Abelian $D$-terms and the superpartner spectrum of anomaly-mediated supersymmetry breaking

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Abstract

We address the tachyonic slepton problem of anomaly mediated supersymmetry breaking using abelian $D$-terms. We demonstrate that the most general extra $U(1)$ symmetry that does not disrupt gauge coupling unification has a large set of possible charges that solves the problem. It is shown that previous studies in this direction that added both an extra hypercharge $D$-term and another $D$-term induced by $B - L$ symmetry (or similar) can be mapped into a single $D$-term of the general ancillary $U(1)_a$. The $U(1)_a$ formalism enables identifying the sign of squark mass corrections which leads to an upper bound of the entire superpartner spectrum given knowledge of just one superpartner mass.
The pure anomaly-mediated supersymmetry breaking (AMSB) spectrum includes negative squared masses for all sleptons, which inflicts an unwanted Higgs mechanism upon electromagnetism. If it were not for this unfortunate feature, AMSB would be an ideal candidate for the standard model of supersymmetry breaking transmission. It has the virtues of being ultraviolet (UV) insensitive, compatible with flavor changing neutral current constraints, and requiring modest model building from the low-energy perspective. Of course, the viability of any AMSB model with two-loop suppressed scalar masses requires significant cooperation from the the Kähler potential, superpotential, and gauge kinetic function, which could be accomplished by either geometric means of an extra dimensional brane world supersymmetry breaking or a purely four dimensional conformal hidden sector.

Salubrious implications to long-standing issues of supersymmetry are not the only alluring aspects of AMSB. AMSB contributions to the masses are not optional – they are present in all broken supergravity theories. We wish, therefore, to make a few contributions to the discussion of AMSB, and in particular regarding the $D$-term solution of the tachyonic slepton problem due to an extra $U(1)$ symmetry in the theory. Our considerations have led us to a general “ancillary $U(1)_a$”, to be defined below, which nicely parametrizes some of the previous studies in this direction, and enables us to make quick work of some interesting theorems regarding the superpartner spectrum.

Recall that the AMSB masses, which are invariant expressions under renormalization group evolution (“RG invariant”), are in the minimal supersymmetric standard model (MSSM) given by

$$M_\lambda = \frac{\beta_g}{g} m_{3/2}$$

$$m_{\bar{Q}}^2 = -\frac{1}{4} \left( \frac{\partial \gamma}{\partial y} \beta_y + \frac{\partial \gamma}{\partial y} \beta_y \right) m_{3/2}^2$$

$$A_y = -\frac{\beta_y}{y} m_{3/2},$$

where $m_{3/2}$ is the gravitino mass. Since the masses are one-loop suppressed compared to the gravitino mass, it will be convenient to define the mass scale $\tilde{m} \equiv m_{3/2}/16\pi^2$. Thus, $\tilde{m}$ is a more direct estimator of the superpartner masses than $m_{3/2}$. Expressions for the pure AMSB masses in terms of $\tilde{m} = m_{3/2}/16\pi^2$, gauge couplings and Yukawa couplings can be found in many places in the literature (see, e.g., Eqs. (44)-(59) of [6]).

Application of Eq. (2) to sleptons shows that their mass-squared is negative. Many groups have proposed ideas to solve this tachyonic slepton problem. Some solutions lose the UV insensitive feature of AMSB, and some solutions preserve it. We wish to keep UV insensitivity, and retain gauge coupling unification, which is one of the more powerful hints that supersymmetry is a necessary element of the full theory of nature. The combination of these requirements leads us to consider low-scale $D$-terms from broken exotic gauge symmetries.

There are many considerations in employing $D$-term slepton solutions. This list includes
The SM charges under this additional $U(1)_a$ can be classified by two free parameters [18] (see Table 1). We know that the slepton masses must have a large, positive contribution from the $D_a$ term, so we know that $r > 0$ and $q > 0$. This gives us freedom to normalize the charge of the right-handed electron chiral superfield to be $r = 1$, and let the overall size of the mass correction be controlled by the $D$-term vev. The remaining charges are then dependent on only one variable $q$:

$$ Q_e = +1, \quad Q_L = 3q, \quad Q_{H_u} = 3q+1, \quad Q_{H_d} = -(3q+1), $$

$$ Q_Q = -q, \quad Q_u = -(2q+1), \quad Q_d = 4q+1. $$

Table 1: Table of charges for the ancillary $U(1)_a$ symmetry. It is defined to be a symmetry that can be made anomaly free without the need of adding any additional states charged under the SM. It can be parametrized by two rational constants [18]. Hypercharge, $B-L$, and the $U(1)_J$ symmetry of Ref. [10] are all special cases of $U(1)_a$ generated by the appropriate two parameter choice ($q, r$).
We do not address here the precise charges of the exotic SM singlet states $s_i$, except to note that their $U(1)_a$ charges must satisfy
\[ \sum s_i = -3(6q + r) \quad \text{and} \quad \sum s_i^3 = -3(6q + r)^3. \] (6)

Although large Yukawa couplings are generically necessary to break the $U(1)_a$, renormalizable operators do not couple the MSSM matter and the SM singlets charged under the $U(1)_a$. Therefore, there will be no additional AMSB contribution of the Yukawa variety. (Pure singlets, which we eschew, would add Yukawa complications when interacting with $H_uH_d$.) Solutions for $s_i$ charges for a particular set of fields and symmetries is a model building question that we will not discuss here, except to say that some simple cases were analyzed in [13, 10]. Consistent with our present goals, we will only discuss results that must hold for all possible model choices.

We consider the general case of a single arbitrary $U(1)_a$, determine what values of $q$ and $D_a$ are needed to solve the slepton problem, and then determine what predictions follow. To more easily keep track of scale factors, we will define $D_a = \eta \tilde{m}^2$. To generate the entire superpartner spectrum we need only specify
\[ \tilde{m}, \tan \beta, \eta, q, \text{ and } b_ag_a^4 \] (input parameters)
where $b_ag_a^4$ characterizes the anomaly mediated contribution due to the $U(1)_a$ and is clarified in Eq. (9) below.

We can then compute the $D$-term contribution to the superpartner masses:
\[
\begin{align*}
(m_L^{D_a})^2 &= 3q\eta\tilde{m}^2 \\
(m_e^{D_a})^2 &= \eta\tilde{m}^2 \\
(m_Q^{D_a})^2 &= -q\eta\tilde{m}^2 \\
(m_u^{D_a})^2 &= -(2q - 1)\eta\tilde{m}^2 \\
(m_d^{D_a})^2 &= (4q + 1)\eta\tilde{m}^2 \\
(m_{H_u}^{D_a})^2 &= (3q + 1)\eta\tilde{m}^2 \\
(m_{H_d}^{D_a})^2 &= -(3q + 1)\eta\tilde{m}^2.
\end{align*}
\] (8)

Additionally, a light, propagating $Z'$ associated with this $D$-term (which is not necessarily required [10, 14]) would alter the anomalous dimensions of the matter fields by $\Delta \gamma_i = 4(Q_a^i)^2 g_a^2$ resulting in an AMSB contribution to the soft masses:
\[
\begin{align*}
(m_L^{\Delta \gamma})^2 &= -2(3q)^2 b_ag_a^4\tilde{m}^2 \\
(m_e^{\Delta \gamma})^2 &= -2b_ag_a^4\tilde{m}^2 \\
(m_Q^{\Delta \gamma})^2 &= -2q^2 b_ag_a^4\tilde{m}^2 \\
(m_u^{\Delta \gamma})^2 &= -2(2q + 1)^2 b_ag_a^4\tilde{m}^2 \\
(m_d^{\Delta \gamma})^2 &= -2(4q + 1)^2 b_ag_a^4\tilde{m}^2 \\
(m_{H_u}^{\Delta \gamma})^2 &= -2(3q + 1)^2 b_ag_a^4\tilde{m}^2 \\
(m_{H_d}^{\Delta \gamma})^2 &= -2(3q + 1)^2 b_ag_a^4\tilde{m}^2.
\end{align*}
\] (9)
where $b_a$ is the beta function coefficient of $U(1)_a$ ($dg_a/dt = b_ag_a^3/16\pi^2$). The effect of this contribution is to drive the sleptons more negative in squared mass in the infrared.

We can immediately construct a few theorems about the spectrum that arise from the ansatz of ancillary $U(1)$ symmetry. Since contributions must be positive for $m_L^2$, we know that $q > 0$ is necessary. Therefore, by mere inspection of Eqs. (8)- (9), we determine that additional D-term mass contributions to the squarks are always of well-defined sign:

$$ (m^D_Q)^2 < 0, \quad (m^D_u)^2 < 0, \quad (m^D_d)^2 > 0. \quad (10) $$

In the event that $U(1)$ breaking is at a low scale and $(m_i^{A\gamma})^2$ mass contributions are present, we can see that $\eta$ must necessarily be above $2b_ag_a^4$ in order to overcome the additional tachyonic contributions $\Delta \gamma$ imposes on the sleptons. Nevertheless, we note that the up-squark mass still must be shifted lower in mass compared to the pure MSSM AMSB mass ($\Delta m_u^2 < 0$), and the down-squark mass still must be shifted higher in mass ($\Delta m_d^2 > 0$). Only the left-squark doublet mass can be shifted either up or down in this case depending on the precise relationship between the parameters.

Determining the mass of the squarks with respect to the pure-AMSB value is another way to test the theory. Precise measurements of heavy squarks are difficult at colliders, but it is profitable and feasible to answer simple binary questions of the form “Is this squark mass we just measured heavier or lighter than the pure AMSB mass we infer for it from the gaugino spectrum?”

We also can state another important implication to the superpartner spectrum. Given the intrinsic mass scale $\tilde{m} = m_{3/2}/16\pi^2$, we can conclude that

$$ \text{mass of any squark or slepton} < 6\tilde{m} \quad (11) $$

under the broad assumptions we have outlined for this study. Measurement of any one superpartner fixes $\tilde{m}$, and the above equation can then be used to cap the maximum value of any other superpartner in the spectrum. For example, if we measure the lightest chargino and determine that it is mostly Wino, we can use the resulting value of $M_2$ ($M_2 \simeq 0.4\tilde{m}$) to place a limit of about $15M_2$ for any other superpartner. The bound comes from the very conservative requirement that all first generation superpartner masses are positive. The constraints from EWSB symmetry breaking and experimental limits on third generation mass eigenstates makes the result even tighter for various values of $q$, $\eta$ and $\tan \beta$ input parameters.

In Fig. 1 we have demonstrated the parameter space for $q$ and $\eta$ given the other two input parameters fixed at $\tilde{m} = 500$ GeV and $\tan \beta = 10$. Here (and for Fig. 2, as well), we have excluded the effects of Eq. (9). For any given value of $\eta$, $q$ cannot be too large, otherwise some of the squark masses will go tachyonic; and $q$ cannot be too low, otherwise the left-sleptons will remain tachyonic. Likewise, for any given value of $q$, $\eta$ cannot be too

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1 This is the case of Refs. [10] and [14].
The parameter space of AMSB with ancillary $U(1)_a$ is spanned by four parameters, $\tilde{m}$, $\tan\beta$, $\eta$ and $q$. We demonstrate the region of allowed $\eta$ and $q$ that insures that all sfermions have acceptable masses and allows proper EWSB. For this example, we have chosen $\tilde{m} = 500$ GeV and $\tan\beta = 10$. The horizontal lines represent the parameter space that compose the spectrums of Fig. 2.

large, otherwise some of the squarks will go tachyonic; and $\eta$ cannot be too small, otherwise both right and left sleptons will remain tachyonic.

If $g_a \sim g_Y$ at the weak scale, $U(1)_a$ would lower all soft masses of scalars charged under it, resulting in less parameter space. Alternatively, one might ponder starting $g_a(m_{\text{GUT}})$ with the gauge unification value and running $g_a$ to the weak scale. To avoid committing to a specific model, the beta function coefficient $b_a$ may be approximated by excluding the contributions due to the unspecified singlet sector, which generally should contribute less. (Without this assumption, one could find, in principle, parameter space allowing implausibly large values of $b_a$ in the hundreds.) At the weak scale, this typically would lead to smaller $g_a(m_Z) < g_Y(m_Z)$, expanding the allowed parameter space all the way to that of Fig. 1 in the limit $g_a(m_Z) \to 0$.

In Fig. 2, we demonstrate how the spectrum scales with $\eta$ for the other three parameters fixed $\tilde{m} = 500$ GeV, $\tan\beta = 10$, and $q = 0.1, 1, 2$. The $\eta$ values terminate before any squark gets dangerously close to zero mass. This is because EWSB begins to fail. The up and down Higgs squared masses are driven in opposite directions due to their charge assignments. If $m_{H_u}^2$ and $m_{H_d}^2$ are nearly equal, $B\mu$ goes to a small, non-zero, positive value. The classical condition that the Higgs potential be bounded below then fails.

We now wish to compare our use of $D$-terms as an AMSB slepton solution directly to previous studies. Our approach is related to the work in Ref. [10, 14]. Our discussion
Figure 2: The parameter space of AMSB with ancillary $U(1)_a$ is spanned by four parameters, $\tilde{m}$, $\tan \beta$, $\eta$ and $q$. We demonstrate how the first generation scalar soft masses (in GeV) scale with $\eta$ for three different choices of $q$. For this example, we have chosen $\tilde{m} = 500$ GeV and $\tan \beta = 10$. The shaded region on the left is excluded by insufficient slepton mass squared. The shaded region on the right is excluded by improper electroweak symmetry breaking. Notice the entire spectrum is bounded below about $4\tilde{m}$, which is an illustration of how a technical analysis of the spectrum for any given set of inputs usually imply a tighter, lower-mass spectrum than the most general expression in (11). Negative soft mass values are defined as $-\sqrt{-m^2}$.

is designed to imagine $U(1)_a$ freely parameterizing the $D$-terms of either scenario. One simplifying feature of the ancillary $U(1)$ assumption is that it allows for the use of a single exotic $D$-term while maintaining gauge coupling unification.

The study in [10] analyzed the single case $(q, r) = (-7/3, 3)$ originating from the condition $\text{Tr}(YQ) = 0$, which zeroed one-loop corrections to kinetic mixing. We call this special case $U(1)_J$. However, this case does not solve the slepton problem by itself, as the mass-squared contribution to the slepton-left $L$ is opposite sign to that of slepton-right $e$. It is then necessary to introduce an additional $D$-term proportional to the hypercharge of each superpartner, such that a judicious combination of the two $D$-terms would yield positive mass-squareds for both slepton-right and slepton-left. Unspecified, additional model building complexity would be needed to explain the origin of these additional hypercharge $D$-terms.

In that framework, the soft masses that compose the FI scale and singlet VEVs presumably originate from AMSB. However, soft mass corrections from physics beyond the cutoff scale $\Lambda$ of the MSSM effective theory are cancelled exactly to one loop order. Higher order corrections begin on the order $m_{3/2}^4/\Lambda^2$. And so without additional dynamics, relevant scales
cannot be too high if they are to be responsible for sufficient $D$-terms.

The authors of Ref. [14] consider the $B - L$ gauge symmetry, which has opposite sign charges for $L$ and $e$ requiring the use of a second $D$-term. They use a hypercharge $D$-term established through an operator of the form $\int d^2 \theta W Y W_{B-L}$, which, in essence, allows free parameterization for the hypercharge $D$-term. As was the case with $U(1)_J$ $D$-terms, $U(1)_{B-L}$ $D$-terms must be augmented by additional $U(1)_Y$ $D$-terms in order to solve the tachyonic slepton problem in this scenario.

We now demonstrate that our simple parametrization of the $D$ terms is a general parametrization that can describe the entire parameter space of both Ref. [10] and Ref. [14]. In both of these analyses, there are two $D$-term contributions to the scalar masses of the squarks and sleptons. One $D$-term is an arbitrary additional contribution to the ordinary hypercharge $D_Y$-term, and the other $D$-term is proportional to the charges of a special $U(1)$ which is $B - L$ in Ref. [14] or $J$ in Ref. [10]:

$$\Delta m^2_i = Y_i D_Y + Q_{B-L}^i D_{B-L}, \quad \text{or} \quad \Delta m^2_i = Y_i D_Y + Q_J^i D_J.$$

Table 1 contains the charges of each superpartner state in terms of hypercharge, $B - L$, $J$ symmetry, and our general $U(1)_a$ symmetry.

We can recast these $D$-term contributions to the mass of sparticle $i$ in terms of our language (Eq. 4) by setting $\eta = 1$ and $r = 1$, and then identifying $D_a$ and $q$ with

$$D_a = r_Y D_Y + r_X D_X, \quad (12)$$

and

$$q = \frac{q_Y D_Y + q_X D_X}{r_Y D_Y + r_X D_X} \quad (13)$$

where $X$ is a label for either $B - L$ or $J$, $(q_Y, r_Y) = (-1/6, 1)$, $(q_{B-L}, r_{B-L}) = (-1/3, 1)$, and $(q_J, r_J) = (-7/3, 3)$.

These simple mappings of the two $D$-terms analyses of Refs. [10] and [14] are made possible because all $U(1)$’s under discussion are special cases of $U(1)_a$. This allows a vector space to be defined on the “basis vectors” of $q$ and $r$. The vector equation of equality between the $U(1)_a$ $D$-term parametrization and the other parametrizations can be written as

$$(\vec{v}_a - \vec{v}_Y - \vec{v}_X) \cdot \vec{X}_i = 0 \quad (14)$$

where $\vec{v}_a = (q, r) D_a$, $\vec{v}_Y = (q_Y, r_Y) D_Y$, $\vec{v}_X = (q_X, r_X) D_X$, and $\vec{X}_i = (a_i, b_i)$ such that $Q_{a_i}^i = a_i q + b_i r$. The solution is therefore independent of $i$ (i.e., independent of particle type), and given by Eqs. (12) and (13).

As a short digression, we remark that the $D$-term solution of Ref. [7], on the other hand, cannot be mapped using our analysis because the symmetries considered are not special
cases of $U(1)_a$. While that study cleverly exploited the higher-order non-decoupling of $D$-term contributions to mass, the examples were not simply consistent with gauge coupling unification. The features of the symmetries that disrupt gauge coupling unification in that example are also what make it incapable of being described as a special case of $U(1)_a$.

In principle, one could tell the difference between a general $U(1)_a$ $D$-term contribution to the supersymmetry masses and contributions that arise from the combination of $D_Y$ and $D_{B-L}$ (or $D_J$). The typical map of the latter to the former (Eqs. 12 and 13) generates an irrational $q$ charge. This means the ratio of $U(1)_a$ charges among the superpartners would be irrational, which is not expected for a sensible $U(1)$ theory. Unfortunately, we see no reasonable way these possibilities could be distinguished experimentally via superpartner measurements alone.

To add to the potential complexity of interpreting experiment, kinetic mixing of the form $\int d^2 \theta W_Y W_a$ can also significantly change the effective charges, even if there is only one extra $U(1)_a$ symmetry in the low-energy theory. If we apply the methods of Ref. [19] to the present case of a not-so-hidden extra $U(1)_a$, we find that a hypercharge-mixed $D$-term contribution arises. Its effects can be parametrized as,

\[ D_{a}^{\text{eff}} \rightarrow (1 + \chi r_Y) D_a \]  \hspace{1cm} (15)

and

\[ q^{\text{eff}} \rightarrow \frac{q + \chi q_Y}{1 + \chi r_Y}. \]  \hspace{1cm} (16)

This is equivalent to the generation of the hypercharge $D$-term of Ref. [14], except we are mixing it with a $U(1)_a$ symmetry that is not necessarily the special case $U(1)_{B-L}$. We then map it back into a single $D$-term of $U(1)_a$ again, albeit with a new charge $q^{\text{eff}}$ which is likely to be irrational.

In the preceeding, we have shown that we can map the entire space of interesting $D$-term induced solutions to the tachyonic slepton problem that do not disrupt supersymmetric gauge coupling unification into the two parameters $q$ and $D_a$ of the ancillary $U(1)_a$. Furthermore, this general result is especially satisfying, because it allows us to prove several interesting features of the superpartner spectrum independent of any specific model details. For example, it manifestly demonstrates that there is a definite sign of the $D$-term mass correction to each superpartner in this general approach. It also allowed us to prove the bound of all superpartners as expressed in Eq. (11) within any specific model consistent with our starting axiom that gauge coupling unification is not disrupted. A $D$-term solution of this motivated type leads to these unambiguous propositions regarding the superpartner spectrum and is therefore subject to experimental refutation.
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