/2} = F^T/2R$ denotes the universal gaugino mass at $M_{\rm KK}$.

The above Yukawa coupling matrices produce well the observed quark masses and mixing angles. Also, after taking into account the renormalization group evolution from $M_{\rm KK}$ to the weak scale M_W, the resulting soft parameters pass all phenomenological constraints (including those from flavor-changing processes) for a reasonable range of $M_{1/2}$, λ^u_{ij} and λ^d_{ij}. For instance, when λ^u_{ij} and λ^d_{ij} are taken to be of order unity for $M_{1/2} = 500\,\text{GeV}$, the resulting SUSY contributions to $B\bar{B}$ mixing, $D\bar{D}$ mixing, and $b \to s\gamma$ are well below the SM values [20]. For the same parameter values, the SUSY contribution to $K\bar{K}$ mixing is comparable to the SM contribution; however, it can be easily made small enough by choosing a bit larger $M_{1/2}$ and/or a bit smaller $\lambda^d_{12,21}$ [20].

One distinctive feature of the radion-mediated SUSY breaking for quasi-localized matter superfields is that it gives a rather large value of the *non-universal* A_{mn}/y_{mn} which is of the order of $M_{1/2}\ln y_{mn}$. As a consequence, the most stringent constraint on the model comes from $\mu \to e\gamma$ which would be induced by the slepton–A coupling: $A^\ell_{ij} H_1 \tilde{\ell}_i \tilde{e}_j$. To satisfy the experimental bound on Br($\mu \to e\gamma$) for generic models of radion mediation, one would need some degree of fine tuning for some of the off-diagonal elements of the lepton Yukawa matrix *unless* the slepton and gaugino masses are heavier than $\mathcal{O}(1)$ TeV. Another distinctive feature of radion mediation is that although the A_{mn}/y_{mn} are highly flavor dependent, they are quantized up to small corrections being a high power of ϵ (see (20)). This feature provides a natural way to minimize the $\mu \to e\gamma$ rate, which is to choose

$$\phi(\ell_1) = \phi(\ell_2) \quad \text{or} \quad \phi(e_1) = \phi(e_2). \quad (26)$$

Let us thus consider the following two examples for the lepton/slepton sector:

$$\begin{aligned} \text{Model I}: \quad & \phi(\ell_i) = (-3,-3,-1), \\ & \phi(e_i) = (-4,-1,-1), \\ \text{Model II}: \quad & \phi(\ell_i) = (-5,-2,-1), \\ & \phi(e_i) = (-2,-2,-1). \end{aligned} \quad (27)$$

The model I gives at $M_{\rm KK}$

$$y^\ell_{ij} = \begin{pmatrix} \epsilon^7\lambda^\ell_{11} & \epsilon^4\lambda^\ell_{12} & \epsilon^4\lambda^\ell_{13} \\ \epsilon^7\lambda^\ell_{21} & \epsilon^4\lambda^\ell_{22} & \epsilon^4\lambda^\ell_{23} \\ \epsilon^5\lambda^\ell_{31} & \epsilon^2\lambda^\ell_{32} & \epsilon^2\lambda^\ell_{33} \end{pmatrix} ,$$

$$\frac{A^\ell_{ij}}{y^\ell_{ij}} = 2M_{1/2}\ln 5 \begin{pmatrix} 7 & 4 & 4 \\ 7 & 4 & 4 \\ 5 & 2 & 2 \end{pmatrix} ,$$

$$m^{2(\tilde{\ell})}_{i\bar{j}} = (2\ln 5)^2 |M_{1/2}|^2 \begin{pmatrix} 9\epsilon^6 & 0 & 0 \\ 0 & 9\epsilon^6 & 0 \\ 0 & 0 & \epsilon^2 \end{pmatrix}$$

$$\approx \left|M_{1/2}\right|^2 \begin{pmatrix} 6\times 10^{-3} & 0 & 0 \\ 0 & 6\times 10^{-3} & 0 \\ 0 & 0 & 0.4 \end{pmatrix}$$

$$m^{2(\tilde{e})}_{i\bar{j}} = (2\ln 5)^2 |M_{1/2}|^2 \begin{pmatrix} 16\epsilon^8 & 0 & 0 \\ 0 & \epsilon^2 & 0 \\ 0 & 0 & \epsilon^2 \end{pmatrix}$$

$$\approx \left|M_{1/2}\right|^2 \begin{pmatrix} 4\times 10^{-4} & 0 & 0 \\ 0 & 0.4 & 0 \\ 0 & 0 & 0.4 \end{pmatrix} , \quad (28)$$

while the model II gives

$$y^\ell_{ij} = \begin{pmatrix} \epsilon^7\lambda^\ell_{11} & \epsilon^7\lambda^\ell_{12} & \epsilon^6\lambda^\ell_{13} \\ \epsilon^4\lambda^\ell_{21} & \epsilon^4\lambda^\ell_{22} & \epsilon^3\lambda^\ell_{23} \\ \epsilon^3\lambda^\ell_{31} & \epsilon^3\lambda^\ell_{32} & \epsilon^2\lambda^\ell_{33} \end{pmatrix} ,$$

$$\frac{A^\ell_{ij}}{y^\ell_{ij}} = 2M_{1/2}\ln 5 \begin{pmatrix} 7 & 7 & 6 \\ 4 & 4 & 3 \\ 3 & 3 & 2 \end{pmatrix} ,$$

$$m^{2(\tilde{\ell})}_{i\bar{j}} = (2\ln 5)^2 |M_{1/2}|^2 \begin{pmatrix} 25\epsilon^{10} & 0 & 0 \\ 0 & 4\epsilon^4 & 0 \\ 0 & 0 & \epsilon^2 \end{pmatrix} \quad (29)$$

$$\approx \left|M_{1/2}\right|^2 \begin{pmatrix} 3\times 10^{-5} & 0 & 0 \\ 0 & 6\times 10^{-2} & 0 \\ 0 & 0 & 0.4 \end{pmatrix} ,$$

$$m^{2(\tilde{e})}_{i\bar{j}} = (2\ln 5)^2|M_{1/2}|^2 \begin{pmatrix} 4\epsilon^4 & 0 & 0 \\ 0 & 4\epsilon^4 & 0 \\ 0 & 0 & \epsilon^2 \end{pmatrix}$$
$$\approx \left|M_{1/2}\right|^2 \begin{pmatrix} 6\times 10^{-2} & 0 & 0 \\ 0 & 6\times 10^{-2} & 0 \\ 0 & 0 & 0.4 \end{pmatrix}.$$

When all λ^ℓ_{ij} are taken to be of order unity, these two models give a too large $\mu \to e\gamma$ rate, though they pass all other phenomenological constraints. To make the models consistent with $\mu \to e\gamma$, one needs some degree of fine tuning for λ^ℓ_{12} and λ^ℓ_{21} defined in (28) and (29). Through a detailed numerical analysis including the renormalization group evolution of the soft parameters from $M_{\mathrm{KK}} \sim 10^{16}\,\mathrm{GeV}$ to M_W, we find that λ^ℓ_{12} and λ^ℓ_{21} are required to satisfy

$$\text{Model I}: \quad \lambda^\ell_{12} \lesssim 2\times 10^{-2}\left(\frac{M_{1/2}}{500\,\mathrm{GeV}}\right)^2,$$
$$\lambda^\ell_{21} \lesssim 10^{-1}\left(\frac{M_{1/2}}{500\,\mathrm{GeV}}\right)^2,$$
$$\text{Model II}: \quad \lambda^\ell_{12} \lesssim 5\times 10^{-2}\left(\frac{M_{1/2}}{500\,\mathrm{GeV}}\right)^2,$$
$$\lambda^\ell_{21} \lesssim 10^{-2}\left(\frac{M_{1/2}}{500\,\mathrm{GeV}}\right)^2. \quad (30)$$

The above constraint from $\mu \to e\gamma$ suggests that the holomorphic brane Yukawa couplings of leptons at $y = 0$ conserve, at least approximately, some of the lepton flavors, e.g. $L_e - L_\tau$, so

$$\lambda^\ell_{ij} = \lambda^\ell_i \delta_{ij}.$$

Still one can achieve large flavor mixing in the neutrino mass matrix by introducing lepton flavor violation at the other fixed point $y = \pi$. For instance, the gauge-singlet right-handed neutrinos N_i can have *flavor diagonal* Yukawa couplings

$$\delta(y)\kappa_i H_2 \ell_i N_i$$

at $y = 0$, but *flavor non-diagonal* Majorana mass masses

$$\delta(y-\pi) M_{ij} N_i N_j$$

at $y = \pi$. In this case, it is straightforward to adjust $\phi(N_i)$ to make the resulting light neutrino mass matrix take a form which explains the observed large neutrino mixings [20].

So far, we have ignored the effects of brane kinetic terms

$$\int \mathrm{d}^5x\,\mathrm{d}^4\theta \quad (31)$$
$$\times \left[\delta(y) L^0_{m\bar{n}} + \delta(y-\pi)\mathrm{e}^{-(q_m\pi kT + q_{\bar{n}}\pi kT^*)} L'^0_{m\bar{n}}\right] \Phi_m \Phi^*_{\bar{n}}$$

under the assumption that the orbifold length is large enough, so that the brane to bulk kinetic term ratios are small enough:

$$\epsilon_{m\bar{n}} \equiv \frac{L^0_{m\bar{n}}}{\sqrt{Y_m Y_n}} \;\text{ or }\; \frac{L'^0_{m\bar{n}}\mathrm{e}^{-(q_m+q_n)\pi kR}}{\sqrt{Y_n Y_m}} \approx 10^{-2}, \quad (32)$$

which is the case if the 5D theory is strongly coupled at the cutoff scale Λ [21]. (See (14).) Once included, the brane kinetic terms can provide additional flavor violation. The precise form of the soft parameters at M_{KK} including the effects of brane kinetic terms is somewhat involved [20]; however, the size of corrections can be easily estimated to be

$$\delta m^2_{m\bar{n}} = \mathcal{O}(\epsilon_{m\bar{k}} m^2_{k\bar{n}}) \quad \text{and/or} \quad \mathcal{O}(\epsilon_{k\bar{n}} m^2_{m\bar{k}}),$$
$$\delta A_{mn} = \mathcal{O}(\epsilon_{m\bar{k}} A_{kn}) \quad \text{and/or} \quad \mathcal{O}(\epsilon_{n\bar{k}} A_{mk}), \quad (33)$$

where $m^2_{m\bar{n}}$ and A_{mn} are the soft parameters of (20) in the limit $L^0_{m\bar{n}} = L'^0_{m\bar{n}} = 0$. If the brane kinetic coefficients are comparable to the bulk kinetic coefficients, the model gives too large $K\bar{K}$ mixing and $\mu \to e\gamma$ rates [20]. However in the large radius (strong coupling) limit with $\epsilon_{m\bar{n}} \approx 10^{-2}$, the brane kinetic coefficients do not lead to dangerous flavor violation for reasonable ranges of the other parameters.

4 Conclusion

In this paper, we have examined the soft SUSY breaking parameters in supersymmetric theories on a slice of AdS_5 in which the hierarchical Yukawa couplings are generated by quasi-localizing the bulk matter fields in the extra dimension. In this class of models, the quasi-localization of the matter zero modes has a common origin with the quasi-localization of the gravity zero mode, i.e. the gauging of the graviphoton which gives a non-zero bulk cosmological constant and hypermultiplet masses. As a result, generically the small Yukawa couplings are given by $y_{mn} \sim \mathrm{e}^{-c\pi kR}$ for a constant c of order unity, and so the radion R is required to be stabilized at $\mathrm{e}^{-\pi kR} \approx 10^{-2}$–$10^{-5}$. Unless there is an additional gauge-singlet 5D field other than the minimal 5D SUGRA multiplet, the prime candidates for the messenger of SUSY breaking is the radion superfield in bulk and/or some brane superfields at $y = 0$ and π. We found that if the brane SUSY breaking takes place *only* at one of the orbifold fixed points from which the Yukawa couplings arise, which is a natural setting; the dangerous flavor-violating soft parameters are suppressed with an appropriate correlation with the Yukawa suppression, thereby the model ameliorates (or avoids) the SUSY flavor problem. This type of models can be considered as the AdS dual of the recently studied 4D models containing a supersymmetric CFT to generate hierarchical Yukawa couplings. If SUSY breaking is dominated by the radion mediation, the model provides a concrete prediction on soft parameters which can be tested by low-energy experiments. We have presented some models in this framework which pass all phenomenological constraints from flavor-violating processes without severe fine tuning of parameters.

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References

1. L. Randall, R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999)
2. T. Gherghett a, A. Pomarol, Nucl. Phys. B **586**, 141 (2000) [hep-ph/0003129]
3. S.J. Huber, Q. Shafi, Phys. Lett. B **498**, 256 (2001) [hep-ph/0010195]
4. Y. Grossman, M. Neubert, Phys. Lett. B **474**, 361 (2000) [hep-ph/9912408]
5. This subject has been discussed briefly in D. Marti, A. Pomarol, Phys. Rev. D **64**, 105025 (2001); K. Choi, Do Young Kim, Ian-Woo Kim, T. Kobayashi, hep-ph/0301131; SUSY breaking in warped geometry in different context has been discussed in Z. Chacko, E. Ponton, hep-ph/0301171; L.J. Hall, Y. Nomura, T. Okui, S.J. Oliver, hep-th/0302192
6. A.E. Nelson, M.J. Strassler, JHEP **0009**, 030 (2000); JHEP **0207**, 021 (2002); T. Kobayashi, H. Terao, Phys. Rev. D **64**, 075003 (2001); T. Kobayashi, H. Nakano, H. Terao, Phys. Rev. D **65**, 015006 (2002); T. Kobayashi, H. Nakano, T. Noguchi, H. Terao, Phys. Rev. D **66**, 095011 (2002); JHEP **0302**, 022 (2003)
7. Z. Chacko, M.A. Luty, JHEP **0105**, 067 (2001)
8. D. Marti, A. Pomarol, Phys. Rev. D **64**, 105025 (2001)
9. D.E. Kaplan, N. Weiner, hep-ph/0108001; G. von Gersdorff, M. Quiros, Phys. Rev. D **65**, 064016 (2002)
10. A. Ceresole, G. Dall'Agata, Nucl. Phys. B **585**, 143 (2000)
11. R. Altendorfer, J. Bagger, D. Nemeschansky, Phys. Rev. D **63**, 125025 (2001); A. Falkowski, Z. Lalak, S. Pokorski, Phys. Lett. B **491**, 172 (2000); J. Bagger, D. Nemeschansky, R.J. Zhang, JHEP **0108**, 057 (2001); A. Falkowski, Z. Lalak, S. Pokorski, Nucl. Phys. B **613**, 189 (2001)
12. K. Choi, H.D. Kim, I.-W. Kim, JHEP **0211**, 033 (2002); **0303**, 034 (2003); K. Choi, I.-W. Kim, Phys. Rev. D **67**, 045005 (2003)
13. N. Arkani-Hamed, T. Gregoire, J. Wacker, JHEP **0203**, 055 (2002)
14. W.D. Linch, M.A. Luty, J. Phillips, Phys. Rev. D **68**, 025008 (2003)
15. N. Arkani-Hamed, M. Schmaltz, Phys. Rev. D **61**, 033005 (2000); E.A. Mirabelli, M. Schmaltz, Phys. Rev. D **61**, 113011 (2000); D.E. Kaplan, T.M. Tait, JHEP **0111**, 051 (2001); M. Kakizaki, M. Yamaguchi, hep-ph/0110266; N. Haba, N. Maru, Phys. Rev. D **66**, 055005 (2002); A. Hebecker, J. March-Russell, Phys. Lett. B **541**, 338 (2002)
16. J.A. Casas, J.R. Espinosa, I. Navarro, Nucl. Phys. B **620**, 195 (2002)
17. A. Hebecker, Nucl. Phys. B **632**, 101 (2002)
18. L. Randall, R. Sundrum, Nucl. Phys. B **557**, 79 (1999); G.F. Giudice, M.A. Luty, H. Murayama, R. Rattazzi, JHEP **9812**, 027 (1998)
19. M.A. Luty, R. Sundrum, Phys. Rev. D **65**, 066004 (2002); **67**, 045007 (2003)
20. K. Choi, K.S. Jeong, K. Okumura (in preparation)
21. Z. Chacko, M.A. Luty, E. Ponton, JHEP **0007**, 036 (2000); Y. Nomura, Phys. Rev. D **65**, 085036 (2002)
22. J. Kubo, H. Terao, Phys. Rev. D **66**, 116003 (2002); K.S. Choi, K.Y. Choi, J.E. Kim, Phys. Rev. D **68**, 035003 (2003)
23. K. Choi, J.S. Lee, C. Munoz, Phys. Rev. Lett. **80**, 3686 (1998)
24. A. Karch, T. Kobayashi, J. Kubo, G. Zoupanos, Phys. Lett. B **441**, 235 (1998)
25. E.J. Chun, A. Lukas, Phys. Lett. B **387**, 99 (1996); K. Choi, E.J. Chun, H.D. Kim, Phys. Lett. B **394**, 89 (1997)