SUSY Flavor Problem and Warped Geometry

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Abstract

We point out that supersymmetric warped geometry can provide a solution to the SUSY flavor problem, while generating hierarchical Yukawa couplings. In supersymmetric theories in a slice of AdS$_5$ with the Kaluza-Klein scale $M_{KK}$ much higher than the weak scale, if all visible fields originate from 5D bulk fields and supersymmetry breaking is mediated by the bulk radion superfield and/or some brane chiral superfields, potentially dangerous soft scalar masses and trilinear $A$ parameters at $M_{KK}$ can be naturally suppressed compared to the gaugino masses by small warp factor. We present simple models yielding phenomenologically interesting patterns of soft parameters in this framework.

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Warped extra dimension can provide a small geometric factor which can be useful to explain 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 warp factor $e^{-\pi kR}$ may be responsible for the 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 the hierarchical quark and lepton masses [2,3], and/or small neutrino masses [4]. In this paper, we wish to examine the implications of warped geometry to the flavor structure of soft SUSY breaking parameters as well as Yukawa couplings. To this end, we consider supersymmetric 5D theory on a slice of AdS$_5$ in which all visible massless 4D modes arise from 5D bulk vector multiplets and hypermultiplets [2]. The Kaluza-Klein (KK) scale $M_{KK} \approx ke^{-\pi kR}$ of the model is assumed to be much higher than $M_W$, e.g. $10^{16} - 10^{13}$ GeV, and then the stability of $M_W$ is achieved by the weak scale supersymmetry (SUSY) as it is in the supersymmetric standard model. If there is no light gauge-singlet 5D field other than the minimal 5D supergravity (SUGRA) multiplet, SUSY breaking is mediated mainly by the radion superfield and/or some brane superfields $^1$. As we will see, in such case AdS$_5$ geometry can naturally suppress the potentially dangerous soft scalar masses and trilinear $A$ parameters at $M_{KK}$, while generating hierarchical Yukawa couplings, for rather generic forms of flavor violation in the underlying theory. This would solve the SUSY flavor problem since the squark and slepton masses at $M_W$ (at least of the first and second generations) are generated mainly by the flavor-independent renormalization group (RG) evolution arising from the unsuppressed gaugino masses. Also this framework using the warped geometry to solve the SUSY flavor problem can provide predictions for the shape of soft parameters at $M_{KK}$, which can lead to interesting phenomenology at $M_W$.

$^1$The 4D SUGRA multiplet always participates in the mediation of SUSY breaking through the conformal anomaly [5]. However as we will see, in the models under consideration its contributions to soft parameters are negligible compared to those from the radion and brane superfields.
Our starting point is generic 5D gauge theory coupled to the minimal 5D SUGRA on $S^1/Z_2$. The action of the model is given by [6,2]

$$S = \int d^5x \sqrt{-G} \left[ \frac{1}{2} \left( R + \Psi_M^i \gamma^{MNP} D_N \Psi_{iP} - \frac{3}{2} C_{MN} C^{MN} + 12k^2 \right) + \frac{1}{g_5^2} \left( -\frac{1}{4} F^{aMN} F^a_{MN} + \frac{1}{2} D_M \phi^a D^a \phi^a + i \lambda^{ia} \gamma^M D_M \lambda^i \right) + |D_M h_I^i|^2 + i \bar{\Psi}_I \gamma^M D_M \Psi_I + i \bar{\tilde{c}}_I k \epsilon(y) \bar{\Psi}_I \Psi_I + ... \right]$$

(1)

where $R$ is the 5D Ricci scalar, $\Psi_M^i$ ($i = 1, 2$) 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 fundamental range $0 \leq y \leq \pi$. Here $\phi^a, A_M^a$ and $\lambda^a$ are 5D scalar, vector and symplectic Majorana spinors constituting a 5D vector multiplet, $h_I^i$ and $\Psi_I$ are 5D scalar and Dirac spinor constituting the $I$-th hypermultiplet with kink mass $\tilde{c}_I k \epsilon(y)$. Note that we set the 5D Planck mass $M_5 = 1$ and all dimensionful parameters, e.g. the 5D gauge coupling $g_5$ and the AdS curvature $k$, are defined in this unit. With appropriate values of the brane cosmological constants at the orbifold fixed points ($y = 0, \pi$), the ground state geometry is given by a slice of AdS$_5$ having the radion ($R$) dependence:

$$G_{MN} dx^M dx^N = e^{-2kRy} \eta_{\mu \nu} dx^\mu dx^\nu + R^2 dy^2.$$  

(2)

It is convenient to write the above 5D action in $N = 1$ superspace [7-9]. For the 5D SUGRA multiplet, we keep only the radion superfield

$$T = \left( R + i B_5, \frac{1}{2} (1 + \gamma_5) \Psi_5^{i=2} \right)$$

and replace 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 [8,10]. After such field redefinition, the relevant piece of the action is given by$^2$ [7,8]

$^2$Note that our superfield basis for hypermultiplets differs from [8] which is related to ours by $H_I \rightarrow e^{-(\frac{3}{2} - \tilde{c}_I)Tk|y|} H_I$ and $H_I^c \rightarrow e^{-(\frac{3}{2} + \tilde{c}_I)Tk|y|} H_I^c$. 

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\[
\int d^5 x \left[ \int d^4 \theta \left\{ \frac{T + T^*}{2} \left( e^{(\frac{i}{2} + \bar{\epsilon}_I)(T + T^*) \frac{1}{2} \psi} H_I H_I^* + e^{(\frac{i}{2} + \bar{\epsilon}_I)(T + T^*) \frac{1}{2} \lambda} \bar{H}_I \bar{H}_I^* \right) \right\} \\
+ \left\{ \int d^2 \theta \frac{1}{4 g_5^2} T W^{a \alpha} W^a \alpha + h.c. \right\} \right],
\]

where \( W^a_\alpha \) is the chiral spinor superfield for the vector superfield \( \mathcal{V}^a \) containing \( (A^a_\mu, \lambda^a) \) with \( \lambda^a = \frac{1}{2} (1 - \gamma_5) \lambda^{a1} \), \( H_I \) and \( H'_I \) are chiral superfields containing \( (h_I^1, \psi_I) \) and \( (h'^2_I, \psi'^2_I) \), respectively, with \( \psi_I = \frac{1}{2} (1 - \gamma_5) \Psi_I \), \( \psi'^2_I = \frac{1}{2} (1 + \gamma_5) \Psi_I \). As the theory is orbifolded by \( Z_2 : y \rightarrow -y \), 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. As for \( \mathcal{V}^a \), one needs \( \mathcal{V}^a(-y) = \mathcal{V}^a(y) \) to obtain massless 4D gauge multiplet. On the other hand, \( H_I(-y) = z_I H_I(y) \) and \( H'_I(-y) = -z_I H_I(y) \) \( (z_I = \pm 1) \), and then massless 4D matter multiplet can originate either from \( H_I \) or \( H'_I \) depending on the sign of \( z_I \).

In addition to the bulk action (1), there can be brane actions at the fixed points \( y = 0, \pi \). The general covariance requires that the 4D metric in brane action should be the 4D component of the 5D metric at 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 brane actions [8,11,12]. For instance, the brane actions which would be relevant for Yukawa couplings and soft parameters are given by\(^3\)

\[
S_{brane} = \int d^5 x \left\{ \delta(y) \left\{ \int d^4 \theta \left( \frac{1}{4} \omega_a(Z)W^{a \alpha}W^a_\alpha + \lambda_{IJK}(Z)\bar{H}_I \bar{H}_J \bar{H}_K \right) + h.c. \right\} \\
+ \delta(y - \pi) \left\{ \int d^4 \theta e^{-c_j \pi k T + c_j \pi k T^*} \left( \lambda_I J(J' \bar{Z}) \bar{H}_I \bar{H}_J \bar{H}_K \right) + h.c. \right\} \right\},
\]

where \( Z \) and \( Z' \) denote generic 4D chiral superfields living only on the brane at \( y = 0 \) and

\(^3\)The chiral anomaly of the \( R \) and \( B_5 \)-dependent field redefinition induces \( T \)-dependent pieces in \( \omega_a \) and \( \omega'_a \)[10]. But they are loop-suppressed and not very relevant for the discussion in this paper.
y = \pi, \text{ respectively, } \tilde{H}_I = H_I \text{ for } z_I = 1, \text{ while } \tilde{H}_I = H_I^c \text{ and for } z_I = -1, \text{ and }

\begin{align*}
c_I = z_I \tilde{c}_I - \frac{1}{2}.
\end{align*}

Here \(L_{IJ}(L'_{IJ})\) are generic hermitian functions of \(Z\) and \(Z^* (Z' \text{ and } Z'^*)\), and \(\omega_a\) and \(\lambda_{IJK}\) (\(\omega'\) and \(\lambda'_{IJK}\)) are generic holomorphic functions of \(Z (Z')\).

The 4D Yukawa couplings and soft parameters can be studied by constructing the effective action of massless 4D superfields. In our superfield basis, the \(y\)-independent modes of 5D superfields correspond to massless 4D superfields. Let \(V^a\) denote the \(y\)-independent mode of \(V^a\), and \(Q_I\) to be the \(y\)-independent mode of \(H_I\) when \(z_I = 1\) or of \(H_I^c\) when \(z_I = -1\). Here we will assume that all visible 4D gauge and matter fields are in \(\{V^a, Q_I\}\), and examine their Yukawa couplings and soft SUSY breaking parameters when the SUSY breaking is mediated by the brane superfields \(Z, Z'\) and the radion superfield \(T\). Those Yukawa couplings and soft parameters at the Kaluza-Klein (KK) scale \(M_{KK} \approx ke^{-\pi kR}\) can be evaluated from the 4D effective action which can be written as

\begin{align*}
\left[ \int d^4 \theta Y_{IJ} Q_I Q_J^* \right] + \left[ \int 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{align*}

where \(Y_{IJ}\) are hermitian wave function coefficients, \(f_a\) are holomorphic gauge kinetic functions, and \(\tilde{y}_{IJK}\) are holomorphic Yukawa couplings. From (3) and (4), we find

\begin{align*}
Y_{IJ} &= \frac{1}{c_I k} \left( 1 - e^{-c_I \pi k (T + T^*)} \right) \delta_{IJ} + L_{IJ}(Z, Z^*) + \frac{L'_{IJ}(Z', Z'^*)}{e^{(c_I + c_J + c_K) \pi k T}}; \\
f_a &= \frac{2\pi}{g_{5a}^2} T + \omega_a(Z) + \omega'_a(Z'), \quad \tilde{y}_{IJK} = \lambda_{IJK}(Z) + \frac{\lambda'_{IJK}(Z')}{e^{(c_I + c_J + c_K) \pi k T}}.
\end{align*}

Note that 5D SUSY in bulk enforces that the Yukawa couplings of \(Q_I\) originate entirely from the brane action (4).

It is straightforward to compute soft parameters for generic forms of \(L_{IJ}, L'_{IJ}\). Here we will focus on

\begin{align*}
L_{IJ}(Z, Z^*) &= -\kappa_{IJ} ZZ^*, \quad L'_{IJ}(Z', Z'^*) = -\kappa'_IJ Z' Z'^*,
\end{align*}

where \(\kappa_{IJ}\) and \(\kappa'_{IJ}\) are generic constants of order one, while keeping \(\omega_a\) and \(\lambda_{IJK}\) (\(\omega'_a\) and \(\lambda'_{IJK}\)) as generic holomorphic functions of \(Z (Z')\). The results for generic forms of \(L_{IJ}\) and
$L_{Ij}$ will be presented elsewhere [13]. In regard to the suppression of soft scalar masses and trilinear $A$ parameters at $M_{KK}$, those generic results show the same behavior as our case. For simplicity, we further assume that $\langle Z \rangle \ll 1$ and $\langle Z' \rangle \ll 1$ (in the unit with $M_5 = 1$), so $\langle L_{Ij} \rangle \ll 1$ and $\langle L'_{Ij} \rangle \ll 1$ though $\kappa_{Ij}$ and $\kappa'_{Ij}$ are generically of order one. However, in regard to SUSY breaking, we consider the most general case in which any of the $F$-components of $T, Z$ and $Z'$ can be the major source of SUSY breaking. Note that here we are not concerned with the dynamical origin of those $F$-components, but with the resulting soft parameters of visible fields for generic values of the $F$-components.

The Yukawa couplings $y_{IJK} \phi^I \psi^J \psi^K$ of canonically normalized superfield $Q_I = \phi^I + \theta \psi^J + \theta^2 F^I$ are easily obtained from (6):

$$y_{IJK} = \left( Y_I Y_J Y_K \right)^{-1/2} \left( \lambda_{IJK} + \frac{\lambda'_{IJK}}{e^{(c_I + c_J + c_K)\pi kT}} \right),$$

(8)

where $Y_I = (1 - e^{-c_I \pi k(T + T^*)})/k c_I$. The gaugino masses $\frac{1}{2} M_a \lambda^a \lambda^a$ for canonically normalized gauginos $\lambda^a$, the soft scalar mass-squares $m_{IJ}^2 \phi^I \phi^J$, and the trilinear $A$-terms $A_{IJK} \phi^I \phi^J \phi^K$ for canonically normalized scalar fields $\phi^I$ are also obtained to be

$$M_a = \frac{1}{2} g_a^2 \left( \frac{2\pi}{2g_a^2} F^T + \frac{\partial \omega_a}{\partial Z} F^Z + \frac{\partial \omega'_a}{\partial Z'} F^{Z'} \right),$$

$$m_{IJ}^2 = (Y_I Y_J)^{-1/2} \left[ \frac{\pi^2 c_I k \delta_{IJ}}{e^{c_I \pi k(T + T^*)} - 1} + \kappa_{IJ} |F^Z|^2 + \frac{\kappa'_{IJ} |F^{Z'}|^2}{e^{(c_I + c_J + c_K)\pi kT}} \right],$$

$$A_{IJK} = (Y_I Y_J Y_K)^{-1/2} \left[ F^T \left( \frac{\partial}{\partial T} \ln \left( \frac{\lambda_{IJK} + \lambda'_{IJK} e^{(c_I + c_J + c_K)\pi kT}}{Y_I Y_J Y_K} \right) \right) \right.$$  
$$\times \left( \lambda_{IJK} + \frac{\lambda'_{IJK}}{e^{(c_I + c_J + c_K)\pi kT}} \right) + F^Z \frac{\partial \lambda_{IJK}}{\partial Z} + \frac{F^{Z'}}{e^{(c_I + c_J + c_K)\pi kT}} \frac{\partial \lambda'_{IJK}}{\partial Z'} \right],$$

(9)

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.

One can now consider two simple models in which the SUSY flavor problem is solved (or ameliorated) by warped geometry. Here we will briefly describe the models, while leaving the details of the models including the phenomenological aspects to the subsequent work [13]. In the model (I), Yukawa couplings exist only at $y = 0$, i.e. $\lambda'_{IJK} = 0$, and there is no SUSY breaking $Z$ at $y = 0$. Then an appropriate radion stabilization mechanism is assumed.
to yield $e^{-\pi k R} \approx 10^{-2} - 10^{-5}$ for which $M_{KK} \approx k e^{-\pi k R} \approx 10^{16} - 10^{13}$ GeV. Note that these conditions are stable against radiative corrections. We further assume that $c_I = 0$ for Higgs superfields $Q_I$, while $c_{J,K} > 0$ with $e^{-c_J \pi k R} \ll 1$ and $e^{-c_K \pi k R} \ll 1$ for $Q_J$ and $Q_K$ denoting the quark and/or lepton superfields. In this model, the quark and lepton Yukawa couplings are given by

$$y_{IJK} = \left( \frac{k^2 c_{JCK}}{2\pi R} \right)^{1/2} \lambda_{IJK},$$

so warped geometry is not responsible for hierarchical Yukawa couplings. As for the soft parameters at $M_{KK}$, $M_a = \mathcal{O}(F_T)$ and/or $\mathcal{O}(F_{Z'})$, and

$$m_{J\bar{K}}^2 = \frac{k(c_J c_K)^{1/2}}{e(c_J + c_K)\pi k R} \left( \frac{\pi^2 k c_J \delta_{JK} \left| F_T \right|^2 + c'_{JK} \left| F_{Z'} \right|^2}{2\pi R} \right) = \mathcal{O}(e^{-(c_J + c_K)\pi k R} M_a^2),$$

$$A_{IJK} = -y_{IJK} \frac{F_T}{2R} = \mathcal{O}(y_{IJK} M_a).$$

Note that the squark and slepton mass-squares $m_{J\bar{K}}^2$ are suppressed by small warp factor. On the other hand, $A_{IJK}/y_{IJK}$ is unsuppressed, but it is universal and generically contains nonzero CP phase. As a result, $m_{J\bar{K}}^2$ of the first and second generations at the weak scale are (approximately) flavor-independent as they are generated mainly by the RG evolution arising from $M_a$. In fact, one can consider the variation of the model (I) in which $e^{-c_I \pi k R} \ll 1$ for Higgs superfields $Q_I$. In such model, both $A_{IJK}/y_{IJK}$ and $m_{J\bar{K}}$ at $M_{KK}$ are exponentially suppressed compared to $M_a$.

In the model (II), $\lambda_{IJK} = 0$ and all $\lambda'_{IJK}$ are of order unity, and there is no SUSY breaking $Z$ at $y = 0$. The radion is stabilized again at $e^{-\pi k R} = 10^{-2} - 10^{-5}$, and we assume that $c_I = 0$ for Higgs superfield $Q_I$, while $c_{J,K} \geq 0$ for quark and/or lepton superfields $Q_{J,K}$. In this model, hierarchical Yukawa couplings are naturally generated by warped geometry:

$$y_{IJK} = \left( \frac{1}{(2\pi R Y_J Y_K)^{1/2}} \right) \frac{\lambda_{IJK}}{e(c_J + c_K)\pi k R} = \mathcal{O}(e^{-(c_J + c_K)\pi k R}),$$

where $Y_J = (1 - e^{-2c_J \pi k R})/k c_J = 2\pi R$ when $c_J = 0$, and $Y_J = 1/k c_J$ when $e^{-c_J \pi k R} \ll 1$. Warped geometry similarly suppresses $m_{J\bar{K}}^2$ and $A_{IJK}$ at $M_{KK}$ as
\[ m_{jk}^2 = (Y_j Y_k)^{-1/2} \left( \frac{\pi^2 c_{jk} \delta_{jk} |F'_T|^2}{e^{2c_{jk} \pi kR} - 1} + \frac{k'_{jk} |F'_Z|^2}{e^{(c_j + c_k)\pi kR}} \right) = \mathcal{O}(e^{-(c_j + c_k)\pi kR} M_a^2), \]

\[ A_{ijk} = \frac{e^{-(c_{ijk}+c_{k}^r)}kR}{(2\pi R Y_j Y_k)^{1/2}} \left[ -\frac{F'_T}{2R} \left( 1 + 2(c_j + c_k)\pi kR + \frac{2c_{jk} \pi kR}{e^{2c_{jk} \pi kR} - 1} \right. \right. \\
\left. \left. + \frac{2c_{k}^r \pi kR}{e^{2c_{jk} \pi kR} - 1} \right) \right] = \mathcal{O}(y_{ijk} M_a), \]  

where again \( M_a = \mathcal{O}(F^T) \) and/or \( \mathcal{O}(F^Z') \). An interesting feature of this model is that \( m_{jk}^2 \) at \( M_{KK} \) have a correlation with Yukawa couplings. To be specific, let us consider the case with \( n_{qi} = (3, 2, 0), n_{ui} = (5, 2, 0), n_{di} = (2, 1, 1) \) where \( n_K = -c_K \pi kR / \ln(0.2) \) for the 3-generations of \( SU(2) \)-doublet quark superfields \( q_i \) and the \( SU(2) \)-singlet antiquark superfields \( u_i, d_i \). This model gives the up-quark Yukawa couplings \( y_{ij}^u \sim 0.2^{n_u_i + n_u_j} \) and the down-quark Yukawa couplings \( y_{ij}^d \sim 0.2^{n_d_i + n_d_j} \), producing the correct quark masses and mixing angles. At the same time, the model predicts the following patterns of squark mass-squares at \( M_{KK} \): \( m_{ij}^2(\tilde{q})/M_a^2 \sim 0.2^{n_u_i + n_u_j} \), \( m_{ij}^2(\tilde{u})/M_a^2 \sim 0.2^{n_d_i + n_u_j} \), and \( m_{ij}^2(\tilde{d})/M_a^2 \sim 0.2^{n_d_i + n_d_j} \). So in the model (II) also, the squark masses at the weak scale of the first and second generations are generated mainly by the flavor-independent RG evolution arising from \( M_a \). Note that this model can give a sizable 23-component of the squark mass-square matrices, which may lead to interesting phenomenology [13].

When \( c_{ijk} \leq 0 \) for the quark and lepton superfields \( Q_{ijk} \), one can obtain similar models yielding exponentially suppressed Yukawa couplings and/or soft parameters by exchanging \((Z, \kappa_{ijk}, \lambda_{ijk})\) and \((Z', \kappa'_{ijk}, \lambda'_{ijk})\), in the previously described models for \( c_{ijk} \geq 0 \). At any rate, the suppression of Yukawa couplings and soft parameters by warp factor has a simple geometric interpretation. If \( e^{-c_j \pi kR} \ll 1 \) \((e^{-c_j \pi kR} \gg 1)\), the wavefunction of \( Q_j \) is localized near at \( y = 0 \) \((y = \pi)\) with an exponentially small tale at \( y = \pi \) \((y = 0)\). As a consequence, its overlaps with the brane Yukawa coupling \( \lambda'_{ijk} \) \((\lambda_{ijk})\) and the brane SUSY breaking \( F^Z' \) \((F^Z)\) at \( y = \pi \) \((y = 0)\) are exponentially suppressed. Also the wavefunction coefficients \( Y_{ijk} \) of such \( Q_{ijk} \) are exponentially insensitive to the distance between two fixed points, so to the value of \( T \), explaining why the contributions of \( F^T \) to \( m_{jk}^2 \) are exponentially suppressed.
independently of the sign of $c_{J,K}^4$. On the other hand, the wavefunction of $V^a$ is constant over the 5-th dimension, so there is no such suppression in gaugino masses.

The results (10)-(13) are obtained for the 4D effective action (5) without including the possible threshold corrections due to massive KK modes. Obviously, the (exponential) suppressions of $y_{IJK}$ and $A_{IJK}$ are stable against KK threshold corrections as they are due to the (exponentially) small holomorphic Yukawa couplings $\tilde{y}_{IJK}$. On the other hand, $m_{j\bar{k}}^2$ generically get threshold corrections of order $M_a^2/8\pi^2$. In fact, the structure of flavor-violating interactions in the models (I) and (II) suggests that the flavor-violating part of the KK threshold corrections to $m_{j\bar{k}}^2$ is further suppressed by a small factor involving $\lambda_{IJK}$ or $\lambda'_{IJK}/e^{(c_J+c_K)\pi k R}$ or $\kappa'_{JK}/e^{(c_J+c_K)\pi k R}$ [13]. There can be additional corrections to soft parameters which are induced by non-renormalizable SUGRA interactions in 4D effective action [14], but they are suppressed by $M_{KK}^2/8\pi^2 M_{Pl}^2 \sim e^{-2\pi k R}/8\pi^2$, so are small enough. There are also the model-independent SUSY breaking effects mediated by the 4D superconformal anomaly [5]. In the models (I) and (II), we obtain the gravitino mass $m_{3/2} = O(e^{-\pi k R} M_a)$, so the anomaly-mediated contributions to soft parameters are $\delta M_a \sim e^{-\pi k R} M_a/8\pi^2$, $\delta m_{j\bar{k}}^2 \sim e^{-2\pi k R} M_a^2/(8\pi^2)^2$ and $\delta A_{IJK} \sim e^{-\pi k R} M_a/8\pi^2$, which are small enough to be ignored. So the leading radiative corrections to Yukawa couplings and soft parameters at $M_W$ in the models (I) and (II) come from the standard RG running down to $M_W$ starting from the boundary values at $M_{KK}$ given by (10)-(13).

We finally note that the idea of AdS/CFT correspondence suggests a CFT framework which would reproduce the main features of our AdS models. Indeed, models involving superconformal (SC) sector have been proposed to generate hierarchical Yukawa couplings as well as exponentially suppressed soft masses [15,16]. It is unclear yet the possible connection between these SC models and the models (I) and (II) presented here, though one can easily

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find the correspondence: $c_I \rightarrow \gamma_I/2$ and $\pi k R \rightarrow \ln(\Lambda/M_c)$, where $\gamma_I$ is the anomalous dimension of $Q_I$ driven by the coupling to the SC sector, and $\Lambda$ and $M_c$ are the cutoff scale and the decoupling scale of the SC sector, respectively. At any rate, the AdS approach discussed in this paper provides interesting perturbative framework to solve the SUSY flavor problem and the Yukawa hierarchy problem.

To conclude, we have noted that supersymmetric 5D theory on a slice of AdS$_5$ with the Kaluza-Klein scale $M_{KK} \approx 10^{16} - 10^{13}$ GeV can provide a solution to the SUSY flavor problem, while generating hierarchical Yukawa couplings. This framework utilizes the AdS warp factor $e^{-\pi k R} \approx 10^{-2} - 10^{-5}$ to suppress the soft scalar masses and trilinear $A$-parameters at $M_{KK}$, and provides phenomenologically interesting prediction for the patterns of soft parameters.

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