When Anomaly Mediation is UV Sensitive

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Abstract

Despite its successes—such as solving the supersymmetric flavor problem—anomaly mediated supersymmetry breaking is untenable because of its prediction of tachyonic sleptons. An appealing solution to this problem was proposed by Pomarol and Rattazzi where a threshold controlled by a light field deflects the anomaly mediated supersymmetry breaking trajectory, thus evading tachyonic sleptons. In this paper we examine an alternate class of deflection models where the non-supersymmetric threshold is accompanied by a heavy, instead of light, singlet. The low energy form of this model is the so-called extended anomaly mediation proposed by Nelson and Weiner, but with potential for a much higher deflection threshold. The existence of this high deflection threshold implies that the space of anomaly mediated supersymmetry breaking deflecting models is larger than previously thought.

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I. INTRODUCTION

A puzzling aspect of the standard model (SM) is the stability of the Higgs mass at the electroweak scale—a mystery which is elegantly resolved by supersymmetry (SUSY). Of course a realistic theory requires SUSY to be broken and constructing viable models for SUSY breaking is not a trivial issue. The challenge, known as the SUSY flavor problem, is to find a scenario which inherently suppresses the contributions to flavor and $CP$ violating operators—while hopefully also significantly reducing the number of SUSY breaking parameters.

One such scenario, particularly successful in terms of the SUSY flavor problem, is anomaly mediated supersymmetry breaking (AMSB)\cite{1,2} (for a review see \cite{3,4}). To realize AMSB requires a true sequestering of the hidden and visible sectors so that only gravitational interactions exist between the two. This was originally achieved via a specific geometry in a model with extra dimensions \cite{1} and then later shown to be possible in four dimensions with a specific type of a superconformally invariant SUSY breaking sector \cite{5,6}. The SUSY breaking is thus intimately linked with the breaking of superconformal invariance, broken at loop level. The soft SUSY breaking terms are related to a mass scale $F_\phi$ (the gravitino mass), the low energy beta-functions, and the anomalous dimensions. They are also renormalization group equation (RGE) scale invariant, ultra violet (UV) insensitive, highly predictive, and absent of new flavor violation.

Regardless of these features, AMSB is problematic since naively applying it to the minimal supersymmetric standard model (MSSM) produces tachyonic slepton masses. The lack of free parameters and the insensitivity of the soft parameters to UV physics makes this problem challenging. Still solutions exist: adding low energy Yukawa couplings to alter the beta functions and anomalous dimensions\cite{7,9}, utilizing $D$-term contributions to soft breaking masses\cite{10,12}, considering low energy threshold effects \cite{13}, and deflecting from the AMSB trajectory \cite{14} (with a variety of models along this theme \cite{15,22}).

In the deflection models, a superconformal violating threshold introduces new SUSY breaking causing a deflection from the anomaly mediated trajectory. Typically this is achieved via messengers that are charged under the SM resulting in mixed anomaly and gauge mediated supersymmetry breaking (for reviews of gauge mediated supersymmetry breaking (GMSB) see for example \cite{23,24}). The deflecting threshold in \cite{14} is achieved
through a shallow potential yielding a high messenger scale (much larger than $F_\phi$) and a light singlet (modulus); therefore, such deflection is sometimes called light singlet deflection.

Here we offer a scenario (which we shall call heavy singlet deflection) that contains a high-scale SUSY threshold which, in the SUSY limit, has no effect on the AMSB trajectories. Turning on SUSY breaking induces a second threshold which is SUSY violating and leads to deflection at that scale. The deflection terms induced are precisely of the extended anomaly mediation (EAM) form\cite{17, 18} making the theory a UV completion of EAM; however, the induced threshold need not be near $F_\phi$ (as is assumed for EAM) and indeed can generically be much higher. As such, we will demonstrate that deflection can occur at thresholds well above $F_\phi$ despite the absence of a light-singlet. The existence of this arbitrarily high deflection threshold indicates that the space of AMSB deflecting models is larger than hitherto expected and that AMSB is not as UV insensitive as commonly believed.

One of the main points of this paper is to clarify that deflection, in general, arises whenever a threshold is controlled by a vacuum expectation value (VEV) stabilized by a SUSY breaking term. These terms break the conformal invariance of AMSB and convey that information through the threshold. Ref. \cite{14} proposed a scenario where this deflection included a light singlet; however, a light singlet is not a requirement of a high-scale deflecting threshold (as the existence of heavy singlet deflection demonstrates). To demonstrate these points on deflection, we employ a simple example often quoted in the literature as a model without deflection, but which actually contains deflection of the heavy singlet type.

Following a review of AMSB and the necessary conditions for deflection in Section II, this paper presents, in Section III, an example model illustrating the ideas discussed above. The latter portion of Section III is devoted to the discussion of a model employing this mechanism to generate non-tachyonic slepton masses, which resembles the first model presented in \cite{17}.

II. AMSB AND DECOUPLING

In AMSB, only gravitational interactions link the SUSY breaking sector to visible sector and therefore gravity communicates the SUSY breaking. A convenient way to describe this transmission of SUSY breaking is to separate out an auxiliary component of the gravity supermultiplet called the conformal compensator, $\phi$. Superconformal invariance then dictates how $\phi$ appears in the lagrangian, and using canonically normalized fields one factor of $\phi$
appears per unit dimension of the coupling; that is \( M \to \phi M \).

Due to the SUSY breaking in the sequestered sector, the conformal compensator picks up a non-zero \( F \) component; this may be parameterized as

\[
\phi = 1 + F_\phi \theta^2.
\]  

With this choice AMSB predicts the form of the SUSY breaking terms:

\[
(m^2)_i^j = -\frac{1}{8}|F_\phi|^2 [\gamma_i^j, \gamma] - \frac{1}{4}|F_\phi|^2 \left[ \frac{1}{2} \partial \gamma_i^j \beta_{gG} + \frac{\partial \gamma_i^j}{\partial Y^m_{\ell m}} \beta_{Y^m_{\ell m}} + \text{h.c.} \right]
\]  

\[ a^{ijk} = F_\phi \beta_{Y^i}^{ijk} \]  

\[ M_G = -\frac{\beta_{gG}}{g_G} F_\phi, \]  

which may be derived using Eq. (1) and promoting both the wave function renormalization constant, \( Z_{ij} \), and inverse gauge coupling, \( 1/\alpha_G \), to the superfields \( Z_{ij} \), \( \tau_G \).

The AMSB expressions of Eqs. (2)–(4) are not just boundary conditions; rather, they also represent solutions to the RGEs. This unique property is a result of their being determined by superconformal invariance and their origin from quantum corrections. Since these equations solve the RGEs, they have been labeled AMSB trajectories—that is, they form a set of equations that may be evaluated at any renormalization scale.

In addition to being solutions of the RGEs, the AMSB trajectories are UV insensitive; in other words, they generically decouple heavy thresholds. One method to see why this decoupling occurs is to consider what additional SUSY breaking results from the introduction of a threshold—if it is not comparable in size to the AMSB contribution then it may be safely neglected.

To obtain what new SUSY breaking might be generated by introduction of an intermediate threshold, consider a mass scale \( M \) with \( \Lambda \gg M \gg F_\phi \) and \( \Lambda \) some cutoff for new physics. Additionally, split the set of superfields \( \{\Phi_i\} \) into two sets: \( q = \{\Phi_i | M_{\Phi_i} \sim M\} \) and \( Q = \{\Phi_i | M_{\Phi_i} \ll M\} \), so that \( q \) are heavy fields integrated out at the threshold \( M \) and \( Q \) are light fields remaining in the theory below \( M \).

\[ \frac{d \ln Z_i^j}{d \ln \mu} = \gamma_i^j \] so that, for example, a yukawa coupling in the superpotential \( W \supset \frac{1}{3!} Y^{ijk} \Phi_i \Phi_j \Phi_k \) has a beta function of \( \beta_{Y^{ijk}} = -\frac{1}{2} Y^{ijp} \gamma_p^{jk} + (i \leftrightarrow k) + (j \leftrightarrow k) \). We also choose all SUSY breaking terms to come with a plus sign in the potential.
The most generic lagrangian above $M$ is
\begin{equation}
\mathcal{L}^+ = \mathcal{L}_Q^+ + \mathcal{L}_{Qq} + \mathcal{L}_q
\end{equation}
where $\mathcal{L}_Q^+$ involves only the light fields, $\mathcal{L}_q$ involves only the heavy fields, and $\mathcal{L}_{Qq}$ contains all interactions between the two. Integrating out the heavy fields yields
\begin{equation}
\mathcal{L}^- = \mathcal{L}_Q^- + M^4 + \ln \left( \frac{\mu}{M} \right) f_0(Q) + \frac{1}{M} f_1(Q) + \frac{1}{M^2} f_2(Q) + \cdots
\end{equation}
which, apart from the cosmological constant term, has at most logarithmic dependence on $M$ (which is induced by quantum corrections). Absorbing the log dependence into $\mathcal{L}_Q^-$ (since it belongs to the wave function renormalization anyway), the lagrangian below the threshold may be recast schematically as
\begin{equation}
\mathcal{L}^- = \mathcal{L}_Q^- + M^4 + \int d^2\theta \left[ \frac{Q^4}{M^\phi} + \frac{Q^5}{M^2\phi^2} + \cdots \right] + \text{h.c.} \nonumber \\
+ \int d^4\theta \left[ \frac{(Q^1Q)^2}{M^2\phi^1\phi} + \frac{(Q^1Q)^3}{M^4(\phi^1\phi)^2} + \cdots \right]
\end{equation}

The additional SUSY breaking introduced by the threshold can now be evaluated. Assuming the threshold does not introduce any relevant\footnote{in the renormalization sense of the word} operators—that is, $\mathcal{L}_Q^+ = \mathcal{L}_Q^-$—the portion involving only $Q$’s must, by assumption, respect the superconformal invariance. Thus, the only SUSY breaking terms that did not exist above $M$ are\footnote{We adopt the notation that an underline indicates the scalar component of a superfield.}
\begin{equation}
\Delta\mathcal{L}_{SB}^- = \left[ \frac{F_Q}{M} Q^3 - \frac{F_\phi Q}{M^2} Q^4 + \text{h.c.} \right] \\
+ \frac{|F_Q|}{M^2} |Q|^2 - \frac{F_\phi Q}{M^2} Q^2 - \frac{F_Q^* Q^4}{M^2} - \frac{|F_\phi|^2}{M^2} |Q|^4 + \cdots
\end{equation}
which may be read off of Eq. (7).

The salient issue is to then determine their size compared to the AMSB contribution. By assumption, $F_\phi \ll M$, $\langle Q \rangle \ll M$, with the latter condition following from the requirement that $Q$ is a light field (that is, if $\langle Q \rangle \sim M$, it would obtain a mass $M$ and be a heavy field). Due to these considerations, the only terms that can potentially contribute SUSY breaking comparable to those of AMSB are those that involve the ratio $F_Q/M$. If it is assumed that

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\[ \langle F_Q \rangle \ll MF_\phi, \text{ then all the terms of Eq. (8) are much smaller than the SUSY breaking contributions due to } F_\phi \text{ so they may indeed be safely neglected.} \]

The argument is then complete: assuming \( \mathcal{L}_Q^+ = \mathcal{L}_Q^- \) and \( F_Q \ll MF_\phi \) implies that no significant SUSY breaking can result from the threshold \( M \) and the only important contribution is due to AMSB. Since the threshold does not provide any significant SUSY breaking, the AMSB expressions must be valid, and as they are correct at any scale, they can just be evaluated at scales below \( M \). This result—that the AMSB form may be employed below \( M \) whilst maintaining ignorance of the theory above \( M \)—is the previously mentioned UV insensitivity of AMSB.

### A. Pomarol and Rattazzi’s Light Singlet Deflection

Given that AMSB’s UV insensitivity is dependent on two assumptions, it is worthwhile to look for any scenarios that may violate these conditions. Since the AMSB expressions are a result of superconformal invariance, these violations must arise from breaking this symmetry; that is, they may only result from fields which acquire a VEV when \( F_\phi \neq 0 \). A subset of this condition was first considered by Pomarol and Rattazzi [14], where they discovered that remnant light singlets can have large \( F \)-term VEVs. The argument proceeds as follows: consider the general form for the VEV of the auxiliary component

\[
\langle F_Q \rangle = L_Q + M_Q \langle Q \rangle + \frac{1}{2!} Y_Q \langle Q \rangle^2 + \cdots.
\]

(9)

Since \( Q \) is a light field, \( \langle Q \rangle \ll M \) so that neither the Yukawa couplings or any higher dimensional operators can yield a VEV of the desired size. This leaves the mass term and the linear term; however, the mass term can not generate an \( F_Q \) VEV of the needed size: while it is possible \( \langle Q \rangle = F_\phi \ll M \), the condition of \( Q \) being light means \( M_Q \ll M \) so that \( M_Q \langle Q \rangle \) is necessarily much smaller than \( MF_\phi \). Thus the only term that can yield a large enough \( F \) VEV is the linear term \( L_Q \). Since any field charged under the residual gauge symmetry below \( M \) would be forbidden from having a linear term, the field whose auxiliary VEV is large must be a singlet of the surviving gauge group.

It is then natural to ask how light this singlet—which shall be called \( S \)—must be; that is, is it possible to raise the mass of \( S \) to the threshold \( M \)? To answer this question, assume
that the field $S$ has a SUSY mass $\mu$. The lagrangian is

$$\mathcal{L}_S = \int d^4\theta \left[ S^\dagger S + cM\phi^\dagger S + cM\phi S^\dagger \right] + \left[ \int d^2\theta \frac{1}{2}\mu\phi S^2 + \text{h.c.} \right]$$

(10)

with $c$ some (real) order one constant. The equations of motion for the auxiliary component gives

$$\langle F_S \rangle = -cMF_\phi - \mu \langle S \rangle$$

(11)

so that the $F$-term is indeed of the correct size. Varying the lagrangian with respect to $S$ yields the constraint

$$\mu (\langle F_S \rangle + F_\phi \langle S \rangle ) = 0$$

(12)

which implies that either $\mu = 0$ or $\langle F_S \rangle = -F_\phi \langle S \rangle$; however, the latter condition cannot be satisfied since $\langle S \rangle \ll M$. A consistent system then demands we take $\mu = 0$. As the singlet is not allowed a SUSY mass term, it must get its mass from the SUSY breaking, so that the singlet’s mass is at or below $F_\phi$.

The result, as argued by Pomarol and Rattazzi[14], is that having light singlets in the low-scale theory permits the threshold to introduce SUSY breaking effects comparable to $F_\phi$ yielding deflection. As Pomarol and Rattazzi’s scheme requires singlets with mass around $F_\phi$, this particular version is called light singlet deflection.

B. Heavy Singlet Deflection

An alternative deflection scenario would arise if the threshold introduces relevant operators into $\mathcal{L}_Q$ which contain a non-superconformal $\phi$ coupling. Such a coupling can not arise from any explicit SUSY mass term in $\mathcal{L}_Q^0$ because superconformal invariance is respected by the SUSY preserving lagrangian. Therefore, the most promising origin of a superconformal violating term would be the VEV of some field which depended on $F_\phi$ (as the SUSY breaking sector does not respect superconformal invariance). Assuming that the $Q$ fields are all charged under the surviving gauge group (to avoid light singlet deflection), as well as insisting this field not break the residual gauge symmetry, means that no $Q$ field can be responsible for this term. Furthermore, as the $Q$s must come in at least pairs (to retain gauge invariance), the field whose VEV imparts a superconformal breaking term in $\mathcal{L}_Q$ would have, at best, a yukawa coupling to the $Q$s. This implies that when said field acquires a VEV, it generates a mass term for the $Q$ fields. This is important because it means that the VEV
can not be of order $M$ (as this would result in the $Q$ fields being heavy), and so whatever field produces the superconformal violation in $\mathcal{L}_Q$ can not be responsible for the threshold $M$.

While the superconformal violating VEV can not be responsible for the threshold $M$, it is possible that the threshold generates the superconformal violating term. The basic idea would be that when the superconformal preserving threshold $M$ is created (however that occurs), it induces a small superconformal violating VEV in one of the other $q$ fields, say $q'$. After the theory is written in the true vacuum (where $\langle q' \rangle = 0$), any field $Q$ having a yukawa coupling to $q'$ would have a corresponding superconformal violating coupling in $\mathcal{L}_Q$ from $q' \rightarrow \langle q' \rangle$. The resulting theory below $M$ would not have any light singlets, but would have deflection from the AMSB trajectories. The next section will discuss an example model to make this idea more concrete.

III. AN EXAMPLE MODEL

Surprisingly, a model commonly used to demonstrate AMSB decoupling\cite{13, 14} also generates non-superconformal couplings in the low-scale theory. The superpotential is

$$\mathcal{W}_0 = (\lambda_X X^2 - M^2 \phi^2) S, \quad (13)$$

where $X$ and $S$ are singlets and an $R$ symmetry is assumed ($S \rightarrow -S$, all other fields invariant) to forbid an explicit mass term or linear term for $X$. A cubic term for $S$ as well as a mass-mixing term for $X$ and $S$ are still allowed; however, they do not alter the discussion and are therefore omitted for simplicity. The resulting lagrangian for Eq. (13) is\cite{13}

$$\mathcal{L} \supset \frac{1}{2} |F_X|^2 + \frac{1}{2} |F_S|^2 + (2\lambda_X FX - 2M^2F_\phi) S + (\lambda_X X^2 - M^2) F_S + \text{h.c.} \quad (14)$$

By solving the equations of motion for the auxiliary fields, the $F$-terms

$$-F_S^* = \lambda_X X^2 - M^2 \quad -F_X^* = 2\lambda_X X S \quad (15)$$

are obtained. In the SUSY limit $\langle F_S \rangle = \langle F_X \rangle = 0$ so that

$$\langle X \rangle = \frac{M}{\sqrt{\lambda_X}} \quad \langle S \rangle = 0 \quad (16)$$

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\footnote{4 We again remind the reader that an underline indicates the scalar component of a superfield.}
After $X$ obtains a VEV, the superfields $X$ and $S$ mix through a SUSY mass term,

$$\mathcal{W} \supset 2\sqrt{\lambda_X} M X S,$$

leaving both fields heavy so that light singlet deflection does not occur. This well-established result implies that the superfield $X$ has a VEV given by $\langle X \rangle = \langle X \rangle \phi$ and superconformal invariance is retained in the theory below $M$.

Activating SUSY breaking shifts the VEVs of the superfields $X$ and $S$ by $O(F_\phi)$. Since $F_\phi \ll M$, the effect on $\langle X \rangle$ is not important; however, $\langle S \rangle \sim F_\phi$ once SUSY is broken. This can be understood from the fact that $\langle X \rangle = \langle X \rangle \phi$ implies $\langle F_X \rangle / \langle X \rangle = F_\phi$ which, combined with the $F$-term expression Eq. (15), implies $\langle S \rangle \sim F_\phi$. Because $F_\phi \neq 0$ breaks superconformal invariance, the induced VEV of $S$ need not preserve the AMSB form. To check this, one may examine the scalar potential before eliminating the auxiliary fields,

$$V = -\frac{1}{2} |F_X|^2 - \frac{1}{2} |F_S|^2 - 2\lambda_X X S F_X + 2M^2 F_\phi S - (\lambda_X X^2 - M^2) F_S + \text{h.c.}$$

Minimizing with respect to $X$ gives

$$\frac{\delta V}{\delta X} = -2\lambda_X F_X S - 2\lambda_X X F_S = 0$$

which can be solved for the ratio of the $F$-term to scalar field VEVs:

$$\frac{\langle F_S \rangle}{\langle S \rangle} = -\frac{\langle F_X \rangle}{\langle X \rangle} = -F_\phi.$$ 

The last equal sign of Eq. (20) follows from the $X$ threshold preserving AMSB. Explicit minimization of the potential after substitution of the $F$-terms produces the same result.

From Eq. (20) it is seen that the threshold set by $\langle S \rangle$ is not AMSB preserving; rather,

$$\langle S \rangle = \frac{\langle S \rangle}{\phi} = -\frac{F_\phi^\dagger}{2\lambda_X \phi}.$$ 

This deflection may be communicated to the low-scale theory by introducing a vector-like pair of fields $Y$ and $\overline{Y}$ (charged under some unspecified gauge group) with the additional term $^5 \mathcal{W}_{\text{Mess}} = \lambda_Y S Y \overline{Y}$. Note that these fields are massless if uncharged under the $R$ symmetry described above. Below $M$ the $Y$’s acquire a mass

$$\mathcal{W}_{\text{Mess}} = \frac{\lambda_Y \langle S \rangle}{\phi} Y \overline{Y} = -\frac{\lambda_Y}{2\lambda_X \phi} F_\phi^\dagger Y \overline{Y}$$

$^5$ As this document was in preparation, [26] appeared with a similar superpotential to the full superpotential, $\mathcal{W}_0 + \mathcal{W}_{\text{Mess}}$, proposed here.
leading to deflection. Recall that both the singlets $X$ and $S$ are heavy, so light singlet deflection does not occur here; rather, the AMSB preserving threshold $M$ induces a super-conformal violating mass term for the light $Y$ fields. These fields then act as messengers of the deflection at a scale naturally around $F_\phi$, but possibly much higher due to the ratio $\lambda_Y/\lambda_X$.

Interestingly, the mass term Eq. (22) is the precise form of the EAM scenario proposed in [17], and the results of the EAM literature[17, 18] can be applied directly to the low-scale theory here. As the original theory had no trace of EAM terms, the heavy singlet deflection scenario is an ultraviolet completion of EAM; however, it is important to emphasize that EAM considers deflection occurring near $F_\phi$ where this scenario permits thresholds far above that scale. This is important because the scale of SUSY breaking in AMSB is $F_\phi$ and it is thus reasonable to expect that physics at this scale would deflect from the AMSB trajectory. Much above $F_\phi$, however, we expect the rules of Pomarol and Rattazzi to be respected: namely no deflection without a remnant light field. This is not what happens for heavy singlet deflection—in fact, this deflection is related to physics far above the $F_\phi$ scale and is not associated with a light remnant field.

A. Deflection

We now explore this deflection in the MSSM. To retain gauge coupling unification, we take the messengers to be in the 5 and \bar{5} representation of $SU(5)$. The effective messenger coupling in Eq. (22) is

$$\lambda \equiv \frac{\lambda_Y}{2\lambda_X}, \quad (23)$$

which sets the messenger scale, $M_{\text{mess}} = \lambda F_\phi$, and as a ratio of two dimensionless couplings can actually be quite large. The strictest constraint is ensuring that the $X$–$S$ sector is tachyon free, which requires $M_{\text{mess}} \lesssim M$; otherwise, $M_{\text{mess}}$ may be arbitrarily high.

Now a general threshold can be parameterized as

$$M_{\text{mess}} = M_{\text{mess}} \left( 1 + F_\phi \theta^2 + r F_\phi \theta^2 \right), \quad (24)$$

where the $M_{\text{mess}}(1 + F_\phi \theta^2) = M_{\text{mess}} \phi$ is the AMSB conserving part of the threshold and $r$ measures deflection. Assuming no visible sector-messenger yukawa couplings, the deflected
gaugino and sfermion masses can be calculated at this threshold as in [18]:

\[
M_{G}\big|_{M_{\text{mess}}} = M_{G}^{\text{AMSB}}\big|_{M_{\text{mess}}} + r F_{\phi} \frac{\Delta \beta_{g_G}}{g_G} \big|_{M_{\text{mess}}}
\]

\[
m^2\big|_{M_{\text{mess}}} = m_{\text{AMS}}^2\big|_{M_{\text{mess}}} + \frac{1}{4} r (r + 2) |F_{\phi}|^2 \frac{\partial \gamma}{\partial g_G} \Delta \beta_{g_G} \big|_{M_{\text{mess}}}
\]

where \(\gamma\) is the anomalous dimension of the scalar (identical above and below the threshold), \(\Delta \beta_{g_G}\) is the difference between beta function above and below the threshold, and the AMSB expressions are given in Eqs. (2) and (4).

From equation Eq. (22) we see that we have \(M_{\text{mess}} = M_{\text{mess}} (1 - F_{\phi} \theta^2)\) so that \(r = -2\). This means

\[
M_{G}\big|_{M_{\text{mess}}} = M_{G}^{\text{AMSB}}\big|_{M_{\text{mess}}} - 2 F_{\phi} \frac{\Delta \beta_{g_G}}{g_G} \big|_{M_{\text{mess}}}
\]

\[
m^2\big|_{M_{\text{mess}}} = m_{\text{AMS}}^2\big|_{M_{\text{mess}}},
\]

indicating that scalar masses will not be directly affected by the deflection but rather are altered through RGE effects of the gaugino masses. It is therefore possible to lift the slepton masses with a large number of messengers \(N_s\) or a large messenger scale (set by \(\lambda\)). We therefore explore this scenario for different values of \(\lambda\) versus the number of messengers.

It is sufficient to check that the right-handed slepton masses are larger than current bounds to ensure that all slepton masses have been lifted. The allowed parameter space is shown in Figure [1]. Here the light purple region labeled \(m^2_{\tilde{\ell}} < 0\) is excluded due to tachyonic sleptons whilst the red region labeled \(\alpha \gg 1\) is excluded because of non-perturbativity of the gauge coupling up to the grand unified theory (GUT) scale.

The spectrum has the general form

\[
m_{\tilde{g}} > m_{\tilde{q}} > m_{\tilde{W}} > m_{\tilde{B}} > m_{\tilde{L}} > m_{\tilde{R}}
\]

with about an order of magnitude difference between the colored fields and right-handed sleptons. The lightest supersymmetric particle (LSP) is the right-handed stau and therefore LHC signals will be dominated by charged tracks and leptonic signals such as described in [27]. The model as-is lacks a dark matter candidate; however, the strong \(CP\) problem may be solved using a Peccei-Quinn symmetry [28–31] which due to SUSY yields an axino in the spectrum. Even if \(R\)-parity is conserved, the stau can decay into the axino with a lifetime of the order of seconds, and is therefore cosmologically innocuous [32]. There are alternative
FIG. 1: $M_{\text{mess}}$ and $\lambda$ versus the number of messengers. The light purple region labeled $m_\ell^2 < 0$ and the red region denoted $\alpha \gg 1$ are excluded due to tachyonic sleptons and non-perturbativity of the gauge coupling up to the GUT scale, respectively.

possibilities such as the inclusion of $R$-parity violating operators\cite{32, 33}, or adding an extra singlet, as in \cite{18, 26}, which could also resolve this issue.

IV. CONCLUSION

We have shown that non-supersymmetric thresholds, which are a necessary ingredient for deflecting off of the AMSB trajectory, are possible even in cases where the threshold is not controlled by a light field (as it is in \cite{14}). The key ingredient is the stabilization of the potential via non-supersymmetric terms, which break the conformal invariance of AMSB. If the corresponding VEV also serves as a threshold for some messengers, then the conformal breaking is broadcast to the low-scale theory resulting in deflection. This can lead to deflection in scenarios previously thought to exhibit decoupling behavior—remarkably even possessing high-scale deflecting thresholds—and indicates that the class of thresholds which lead to AMSB deflection is larger than commonly believed. The simplest example of this mechanism is a UV-completion of the EAM scenario which requires non-WIMP dark matter and sleptons significantly lighter than the rest of the SUSY particles.
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