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Yael Shadmi

Fermi National Accelerator Laboratory
P.O. Box 500, Batavia, Illinois 60510

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Yael Shadmi

Fermi National Accelerator Laboratory
P.O.Box 500, Batavia
IL 60510, USA
email: *yshadmi@fnal.gov*

Abstract

The messenger sector of existing models of gauge-mediated supersymmetry breaking may be simplified by using a non-renormalizable superpotential term to couple the vector-like quark and lepton messenger fields to a chiral gauge-invariant of the supersymmetry-breaking sector. This eliminates the need for a fundamental singlet and for an additional gauge sector needed to generate appropriate expectation values for the singlet component fields. This scenario is more natural if the supersymmetry-breaking sector itself involves a non-renormalizable superpotential. Several examples are constructed based on non-renormalizable $SU(n) \times SU(n-1)$ supersymmetry-breaking theories.

1 Introduction

Most models of gauge-mediated supersymmetry breaking rely on a singlet field S with A- and F-type expectation values to generate supersymmetry breaking masses for a pair of "messenger fields", f and $\bar{f}$, through the superpotential coupling

$$W = S\, f \cdot \bar{f}\ . \tag{1.1}$$

With the fields f, $\bar{f}$ transforming as a vector multiplet of the Standard Model (SM) gauge group, supersymmetry breaking is then communicated to the SM fields through the SM gauge interactions [1].

It is usually non-trivial to generate appropriate expectation values for the singlet. To do that, the most economical models employ a $U(1)$ gauge symmetry sector with superpotential couplings to the singlet, in addition to the basic supersymmetry-breaking sector [1]. Also, an S^3 term must be included in the superpotential to avoid runaway behavior.

But a generic supersymmetry-breaking theory contains different gauge-invariants with different A- and F-type vevs. It is therefore natural to try to use these to replace the fundamental singlet. The field S of eqn. (1.1) is then a composite, and the term (1.1) is a higher-dimension term, suppressed by an appropriate power of some scale M. While the appearance of this scale is in general ad hoc, some supersymmetry-breaking models inherently involve such a scale, since they rely on non-renormalizable superpotentials to achieve supersymmetry breaking [2],[3],[4]. Furthermore, as we will see below, with a renormalizable theory as the supersymmetry-breaking sector, these models are only viable when some dimensionless coupling is taken to be extremely small, on the order of 10^{-9}. This constraint can be alleviated if the supersymmetry-breaking sector involves a non-renormalizable superpotential. The reason for this the following. Since the term (1.1) is suppressed by some power of the scale M, and since we would like M to be large, the fields making up the composite singlet should have large expectation values for the messenger mass scale to be of the correct order. If the supersymmetry-breaking sector is non-renormalizable, with terms suppressed by M, the typical expectation values are naturally large. However, in a renormalizable model, this requires some small coupling. It should be stressed that all our examples do require some small coupling, between 10^{-4} and 10^{-1} depending on the model we consider.

We present several examples of gauge-mediated supersymmetry-breaking models in which the singlet field S is replaced by a composite field of the supersymmetry-breaking sector. As the supersymmetry-breaking sector we use a class of $SU(n) \times SU(n-1)$ gauge theories described in [2].

There are several motivations for using these particular theories. First, the $SU(n) \times SU(n-1)$ theories involve non-renormalizable superpotentials for $n > 4$, and therefore provide a natural setting for introducing the non-renormalizable term (1.1) as explained above.

Second, these theories have supersymmetry-breaking, calculable minima that may be studied through a simple sigma model. In fact, it is possible to study many features of the minimum analytically, and this will prove useful for the present analysis.

As an added bonus, the superpotentials of these theories do not conserve any R-symmetry. Hence, the models we construct are probably the only phenomenological examples with dynamical supersymmetry-breaking that do not resort to supergravity considerations in order to avoid a massless R-axion.

A potentially problematic feature of these models is that they contain massless fermions. However, as we will see in section 3, the massless fermions do not pose any cosmological problem if the scale M is sufficiently big, as is the case in the examples we construct. It should be stressed that the existence of the massless fermions is not related to the focus of this paper, namely, the possibility of eliminating the fundamental singlet. For example, we expect our qualitative results to hold for models based on the analogous $SU(n) \times SU(n-2)$ supersymmetry-breaking theories of [3], which do not contain massless fermions. In fact, our main results probably apply to a much larger class of theories, since they follow from simple dimensional analysis.

We discuss the general requirements on the models, and derive some general results based on dimensional analysis in section 2. In section 3 we study some examples based on $SU(n) \times SU(n-1)$ supersymmetry-breaking theories. Some technical details concerning the $SU(n) \times SU(n-1)$ minimum we consider are collected in the Appendix.

2 Communicating supersymmetry breaking to the standard model

As outlined in the introduction, our models consist, apart from the fields of the supersymmetric Standard Model, of a supersymmetry-breaking sector (SB), and of the vector-like quark and lepton messenger fields [1] f and $\bar{f}$, with the superpotential

$$W \;=\; W_{SB} \;+\; S\, f \cdot \bar{f}\ . \tag{2.1}$$

Here W_{SB} is the superpotential of the supersymmetry-breaking model, and

$$S \;=\; \frac{1}{M^{d-1}} S'\ , \tag{2.2}$$

where S' is a gauge-invariant combination of the fields of the supersymmetry-breaking sector, of dimension d. The field S is chosen so that it has both A- and F-type vacuum expectation values. In the following, we will sometimes refer to these vevs as S and F_S.

Let us now summarize the requirements on the expectation values S and F_S.

First, for the scalar messengers to have positive masses [1] [1],

$$F_S < S^2 \ . \tag{2.3}$$

Second, to generate appropriate masses for the SM superpartners we need [1, 5]

$$\frac{F_S}{S} \sim 10^4 - 10^5 \text{ GeV} \ . \tag{2.4}$$

For brevity, we will require $F_S/S \sim 10^{4.5}$ GeV.

Third, the most serious constraint on these models comes from the requirement that the supersymmetry-breaking scale is low enough. In principle, the Kähler potential may contain higher dimension terms, suppressed by some power of M, that couple either the standard-model fields, or the messenger fields, to the fields of the supersymmetry-breaking sector. Such terms could induce contributions of the order of

$$m_0 \ = \ \frac{F_0}{M} \ , \tag{2.5}$$

to the masses of the scalar messengers, or to the masses of the SM scalar superpartners. Here F_0 is the supersymmetry-breaking scale squared. As we will see shortly, when combined with (2.3), the requirement

$$m_0 \ = \ \frac{F_0}{M} \sim 1 \text{ GeV} \ , \tag{2.6}$$

which would avoid problems with flavor-changing neutral currents, can only be satisfied in the type of models we are considering by taking one of the dimensionless couplings that appear in the superpotential to be extremely small, on the order of 10^{-9}. Although not unnatural in the 't Hooft sense, since taking any of these couplings to zero typically restores some global symmetry, we find this unacceptably small. Instead, we must assume that no higher-dimension terms that couple the SM fields and the supersymmetry-breaking sector fields appear in the Kähler potential at the tree level. We will therefore take $M < M_{Planck}$. Below we choose $M \leq M_{GUT}$.

The scenario we envision is that some new physics takes place above the scale M. This new dynamics involves the fields of the supersymmetry-breaking model (or just some of them) and the messengers, and gives, as its low energy theory, the theory we describe with the superpotential (2.1). It would of course be nice to have an actual microscopic theory that does this, but at present we do not know of such an example.

While it is perhaps not unreasonable to assume that no terms coupling the SM fields to the fields of the supersymmetry-breaking sector appear in the Kähler potential, one cannot

[1] Note that the messenger masses only depend on the absolute value of F_S, but for simplicity, we omit the absolute-value sign throughout this paper.

assume the same for the messenger fields, since these couple directly to the fields of the supersymmetry-breaking sector through the superpotential. It is therefore necessary to ensure that contributions to the messenger masses from possible Kähler-potential terms, of the order F_0/M, are negligible compared to contributions coming from (1.1). In fact, F_0/M should be small compared to both the messenger masses and their mass splittings, in order to generate acceptable masses for the SM superpartners. A non-zero value of $\mathrm{Str}M^2$, taken over the messengers, may lead to negative masses squared for the SM squarks and sleptons, especially in models of the type we are considering, in which a large hierarchy of scales exists due to the presence of non-renormalizable terms suppressed by a large energy scale M [6], [7], [8], (see also [9]). We therefore require[2, 3]

$$\frac{F_0}{M} \leq 10^{-1} \frac{F_S}{S} \ . \tag{2.7}$$

Finally, one would like to have $\sqrt{F_0} \leq 10^9$ GeV, so that supergravity contributions to the superpartner masses are at most at the order of 1 GeV. With (2.4), (2.7), this is automatically satisfied for $M \leq 10^{15}$ GeV. However, for $M = M_{GUT}$, the stricter bound,

$$\frac{F_0}{M} \leq 10^{-2} \frac{F_S}{S} \tag{2.8}$$

is needed, instead of (2.7).

Let us now see what the requirements (2.3), (2.4) and (2.7) imply for our models. Here we will only present rough order-of-magnitude estimates. A more quantitative analysis is undertaken in section 3 where specific examples are studied.

Since the field S is a composite field of dimension d,

$$S \sim \ M \left(\frac{v}{M}\right)^d \ , \tag{2.9}$$

where v is the typical expectation value in the problem. We also have,

$$\frac{F_S}{S} \ \sim \ \left(\frac{v}{M}\right)^{-1} \frac{F_0}{M} \ . \tag{2.10}$$

If no large numerical factors appear in (2.10), eqn. (2.7) (2.8), then imply

$$\frac{v}{M} \leq 10^{-1} \ or \ 10^{-2} \ . \tag{2.11}$$

Now let us assume that the highest-dimension term appearing in the superpotential of the supersymmetry-breaking model, W_{SB}, is also of dimension d. In particular, for $d > 3$, we

[2]The dangerous contribution to the supertrace is of the order $\left(\frac{F_0}{M}\right)^2 \log\left(\frac{M^2}{S^2}\right)$ and for $M = M_{GUT}$ the logarithm is approximately 5.

[3]I thank S. Trivedi for a discussion of this estimate.

assume that this highest-dimension term is necessary for supersymmetry breaking to occur. Then the supersymmetry-breaking scale will typically be of the order

$$F_0 \;\sim\; W/v \;\sim\; \alpha M^2 \left(\frac{v}{M}\right)^{d-1} \tag{2.12}$$

where α is the dimensionless coefficient of the highest-dimension term in the superpotential W_{SB}. We then have

$$F_S \;\sim\; \alpha \;\; M^2 \left(\frac{v}{M}\right)^{2(d-1)} \;, \tag{2.13}$$

and

$$\frac{F_S}{S^2} \;\sim\; \alpha \left(\frac{v}{M}\right)^{-2} \;. \tag{2.14}$$

If no large numerical factors appear in (2.14), we see from (2.3) and (2.11),

$$\alpha \;\leq\; \left(\frac{v}{M}\right)^2 \leq 10^{-2} \; or \; 10^{-4} \;. \tag{2.15}$$

Thus, generically, some of the dimensionless couplings appearing in the supersymmetry-breaking superpotential W_{SB} need to be small in order to satisfy both (2.3), (2.7).

It is worth noting that, whereas the requirement (2.11) holds quite generally in the absence of large numerical factors, (in fact, it is not much of a constraint, since v should be much smaller than M for the analysis to be valid), the condition (2.15) depends sensitively on the assumption that the highest dimension term in W_{SB} is of the same dimension as the composite S. In particular, if the dimension of the composite S is smaller than the highest-dimension term in W_{SB}, the condition (2.15) may be avoided altogether. However, as the examples we discuss in the next section demonstrate, chiral gauge-invariant fields, or moduli, may scale in the same way with v/M even when they have different dimensions. The reason for this is simple–the different terms appearing in W_{SB} are nothing but gauge-invariants, and at a generic minimum these terms are comparable, so that the expectation values of the corresponding gauge invariants only differ by dimensionless couplings.

Finally, it would seem that (2.15) may be avoided if F_S is suppressed compared to S^2. But that typically means that F_S is also suppressed relative to F_0, so that the RHS of (2.10) contains a small factor, which then enters squared in (2.15), making matters worse. One is therefore led to consider regions in which F_S is not particularly suppressed with respect to the other F components in the problem.

At this stage, both α and v/M are determined. The messenger scale

$$\frac{F_S}{S} \;\sim\; \alpha \, M \left(\frac{v}{M}\right)^{d-2} \;\sim\; M \left(\frac{v}{M}\right)^{d} \;. \tag{2.16}$$

is then completely fixed in terms of the scale M. Here we have used (2.9), (2.13), (2.15). For example, for M_{GUT}, with (2.8), one needs $d = 6$ or 7 to obtain the desired messenger scale.

For $M = 10^{15}$ GeV, with (2.7), one needs instead $d = 10$ or 11. Thus, for these models to be viable, the supersymmetry-breaking model must involve a non-renormalizable superpotential.

To summarize, the conditions (2.3), (2.7) imply a specific relation between the coupling α and v/M (see eqn. (2.15)). Then, to generate the correct hierarchy between the messenger scale and the scale M, it is necessary to have, for large M, either a very small coupling, or a non-renormalizable superpotential W_{SB}. Thus, by using a non-renormalizable supersymmetry-breaking sector, one can avoid dimensionless couplings that are extremely small. Indeed, for a renormalizable model, eqn. (2.16) gives, with $d = 3$ and $M = M_{GUT}$, $\alpha \sim 10^{-8}$.

In the next section we will therefore turn to specific examples with a non-renormalizable $SU(n) \times SU(n-1)$ supersymmetry-breaking sector.

3 Models with an $SU(n) \times SU(n-1)$ supersymmetry-breaking sector

3.1 The $SU(n) \times SU(n-1)$ theories

As our supersymmetry-breaking sector we use the $SU(n) \times SU(n-1)$ gauge theories of [2]. These theories have the matter content, $Q \sim (\Box, \Box)$, $L_I \sim (\bar{\Box}, \mathbf{1})$, with $I = 1 \ldots n-1$ and $R_A \sim (\mathbf{1}, \bar{\Box})$, with $A = 1 \ldots n$, and the superpotential

$$W_{SB} = \lambda \; \Sigma_I Y_{II} \; + \; \alpha \; \frac{b^1}{M^{n-4}} \; + \; \beta \; \frac{b^n}{M^{n-4}} \; . \tag{3.1}$$

where $Y_{IA} = L_I \cdot Q \cdot R_A$, and $b^A = (R^{n-1})^A$ are the baryons of $SU(n-1)$. (When appropriate, all indices are contracted with ϵ-tensors).

In the presence of the superpotential (3.1), the original $SU(n-1) \times SU(n) \times U(1) \times U(1)_R$ global symmetry is broken to $SU(n-1) \times U(1)_R$ for $\lambda \neq 0$, which is further broken to $SU(n-2) \times U(1)_R$ for $\alpha \neq 0$. Finally, the last term in (3.1) breaks the $U(1)_R$ symmetry, so that the remaining global symmetry is $SU(n-2)$.

As was shown in [2], these theories break supersymmetry as long as $\alpha \neq 0$. For $\alpha = 0$, the theories have runaway supersymmetric minima along the baryon flat directions, and far along these flat directions, the light degrees of freedom are weakly coupled [10]. Therefore, for large M, the properties of the minimum can be reliably calculated [6]. In [6], this was used to study the minimum of the analogous $SU(n) \times SU(n-2)$ theory (see also [8] for the case of $SU(n) \times SU(n-1)$). We will therefore only outline the main points of the argument here, and refer the reader to [6] for details.

Consider then a D-flat direction with the fields R_A, with $A = 1 \ldots n-1$, obtaining expectation values of order v. The gauge group $SU(n-1)$ is then completely broken at the scale v, while the $SU(n)$ group remains unbroken. However, as a result of the first term in (3.1), all $SU(n)$ fields now get masses of order λv. For large enough v, these fields can be integrated out, leaving, at low energies a pure $SU(n)$ which confines at the scale

$$\Lambda_L = \left((\lambda v)^{n-1}\Lambda^{2n+1}\right)^{\frac{1}{3n}} \ . \tag{3.2}$$

Below this scale, one is then left with the light components of the fields R, with the $SU(n-1)$ dynamics negligibly weak, and the (strong) $SU(n)$ dynamics decoupled, except that its non-perturbative contribution to the superpotential

$$\Lambda_L^3 = \left(\lambda^{n-1}b^n\Lambda^{2n+1}\right)^{\frac{1}{n}} \ , \tag{3.3}$$

arising from gaugino condensation in the pure $SU(n)$, involves the fields R (recall $b^n \sim (R_1 \ldots R_{n-1})$). As was argued in [10], quantum corrections to the Kähler potential for the fields R are very small, so that it is of the form

$$K \ = \ R^{\dagger A}\, R_A \ . \tag{3.4}$$

Thus, all the properties of the vacuum may be calculated.

As in [6], we will find it convenient to work in terms of the baryons. Our low energy theory is then a theory of the n baryons b^A, with the superpotential

$$W_{SB} \ = \ \left(\lambda^{n-1}\,\Lambda^{2n+1}\,b^n\right)^{\frac{1}{n}} \ + \ \alpha\,\frac{b^1}{M^{n-4}} \ + \ \beta\,\frac{b^n}{M^{n-4}} \ , \tag{3.5}$$

and the Kähler potential obtained from (3.4) as in [11], [12],

$$K \ = \ (n-1)\,(b_A^\dagger b^A)^{\frac{1}{n-1}} \ . \tag{3.6}$$

At the minimum we consider, the only baryons with non-zero vevs are b^1 and b^n. It is convenient to define r and v such that

$$b^1 = r\,b^n \quad \text{and} \quad b^n = v^{n-1} \ , \tag{3.7}$$

where, as in the subsequent discussion, b^A stands for the expectation value rather than the field. The ratio r is determined by the ratio of dimensionless coupling β/α, and is given in the appendix, where various details of the minimum are summarized.

We can then write the F-type expectation-values of b^1, b^n as,

$$F_{b^A} \ = \ F_A\,\alpha\,v^n\,\left(\frac{v}{M}\right)^{n-4} \ , \tag{3.8}$$

where $F_{A=1,n}$, which are dimensionless functions of n and r, are given in the appendix.

This is in fact all we need if we only wish to use the baryon operators as our composite singlets. It would also be useful however to consider the trilinears Y_{IA} for this purpose. Their vevs are given by (see appendix),

$$Y_{II} \;=\; \frac{n\,\alpha}{q\,\lambda}\,M^3\left(\frac{v}{M}\right)^{n-1} \quad , \quad Y_{1n} \;=\; r\,Y_{II} \tag{3.9}$$

with $I = 1 \ldots n-1$, and where q is a function of n and r, given in the appendix.

Finally the F-type vevs of the fields R, may be written as

$$F_{R_A} \;=\; f_{r_A}\,\alpha\,M^2\,\left(\frac{v}{M}\right)^{n-2} \quad , \tag{3.10}$$

where again, f_{r_A} are dimensionless functions of n and r and are given in the appendix.

3.2 $S = b^n/M^{n-2}$

Choosing $S = b^n/M^{n-2}$ we have,

$$S \;=\; M\,\left(\frac{v}{M}\right)^{n-1} \quad , \tag{3.11}$$

$$\frac{F_S}{S} \;=\; F_n\,\alpha\,M\,\left(\frac{v}{M}\right)^{n-3} \quad , \tag{3.12}$$

$$\frac{F_S}{S^2} \;=\; F_n\,\alpha\,\left(\frac{v}{M}\right)^{-2} \quad . \tag{3.13}$$

To satisfy the requirements (2.3) and (2.7) without having very small couplings, it is best to choose a region in which F_S is not suppressed compared to the other F components in the problem, so that F_n is order 1. To see this, note that

$$\frac{F_S}{S} \;\sim\; F_n\,\frac{F_0}{v}\;.$$

Therefore, the smaller the factor F_n gets, the smaller the value of v that is needed to keep F_0 low. Since v enters squared in (3.13), this would require a smaller coupling α as well.

We find that the optimal choice is $r \sim 0.5$ (corresponding to β/α between 0.5 and 0.74 for $n = 4 \ldots 20$). Taking $M = M_{GUT}$ and $n = 8$, the different requirements on F_S and S can be met with $\alpha = 3.2 \cdot 10^{-4}$ and $v/M = 2.4 \cdot 10^{-2}$. Alternatively, for $n = 7$, one can take $\alpha = 9 \cdot 10^{-5}$, with $v/M = 1.3 \cdot 10^{-2}$.

Note that since $M = M_{GUT}$, we use the stronger constraint (2.8). Taking instead $M = 10^{15}$ GeV, for which the less-stringent constraint (2.7) can be used, we find that for $n = 12$ $\alpha = 7 \cdot 10^{-3}$ and $v/M = 0.11$. Raising n to $n = 13$, one can take $\alpha = 1 \cdot 10^{-2}$ with $v/M = 0.14$. For all these choices, and in the following section, $F_S/S = 10^{4.5}$ GeV, $F_S/S^2 = 0.75$ and $F_0 = 1 - 2 \cdot 10^{18}$ GeV^2.

Choosing the baryon b^1, instead of b^n, to play the role of the singlet leads to similar results.

It is amusing to note that these models contain gauge-invariant operators that are natural candidates for generating a μ-term. Consider for example the $SU(8) \times SU(7)$ model with $S = b^8/M^6$, and add the superpotential term

$$\frac{1}{M^2} Y_{22}\, H_U\, H_D \ , \tag{3.14}$$

where H_U and H_D are the two Higgs doublets. The F-vev of Y_{22} vanishes for $r = 0.57$. For this choice then, (3.14) generates a μ-term but no $B\mu$-term. Also note,

$$\frac{Y_{22}}{S} = nq^{-1}\,\frac{\alpha}{\lambda} \ , \tag{3.15}$$

so taking $\lambda = 1$, we get $Y_{22} \sim 10^2$ GeV (where we also used the fact that $q^{-1} \sim 0.8$).

However, we have assumed throughout that the Kähler potential does not contain any terms that couple the SM fields to the fields of the supersymmetry-breaking sector. Such terms, if present, would contribute masses of order $10^{2.5}$ GeV to the SM scalars superpartners. But this assumption would be quite implausible if we allowed superpotential terms of the form (3.14).

3.3 $S = Y/M^2$

We can also take the trilinear invariants, Y_{IA}, to replace the singlet. Here we take $S = Y_{1n}/M^2$, which turns out to be the optimal choice. For this choice we have

$$S \ = \ \frac{n}{q\,\lambda}\, M \left(\frac{v}{M}\right)^{n-1} \ , \tag{3.16}$$

where we used (3.9), and,

$$\frac{F_S}{S} \ = \ \frac{(1+r^2)^{\frac{n-2}{2(n-1)}}}{r}\, f_{R_n}\, M \left(\frac{v}{M}\right)^{n-3} \ . \tag{3.17}$$

Here and throughout this section, we set the dimensionless baryon coupling, α, to 1. As we will see, the "small coupling" in this case is the Yukawa coupling λ, multiplying the trilinear terms in W_{SB}. Note that this coupling drops out in the ratio F_S/S, but appears in the ratio F_S/S^2.

Again, it is best to consider regions in which f_{R_n} is not small, and we choose $r = 0.5$. For $M = M_{GUT}$, one can take $n = 10$, with $\lambda = 4 \cdot 10^{-3}$, and $v/M = 2.2 \cdot 10^{-2}$. Choosing instead $M = 10^{15}$ GeV, we need $\lambda = 1 \cdot 10^{-2}$, $v/M = 8.8 \cdot 10^{-2}$ with $n = 13$, and $\lambda = 0.12$ with $v/M = 0.1$ for $n = 14$.

Recall that to get the low-energy theory we are using, we have integrated out the fields Q and L, assuming their masses, λv, are much bigger than Λ, the scale of the $SU(n)$ group. Since we are now considering small values of λ, we must make sure that the ratio (see appendix)

$$\frac{\Lambda}{\lambda v} = \left(n\, q^{-1}\alpha\lambda^{-3}\left(\frac{v}{M}\right)^{n-4}\right)^{\frac{n}{2n+1}} \tag{3.18}$$

is still small. It is easy to see that for sufficiently high values of n, this is indeed the case. Setting $\alpha = 1$ and neglecting q^{-1}, which is order 1 for $r \sim 0.5$, one can check that it is acceptably small in all our examples.

Note that the "small coupling" in this case is around 10^{-3} for M_{GUT}, and 10^{-1} for M^{15} GeV, an order of magnitude bigger than the "small coupling" that is required when using the baryon as the singlet. The difference is due to a numerical factor–essentially a factor of n that enters the ratio F_S/S^2.

Finally, we note that for $M < 10^{15}$ GeV, the typical size of the "small coupling" remains the same (see eqn. (2.15)), but the value of n goes down.

3.4 Discussion

Throughout this section, we have assumed only one term of the form (1.1). This cannot be justified by any symmetry arguments, since the only global symmetry we have left is an $SU(n-2)$ global symmetry, which can be invoked to rule out terms such as $b^A f\bar{f}$ with $A = 2\ldots n-1$. However, our qualitative results remain unaffected even if several terms of the form (1.1), with different composites appear, unless some special cancellation occurs. First, we note that vevs of the baryons and trilinears differ by the "small coupling", either α or λ, which gives at least an order of magnitude difference. Thus, if we use a baryon to generate the messenger masses, through the term $bf\cdot\bar{f}/M^{n-2}$, additional terms such as $Yf\cdot\bar{f}/M^2$, are negligible, and vice versa. Furthermore, in the examples we constructed with $S = b^n$, b^1 had comparable, or smaller vevs. Its presence in the superpotential would thus not affect the results dramatically, unless its coupling to the messengers appears with a different coefficient, such that some combination of expectation values conspires to cancel. The same is true for the trilinears.

Let us now summarize the different energy scales that appear in these models. For concreteness, take the $SU(7)\times SU(6)$ model with $M = M_{GUT}$ and $S = b^7/M^3$. The $SU(6)$ group is broken at the scale $v \sim 10^{14}$ GeV, which is also the mass scale of the fields Q and L. $SU(7)$ then confines at the scale $\Lambda \sim 10^{10}$ GeV.

The light fields of the supersymmetry-breaking sector are the baryons b^1, b^7, whose scalar and fermion components have masses of $10^{4.5}$ GeV, and the baryons b^A with $A = 2\ldots 6$, which make up one fundamental of the unbroken global $SU(5)$, whose fermion components are massless (as required by anomaly matching), and whose scalar components have masses of order $10^{4.5}$ GeV. Finally, the messenger masses are also of order $10^{4.5}$ GeV, and the supersymmetry-breaking scale is $\sqrt{F_0} \sim 10^9$ GeV.

New massless fermion species, beyond those present in the Standard Model, may spoil the predictions of standard nucleosynthesis theory, if they contribute significantly to the entropy at the time of nucleosynthesis ($T \sim 1$ MeV) [15]. However, the massless fermions of our models–the fermion components of $b^{2\dots(n-2)}$–interact extremely weakly, so that their decoupling temperature is very high. Consequently, their contribution to the entropy at the time of nucleosynthesis is negligible. To see this, note that at sufficiently low temperatures, the intreractions of these fermions are described by the low-energy Lagrangian derived from (3.5), (3.6) (see [16]). Their dominant interaction comes from a 4-fermion term suppressed by v^{-2}. The rate of this interaction is therefore $\Gamma \sim v^{-4}T^5$, which is comparable to the expansion rate $H \sim T^2/M_{Planck}$ only for $T \geq 10^{13}$ GeV[4].

As mentioned above, our models also contain exotic scalars and fermions with masses around $10^4 - 10^5$ GeV. These would be present in generic models of the type we consider, whereas the existence of the massless fermions is a specific feature of the $SU(n) \times SU(n-1)$ supersymmetry-breaking sector. The interactions of this exotic matter are again extremely weak. The dominant fermion interaction is the 4-fermion interaction mentioned above. The scalar-interaction Lagrangian derived from (3.5), (3.6) contains couplings involving only scalar baryons, as well as couplings of scalar baryons to scalar messengers. (Note that all scalar baryons have couplings to the messengers through the Kähler potential (3.6)). These couplings are very small. The typical 4-scalar term has a coefficient $(v/M)^{2n-4}$, and higher order terms are further suppressed by negative powers of v. Therefore, the interactions of these fields are not thermalized at temperatures for which the low-energy effective theory is valid.

In fact, the maximum reheating temperature after inflation is constrained by requiring that the decay of the LSP to the gravitino does not overclose the universe [17]. In our case, the gravitino mass is $\mathcal{O}(1)$ GeV, for which the authors of ref. [17] conclude that the reheating temperature cannot exceed $\mathcal{O}(10^8)$ GeV. Therefore, once diluted by inflation, the baryon fields are not produced thermally[5].

Finally we note that the superpartner spectrum of our models is identical to that of the

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[4]At this temperature the low-energy theory is no longer valid for specific models. In all the examples we considered however, the low energy theory is valid below, say, 10^9 GeV, where the rate is even smaller.

[5]The baryons are reminiscent of the moduli of Hidden Sector models in that they interact very weakly and have large vevs. One may therefore worry about the analog of the Polonyi problem [18]. However, here v is at most 10^{14} GeV and the mass of the baryons is $10^4 - 10^5$ GeV so the ratio of baryon density to entropy, which scales like $v^2 m_b^{-1/2}$ is about nine orders of magnitude smaller than in the supergravity case.

models of [1], since the masses of the Standard Model superpartners only depend on the messenger masses. The only different feature, from the point of view of phenomenology, is that the supersymmetry-breaking scale is relatively high, $\sqrt{F_0} \sim 10^9$ GeV, so that the decay of the LSP to the gravitino would not occur inside the detector. In this respect our models are similar to the models of [6], [8].

4 Conclusions

In this paper we explore the possibility of eliminating the fundamental singlet of existing models of gauge-mediated supersymmetry breaking, by introducing a non-renormalizable superpotential term that couples the messengers to a chiral gauge-invariant of the supersymmetry-breaking theory.

We show that to obtain viable models without $\mathcal{O}(10^{-9})$ couplings, the theory used as the supersymmetry-breaking sector should have a non-renormalizable superpotential.

We then construct several examples with non-renormalizable $SU(n) \times SU(n-1)$ theories as the supersymmetry-breaking sector, taking different gauge-invariants to replace the fundamental singlet. These examples only require couplings of order $10^{-4} - 10^{-3}$ for $M = M_{GUT}$, and of order $10^{-2} - 10^{-1}$ for $M \leq 10^{15}$ GeV, where M is the suppression scale of the non-renormalizable terms.

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Note: After this work was completed I became aware of the work of Haba, Maru and Matsuoka [14], hep-ph/9703250 and hep-ph/9612468, where the same idea is considered, but with different supersymmetry-breaking theories. The model of hep-ph/9612468 uses the renormalizable 3-2 model with a coupling of order 10^{-9}, in agreement with the discussion of section 2.

Appendix

This appendix summarizes some details regarding the minimum we consider.

It is convenient to work in terms of the baryon fields b^A. The Kähler metric can be derived from (3.6),

$$g_{AB} \;=\; k^{\frac{1}{n-1}-2}\left(-\frac{n-2}{n-1} b_B^\dagger b_A \,+\, k\,\delta_{AB}\right) \;, \tag{A.1}$$

with $k = b_A^\dagger b^A$, and can be easily inverted to get,

$$g^{-1}_{AB} \;=\; k^{-\frac{1}{n-1}}\left((n-2)\, b_A^\dagger b_B \,+\, k\,\delta_{AB}\right) \;. \tag{A.2}$$

The potential is then given by,

$$g^{-1}_{AB}\, W_A\, W_B \;, \tag{A.3}$$

with, using (3.5),

$$W_1 \;=\; \alpha \;, \tag{A.4}$$

$$W_n \;=\; \frac{1}{n}\left(\lambda^{n-1}\,\Lambda^{2n+1}\right)^{\frac{1}{n}}\,(b^n)^{\frac{1}{n}-1} \,+\, \beta \;, \tag{A.5}$$

$$W_A \;=\; 0 \quad for\ A = 2\ldots n-1 \;. \tag{A.6}$$

One can then show analytically that the potential is minimized for $b^1 = rb^n \equiv rv^{n-1}$, where $0 < r < \sqrt{n-1}/2$ is determined from

$$\frac{\beta}{\alpha} \;=\; \frac{(r^2+1)\,(n\,r^2 - r^2 + 2)\,P - 3\,r^4 + 8n\,r^4 - 2n + 5r^2 + n^2 r^2 - 4nr^2 - 3n^2 r^4 + 2}{2(n-1)\,(2n-2+r^2)\,r^3} \;, \tag{A.7}$$

where

$$P \;=\; \sqrt{n-1}\,\sqrt{n-1-4r^2} \;, \tag{A.8}$$

and with v given by,

$$\frac{\Lambda}{v} \;=\; \left(nq^{-1}\alpha\lambda^{-\frac{n-1}{n}}\left(\frac{v}{M}\right)^{n-4}\right)^{\frac{n}{2n+1}} \;, \tag{A.9}$$

where

$$q \;=\; -\frac{(n-1)}{n\,(n-2)}\,\frac{r\,(P+n-r^2-1)}{(r^2+1)} \;. \tag{A.10}$$

Note that $q < 0$. We therefore take α to be negative. The bounds on α appearing in the text refer to its absolute value.

The functions F_1, F_n defined in (3.8) are then given by,

$$F_1 \;=\; 1 \,+ (n-1)\,r^2 \,+\, (n-2)\,r\left(\frac{1}{q} \,+\, \frac{\beta}{\alpha}\right) \;, \tag{A.11}$$

$$F_n \;=\; (n-2)\,r \,+\, (n-1+r^2)\left(\frac{1}{q} \,+\, \frac{\beta}{\alpha}\right) \;. \tag{A.12}$$

The simplest flat direction that results in the baryon configuration (3.7) is of the form $R_{Ai} = av\delta_{Ai}$ for $A = 2\ldots n-1$, $R_{11} = a^{-(n-2)}v$, and $R_{n1} = ra^{-(n-2)}v$, where $a = (1+r^2)^{\frac{1}{2(n-1)}}$, and the second index on R is the $SU(n-1)$ gauge index.

We then have

$$f_{R_1} \;=\; \frac{1}{n-1}\,(1+r^2)^{-\frac{n-2}{2(n-1)}}\,\left(1\;-\;\frac{(n-2)\,F_1}{r\,F_n}\right)\,F_n\;, \tag{A.13}$$

$$f_{R_n} \;=\; -\,\frac{n-2}{n-1}\,(1+r^2)^{-\frac{n-2}{2(n-1)}}\,r\,\left(1\;-\;\frac{F_1}{(n-2)\,r\,F_n}\right)\,F_n\;, \tag{A.14}$$

$$f_{R_A} \;=\; \frac{1}{n-1}\,(1+r^2)^{\frac{1}{2(n-1)}}\,\left(1\;+\;\frac{F_1}{r\,F_n}\right)\,F_n\quad for\;\,A=2\ldots n-1\;. \tag{A.15}$$

Finally, to obtain the expectation values of the trilinears Y_{IA}, recall that $Y_{IA} = L_I \cdot Q \cdot R_A$, and that L and Q are the heavy flavors of $SU(n)$ with a mass matrix $m = diag(R_{11}, .., R_{(n-1)(n-1)})$. Therefore, using [13]

$$\langle Q\cdot L\rangle \;=\; \left(\Lambda^{2n+1}\det m\right)^{\frac{1}{n}}\;m^{-1}\;, \tag{A.16}$$

one finds (3.9).

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