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Viable ultraviolet-insensitive supersymmetry breaking*

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Abstract: It is known that one can add D-term contributions for $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_{B-L}$ to the anomaly-mediated supersymmetry breaking to make the superparticle spectrum phenomenologically viable. We point out that this can be done without spoiling its important virtue, namely the ultraviolet insensitivity. This framework can be derived from supersymmetry breaking and $\mathrm{U}(1)_{B-L}$ breaking on hidden brane(s).

Keywords: Anomalies in Field and String Theories, Supersymmetry Breaking, Beyond Standard Model, Supersymmetric Standard Model.

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Contents

1. Introduction

Supersymmetry is one of the most motivated candidates for physics beyond the standard model because it makes the hierarchy between the electroweak scale and a fundamental scale, such as the Planck scale, stable against radiative corrections. Of course, supersymmetry needs to be softly broken in Nature as we have not yet observed any superparticle. The main challenge in supersymmetric model building is probably the flavor problem: arbitrary soft breaking of supersymmetry would induce too-large flavor-changing neutral-current effects. This problem can be swept under the rug by making specific assumptions about the superparticle mass spectrum at the electroweak scale. However, justifying such assumptions in a consistent framework which includes the explanation for the fermion mass hierarchy proves difficult. The main reason for the difficulty is that any models which explain the fermion mass hierarchy must distinguish three generations of quarks and leptons. Their structures should also be sufficiently complicated to generate the observed fermion mass hierarchy. For instance, in the so-called Froggatt-Nielsen mechanism [1], there are a bunch of heavy fields which distinguish different flavors and generate the effective Yukawa couplings when they are integrated out. These dynamics generically make the masses of squarks and sleptons non-degenerate, and hence cause dangerous flavor-changing effects [2, 3].

One way to avoid such a problem is to assume that superparticle masses are generated at an energy scale much lower than the energy scale of flavor physics which explains the fermion mass hierarchy. This is indeed the idea behind the gauge mediation of supersymmetry breaking [4, 5] (and recently revived by [6]). No matter how complicated the flavor physics is, its only low-energy consequences are the Yukawa matrices of quarks and leptons below the scale of flavor physics. The superparticle masses are generated at much lower energy scales where the only flavor-dependent parameters in the lagrangian are the Yukawa matrices. This makes the flavor physics completely hidden from the superparticle masses and hence safe phenomenologically. The same is true of recently proposed gaugino mediated supersymmetry breaking [7, 8]. The scale of flavor physics should be above the compactification scale where the superparticle masses are generated, and there is barely room for flavor physics above the compactification scale. However, flavor-mechanisms of a purely extra-dimensional nature [9, 10] should work well in this context [11].

Anomaly mediation of supersymmetry breaking [12, 13] is another possibility. If soft supersymmetry-breaking parameters are generated from superconformal anomalies, their pattern at an energy scale is determined only by the physics at that energy scale. The only source of supersymmetry breaking in this scheme is in the auxiliary component of the supergravity multiplet, and hence there is only one free parameter: the overall scale of supersymmetry breaking. Then the pattern of supersymmetry breaking does not depend at all on physics at higher energy scales. This is the extremely interesting property: a complete ultraviolet (UV) insensitivity. Even though the flavor physics occurs at the energy scale below the scale where supersymmetry-breaking parameters are generated, the resulting low-energy parameters do not remember the flavor physics at all. This happens in a rather non-trivial fashion. Above the scale of flavor physics, there are additional interactions which distinguish different generations, and the soft supersymmetry-breaking parameters are also highly flavor-dependent. But when heavy particles, such as Froggatt-Nielsen fields, are integrated out, the threshold corrections due to the loops of heavy particles precisely cancel the flavor-dependence of the supersymmetry-breaking parameters.

Despite the appeal of UV insensitivity, anomaly mediation of supersymmetry breaking is excluded because of its high predictivity: sleptons masses-squared are negative and hence the theory would break electromagnetism. There have been proposed various fixes to this problem [14, 15, 16]. Perhaps the most elegant idea is to use non-decoupling effects [14]: if a heavy threshold arises due to supersymmetry-breaking effects, integrating out the heavy particles would modify the low-energy breaking parameters off the anomaly-mediated trajectory. Unfortunately, this casts doubt on the very virtue of the UV insensitivity; if the non-decoupling physics is flavor-dependent, the absence of flavor-changing effects cannot be guaranteed.

In this paper, we point out that there is a different way to modify the anomaly-mediation of supersymmetry breaking in such that the complete UV insensitivity is preserved. In addition, the scalar masses can be modified to make slepton masses positive. This is done by introducing more sources of supersymmetry breaking, namely the auxiliary components in the abelian gauge multiplets. The gauge multiplets are not dynamical; their only raison d'être is to break supersymmetry in their D-components. If the associated fictitious gauge symmetry is non-anomalous under the physical gauge groups, the D-components renormalize among themselves and their contributions to the scalar masses are completely determined by the fictitious gauge charges of the matter fields. Because in a given theory there are only a finite number of fictitious anomaly-free abelian gauge symmetries one can consider, the number of parameters is still small. There are only two additional free parameters in the case of the Minimal Supersymmetric Standard Model (MSSM) or Next-to-MSSM (NMSSM): D-components of $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_{B-L}$.

The model we present is distinct from that of [16] in the following way. The authors of [16] consider adding $\sim (\mathrm{TeV})^2$ Fayet-Iliopoulos terms to the theory to make slepton masses-squared positive. This suggests the existence of an additional U(1) gauge symmetry at the weak scale (along with standard-model singlets to cancel anomalies). In our approach, only the MSSM or NMSSM fields appear at the weak scale while the weak-scale D-terms can be generated by breaking $\mathrm{U}(1)_{B-L}$ at an arbitrarily high scales. In addition, there is no discussion in [16] of the origin of the Fayet-Iliopoulos terms, nor of the complete UV insensitivity of introducing supersymmetry breaking through background D-terms.

This paper is organized as follows. In section 2, we review the basics of anomaly mediation. We then show that adding D-terms for non-anomalous abelian symmetries to the anomaly mediation still gives a consistent trajectory for renormalization-group equations (RGEs) in section 3. In particular, we show that the pattern of soft parameters is determined completely by physics at the energy scale of interest (UV insensitive) with a small number of independent parameters. In section 4, we present a model where one obtains anomaly-mediated supersymmetry breaking together with D-term contributions, by breaking supersymmetry and the abelian gauge groups on hidden brane(s). We derive the superparticle spectrum in section 5. In section 6 we show that small Dirac neutrino masses are generated naturally within our model with order of magnitudes remarkably consistent with the atmospheric neutrino data. Finally we conclude in section 7.

2. Anomaly mediation

Anomaly mediation of supersymmetry breaking assumes that the only source of supersymmetry breaking is in the F-component of the Weyl compensator field ϕ in supergravity: $\phi = 1 + \theta^2 m_{3/2}$. Here and below, we follow the discussions in [14]. Because the Weyl compensator appears in front of every dimensionful parameter, it appears in particular in front of the UV cutoff of the theory, or equivalently, of the

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renormalization scale μ. Therefore, the coefficient of the matter kinetic term

$$\int d^4\theta \mathcal{Z}_i(\mu) Q_i^\dagger Q_i \,, \tag{2.1}$$

where Q denotes a matter chiral multiplet in general, is given by

$$\mathcal{Z}_i(\mu) = Z_i \left(\frac{\mu}{\sqrt{\phi\phi^\dagger}} \right) . \tag{2.2}$$

Here, Z_i is the wave function renormalization factor in the *supersymmetric* theory. In general, the expansion of the matter kinetic term gives the trilinear couplings and the soft scalar masses-squared

$$\ln \mathcal{Z}_i(\mu) = \ln Z_i(\mu) + (\theta^2 A_i(\mu) + \text{h.c.}) - \theta^2\bar{\theta}^2 m_i^2(\mu). \tag{2.3}$$

The trilinear couplings for the superpotential term $\lambda Q_i Q_j Q_k$ is given by $\lambda(A_i + A_j + A_k)\tilde{Q}_i\tilde{Q}_j\tilde{Q}_k$. Comparing eqs. (2.2), (2.3), we find

$$A_{ijk} = -\frac{1}{2}(\gamma_i + \gamma_j + \gamma_k) m_{3/2} \,, \tag{2.4}$$

$$m_i^2 = -\frac{\dot{\gamma}_i}{4} |m_{3/2}|^2 \,. \tag{2.5}$$

Here, γ_i is the anomalous dimension and $\dot{\gamma}_i = d\gamma_i / d\ln\mu$.

The gauge kinetic term

$$\int d^2\theta \, S(\mu) W^\alpha W_\alpha \,, \tag{2.6}$$

is renormalized only at the one-loop level due to holomorphy [17, 18], while the physical (canonical) coupling corresponds to the real superfield $R(\mu)$ with

$$R(\mu) - \frac{T_G}{8\pi^2} \ln R(\mu) = S(\mu) + S^\dagger(\mu) - \frac{1}{8\pi^2} \sum_i T_i \ln \mathcal{Z}_i(\mu), \tag{2.7}$$

where

$$R(\mu) = \frac{1}{g^2(\mu)} + \left(\theta^2 \frac{2m_\lambda(\mu)}{g^2(\mu)} + \text{h.c.} \right) + \theta^2\bar{\theta}^2 R_D. \tag{2.8}$$

We are not interested in R_D in this paper. Again using the fact that the renormalization scale comes together with the Weyl compensator

$$R(\mu) = g^{-2} \left(\frac{\mu}{\sqrt{\phi\phi^\dagger}} \right) , \tag{2.9}$$

we find

$$m_\lambda = \frac{\beta(g^2)}{2g^2} m_{3/2} \,. \tag{2.10}$$

The one unfortunate feature of the above spectrum of soft parameters is the sign of the sleptons squared masses. This comes from the fact that, for fields Q_i without large Yukawa couplings, the dominant contributions to $\dot{\gamma}_i$ are proportional to the beta functions of gauge couplings. Since the SU(2) and U(1) gauge groups of electroweak theory are both asymptotically nonfree in the MSSM, the color-blind sleptons receive large negative contributions to their squared masses. The minimum of the potential in the above theory breaks electromagnetism near the scale of the soft masses ruling out the model.

One otherwise remarkable feature of the spectrum which compounds the slepton mass problem is its UV insensitivity. Though the soft parameters are generated at the Planck scale, the functional form of eqs. (2.4), (2.5), (2.10) are valid at the weak scale. The fact that they indeed solve the RGEs at all orders in perturbation theory provides a non-trivial check. No matter how complicated physics is at high energies, which presumably distinguish flavor in order to generate fermion mass hierarchies via, e.g., the Froggatt-Nielsen mechanism [1], the low-energy soft supersymmetry-breaking parameters are uniquely determined by a single parameter $m_{3/2}$ and hence the flavor breaking is only given by Yukawa couplings. This makes the constraints from the lack of flavor-changing processes essentially automatically satisfied.

Note that the term UV insensitivity refers to the lack of dependences on the particles and interactions in the ultraviolet in the observable sector. The way the supersymmetry breaking effects are induced in the observable sector definitely depends on the coupling between the hidden and observable sectors. The point is that once a particular coupling between two sectors is specified, in this case by putting them on different branes with no light moduli mediating supersymmetry breaking in the bulk [12], we obtain supersymmetry breaking effects reviewed in this section no matter what particles and interactions there are in the observable sector. This is the property preserved in the modification of the anomaly-mediated supersymmetry breaking which will be discussed in the next section.

3. D-term contributions

In the previous section, we reviewed how the supersymmetry breaking in the Weyl compensator field $\phi = 1 + \theta^2 m_{3/2}$ induces supersymmetry-breaking effects via superconformal anomalies. Here we extend the discussion by including a fictitious abelian gauge field $V = \theta^2\bar{\theta}^2 D$ with $D \neq 0$. We show that the low-energy supersymmetry-breaking parameters are completely insensitive to physics in the UV analogous to the purely anomaly-mediated case. The RGE property under the D-term background was also discussed in [16, 19]. Here we give an all-order proof using superfield spurion formalism and also discuss the complete UV insensitivity. We use the notations of [20].

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The idea is to identify an anomaly-free U(1) symmetry in the theory and to pretend that it is gauged. This introduces a fictitious gauge multiplet without a kinetic term. We assume that the gauge multiplet has a supersymmetry-breaking spurion in its D-component:

$$V = \theta^2 \bar{\theta}^2 D\,. \tag{3.1}$$

If needed, we can further introduce a "spectator" field to cancel the $\mathrm{U}(1)^3$ anomaly as well, which has no interactions with physical fields in the superpotential. But because this gauge multiplet is non-dynamical, the spectator fields then do not have any interactions from the rest, and hence do not modify the discussions at all except for justifying the "gauging" of the U(1) symmetry fictitiously.

Because we know the U(1) charges of the matter fields under this fictitious U(1) gauge symmetry, it uniquely determines how the D-term couples to the matter fields, and the coefficients $\mathcal{Z}_i$ are changed to

$$\mathcal{Z}_i \to \mathcal{Z}_i e^{q_i V}\,, \tag{3.2}$$

where q_i is the U(1) charge of the chiral multiplet Q_i. This coupling contributes to the scalar masses

$$m_i^2 \to \bar{m}_i^2 = m_i^2 - q_i D\,. \tag{3.3}$$

It is easy to see that if $m_i^2(\mu)$ is a RGE trajectory, so is $\bar{m}_i^2(\mu)$. The point is that the counter term to the kinetic term takes the form $\mathcal{Z}_0 = \mathcal{Z}(\mu) + \delta\mathcal{Z}$, where

$$\delta\mathcal{Z}_i = \mathcal{Z}_i C\left(\frac{|\lambda_{jkl}|^2}{\mathcal{Z}_j\mathcal{Z}_k\mathcal{Z}_l}, \frac{|\Lambda|}{\mu}, S + S^\dagger - \frac{1}{4\pi^2}\sum_j T_j \ln \mathcal{Z}_j\right), \tag{3.4}$$

because of the reparameterization invariances

$$\begin{aligned} Q_i &\to e^{A_i} Q_i, & \mathcal{Z}_i &\to e^{-(A_i + A_i^\dagger)}\mathcal{Z}_i\,, \\ S &\to S - \frac{1}{4\pi^2}\sum_j T_j A_j, & \lambda_{ijk} &\to e^{-(A_i + A_j + A_k)}\lambda_{ijk}\,. \end{aligned} \tag{3.5}$$

The shift in S is due to the Konishi anomaly [21, 18]. In the function C, the dependence on the superpotential couplings have the product of three kinetic coefficients $\mathcal{Z}_j\mathcal{Z}_k\mathcal{Z}_l$. Under the modification eq. (3.2), this product changes by $e^{(q_j+q_k+q_l)V}$. However, the fictitious U(1) gauge invariance demands that $q_j + q_k + q_l = 0$ for the superpotential coupling $\lambda_{jkl}Q_jQ_kQ_l$ to be allowed, and hence there is no dependence on V here. Another possible dependence is in $S+S^\dagger - \frac{1}{4\pi^2}\sum_j T_j \ln \mathcal{Z}_j$, where the possible shift is given by $-\frac{1}{4\pi^2}\sum_j T_j q_j V$. However, the condition that the fictitious U(1) symmetry is anomaly-free under the true gauge groups $\sum_j q_j T_j = 0$ tells us that this possible shift actually vanishes. Thus, an RGE trajectory remains an RGE trajectory even after the change eq. (3.2). Therefore, the anomaly-mediated supersymmetry-breaking parameters, together with the D-term contributions, would solve the RGE and are hence consistent.

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The main result of this paper is that the anomaly-mediated supersymmetry breaking together with the D-term contributions preserves the UV insensitivity of the purely anomaly-mediated case. The argument is very simple. No matter what the high-energy theory may be, it should preserve the fictitious U(1) gauge symmetry whose gauge multiplet contains the D-term background. After integrating out the heavy fields, the low-energy fields should still satisfy the same fictitious U(1) gauge symmetry, which uniquely determines the D-term contributions. The necessary and sufficient condition for this to hold is that the fictitious U(1) gauge symmetry is not spontaneously broken by high energy physics.

It is important that there is an additional operator in the low-energy theory consistent with the fictitious U(1) gauge invariance if there are U(1) gauge fields, such as $\mathrm{U}(1)_Y$ in the standard model. From the fictitious e^V, one can construct the fictitious $W_\alpha = \theta_\alpha D$. Then the operator

$$c \int d^2\theta \, W_\alpha W_Y^\alpha = cDD_Y = \xi_Y D_Y \, , \tag{3.6}$$

is allowed, and this is nothing but the Fayet-Iliopoulos D-term $\xi_Y \equiv cD$ for the $\mathrm{U}(1)_Y$ gauge multiplet. c is an arbitrary coefficient for this operator. Therefore, in the low-energy theory, all the scalar masses-squared are shifted by $q_i D$, and there is a possible Fayet-Iliopoulos D-term for each of the U(1) gauge factors, with arbitrary coefficients.

Note that the spurion parameter D does not run, while ξ_Y does being a renormalizable coupling. For example, ξ_Y receives renormalization proportional to $\mathrm{Tr}\, q_i Y_i$ at the one-loop level. An interesting point to check is the reparameterization invariance of the D-terms. Since the total masses-squared including the Fayet-Iliopoulos term is

$$m_i^2(\mu) \to \bar{m}_i^2(\mu) = m_i^2(\mu) - q_i D + g_Y^2(\mu) Y_i \xi_Y(\mu) \, , \tag{3.7}$$

one can change the definition of the U(1) charge q_i as $q_i' = q_i + \eta Y_i$ while changing the Fayet-Iliopoulos term as $\xi_Y' = \xi_Y + \eta D/g_Y^2$. The renormalization of ξ_Y' also changes accordingly including now a term proportional to $\mathrm{Tr}\, q_i' Y_i$. The running of soft parameters is unaffected iff D is unrenormalized in both bases because the renormalization of ξ_Y' changes and cancels the redefinition of other parameters.

4. Model

It is easy to justify the assumptions of our picture, namely the controlled introduction of supersymmetry breaking through the Weyl compensator and the D-term of a fictitious gauge multiplet.

It was pointed out in [12] that a physical separation between the observable and the hidden sectors by putting them on different branes results in their separation in

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the Kähler density, and hence the Kähler potential takes the form

$$K = -3 \ln \left(1 - \frac{1}{3} f_{\rm obs}(Q, Q^\dagger) - \frac{1}{3} f_{\rm hid}(H, H^\dagger) \right), \tag{4.1}$$

where H denotes the hidden sector fields and $f_{\rm obs}$, $f_{\rm hid}$ are arbitrary real functions. The superpotential has also a similar separation between observable and hidden pieces. This separation was shown to give vanishing soft masses-squared at the tree-level and the supersymmetry breaking is given purely by anomalies.

We propose an additional U(1) gauge field in the bulk, which is broken on the hidden brane. The anomaly cancellation on the observable brane is obviously required. Because the Kähler potential has two independent functions $f_{\rm obs}$ and $f_{\rm hid}$, even though the true gauge invariance is only a single U(1), there are two *global* U(1) symmetries acting on observable and hidden fields separately. This $U(1) \times U(1)$ is reduced to the diagonal subgroup only through charged bulk fields coupling to both branes. If these charged fields are heavy, therefore, the two global U(1)'s are preserved up to exponentially small effects. After the U(1) group is broken on the hidden brane, the zero-mode gauge multiplet can pick up a non-vanishing D-component due to the supersymmetry-breaking effects on the hidden brane. This effect introduces e^V with $V = \theta^2 \bar\theta^2 D$ as discussed in the previous sections.

A simple toy model is given as follows. Let us consider three chiral superfields X, ψ and $\bar\psi$ living on the hidden brane with the U(1) charges $X(0)$, $\psi(+1)$ and $\bar\psi(-1)$. We introduce the superpotential

$$W = X(\psi\bar\psi - \mu^2), \tag{4.2}$$

which forces ψ, $\bar\psi$ to acquire expectation values. Now, since ψ and $\bar\psi$ fields are living on the hidden brane, they can directly couple to the field z which breaks supersymmetry by its F-component in the Kähler potential, i.e., $K \ni \frac{c}{M_*^2} z^\dagger z \psi^\dagger \psi + \frac{\bar c}{M_*^2} z^\dagger z \bar\psi^\dagger \bar\psi$. These couplings generate soft masses for the ψ and $\bar\psi$ fields in addition to the anomaly-mediated ones. If c and $\bar c$ are different, then, the two soft masses are different, $m_\psi^2 \neq m_{\bar\psi}^2$, resulting in different expectation values for the ψ and $\bar\psi$ fields. This causes nonvanishing D-term expectation value for the gauge multiplet at low-energy, giving soft masses for the observable fields proportional to their U(1) charges. Since the D-term expectation value is of order $m_\psi \simeq c F_z / M_*$, however, we need somewhat small coefficients $c \sim \bar c \sim 1/((16\pi^2)^2 M_* R)$ to make the D-term contribution comparable to the anomaly mediated one. Here, R is the compactification radius for the extra dimension.

The above dynamics can also be understood in the following way using superfield spurion language. The Käher potential for the ψ and $\bar\psi$ fields can be written as

$$K = \psi^\dagger e^{U+V} \psi + \bar\psi^\dagger e^{\bar U - V} \bar\psi, \tag{4.3}$$

where V is the U(1) gauge multiplet, while $U = -\theta^2\bar{\theta}^2 m_\psi^2$ and $\bar{U} = -\theta^2\bar{\theta}^2 m_{\bar{\psi}}^2$ parameterize the soft masses for the ψ and $\bar{\psi}$ fields coming from direct couplings to the z field. Now we integrate out all the heavy fields in the theory in a supersymmetric manner. When ψ and $\bar{\psi}$ acquire expectation values, we can go the unitary gauge $\psi = \bar{\psi} = \mu$ and find

$$K = \mu^2 (e^{U+V} + e^{\bar{U}-V}) . \tag{4.4}$$

Minimizing it with respect to V, we find $V = -(U - \bar{U})/2 = \theta^2\bar{\theta}^2(m_\psi^2 - m_{\bar{\psi}}^2)/2$, and hence the gauge multiplet acquires an expectation value in its D-component. Since the gauge multiplet couples to matter fields Q_i on the observable brane such as

$$K = Q_i^\dagger e^{q_i V} Q_i , \tag{4.5}$$

it induces D-term contribution to the matter scalar masses according to their U(1) charges q_i.

An important point is that the *global* U(1) symmetry in the effective theory of the observable brane is not broken even after ψ and $\bar{\psi}$ fields have expectation values, because of the accidental global U(1) symmetry in the Kähler potential at the tree-level. Therefore, the interactions among the observable fields respect the U(1) invariance which is necessary for the soft masses to be on the trajectory defined by eq. (3.3). One may worry that loop diagrams by the bulk U(1) gauge field can break the U(1) symmetry on the observable brane; however such effects can be made arbitrarily small by making the U(1) gauge coupling constant small, while the D-term contribution to the scalar masses is independent of the size of the U(1) gauge coupling constant, as we have seen already explicitly and has been known for some time [3].

One can also make a more realistic model in which the anomaly-mediated and D-term contributions appear naturally at the same order. Let us consider that the U(1)-breaking fields ψ and $\bar{\psi}$ live on the third brane which is different from the observable and the supersymmetry-breaking branes. Suppose that $\psi = \bar{\psi}$ direction is a flat direction, and ψ and $\bar{\psi}$ have different interactions in the superpotential. Then, the anomaly-mediated contribution generates expectation values for these fields but two expectation values are slightly different, so that there remains nonvanishing D-term of the order of the soft masses generated by anomaly mediation. This D-term contribution is nondecoupling and gives soft supersymmetry-breaking masses to the observable fields proportional to their U(1) charges [14]. A crucial difference from the previous toy model is that the resulting D-term contribution is automatically the same order with the anomaly-mediated one. Once again, an important point here is that since ψ and $\bar{\psi}$ fields live on a different brane than the observable one, they do not directly couple to the standard-model fields. Therefore, no matter how large the U(1) gauge symmetry is broken by the expectation values of the ψ and $\bar{\psi}$ fields, the interactions among the standard-model fields still respect the U(1) invariance. The only assumption needed is that there is no light bulk field which carries nonvanishing U(1) charge.

In order to understand the importance of breaking the bulk U(1) on a different brane, consider the following simple toy model where direct couplings between the U(1) breaking fields and the standard-model fields are allowed. In addition to the light fields, there are superheavy fields $\Psi, \bar{\Psi}$ with U(1) charges of $-2, +2$ and a field φ with charge $+1$ which gets an expectation value to break the U(1) symmetry. Suppose a light field l also has charge $+1$. Then we can write the relevant interactions

$$W = l\Psi\varphi + M_\Psi \Psi\bar{\Psi}\,. \tag{4.6}$$

Upon integrating out Ψ and $\bar{\Psi}$, we generate the following term in the Kähler potential:

$$\frac{1}{M_\Psi^2}(\varphi l)^\dagger e^{2V}(\varphi l)\,. \tag{4.7}$$

Once φ acquires its expectation value, and inserting the D-term for V, this gives a contribution to the l soft mass which is not proportional to its U(1) charge and is off the trajectory. Of course, M_Ψ preserves U(1) and hence there is no reason for it to be as low as φ; if $\varphi \ll M_\Psi$ the deviation from the trajectory can be suppressed. But $\varphi \ll M_\Psi$ is an additional assumption spoiling the spirit of UV insensitivity we are seeking.

Note that the U(1) hypercharge gauge field is localized on the observable brane and does not get a D-term in the above way. However, once the bulk U(1) gauge multiplet obtains the D-term expectation value, the Fayet-Iliopoulos D-term for other U(1) gauge multiplets on the observable brane can be generated either from the tree-level kinetic mixing terms $\int d^2\theta\, W_{i\alpha} W_j^\alpha$ or from loop corrections. Specifically, the Fayet-Iliopoulos term for the U(1) hypercharge can be generated in these ways.

Therefore, the features necessary to introduce supersymmetry breaking in the controlled way we desire, through the F-component of the Weyl compensator field as well as the D-term of a spurious background gauge field, can be naturally obtained by separating the observable and hidden sectors on different branes with a bulk U(1) gauge field broken on a brane other than the observable one.

5. The (N)MSSM

In the MSSM, the squarks and sleptons have known gauge quantum numbers and Yukawa couplings and hence their masses from anomaly mediation are completely calculable. This high degree of predictivity results in disastrous negative masses-squared for the sleptons and hence the model is excluded. However, as shown in the previous section, we can introduce D-terms to modify the scalar masses-squared and it is interesting to ask if such a modification can make all slepton masses-squared positive.

The first question is what non-anomalous U(1) symmetries are there in the MSSM. Assuming there are also neutrino masses (Dirac masses for the sake of the discussions here) and mixing among the leptons, there are only two non-anomalous

U(1) symmetries: $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_{B-L}$. Therefore we can introduce two additional parameters to the anomaly mediated supersymmetry breaking, D_Y and D_{B-L}. Here and below, we use the notation $D_Y = -g_Y^2 \xi_Y$. In the framework of higher dimensional models discussed in the previous section, the $\mathrm{U}(1)_{B-L}$ gauge field is identified with the bulk U(1) gauge field and D_Y is generated through kinetic mixing between $\mathrm{U}(1)_{B-L}$ and $\mathrm{U}(1)_Y$. The slepton masses are then modified as

$$
\begin{aligned}
m_{\tilde{L}}^2 &\to \bar{m}_{\tilde{L}}^2 = m_{\tilde{L}}^2 + \frac{1}{2} D_Y + D_{B-L} \,, \\
m_{\tilde{e}}^2 &\to \bar{m}_{\tilde{e}}^2 = m_{\tilde{e}}^2 - D_Y - D_{B-L} \,,
\end{aligned} \tag{5.1}
$$

where $m_{\tilde{f}}^2$ represents the pure anomaly-mediated contribution. A necessary condition for a viable spectrum is

$$
D_Y < -D_{B-L} < \frac{1}{2} D_Y < 0 \,. \tag{5.2}
$$

It has been shown that the spectrum can in fact be made viable in some region of parameter space [16].

As for the particles in the supersymmetric standard model, we find the gaugino masses

$$
\begin{aligned}
M_1 &= 1.43 M \,, \\
M_2 &= 0.414 M \,, \\
M_3 &= -3.67 M \,,
\end{aligned} \tag{5.3}
$$

where $M = m_{3/2}/(16\pi^2)$. The masses for the first and second generation scalars with Yukawa couplings neglected are

$$
\begin{aligned}
m_{\tilde{L}}^2 &= -0.351 M^2 + \frac{1}{2} D_Y + D_{B-L} \\
m_{\tilde{e}}^2 &= -0.374 M^2 - D_Y - D_{B-L} \\
m_{\tilde{Q}}^2 &= 11.7 M^2 - \frac{1}{6} D_Y - \frac{1}{3} D_{B-L} \\
m_{\tilde{u}}^2 &= 11.8 M^2 + \frac{2}{3} D_Y + \frac{1}{3} D_{B-L} \\
m_{\tilde{d}}^2 &= 11.9 M^2 - \frac{1}{3} D_Y + \frac{1}{3} D_{B-L} \,,
\end{aligned} \tag{5.4}
$$

while for the third generation scalars and Higgs bosons

$$
\begin{aligned}
m_{\tilde{L}_3}^2 &= -0.352 M^2 + \frac{1}{2} D_Y + D_{B-L} \\
m_{\tilde{e}_3}^2 &= -0.377 M^2 - D_Y - D_{B-L} \\
m_{\tilde{Q}_3}^2 &= 9.55 M^2 - \frac{1}{6} D_Y - \frac{1}{3} D_{B-L} \\
m_{\tilde{u}_3}^2 &= 7.54 M^2 + \frac{2}{3} D_Y + \frac{1}{3} D_{B-L} \\
m_{\tilde{d}_3}^2 &= 11.9 M^2 - \frac{1}{3} D_Y + \frac{1}{3} D_{B-L}
\end{aligned}
$$

$$
\begin{aligned}
m_{H_u}^2 &= -6.72M^2 - \frac{1}{2}D_Y \\
m_{H_d}^2 &= -0.402M^2 + \frac{1}{2}D_Y\,,
\end{aligned} \tag{5.5}
$$

for $\tan\beta = 3$ as an example.[1]

An interesting feature of the above spectrum is that while the scalar masses receive contributions from D-terms, the gaugino masses remain those of anomaly mediation. Thus, in a broad region of the parameter space, the lightest supersymmetric particle is the neutral wino nearly degenerate with the charged one, giving rise to signatures discussed in [22]. But the sleptons can also be the lightest supersymmetric particle depending on the parameters M, D_Y and D_{B-L}. Furthermore, since the scalar masses are given by only three parameters, there are a number of sum rules for the scalar masses as derived in [16].

Note that the right-handed neutrino superfield N, which in this minimal model only picks up soft masses from the D-term and the (miniscule) anomaly-mediated contribution from the tiny neutrino Yukawa coupling, gets a *negative* mass squared. However, this difficulty is easily remedied by e.g. adding additional couplings to new standard-model singlets charged under $\mathrm{U}(1)_{B-L}$. For instance, consider new fields $X, \bar{X}$ with charges $-2, +2$ and a singlet Y with zero charge under $\mathrm{U}(1)_{B-L}$. Now add the couplings $YX\bar{X} + N^2X$. Then, the scalar components of $N, Y, X, \bar{X}$ can pick up net positive soft masses from anomaly and D-term mediations.

Now it is well-known that the MSSM probably does not work within the anomaly-mediated supersymmetry breaking because the μ term is associated with a too-large $B\mu$ term because $B \sim m_{3/2}$ rather than $m_{3/2}/(16\pi^2)$. The next simplest possibility is the NMSSM with the superpotential

$$
\lambda S H_u H_d + hS^3\,, \tag{5.6}
$$

where S is a new standard model singlet chiral superfield. Because $m_{H_u}^2$ is negative and $m_{H_d}^2$ less so with the above spectrum, it is reasonable to expect that one can find a correct electroweak symmetry breaking. The details of the phenomenological analysis is beyond the scope of this paper and will be discussed elsewhere.

6. Neutrino masses

Since we have introduced a fictitious $\mathrm{U}(1)_{B-L}$ gauge symmetry, a natural question is how we can explain small neutrino masses. As mentioned earlier, Dirac neutrinos with small Yukawa couplings, possibly as consequences of flavor symmetries, is definitely a natural possibility within our framework. However, one can resort to other mechanisms to explain small neutrino masses consistent with our framework as well.

[1] These mass parameters have been evaluated at a scale $\mu = 500\,\mathrm{GeV}$, for simplicity. To obtain the physical masses, superparticle threshold corrections should be taken into account appropriately.

For instance, we can consider the situation where the Kähler potential contains a term $\frac{1}{M_*}LH_uN$ + h.c. while the superpotential coupling LH_uN is absent. This possibility is naturally realized using a global $\mathrm{U}(1)_R$ symmetry, for example. Then, the Kähler potential term picks up the Weyl compensator

$$\frac{\phi^\dagger}{\phi^2}\frac{1}{M_*}LH_uN\,, \tag{6.1}$$

and by expanding $\phi^\dagger = 1+\bar{\theta}^2 m_{3/2}$, we find a term which is effectively a superpotential LH_uN with a coupling constant of the order of $m_{3/2}/M_* {\sim} 16\pi^2 v(M_{\mathrm{Planck}}R)^{1/3}/M_{\mathrm{Planck}}$ [23, 24]. This size for the Dirac neutrino Yukawa coupling is remarkably consistent with the size expected from the atmospheric neutrino data.

Another possibility is to write down the mass term for the right-handed neutrino by hand: $\frac{M_R}{2}NN$ in the superpotential, giving the standard seesaw mechanism [25]. This of course breaks $\mathrm{U}(1)_{B-L}$ explicitly, but only softly, and does not change the RGEs of soft supersymmetry-breaking parameters at all. The only concern is the threshold corrections when the right-handed neutrinos are decoupled. It is interesting that the threshold corrections are still exactly the same with supersymmetric thresholds which keep the supersymmetry-breaking parameters on the trajectories of our framework at the leading order. One may imagine that such $\mathrm{U}(1)_{B-L}$-breaking mass parameter was induced on our brane by a field carrying $\mathrm{U}(1)_{B-L}$ charge in the bulk, or even simply by a spontaneous breakdown occurring on our brane since dangerous operators such as those given in eq. (4.7) are generically generated only at the loop level. In this framework, however, the supersymmetry-breaking parameters are not exactly on the trajectories below the scale of $\mathrm{U}(1)_{B-L}$ breaking, and one expects corrections at higher orders in $1/(16\pi^2)$.

7. Conclusion

Anomaly mediation of supersymmetry breaking is attractive because the soft supersymmetry breaking parameters at a given energy scale are determined only by physics at that energy scale (UV insensitivity) and hence is highly predictive (only one parameter). The resulting supersymmetry breaking parameters are automatically safe from generating too-large flavor-changing neutral current effects no matter how complicated flavor physics at high-energy scales is. However the anomaly mediation failed phenomenologically because of its high predictivity: slepton masses-squared are negative. It is known that one can add D-term contributions for $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_{B-L}$ to the anomaly-mediated supersymmetry breaking to make the superparticle spectrum phenomenologically viable. In this paper we have shown that one can indeed add these D-terms keeping the complete UV insensitivity no matter how complicated thresholds there are. The only assumption is that the $\mathrm{U}(1)_{B-L}$ is a global symmetry in the supersymmetric standard model.

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This framework can be derived from supersymmetry breaking and $\mathrm{U}(1)_{B-L}$ breaking on hidden brane(s). The $\mathrm{U}(1)_{B-L}$ gauge multiplet lives in the bulk but it is broken at a high scale unlike in the gaugino mediation. Assuming no light bulk fields charged under $\mathrm{U}(1)_{B-L}$ (similar to the usual assumption in anomaly mediation), the physics on the observable brane has a global $\mathrm{U}(1)_{B-L}$ symmetry up to exponentially suppressed effects. The supersymmetry breaking effects are described in terms of the Weyl compensator with an F-component expectation value, a non-dynamical $\mathrm{U}(1)_{B-L}$ gauge multiplet with a D-component expectation value, and a Fayet-Iliopoulos term for $\mathrm{U}(1)_Y$. Because the low-energy supersymmetry breaking is constrained by the superconformal and $\mathrm{U}(1)_{B-L}$ invariances, the soft parameters are determined completely by the physics at the energy scale of interest and hence UV insensitive.

Despite the $\mathrm{U}(1)_{B-L}$ invariance, one can still generate neutrino masses at the correct energy scale for atmospheric neutrino oscillations naturally due to an operator at the fundamental scale.

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