

Radius stabilization and anomaly-mediated supersymmetry breaking

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We analyze in detail a specific 5-dimensional realization of a “brane-universe” scenario where the visible and hidden sectors are localized on spatially separated 3-branes coupled only by supergravity, with supersymmetry breaking originating in the hidden sector. Although general power counting allows order $1/M_{\text{Planck}}^2$ contact terms between the two sectors in the 4-dimensional theory from exchange of supergravity Kaluza-Klein modes, we show that they are not present by carefully matching to the 5-dimensional theory. We also find that the radius modulus corresponding to the size of the compactified dimension must be stabilized by additional dynamics in order to avoid runaway behavior after supersymmetry breaking and to understand the communication of supersymmetry breaking. We stabilize the radius by adding two pure Yang-Mills sectors, one in the bulk and the other localized on a brane. Gaugino condensation in the 4-dimensional effective theory generates a superpotential that can naturally fix the radius at a sufficiently large value that supersymmetry breaking is communicated dominantly by the recently discovered mechanism of anomaly mediation. The mass of the radius modulus is large compared to $m_{3/2}$. The stabilization mechanism requires only parameters of order one at the fundamental scale, with no fine-tuning except for the cosmological constant.

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I. INTRODUCTION

Supersymmetry (SUSY) breaking communicated by supergravity (SUGRA) is a very natural and attractive solution to the hierarchy problem. In its usual incarnation, this mechanism requires only a hidden sector that breaks SUSY, and the presence in the effective theory below the Planck scale of the following higher-dimension operators connecting the hidden and visible sector fields:

$$\begin{aligned} \mathcal{L}_{\text{eff}} \sim & \int d^4\theta \frac{1}{M_4^2} \Sigma^\dagger \Sigma [Q^\dagger Q + (H_u H_d + \text{H.c.})] \\ & + \int d^2\theta \frac{1}{M_4} [\Sigma W^\alpha W_\alpha + (\Sigma^\dagger H_u H_d + \text{H.c.})]. \end{aligned} \quad (1.1)$$

Here, M_4 is the 4-dimensional Planck scale, Σ is a field in the hidden sector with $\langle F_\Sigma \rangle \neq 0$, Q is a matter field in the visible sector, $H_{u,d}$ are Higgs fields, and W_α is a field strength for the standard model gauge group. This simple setup generates all required soft SUSY breaking terms of order $\langle F_\Sigma \rangle / M_4 \sim m_{3/2}$ (including the μ term [1]).

The main drawback of this scenario is that it does not explain why the squark masses generated from the term $\int d^4\theta \Sigma^\dagger \Sigma Q^\dagger Q$ approximately conserve flavor, as required to avoid excessive flavor-changing neutral currents. Reference [2] proposed an elegant solution to this problem in the context of higher-dimensional theories. It was pointed out that if

the visible and hidden sectors are localized on spatially separated “3-branes,” then contact terms of the form Eq. (1.1) can be suppressed even though they are not forbidden by any symmetry of the low-energy theory. This can be easily understood by focusing on the effective D -dimensional theory ($D > 4$) below the string scale M , but above the compactification scale $1/r$. If the hidden and visible branes are separated by a distance of order r , then the contribution from the exchange of bulk fields of mass $M \gg 1/r$ is suppressed by the Yukawa factor e^{-Mr} . We also expect the contributions of extended objects with string-scale tensions to be exponentially suppressed by e^{-Mr} . We see that stringy physics generates only exponentially small contact terms in the D -dimensional effective theory. When we match the D -dimensional theory to the 4-dimensional low-energy effective theory, a more careful analysis is required to show that the exchange of supergravity Kaluza-Klein (KK) excitations does not lead to contact Kähler terms of order $1/M_4^2$. We perform this analysis for a specific model in this paper, with the result that no such terms are generated. Thus, all of the effective interactions of the form of Eq. (1.1) are highly suppressed, and SUSY breaking must be communicated in a different way.

Reference [2] further argued that, given this suppression of the terms in Eq. (1.1), the leading contribution to SUSY breaking in the visible sector arises at loop level, and is directly related to the conformal anomaly. This mechanism applied to gaugino masses and A terms was independently discovered in Ref. [3], which also gave a detailed discussion of the exactness of the result. In this “anomaly-mediated supersymmetry breaking” (AMSB) scenario, all soft SUSY breaking parameters are completely predicted up to an overall scale by anomalous dimensions and conserve flavor to a high degree. This leads to interesting testable predictions for

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the gaugino masses [2–4]. Unfortunately, the slepton mass-squared terms are predicted to be negative if the visible sector is the minimal supersymmetric standard model. There have been several suggestions in the literature for natural solutions to this problem [5].

In this paper we investigate the basic features of the AMSB scenario in detail in a specific 5-dimensional effective field theory. The theory consists of minimal 5-dimensional SUGRA compactified on a S^1/Z_2 orbifold. The two $(3+1)$ -dimensional boundaries of this space corresponding to the orbifold fixed-points serve as the “3-branes” on which the hidden and visible sectors are localized. The higher SUSY of the 5-dimensional theory is broken explicitly down to $\mathcal{N}=1$ in 4 dimensions by the orbifold projection. This setup is very similar to the five-dimensional effective theory arising from heterotic M theory after Calabi-Yau compactification of six of the eleven dimensions [6–9], but our field content is the minimal one required for consistency of the five-dimensional effective theory. In particular, the Calabi-Yau moduli do not appear as light fields in our five-dimensional model. This accounts for the substantial differences between AMSB and other analyses of supersymmetry breaking in the heterotic M-theory scenario. We defer consideration of non-minimal field content for later work. Our final result is that anomaly mediation is the leading source of SUSY breaking in the visible sector if the radius is sufficiently large, but it is crucial to take into account the dynamics of the radius of the compactified dimension. While our analysis is limited to a specific 5-dimensional theory with a particular mechanism for stabilizing the radius, we believe that these features are more general.

Starting with the 5-dimensional theory described above, we construct the 4-dimensional effective theory below the compactification scale to analyze SUSY breaking. As already mentioned, a crucial feature of the effective theory is the presence of a radius modulus corresponding to the size of the compactified dimension. In particular, if this modulus is not stabilized we will show that its equations of motion set to zero the supersymmetry breaking order parameter for AMSB, namely the four-dimensional SUGRA auxiliary scalar. This agrees with a direct five dimensional SUGRA analysis, where there are no bulk fields which can transmit the effect of such an order parameter to the visible sector. This naturally raises doubts as to whether AMSB occurs in this scenario [9]. A related issue is the fact that the radius modulus must be stabilized in order to cancel the cosmological constant in the presence of SUSY breaking. We show that if the bare bulk cosmological constant is zero, there is no potential for the radius modulus, but the low-energy cosmological constant cannot be cancelled. In the presence of a bulk cosmological constant, SUSY breaking gives this modulus a runaway potential.

The picture changes completely when a stabilization mechanism is introduced for the radius. We propose a stabilization mechanism for the radius modulus that relies entirely on gaugino condensation and SUSY breaking. The mechanism requires two super-Yang-Mills (SYM) sectors, one in the bulk and one localized on a 3-brane, as well as a SUSY

breaking sector localized on the hidden 3-brane. Upon matching to the 4-dimensional theory, the bulk SYM sector gives rise to a 4-dimensional SYM sector with a gauge coupling that depends on the radius r . This gives rise to an r -dependent gaugino condensate which, together with the brane-localized gaugino condensate, gives a stabilizing potential for the radius modulus.¹ The radius is naturally large compared M_5 if the condensation scale of the 3-brane super-Yang-Mills sector Λ_{bdy} is small compared to M_5 . The radius depends only logarithmically on Λ_{bdy} , we can obtain a sufficiently large radius for anomaly mediation to dominate if the theory is strongly coupled near the scale M_5 .² Since the condensation scale of the 3-brane super-Yang-Mills sector is naturally exponentially small compared to the fundamental scale, this mechanism does not require the introduction of small parameters at the fundamental scale. SUSY breaking (and fine-tuning) is required to cancel the net low-energy cosmological constant. The mass of the radius modulus is large compared to $m_{3/2}$, and the effective theory below this scale is of the “sequestered” form proposed in Ref. [2]. The general lesson we draw from this is that AMSB works provided that moduli are stabilized. Our stabilization is similar in spirit to the racetrack mechanism [12], but it does not require large gauge groups and our results follow from a completely systematic effective field theory analysis.

This paper is organized as follows. In Sec. II we describe the 5-dimensional model and carry out the matching to the 4-dimensional effective theory. We show that there are no $O(1/M_4^2)$ contact terms between the hidden and visible sectors, and that the cosmological constant cannot be cancelled in the absence of a mechanism for radius stabilization. In Sec. III, we show how gaugino condensation can fix the radius, and show that anomaly mediation works in this scenario. Section IV contains our conclusions.

II. FROM 5 TO 4 DIMENSIONS

A. The 5-dimensional model

We consider minimal (ungauged) 5-dimensional SUGRA compactified on a S^1/Z_2 orbifold with matter and gauge fields localized on the two orbifold boundaries. This system is relatively simple to study because the orbifold projection explicitly breaks the supersymmetry of the 5-dimensional theory (8 real supercharges) down to $\mathcal{N}=1$ in 4 dimensions (4 real supercharges).

The on-shell Lagrangian for the bosonic fields of 5-dimensional SUGRA is [13]

¹A similar mechanism can be used to stabilize the radius in non-supersymmetric theories [10]. This may be interesting for solutions of the hierarchy problem involving extra dimensions that are only slightly larger than the fundamental scale [11].

²It is interesting that even for strong coupling, this mechanism gives a radius that is naturally close to the scale where bulk gravitational loops give a contribution to soft masses comparable to the contributions of anomaly mediation. We will not pursue this possibility here.

TABLE I. Bosonic fields of 5-dimensional SUGRA with their Z_2 parity assignments. The parity assignments of the graviphoton fields are fixed by the Chern-Simons term.

Field	Z_2 parity
$g_{\mu\nu}$	+
$g_{5\mu}$	-
g_{55}	+
B_μ	-
B_5	+

$$\mathcal{L}_{\text{SUGRA},5} = -M_5^3 \left[\sqrt{-g^{(5)}} \left(\frac{1}{2} \mathcal{R}^{(5)} + \frac{1}{4} H^{MN} H_{MN} \right) + \frac{1}{6\sqrt{6}} \epsilon^{MNPQR} B_M H_{NP} H_{QR} + \text{fermion terms} \right], \quad (2.1)$$

where $M, N, \dots = 0, \dots, 3, 5$, are 5-dimensional spacetime indices, and $H_{MN} = \partial_M B_N - \partial_N B_M$ is the field strength for the graviphoton B_M . Under the Z_2 parity, the fields transform as $\phi(x^5) \mapsto \pm \phi(-x^5)$, where the parity assignments of the bosonic fields are given in Table I. The orbifold projection keeps only those field configurations that are even under Z_2 .

We assume that there are fields localized on the orbifold boundaries, so these must be coupled to SUGRA. The Lagrangian has the form

$$\mathcal{L}_5 = \mathcal{L}_{\text{SUGRA},5} + \delta(x^5) \mathcal{L}_{\text{vis}} + \delta(x^5 - \pi r) \mathcal{L}_{\text{hid}}. \quad (2.2)$$

We will not need the details of the bulk-boundary couplings in \mathcal{L}_{vis} and \mathcal{L}_{hid} , but it is important for us to know that such couplings exist and preserve $\mathcal{N}=1$ SUSY. As shown in Ref. [14] for 5-dimensional gauge- and hypermultiplets, the couplings of bulk and boundary fields can be worked out in a straightforward fashion if the auxiliary fields of the bulk theory are known. Building on earlier work [15], an explicit off-shell formulation for 5-dimensional SUGRA was recently given by Zucker [16]. Following Ref. [14], one first decomposes the 5-dimensional SUGRA multiplet into off-shell multiplets of the unbroken 4-dimensional $\mathcal{N}=1$ SUSY. In addition to the $\mathcal{N}=1$ SUGRA multiplet, this yields two vector multiplets (with vector fields $g_{5\mu}$ and B_μ) with odd orbifold parity, and one chiral multiplet (with real scalar fields g_{55} and B_5) with even parity. It should then be possible to couple these multiplets to $\mathcal{N}=1$ fields localized on the boundaries using the usual $\mathcal{N}=1$ superfield calculus.

B. Matching to 4 dimensions

We now consider integrating out the KK modes of the 5-dimensional SUGRA multiplet at the scale r to obtain a 4-dimensional effective theory. We are interested in effects of order $1/M_4^2 \sim 1/(rM_5^3)$, which means that we can restrict attention to tree-level effects in the SUGRA fields. [In the normalization of the SUGRA fields given in Eq. (2.1), the propagator for all bosonic SUGRA fields is of order $1/M_5^3$.]

There are SUGRA loop effects suppressed by additional powers of $1/(M_5 r)^3 \sim 1/(M_4 r)^2$. For some values of r these effects could be interesting [2]. Here we will simply assume that r is sufficiently large that these loop effects can be neglected. The matching of the SUGRA fields at the tree level is performed simply by using the metric

$$ds^2 = g_{\mu\nu}(x) dx^\mu dx^\nu + r^2(x) d\vartheta^2, \quad (2.3)$$

where $\vartheta \in [0, \pi]$ is a coordinate for the compact dimension, and $g_{\mu\nu}(x)$, $r(x)$ parametrize the massless metric and radius modulus fields. (We are implicitly expanding about a flat metric, so the zero-mass KK modes are independent of ϑ .) Ignoring the boundary fields for the moment, the bosonic terms in the 4-dimensional effective theory are

$$\mathcal{L}_4 = -2\pi M_5^3 \sqrt{-g^{(4)}} \left[\frac{r}{2} \mathcal{R}^{(4)} + \frac{1}{2r} \partial^\mu B_\vartheta \partial_\mu B_\vartheta \right]. \quad (2.4)$$

Note that there is no explicit kinetic term for the radius modulus. After an r -dependent Weyl rescaling of the metric, a kinetic term for the radius modulus is generated (with the correct sign). The couplings of the radius modulus to boundary fields is very different in the two bases. Before Weyl rescaling, there are no couplings of r to boundary fields at leading order in the low-energy expansion. This is because r arises from fluctuations of g_{55} , which by general covariance can only couple to the 55 component of the matter stress tensor. This component vanishes for matter confined to 3-branes, at leading order in $1/M_5$. At higher order in derivatives and $1/M_5$, we can write terms containing the curvature tensor that depend on derivatives of r but these will be a small correction. In the rescaled basis the radius modulus has non-derivative couplings to fields localized on the branes.

Equation (2.4) is to be matched to the most general Lagrangian describing 4-dimensional SUGRA coupled to a modulus T . Using the superconformal approach to SUGRA [17], this can be written as

$$\mathcal{L}_{\text{SUGRA},4} = \int d^4 \theta \phi^\dagger \phi f(T^\dagger, T), \quad (2.5)$$

where

$$\phi = 1 + \theta^2 F_\phi \quad (2.6)$$

is the conformal compensator. We do not include a superpotential in Eq. (2.5) because T has no potential in this approximation. (Recall that we are not including a bulk cosmological constant.) After integrating out the auxiliary fields, the bosonic terms of Eq. (2.5) are

$$\begin{aligned} \mathcal{L}_{\text{SUGRA},4} = & \sqrt{-g^{(4)}} \left[\frac{1}{6} f \mathcal{R}^{(4)} - \frac{1}{4f} (f_T \partial^\mu T - \text{H.c.}) \right. \\ & \times (f_T \partial_\mu T - \text{H.c.}) - f_{T^\dagger T} \partial^\mu T^\dagger \partial_\mu T \\ & \left. + \text{fermion terms} \right], \end{aligned} \quad (2.7)$$

where $f_T = \partial f / \partial T$, etc. An important point is that Eq. (2.7) must be matched to Eq. (2.4) *without* Weyl rescaling. The reason is that if boundary fields are included, the theory expressed in terms of the Weyl-rescaled metric contains only non-derivative couplings to the radius modulus. Kahler terms involving both T and boundary terms necessarily contain derivative interactions of T , the only consistent way to match is if the Kahler terms are T -independent. Equation (2.4) then shows that there is an explicit kinetic term for only one of the real scalar fields in T , so we must have $f_{T^\dagger T} = 0$. This implies that f is the sum of a holomorphic plus antiholomorphic function, so we can make a field redefinition so that $f = -M_5^3 \cdot (T + T^\dagger)$. Writing $T = T_1 + iT_2$, we have

$$\begin{aligned} \mathcal{L}_{\text{SUGRA},4} = & -M_5^3 \int d^4 \theta \phi^\dagger \phi (T + T^\dagger) \\ & = -M_5^3 \sqrt{-g^{(4)}} \left[\frac{T_1}{3} \mathcal{R}^{(4)} + \frac{1}{2T_1} \partial^\mu T_2 \partial_\mu T_2 \right. \\ & \left. + \text{fermion terms} \right]. \end{aligned} \quad (2.8)$$

Comparing this with Eq. (2.4), we can identify

$$\text{Re}(T) = 3\pi r, \quad \text{Im}(T) = \sqrt{6} \pi B_\vartheta. \quad (2.10)$$

Equation (2.8) has the ‘‘no-scale’’ form considered long ago [18]. The essential new ingredient in the present case is that the no-scale form is stable under radiative corrections because the cutoff of the 4-dimensional theory is of order $1/r \ll M_4$.

We now consider the fields localized on the orbifold boundaries. We are particularly interested in contact interactions between the hidden and visible sectors. The only contact interaction of order $1/M_4^2$ in the 4-dimensional effective theory that is not forbidden by symmetries is

$$\frac{1}{M_4^2} \int d^4 \theta (\Sigma^\dagger \Sigma) (Q^\dagger Q) = \frac{4}{M_4^2} (\psi_\Sigma \psi_Q) (\bar{\psi}_\Sigma \bar{\psi}_Q) + \dots \quad (2.11)$$

where we have explicitly shown the 4-fermion component. The only diagrams that can contribute to the 4-fermion term in Eq. (2.11) at order $1/M_4^2$ consist of tree-level exchange of bosonic SUGRA fields. The bulk-boundary couplings cannot involve any suppression by $1/M_5$, otherwise the final result will be less than $1/M_4^2$. It may appear that these conclusions are invalidated by power-divergent loop graphs with a cutoff of order M_5 . However, general renormalization theory tells

us that the divergent contributions will have the same structure as local terms in the effective field theory, and therefore do not give new effects.

Now, the exchange of Kaluza-Klein excitations of the graviton couple to derivatives of the fermion fields and therefore cannot yield a term of the form Eq. (2.11). Couplings of the graviphoton to boundary fields are restricted by the orbifold projection and graviphoton gauge invariance

$$\delta B_M = \partial_M \alpha, \quad \alpha(-x^5) = -\alpha(x^5). \quad (2.12)$$

Boundary fields cannot be charged under this symmetry because B_μ vanishes on the boundary. The only term consistent with these constraints that can give rise to the 4-fermion term in Eq. (2.11) has the form

$$\Delta \mathcal{L}_5 = \delta(x^5) H_{5\mu} K_{\text{vis}}^\mu + \delta(x^5 - \pi r) H_{5\mu} K_{\text{hid}}^\mu \quad (2.13)$$

where K^μ is a dimension-3 current constructed from boundary fields; its precise form will be determined by matching to the 4-dimensional theory.

The power-counting argument above shows that Eq. (2.13) will give rise to contact terms of order $1/M_4^2$ from tree-level exchange of B_μ fields. We can determine these terms by integrating out B_μ using its classical equations of motion. Imposing periodicity and consistency with the orbifold projection, we obtain

$$\partial_5 B^\mu = \frac{1}{M_5^3} \left[\delta(y) K_{\text{vis}}^\mu + \delta(y - \pi r) K_{\text{hid}}^\mu - \frac{1}{2\pi r} (K_{\text{vis}} + K_{\text{hid}})^\mu \right]. \quad (2.14)$$

In this computation it was important that we considered the B_ϑ field to be independent of x_5 , corresponding to the zero-mode ($\text{Im} T$) of the five-dimensional field. Substituting back into the Lagrangian and integrating over the compact dimension to obtain the 4-dimensional effective theory, we obtain the contact terms³

$$\begin{aligned} \Delta \mathcal{L}_4 = & -\frac{1}{r} \partial_\mu B_\vartheta (K_{\text{vis}} + K_{\text{hid}})^\mu \\ & - \frac{1}{4\pi M_5^3 r} (K_{\text{vis}} + K_{\text{hid}})^\mu (K_{\text{vis}} + K_{\text{hid}})_\mu. \end{aligned} \quad (2.15)$$

We compare this with the contact terms in the 4-dimensional SUGRA with matter fields:

$$\mathcal{L}_4 = \int d^4 \theta \phi^\dagger \phi [-M_5^3 (T + T^\dagger) + f_{\text{vis}} + f_{\text{hid}}]. \quad (2.16)$$

³This procedure also gives rise to terms proportional to $\delta(0) \cdot (K_{\text{vis}}^2 + K_{\text{hid}}^2)$ in the 4-dimensional effective theory; these are cancelled by boundary terms proportional to $\delta(0)$ in the 5-dimensional theory. For a discussion of the origin of these terms, see Refs. [7,14].

As argued above, f_{vis} and f_{hid} are independent of T because any dependence would imply a coupling of r (and hence g_{55}) to brane fields (without Weyl rescaling). We therefore obtain

$$\mathcal{L}_4 = -\frac{1}{2T_1} \partial_\mu T_2 J^\mu - \frac{1}{8M_5^3 T_1} J^\mu J_\mu + \dots \quad (2.17)$$

Here $J^\mu = J_{\text{vis}}^\mu + J_{\text{hid}}^\mu$ with

$$J^\mu = i(f_a \partial^\mu \phi^a - \text{H.c.}) + f^a_b \psi_a \sigma^\mu \bar{\psi}^b, \quad (2.18)$$

where $f^a_b = \partial^2 f / (\partial \Phi_a^\dagger \partial \Phi^b)$, etc.

Comparing Eqs. (2.15) and (2.17) and using Eq. (2.10), we see that matching the $\partial_\mu B_\nu K^\mu$ term requires

$$K^\mu = \frac{1}{\sqrt{6}} J^\mu. \quad (2.19)$$

With this identification, the $J^\mu J_\mu$ contact terms also match. This matching would be spoiled by additional contact terms of the form Eq. (2.11), so we conclude that these operators are absent in the 4-dimensional effective theory.

Putting together the various pieces, the four-dimensional effective theory below the compactification scale has the general form

$$\mathcal{L}_4 = -M_5^3 \int d^4 \theta \phi^\dagger \phi (T + T^\dagger) + \mathcal{L}_{\text{hid}} + \mathcal{L}_{\text{vis}}, \quad (2.20)$$

where \mathcal{L}_{hid} is made out of only hidden sector and four-dimensional supergravity (off-shell) multiplets and \mathcal{L}_{vis} is made out of only visible sector and four-dimensional supergravity multiplets. Both are independent of the T chiral multiplet.

C. The role of the radius

We now consider SUSY breaking on the hidden-sector boundary in the theory above. We will show that the presence of an unstabilized radius modulus gives rise to severe difficulties in this scenario when SUSY is broken in the hidden sector.

Independently of how SUSY is broken, it is easy to see from Eq. (2.20) that the F_T equation of motion sets $F_\phi = 0$. This implies that there are no contact terms between the visible and hidden sectors in this theory, consistent with the fact that there is no propagating bulk scalar field in the SUGRA multiplet that could mediate such terms. This makes it rather mysterious how SUSY breaking can be communicated from the hidden to the visible sector [9], especially since F_ϕ is the order parameter for AMSB in the visible sector [2,3].

This feature also gives rise to difficulties in cancelling the cosmological constant. SUSY breaking on the hidden-sector boundary gives rise to a nonzero vacuum energy independent of the radius modulus T . In generic four-dimensional SUGRA models this positive contribution to the cosmological constant can be cancelled by negative SUGRA contributions arising from $F_\phi \neq 0$, but this mechanism is clearly not available here. One can attempt to remedy this by adding a

SUSY-preserving five-dimensional cosmological constant to the theory. To linear order in the cosmological constant, the effect of this is to add a superpotential term linear in T to Eq. (2.20). The potential arising from this theory is now

$$V = -\frac{k}{M_5^3} \text{Re}(T) + V_{\text{hid}}, \quad (2.21)$$

where V_{hid} is the vacuum energy from hidden sector SUSY breaking, and k sets the size of the bulk cosmological constant. However, this introduces a new problem, namely runaway behavior for the radius modulus.⁴ We see that we cannot obtain an appropriate setting for AMSB without adding new physics to stabilize the modulus.

We mention that another means of breaking SUSY is to *not* have a hidden sector which breaks SUSY by itself but rather to simply have a constant superpotential on a brane (and no bulk cosmological constant). Then one finds that $F_T \neq 0$, but $F_\phi = 0$, so SUSY is broken but the cosmological constant vanishes. This is the basic no-scale mechanism of SUSY breaking [18]. We do not pursue this scenario here because it involves the vanishing of the AMSB order parameter F_ϕ .

III. RADIUS STABILIZATION

We now show that the problems found above are solved by dynamically stabilizing the modulus. This modulus must be stabilized in any case for phenomenological reasons. (The radius modulus must have a mass larger than of order 1 cm^{-1} to avoid conflict with post-Newtonian tests of gravity [19].) We will focus on a specific mechanism for stabilizing the radius modulus that requires only a super-Yang-Mills (SYM) sector in the bulk, and another SYM sector on one of the boundaries. We assume that the bulk cosmological constant is negligible; this is natural because of the presence of bulk SUSY.

A. Bulk super-Yang-Mills sector

We begin by discussing the bulk SYM sector. At the compactification scale $1/r$, this theory matches onto a 4-dimensional SYM theory with a gauge coupling that depends on r . The scale where the effective 4-dimensional SYM theory becomes strong therefore depends on r , and gaugino condensation generates a dynamical superpotential that depends on the modulus T . The fact that the dynamical superpotential for T is generated by supersymmetric dynamics rather than induced by SUSY breaking in the hidden sector allows the mass of the modulus to be large compared to $m_{3/2}$. This means that below the scale of the radius modulus, the effective theory has the ‘‘sequestered’’ form dis-

⁴When T becomes sufficiently large, the linearized approximation for the effect of a bulk cosmological constant is no longer valid. We have checked that including the full non-linear effects does not stop the runaway behavior.

TABLE II. Fields of 5-dimensional super-Yang-Mills sector with their Z_2 parity assignments.

Field	Z_2 parity
A_μ	+
A_5	-
Φ	-
λ^1	+
λ^2	-

cussed in Ref. [2], and the leading contribution to SUSY breaking in the visible sector comes from anomaly mediation.

The bulk SYM multiplet consists of a vector field A_M , a real scalar Φ , and a symplectic Majorana gaugino λ^j ($j = 1, 2$). These fields are taken to transform under the orbifold projection as shown in Table II. The even fields form an $\mathcal{N} = 1$ SYM multiplet \mathcal{V} , while the odd fields form an $\mathcal{N} = 1$ chiral multiplet Ψ . These fields can be coupled to the boundary fields using the usual rules for constructing $\mathcal{N} = 1$ invariants. (For more details, see Ref. [14].)

We assume that the fields on the boundaries are uncharged under the bulk SYM sector. However, there are in general higher-dimension operators coupling the bulk SYM fields to the boundary fields. Using a normalization of the fields where the gauge coupling is factored out of the kinetic terms

$$\mathcal{L}_5 = \frac{1}{g_5^2} \text{tr} \left[-\frac{1}{4} F^{MN} F_{MN} + \partial^M \Phi \partial_M \Phi + \dots \right], \quad (3.1)$$

the bulk SYM propagator is proportional to $g_5^2 \sim 1/M_5$. Therefore, exchange of SYM fields between the boundaries can give rise to contact terms of order $1/M_4^2 \sim 1/(rM_5^3)$ only if there are boundary couplings of order $1/M_5$. However, it is easy to see that no such terms are possible unless there is a singlet S on the boundary, in which case we can write

$$\Delta \mathcal{L}_5 = \delta(y) \int d^2 \theta \frac{1}{M_5} S \text{tr}(\mathcal{W}^\alpha \mathcal{W}_\alpha) + \text{H.c.}, \quad (3.2)$$

where \mathcal{W}^α is the field strength of the $\mathcal{N} = 1$ SYM field \mathcal{V} . [Note that boundary couplings involving the $\mathcal{N} = 1$ chiral multiplet Ψ are restricted by gauge invariance $\delta \Psi = i \partial_5 \alpha$, $\alpha(-x^5) = +\alpha(x^5)$.] If there are singlets in both the hidden and visible sector, this will induce contact terms between them only at the 1-loop level, and the presence of two SYM propagators in the leading diagram means that the effects are suppressed by $1/M_5^4$, and therefore negligible. [The contact terms are Kahler terms by $U(1)_R$ invariance.] We conclude that introducing the bulk SYM sector does not introduce new contact terms into the effective 4-dimensional theory.

We now construct the 4-dimensional effective theory for the bulk SYM sector. When we perform the KK decomposition, the odd fields have KK masses starting at $1/r$ and are therefore integrated out. The even fields have a massless zero mode, which becomes a 4-dimensional SYM sector in the

effective theory. The tree-level matching condition for the effective 4-dimensional gauge coupling is

$$\frac{1}{g_4^2} = \frac{2\pi r}{g_5^2}. \quad (3.3)$$

Because g_4 depends on r , gaugino condensation in the effective 4-dimensional SYM sector will give rise to KK masses starting at $1/r$, and are therefore integrated T -dependent dynamical superpotential.

The T dependence of the dynamical superpotential can be determined exactly using holomorphy arguments [20]. The holomorphic 4-dimensional gauge coupling $S = 1/(2g_4^2) + \dots$ is given exactly by

$$S(\mu = 1/g_5^2) = \frac{2T}{3g_5^2} + c, \quad (3.4)$$

where c is a real constant that parametrizes the scheme dependence. It may appear that cancelling large logs requires us to match at a scale $\mu \sim 1/r$:

$$S(\mu = 1/T) = \frac{2T}{3g_5^2} + c. \quad (3.5)$$

However, for $\mu < 1/r$ this leads to [for an $SU(N)$ gauge group]

$$S(\mu) = \frac{2T}{3g_5^2} + \frac{3N}{16\pi^2} \ln(\mu T) + c. \quad (3.6)$$

The logarithmic dependence on T implies that $1/g_4^2 \propto \text{Re}(S)$ depends on $\text{Im}(T) \propto B_\vartheta$. But from the 5-dimensional theory we know that B_ϑ is derivatively coupled, so this is impossible. It is easy to see that the only way to avoid this contradiction consistent with holomorphy is Eq. (3.4). We have also checked that carefully evaluating the threshold corrections due to the infinite tower of SYM KK states also reproduces Eq. (3.4). The dynamical scale of the theory is therefore

$$\Lambda_{\text{bulk}} \propto \frac{1}{g_5^2} e^{-32\pi^2 T / (9Ng_5^2)}. \quad (3.7)$$

In order to obtain believable numerical estimates we need to estimate the constant of proportionality in Eq. (3.7). This can be done using ‘‘naive dimensional analysis’’ (NDA) [21,22].⁵ The principle of NDA is that in a strongly coupled theory with no small parameters, both the fundamental and the effective theory become strongly coupled (in the sense that loop corrections are order 1) at the same scale. To estimate Λ , note that NDA implies that if the gauge coupling and the radius are chosen so that the fundamental theory is

⁵NDA is applied to higher-dimension theories with branes in Ref. [23].

strongly-coupled at a scale Λ_0 , then $\Lambda \sim \Lambda_0$. The strong-coupling value of the 5-dimensional gauge coupling is

$$g_5^2|_{\text{strong}} \sim \frac{l_5}{N\Lambda_0}, \quad (3.8)$$

where $l_5 = 24\pi^3$ is the (inverse of the) 5-dimensional loop counting parameter, and we have taken into account the N dependence appropriate for the large N limit. The strong-coupling value of the radius is where the KK modes have mass of order Λ_0 :

$$r|_{\text{strong}} = \frac{1}{3\pi} T|_{\text{strong}} \sim \frac{1}{\Lambda_0}. \quad (3.9)$$

This implies that the strong-coupling value of the exponential in Eq. (3.7) is order 1, and we obtain

$$\Lambda_{\text{bulk}} \sim \frac{l_5}{Ng_5^2} e^{-32\pi^2 T/(9Ng_5^2)}. \quad (3.10)$$

The dynamical superpotential generated in the 4-dimensional effective theory is therefore

$$W_{\text{bulk,dyn}} \sim \frac{1}{Nl_4} \Lambda_{\text{bulk}}^3 \sim \frac{l_5^3}{l_4 N^4 g_5^6} e^{-32\pi^3 T/(3Ng_5^2)}. \quad (3.11)$$

Using $l_5 = 24\pi^3$ and $l_4 = 16\pi^2$, the dimensionless prefactor is $l_5^3/l_4 = 864\pi^7 \approx 3 \times 10^6$. However, this estimate depends sensitively on the value used for l_5 , and should be regarded as very uncertain. Nonetheless, it is clear that the prefactor will be large unless NDA is completely misleading.⁶

B. Boundary super-Yang-Mills sector

In addition to the bulk SYM sector, we assume that the theory contains a SYM sector localized on one of the boundaries. As with the bulk SYM, we assume that there are no matter fields charged under the SYM gauge group. If this SYM sector is in the hidden sector, there is no danger from flavor-violating higher-dimension contact terms. If it is in the visible sector, the lowest-dimension potentially flavor-violating operator is

$$\Delta\mathcal{L}_5 \sim \delta(y) \int d^4\theta \frac{1}{M_5^3} Q^\dagger Q \text{tr}(W^\alpha W_\alpha) + \text{H.c.}, \quad (3.12)$$

where W^α is the field strength of the boundary SYM multiplet. This gives flavor-violating interactions suppressed by $(\Lambda_{\text{bdy}}/M_5)^3$, where Λ_{bdy} is the dynamical scale. This is negligible for the values of Λ_{bdy} we will be interested in (see below), and we conclude that the boundary SYM sector may be either in the hidden or the visible sector.

⁶In this connection, it may be worthwhile to point out that exact results obtained in $\mathcal{N}=2$ theories spectacularly confirm the expectations of NDA [24].

C. 4-dimensional effective theory

Now we are ready to analyze the 4-dimensional effective theory, including all sectors. Below the scale $1/r$, the 4-dimensional theory consists of 4-dimensional SUGRA and the modulus T coupled to the bulk and boundary SYM sectors. In addition, the theory contains the visible and hidden sectors, which we do not specify explicitly. We now write the effective Lagrangian below the scales Λ_{bulk} and Λ_{bdy} where the SYM sectors become strong, and below the scale of SUSY breaking in the hidden sector. In this regime, the only light fields are the SUGRA fields, the modulus T , the Goldstino from the SUSY breaking sector, and the visible sector fields. The effective Lagrangian is

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & -M_5^3 \int d^4\theta \phi^\dagger \phi (T + T^\dagger) \\ & + \left(\int d^2\theta \phi^3 [c + a e^{-bT}] + \text{H.c.} \right) - V_{\text{hid}} + \dots \end{aligned} \quad (3.13)$$

Here

$$c \sim \frac{1}{l_4} \Lambda_{\text{bdy}}^3 \quad (3.14)$$

arises from gaugino condensation in the boundary SYM theory (we neglect N dependence in the boundary SYM theory);

$$a \sim \frac{l_5^3}{l_4 N^4 g_5^6}, \quad b = \frac{32\pi^2}{3Ng_5^2} \quad (3.15)$$

arise from gaugino condensation in the bulk SYM theory; and $V_{\text{hid}} > 0$ is the vacuum energy generated by the SUSY breaking sector. We have chosen not to add a 5-dimensional cosmological constant. The constant c can be chosen real by a $U(1)_R$ rotation, but a is in general complex. The terms omitted in Eq. (3.13) contain the interactions of the visible sector fields and a Goldstino from SUSY breaking in the hidden sector (which will eventually become the longitudinal components of the massive gravitino). The terms involving the Goldstino can be included using a non-linear realization of SUSY coupled to SUGRA [25], but are not relevant for computing the effective potential for T ; the same is true for the visible sector interactions.

The superpotential in Eq. (3.13) is exact, but the Kahler potential contains unknown $\mathcal{O}(1/M_4^4)$ corrections from loop corrections and higher-dimension operators. These will be shown to give small corrections below.

We now turn to the minimization of the scalar potential, neglecting corrections to the Kähler potential. The scalar potential obtained from Eq. (3.13) is

$$\begin{aligned} V = & \frac{1}{M_5^3} \{ (3c^* b a e^{-bT} + \text{H.c.}) \\ & + b [b(T + T^\dagger) + 6] |a|^2 e^{-b(T + T^\dagger)} \} + V_{\text{hid}}. \end{aligned} \quad (3.16)$$

Note that the first term is proportional to the boundary SYM gaugino condensate. Only the first term in Eq. (3.16) depends on $\text{Im}(T)$. Minimizing with respect to $\text{Im}(T)$, we obtain the effective potential for $T_1 = \text{Re}(T)$:

$$V = \frac{1}{M_5^3} \{ -6b|a||c|e^{-bT_1} + 2b(bT_1 + 3)|a|^2e^{-2bT_1} \} + V_{\text{hid}}. \quad (3.17)$$

The term in brackets is a sum of two different exponentials with opposite signs, the negative sign in the first term arising from the minimization with respect to $\text{Im}(T)$. As $T_1 \rightarrow \infty$ the first term dominates, and the potential approaches $+V_{\text{hid}}$ from below. Provided the second term dominates for small T_1 there will be a nontrivial minimum with vacuum energy below $+V_{\text{hid}}$. This means that the parameters can be adjusted to give a vanishing cosmological constant.

We look for a minimum with $b\langle T_1 \rangle \gg 1$. Explicitly carrying out the minimization we find that

$$b\langle T_1 \rangle e^{-b\langle T_1 \rangle} \approx \frac{3|c|}{2|a|} \sim \frac{N^4 \Lambda_{\text{bdy}}^3 g_5^6}{l_4 l_5^3}, \quad (3.18)$$

where we have neglected terms suppressed by powers of $1/(b\langle T_1 \rangle)$. Note that $\langle T_1 \rangle$ can be made arbitrarily large by making Λ_{bdy} small compared to $1/g_5^2$. (The loop suppression factors also tend to increase $\langle T_1 \rangle$.) The vacuum energy at the minimum is

$$\langle V \rangle \approx -\frac{3|c|^2}{M_4^2} + V_{\text{hid}}, \quad (3.19)$$

where $M_4^2 = M_5^3 \pi r$. The fact that the first term is negative allows us to choose the parameters to fine-tune the cosmological constant to zero.

Because the superpotential has non-trivial T dependence, the F_T equation of motion no longer sets $F_\phi = 0$. Instead we have

$$\langle F_\phi \rangle \approx \frac{|c|}{M_4^2}. \quad (3.20)$$

SUSY is broken, and the gravitino mass is

$$m_{3/2} \sim \frac{V_{\text{hid}}^{1/2}}{M_4} \sim \frac{|c|}{M_4^2}, \quad (3.21)$$

so that $\langle F_\phi \rangle \sim m_{3/2}$.

The mass of the radius modulus is computed from

$$\langle V'' \rangle \approx \frac{6b^2|c|^2}{M_4^2}, \quad (3.22)$$

where the primes denote differentiation with respect to T_1 . The kinetic term for T_1 arises from mixing with the metric; it can be made manifest by making a T_1 -dependent Weyl trans-

formation. This gives a kinetic term $\sim M_4^2 (\partial T_1)^2 / T_1^2$, and the physical mass of the radius modulus is

$$m_r^2 \sim \frac{b^2|c|^2}{M_4^4} \langle r \rangle^2. \quad (3.23)$$

It is easy to see that the other real scalar and the fermion component of T also get a mass of this order. Comparing with Eq. (3.21), we see that

$$\frac{m_r}{m_{3/2}} \sim b\langle r \rangle \gg 1. \quad (3.24)$$

Since the modulus is heavy we can integrate it out of our effective theory. The different component fields in T have mass differences of order m_r , so this is not an approximately supersymmetric threshold; also it is easy to see that F_T does not vanish ($\langle F_T \rangle \sim \langle r \rangle \langle F_\phi \rangle$). However, T couples to visible sector only through higher-dimension derivative interactions (recall that the modulus is the zero mode of the five-dimensional graviton polarized transverse to the branes), so this does not give a contribution to SUSY breaking in the visible sector at order $1/M_4^2$. We conclude that at order $1/M_4^2$, the effective theory below the modulus mass is precisely the ‘‘sequestered form’’ proposed in Ref. [2]: the visible sector is coupled only to a geometrically flat four-dimensional SUGRA background with broken SUSY ($F_\phi \neq 0$).

We now return to the question of the corrections to the Kähler potential in Eq. (3.13). The Kähler potential contains unknown $\mathcal{O}(1/M_4^4)$ corrections from loop corrections and higher-dimension operators, and one might worry that these are more important than the exponentially (in T) suppressed effects in the superpotential. This does not occur because the potential vanishes in the limit where the superpotential vanishes, so the Kähler corrections enter multiplicatively. This ensures the stability of the results above, in that the Kähler corrections to the modulus potential are of order $1/(rM_5)$ smaller than the leading potential we computed.

We now show that this scenario for radius stabilization can give rise to a sufficiently large radius without introducing small numbers or fine tuning. From Eq. (3.18), the stabilized value of the radius is

$$r \sim \frac{Ng_5^2}{l_5} \ln \left(\frac{l_5^3}{M_5^2 \langle F_\phi \rangle N (Ng_5^2)^3} \right). \quad (3.25)$$

Because the radius depends logarithmically on the fundamental parameters, we cannot obtain hierarchies of many orders of magnitude. In fact, because of the factor $1/l_5 \sim 10^{-3}$ multiplying the logarithm in Eq. (3.25), the bulk SYM gauge coupling g_5 must be large, and the fundamental theory must be close to strongly coupling.

The simplest assumption is that both gravity and the bulk SYM sector become strong at a single scale Λ_0 . NDA gives $\Lambda_0 \sim (l_5)^{1/3} M_5 \sim 10M_5$, and we will take this scale to be the fundamental scale of the theory (e.g. the scale of string–M-theory excited states). Using the NDA estimates for Λ_0 and g_5^2 , we obtain

$$r \sim \frac{1}{\Lambda_0} \ln \left(\frac{\Lambda_0}{N \langle F_\phi \rangle} \right). \quad (3.26)$$

Using $\langle F_\phi \rangle \sim 100$ TeV and $l_4 M_4^2 \sim l_5 M_5^3 r$, we obtain (for $N = 2$)

$$r \sim \frac{30}{\Lambda_0}, \quad \Lambda_0 \sim 2 \times 10^{18} \text{ GeV}. \quad (3.27)$$

This is sufficient to suppress FCNC effects from massive string states, but bulk gravitational loops give contact terms suppressed by [2]

$$\frac{1}{l_4 M_4^2 r^2} \sim \frac{1}{l_5 M_5^3 r^3} \sim \frac{1}{\Lambda_0^3 r^3} \sim 4 \times 10^{-5}. \quad (3.28)$$

This gives a contribution to soft scalar mass-squared terms of order $\langle F_\phi \rangle^2 / (\Lambda_0 r)^3 \sim (600 \text{ GeV})^2$, which is comparable to the contribution from anomaly mediation. It is interesting that this mechanism for radius stabilization can naturally stabilize the radius at a value where loop effects are important. This may give a solution to the problem of negative slepton masses, but we will not pursue this point here.

Another possibility is that the bulk SYM sector becomes strong at a scale Λ_{gauge} that is smaller than the scale Λ_{grav} . Here, Λ_{gauge} is a fundamental scale of new strong physics, while Λ_{grav} is not directly a physical scale, but corresponds to a weak gravitational coupling at the fundamental scale Λ_{gauge} . This occurs naturally if the gauge interactions propagate in fewer dimensions than gravity in the fundamental theory. For $\Lambda_{\text{gauge}} / \Lambda_{\text{grav}} \sim \frac{1}{10}$, we obtain $r \Lambda_{\text{grav}} \sim 160$. This is sufficient to suppress gravitational loop effects, and also suppresses flavor-changing contributions from string–M-theory states at the scale Λ_{gauge} . We have also checked that the contact terms from bulk gauge fields are negligible. These estimates are quite rough, but we conclude that it is very plausible that this mechanism can give a sufficiently large radius so that anomaly mediation dominates.

IV. CONCLUSIONS

We have studied a five-dimensional model with brane-localized visible and hidden sectors localized on ‘‘3-branes’’ and shown that when the compactification radius is properly stabilized, the transmission of supersymmetry breaking to

the visible sector proceeds by the mechanism of anomaly-mediation. Although the radius modulus participates strongly in the supersymmetry breaking, it does not contribute to soft visible sector masses at order $1/M_4^2$ because it does not directly couple to the visible brane. The stabilization mechanism for the radius modulus employed in this paper is very simple, involving gaugino condensates in the bulk and on a brane. The bulk gauge fields do not give additional contributions to visible soft masses due to the constraints of gauge invariance. The advantage of this mechanism is that it gives a non-perturbative superpotential for the modulus arising from field-theoretic mechanisms that are under theoretical control. It is also possible that such a superpotential could also arise from non-perturbative string–M-theory effects due to extended states.

This work is evidence that anomaly-mediated supersymmetry breaking gives a model-independent contribution to soft supersymmetry breaking in the visible sector at order $1/M_4^2$ in *any* model with SUSY breaking on a hidden-sector brane, and stabilized moduli. If there are no additional light bulk fields that give a larger contribution, anomaly-mediation dominates, giving a natural solution to the supersymmetric flavor problem as well as potentially testable predictions. These features can be upset by the presence of additional bulk fields with significant couplings to the visible sector.⁷ Knowledge of the true string theory vacuum, or experiment, is required to find out if such light non-minimal bulk fields are present.

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⁷For a phenomenologically interesting example, see Ref. [23].

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