The NMSSM, Anomaly Mediation and a Dirac Bino

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Abstract

We introduce a new model of supersymmetry breaking dominated by anomaly mediation. It has a viable spectrum, successful electroweak symmetry breaking, solves the mu-problem and maintains the anomaly-mediated form for soft-masses down to low energies thus solving the flavor problem. The model consists of the minimal supersymmetric standard model plus a singlet, anomaly-mediated soft masses and a dirac mass which marries the bino to the singlet. We describe a large class of models in the UV which can produce such boundary conditions. The dirac mass does not affect the so-called “UV insensitivity” of the other soft parameters to running or supersymmetric thresholds and thus flavor physics at intermediate scales would not reintroduce the flavor problem. The dirac bino is integrated out at a few TeV and produces finite and positive contributions to all hypercharged scalars at one loop thus producing positive squared slepton masses. The theory predicts some CP violation in the Higgs sector leading to a correlation between the spectra to be seen at the LHC and electric dipole moments within experimental reach in the near future.
1 Introduction

Anomaly Mediated supersymmetry breaking (AMSB) [1] (see also [2]) is a very predictive form of supersymmetry breaking with many desirable features, chief amongst these is its insensitivity to intermediate scales (so-called “UV insensitivity”). Soft masses and scalar couplings are determined by low energy couplings, independent of the details of the running from the high scale. Unfortunately this predictivity is also its major drawback. Sleptons, being charged under non-asymptotically free gauge groups, have negative squared masses. Attempts to fix this problem [3–22] often result in reintroducing a dependence on the UV scale. Here we propose a modification of pure anomaly mediation that fixes the slepton mass problem whilst retaining UV insensitivity.

Anomaly mediation also has a $\mu$-problem. Including a $\mu$ term in the superpotential explicitly breaks the conformal symmetry and generates a $B\mu$ term that is a loop factor too large. If instead one works with the next-to-minimal supersymmetric standard model (NMSSM) then conformal invariance is not explicitly broken at tree level and it is possible to get correct EWSB. So, we will work with the NMSSM which includes a SM singlet, $S$. The same singlet will also be used to fix the tachyonic slepton problem in a UV insensitive way.

The feature we add to AMSB is Supersoft Supersymmetry breaking (SSSB) [23]. SSSB is a way of generating Dirac gaugino masses which, when integrated out, produce finite positive squared scalar masses a loop factor smaller than the gaugino mass. SSSB is UV insensitive since the scalar masses only run once the gauginos have been integrated out. In our model, the bino marries a singlet and we show that this singlet can be the same singlet that appears in the NMSSM. We have a supersoft contribution for fields charged under $U(1)'$ from a Dirac mass for (only) the bino.

The negative squared masses are generated at two loops by anomaly mediation whereas the positive SSSB contributions are generated at one loop. The scale of the AMSB contribution is set by the gravitino mass and the SSSB contribution by the D-term vev of a hidden-sector $U(1)'$. We will demonstrate how these two scales arise in a model of dynamical SUSY breaking and can be comparable. The result is that the slepton mass squareds are pushed positive and the spectrum becomes viable.

The rest of this paper is laid out as follows. In Section 2 we review AMSB and SSSB and set conventions for the remainder of the paper. We also discuss their mutual UV insensitivity. We then address the issue of EWSB, in
Sections 3 and 3, and calculate the spectrum for the squarks and sleptons. We treat the relative sizes of the two contributions as a free parameter and show that varying this parameter produces viable spectra. In Section 4 we discuss CP violation and show that while CP is maximally broken in the Higgs sector, our model lives at the boundary of current experimental probes. In Section 5 we present possibilities for the UV physics that reproduces our model. For example, until this stage we have treated the two supersymmetry-breaking scales as independent. In 5.2 we discuss an explicit model of dynamical SUSY breaking that relates these two scales and we show how they typically have the right ratio. In 5.3 we present extra-dimensional realizations that naturally suppress all other sources of supersymmetry breaking. In 5.1 we address the issue of the one remaining dangerous operator, kinetic mixing between hidden and visible sector $U(1)$s, and discuss possibilities for natural suppression. In Section 6 we conclude.

2 Anomaly and Supersoft Mediation

In order for AMSB to dominate the hidden sector which breaks SUSY must be sequestered from the MSSM. This forbids contact interactions that would otherwise dominate. The sequestering was done originally via a five-dimensional setup [1], but recently was realized entirely in four dimensions using a strongly coupled CFT [24,25]. We will need the same sequestering with the exception of couplings between a hidden sector gauge field and visible sector fields. This will be accomplished by putting the hidden sector gauge fields in the bulk. We will show explicitly in Section 5 how to generate only the desired operators.

In such a setup the SUSY breaking is a result of the F-term of the conformal compensator, $\Phi = 1 + \theta^2 m_{3/2}$, whose value is determined by tuning the cosmological constant to zero. After rescaling fields the conformal compensator appears with the cutoff of the theory, as well as any explicit mass scales in the Lagrangian. When regulating loops this $\theta$ dependence of the cutoff leads to superpartner masses. For instance, consider the gauge kinetic term,

$$\int d^2 \theta \frac{1}{g^2(\mu)} W_\alpha W^\alpha,$$  \hspace{1cm} (2.1)
where the holomorphic gauge coupling, \(g_h^2(\mu)\), runs only at 1-loop, i.e.

\[
\frac{1}{g_h^2(\mu)} = \frac{1}{g_h^2(\Lambda)} + b_0 \log \left( \frac{\mu}{\Lambda \Phi} \right). \tag{2.2}
\]

The cutoff scale comes with powers of the conformal compensator, expanding the logarithm leads to a mass for the gaugino,

\[
m_{\lambda_i} = \frac{\beta(g_i)}{g_i} m_{3/2}. \tag{2.3}
\]

A similar calculation goes through for the wavefunction renormalization of the chiral fields leading to soft masses and A-terms of,

\[
m_i^2 = -\frac{1}{4} \dot{\gamma} m_{3/2}^2 \quad A_{ijk} = -\frac{1}{2} (\gamma_i + \gamma_j + \gamma_k) m_{3/2} \tag{2.4}
\]

where the dot corresponds to \(d/dt\) and \(t \equiv \log \mu\) identifies the renormalization scale. Here we see the UV insensitivity, these results are true at all scales – in order to calculate the masses in the IR we need only know the value of couplings at that scale. Unfortunately this predicts negative squared masses for any fields charged under non-asymptotically free gauge groups. In particular the sleptons are tachyonic. This is one of the shortcomings of AMSB.

SSSB requires an auxiliary \(U(1)\)' gauge symmetry which is broken at a high scale where the auxiliary component of the vector superfield gets a D-term vev. Including the singlet in the MSSM makes it is possible to write down supersymmetry-breaking operators,

\[
\lambda_0 \int d^2 \theta \sqrt{2} \frac{W'_\alpha W'^\alpha S}{M} \quad \lambda_1 \int d^2 \theta \frac{W'_\alpha W'^\alpha S^2}{M^2}. \tag{2.5}
\]

The factor of \(\sqrt{2}\) is to simplify normalisations below. The first type of operator marries gauginos with adjoint fermions giving Dirac masses. The Dirac bino mass generated by (2.5) is,

\[
m_1 = \frac{\lambda_0 D'}{M} \equiv \lambda_0 m_D \tag{2.6}
\]

As well as generating Dirac gaugino masses the first operator gives a gaugino scale mass to the real part of the adjoint scalar leaving the pseudoscalar massless\(^1\). It also generates new scalar trilinear vertices involving

\(^1\)Here we take \(\lambda_0 \) real for simplicity. Later we will allow for arbitrary phases when we discuss CP violation in Section 3.1.
the singlet scalar and charged MSSM fields. As a result, in addition to the usual one loop corrections to scalar masses coming from gauge interactions there are now new diagrams involving the singlet scalar. These diagrams cut off the loop integrals above the gaugino mass, resulting in scalar masses that are finite and thus UV insensitive. Scalar masses are only generated once the gauginos have been integrated out. A simple way to see this is to note that there are no counterterms allowed by the symmetries. The lowest dimension operator which produces a scalar mass for MSSM fields would be

\[
\int d^4 \theta W' W' W' Q Q^\dagger \frac{1}{M^6},
\]

which is proportional to \( D'^4 \) and is not a counterterm for divergent contributions which would be proportional to \( D'^2 \). The supersoft scalar mass for a scalar of \( U(1)_Y \) charge \( Q \) is given by,

\[
m_{ss}^2 = Q^2 \frac{\alpha_Y}{\pi} m_1^2 \log 4(1 + \frac{\lambda_1}{\lambda_0^2}).
\]

The second operator in (2.5) splits the real and imaginary parts of the adjoint scalar, giving one a positive mass squared and the other a negative mass squared. Depending on the relative sign and size of these two operators one component of the scalar singlet may acquire a vev, this will be discussed further when we discuss the breaking of electroweak symmetry.

We have described how both supersoft and AM are, by themselves, UV insensitive. There is still the question of whether in combination they remain UV insensitive. The gaugino now has both a Majorana mass term generated by anomaly mediation and a Dirac mass term from supersoft. Above the Dirac mass the only contribution to the running of scalar masses is from the gaugino’s Majorana mass and this is precisely the contribution necessary to keep the scalars running on the anomaly mediated trajectory. The Dirac and Majorana mass log divergences at all loop orders renormalize independently due to an R-symmetry. Even (supersymmetric) thresholds to leading order in \( D'/M \) maintain the structure of the soft terms in the infrared, as long as there are no particles which are charged under both \( U(1) \)'s. This is clear for scalars as the lowest order operator which can be written is that in (2.7).

Thus, above the Dirac gaugino mass the two SUSY breaking mechanisms are decoupled and the low energy physics is insensitive to all physics above the Dirac mass scale\(^2\). At the Dirac mass scale the bino is integrated out

\(^2\)This was shown explicitly at the one and two loop level in [26].
and generates a finite positive mass squared for all scalars charged under hypercharge. At this scale ($O(10 \text{ TeV})$) the running is pushed off the anomaly mediated trajectory. The combination of supersoft and anomaly mediation is as insensitive as supersoft alone and the running above the Dirac mass scale is purely that of AMSB.

As far as IR phenomenology is concerned there are two SUSY breaking mass scales that determine superpartner masses at the Dirac mass scale, the AM scale ($m_{3/2}$) and the supersoft scale ($m_D$). We will parametrize these as $m_{3/2}$ and their ratio, $r = m_D/m_{3/2}$. We will show later how these two scales can be related through dynamics, making $r$ a derived quantity. When we give examples of superpartner spectra we will ignore the small effects of running from the bino mass scale.

3 The NMSSM: model and spectra

As discussed in Section 2 if there are any explicit mass scales in the superpotential they appear with powers of the conformal compensator upon rescaling of fields. Therefore in the MSSM with a $\mu$ term, a $B\mu$ term will be generated of size $\mu m_{3/2}$, much larger than the soft Higgs mass squared which is $O((m_{3/2}^2/16\pi)^2)$. Instead we combine the NMSSM and AMSB thus removing all renormalizable operators involving explicit mass parameters. The relevant piece of the superpotential (including supersoft operators) is,

$$W = \lambda S H_u H_d + \frac{\kappa}{3} S^3 + \frac{\lambda_1}{M^2} W' W' S^2 + \frac{\lambda_0}{M} W' WS + y_t Q H_u \bar{U} + \text{h.c.}, \quad (3.1)$$

where we ignore the effects of all Yukawas other than the top Yukawa, a good approximation at low $\tan \beta$. The resulting potential contains both superpotential terms and soft terms, $V = V_{\text{susy}} + V_{\text{soft}}$. With,

$$V_{\text{susy}} = \frac{1}{8} \left( g_Y (|H_u|^2 - |H_d|^2) + 4 \lambda_0 m_D Re S \right)^2 + \frac{g_2^2}{8} \left( |H_u|^2 - |H_d|^2 \right)^2 + \frac{g_2^2}{2} \left| H_u^0 H_d^{0*} + H_u^0 H_d^{-*} \right|^2 \quad (3.2)$$
and

\[ V_{\text{soft}} = m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + m_S^2 |S|^2 + \lambda_S A \lambda (S H_u H_d + h.c.) + \frac{\kappa}{3} A \kappa (S^3 + h.c.). \]  

(3.3)

These soft parameters get contributions from both anomaly mediation and supersoft, in the case of Higgs soft masses, and purely anomaly mediation, for A-terms and singlet masses. The AMSB contributions are given by (2.4) specialised to the case for our superpotential, see Appendix A.

In Section 3 we will evaluate the soft masses in the IR and give examples of UV parameters that result in a viable spectrum. We can limit the parameter space we need to analyse by requiring that we get correct electroweak symmetry breaking.

We now calculate the superpartner spectrum in the model. The effect of the supersoft operators (2.5) is to give the bino a large Dirac mass and to add a positive, finite contribution to the squared mass of all fields, proportional to their hypercharge squared. There are also positive supersoft contributions to the Higgses and singlet scalar masses through their superpotential couplings.

In pure AMSB the superpartner masses can be calculated using (2.3) and (2.4) and the formulae in Appendix A. The one-loop supersoft corrections are given by (2.6) and (2.8), and for simplicity we normalise \( \lambda_0 \) to 1. The relative size of the supersoft and the AMSB contributions is set by the ratio of \( r \equiv m_D/m_{3/2} \). In Section 4 we will demonstrate how it is possible to dynamically generate \( r \lesssim 1/2 \) but for now we leave \( r \) as a free variable.

To see how the addition of supersoft fixes the tachyonic slepton problem we list the masses of the gauginos, the sleptons, the down type squarks from all generations and the up type squarks from the first two generations. The soft masses of these fields depend only upon gauge couplings whereas the Higgs scalar and third generation squark masses depend on the particular choice of \( y_t, \kappa, \) and \( \lambda_S \).

The winos, gluinos and bino/singletino system have masses

\[
M_1 = \begin{pmatrix} 1.40 M & 158 r M \\ 158 r M & \kappa |\langle S \rangle| \end{pmatrix} \\
M_2 = 0.427 M \\
M_3 = -4.32 M,
\]

(3.4) (3.5) (3.6)

\(^3\text{Note } H_u H_d = H_u^+ H_d^- - H_u^0 H_d^0.\)
where $M = m_3/\sqrt{16\pi^2}$. The scalar masses that are independent of the top Yukawa are

$$m^2_{L_{1,2,3}} = (-0.363 + 28.1 r^2) M^2 \quad (3.7)$$
$$m^2_{\tilde{e}_{1,2,3}} = (-0.358 + 112 r^2) M^2 \quad (3.8)$$
$$m^2_{\tilde{Q}_{1,2}} = (16.3 + 3.12 r^2) M^2 \quad (3.9)$$
$$m^2_{\tilde{u}_{1,2}} = (16.4 + 49.9 r^2) M^2 \quad (3.10)$$
$$m^2_{\tilde{d}_{1,2,3}} = (16.5 + 12.5 r^2) M^2. \quad (3.11)$$

It is clear that a large enough $r$ will fix the tachyon problem. From (2.8), we see that the exact size of the supersoft correction depends not only upon $r$ but also $\lambda_1$. However, as we will see later at viable points $\lambda_1 \ll 1$ so the log in (2.8) is well approximated by log 4, we have used this approximation when calculating the masses above.

In (3.6) we have included the Majorana mass for the singletino that arises, when $S$ gets a vev, from the cubic term in the superpotential. We ignored this effect when calculating the supersoft contribution (2.8), the correction is small—of order $\kappa^2 S^2/m^2_D$. In addition, there is a small Majorana mass for the gaugino coming from AMSB. Both Majorana masses are much smaller than the Dirac mass and we ignore the small splitting in the bino-singletino system, quoting just the Dirac mass when we give spectra.

All that remains is to find appropriate values for the parameters $\lambda_S, \lambda_1, \kappa, r$ and the one mass scale $M$ that lead to superpartner masses consistent with present experimental constraints and viable EWSB.

### 3.1 EWSB

The Higgs-singlet potential that must break electroweak symmetry is given by (3.2) and (3.3). The singlet vev is controlled by the $\lambda_1$ coupling (having normalised $\lambda_0$ to 1). We now describe what happens as $\lambda_1$ is varied.

We first consider $\lambda_1 < 0$. For $-3/4 < \lambda_1 < 0$ the vev of $S$ is tiny, consequently so is the $\mu$-term, and the supersoft corrections to the slepton masses are positive, (2.5). At $\lambda = -3/4$ the supersoft contributions turn off. For $-1 < \lambda_1 < -3/4$ the supersoft contributions change sign since the scalar $S$ is now heavier than the bino. As we approach $\lambda_1 = -1$ the mass of the real part of $S$ goes to zero and the slepton mass contributions go to $-\infty$, an IR divergence. Lowering $\lambda_1$ further increases the supersoft mass from $-\infty$, 


to 0 at $\lambda_1 = -5/4$, at the same time the mass for the real part of $S$ decreases from 0 to $-m_D^2$, leading to a large real $S$ vev. Below $\lambda = -5/4$ the supersoft contributions are again positive but the vev of the real part of $S$ is greater than $m_D$. However, the appearance of $S$ in the hypercharge $D$-term forces the Higgses to acquire unacceptably large vevs. We see that the whole region of $\lambda_1 < 0$ is ruled out, leaving only positive $\lambda_1$.

The region $\lambda_1 > 0$ gives reasonable supersoft slepton masses and a negative squared mass for the imaginary part of $S$, of magnitude $\lambda_1 m_D^2$. The vev of $S$ determines the $\mu$ and $B\mu$ terms and so we take $\lambda_1$ small. However, it can not be taken arbitrarily small due to the fact that we now have a CP violating vev and too small a $\lambda_1$ leads to a large contribution to EDMs, discussed below. From now on we consider $\lambda_1$ small and positive.

While the Higgs vev of 175 GeV and a reasonable $\mu$ term are relatively easy to achieve by the correct choice of the overall scale of the potential, one issue that persists over a very large range of parameters is that tan $\beta$ is low ($\sim 1$) and $\text{Im}(H_d) \approx \text{Re}(H_d)$ since the potential is relatively symmetric with respect to both the real and imaginary parts of the field. The value tan $\beta \sim 1$ is not a problem for the physical Higgs mass, but requires a too-large top Yukawa coupling if one wants to avoid Landau poles below the GUT scale.

However, there is one more operator that can be added to the Higgs potential,

$$\frac{\lambda_c}{M^2} W'W'H_uH_d$$

When the $W'$ gets a vev, this is a $B\mu$ term for the Higgs which is $\lambda_c m_D^2$. This term also splits the real and imaginary parts of the vev of $H_d$, if we allow the coupling constant $\lambda_c$ to be complex.

There is a portion of parameter space in which $\mu$ is of moderate size ($\sim 100$ GeV) and the large negative contribution to the (up-type) Higgs soft mass is partly canceled by a large positive contribution from the finite bino loop (i.e., the supersoft contribution), and the new $B\mu$ term. In this range (see Table 1 below) it is possible to get a reasonably-sized tan $\beta$ and a viable spectrum. Even the (standard-model-like) Higgs mass is typically well above the LEP bound.

There are several general features worthy of note in Table 1. The mass of the lightest Higgs quoted in the table includes the dominant loop corrections and has been calculated \footnote{We thank Carlos Wagner for assistance with this.} for a top quark pole mass of 178 GeV but it is un-
## Table 1: Typical solutions, all masses in GeV. Here $m_{\tilde{B}}$ denotes to the Dirac mass between the bino and the singletino.
doubtedly subject to further corrections. There is spontaneous CP violation so the Higgs states can not be split into a CP-even and a CP-odd sector – they mix. The singlet state is heavy and the low energy Higgs phenomenology has similar aspects to that of the MSSM with explicit CP violation [27]. This CP violation predicts electric dipole moments not far beyond present experimental bounds. The LSP is a neutralino, a linear combination of a neutral Higgsino and Wino.

For our model, $y_t$ is typically the most sensitive parameter. Normally we would define the sensitivity as $\partial \log v / \partial \log y_t$. However, our points lie on the edge of parameter space; while decreasing $y_t$ produces only a moderate change in the Higgs vev, increasing $y_t$ by a minute amount causes a very large change. Instead we define a parameter comparing the region of parameter space over which $y_t$ assumes it’s natural value, where the Higgs vev wants to live around the scale $m_{3/2}$, to the region on the edge of parameter space where the Higgs vev is around $5 \times 175 \text{GeV}$. For our points, this naturalness parameter is of order $10^{-3}$.

We have attempted to use the singlet already present in the NMSSM and an integral part of EWSB to also solve the tachyon problem of AMSB. This is the minimal model and it turns out to be tightly constrained, as can be seen from the fraction of a percent the naturalness parameter assumes. We will discuss in the conclusion possible ways to alleviate this tuning.

4 CP Violation

In order to calculate the amount of CP violation predicted by the model we must first identify all the CP violating phases. In principle many parameters in the model could contain a phase but because the soft parameters are generated through anomaly mediation many of these phases are related. In particular, the Majorana gaugino mass ($m_i$) and all the A-terms ($A_h$, $A_\lambda$ and $A_\kappa$) are proportional to $m_{3/2}$ with no new phase entering through the running. In addition to the phases present in parameters there is also spontaneous CP violation coming from the vevs for $S, H_u$ and $H_d$.

The simplest way to identify the physical CP violating phases is to notice that the superpotential and soft parameters break three $U(1)$ symmetries.

\footnote{More precisely we find the region of parameter space where the Higgs vev varies by at most 100\% from the Standard Model value.}
By allowing both fields and parameters to transform we can restore the symmetries. The charge assignments are given in Table 2. The physical CP violating phases will be invariant under the spurious symmetries.

The superpotential generically has five couplings with phases \( \lambda_0, \lambda_1, \lambda_c, \lambda_S, \kappa \). (4.1)

In addition, there are four phases that come from the soft SUSY breaking terms,

\[ m_i, A_y, A_\lambda, A_\kappa \],

all of which have the same phase, as was explained above. Phases can also come from the vevs of \( S, H_u, \) and \( H_d \) but one combination of these phases can be removed by an \( SU(2) \) gauge transformation. We have a total of eight phases, however due to the symmetries only five are physical CP violating phases.

There are three independent, CP violating combinations of parameters allowed by the charge assignments in Table 2

\[ \lambda_0^2 \lambda_1^*, \lambda_0^3 A_y^* \kappa^*, \lambda_0 \lambda_S^* A_i^* \lambda_c, \]

where \( A_i \) can be any of \( A_\kappa, A_\lambda, A_y \) and \( m_i \). Including vevs of fields, we can write down two more combinations which are both invariant under the

\[ \text{Table 2: } U(1) \text{ charges of parameters and fields.} \]

| \( U(1)_P \) | \( U(1)_R' \) | \( U(1)_S \) |
|---|---|---|
| \( \lambda_0 \) | 0 | 0 | -1 |
| \( \lambda_1 \) | 0 | 0 | -2 |
| \( \lambda_S \) | -2 | 2 | -1 |
| \( \lambda_c \) | -2 | 0 | 0 |
| \( \kappa \) | 0 | 2 | -3 |
| \( m_i, A_y, \lambda_\kappa \) | 0 | -2 | 0 |
| \( H_u, d \) | 1 | 0 | 0 |
| \( S \) | 0 | 0 | 1 |

\[ \text{Note that we are only interested in phases beyond those already present in the SM, so we ignore the QCD vacuum angle as well as the Yukawas and their associated CKM phase.} \]
\[ U(1)'s \text{ above and also gauge invariant,} \]
\[ \lambda_0 \langle S \rangle \text{ and } \lambda_S A_i \langle S \rangle \langle H_u H_d \rangle. \quad (4.4) \]

Thus any CP violating physical amplitude must be a function only of the combinations of parameters and vevs given in (4.3) and (4.4). At the one-loop level we will see this explicitly.

The strongest constraints on new CP violating phases come from attempts to measure electric dipole moments (EDMs) of the electron and the neutron. The effective dipole operator coupling fermions to the photon is,
\[ -\frac{1}{2} D_\psi \bar{\psi} \sigma^{\mu\nu} \psi \gamma_5 F_{\mu\nu} + h.c. \quad (4.5) \]

The coefficient of the operator, \( D_\psi \), is related to the electric dipole moment, \( d_\psi \) as \([28, 29]\),
\[ d_\psi = |D_\psi| \sin \phi \quad \text{with} \quad \phi = \text{Arg}(m_\psi^* D_\psi). \quad (4.6) \]

For the electron the present bound on its EDM is derived from the measurement of the EDM of \(^{205}\text{Tl}\) and is \([30]\) \( |d_e| \leq 1.6 \times 10^{-27} \text{ cm} \), the corresponding limit for the neutron \([31]\) is \( |d_n| \leq 6.3 \times 10^{-26} \text{ cm} \) both with 90% confidence.

First the case of electron EDMs. On the face of it this operator appears to be dimension 5 but it is chirality violating and so is proportional to the mass of the fermion, making it effectively dimension 6. At leading order in the electron Yukawa one-loop diagrams involving superpartners contribute to the EDM, for a review of the calculation in the MSSM see, for example, \([28]\) and \([29]\) and references therein. The calculation in our case is simplified because diagrams involving the bino are now suppressed since the Dirac mass of the bino is order many TeV. In the NMSSM there is an additional contribution but it again involves exchange of a bino and can be ignored. Using the physical points from Table 1 as a guide to the form of the spectrum we find,
\[ D_e \sim \frac{g_2^2}{16\pi^2} m_e \left( \frac{2}{|m_l|^2} + \frac{v_u^* m_H^* m_W^*}{v_d |m_l|^4} \right), \quad (4.7) \]

and
\[ \phi = \text{Arg} \left( 2 + \frac{v_u^* m_H^* m_W^*}{v_d |m_l|^2} \right). \quad (4.8) \]

As promised, this phase is only a function of those combinations of parameters and vevs (4.3), (4.4) allowed by the spurious symmetries. We list the
resulting electron EDMs for each point at the bottom of Table 1; they are close to but below the experimental bound. There are similar diagrams that contribute to quark EDMs which in turn lead to a neutron EDM. As for the leptons we are close to but below the neutron EDM bound.

In addition, the neutron EDM receives a contribution from the dimension 6 operator first discussed by Weinberg [32],

\[-\frac{1}{6} C f_{abc} G_{\alpha\mu}^a G_{\beta\mu}^b G_{\gamma\delta}^c \epsilon^{\alpha\beta\gamma\delta}, \tag{4.9}\]

where $f_{abc}$ are the structure constants. This operator is not proportional to a light quark mass and so is the dominant contribution to the neutron EDM, it is not present for the electron as it requires the fermion to be charged under a non-abelian gauge group. The largest contribution to the coefficient $C$ comes from a two loop diagram containing 3 external gluons with the quarks in the loop exchanging a Higgs boson. For this contribution to be non-zero the Higgs sector must have mixing between at least three scalar fields e.g. the NMSSM.

Using naive dimensional analysis the effects of (4.9) on the neutron EDM operator of (4.5) can be estimated. The dominant contribution requires two mass insertions on the Higgs propagator. If the scalar states inserted are too light, the contribution to the neutron EDM could be huge. Fortunately we see from electroweak symmetry breaking that the singlet mass is multi TeV and the coupling $\lambda_S$ is small, of order $10^{-2}$, and the contribution to the neutron EDM (see Table 1) is below, but relatively close to the experimental bound.

5 UV Models

We now discuss several natural UV realisations of our model that result in the interesting IR physics discussed above.

5.1 Symmetries, Forbidden Operators and GUTs

Before we discuss the mediation of supersymmetry breaking we comment on the remaining relevant operators that can affect weak scale physics.

The most dangerous operator\(^7\) which could be introduced is kinetic mixing

\(^7\)We thank Hitoshi Murayama for reminding us of this fact.
between the hidden sector $U(1)$ and hypercharge, \textit{i.e.},

$$\int d^2 \theta W_\alpha^0 W'_\alpha. \quad (5.1)$$

If this term is present the D-term for the hidden sector $U(1)$ would be a tadpole for the hypercharge D-term, destabilising the gauge hierarchy. While there is no symmetry that could forbid this operator as the theory stands, it is technically natural not to include this term at the cutoff as it won’t be generated through loops since there is no matter in the theory that is charged under both groups.

Ultimately one would like to embed the whole model in a GUT, such as $SU(5)$ [33], and doing so may allow us to suppress or forbid entirely the kinetic mixing due to symmetries. The GUT would be broken down to $SU(3) \times SU(2) \times U(1)$ by an adjoint chiral field, $\Sigma$, acquiring a GUT scale vev. In addition to using this field to break the GUT group we use it to pick out the $U(1)_Y$ direction so that the only supersoft operator is for $U(1)_Y$. At the GUT scale the supersoft operator of (2.5) becomes,

$$\int d^2 \theta \sqrt{2} W'_\alpha W^\alpha \Sigma S \Sigma M^2. \quad (5.2)$$

The low energy supersoft operator would then be suppressed by a factor of $\langle \Sigma \rangle / M$ which could be $\sim 10^{-2}$ or so, but of course depends on the physics at the GUT/string/Planck scales.

In the context of a GUT, one could potentially use a symmetry to forbid the kinetic mixing. One possibility is to impose a $U(1)$ R-symmetry which is spontaneously broken by the GUT. The singlet and all the MSSM fields would have charge $2/3$ under the R-symmetry while $\Sigma$’s charge would be $-2/3$. This forbids the kinetic mixing term even below the GUT scale, whilst allowing all the previous operators discussed once appropriate powers of $\Sigma/M$ are included. However, it is an open question how to arrange for a field carrying non-zero $U(1)_R$ charge to have a GUT breaking (super)potential. More importantly, there will be a modulus produced, the superpartner of the R-axion, that could cause the GUT scale to act effectively as a non-supersymmetric threshold and take us off the AMSB trajectory [4, 8, 34]. Instead of a continuous symmetry one could charge fields under a discrete R-symmetry by adding terms in the GUT theory which explicitly break the $U(1)_R$ but preserve the discrete subgroup. As an example, $Z_{12}$ with charges -6 for $d^2 \theta$, +2 for $H_u$, $H_d$ and $S$, +3 for $W_\alpha$ and $W'_\alpha$, and -2 for $\Sigma$. The kinetic mixing term would thus be suppressed by a factor of $(\langle \Sigma \rangle / M)^6$. 

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5.2 Supersymmetry Breaking: The 4 − 1 Model

What we want is a situation where the $D$-term breaking in the $U(1)$ is comparable to the overall scale of supersymmetry breaking in the hidden sector. This can happen in any supersymmetry-breaking sector in which the $U(1)$ is required for supersymmetry breaking, i.e., if the $U(1)$ gauge coupling is turned off, a supersymmetry-preserving minimum is restored. While there are other examples in the literature of this kind, we will describe a particularly simple one [35].

The model has an $SU(4) \times U(1)$ gauge group. The matter content consists of the following $SU(4)$ representations: an antisymmetric tensor $A_2$, a fundamental $F_{−3}$, an anti-fundamental $\bar{F}_{−1}$ and a singlet $S_4$. The subscripts are the charges under the $U(1)$. The only allowed superpotential term is

$$W = \lambda S_4 F_{−3} \bar{F}_{−1}$$

However $SU(4)$ will confine and the gauginos will condense leading to a non-perturbatively generated superpotential. The complete superpotential becomes,

$$W = \lambda S_4 F_{−3} \bar{F}_{−1} + \frac{\Lambda^5_4}{(\bar{F}_i F_j A^{ik} A^{lm} \epsilon_{jklm})^{1/2}}.$$  

We consider a regime where $g_4 \gg \lambda \gg g_1$. The minimum of the potential will occur along the $SU(4)$ D-flat direction. We can make a gauge rotation so that the vevs along the $SU(4)$ D-flat directions have the form,

$$A_2 = \begin{pmatrix} a \sigma_2 \\ a \sigma_2 \end{pmatrix}, \quad F = \bar{F} = \begin{pmatrix} b \\ 0 \\ 0 \\ 0 \end{pmatrix}, \quad S = c.$$  

For convenience we rescale the fields, $\phi \rightarrow \frac{A}{\Lambda^{1/5}} \phi$. In this paramatrization the $U(1)$ D-term is,

$$D_1 = g_1 \frac{\Lambda^2}{\lambda^{2/5}} (2|a|^2 - 4|b|^2 + 4|c|^2)$$

whilst the contribution to the potential is $V_D = \frac{D^2}{2}$ and the F-term contribution to the potential is,

$$V_F = \lambda^{6/5} A^4 \left( |b|^4 + \left| 2bc - \frac{1}{a b^2} \right|^2 + \left| \frac{1}{a^2 b} \right|^2 \right).$$
The minimum of the potential has non-zero D-term. If the \( U(1) \) gauge coupling is turned off it is possible to satisfy the F-term constraint, \( b \to 0 \) while \( a \) and \( c \) scale like inverse powers of \( b \) and SUSY is unbroken. If we turn on the gauge coupling the minimum of the potential moves in from infinity. In order to get a non-zero D-term it is crucial that the gauge group be involved in the SUSY breaking. Numerically we find that for \( \lambda = 10 g_1 \) \( V_D \approx 0.5 V_F \). Thus, we have a non-perturbative mechanism whereby the D-term and F-term contributions to the vacuum energy are comparable.

The model above is not unique. There are other supersymmetry breaking models where a gauged \( U(1) \) plays an important role in the dynamics, see for example [35].

### 5.3 D-term transmission

Now we describe how to transmit supersymmetry breaking in the form of anomaly mediation plus a dirac bino mass while sequestering all other contributions. This can be done via a fifth dimension [1] or via conformal sequestering [24, 25, 36] with the hidden sector \( U(1) \) in the “bulk”.

#### 5.3.1 Flat Extra Dimension

We place the \( SU(4) \) gauge fields and all of the matter of the 4–1 model on the hidden brane while the MSSM matter and gauge fields are restricted to lie on a separate brane at the other orbifold fixed point. The \( U(1) \) gauge field of the 4–1 model propagates in the bulk. The dynamics discussed in the previous section generate a non-zero D-term on the hidden brane, \( D_b \).

In the bulk the vector multiplet is part of an \( N = 2 \) gauge multiplet that decomposes in \( N = 1 \) language as a vector multiplet and a chiral multiplet, also in the adjoint of the gauge group [37]. In Wess Zumino gauge,

\[
V' = -\theta \sigma^\mu \bar{\theta} A'_\mu + i \bar{\theta}^2 \theta \lambda_1 - i \bar{\theta}^2 \bar{\theta} \lambda_1 + \frac{1}{2} \bar{\theta}^2 \theta^2 D'
\]

\[
\Phi = \frac{1}{\sqrt{2}} (\Sigma + i A_5) + \sqrt{2} \bar{\theta} \lambda_2 + \theta^2 F.
\]

We take the vector superfield to be even under the orbifold boundary conditions and the chiral superfield to be odd. The relevant part of the action
is,
\[
\int d^4xdy \left[ \int d^2\theta \frac{1}{4} W^{\alpha} W_\alpha' \right. + h.c. + \int d^4\theta (\partial_5 V' - \frac{1}{\sqrt{2}} (\Phi + \Phi^\dagger))^2 \\
+ \delta(y - \pi R) \int d^4\theta X^\dagger e^{\phi_5 V'} X + \delta(y) \int d^2\theta \frac{W^{\alpha} W_{Y\alpha} S}{M^{3/2}} \right].
\] (5.8)

This leads to equations of motion
\[
D' + \partial_5 \Sigma + \delta(y - \pi R) \sqrt{2\pi R} D_b + \delta(y) \frac{D_Y S}{M^{3/2}} = 0 \quad \text{and} \quad \frac{1}{2} \eta^{\mu\nu} \partial_\mu \partial_\nu \Sigma - \partial_5 D' = 0.
\] (5.9)

In the bulk the zero mode of the scalar, \( \Sigma \), has constant slope. Without the boundary source terms the zero mode of \( \Sigma \) would be projected out and \( \Sigma \) would be zero across the whole space. Here, however, there are non-zero source terms at the branes causing \( \Sigma \) to jump at the branes. Taking \( D_Y = 0 \) there is only a jump at \( y = \pi R \), this causes \( \Sigma \) to have a gradient across the whole space. Although its value at \( y = 0 \) is unaltered from the case with no sources it now has gradient there, leading to a non-zero D-term at \( y = 0 \).

Using the periodicity of \( \Sigma \) we find,
\[
D' = -\frac{1}{\sqrt{2\pi R}} D_b
\] (5.10)

and
\[
\Sigma = \frac{D_b}{\sqrt{2\pi R}} (y - 2\pi R \theta(y - \pi R)),
\] (5.11)

where \( \theta(y) \) is the Heaviside function.

A concern one should have raised by now is the fact that there are light fields in the bulk (the \( U(1) \) gauge multiplet) which could now in principle communicate supersymmetry breaking directly to the visible sector by generating scalar masses through higher dimensional operators involving bulk fields. The gauge invariant combination of bulk fields which could couple to the visible sector is \( (\partial_5 V' - \frac{1}{\sqrt{2}} (\Phi + \Phi^\dagger)) \). This combination has odd boundary conditions, but could couple through a partial derivative, as
\[
\delta(y) \int d^4\theta \frac{\partial_5 (\Phi + \Phi^\dagger - \sqrt{2} \partial_5 V')}{M^{5/2}} Q^\dagger Q.
\] (5.12)

A scalar mass for \( Q \) would come from the \( D \)-component of the coefficient and thus proportional to \( \partial_5^2 D' \). However, \( D' \) is constant in the bulk and through
$y = 0$, so this potentially dangerous contribution vanishes *dynamically*. In fact all such contributions cancel because gauge invariance requires the gauge field to appear as $\partial_5 V'$.

So the D-term on the hidden brane is transmitted through the bulk to our brane. Any F-term generated on the hidden brane is still sequestered from the MSSM, forbidding SUSY breaking contact interactions.

### 5.3.2 Warped Extra Dimension

We now wish to generalise the discussion of Section 5.3 to a warped space. In particular we consider a slice of five dimensional AdS space [38] whose fifth dimension is compactified on an orbifold $S^1/Z_2$, thus $0 \leq y \leq \pi R$. The orbifold has 2 fixed points $y = 0$ and $y = \pi R$ at which we place 3 branes. The metric on the space is given by,

$$ds^2 = g_{\mu\nu} dx^\mu dx^\nu + dy^2 = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2,$$  \hspace{1cm} (5.13)

where $\mu, \nu$ run over our 4 spacetime dimensions and $\sigma(y) = k|y|$ where $k^{-1}$ is the AdS curvature length.

The equations of motion now contain factors of the metric and are,

$$\sqrt{-g} D' + \partial_5 (\sqrt{-g} \Sigma) + \sqrt{-g} \delta(y - \pi R) \sqrt{2\pi R} D_{\pi R} = 0$$

$$\partial_5 D' = 0. \hspace{1cm} (5.15)$$

As before these can be solved by using the periodicity of $\Sigma$. Due to the factors of the metric $\Sigma$ no longer has constant slope in the bulk, instead it

\[ \text{Figure 1: The profile of } \Sigma \text{ is shown for a flat extra dimension on the left and for the warped case on the right.} \]
has an exponential profile (see Figure 1),

\[ D' = -\frac{4k\sqrt{2\pi R}}{e^{8k\pi R} - 1}D_{\pi R} \] (5.16)

\[ \Sigma(y) = \frac{D'}{4k}(1 - e^{4k|y|})\frac{y}{|y|} \] (5.17)

Note that just as in the flat case, the dangerous operator (5.12) which would transmit soft masses to scalars dynamically vanishes by virtue of the equations of motion, specifically (5.15).

Using the inspiration of the AdS/CFT correspondance one can attempt to construct models in which the hidden sector are the IR dynamics of a CFT and the visible and hidden sectors are conformally sequestered. The \( U(1) \) in the bulk would then correspond to an exact global symmetry of the conformal symmetry which is weakly gauged.

## 6 Conclusions

For AMSB to be the dominant source of SUSY breaking the Kahler potential has to take on a “sequestered” form. This can be achieved by separating the SUSY breaking dynamics from the SM, either geometrically or through large anomalous dimensions. We presented explicit models where the sequestered SUSY breaking dynamics involve a gauged \( U(1) \) leading to a supersoft operator for hypercharge. We demonstrated how through this operator the singlet introduced in the NMSSM to solve the \( \mu \) problem may be used to solve the tachyon problem. These two competing contributions to SUSY breaking are the same size not through choice but dynamics. In the visible sector the field content is simply that of the NMSSM. The spectrum is similar to that of AMSB with an additional positive shift of all the scalar masses proportional to the square of their hypercharge. The theory remains UV insensitive and the bino has a large Dirac mass.

The supersoft operator treats the real and imaginary parts of the singlet differently and in our minimal model this leads to spontaneous CP violation in the Higgs sector. We found that the predicted size of the electron EDM is right at the present day constraints. Along with a Dirac bino this is one of the most interesting signatures of our model. The significant CP violation in the Higgs sector will cause significant deviations from conventional Higgs phenomenology and warrants further investigation.
The minimal model we present is very tightly constrained since the same field that is involved in EWSB is also involved in fixing the tachyon problem of AMSB. This results in the amount of viable parameter space being quite small and it appears to be tuned below the percent level. However, there are generalisations of the model with less minimal field content where the two effects are separated, we intend to investigate these further elsewhere [39]. For instance\(^8\), by adding an additional singlet and gauging B-L the supersoft operator could again be used to lift the slepton squared masses while the \(\mu\) problem is solved in an independent sector.

Acknowledgements

We thank Hitoshi Murayama, Ann Nelson, Raman Sundrum and Neal Weiner for conversations. P.J.F. would like to thank the JHU theory group for hospitality while this work was initiated. The work of P.J.F. is supported in part by the U.S. Department of Energy. The work of D.E.K. is supported by the National Science Foundation and the Department of Energy’s Outstanding Junior Investigator Program.

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\(^8\)We thank Neal Weiner for discussions on this issue.

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**APPENDIX**

Here we list the anomalous dimensions and beta functions relevant for calculating anomaly mediated soft masses in the NMSSM.
These set of equations completely determine the one-loop anomaly mediated contributions to the soft parameters.