Published by Institute of Physics Publishing for SISSA/ISAS
Received: *January 11, 2002*
Accepted: *May 27, 2002*

Realistic anomaly mediation with bulk gauge fields

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Abstract: We present a simple general framework for realistic models of supersymmetry breaking driven by anomaly mediation. We consider a 5-dimensional 'brane universe' where the visible and hidden sectors are localized on different branes, and the standard model gauge bosons propagate in the bulk. In this framework there can be charged scalar messengers that have contact interactions with the hidden sector, either localized in the hidden sector or in the bulk. These scalars obtain soft masses that feed into visible sector scalar masses at two loop order via bulk gauge interactions. This contribution is automatically flavor-blind, and can be naturally positive. If the messengers are in the bulk this contribution is automatically the same order of magnitude as the anomaly mediated contribution, independent of the brane spacing. If the messengers are localized to a brane the two effects are of the same order for relatively small brane spacings. The gaugino masses and A terms are determined completely by anomaly mediation. In order for anomaly mediation to dominate over radion mediation the radion must be is stabilized in a manner that preserves supersymmetry, with supergravity effects included. We show that this occurs in simple models. We also show that the μ problem can be solved by the vacuum expectation value of a singlet in this framework.

Keywords: Supersymmetry Breaking, Supersymmetric Standard Model.

Contents

1. Introduction

Anomaly mediated supersymmetry breaking (AMSB) [1] is an attractive framework for solving the supersymmetric flavor problem. AMSB requires only that non-gravitational interactions (including Planck-suppressed contact interactions between the visible and hidden sectors) are negligible compared to the interactions of low-energy supergravity. This can naturally occur in models with extra dimensions in which the visible and hidden sectors are localized on different branes [1], or in models where a conformal sector dynamically suppresses the contact terms [2]. In AMSB, the leading contribution to supersymmetry (SUSY) breaking parameters is a model-independent supergravity effect closely related to the conformal anomaly [1, 3]. The SUSY breaking masses are automatically flavor-blind, and are of order $m_{3/2}/16\pi^2$.

If the visible sector is the minimal supersymmetric standard model (MSSM), and there are no interactions between the visible and hidden sectors other than supergravity, then the slepton mass-squared terms are negative [1]. In this paper we discuss a very simple and general framework that can naturally explain why the slepton masses are positive and of the correct magnitude, while preserving the gaugino mass predictions of minimal AMSB. Our proposal is closely related to that of ref. [4], but is more general and robust, as we will explain. For other proposals for realistic models based on AMSB, see refs. [5]. For potential experimental signals of AMSB, see for example refs. [6].

The idea is very simple. Consider a 5D model in which the visible sector gauge bosons propagate in the bulk. In this scenario contact terms between the gauge fields and the hidden sector are not suppressed by higher-dimensional locality, but they are naturally suppressed if there are no singlets in the hidden sector. We also assume that the extra

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dimension is stabilized in a manner that does not break SUSY. (We will show that this occurs in simple models.) The leading contribution to the gaugino masses and A terms then comes from anomaly mediation. This framework for anomaly mediation was previously discussed in ref. [4], where it was shown that contact terms between the bulk gauge fields and the hidden brane can generate visible scalar masses of the same order as anomaly mediation.

In this paper we point out an additional contribution to the visible scalar masses in this framework. Because the standard model gauge bosons propagate in the bulk, there can be additional charged scalars, either in the bulk or on the hidden brane, that have contact terms with fields localized on the hidden brane. One particularly natural candidate is the scalar component of the bulk gauge supermultiplet. Up to volume factors, these scalars will get a mass of order $m_{3/2}$ from contact interactions with the SUSY breaking sector.[1] These in turn contribute to visible scalar masses at 2 loops. The contribution from these scalar messengers comes from gauge loops, and is therefore automatically flavor-blind. The sign of the visible mass-squared depends on the coefficient of a higher-dimension operator on the hidden brane, and can naturally be positive. For scalar messengers localized on the hidden brane, the messenger contribution to visible scalar masses is the same size as the AMSB contribution only for a special value of the compactification scale, as in ref. [4]. However, if the messengers propagate in the bulk, the messenger contribution to visible scalar masses is naturally the same as the AMSB contribution. This is independent of the compactification scale, and independent of whether the fundamental theory is strongly or weakly coupled. We therefore regard the framework of bulk scalar messengers as a very attractive and robust framework for realistic models of AMSB.

We also show that in this framework the μ problem can be naturally solved in the context of the 'next-to-minimal supersymmetric standard model' (NMSSM). This is a model with a singlet at the weak scale whose VEV gives rise to an effective μ term. The singlet mass-squared can be driven negative at the weak scale by radiative corrections, so all weak scale VEV's are explained by radiative symmetry breaking.

This paper is organized as follows. In section 2, we introduce models of scalar messengers and estimate the SUSY breaking parameters. In section 3, we address the μ problem. In section 4, we discuss radion stabilization. Section 5 contains our conclusions.

2. Scalar messengers

As discussed above, we assume that the standard-model gauge multiplets propagate in the bulk. A 5D gauge multiplet consists of the gauge field A_M ($M = 0, \ldots, 4$), a real adjoint scalar σ, and a fermion λ. We assume that the 5^{th} dimension is compactified on a S^1/Z_2 orbifold of radius r. The fixed points of the orbifold are 'branes' on which the hidden and visible sectors can be localized. The Z_2 parity assignments of the gauge field assigned so that A_5, σ, and half of the λ components are odd. These states will then get masses of order $1/r$, and the surviving degrees of freedom make up an $\mathcal{N} = 1$ gauge multiplet. (See e.g. refs. [7, 8] for details.)

[1] We assume that the bulk spacetime is nearly flat, i.e. the warp factor is not significant.

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In order to proceed, we must estimate higher-dimension contact terms that come from new physics at the fundamental scale. In the spirit of string unification, we will assume that there is a fundamental scale Λ such that all quantum corrections (whether from brane or bulk loops) are of order ϵ. The value $\epsilon \sim 1$ corresponds to strong coupling at the fundamental scale, while $\epsilon \sim 1/\ell_4 = 1/16\pi^2$ corresponds to the coupling strength of a 4D theory with all couplings of order unity.[2] The power counting of factors of ϵ and loop factors is explained in ref. [9].

The 4D gauge couplings are given by

$$g_4^2 \sim \frac{g_5^2}{r} \sim \frac{\epsilon \ell_5}{\Lambda r}, \tag{2.1}$$

where $\ell_5 = 24\pi^3$ is the 5D loop factor. From the fact that $g_4^2 \sim 1$ we can fix the radius:

$$\Lambda r \sim \epsilon \ell_5 . \tag{2.2}$$

Since Λ naturally sets the scale for heavy states in the bulk, so we require $\Lambda r \gtrsim 10$ for the absence of flavor-changing neutral currents (FCNC's). The acceptable range for ϵ is therefore

$$10^{-2} \lesssim \epsilon \lesssim 1 . \tag{2.3}$$

With this estimate, we find (somewhat amusingly)

$$\Lambda \sim M_4 , \tag{2.4}$$

where $M_4 \simeq 2 \times 10^{18}$ GeV is the 4D Planck scale.

We find the strong coupling scenario $\epsilon \sim 1$ particularly attractive, since the small value of the observed gauge couplings is explained by the same large radius that suppresses SUSY flavor violation. Also, strong coupling at the fundamental scale may be required to solve general difficulties with weakly-coupled string vacua [12].

2.1 Hidden scalar messengers

Since the visible gauge fields are in the bulk, we can have charged chiral fields Φ localized on the hidden brane. We now explore the possibility that these play the role of the hidden scalar messengers.

The hidden fields Φ will in general have couplings to the field X whose F component breaks SUSY:

$$\Delta \mathcal{L}_5 \sim \delta\left(x^5 - x^5_{\text{hid}}\right) \int d^4\theta \, \frac{\epsilon \ell_4}{\Lambda^3} X^\dagger X \Phi^\dagger \Phi . \tag{2.5}$$

These give a scalar mass-squared term

$$m^2_{\tilde{\Phi}} \sim \epsilon \ell_4 m^2_{3/2} . \tag{2.6}$$

[2]The counting used e.g. in ref. [4] assumes that all couplings are order 1 in units of the 5D Planck scale M_5, and that heavy bulk states have mass of order M_5. However the 5D gauge couplings must be larger than this estimate by a factor of 10 to account for the fact that $g_4 \sim 1$.

The sign of $m^2_{\tilde{\Phi}}$ is determined by the sign of the operator eq. (2.5), which arises from unknown short-distance physics. We also require the field Φ to have a supersymmetric mass M_Φ to give mass to the fermion components of Φ. This supersymmetric mass can arise from a brane-localized mass term on the hidden brane.

The effect of the supersymmetry breaking messenger mass on the visible sector scalar mass via 2-loop diagrams. This effect is analogous to D-type gauge mediation in four dimensions [11]. Since it arises from gauge interactions it is flavor blind. The size of the visible scalar mass depends on the relative magnitude of M_Φ and the compactification scale $1/r$. Let us first consider the case $M_\Phi \gg 1/r$. Then we can integrate out the field Φ in the 5D theory. In the effective lagrangian below the scale M_Φ the leading terms that depends on $m^2_{\tilde{\Phi}}$ have the form

$$\Delta\mathcal{L}_5 \sim \delta\left(x^5 - x^5_{\text{hid}}\right) \int d^4\theta \, \frac{\epsilon\ell_5}{\ell_4} \frac{1}{(M_\Phi/Z_\Phi)^2\Lambda} \bar{W}\bar{D}DW \,, \tag{2.7}$$

where D is the SUSY covariant derivative and $Z_\Phi = 1 - \theta^4 m^2_{\tilde{\Phi}}$. Expanding out the $m^2_{\tilde{\Phi}}$ term in Z_Φ, we obtain precisely the 'gaugino assisted' operators of ref. [4], but here they are suppressed by M_Φ rather than the fundamental scale. This operator gives rise to a 1-loop contribution to visible scalar masses of order

$$\Delta m^2_{\tilde{Q},\tilde{L}} \sim \frac{\epsilon\ell_4}{(M_\Phi r)^2} \frac{m^2_{3/2}}{\ell_4^2} \,. \tag{2.8}$$

(See ref. [4] for a detailed calculation.) This can be the same order as the anomaly mediated contribution for the right value of M_Φ.

We obtain a somewhat more robust result in the case $M_\Phi \ll 1/r$. In this case the 4D effective theory below the compactification scale includes Φ. The 2-loop running between the compactification scale and the mass M_Φ gives a log-enhanced contribution to the visible scalar masses

$$\Delta m^2_{\tilde{Q},\tilde{L}} \sim \epsilon\ell_4 \ln(M_\Phi r) \frac{m^2_{3/2}}{\ell_4^2} \,. \tag{2.9}$$

This naturally gives a flavor-blind positive contribution to the squark and slepton mass-squared terms that is somewhat larger than the AMSB contribution. (Remember $\epsilon\ell_5 \sim \Lambda r \gtrsim 10$.) To get a positive log-enhanced contribution to the visible scalar mass, we require $m^2_{\tilde{\Phi}} < 0$, so we require $|M_\Phi| > |m_{\tilde{\Phi}}|$ in order for the potential to be stable at $\Phi = 0$. In this case, the messenger contribution is the same size as the AMSB contribution for $\epsilon \sim 1/\ell_4$, which is what we obtain if we extrapolate the weak coupling of the standard model to the fundamental scale. The results are only logarithmically sensitive to the SUSY mass M_Φ.

2.2 Bulk scalar messengers

The mechanism discussed above is very simple and general, but it requires the introduction of a new supersymmetric mass scale M_Φ. Also, the new contribution to the visible scalar masses is larger than the AMSB contribution, especially for the attractive case where the fundamental theory is strongly coupled. Both of these potential difficulties can be elegantly removed if the scalar messengers propagate in the bulk. We then obtain a very robust framework for realistic AMSB that is the central result of this paper.

We can naturally obtain a SUSY mass for a bulk scalar of order the compactification scale $1/r$, in at least two ways. One way is to impose enough orbifold boundary conditions such that all bulk scalars are odd under at least one orbifold projection. Another way is to add a Planck-scale mass term localized on one of the branes; the lightest scalar KK mode is then 'repelled' from the brane with the mass term and gets a mass of order $1/r$ [9, 10]. With either mechanism, bulk scalar messengers do not require the introduction of an additional supersymmetric mass scale.

Let us consider these mechanisms in more detail for the the case of a bulk hypermultiplet. Using the formalism of ref. [8] this can be parameterized by two independent $\mathcal{N} = 1$ chiral superfields Φ and $\bar{\Phi}$ that depend on x^5.

Let us first consider an orbifold symmetry of the type $x_0^5 + x^5 \mapsto x_0^5 - x^5$ to give mass to the bulk hypermultiplets. The kinetic term is given by

$$\mathcal{L}_5 = \int d^4\theta \left(\Phi^\dagger \Phi + \bar{\Phi}^\dagger \bar{\Phi} \right) + \left(\int d^2\theta \, \bar{\Phi} \partial_5 \Phi + \text{h.c.} \right). \tag{2.10}$$

The second term implies that if Φ is odd under such a parity, then $\bar{\Phi}$ must be even, and *vice versa*. We can therefore compactify on a $S^1/(Z_2 \times Z_2)$ orbifolds, where the two Z_2 reflections leave invariant x^5_{vis} and x^5_{hid}, respectively. We choose the fields to have the following parity assignments: $\Phi \sim (-,+)$, $\bar{\Phi} \sim (+,-)$. This gives all components of the hypermultiplet a SUSY mass of order $1/r$.

Another way to give a SUSY mass to the even hypermultiplet when we have only one Z_2 symmetry is to add a brane-localized 'mass' term of the form

$$\Delta\mathcal{L} \sim \delta\left(x^5 - x^5_{\text{vis}}\right) \int d^2\theta \, \frac{\ell_5}{\ell_4} \Lambda \, \text{tr}(\Phi^2). \tag{2.11}$$

This is allowed by gauge symmetry e.g. if Φ is an adjoint. The effect of this term is to 'repel' the wavefunction of Φ and away from the visible brane by a factor of $\ell_4/(\ell_5 \Lambda r)$ and give the lowest-lying KK modes of Φ a mass of order $1/r$ [9].[3]

We assume that Φ has couplings to the hidden brane of the form

$$\Delta\mathcal{L} \sim \delta\left(x^5 - x^5_{\text{hid}}\right) \int d^4\theta \, \frac{\epsilon \ell_5}{\Lambda^3} \, X^\dagger X \Phi^\dagger \Phi. \tag{2.12}$$

(This coupling is allowed in the orbifold scenario, and is unsuppressed in the scenario with large visible 'mass' term.) Remarkably, the 2-loop D-type gauge mediation diagrams with one insertion of the operator above give rise to visible scalar masses of order

$$m^2_{\tilde{Q},\tilde{L}} \sim \frac{m^2_{3/2}}{\ell_4^2}, \tag{2.13}$$

the same as the AMSB contribution. Note that this is independent of ϵ, which tells us how strongly coupled the fundamental theory is. Since $\epsilon \ell_5 \sim \Lambda r$, this means the result is also independent of the compactification radius. The result is therefore very robust.

[3]This result can be understood from an elementary electrostatic argument [9].

There are also are couplings of the bulk messengers to the visible fields that violate flavor. We now show that the flavor changing neutral currents induced by these operators are consistent with experimental constraints.

In the orbifold model, the coupling of the Φ fields to the visible sector is suppressed because Φ is odd under the Z_2 reflection about the visible brane. The leading direct coupling to visible fields has the form

$$\Delta\mathcal{L} \sim \delta\left(x^5 - x^5_{\rm vis}\right) \int d^4\theta \, \frac{\epsilon\ell_5}{\Lambda^5} \, Q^\dagger Q \partial_5 \Phi^\dagger \partial_5 \Phi \,. \tag{2.14}$$

This gives rise to flavor-changing scalar masses from a 1-loop diagram connecting this coupling to the coupling eq. (2.12). The result is

$$\Delta m^2_{\tilde{Q}} \sim \frac{\ell_4}{(\epsilon\ell_5)^4} \, \frac{m^2_{3/2}}{\ell_4^2} \,. \tag{2.15}$$

This gives unobservably small FCNC's for strong coupling, and can give FCNC's near the experimental couplings for weak coupling. (Recall that $\epsilon\ell_5 \sim \Lambda r \gtrsim 10$.) There is a contribution of the same size coming from the $\bar{\Phi}$ field, which has orbifold suppressed couplings to the hidden brane, but unsuppressed couplings to the visible brane.

In the 'large brane mass' scenario, the 1-loop contribution to the flavor-changing mass of squarks is suppressed by the KK wavefunction factor. The resulting mass is of order

$$\Delta m^2_{\tilde{Q}} \sim \frac{\ell_4^3}{\epsilon^4 \ell_5^6} \, \frac{m^2_{3/2}}{\ell_4^2} \,, \tag{2.16}$$

which gives unobservably small FCNC's for for strong coupling, and can give FCNC's near the experimental couplings for weak coupling.

Another very natural candidate for the bulk messenger is the adjoint scalar of the bulk gauge multiplet. Recall that the bulk gauge multiplet contains 2 4D scalars A_5 and σ. The field A_5 cannot have non-derivative couplings by 5D gauge invariance, but non-derivative couplings of σ are allowed. To analyze the couplings of the bulk gauge multiplet to the branes, we use the formalism of ref. [8]. The bulk gauge multiplet is parameterized by $\mathcal{N} = 1$ superfields V, Σ that depend explicitly on x^5. Here V is a vector superfield and Σ is an adjoint chiral superfield with

$$V = \cdots + \bar{\theta}\sigma^\mu\theta A_\mu + \cdots \,, \qquad \Sigma = \frac{1}{2}(\sigma + iA_5) + \cdots \,. \tag{2.17}$$

Under holomorphic gauge transformations g these transform as

$$e^V \mapsto g^{-1\dagger} e^V g^{-1} \,, \qquad \Sigma \mapsto g\Sigma g^{-1} - g\partial_5 g^{-1} \,. \tag{2.18}$$

We can then form the combination

$$\tilde{\Sigma} = \Sigma + e^{-V}\Sigma^\dagger e^V - e^{-V}\partial_5 e^V = \sigma + \cdots \,, \tag{2.19}$$

which transforms as a gauge adjoint:

$$\tilde{\Sigma} \mapsto g\tilde{\Sigma}g^{-1} \,. \tag{2.20}$$

Note that $\tilde{\Sigma}$ is not holomorphic, and therefore cannot appear in superpotential terms.

We can therefore write couplings to the hidden sector of the form

$$\Delta\mathcal{L}_5 \sim \delta\left(x^5 - x^5_{\rm hid}\right) \int d^4\theta \, \frac{\epsilon \ell_5}{\Lambda^3} \, X^\dagger X \, {\rm tr}\left(\tilde{\Sigma}^\dagger e^V \tilde{\Sigma} e^{-V}\right) . \tag{2.21}$$

The estimate of the visible scalar masses is the same as for the bulk hypermultiplet case, so we again obtain scalar masses of the same order as the AMSB contribution. We must also worry about flavor-violating operators similar to eq. (2.14), and the estimates for FCNC's are again the same.

The discussion above assumes that the bulk gauge fields must have a SUSY mass of order $1/r$. As above, this can be acheived either by an orbifold projection, or by using a brane-localized 'mass' term for Σ. The estimates for FCNC's are the same as above. If we impose an orbifold projection, then A_5 (and hence Σ) must be odd since the covariant derivative $D_5 = \partial_5 + iA_5$ is odd. In this case we require that the hidden sector be on a brane that is located away from orbifold fixed points. This means that the position of the hidden brane is dynamical, and gives rise to a modulus that must be stabilized. We will not address this problem here, but it does not appear unnatural for such a brane to be stabilized between the fixed points of a S^1/Z_2 orbifold.

3. The μ problem

We now consider the μ problem in the present framework. In AMSB an explicit μ term gives $B \sim m_{3/2}$, which is much larger than anomaly-mediated masses of order $m_{3/2}/16\pi^2$. Therefore, in AMSB models the μ problem cannot be circumvented by simply adding a tree-level μ term. Here we point out that a μ term of the correct size can naturally be generated by the vacuum expectation value of a singlet S on the visible brane, with superpotential couplings

$$\Delta\mathcal{L}_5 = \delta\left(x^5 - x^5_{\rm vis}\right) \int d^2\theta \left[\lambda S H_u H_d + \frac{k}{3} S^3\right] . \tag{3.1}$$

Note that the AMSB contribution to m_S^2 comes from the superpotential couplings λ and k, and is therefore positive. The hidden messengers do not contribute to the S mass (at leading 2 loop order) because S is uncharged, so $m_S^2 > 0$ at the compactification scale. However, the 4D RG evolution to the weak scale can make $m_S^2 < 0$ at the weak scale. Also the A terms generated by AMSB favor a nonzero VEV for S.

We assume that the compactification scale is higher than the GUT scale, and that GUT threshold corrections are negligible. In this case the hidden messenger corrections consist of a universal scalar mass-squared Δm_0^2 for all squarks and sleptons (assuming a SO(10) GUT) and a correction Δm_H^2 for the electroweak Higgs scalars. In simple models, the ratio $\Delta m_H^2/\Delta m_0^2$ is a ratio of GUT Casimirs. However, in realistic GUT's the electroweak Higgs is generally a mixture of GUT representations, and we treat the ratio $\Delta m_H^2/\Delta m_0^2$ as a free parameter. With these simplifying assumptions, the parameters of this model are therefore the couplings λ, k, the top Yukawa coupling y_t, and the mass parameters Δm_0^2 and Δm_H^2, all renormalized at the GUT scale, and the SUGRA order parameter $\langle F_\phi \rangle \sim m_{3/2}$ that sets the size of the AMSB SUSY breaking. (We do not consider large $\tan\beta$ solutions, so we can neglect y_b and y_τ.) These parameters are constrained by the observed value of the W and Z masses, and the top quark mass.

In this framework, it is not difficult to find solutions which satisfy all experimental constraints with fine tuning of order 1%. (The largest fine-tuning is in α_3.) We have not performed a systematic analysis of the parameter space, but we make a few comments on the solutions we have found. The solutions we find have $\Delta m_H^2 > \Delta m_0^2$, which appears to be required in order to obtain $m_S^2 < 0$ at the weak scale (radiative symmetry breaking). This can be natural if the electroweak Higgs has an admixture of a large GUT representation. The presence of the S field means that this model has an additional neutralino that can mix with the other neutralinos and spoil the minimal AMSB prediction of wino LSP. However, it is not difficult to find realistic solutions in which the LSP is mostly wino, so this is still a potential signal for this class of models.

An illustrative example of a realistic point in parameter space has spectrum as follows (all masses are in GeV): $\chi_1^0 = 103$, $\chi_1^\pm = 108$, $h^0 = 128$, $\chi_2^0 = 194$, $\chi_3^0 = 205$, $\chi_2^\pm = 220$, $\tilde{\ell} \simeq 240$, $H^0 = 266$, $\chi_4^0 = 445$, $\tilde{t} = 570, 700$, $A^0 = 870$, $\tilde{q} \simeq 1010$, $\tilde{g} = 1030$, $\tan\beta = 9.1$.

4. Supergravity and SUSY radion stabilization

For the scenario above to work, it is crucial that the radion is stabilized without SUSY breaking by a radion F term. Otherwise the leading contribution to visible scalar masses comes from radion mediated SUSY breaking [13]; this also leads to viable models with a phenomenology similar to gaugino mediated SUSY breaking [14].

We therefore analyze radion stabilization without SUSY breaking. In the 4D effective theory, the radion r is parameterized by a chiral superfield with $\mathrm{Re}(T) \propto r$. It is important to discuss the radion effective theory in the presence of supergravity, since the Kähler term for the radion field is proportional to $T + T^\dagger$, which is trivial if supergravity effects are ignored.

We will consider theories in which the radion is stabilized in the 4D effective theory, i.e. $m_r \ll 1/r$. In this case, we can write an effective 4D field theory of the radion coupled to SUGRA:

$$\mathcal{L}_{\text{eff}} = \int d^4\theta\, \phi^\dagger \phi f + \left(\int d^2\theta\, \phi^3 W + \text{h.c.} \right), \tag{4.1}$$

where $\phi = 1 + \theta^2 F_\phi$ is the superconformal compensator and $f = -3M_4^2 + \cdots$. For the moment we consider a Kähler function f and superpotential W depending on a general set of 4D fields. We will consider specific models below.

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4.1 Supersymmetric vacua in supergravity

We first review the conditions to have a vacuum with SUSY breaking and vanishing cosmological constant. The supergravity potential is[4]

$$V = \frac{f_0^2}{f^2}\left[\left(\tilde{f}^{-1}\right)^a{}_b\left(W_a - \frac{3f_a W}{f}\right)\left(W^{\dagger b} - \frac{3f^b W^\dagger}{f}\right) + \frac{9|W|^2}{f}\right]. \tag{4.2}$$

where $f_0 = -3M_4^2$ and

$$\tilde{f}^a{}_b = f^a{}_b - \frac{f^a f_b}{f}. \tag{4.3}$$

The auxiliary fields are given by

$$F_a^\dagger = \left(\tilde{f}^{-1}\right)_a{}^b\left[W_b - \frac{3f_b W}{f}\right], \qquad F_\phi = -\frac{1}{f}\left[3W + f^a F_a^\dagger\right]. \tag{4.4}$$

A vacuum has unbroken SUSY and vanishing cosmological constant provided (assuming the matrix $\tilde{f}^a{}_b$ is nonsingular)

$$\frac{W}{f^3} = \text{stationary}, \qquad W = 0. \tag{4.5}$$

This is satisfied provided that W is stationary and $W = 0$ at the minimum. This is the well-known result that there is a one-to-one correspondence between supersymmetric vacua in global SUSY and supergravity [15].

We now apply this to a theory where the only light field is the radion T. The only possible subtlety is that the radion kinetic function

$$f = -3M_5^3(T + T^\dagger) \tag{4.6}$$

is unconventional. However, $\tilde{f}^T{}_T = M_5^3/(T+T^\dagger)$ is nonsingular for nonzero r, and therefore the arguments above go through. The conditions for radius stabilization with unbroken SUSY and vanishing cosmological constant are therefore

$$\frac{\partial W}{\partial T} = 0, \qquad W = 0. \tag{4.7}$$

4.2 Bulk super Yang-Mills

A simple way of generating radion superpotentials is to consider a theory with a super Yang-Mills (SYM) sector in the bulk. Below the compactification scale, this becomes a $\mathcal{N} = 1$ SYM theory (assuming an appropriate orbifold projection) with a 4D gauge coupling that depends on T. This theory becomes strong below the compactification scale and generates a T-dependent dynamical superpotential [16]

$$W_{\text{dyn}} = ae^{-bT}, \tag{4.8}$$

[4] Conventionally the supergravity potential is expressed in terms of the Kähler potential K, given by $f = -3M_4^2 e^{-K/3M_4^2}$.

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where (for gauge group SU(N))

$$a \sim \frac{1}{N^4 g_5^6}\,, \qquad b = \frac{16\pi^2}{3N g_5^2}\,. \tag{4.9}$$

With a single SYM sector in the bulk, we cannot satisfy the conditions eq. (4.7) for a SUSY vacuum.[5] However, if there are two SYM sectors in the bulk, the dynamical superpotential is

$$W_{\rm dyn} = a_1 e^{-b_1 T} + a_2 e^{-b_2 T}\,, \tag{4.10}$$

which has stationary points corresponding to SUSY vacua. The condition $W = 0$ can be satisfied by adding a constant to the superpotential (which itself may arise dynamically from gaugino condensation). The radion is stabilized at

$$\langle T \rangle = \frac{1}{b_1 - b_2} \ln\left(-\frac{a_1 b_1}{a_2 b_2}\right). \tag{4.11}$$

$\langle T \rangle$ is real provided that a_1 and a_2 are real with opposite sign.

5. Conclusions

We have presented a simple and attractive mechanism for obtaining realistic models of anomaly-mediated supersymmetry breaking. The basis of the mechanism is that in four dimensions, D-type gauge mediation gives flavor-blind, 2-loop contributions to the scalar masses that are naturally of the same size as the contribution from anomaly mediation. However in four dimensions it is generally problematic to suppress direct contact terms between the hidden and visible sectors which then give supersymmetry breaking contributions larger than both these effects. (See however ref. [2].) We therefore embed the theory in a higher dimensional space, with the visible and hidden sectors on different branes. Now contact terms between the visible scalars and the hidden sector are suppressed by higher-dimension locality, while contact terms between the bulk gauge multiplets and the hidden sector are suppressed if the hidden sector contains no singlets. In this case, the visible scalars can obtain 2-loop D-type contribution from 'messenger' scalars localized on the hidden brane or in the bulk. This mechanism preserves many of the attractive features of anomaly mediation, in particular the predictions for the gaugino masses and A-terms.

If the messengers are localized on the hidden brane, the D-type contributions are the same size as the anomaly-mediated contributions only for a special choice of compactification radius. This is similar (and in fact closely related) to the mechanism of ref. [4].

We obtain a more robust and attractive model if the scalar messengers are in the bulk. In this case, the contributions to the soft scalar masses from D-type gauge mediation and anomaly mediation are automatically the same size, independent of the size of the extra dimension, and independent of whether the theory is strongly or weakly coupled at the fundamental scale. It is possible for the components of the supersymmetric gauge multiplet in higher dimensions to play the role of the bulk messenger, so this mechanism does not require the introduction of additional multiplets.

[5] Ref [16] showed that this model with the addition of a constant superpotential term and brane-localized SUSY breaking stabilizes the radius with $\langle F_T \rangle \neq 0$.

We have also shown that this framework admits a solution to the μ problem based on the next-to-minimal supersymmetric standard model.

In order for the effects we are considering to dominate over radion mediation, we require that the brane spacing be stabilized in a supersymmetric way. We construct a very simple model in which this naturally occurs.

We believe that this framework is the simplest possibility for realistic models of anomaly mediated supersymmetry breaking.

Acknowledgments

Z.C. was supported in part by the U.S. Department of Energy under Contract DE-AC03-76SF00098, and in part by the NSF under grant PHY-00-98840. M.A.L. was supported by the NSF under grant PHY-98-02551. Z.C. would like to thank Y. Nomura for a very useful discussion.

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