Non-universal Soft Parameters in Brane World and the Flavor Problem in Supergravity

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\textbf{Abstract}

We consider gravity mediated supersymmetry (SUSY) breaking in 5D spacetime with two 4D branes B1 and B2 separated in the extra dimension. Using an off-shell 5D supergravity (SUGRA) formalism, we argue that the SUSY breaking scales could be non-universal even at the fundamental scale in a brane world setting, since SUSY breaking effects could be effectively localized. As an application, we suggest a model in which the two light chiral MS SM generations reside on B1, while the third generation is located on B2, and the Higgs multiplets as well as gravity and gauge multiplets reside in the bulk. For SUSY breaking of the order of 10–20 TeV caused by a hidden sector localized at B1, the scalars belonging to the first two generations can become sufficiently heavy to overcome the SUSY flavor problem. SUSY breaking on B2 from a different localized hidden sector gives rise to the third generation soft scalar masses of the order of 1 TeV. Gaugino masses are also of the order of 1 TeV if the size of the extra dimension is $\sim 10^{-16}$ GeV$^{-1}$. As in 4D effective supersymmetric theory, an adjustment of TeV scale parameters is needed to realize the 100 GeV electroweak symmetry breaking scale.

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1 Introduction

The minimal supersymmetric standard model (MSSM) is an attractive avenue for physics beyond the standard model (SM). The MSSM not only resolves the naturalness problem regarding the small Higgs scalar mass that SM suffers from, but also improves the scenario of gauge couplings unification and electroweak symmetry breaking. However, the MSSM has numerous (> 100) theoretically undetermined parameters. To understand the origin of many soft supersymmetry (SUSY) breaking parameters in the MSSM simply and consistently, one often invokes the supergravity (SUGRA) models. However, models relying on the conventional gravity-mediated SUSY breaking scenario generally contain flavor changing processes that are strongly constrained by experiments. Approaches for resolving this notorious problem include gauge-mediated SUSY breaking and anomaly mediated SUSY breaking. On the other hand the effective supersymmetric theory (ESUSY) provides an alternative scenario. In ESUSY, the superpartners of the first two generations are required to be sufficiently heavy (~ 20 TeV) to suppress flavor changing processes in SUGRA models. (This also could provide a mechanism for suppressing dimension five proton decay operators.) Because of the relatively small Yukawa couplings of the first two generations, the Higgs boson masses are radiatively stable despite these heavy masses. On the other hand, the left- and right-handed top squarks are constrained to be not much heavier than 1 TeV or so to preserve the gauge hierarchy solution because of the relatively large top quark Yukawa coupling. One expects from $SU(2)$ symmetry that the left-handed sbottom mass also does not exceed a TeV or so. In addition, the B-ino, W-ino, and also Higgsino, which are coupled to the Higgs scalars with “sizable” gauge couplings, should remain lighter than a TeV or so.

However, with only a unique SUSY breaking scale at $M_{\text{Planck}}$ as in ordinary 4D SUGRA models, such a large hierarchy between the first two and the third genera-

\footnote{In Ref. \cite{9}, an ESUSY idea is realized by introducing an anomalous $U(1)_A$.}
ations’ soft masses is not easily derived radiatively at the electroweak scale, even though the first two generations’ Yukawa couplings are relatively small. Moreover, the first two generations’ heavy soft mass squareds tend to drive the third ones to negative values at low energy through two loop renomalization group equations, so that the third generation squark masses should be raised to several TeV at the grand unified theory (GUT) scale. This makes the scenario perhaps somewhat less attractive.

In this paper we will introduce two (or more) SUSY breaking scales at the GUT or some other fundamental scale within a gravity-mediated SUSY breaking scenario without invoking any flavor symmetry, deploying two (or more) 4 dimensional branes or orbifold fixed points in 5 dimensional spacetime. In 4D gravity-mediated SUGRA models using only hidden sector SUSY breaking, there is no “easy” way to provide several SUSY breaking effects that depend on the flavor; once SUSY breaks down in a hidden sector, the consequences typically spread impartially among particles in the visible sector through direct gravity couplings appearing in the SUGRA Lagrangian. However, if we could somehow couple each visible sector field gravitationally exclusively to a specific hidden sector among several hidden sectors with different SUSY breaking scales, it would be possible to introduce flavor dependent SUSY breaking scales even within the gravity-mediated scenario. This is where the extra dimension and branes can play an important role.

The SUSY breaking soft scalar masses and the “A” terms are derived from the scalar potential in SUGRA models. In 5D brane world, the scalar potentials as well as the Yukawa interaction terms are usually constructed only on the branes due to supersymmetry. Thus, as shown in [12] for global SUSY, the scalar potentials are localized by the delta function on the brane and separated from each other at tree level. In this paper we will show that the SUSY breaking effects from the hidden

\footnote{Reference [10] shows that the 10 TeV–sub-TeV hierarchy between the first two and the third generations soft scalar masses at the electroweak scale can be derived radiatively from 10 TeV soft masses at the GUT scale.}
sectors localized on the branes are transmitted to the visible sector fields located on the same brane and bulk gauginos through direct gravitational couplings, but remain shielded at least at tree level from the other brane fields \([7, 13]\) and bulk scalars. Thus if the SUSY breaking scales are \(m_1\) and \(m_2\) at two separated branes B1 and B2, respectively, the localized chiral multiplets at B1 (B2) could get soft SUSY breaking effects of order \(m_1\) (\(m_2\)). The mixing terms from one loop contribution among fields from two distinct branes can be sufficiently suppressed if the interval size is sufficiently large compared to the fundamental scale \([14]\).

Let us assume that the SUSY breaking scales at B1 and B2 are \(m_1\) (10–20 TeV) > \(m_2\) (\(\sim\) 1 TeV). Then, we can easily make the superpartners of the first two generations sufficiently heavy by locating them on B1. In order to keep the radiative corrections to the Higgs masses under control, the third family resides on B2 (or in the bulk). The two Higgs multiplets should reside in the bulk, so that the first two and the third generations of the quarks and leptons can couple to them at B1 and B2, respectively. Note that matter fields in the bulk are accompanied by bulk gauge fields for the consistency of the gauge theory.

The gaugino mass term can be generated at tree level in some D=5 off-shell SUGRA formalism \([17]\) if SUSY is spontaneously broken, even when a singlet field does not couple to the gauge kinetic term. Thus heavy gaugino masses from SUSY breaking at B1 could give rise to large radiative corrections to the Higgs mass and spoil the naturalness solution. However, since the bulk gaugino mass is given by \(m_{1/2} \sim m_{3/2} \times (M_*/M_P)\), where \(M_*\) is a 5D SUGRA fundamental scale and \(M_P\) is the reduced 4D Planck scale, a not so “large” extra dimension, for example, of order \(10^{-16} \text{ GeV}^{-1}\) (or \(M_* \sim 10^{17} \text{ GeV}\)), suppresses the gaugino masses to 1 TeV scale. Moreover, by introducing an additional set of the localized gauge multiplet at B2 with large coefficient (\(\sim 10\)) (i.e., in addition to the bulk gauge multiplet), the gauginos mass can be more suppressed to the 100 GeV scale as we will see. In such a way, the masses of the various gauginos in the MSSM also could be made
2 5D off-shell SUGRA and 4D brane matter

In order to couple 5D SUGRA to 4D brane matter it is convenient to employ the 5D off-shell SUGRA formalisms recently developed in [15, 16, 17, 18, 19]. Although we will employ the formalism in [15, 16, 17, 18], our essential results such as localized SUSY breaking effects and realization of ESUSY should also hold in other formalisms and/or in other spacetime dimensions.

The 5D formalism in [15] is an extension of the 4D SUGRA off-shell formalism of [20]. The latter contains (16+16) bosonic and fermionic degree of freedoms, namely, $e^a_m$ (vierbein), $\psi_m$ (gravitino), $a_m$ ($U(1)_R$ axial gauge vector), $b_a$ (axial vector), $\xi$ (spinor), $S$ (scalar), $t^1$ (pseudoscalar), and $t^2$ (scalar), where every field is auxiliary except for the vierbein and the gravitino. In 5D, N=2 SUGRA, the familiar N=2 SUGRA on-shell fields, namely the fünfbein ($e^A_M$), gravitino ($\psi^j_M$) ($j = 1, 2$) and graviphoton ($A_M$) are supplemented by the auxiliary fields,

\[ \vec{e}, \ v_{AB}, \ \vec{V}_M, \ \zeta, \ \text{and} \ C, \]

an isotriplet, an antisymmetric tensor, a $SU(2)_R$ gauge vector, a spinor, and a scalar, respectively. In the gravitino $\psi^j_M$, the index “$j$” (=1,2) is a $SU(2)$ index. Although the closure of SUSY algebra only requires the above (40+40) bosonic and fermionic degrees of freedom (“minimal multiplet”), the corresponding action turns out to be unphysical; the equations of motion give rise to $\det e^A_M = 0$. (This problem is also observed in 4D SUGRA formalism with a single auxiliary fermion [20].) To resolve it, additional (8+8) auxiliary fields are needed, either from a nonlinear multiplet (version I) [15], a hypermultiplet (version II) [18], or a tensor multiplet (version III) [16]. In this paper we prefer the first choice. The nonlinear multiplet contains the auxiliary fields, $\phi^\alpha$ (scalar), $\chi$ (spinor), $\varphi$ (scalar), and $V_A$ (vector), where $\alpha$ is the index of an additional $SU(2)$ and $a$ is the Lorentz index. Closure of the algebra on $\varphi$ and

non-universal.
$C - \frac{1}{16\kappa} R_{AB}^{\ AB} - \frac{\kappa}{24} F_{AB} F^{AB} + 5\kappa t^2 + \frac{\kappa}{4} v_{AB} v^{AB} - \frac{1}{8\kappa} \hat{D}_M \phi^i \hat{D}^M \phi^i + \cdots = 0 \ , \ (2)$

where $\hat{D}_M$ is the supercovariant derivative defined in [15], and “…” contains the fermionic contributions. The bosonic part of the pure SUGRA action is given by [15]

$$S_{SUG} = \int d^4x \int_{-L}^{L} dy \ e^{\frac{x}{2\kappa^2}} \left[ -\frac{1}{4\kappa^2} R(\phi)_{AB}^{\ AB} - \frac{1}{6} F_{AB} F^{AB} - \frac{1}{\sqrt{3}} F_{AB} v^{AB} 
- \frac{\kappa}{6\sqrt{3}} \epsilon^{ABCDE} A_A F_{BC} F_{DE} + v_{AB} v^{AB} - 12t^2 - \frac{1}{4} V_A V^A \right] \ , \ (3)$$

where the factor $2L$ for the size of the extra dimension ensures that the mass dimensions of all fields correspond to their canonical 4D ones, and the induced SUGRA parameter $\kappa$ in Eq. (3) and the original 5D SUGRA parameter $\kappa_*$ are defined as

$$\kappa^2 \equiv \frac{8\pi G}{M_P^2} \ , \ (4)$$

$$\kappa_*^3 \equiv \frac{1}{M_*^3} \equiv 2L \times \kappa^2 = \frac{2L}{M_P^2} \ , \ (5)$$

where $M_P$ and $M_*$ are the Planck and the fundamental scale mass parameters, respectively. For future convenience, we define $\kappa_*$ such that it has length dimension unlike the definition in ordinary 5D SUGRA.

With the fifth dimension compactified on a $S^1/Z_2$ orbifold, $Z_2$ parities assigned to fields in 5D N=2 theory, consistent with SUSY transformation, are given in Table I. Since only the even parity particles contain massless modes, truncation of the heavy Kaluza-Klein (KK) modes reduces the 5D, N=2 SUGRA to 4D, N=1 SUGRA of [20]. The relations between the auxiliary fields in 4D N=1 SUGRA and those in the

| $e_m^a$ | $e_5^a$ | $A_5$ | $e_{55}$ | $A_m$ | $e_{m5}$ | $\psi_{mR}^i (\psi_{mL}^i)$ | $\psi_{SR}^i (\psi_{SL}^i)$ | $v_a^5$ | $\zeta$ | $C$ | $V_{m5}^5$ | $V_{51^2}$ | $t_{1,2}^i$ | $V_5$ | $\varphi$ | $\chi$ |
|-------|-------|-------|-------|-------|-------|----------------|----------------|-------|-------|-------|-------|-------|-------|-------|-------|-------|-------|

Table I. $Z_2$ parity assignments for the fields in D=5, N=2 pure SUGRA.

$V_A$ gives rise to the constraint [15]
The 5D N=2 hypermultiplet \[18\] could be extended from the 4D N=2 hypermultiplet \[22\]. Since there is no field carrying vector indices in the hypermultiplet, the structures of the hypermultiplets in D=4,5,6 are almost the same. The bosonic part of the corresponding 5D N=2 off-shell Lagrangian is \[15\]

\[
\mathcal{L}_{HYP} = \frac{1}{2L} d^\alpha_\beta \left[ \frac{1}{2} D_A A^j_\alpha A^j_\beta + \frac{1}{2} F^j_\alpha F^j_\beta \left( 1 - \frac{4}{3} A_A A^A \right) \right] + A^j_\alpha A^j_\beta \left( \frac{1}{8} R(\tilde{\omega})_{AB}^{AB} + \frac{\kappa}{12} F_{AB} F^{AB} - 2\kappa C - 10\kappa^2 \tilde{t}^2 - \frac{\kappa^2}{2} v_{AB} v^{AB} \right) ,
\]

where \( j(=1,2) \) is the \( SU(2)_R \) index as before, and \( \alpha(=1,2,\cdots,2r) \) is a \( USp(p,q) \) \((p+q=2r)\) index, and \( A, B \) are 5D Lorentz indices. As in the pure SUGRA case, the linear dependence of the Lagrangian on \( C \) must be resolved using the nonlinear multiplet. Equation (2) eliminates the couplings of \( A^j_\alpha A^j_\beta \) with \( \tilde{t}^2, R_{AB}^{AB} \), and so on in the Lagrangian (3).

For \( r = 1, USp(2,0) = SU(2), \) \( d^\alpha_\beta = -\epsilon^{\alpha\gamma} \epsilon_{\beta\gamma} = \delta^\alpha_\beta. \) The component fields in the hypermultiplet fulfill the reality constraints \[15, 22\]. For the scalar field, for instance,

\[
A^j_\alpha \equiv A^{\alpha*}_j = e^{jk} \rho_{\alpha\beta} A^\beta_k ,
\]

where \( \rho_{\alpha\beta} \) is proportional to \( \epsilon_{\alpha\beta} \) and satisfies

\[
|\det \rho|^2 = 1 , \quad \text{and} \quad \rho^{\alpha\beta} = \rho^{\alpha\beta*} .
\]

The auxiliary field \( F^j_\alpha \) is similarly constrained. The fermion \( \xi^\alpha \) satisfies the symplectic Majorana condition,

\[
\overline{\xi}_\alpha \equiv (\xi^\alpha)^\dagger \gamma^0 = \rho_{\alpha\beta} \xi^{\beta T} C ,
\]
where \( C = \text{diag}(-i\tau^2, -i\tau^2) \).

The hypermultiplet is split under \( Z_2 \) symmetry into two chiral multiplets with even and odd parities, respectively. The (bulk) chiral multiplet with even parity can take part in the superpotential on the 4 dimensional branes. On the other hand, the (bulk) chiral fields with odd parity vanish on the branes and do not possess massless modes. Consequently, they are neglected in the low energy physics. The surviving fields at low energy are \( A_1^1 \) (or \( A_2^{2*} \)), \( \xi_{1L} \) (or \( \xi_{2R}^* \)) and \( F_1^1 \) (or \( F_2^{2*} \)).

In the more general case, the Lagrangian could be written as a non-linear sigma model, as shown in the on-shell formalisms of [23]. The kinetic term is then \( \mathcal{L}_{\text{kin}} = h_{uv}(\phi) \partial_M \phi^u \partial^M \phi^v \), where \( h_{uv}(\phi) \) \((u, v = 1, 2, \ldots, 4r)\) is a metric on a quaternionic manifold with coordinates \( \phi^u \).

A 4D chiral multiplet consists of scalar and fermion fields with

\[
A = (A, B; \psi; F, G),
\]

where \( A, B \) are dynamical real scalars, \( \psi \) is a chiral fermion, and \( F, G \) are real auxiliary fields. Thus bosonic degrees of freedom are the same as those of fermions. In this paper, the first member in a chiral multiplet such as \( A \) in Eq. (13) will be used, if necessary, to designate the chiral multiplet or superfield to which it belongs.

To each chiral multiplet, a weight \( w_i \) \((U(1)_R \) charge) is assigned.

According to the local tensor calculus [20, 21], the canonical kinetic part of the off-shell Lagrangian for the chiral multiplet is given by the “D-density” of the product, \(-[A_i \times A_i]_D/4\), and its bosonic part is [11, 18, 20]

\[
\mathcal{L}_{4dCHI} = \sum_i \left[ \frac{1}{2} \left( \hat{D}_m A_i \hat{D}^m A_i + \hat{D}_m B_i \hat{D}^m B_i + F_i^2 + G_i^2 \right) \right. \\
-4\kappa t^2 (F_i A_i - G_i B_i) - 4\kappa t^1 (F_i B_i + G_i A_i) + 2\kappa b^m (A_i \hat{D}_m B_i - B_i \hat{D}_m A_i) \\
-\frac{\kappa}{2} (A_i^2 + B_i^2) \left[ 12 w_i S - \frac{w_i}{4\kappa} R_{ab}^{ab} + 48\kappa (w_i - 1) \left[ (t^1)^2 + (t^2)^2 \right] \\
-6\kappa w_i b_m b^n - 8 S \right],
\]

where \( i \) labels chiral multiplets, and \( \hat{D}_m \) is the supercovariant derivative, defined as
\[ \hat{D}_m A \equiv \partial_m A - \bar{\psi}_m \tau^2 \psi - w B a_m , \]  
\[ \hat{D}_m B \equiv \partial_m B - \bar{\psi}_m \tau^2 \gamma^5 \psi + w A a_m . \]  

We note that in the Lagrangian (14) several weight \((w_i)\) dependent terms appear. This can be generalized to non-trivial cases, \([f(A) \times A]_D\), where \(f\) is an arbitrary function of \(A\) \[^2\]. The Lagrangian in Eq. (14) can be written in terms of 5D SUGRA auxiliary fields using Eqs. (6)–(8),

\[ L_{5dCHI} = \sum_{I,i} \delta(y - y_I)[|\partial \phi_i|^2 + |F_i|^2 - 4\kappa_* \phi_i^* F_i^* M - 4\kappa_* \phi_i^* F_i^* M^* \]  
\[ - i \kappa_* \phi_i^2 \left( (3w_i - 2)(2\partial_3 t^3 + \kappa_* \mathcal{M} \mathcal{V}^* + \kappa_* \mathcal{M}^* \mathcal{V}) - 48\kappa_* (w_i - 1)|\mathcal{M}|^2 \right) \]  
\[ - 12w_i C + 8C + \frac{w_i}{4\kappa_*} R^{ab}_{\alpha \beta} - 6\kappa_* w_i v_{a_5} v_{a_5} \]  
\[ + \frac{\kappa_*}{4} w_i^2 (V_3^2 - \frac{2}{\sqrt{3}} F_{m5} + 4v_{m5}^2)(V_{m5}^3 + \frac{2}{\sqrt{3}} F_{m5}^3 - 4v_{m5}^2) \]  
\[ + 2\kappa_* w_i v_{m5}^2 (w_i V_3^2 - \frac{2}{\sqrt{3}} F_{m5} + 4v_{m5}^2) \right] , \]

where we complexified some bosonic degrees for future convenience,

\[ \phi \equiv \frac{1}{\sqrt{2}}(A + iB) , \]  
\[ F \equiv \frac{1}{\sqrt{2}}(F - iG) , \]  
\[ M \equiv t^2 + it^1 , \]  
\[ V \equiv V_5^1 - iV_5^2 . \]

Note that in Eq. (17), the expansion parameter is not \(\kappa\) but \(\kappa_*\) defined in Eq. (5). Since the Lagrangian (14) describes the dynamics of scalar fields localized on the brane, a delta function appears as an overall factor. \(I = 1, 2\) indicate the two brane locations at \(y = 0\) and \(y = L\), respectively.

Next we briefly review how the superpotential is constructed in the 4D SUGRA, which would enable us to derive the relevant scalar potentials. In the local tensor
calculus \([20, 21]\), a product "\(\cdot\)" is defined such that the product, \(A_3 = A_1 \cdot A_2\) of two chiral multiplets of weights \(w_1\) and \(w_2\), respectively, yields another chiral multiplet with weight \(w_3 = w_1 + w_2\), and components

\[
\begin{align*}
A_3 &= A_1 A_2 - B_1 B_2, \quad (22) \\
B_3 &= A_1 B_2 + A_2 B_1, \quad (23) \\
\psi_3 &= (A_1 - \gamma_5 B_1) \psi_2 + (A_2 - \gamma_5 B_2) \psi_1, \quad (24) \\
F_3 &= A_1 F_2 + B_1 G_2 + F_1 A_2 + G_1 B_2 + \bar{\psi}_1 \psi_2, \quad (25) \\
G_3 &= A_1 G_2 - B_1 F_2 + G_1 A_2 - F_1 B_2 - \bar{\psi}_1 \gamma_5 \psi_2. \quad (26)
\end{align*}
\]

These product rules are valid in the locally supersymmetric case as well as in global supersymmetry. Using the rules, we can define superpotentials such as \(A_1 \cdot A_2\), \(A_1 \cdot A_2 \cdot A_3\), \(\cdots\).

By examining the SUSY transformations of members in a chiral multiplet, one can find a supersymmetric invariant, "F-density," which is useful for constructing the Lagrangian. A supersymmetric invariant up to total derivative terms in SUGRA is given by \([16, 20, 21]\)

\[
[A]_F = F + i \kappa \bar{\psi}_m \gamma^m \psi + \frac{\kappa^2}{2} \bar{\psi}_m \tau^2 \gamma^{mn}(A + \gamma^5 B) \psi_n - 12 \kappa t^2 A - 12 \kappa t^1 B, \quad (27)
\]

with weight equal to 2 \([20]\). The Pauli matrix \(\tau^2\) contracts the \(SU(2)_R\) indices of the gravitino. The first term in Eq. (27) is present also in the globally supersymmetric case, while the other terms are the SUGRA corrections. In particular, if \(A\) and/or \(B\) acquire vacuum expectation values (VEVs), the third term could yield the gravitino mass term by absorbing the chiral fermion’s degree of freedom of the second term. Thus, for a given superpotential \(W(\phi_1, \phi_2, \cdots)\), the Yukawa interactions and scalar potentials are read off using Eqs. (22)–(26) and (27), and it can be checked that the bosonic part is given by

\[
\mathcal{L}_{4dYUK} = \left[ \sum_z W_z F_z - 12 \kappa W \mathcal{M} + \text{h.c.} \right], \quad (28)
\]
where every bosonic field is complexified using Eqs. (18)–(21), $F_z$ is the auxiliary field involved in $\phi_z$, and
\[ W_z \equiv \frac{\partial W}{\partial \phi_z}. \] (29)

3 Effective 4D scalar potential

To illustrate the emergence of non-universal soft masses, let us consider a model with a visible (V) and a hidden (H) hypermultiplet in the bulk, and two visible and some hidden chiral multiplets localized at each brane. The chiral multiplet from the bulk hypermultiplets with even parity $\Phi^V$ comprises a trilinear superpotential together with two visible chiral multiplets $\phi_1^V$ and $\phi_2^V$, localized at B1, while the hidden sector multiplets $\Phi^H$ and $\phi_x^H (x = a, b, \cdots)$ make up another superpotential at B1,
\[ W_{1V} = Y_{12}^V \Phi^V \phi_1^V \phi_2^V, \quad W_{1H} = W_{1H}(\Phi^H, \phi_x^H, \phi_x^H, \cdots). \] (30)

Similarly, the two (bulk) chiral multiplets, the two localized visible multiplets, $\phi_1^V, \phi_2^V$, and hidden chiral multiplets $\phi_a^H (x = a, b, \cdots)$ constitute the following superpotentials at B2,
\[ W_{2V} = Y_{12}^{2V} \Phi^V \phi_1^{2V} \phi_2^{2V}, \quad W_{2H} = W_{2H}(\Phi^H, \phi_a^{2H}, \phi_b^{2H}, \cdots). \] (31)

We assign weights 2, 0, and 2/3 to $\Phi^V$, $\phi_i^V (I, i = 1, 2)$, and $\phi_x^H$, respectively, for simplicity of the model. Other weight assignments that do not change our essential conclusions are also possible.

The bosonic part of the locally supersymmetric off-shell Lagrangian for the Yukawa interactions is
\[ \mathcal{L}_{YUK} = \sum_{I,S,z} \delta(y - y_I) \left[ W_{z IS}^V \mathcal{F}_z^{IS} - 12\kappa_S W_{z IS}^V \mathcal{M} + \text{h.c.} \right], \] (32)
where $I (= 1, 2)$ denotes the brane locations ($y_1 = 0, y_2 = L$), $S$ stands for visible or hidden sectors ($S = V, H$), $z = \Phi^V, 1, 2$ for $S = V$, and $z = \Phi^H, a, b, \cdots$ for $S = H$. To get a scalar potential for the dynamical fields, we replace the auxiliary fields in the
Lagrangians (3), (9), (17), and (32), using the equations of motion. The equations of motion for $F$ fields give

$$\mathcal{F}_{iS} = -(W_{iS}^I)^* + 4\kappa_* \phi_i^{I*} \mathcal{M}^*,$$

(33)

$$\mathcal{F}_\Phi = -\sum_I 2\Delta(y_I)(W_{\Phi}^{IS})^*(1 - 4\kappa^2 A_5 A_5)^{-1},$$

(34)

where $i$ is 1, 2 for $S = V$, and $a, b, \cdots$ for $S = H$, and

$$\Delta(y_I) \equiv L \times \delta(y - y_I).$$

(35)

In Eq. (34), we see that the bulk field $\mathcal{F}_\Phi^S$ gets contributions from both localized brane “sources.” Let us insert the above expressions into the original Lagrangian and also eliminate the auxiliary field $\mathcal{M}$, whose equation of motion gives

$$\mathcal{M} = -\sum_{I=1,2} 2\kappa_* \Delta(y_I) \left[1 + \frac{4\kappa^2}{3} \Delta(y_I) \sum_i |\phi_i^{IV}|^2\right]^{-1}$$

$$\times \left[\left(W_{IV} + W_{IH}\right) + \frac{\kappa_*}{2} \sum_i |\phi_i^{IV}|^2\right].$$

(36)

In Eq. (36), the delta function $\Delta(y_I)$ is present in the denominator as well as the numerator, which seems surprising at first. However, to couple the 5D gravity multiplet supersymmetrically to 4D chiral fields, the delta function couplings are inevitable. In this paper, we will take a pragmatic approach to the delta function; we regard the $\delta(y - y_I)$ functions as finite walls of height $M_*$ with a small but non-vanishing “thickness” $1/M_*$. (The exact shape of the delta function would be determined by a fundamental theory, but it would not affect the low energy physics.) Despite their thickness, both walls are assumed to be sufficiently far apart from each other, so that

$$\Delta(y_I) \times \Delta(y_J) = 0 \quad \text{for} \quad I \neq J.$$

(37)

Due to Eq. (37), $\mathcal{M}$ in Eq. (36) is given by the sum of the contributions from the two branes. In a more fundamental theory, the “blow-up” of the delta function could be cured and excitations of localized fields on the branes in the $y$ direction would be
possible. The apparently divergent operator in the denominator in Eq. (36) is now suppressed,

\[ \kappa_s^2 \Delta(y_I) |\phi_i^I|^2 \sim \frac{L}{M_*} |\phi_i^I|^2, \]

and we can treat it perturbatively unless \( \langle \phi_i^I \rangle \sim M_* \), as in ordinary 4D supergravity. In our case, we take

\[ L \sim 10^{-16} \text{ GeV}^{-1}. \]  

For the hidden sector fields and superpotentials we assume the VEVs [4]

\[ \langle W_{IH} \rangle = -m_I M_*^2, \]  

\[ \langle W_{zIH} \rangle = (b^*_I z) m_I M_*^2. \]  

For convenience, we define the following expressions:

\[ W^I \equiv W^I, \quad W^IV \equiv W^I, \quad \sum_i |W_i^IV|^2 \equiv |W_i^I|^2, \]  

\[ \sum_i |\phi_i^IV|^2 \equiv |\phi_i^I|^2, \quad \sum_i |b_i^I|^2 \equiv |b_i^I|^2. \]

After inserting Eq. (36) and Eqs. (40) and (41) into the original Lagrangian, assembling the results from the auxiliary fields \( F_{IS} \) and \( F_\Phi \) in Eqs. (33) and (34), we obtain the localized F term scalar potentials for \( \Phi \) and \( \phi_i^I \) in 5D supergravity,

\[ V_{5d}(\Phi, \phi_i^I) = \sum_{I=1,2} \frac{\Delta(y_I)}{L} \left[ |W_i^I|^2 + 2 \Delta(y_I) |W_i^I|^2 + |b_i^I|^2 (m_I M_*^2)^2 \right] - \frac{8\kappa_s^2}{3} \Delta(y_I) |W_i^I|^2 - m_I M_*^2 \frac{8\kappa_s^2}{3} \Delta(y_I) |\phi_i^I|^2 \right] + O(\kappa_s^6), \]

where we retain only terms up to \( O(\kappa^4) \) for examining the soft scalar masses and A terms. Elimination of \( \gamma \) gives rise to corrections of order \( O(\kappa^6) \) to the potential (44), and from the equations of motion for \( V_m^3 \) and \( v_m5 \), the derivative interaction terms between \( \phi_i^I \)s and \( F_m5 \) are derived. We should note here that in the localized potential, there is no cross term between the two brane contributions due to the delta function property in Eq. (37). Thus the potential at B1 is associated only with bulk scalar \( \Phi \)
and B1 brane fields $\phi^1_1$ and $\phi^1_2$, while the potential at B2 contains $\Phi$ and B2 brane fields $\phi^2_1$ and $\phi^2_2$.

We note the appearance of $\Delta(y_I)$ dependent couplings in the potential (44). In momentum space,

$$
\delta(y - y_I) = \frac{1}{2L} \sum_{n=-\infty}^{\infty} e^{i(n\pi/L)(y-y_I)} = \frac{1}{2L} \sum_{k_5=-\infty}^{\infty} e^{ik_5(y-y_I)},
$$

(45)

which is used in the Feynman rule for the delta function coupling for calculating certain processes [12, 14]. To obtain an effective 4D theory, we truncate the KK modes and introduce the cuttoff, $\Lambda_c \lesssim 1/L$ ($\sim 10^{16}$ GeV). Thus the summation in Eq. (45) should be implemented only for $k_5 = 0$ or $n = 0$, so that $\Delta(y - y_I)$ in the potential (44) reduces to 1/2 in low energy physics, independent of regularization schemes for the delta function. Cross interaction terms between fields localized on two different branes could be generated at loop level below $\Lambda_c \sim 1/L$, while above $\Lambda_c \sim 1/L$, the effective theory is 5D SUGRA and the “finite” delta function (or the finite wall) properties like Eq. (37) are restored.

Before deriving the effective 4D potential, let us discuss the delta function coupling for a moment. In Ref. [12], the supersymmetric coupling between 5D bulk super Yang-Mills and charged chiral matter localized on 4 dimensional brane was studied. In this model too, a delta function squared coupling appears after eliminating the auxiliary field $X^{3a}$. The 5D gauge multiplet contains a dynamical scalar $\Phi^a$ with odd parity under $Z_2$ symmetry, so that $X^{3a} - \partial_5 \Phi^a$ should be identified, from their SUSY transformation properties, with the auxiliary field $D^a$ appearing in the 4D gauge multiplet. The delta function coupling is interpreted as a counterterm to divergences arising from the derivative interactions $\partial_5 \Phi^a$ of an odd parity scalar with other chiral fields. However, such a dynamical odd parity scalar field is not found in the gravity multiplet as seen from Table I. Moreover, even if present, the scalar with odd parity would be decoupled at low energy [1]. In Ref. [14], an additional odd

\textsuperscript{5}In Ref. [14], a 5D SUGRA coupled to a 4D super Yang-Mills was considered, where a delta function squared coupling appears also in the 5D Lagrangian. Since a delta function squared coupling
parity scalar included in a hypermultiplet plays the role of $\Phi^a$ in Ref. [12]. But SUSY transformations of the 5D SUGRA auxiliary fields do not necessarily require additional fields for proper identification with the 4D auxiliary fields. Most pressing of all, in our model, such a method appears to be unable to control every delta function coupling.\footnote{This results from the fact that in general the Lagrangian of the 4D chiral multiplet contains quadratic terms in $\mathcal{M}$, while in previous examples the associated auxiliary fields, namely $D^a$ in Ref. [12] and $b_m$ in Ref. [14], appear in only linear form in the corresponding brane Lagrangian.} Thus we have adopted a different approach to the delta function coupling in which it is regarded as a wall with height $M_*$ and width $1/M_*$. Then, even if the wall of the delta function is as high as the fundamental scale $M_*$ or even the Planck scale, all delta function couplings are in turn small since they are always accompanied by the SUGRA parameter $\kappa_*$. However, in the low energy effective theory, the delta function couplings are not problematic in any case if the introduced cutoff ($1/L \sim 10^{16}$ GeV) is high enough to suppress any induced non-renormalizable operators.

The effective 4D potential is obtained by integrating $V_{5d}$ over the fifth direction $-L \leq y \leq L$,

$$V_{4d} = \sum_{I=1,2} (m_I M_P)^2 \left( |b^I|^2 - \frac{4}{3} \right) + |W^I|^2 + |W^I_t|^2$$

$$+ \frac{4m_I}{3} (W^I + W^I_t) + \frac{8(m_I)^2}{9} |\phi_{i_*}|^2$$

$$+ \text{non-renormalizable interactions} \quad (46)$$

In Eq. (46) the cosmological constant term and the ordinary F term potential of globally supersymmetric theory appears in the first line, the SUSY breaking $A$ terms and soft scalar mass terms are in the second line. The non-renormalizable terms are suppressed by powers of $1/M_*$. As in ordinary 4D supergravity, a vanishing cosmological constant requires fine-tuning. Thus we get a theory with softly broken SUSY but a vanishing cosmological constant through the fine-tuning, which is impossible in
globally supersymmetric theory. In contrast to the 4D case, the global SUSY F term potential is divided into two parts that depend on the locations of the associated fields \[25\]. Note that the potential in Eq. (46) does not allow a non-zero VEV for scalar fields, so the internal symmetry cannot be spontaneously broken at the fundamental scale. With \[\sum_i W^I_i \phi_i^I = 2W^I\], this can be proved by showing that while the minimum value of the potential always becomes non-zero and positive when a non-zero VEV for any scalar is assumed, it is vanishing for zero VEVs of all of the scalar fields \[1\]. It is possible that the internal symmetry breaks down radiatively at low energy, as in ordinary 4D SUGRA models.

From the scalar curvature terms in Eqs. (3) and (17), we see that the effective 4D potential in Eq. (46) is not yet in the Einstein frame. However, the canonical normalization of the effective 4D gravity just rescales the Yukawa coefficients and soft masses in Eq. (46) by a small amount, and adds non-renormalizable interactions.

As seen in Eq. (46), a soft mass term of the bulk scalar is not generated at tree level under the formalism employed in this paper. It is basically because the terms \[\kappa^2 A^\alpha_j A^\beta_j |t|^2\] (or \[\kappa^2 |\Phi|^2 |\mathcal{M}|^2\]) and \[A^\alpha_j A^\beta_j R_{AB}^{AB}\] (or \[|\Phi|^2 R_{AB}^{AB}\]) are cancelled out from the Lagrangian \[9\]. If the term \[\kappa^2 |\Phi|^2 |\mathcal{M}|^2\] exists in the Lagrangian, the bulk scalar could couple to the superpotentials \[W^I\] on the branes by Eq. (32) when eliminating the auxiliary field \[\mathcal{M}\]. However, a more general formalism might allow the soft mass term of the bulk scalar to be generated at tree level, under the condition that the term \[|\Phi|^2 R_{AB}^{AB}\] is present (generally when the metric on a quaternionic manifold with coordinates \[A^I_{\alpha}\] has a non-trivial form), a soft mass term of the bulk scalar can be also generated through the coupling between the bulk scalar and the F term of a hidden sector superfield like \[\kappa^2 |\Phi|^2 |\mathcal{F}_z^{IH}|^2\] (\[\sim |\Phi|^2 m^2_t M^2_s / M^2_P\]) in the Lagrangian. Such a coupling could be obtained after canonically normalizing the 5D gravity kinetic term. The coupling could be important in phenomenology \[26\]. Here we will just follow the formalism of Ref. \[15\]. In this formalism, a soft mass term of the bulk scalar can be generated at
one loop. We will discuss it later.

4 Non-universal SUSY breaking soft parameters

In the potential (46), we note that the soft mass term of the bulk scalar $\Phi$ is not generated by localized SUSY breaking at tree level. On the other hand, fields localized on the branes $\phi^I_i$ obtain their soft masses only from their associated branes. Since the mass parameters $m_I$ are different in general,

$$ m_I \neq m_J \quad \text{for} \ I \neq J , \quad (47) $$

$A$ terms and soft scalar masses of brane fields are also different at tree level, unless the relevant brane scalar fields live on the same brane. This consequence is definitely different from that of the ordinary 4D SUGRA scenario. The latter corresponds to taking the limit $L \rightarrow 0$ from the start and neglecting the delta functions in our model. Then, the auxiliary field $\mathcal{M}$ does not discern the location of fields, and cross terms between the two brane contributions appear, for example, from $|\mathcal{M}|^2$ as well as from $\mathcal{M}$. Thus all of the $A$ terms and soft scalar masses receive contributions from both $\langle W^{1H} \rangle$ and $\langle W^{2H} \rangle$, and universality of SUSY breaking soft parameters is observed.

The different SUSY breaking masses are confirmed by cross-checking the gravitino mass terms. With Eqs. (22)–(26), the third term in Eq. (27) provides the localized gravitino masses,

$$ m_{3/2,I} \sim k^2_+ \langle W^{1H} \rangle \sim m_I . \quad (48) $$

Therefore the gravitino acquires mass $m_1$ from B1 and $m_2$ from B2, so that the SUSY breaking scales are $O(m_1)$ and $O(m_2)$ at B1 and B2, respectively.

Although the SUSY breaking effects are generated through the direct coupling appearing in SUGRA at tree level and they are localized on the corresponding branes, they can be transmitted through “loops” to the bulk and even to the other branes below the compactification scale $[14]$. The off-shell Lagrangian of the hypermultiplet
contains the interaction term with the gravitino and $t^2$ \[^{18}\],

$$L_{HYP} \supset -\frac{\kappa^2}{L} A^\alpha \tau^2 \bar{\zeta}_\alpha \gamma^m \psi_m \; t^2,$$  \tag{49}

which survives under the $S^1/Z_2$ compactification. The VEV $\langle t^2 \rangle = \langle M \rangle \sim \Sigma_I \Delta (y_I) m_I M_s$ from SUSY breaking at the branes makes a bulk scalar’s soft mass generated at one loop, so that \[^{14}\]

$$\delta m^2_\Phi \sim \sum_I \frac{1}{16\pi^2} \left( \frac{m_I M_s}{M_p^2} \right)^2 \frac{1}{L^2},$$  \tag{50}

where $1/L$ is the cutoff $\Lambda_c$. Note that for $1/L \sim 10^{16}$ GeV, $M_s \sim 10^{17}$ GeV, and $m_1 \sim 10$ TeV, $\delta m_\Phi$ is of order a few GeV, which is quite small compared to other SUSY breaking soft masses.

We could exploit these phenomena for resolving the notorious flavor changing problem in SUGRA. We assume that

- the first two MSSM generations reside at B1;
- the third generation is at B2;
- the two Higgs multiplets as well as gravity and gauge multiplets are in the bulk;
- SUSY breakings arise from two hidden sectors localized on branes, and are subsequently transmitted by gravity. The SUSY breaking scale at B1 should be suitably higher than the breaking scale at B2.

Let us now consider gaugino masses. In Ref. \[^{17}\], the gaugino can become massive when SUSY is broken in the hidden sector,

$$L_{SYM}^{\text{bulk}} = \frac{1}{2L} \text{Tr} \left[ -\frac{1}{4} F_{MN} F^{MN} + \cdots - 2\kappa \bar{\lambda} \tau \bar{\tau} \lambda \; t + \cdots \right],$$  \tag{51}

where $\lambda$ is the gaugino which satisfies the symplectic Majorana condition. The coupling of the auxiliary field $\bar{t}^7$ to $\bar{\lambda} \tau \lambda$ means that the visible sector gaugino becomes massive either from the hidden sector gaugino condensates $\langle \bar{\lambda}^H \tau \lambda^H \rangle \neq 0$, or if the hidden sector superpotential on brane obtains a VEV $\langle W^H \rangle \neq 0$. After orbifolding, only $t^1$ and $t^2$ (or $\tau^1$ and $\tau^2$) contributions survive at low energy, and yield the 4D
Majorana mass term. With \( \langle t^2 \rangle = \langle M \rangle \sim \sum_I \Delta(y_I) m_I M_\star \) from Eqs. (36) and (40), the gaugino mass at low energy is given

\[
m_{1/2} \sim \frac{m_1 M_\star}{M_P}.
\]

Note that the gaugino mass is lighter for a larger extra dimension. With \( m_1 = 10 \) TeV and \( M_\star = 10^{17} \) GeV (or \( 1/L \sim 10^{16} \) GeV), the gaugino mass is of order 1 TeV, which manages to preserve the gauge hierarchy solution. This could be decreased to a few hundred GeV for \( 1/L \sim 10^{12} \) GeV.

Let us consider another possibility for lowering the gaugino mass. We can introduce an additional set of the localized gauge multiplet \( (A_m, \lambda, D) \) at B2, which respects the same gauge symmetry as the bulk gauge multiplet does. Consider the Lagrangian

\[
\mathcal{L}_{SYM} = \left[ \mathcal{L}_{SYM}^{\text{bulk}} + N \delta(y - L) \left( -\frac{1}{4} F_{mn} F^{mn} + i \bar{\lambda} \gamma^m D_m \lambda + \cdots \right) \right],
\]

with the same gauge coupling. Below the compactification scale \( 1/L \), the coefficient of the gauge kinetic term is given by \(- (1 + N)/4\). Thus canonical normalization of the kinetic term results in the suppressed (bulk) gauge coupling \( g/\sqrt{1 + N} \), as well as a suppressed gaugino mass \( m_{1/2}/(1 + N) \). Hence with \( N \sim O(10) \) (and strong coupling \( g \sim O(1) \)) and \( 1/L \sim 10^{16} \) GeV, a gaugino mass in the 100 GeV range can be achieved. In such a way, the masses of the various gauginos in MSSM also could be made non-universal.

To realize electroweak symmetry breaking in the MSSM, both the \( \mu \) term and the “\( B \)” term (\( \sim \mu m_{3/2} \)) of the right magnitudes must be present. Hence it is desirable that the \( \mu \) term is located at B2. From the weight assignments for hypermultiplets, the ordinary bilinear \( \mu \) term is not allowed in our example, so we should assume that the \( \mu \) term derives either from some trilinear or from a non-renormalizable interaction between two Higgs and some singlet field localized at B2 which develops a suitable non-zero VEV. We will not pursue this further in this paper.
The electroweak Higgs doublets reside in the bulk, and their soft mass squareds at $M_*$ arise predominantly through loops involving quarks and squarks rather than loops involving fields in the gravity multiplet as in Eq. (50). But their values at $M_*$ are also much smaller than the weak scale. The large top Yukawa coupling and the heavy first two generations’ soft mass squareds will drive this to $-(\text{a few TeV})^2$ at $M_Z$ as in the 4D ESUSY. Consequently, an adjustment of TeV scale parameters may be needed to achieve the desired 100 GeV electroweak symmetry breaking scale.

Finally, mixing of the first two generations with the third generation can arise effectively through one loop interactions below the compactification scale by introducing suitable heavy bulk fields belonging to vector-like representations of the SM gauge symmetry that couple to the ordinary MSSM particles on the branes. For neutrino mixings, we could introduce right-handed neutrinos in the bulk. We will not go into detail here.

5 Conclusion

In conclusion, through the use of off-shell SUGRA formalism and brane world framework, we have shown that gravity mediated SUSY breaking effects can be non-universal. With SUSY breakings in two hidden sectors localized on distinct branes, the gravity mediated soft masses of the bulk visible sector, gauginos, and the gravitino result from both branes, while SUSY breaking tree level effects in the localized visible sector fields arise only from the hidden sector living at the same brane. The soft masses for the bulk scalars are generated radiatively. This enables us to overcome the SUSY flavor problem geometrically by generating soft scalar masses $\sim$10–20 TeV for the first two generations, while the third generation masses are of order a TeV. Radiative electroweak breaking can be realized in this approach, with TeV scale quantities conspiring to realize the 100 GeV electroweak scale.

Acknowledgments

We thank A. Riotto and G. Dvali for useful discussions. The work is partially sup-
ported by the DOE under contract number DE-FG02-91ER40626.

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