Published by Institute of Physics Publishing for SISSA/ISAS
Received: *July 29, 2002*
Accepted: *August 21, 2002*

Purely four-dimensional viable anomaly mediation

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Abstract: Anomaly mediation of supersymmetry breaking solves the supersymmetric flavor problem thanks to its ultraviolet-insensitivity. However, it suffers from two problems: sleptons have negative masses-squared, and there are likely bulk moduli that spoil the framework. Here, we present the first fully ultraviolet-insensitive model of anomaly mediation with positive slepton masses-squared in a purely four-dimensional framework. Our model is based on the additional D-term contributions to the sparticle masses, and the conformal sequestering mechanism.

Keywords: Supersymmetry Breaking, Supersymmetry Phenomenology.

http://jhep.sissa.it/archive/papers/jhep082002034/jhep082002034.pdf

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1. Introduction

The absence of superpartners with masses degenerate to those of the Standard Model particles assures us that supersymmetry (SUSY) is broken if it is present in nature. However, arbitrary patterns of supersymmetry breaking masses give flavor-changing neutral current rates in gross violation of current experimental bounds. This is known as the SUSY flavor problem.

Anomaly-mediated supersymmetry breaking [1, 2] represents an attractive solution to this problem. In anomaly mediation, the breaking of scale invariance mediates supersymmetry breaking between the visible and the hidden sectors. (For an expanded review of this point, see [3]). As a result, the supersymmetry breaking terms can be written entirely in terms of beta functions and anomalous dimensions *at any given energy*. This means that the SUSY breaking terms are independent of any flavor-dependent physics that might occur at high scales. This is the anomaly-mediated solution to the SUSY flavor problem. While the anomaly-mediated contribution to supersymmetry breaking is always present, its effect can be sub-dominant. For the anomaly mediation to solve the SUSY flavor problem, other contributions to supersymmetry breaking must be somehow suppressed. In its original incarnation, this suppression was accomplished through a spatial separation of the visible and hidden sectors in an extra-dimension, which was thought to have the effect of "sequestering" the hidden sector [1].

The theory of anomaly mediation as described above suffers from two difficulties. First, as was shown in detail in [4], the geometric separation of the visible and hidden sectors is insufficient to sequester the hidden sector. In fact in most cases, the exchange of fields in the bulk supergravity multiplet is sufficient to generate a contribution to universal scalar masses at the tree level, completely swamping the anomaly-mediated piece. Perhaps the simplest counter-example to geometric sequestering appeared in the model of radion-mediated supersymmetry breaking [5]. The second problem facing anomaly mediation is a direct result of its ultraviolet (UV) insensitivity. This property renders anomaly mediation an extremely predictive framework, considered to be a desirable feature. However, one prediction is a negative mass squared for the sleptons, rendering the theory phenomenologically inviable.

Neither of the above problems are insurmountable. In [6, 7], it was shown that the sequestering could take place in a completely four dimensional framework where the problem of bulk moduli is absent, using conformal field theories. In this case, the sequestering remains effective as long as there are no moduli coupled to both the hidden and the MSSM sectors. This four dimensional solution is reviewed in section 2. The problem of tachyonic sleptons has also been overcome [1] [8]–[11]. However, no explicit model exists that both effectively sequesters the supersymmetry breaking sector and provides a phenomenologically viable spectrum. In this paper we provide the first such model, a viable model of anomaly mediation in four dimensions.

To provide the masses for the sleptons, we select the mechanism of [10, 11], which gives masses through the addition of D-terms. We find this mechanism particularly attractive, as it gives masses without sacrificing the property of UV insensitivity. We review this mechanism in section 3, before moving on and presenting our model in section 4.

2. 4D anomaly mediation

The concerns about geometric sequestering raised in [4] can be addressed by utilizing the framework of Luty and Sundrum [6, 7], in which the sequestering of anomaly mediation is realized in a completely four-dimensional way. This framework is inspired by the AdS/CFT correspondence [12] which conjectures that a 5D theory of gravity in anti-deSitter space is dual to a four dimensional conformal field theory. One can hope that this duality can be stretched further to a duality between the five dimensional brane world scenario and a four dimensional conformal field theory. With this in mind we can expect that the sequestering attempted in [1] with an extra dimension might be realized in a 4D model in which the hidden sector is conformal. Luty and Sundrum showed this is indeed the case by looking at a 4D theory with the hidden sector being SUSY QCD with $\frac{3}{2}N_c \leq N_f \leq 3N_c$ (i.e. the theory is in the Seiberg conformal window [13]).

The main goal of sequestering is to suppress flavor violating operators of the form

$$\int d^4\theta \frac{c^i_j}{M_*^2} T^\dagger_J T^J Q^\dagger_i Q^j \tag{2.1}$$

that might lead to large flavor-changing effects. Here T^J are fields in the (conformal) hidden sector, Q^i are fields in the visible sector, and M_* can be the reduced Planck scale. It is convenient to treat such a term as a correction to the hidden sector UV Kähler potential

$$\mathcal{L}_{hid} \supset \int d^4\theta \mathcal{Z}_0 T_J^\dagger T^J \,, \qquad \mathcal{Z}_0 = 1 + \frac{c_j^i}{M_*^2} Q_i^\dagger Q^j \,. \tag{2.2}$$

Below a certain energy, the gauge coupling, g, of the hidden sector nears its fixed point value, g_*, allowing us to keep only the leading order in $(g^2 - g_*^2)$ in the renormalization group equations for the coupling and the wavefunction renormalization. Defining $\gamma \equiv \partial \log \mathcal{Z} / \partial \log \mu$, the expansion yields:

$$\beta = \beta_*'(g^2 - g_*^2) \,, \qquad \gamma = \gamma_* + \gamma_*'(g^2 - g_*^2) \,, \tag{2.3}$$

where β_*' and γ_*' are positive critical exponents, due to the fact that they describe dynamics near an infrared fixed point. For simplicity we will assume eq. (2.3) is valid already at M_*; this assumption will not change the final result of sequestering. The conformal symmetry will eventually be broken by one of the fields in the hidden sector getting a vacuum expectation value (VEV), $\langle X \rangle$, away from the origin of the moduli space.

Now, we can proceed to demonstrate the sequestering. If at the scale M_*, we were already at the fixed point, $g = g_*$, integrating eq. (2.3) would give

$$\mathcal{Z}(\mu) = \left(\frac{\mu}{M_*} \right)^{\gamma_*} . \tag{2.4}$$

When considering small perturbations about the fixed point, $\Delta g^2 = g^2 - g_*^2$, it is convenient to factor out the above fixed point running by defining

$$\Delta \ln \mathcal{Z} = \ln \mathcal{Z} - \gamma_* \ln \left(\frac{\mu}{M_*} \right) . \tag{2.5}$$

After integrating eq. (2.3) down to the scale of conformal symmetry breaking, $\langle X \rangle$, we find that the dependence of $\Delta \ln \mathcal{Z}(\mu)$ on $\mathcal{Z}_0$ is highly suppressed. This is simply a restatement of the fact that until the symmetry is broken, the theory is nearly conformal, so that deformations introduced at the high scale quickly become irrelevant as the theory is driven toward its fixed point.

In particular, if we make the choice that the value of the holomorphic gauge coupling[1] at the Planck scale is equal to its value at the fixed point, which can be done without a loss of generality, we can write:

$$\Delta \ln \mathcal{Z}(\langle X \rangle) = \left(\frac{\langle X \rangle}{M_*} \right)^{\beta_*'} \Delta \ln \mathcal{Z}_0 . \tag{2.6}$$

The flavor violating operators in eq. (2.1) are therefore power suppressed and the hidden sector is sequestered. For details of this derivation, see References [6, 7].

[1] This is the gauge coupling whose running is exhausted at one loop [14], and is related to the canonical gauge coupling by $1/g_{can}^2 = Re(1/g_{hol}^2) - F \ln Z / 8\pi^2 - N \ln g_{can}^2 / 8\pi^2$ in the NSVZ scheme.

In [7] a model of a conformally sequestered hidden sector is achieved completely naturally, generating all hierarchies dynamically. In the sections that follow, we will assume that the hidden sector is sequestered by this mechanism, leaving anomaly mediation as the leading contribution to visible sector soft masses.

3. Viable anomaly mediation

For anomaly mediation to be viable, the tachyonic sleptons must be eliminated. This may be accomplished by adding D-terms for $B-L$ and hypercharge. In Reference [11] the generation of a D-term for $\mathrm{U}(1)_{B-L}$ is accomplished by appealing to an extra-dimensional framework. We have already mentioned that the extra-dimensional framework is problematic as a realization of anomaly mediation, so we will have to modify this mechanism, which we do in section 4.

The extra-dimensional set-up of [11] contains three branes: a visible sector, a sequestered sector, and a sector responsible for the breaking of $\mathrm{U}(1)_{B-L}$. The $\mathrm{U}(1)_{B-L}$ gauge field is allowed to propagate in the bulk, but is broken at a high scale unlike in the gaugino-mediation models [15]. We now review the process by which the $\mathrm{U}(1)_{B-L}$ breaking sector provides a D-term. The $\mathrm{U}(1)_{B-L}$ breaking sector consists of the superpotential

$$\mathcal{W} = \lambda X(\phi\bar{\phi} - \Lambda^2)\,. \tag{3.1}$$

Here, the ϕ and $\bar{\phi}$ fields have $\mathrm{U}(1)_{B-L}$ charge $+1$ and -1, respectively, while X is $\mathrm{U}(1)_{B-L}$ neutral. In the supersymmetric limit, $\langle\phi\rangle = \langle\bar{\phi}\rangle = \Lambda$. As long as ϕ and $\bar{\phi}$ have different soft supersymmetry breaking masses, perhaps from additional Yukawa couplings not shown in eq. (3.1), their VEVs will shift by different amounts after the effects of supersymmetry breaking are included [16]. The D-term is found to be proportional to the difference in the VEVs, so

$$D_{B-L} \sim \langle\phi\rangle^2 - \langle\bar{\phi}\rangle^2 \sim \tilde{m}^2 - \tilde{\bar{m}}^2\,. \tag{3.2}$$

Generically, the D-term will be of the same order as the soft masses-squared, which is the necessary condition for a viable spectrum. The D-term contributes to the scalar masses-squared for all fields charged under the relevant symmetry, in this case, $B-L$.

The unique feature of this solution to the problem of tachyonic sleptons is its insensitivity to physics in the ultraviolet. More precisely, the new expressions for the soft masses:

$$\tilde{m'}_i^2 \equiv \tilde{m}_i^2 - q_i D\,, \tag{3.3}$$

remain on renormalization group trajectories. This can be shown by either explicitly examining the form of the renormalization group equations [10], or by a spurion argument [11]. The spurion argument also shows clearly that soft parameters remain on RGE trajectories even as heavy particle thresholds are crossed.

As mentioned above, D-terms for both $\mathrm{U}(1)_{B-L}$ and $\mathrm{U}(1)_Y$ are needed to generate a viable spectrum [10]. However, once a D-term for $\mathrm{U}(1)_{B-L}$ is generated, a D-term for $\mathrm{U}(1)_Y$ is naturally generated by including a term for kinetic mixing in the lagrangian, $\mathcal{L} \ni \int d^2\theta\, \mathcal{W}_{B-L}\mathcal{W}_Y$. Therefore, it is sufficient to consider a mechanism to generate D_{B-L}.

The "anomaly mediation plus D-terms" scenario above was originally envisaged in an extra-dimensional framework. If we transport this scenario to a four-dimensional one, we require that the scale of the $\mathrm{U}(1)_{B-L}$ breaking be much less than the Planck scale. Otherwise, it is possible to generate phenomenologically dangerous contributions to the soft masses of the MSSM field. To be concrete, consider the lepton doublet, L. Operators in the Kähler potential of the form

$$\frac{(\phi L)^\dagger e^{qV}(\phi L)}{M_*^2}, \tag{3.4}$$

will generate soft masses for the sleptons once the VEV for the D-term is substituted for V. A key point is that q generically will not be equal to the charge of the L field. These contributions will take the masses off of the anomaly-mediated trajectory, and spoil the UV insensitivity. These operators could arise from integrating out heavy particles charged under the $\mathrm{U}(1)_{B-L}$ symmetry, whose masses are at the Planck scale.[2] Therefore, a requirement for a viable four dimensional model is that we dynamically generate the scale of the $\mathrm{U}(1)_{B-L}$ breaking. We present such a model in the next section.

4. The model

Our goal, then, is to produce a model that naturally generates a D-term of the correct scale without introducing operators that spoil the UV-insensitivity. Looking at sections 2 and 3, it is clear what we need to do. Again, we will have three sectors: a visible sector containing the MSSM, a nearly conformal hidden sector, and a $\mathrm{U}(1)_{B-L}$ breaking sector. The nearly conformal sector can be neatly lifted from Reference [7]. Our task is to generate a superpotential similar to that in eq. (3.1) dynamically. This will generate the D-terms, allowing for a phenomenologically viable spectrum.

The sector that we use to break the U(1) gauge symmetry and produce the D-term is given in table 1. It is a theory with a "strong" Sp(2) gauge group and 3 flavors, corresponding to 6 fundamentals, which we denote by Q. The theory exhibits an SU(6) flavor symmetry. We gauge the $\mathrm{SU}(4) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ subgroup of the flavor symmetry, and this U(1) is the $\mathrm{U}(1)_{B-L}$ that develops a D-term.[3] To cancel the anomalies for the $\mathrm{SU}(4) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ symmetry, we introduce four copies of anti-flavors $\bar{Q}$ transforming as a $\bar{4}$ under a flavor $\mathrm{SU}(4)_g$ global symmetry but not charged under the strong Sp(2) gauge group.

The superpotential of the model is

$$W = \frac{1}{M_*}(Q^i Q^j)\bar{Q}_i \bar{Q}_j + \frac{1}{M_*^3}Q^6 + \frac{1}{M_*^3}\bar{Q}^6\,, \tag{4.1}$$

[2] Alternatively, these could simply be additional fields charged under the U(1) symmetry whose Yukawa couplings to ϕ and $\bar{\phi}$ give those fields different soft masses.

[3] A similar attempt using Sp(1) with 2 flavors runs into the difficulty that the mesons have identical charges under the $\mathrm{SU}(3) \times \mathrm{U}(1) \subset \mathrm{SU}(4)$ flavor symmetry, so it is not entirely clear that a D-term is generated.

	Sp(2)	SU(6) $\supset$ SU(4) $\times$ SU(2) $\times$ U(1)	$\mathrm{SU}(4)_g$
Q	4	$6 = (4,1)_1 + (1,2)_{-2}$	1
$\bar{Q}$	1	$\bar{6} = (\bar{4},1)_{-1} + (1,2)_2$	4
M	1	$15 = (6,1)_2 + (4,2)_{-1} + (1,1)_{-4}$	1

Table 1: Particle content of the U(1) breaking sector and associated mesons.

which is the most general superpotential up to order $1/M_*^3$ consistent with the symmetries described above. The indices run from 1 to 6 over all three flavors. However, the terms do not have to be fully SU(6) symmetric, but symmetric only under SU(4) $\times$ SU(2) $\times$ U(1), i.e. different coefficients for independent invariants are allowed in the last two terms.

The dynamics of the Sp(2) gauge group leads to the quantum modified constraint [17]

$$\mathrm{Pf}\ M = \Lambda^6 \,, \tag{4.2}$$

where the meson is made of quarks $M^{ij} = Q^i Q^j$. Because the quarks are in $(4,1)_1 + (1,2)_{-2}$ representations, the mesons transform as $(6,1)_2 + (4,2)_{-1} + (1,1)_{-4}$. The quantum modified constraint can be satisfied consistently with vanishing SU(4) $\times$ SU(2) $\times$ U(1) D-terms with

$$M_{(6,1)_2} = \left(\sqrt{2}\, V^2, 0, 0, 0, 0, 0\right) , \qquad M_{(1,1)_{-4}} = V^2 \,, \tag{4.3}$$

where $V^6 = \Lambda^6$. The first meson VEV breaks SU(4) to Sp(2) (or equivalently SO(6) to SO(5)), and both of them break U(1). Most of the mesons in the $(6,1)_2 + (1,1)_{-4}$ representations are eaten by the broken gauge multiplets, and one of them is eliminated due to the quantum modified constraint. The $(4,2)_{-1}$ mesons do not acquire mass because of the "accidental" SU(6) symmetry in the Sp(2) dynamics and are pseudo-Nambu–Goldstone fields. The second term in eq. (4.1) then breaks the "accidental" SU(6) and gives the $(4,2)_{-1}$ mesons masses of order Λ^4/M_*^3. On the other hand, the first term in eq. (4.1) gives mass of order Λ^2/M_* to all of the anti-flavors. As a result, all fields that we have introduced acquire supersymmetric masses. There are no new light fields in our model.

We claim that, in the presence of soft supersymmetry breaking induced by anomaly mediation, this model develops a D-term for the U(1).

Unfortunately, this model is incalculable because the mesons become composite and condense at the same scale, and we do not know if the description in terms of meson degrees of freedom is appropriate to work out the soft supersymmetry breaking effects. Of course, once the dynamical degrees of freedom are identified, anomaly mediation gives a unique prediction for their soft parameters. However, we do not know if mesons can be regarded as dynamical degrees of freedom or composite order parameters at this stage.

We make this model calculable, in the same spirit as in [18], by deforming the theory while remaining in the same universality class. The model becomes calculable if the two energy scales are separated: the scale where mesons become composite Λ, and the scale where U(1) is broken V. We identify two deformations which achieve $\Lambda \gg V$ or $\Lambda \ll V$. We will find that a D-term of the correct order of magnitude is successfully generated in either case. We view this as compelling evidence for the generation of a D-term when $\Lambda \sim V$, which is the actual situation in the model we have presented.

4.1 $\Lambda \gg V$

The first limit, $\Lambda \gg V$, is achieved by introducing an additional flavor, Q_7, Q_8. No corresponding anti-flavor is necessary as the additional flavor is not charged under any gauge group except the strong Sp(2). The extra flavor is massive. When its mass $m \gg \Lambda$, we can integrate out the extra flavor first, and the theory goes back to the original model. On the other hand, by taking $m \ll \Lambda$, the Sp(2) dynamics changes at the dynamical scale. The theory has 4 flavors and hence confines with a dynamical superpotential [17]. Together with the mass term,

$$\Delta W = -m\mathcal{M}^{78} + \frac{\mathrm{Pf}\,\mathcal{M}}{\Lambda^5}\,. \tag{4.4}$$

The meson $\mathcal{M}$ includes the 4th flavor, while we keep the notation M for the first three flavors. Below the scale Λ, the confined mesons are dynamical degrees of freedom, as suggested by the non-trivial anomaly matching conditions.

At the scale $V = (m\Lambda^5)^{1/6}$, we solve for $\mathcal{M}^{78}$ and find

$$\mathrm{Pf}\,M = m\Lambda^5 = V^6\,, \tag{4.5}$$

which sets $\Lambda \gg V$ because $m \ll \Lambda$. The mesons acquire expectation values in the same fashion as in eq. (4.3), breaking U(1) and SU(4) $\to$ Sp(2).

The superpotential of eq. (4.1) becomes Yukawa couplings among mesons and anti-flavor quarks

$$W = \frac{1}{M_*} M^{ij} \bar{Q}_i \bar{Q}_j + \frac{1}{M_*^3} M^3\,. \tag{4.6}$$

Once mesons acquire VEVs, the first term gives the anti-flavors masses, and the second term gives masses to the $(4,2)_{-1}$ mesons. The size of Yukawa couplings can be seen by scaling the meson fields to canonical normalization $M \approx \Lambda \hat{M}$ up to unknown $O(1)$ constants.

The main point of this deformation is that physics between the compositeness scale, Λ, and the U(1)-breaking scale, V, is given in terms of composite mesons and their soft supersymmetry breaking can be obtained by the standard formulae of anomaly mediation. Yukawa couplings from the non-renormalizable superpotential are suppressed by Λ/M_* and we ignore them. Then the only contribution to their soft masses come from SU(4) $\times$ SU(2) $\times$ U(1) gauge interactions. The relevant soft masses for the mesons are determined by the expression $\tilde{m}^2 = (-\dot{\gamma}/4)|m_{3/2}|^2$, where $\gamma \equiv (\partial \log Z/\partial \log \mu)$. Calculating in the gauge theory, we find:

$$\tilde{m}^2_{(6,1)_2} = \frac{|m_{3/2}|^2}{(16\pi^2)^2}\left(20g_4^4 - 384g_1^4\right), \tag{4.7}$$

$$\tilde{m}^2_{(1,1)_{-4}} = \frac{|m_{3/2}|^2}{(16\pi^2)^2}\left(-1536g_1^4\right). \tag{4.8}$$

In the presence of soft terms, we have to re-minimize the potential obtained from the superpotential eq. (4.4). The vacuum configuration for the new potential shifts slightly from that in eq. (4.3). And the D-term, which is given in terms of the shifts in the meson VEVs, can be given by:

$$D = 8g_1(\delta M_{(6,1)_2} - \delta M_{(1,1)_{-4}}) = \frac{\tilde{m}^2_{(1,1)_{-4}} - \tilde{m}^2_{(6,1)_2}}{6g_1}\,. \tag{4.9}$$

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As an aside, we note that the potentially large A-terms that result from inserting powers of the chiral compensator in eq. (4.4), do not contribute to the D-term. We have checked that this is the case using the methods of [16]. Now, substituting in the expressions of eqs. (4.7) and (4.8), the final value for the D-term in this case are given by:

$$D = -\frac{|m_{3/2}|^2}{(16\pi^2)^2}\frac{192g_1^4 + \frac{10}{3}g_4^4}{g_1}\,, \tag{4.10}$$

so a D-term of the correct order of magnitude is successfully generated in this limit.

4.2 $\Lambda \ll V$

The second limit $\Lambda \ll V$ is achieved by introducing an additional flavor, Q_7, Q_8, and extending the strong gauge group to Sp(3). Because we are extending the gauge group to Sp(3), now six copies of the anti-flavors are necessary. We assign the extra fourth flavor the superpotential:

$$\Delta W = \lambda Y(Q^7 Q^8 - v^2) + \lambda' Y_{i\alpha} Q^i Q^\alpha\,. \tag{4.11}$$

Here, $i = 1, \ldots, 6$ and $\alpha = 7, 8$. We also extend the original superpotential by letting the indices in eq. (4.1) run over $1, \ldots, 8$. When $v \gg \Lambda$, the extra flavor condenses and breaks the gauge group $\mathrm{Sp}(3) \to \mathrm{Sp}(2)$, while it is eaten by the gauge multiplet. The extra color degrees of freedom of the remaining 3 flavors are made massive by the second term in eq. (4.11). The first term in eq. (4.1) makes the extra anti-flavors massive. Therefore the theory reduces to the original model.

When $v \ll \Lambda$, on the other hand, it yields a calculable model of D-term generation with the scale of U(1) breaking $V \gg \Lambda$. The soft masses can be calculated from anomaly mediation, the corresponding shift in the VEVs can be computed, and the D-term can be identified. The quick way to see that the model is still in the same universality class is by using the dynamical superpotential. Together with the quantum modified constraint,

$$\Delta W = \lambda Y(\mathcal{M}^{78} - v^2) + \lambda' Y_{i\alpha}\mathcal{M}^{i\alpha} + X(\mathrm{Pf}\ \mathcal{M} - \Lambda^8)\,. \tag{4.12}$$

The meson $\mathcal{M}$ includes the 4th flavor, while we keep the notation M for the first three flavors. We can immediately conclude that $\mathcal{M}^{78} = v^2$, $\mathcal{M}^{i\alpha} = 0$. Therefore, $\mathrm{Pf}\ \mathcal{M} = (\mathrm{Pf}\ M)\mathcal{M}^{78} = (\mathrm{Pf}\ M)v^2$. The last term in the superpotential of eq. (4.12) then determines $(\mathrm{Pf}\ M) = \Lambda^8/v^2$. Therefore the meson VEV breaks $\mathrm{SU}(4)\times\mathrm{SU}(2)\times\mathrm{U}(1) \to \mathrm{Sp}(2)\times\mathrm{SU}(2)$ just as before. What is important is the scale of the VEVs: when $v \ll \Lambda$, $\mathrm{Pf}M \gg \Lambda^6$, and hence the meson VEVs corresponding to the U(1) breaking scale, V, can be regarded as classical expectation values of quark fields. Therefore we can calculate the effects of soft parameters using the anomaly mediation formula for quark fields instead of composite fields.

At the scale $(\Lambda^4/v)^{1/3}$, we look at classical flat direction

$$Q = \left(\begin{array}{cccccc|cc} V & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & V & 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & V & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & V & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & V' & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & V' & 0 & 0 \end{array}\right). \tag{4.13}$$

In this notation, the Sp(3) gauge group acts from the left, while the $SU(4) \times SU(2) \times U(1)$ gauge groups act from the right on the quark multiplets. Because $V, V' \sim (\Lambda^4/v)^{1/3} \gg \Lambda$, it is a conventional Higgs mechanism that breaks $Sp(3) \times SU(4) \times SU(2) \times U(1) \to Sp(2) \times SU(2)$. Out of $6 \times 8 = 48$ components of the quark multiplets, $(21+15+3+1)-(10+3) = 27$ components are eaten. The 12 components in the fourth flavor become massive together with $Y_{i\alpha}$ due to the second term in eq. (4.11). The $(4,2)$ components acquire mass from the second term in eq. (4.1). This leaves only one light degree of freedom, namely the direction of V itself. We, however, leave the possibility that $V \neq V'$ in the above expression. This is justified when the U(1) coupling is small. The D-flatness is imposed because the D-non-flat direction acquires a mass of order $g_1 V$. When $g_1 \ll 1$ we can consider a low-energy effective theory where all particles of mass $\sim V$ are integrated out but those of mass $g_1 V$ are not. We can minimize the potential with respect to V and V' later.

Along this flat direction, we need to consider the instanton effect of partially broken gauge groups [19]. Because $\pi_3((Sp(3) \times SU(4) \times SU(2) \times U(1))/(Sp(2) \times SU(2))) = \mathbb{Z}$, there are instanton effects that are present in the $Sp(3) \times SU(4) \times SU(2) \times U(1)$ which are not part of the low-energy $Sp(2) \times SU(2)$ theory. Such effects have to be included in writing down the low-energy theory. In our case, it is

$$\mathcal{M}^{78} = \frac{\Lambda^8}{V^4 V'^2}. \tag{4.14}$$

This makes the first term in the superpotential eq. (4.11)

$$W = \lambda Y \left(\frac{\Lambda^8}{V^4 V'^2} - v^2 \right) \tag{4.15}$$

while the second term had already been used to integrate out the fourth flavor.

Therefore, the low-energy effective theory we need to solve is given by the directions V and V' of the elementary quarks together with the superpotential eq. (4.15) and soft supersymmetry breaking effects calculated in the quark language.

Calculating in this framework, we find that the D-term can be given in terms of the quark soft masses as

$$D = \frac{\tilde{m}^2_{(1,2,-2)} - \tilde{m}^2_{(4,1,1)}}{3 g_1} \tag{4.16}$$

The quark soft masses themselves arise from anomaly mediation and are given by:

$$\tilde{m}^2_{(4,1,1)} = \frac{|m_{3/2}|^2}{(16\pi^2)^2} \left(\frac{45}{4} g_4^2 - 144 g_1^4 \right) + m^2_{Sp(3)}, \tag{4.17}$$

$$\tilde{m}^2_{(1,2,-2)} = \frac{|m_{3/2}|^2}{(16\pi^2)^2}(-576g_1^4) + m^2_{\mathrm{Sp}(3)}\,. \qquad (4.18)$$

Where we have denoted the universal contribution coming from the Sp(3) gauge interactions as $m^2_{\mathrm{Sp}(3)}$. Note that the SU(2) theory is conformal at one-loop, so its coupling does not contribute to $\tilde{m}^2_{(1,2,-2)}$. The final expression for the D-term in this case is then

$$D = -\frac{|m_{3/2}|^2}{g_1(16\pi^2)^2}\left(\frac{15}{4}g_4^4 + 144g_1^4\right). \qquad (4.19)$$

We view the fact that a D-term is generated in both deformations (with the same sign, no less) as compelling evidence for existence of a D-term in the case $\Lambda \sim V$, which is the situation in our model for D-term generation.

4.3 UV insensitivity

Let us return to the questions of UV-insensitivity. There might be a concern that additional contributions to the soft masses of the MSSM fields might arise from integrating out other heavy fields charged under the U(1). However, because of the charges of the Q's under Sp(2) and $\mathrm{SU}(4) \times \mathrm{SU}(2) \times \mathrm{U}(1) \subset \mathrm{SU}(6)$, it is impossible to construct a non Planck-suppressed invariant with these fields that could couple to the MSSM. As long as the dynamically generated scale is much less than the Planck scale, these contributions are small. One might imagine generating non-Planck-suppressed operators if there were further fields charged under the Sp(2) gauge symmetry. However, this would upset the Sp(2) dynamics that breaks U(1). So, once we have specified our mechanism for the breaking of U(1), there is no room for additional contributions to the soft-masses that will displace them from their renormalization group trajectories.

5. Phenomenology

The spectrum of the superpartners should be nearly identical to that presented in references [10, 11], as we have merely presented a field theoretic realization of that scenario. Electroweak symmetry breaking has not been fully investigated in this scenario, and it would be interesting to explore the possibility of incorporating the NMSSM into this framework. We note that the sequestering mechanism of [7] is not perfect, and it is possible that interesting flavor changing effects could be observed soon in the context of this framework.

There are no new light fields in our model. All fields receive supersymmetric masses. That said, we prefer that dynamically generated scale not be too low. The mesons transforming as $(4,2)_{-1}$ only get a mass from the last term of eq. (4.1), so their masses are expected to be $m_{(4,2)_{-1}} \sim \Lambda(\Lambda/M_*)^3$. In the model as written, these fields are absolutely stable, so there could be a concern that they might over-close the universe. However, as long as $\Lambda > 10^{16}$ GeV, they place no constraint on the model above and beyond the usual constraint coming from not producing an over-abundance of gravitinos [20]. This value for Λ is sufficiently small that the Planck-suppressed contributions to the Kähler

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potential of the form of eq. (3.4) are negligible. Of course, lower values for Λ would be possible for smaller values of the re-heating temperature. If rather than being absolutely stable, these particles decayed through sufficiently Planck suppressed operators, it is possible that their lifetime could be roughly the age of the universe. In this case, they could provide a mechanism to explain the existence of ultra-high energy cosmic rays.

It is interesting to note that this mechanism accommodates the see-saw mechanism [21] for neutrino masses. We can add the couplings $W \supset (N^2 M_{(1,1)_{-4}})/M_* + hLNH_u$ to the superpotential (where we normalize our U(1) so that N has charge 2). Then, as the mesons acquire VEVs, the first term provides a mass of the right scale for atmospheric neutrino masses, assuming the Dirac masses are roughly the top mass, as suggested by the SO(10) mass relation. Note that there are corrections to the RGE trajectories at higher orders in this case, but the leading order result remains unaffected [11]. In particular, one expects that the operator

$$\mathcal{K} \ni \frac{|h|^2}{16\pi^2} \frac{(L^\dagger e^{q_L V} L)(M^\dagger e^{q_M V} M)}{M_*^2 M_R^2}, \tag{5.1}$$

will be generated in the Kähler potential by integrating out the N field. Here, $M_R \sim M/M_* = \Lambda^2/M_*$ is the right handed neutrino mass. Inserting the vacuum expectation values for the mesons and the D-term, this represents a perturbation away from the anomaly mediated trajectory at the percent level, and it is not entirely automatic that lepton flavor-changing neutral current effects will be safe. On the other hand, this might lead to interesting signals, for example, in processes such as $\mu \to e\gamma$.

Another option to get neutrino masses is to forbid the usual LH_uN term in the superpotential, but include the non-renormalizable operator LH_uN/M_* in the Kähler potential. After inserting appropriate powers of the chiral compensator, and expanding $\langle\phi^\dagger\rangle = 1 + \bar{m}_{3/2}\bar{\theta}^2$, a term LH_uN is generated in the superpotential, with Yukawa coupling of order $m_{3/2}/M_*$, allowing for Dirac neutrinos of the correct scale [22].

6. Conclusion

We have presented the first fully viable model of anomaly mediation. Using the sequestering mechanism of Luty and Sundrum, we are assured that anomaly-mediated supersymmetry breaking dominates, thereby solving the flavor problem. Supplementing their sequestering mechanism with a mechanism for generating D-terms, we have given an example of a UV-insensitive four dimensional supersymmetry breaking.

Acknowledgments

This work was supported in part by the DOE Contract DE-AC03-76SF00098 and in part by the NSF grant PHY-0098840.

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