One-scale model of dynamical supersymmetry breaking

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A model of gauge-mediated supersymmetry breaking is constructed in which the low-energy physics depends on a single dynamical scale. The strong coupling dynamics of gauge theories plays an important role, in particular through its effects on beta functions and through confinement. The model does not have distinct messenger and supersymmetry-breaking sectors. The scale of supersymmetry breaking is of the order of 10–100 TeV, implying that the decay of the next-to-lightest superpartner into the gravitino is prompt. Super-oblique corrections are enhanced. A Dirac fermion and one complex scalar, in a 10 or $\overline{10}$ of (global) SU(5), are predicted to be relatively light and to satisfy certain mass relations with the standard model squarks and sleptons. [S0556-2821(98)01423-4]

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I. INTRODUCTION

Supersymmetry, broken dynamically, solves the gauge hierarchy problem [1]. Communicating supersymmetry breaking to the superpartners of the minimal supersymmetric standard model (MSSM) via the ordinary SU(3)×SU(2)×U(1) gauge interactions provides a natural explanation for degeneracy of the squarks and sleptons, avoiding large contributions to quark and lepton flavor violation from the superpartners. Such gauge-mediated supersymmetry-breaking (GMSB) models [2–9] are therefore very attractive and have received much attention recently [10–24]. In principle, with GMSB, the scale of supersymmetry breaking $\sqrt{F}$ could be as low as $\sim 4\pi m_W/\alpha_W \sim 30$ TeV, where m_W is the weak scale and α_W is the weak fine structure constant. With R parity conservation and with $\sqrt{F} \lesssim 1000$ TeV, there is the exciting prospect of observing the decay of the next-to-lightest superpartner (NLSP) into the gravitino $\tilde{G}$ and ordinary particles in a typical particle physics detector [10–24]. Furthermore, the decay rate into $\tilde{G}$ scales as F^2 and gives a sensitive probe of the supersymmetry-breaking sector. One candidate for such an event, with two leptons, two photons, and missing energy, has already been reported by the Collider Detector at Fermilab (CDF) Collaboration [25]. Interpreting the photons as coming from the prompt decay of the lightest neutralino into a photon and gravitino requires $\sqrt{F} \lesssim 100$ TeV.

Nearly all explicit models of dynamical supersymmetry breaking with GMSB predict a supersymmetry breaking scale $\sqrt{F}$ which is in the range $10^3 - 10^8$ TeV—too high to allow the prompt decay of the NLSP into the gravitino [7–9,26–41]. (The only reported exceptions [42–44] may or may not have a strongly-coupled local minimum with broken supersymmetry, and also have supersymmetric vacua.) There are several reasons why GMSB models typically have such a large supersymmetry breaking scale. In all GMSB models, the ordinary superpartners gain mass though loops involving ''messenger'' particles which carry SU(3)×SU(2)×U(1) quantum numbers and have a nonsupersymmetric spectrum. In some models $\sqrt{F}$ must be high because supersymmetry breaking is communicated to the messengers via some weak coupling from a dynamical supersymmetry breaking (DSB) sector [7–9,26–33]. In other, more aesthetically pleasing, models the messengers are an integral part of the DSB sector [34–41,45], as first suggested by Affleck, Dine, and Seiberg [46]. The latter models typically have many particles carrying SU(3)×SU(2)×U(1) interactions, with perturbative unification of the standard model gauge couplings only possible if the additional particles are very heavy. Most such models constructed to date achieve SU(3)×SU(2)×U(1) unification by having two or more scales involved in the supersymmetry-breaking dynamics, with the majority of the new particles at a messenger mass scale M which is much heavier than the supersymmetry breaking scale. Since the ordinary superpartner masses are proportional to F/M, both scales are required to be rather high.

In this paper we present a model of dynamical supersymmetry breaking whose low-energy physics is determined by a single energy scale of order 100 TeV. To our knowledge, this is the first explicit, natural model with no supersymmetric minima, all scales generated dynamically, and prompt decay of the NLSP into the gravitino. The model has no segregation of DSB and messenger sectors, is completely chiral and contains no fundamental gauge singlets. It has limits in which one can show that the global minimum of the potential breaks supersymmetry but leaves color and electromagnetism unbroken. Perturbative unification of the SU(3)×SU(2)×U(1) interactions is possible and gives the usual successful prediction for $\sin^2\theta_W$. As our example is strongly coupled at the supersymmetry breaking scale, it is somewhat less predictive than most explicit GMSB models. Still, many of the usual GMSB predictions survive. Unfortunately it is difficult to solve the μ problem in this model.

In Sec. II we give an quick overview of the model. We prove the model breaks supersymmetry in Sec. III. Since the

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TABLE I. Quantum numbers of chiral superfields in the model. $SU(5)_G$ is a global symmetry containing the standard model.

	Sp(4)	SU(3)	SU(2)	$SU(5)_G$
q	1	$\square$	$\square$	1
$\bar{u}$	1	$\bar{\square}$	1	1
$\bar{d}$	1	$\bar{\square}$	1	1
l	1	1	$\square$	1
$\bar{T}$	$\square$	$\bar{\square}$	1	1
$\bar{V}$	$\square$	1	1	$\bar{\square}$
A	1	1	1	$\begin{smallmatrix}\square\\\square\end{smallmatrix}$
B	1	$\square$	1	$\square$
C	1	$\bar{\square}$	1	1

model has complicated behavior due to strong coupling, we review various facts about strong dynamics in $\mathcal{N}=1$ supersymmetric gauge theories in Sec. IV. In Sec. V, we give a more detailed discussion of the model's dynamics, and then justify our claims carefully. The low-energy properties of the model are explained in Sec. VI; the reader who is mainly interested in the implications for experiment can skip to this section. The conclusion contains a summary of our results.

II. THE MODEL: A FIRST SKETCH

The model we consider contains, in addition to the standard model, a dynamical supersymmetry-breaking sector with gauge group $Sp(4)\times SU(3)\times SU(2)$. We will refer to the coupling of the standard model color and weak interactions as $g_3^{\rm SM}$ and $g_2^{\rm SM}$ to distinguish them from g_3 and g_2 of the 3-2 supersymmetry-breaking sector. The matter content of the model is given in Table I. The SU(5) in the last column is the usual grand unification group containing the standard model. Although we do not require gauge group unification and treat SU(5) merely as an approximate global symmetry, we consider only complete multiplets of SU(5) in order that standard model gauge coupling unification be maintained. [The SU(5) assignments could be charge conjugated without changing the model.]

First, we give a brief motivation for the model. The fields $q,\bar{u},\bar{d},l$ make up the matter content of the famous supersymmetry-breaking $SU(3)\times SU(2)$ model of Affleck, Dine, and Seiberg [47–49] (the "3-2 model"). The Sp(4) gauge group has a total of eight fields in its fundamental representation, coming from $\bar{T}$ and $\bar{V}$, and consequently will confine at low energies [50], at a scale $\Lambda_{\rm cont}$. The resulting massless bound states $\bar{A}=(\bar{V}\bar{V})$, $\bar{B}=(\bar{T}\bar{V})$, $\bar{C}=(\bar{T}\bar{T})$, have the correct quantum numbers to pair up with the fields A,B,C and become massive. Thus, below the Sp(4) confining scale the theory will consist of the standard model, the massless fields of the 3-2 model, and massive fields which couple to both sectors and act as messenger fields by communicating the supersymmetry breaking of the 3-2 model to the standard model.

For the model to behave in this way requires a superpotential

$$W=W_{\rm SM}+W_{3\text{-}2}+W_m+W_s\,, \tag{1}$$

where $W_{\rm SM}$ is the standard model superpotential,

$$W_{3\text{-}2}=\lambda_0 q\bar{u}l \tag{2}$$

is the usual superpotential of the 3-2 model, and

$$W_m=y_A A\bar{V}\bar{V}+y_B B\bar{T}\bar{V}+y_C C\bar{T}\bar{T} \tag{3}$$

serves to give masses to the messenger fields. Finally, the couplings

$$W_s=h_1 Cql+h_2 C\bar{u}\bar{d} \tag{4}$$

are not needed to ensure an acceptable pattern of symmetry breaking but will help avoid having stable heavy messenger particles. Note that this is the most general renormalizable superpotential consistent with gauge symmetry which does not couple the MSSM to the DSB sector. For simplicity in the following discussion, we will assume the $h_{1,2}$ couplings are small and have little effect on the dynamics, although this is not essential or even likely, since they get enhanced by strong coupling effects.

The dynamics of the theory is intricate. The SU(3) gauge coupling is expected to flow slowly due to higher loop effects, and approach a fixed point at extreme low energy. As a result, the scale Λ_3 is washed out by the dynamics. By taking $\Lambda_4\ll\Lambda_3$, we can arrange that Sp(4) confinement, at the scale $\Lambda_{\rm conf}$, occurs when the coupling g_3 is substantial. Associated masses of order $\Lambda_{\rm conf}$ for the $B,\bar{B}$ and $C,\bar{C}$ fields remove all SU(3)-charged fields except those of the 3-2 model. The SU(3) beta function then becomes large, causing g_3 immediately to blow up, breaking supersymmetry. We therefore expect the scale $\sqrt{F}$ of supersymmetry breaking to be of order the dynamical scale $\Lambda_{\rm conf}$.

Thus, in this model the messengers and the supersymmetry breaking lie at or near the same scale, which we take to be of order 10–100 TeV. Note that the model has neither vectorlike matter nor nondynamical mass scales. The gravitino is light, and its properties are similar to those of other low-scale GMSB models; it can serve to explain the $e^+e^-\gamma\gamma$ event observed at CDF [25].

The model has another feature which appears in certain regions of parameter space. As we will see below, the fact that the field A is a 4-3-2 gauge singlet tends to make the $A\bar{A}$ dynamical mass smaller than $\Lambda_{\rm conf}$. As a result, the Dirac fermion $\psi_A,\psi_{\bar{A}}$ and the complex scalar A might (but need not) be much lighter than $\Lambda_{\rm conf}$. (The scalar $\bar{A}$ is a composite of strongly interacting fields and will get a large supersymmetry-breaking mass.) The mass spectrum of these fields is interesting and will be discussed in detail in Sec. VI B.

The effects on the standard model superpartners resemble those in usual GMSB models that have heavy messengers charged in both the supersymmetry-breaking and standard model groups. However, because the supersymmetry-breaking sector is strongly interacting, and because the messengers have masses near $\Lambda_{\rm conf}$, there is no separation of scales in this model. In fact one cannot really talk of a ''messenger sector''; the strong dynamics as a whole is responsible for the message. The effective action below the scale $\Lambda_{\rm conf}$ is already far from supersymmetric. This can make some aspects of the model quite different from GMSB models with weakly coupled messenger sectors. For example, the overall scale of the gaugino masses is unrelated to that of the sfermion masses because of the strong dynamics at the scale $\Lambda_{\rm conf}$. Also, there are relatively large ''superoblique'' corrections to gaugino couplings, of order $\alpha_i^{\rm SM}\log(\Lambda_{\rm conf}/m_{\psi_A})/(4\pi)$. Still, the strong couplings of the model preserve an approximate SU(5) global symmetry, which ensures that masses of different gauginos are related by standard model gauge couplings, and similarly for sfermion masses.

TABLE II. Quantum numbers of chiral superfields after Sp(4) confines.

	SU(3)	SU(2)	$SU(5)_G$
q	$\square$	$\square$	1
$\bar{u}$	$\bar{\square}$	1	1
$\bar{d}$	$\bar{\square}$	1	1
l	1	$\square$	1
$\bar{A}$	1	1	$\overline{\begin{smallmatrix}\square\\\square\end{smallmatrix}}$
$\bar{B}$	$\bar{\square}$	1	$\bar{\square}$
$\bar{C}$	$\square$	1	1
A	1	1	$\begin{smallmatrix}\square\\\square\end{smallmatrix}$
B	$\square$	1	$\square$
C	$\bar{\square}$	1	1

III. BREAKING OF SUPERSYMMETRY

In this section we will demonstrate that the model breaks supersymmetry, first showing the model has no flat directions at the quantum level, and then demonstrating that supersymmetry is broken for a generic choice of parameters.

A. Absence of flat directions

Our model has no flat directions at the quantum mechanical level, and hence no supersymmetric minima infinitely far away in field space. Here we study the classical flat directions, which are labeled by holomorphic invariants built from the chiral fields. To simplify the discussion, we rescale all Yukawa couplings to 1.

Any holomorphic invariant involving fields charged under Sp(4) must involve one of $\overline{VV}$, $\overline{TV}$ or $\overline{TT}$. The first two are set to zero by the F-flatness conditions $\partial W/\partial A=0$ and $\partial W/\partial B=0$, while $\partial W/\partial C=0$ assures $\overline{TT}=-(ql+\overline{ud})$. Using $\partial W/\partial\bar{u}=0=\partial W/\partial\bar{d}$ and antisymmetry, one can show the operators $\overline{TTC},\overline{TTu},\overline{TTd}$ are all zero. The operator $\bar{T}qq\bar{T}$ also vanishes; since $qqql$ is identically zero, the F-flatness conditions $\partial W/\partial C=0=\partial W/\partial l$ imply that $\bar{T}qq\bar{T}\propto\bar{u}qq\bar{d}$ $\propto Cqq\bar{d}$, which in turn is zero by $\partial W/\partial\bar{u}=0$. All operators which involve only the 3-2 fields q, $\bar{u}$, $\bar{d}$, and l must be zero. Finally there are some classical flat directions which combine A, B, and C with the 3-2 fields q, $\bar{u}$, $\bar{d}$, l. However, as we now show, even these are removed by quantum mechanical effects.

Along any classically flat direction with expectation values for A, B, or C, some of the fields in the fundamental representation of Sp(4) (components of $\bar{V}$ and/or $\bar{T}$) will be massive. The number of remaining massless fundamentals may be six, four, two or zero. In each case, strong-coupling dynamics of the Sp(4) group [50] then generates a potential energy. If the number remaining is six, then the classical moduli space of the Sp(4) theory is modified by the constraint that $V^5T^3\sim\Lambda_L^8$ [here Λ_L is the *low-energy* Sp(4) strong-coupling scale]. The requirements $\partial W/\partial A=\partial W/\partial B$ $=\partial W/\partial C=0$ imply that $V^5T^3=0$ for a zero-energy vacuum, in contradiction to the previous condition. If the number remaining is 4 (2), the Sp(4) theory generates an Affleck-Dine-Seiberg superpotential via instantons (gaugino condensation) which again lifts the flat directions. And if all of the fields $\bar{V}$ and $\bar{T}$ are massive, then gaugino condensation generates a superpotential $W=\Lambda_L^3$, where again Λ_L is the *low-energy* Sp(4) strong-coupling scale, related by $\Lambda_L^9=A^2BC\Lambda_4^5$ to the high-energy Sp(4) strong-coupling scale Λ_4. Thus, in terms of the original fields, the low-energy superpotential is W $=(A^2BC\Lambda_4^5)^{1/3}$, and the equations $\partial W/\partial A=\partial W/\partial B$ $=\partial W/\partial C=0$ then require the expectation values of A,B,C to vanish for a zero-energy vacuum.

B. Dynamical supersymmetry breaking

Having established that there are no flat directions in the model, we now proceed to show that the model breaks supersymmetry. We need only show this in a particular range of Λ_4, Λ_3 and Λ_2 (these are the strong coupling scales for the three new gauge groups.) Holomorphy in Λ_i/Λ_j ensures there will be no phase transitions as these couplings are varied; at worst there may be singular points for special values of Λ_i/Λ_j, which we can choose to avoid. Thus, if supersymmetry is broken for any open set of values for Λ_i/Λ_j, then it will be broken for most values of Λ_i/Λ_j.

Although we will eventually construct a model in which $\Lambda_3>\Lambda_4\gg\Lambda_2$, it is easiest to show supersymmetry is broken in the regime $\Lambda_4\gg\Lambda_3\gg\Lambda_2$. The large separation of scales allows us to treat the strong dynamics of the gauge groups one at a time, with the remaining weakly coupled groups (including the standard model) serving as spectators. First, the Sp(4) gauge group becomes strongly coupled at the scale Λ_4. It confines, and the low-energy dynamics is given in terms of the mesons $\bar{A}=(\overline{VV})/\Lambda_4$, $\bar{B}=(\overline{TV})/\Lambda_4$, $\bar{C}$ $=(\overline{TT})/\Lambda_4$, which are massless in the absence of a tree-level superpotential. Note that we have normalized the mesons to have dimension one, as is appropriate at low energy. The resulting matter content (aside from the standard model fields) is given in Table II.

The strong dynamics of the theory generates a dynamical superpotential [50]

$$W_{\rm dyn}=\frac{\overline{AABC}}{\Lambda_4}. \tag{5}$$

The superpotential of the theory is now

$$W=W_{\rm SM}+W_{3\text{-}2}+W_M+W_s+W_{\rm dyn}, \tag{6}$$

where $W_{\rm SM}$, $W_{3\text{-}2}$, and W_s are the same as before and

$$W_M=y_A\Lambda_4 A\overline{A}+y_B\Lambda_4 B\overline{B}+y_C\Lambda_4 C\overline{C} \tag{7}$$

is W_m reexpressed in terms of the composite fields. This last set of interactions results in masses of order Λ_4 for the fundamental fields A, B, and C and the composite mesons $\overline{A}$, $\overline{B}$, and $\overline{C}$. Without changing the infrared dynamics, we may integrate out the massive fields, eliminating W_M and W_s from the superpotential and changing the Kähler potential by high dimension operators. The 3-2 sector and standard model sector are then connected only by irrelevant interactions, and the former, as shown by Affleck, Dine and Seiberg, breaks supersymmetry at a scale determined by λ_0, Λ_2 and Λ_3.

IV. BETA FUNCTIONS AND ANOMALOUS DIMENSIONS

Our model exhibits a number of interesting and subtle strong-coupling phenomena, which we will discuss carefully. Because of this, we begin with a review of some dynamical relations in supersymmetric theories, which provide tools for semiquantitative analysis of strongly coupled theories. These tools will not be powerful enough to make our results unambiguous, but they will provide evidence that the model exhibits the qualitative features which we need to make use of.

In $\mathcal{N}=1$ supersymmetric theories there are relationships between beta functions and anomalous dimensions. A Yukawa coupling y_0 in the superpotential $W_Y=y_0\phi_1\phi_2\phi_3$ has a beta function which is a function of all the other Yukawa couplings y_i and gauge couplings g_k in the theory. Non-renormalization theorems in $\mathcal{N}=1$ supersymmetric theories ensure that all vertex functions are trivial and that all running of couplings comes through wave function renormalization. Consequently, the beta function of y_0 is related in a simple way to the anomalous mass dimensions $\gamma_n(y_0,y_i,g_k)$ of the fields ϕ_n:

$$\beta_{y_0}=\frac{1}{2}\,y_0[\gamma_1(y_0,y_i,g_j)+\gamma_2(y_0,y_i,g_j)+\gamma_3(y_0,y_i,g_j)]. \tag{8}$$

The running of gauge couplings is slightly more complicated, but still related linearly to the anomalous dimensions of the fields. According to [51,52] the coupling g_k runs as

$$\beta_{g_k}=-\frac{g_k^3}{16\pi^2}\,\frac{3C_2(G_k)-\Sigma_pT(\phi_p)[1-\gamma_{\phi_p}]}{1-(g_k^2/8\pi^2)C_2(G_k)}. \tag{9}$$

Here $C_2(G_k)$ is the second Casimir operator of the gauge group G_k for which g_k is the coupling, the sum in the numerator is over all matter fields ϕ_p, $T(\phi_p)$ is half the index of the representation of ϕ_p under G_k, and γ_{ϕ_p} is the anomalous dimension of ϕ_p. Note that to leading order in the couplings this expression gives

$$\beta_{g_k}=-\frac{g_k^3}{16\pi^2}\,b_0;\quad b_0\equiv 3C_2(G_k)-\sum_p T(\phi_p), \tag{10}$$

where b_0 is the well-known coefficient of the one-loop correction to the gauge coupling.

There are also conditions which follow from the $\mathcal{N}=1$ superconformal algebra, which constrains the properties of theories at exact or approximate fixed points. One of these is the ''unitarity condition.'' In any four-dimensional conformal field theory, unitarity implies that no gauge invariant operator (except the unit operator) can have dimension less than 1 [53–55]. Any operator whose dimension is exactly one must satisfy the Klein-Gordon equation, and cannot interact with any other fields. These facts will apply also to any operator which is gauge variant only under a very weakly coupled gauge group: as it must have dimension greater than one in the limit of zero gauge coupling, perturbation theory in the small gauge coupling ensures that its dimension cannot be much below 1. A related consequence is that when all gauge couplings are small, a field with a large Yukawa coupling always has a positive anomalous dimension.

Another condition relates the R charge of a chiral operator and its anomalous dimension [53–55]. At a superconformal fixed point there is a special R current that appears in the same superconformal multiplet as the energy-momentum tensor and the supersymmetry current. At a fixed point, the dimension of a chiral operator is $\frac{3}{2}$ times its R charge, from which its anomalous dimension may be calculated. (The result always agrees with results which follow from the beta functions above.) An important implication of this result is that, since R charges are additive, the dimension of a composite chiral operator is equal to the sum of the dimensions of its chiral constituents. This can be restated as resulting from the absence of short-distance singularities when any two chiral operators are brought to the same point. (Similar statements of course apply to antichiral operators.) Unfortunately, the general theory has an infinite set of R symmetries, and often it is impossible to determine which of them appears in the multiplet of currents.

An important corrolary of the above results involves the running of Yukawa couplings. Consider a set of fields with gauge couplings g and Yukawa couplings y with anomalous dimensions $\gamma(y,g)$. Unitarity ensures that $\gamma(y,0)$ is positive. It is easy to check that $\gamma(0,g)$ is negative for a charged field for small g. Together with Eq. (8), these imply that a Yukawa coupling involving charged fields will be irrelevant for $g\ll y$ but may become relevant as g becomes of order y. By contrast, a Yukawa coupling for three gauge-neutral fields is always irrelevant.

V. THE MODEL: A CAREFUL RENDERING

We now provide an detailed overview of the model, making claims about the dynamics which we justify later in this and in the following section. Our approach is semiquantita-

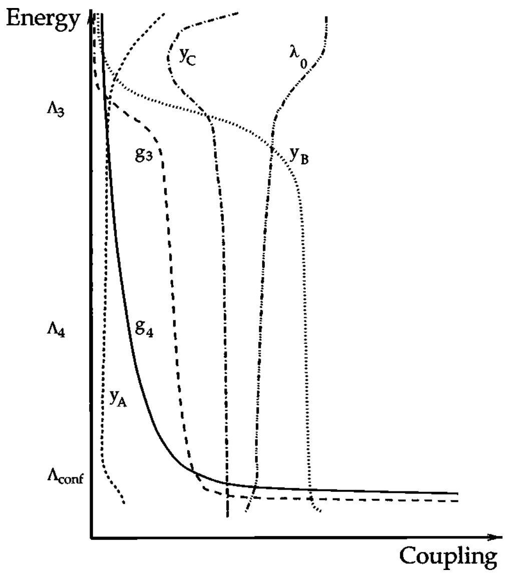

FIG. 1. Possible renormalization group flow for the most important couplings. The actual flow could be very different; this figure is for illustration only.

tive and relies on the dynamical relations described in Sec. IV.

We consider the model in the regime $\Lambda_3 > \Lambda_4 \gg \Lambda_2$. The standard model couplings and g_2 are smaller than the gauge couplings g_3, g_4, and can be neglected in most of our analysis. The Yukawa couplings $\lambda_0, y_\Lambda, y_B, y_C$ may not be small (though we do assume for simplicity that h_1, h_2 are small.) The renormalization-group flow of the gauge and Yukawa couplings cannot be known exactly, but can be analyzed using Sec. IV. A possible form for the behavior of the couplings is sketched in Fig. 1.

Before any gauge couplings are strong (assuming there is such a range) the Yukawa couplings are all irrelevant. Once g_3 becomes strong, however, λ_0, y_B, and y_C may become relevant; they are certainly relevant when they are small, and therefore they become or remain large. By contrast, y_A is irrelevant and may become small. A theory of SU(3) with seven triplets and antitriplets, along with some SU(3) singlets and Yukawa couplings, is expected to flow to a conformal field theory in the far infrared (see Sec. V A). The gauge coupling g_3 and the relevant Yukawa couplings will run slowly as they gradually approach an infrared fixed point. Our knowledge of the properties of this fixed point is limited. We know that it should preserve the global SU(5) symmetry. We also know that it occurs outside of perturbation theory, and so the theory begins to look conformal only at scales far below Λ_3. In fact the theory is unlikely to be extremely close to the fixed point when supersymmetry breaks, although SU(5) will still be a good approximate symmetry (see Sec. V F).

The Sp(4) group has a negative beta function when g_3 is small. The strong coupling effects of g_3 might in principle change the sign of β_{g_4}, but they do not, as shown in Sec. V B. Consequently, the Sp(4) coupling grows. As it becomes strong the coupling y_A will become relevant, though it may not have a large energy range in which to grow. The other relevant Yukawa couplings are expected to remain large. The coupling g_4 is not expected to reach a fixed point (a conspiracy would be needed to allow for this possibility, see Sec. V B) and so at some scale Λ_{conf}—a real physical scale, not to be confused with the holomorphic scale Λ_4—Sp(4) confines. Below this scale, meson degrees of freedom $\bar{A} = \overline{VV}/\Lambda_4$, $\bar{B} = \overline{VT}/\Lambda_4$, $\bar{C} = \overline{TT}/\Lambda_4$ best describe the long distance physics, and the theory possesses the fields of Table II. The factors of Λ_4 are convenient in order to have a holomorphic definition of these fields which has mass dimension one. However, an additional nonholomorphic dimensionless factor is necessary for these composite fields to be canonically normalized. While this factor cannot be computed, it can be estimated on physical grounds to be such that if y_A is large (of order 4π), then the mass of A is of order Λ_{conf}. Similar considerations normalize B and C.[1]

Since the couplings y_B and y_C are strong, the fields $B, \bar{B}$ and $C, \bar{C}$ have masses of order Λ_{conf}. This leaves the SU(3) group below this scale with only two triplets and a large beta function; g_3 blows up at once, at a scale of order Λ_{conf}. Since λ_0 is large, supersymmetry is broken immediately by SU(3) strong coupling effects, with $\sqrt{F}$ close to Λ_{conf} (see Sec. V C).[2] This vacuum is likely to preserve color and electromagnetism, as argued in Sec. V D.

The Yukawa coupling y_A may be driven small, as explained in Sec. V E, and the fields $A, \bar{A}$ therefore may have a supersymmetric mass somewhat smaller than Λ_{conf}. There are also light particles from the 3-2 sector: the goldstino and one other particle whose presence is required by anomaly matching [58]. Finally, there are the light fields of the standard model. We now would like to integrate out the dynamics above the scale Λ_{conf} and write an effective theory for the light degrees of freedom valid below this scale. However, the dynamics of the supersymmetry breaking is strongly coupled, and reliable quantitative analysis is impossible. A qualitative approach to this effective Lagrangian is therefore necessary. We will use naive dimensional analysis to estimate various quantities which cannot be computed [56,57]. Although there is no empirical evidence that this works for strongly coupled theories other than QCD, we expect such estimates to be off by no more than an order of magnitude.

The act of integrating out the supersymmetry-breaking sector introduces soft supersymmetry-breaking terms at the scale Λ_{conf}. Fields which couple strongly to the supersymmetry breaking, such as the composite scalars $\bar{A}$, will have supersymmetry-breaking masses of order $4\pi F/\Lambda_{conf}$ (see Sec. VI A). The $\bar{A}, A$ fermions obtain supersymmetric masses

[1]Note that the superpotential (7) is misleading; the masses of A, B, and C are of order Λ_{conf}, not Λ_4, as a result of this normalization factor, which appears in the Kähler potential.

[2]In an entirely strongly coupled theory, with no small dimensionless parameters, naive dimensional analysis [56,57] would give $F \sim \Lambda_{conf}^2/(4\pi)$. Since in our theory the weak coupling g_2 can suppress supersymmetry breaking, as shown in Sec. V D, we will keep F a free parameter.

of order $y_A\Lambda_{\rm conf}/(4\pi)$—chiral symmetry protects them against supersymmetry-breaking masses. The A scalars discover supersymmetry breaking only via the weak couplings y_A and the standard model gauge couplings $\alpha_k^{\rm SM}$, and hence (see Sec. VI B) have masses of order $\max[y_A\Lambda_{\rm conf}/(4\pi),\alpha_k^{\rm SM}F/\Lambda_{\rm conf}]$. All the MSSM fields couple to the supersymmetry-breaking sector via standard model gauge couplings, as in usual GMSB models. A modified version of the usual results [7–9] applies—both gauginos and scalars acquire masses of order $\alpha_k^{\rm SM}F/\Lambda_{\rm conf}$ (see Sec. VI B). Since $\sqrt{F}$ lies close to $\Lambda_{\rm conf}$, we take $\Lambda_{\rm conf}\sim 30-1000$ TeV and $\sqrt{F}\sim 10-100$ TeV. Such low values of $\sqrt{F}$ lead to rapid decays of the next-to-lightest superpartner to a gravitino, as discussed in Ref. [10–24].

The large splittings in the messenger multiplets lead to substantial superoblique corrections [59,60]; these are estimated in Sec. VI B. Trilinear scalar couplings and the μ problem—a difficulty in this model—are discussed in Sec. VI C. Finally, in Sec. VI D, we explain a couple of ways to avoid stable messenger fields, as required to prevent conflict with experiment.

A. The SU(3) physics

It has been argued that the theory of SU(3) with seven triplets and antitriplets reaches an infrared fixed point [61]. At this fixed point the charged fields have negative anomalous dimensions, so Yukawa couplings involving either two charged fields and a singlet or three charged fields (such as λ_0, y_B, y_C) are relevant and drive the theory away from this fixed point. Does the theory flow to another fixed point in the infrared, or behave altogether differently? It is possible to construct theories, with relevant Yukawa couplings, which preserve no R charges that could be part of a unitary conformal field theory; in such cases a low-energy fixed point is unlikely (though not impossible, since accidental R symmetries may arise in the infrared.) However, in our theory, there are R charges, and corresponding candidate infrared fixed points, which are consistent with unitarity and with the Yukawa couplings of the superpotential of Eq. (1). These would preserve the SU(5) global symmetry which contains the standard model. On the other hand, there is insufficient flavor symmetry to determine the R charges at the fixed point and confirm that unitarity is not violated. Thus we cannot provide an argument that the couplings we have chosen do or do not lead to a fixed point.

However, for our present purposes, such an argument is not really necessary. Even if the SU(3) theory does not reach an approximate fixed point, it is likely that its beta function will be very small, much smaller than the one-loop estimate. [This is certainly true if the Yukawa couplings are small, due to SU(3) two-loop corrections.] A slow-running coupling constant tends to wash out physical effects involving the scale Λ_3. Meanwhile, in the effective theory below the scale of Sp(4) confinement, the SU(3) beta function is rather large, so for a substantial range of Λ_3 and Λ_4, the physical scales associated with strong Sp(4) and SU(3) dynamics are close together, and close to the supersymmetry breaking scale. Furthermore, if SU(5) is a roughly approximate symmetry at high energy, it will be preserved even when the SU(3) coupling becomes large. We will explain how this occurs in Sec. V F.

B. The Sp(4) beta function

We need to show that $\beta_{g_4}<0$ at all scales, so that Sp(4) confines rather than reaching a conformal fixed point. When g_4 is small, this can be proven. When g_4 is large, a proof is not possible, but we will show it is unlikely that β_{g_4} reaches a zero.

The beta function for Sp(4) is given by

$$\beta_{g_4}=-\frac{g_4^3}{16\pi^2}\,\frac{9-(3/2)[1-\gamma_{\bar T}]-(5/2)[1-\gamma_{\bar V}]}{1-g_4^2/2\pi^2}$$
$$=-\frac{g_4^3}{16\pi^2}\,\frac{5+(3/2)\gamma_{\bar T}+(5/2)\gamma_{\bar V}}{1-g_4^2/2\pi^2}. \qquad (11)$$

Above Λ_3 the one-loop formula is approximately correct,

$$\beta_{g_4}=-5\,\frac{g_4^3}{16\pi^2}, \qquad (12)$$

and g_4 grows logarithmically with coefficient 5. However, once g_3 is large we expect that $\bar T$ has a negative anomalous dimension, and once y_B is large then $\bar V$ will have a positive anomalous dimension (by unitarity). The effect of $\gamma_{\bar T}$ will tend to make g_4 run more slowly, while that of $\gamma_{\bar V}$ will tend to make it run more quickly.

The operator $\bar T\bar T\bar T$ is charged only under Sp(4). When $g_4\ll 1$, the unitarity condition demands that $\bar T$ have dimension close to or greater than 1/3, and thus $\gamma_{\bar T}\geq -4/3-\mathcal{O}(g_4^2)$. It follows that the function $5+(3/2)\gamma_{\bar T}+(5/2)\gamma_{\bar V}$, which appears in the numerator of the beta function, is greater than or of order 3 in this regime. Thus g_4 remains asymptotically free, but may run more slowly than the one-loop estimate, leading $\Lambda_{\rm conf}$ to be much lower than Λ_4.

For large g_4 a different argument is necessary. Once g_4 becomes large, then, in the limit $g_2=g_i^{\rm SM}=0$, $\overline{VV}$ and $\overline{VTTT}$ are gauge invariant, implying $\gamma_{\bar V}>-1$, $3\gamma_{\bar V}+5\gamma_V>-10$. Unfortunately this allows the Sp(4) beta function to be arbitrarily close to zero, and it could reach zero when the gauge couplings $g_2, g_i^{\rm SM}$ are nonzero or when other small SU(5) violation is accounted for. While we cannot rule this out, we note that it requires extreme values of $\gamma_{\bar T}$ which are unlikely to be attained. It is more likely that the Sp(4) beta function remains negative, though rather smaller in magnitude than normally expected for a theory undergoing confinement. This suggests that $\Lambda_{\rm conf}\ll\Lambda_4$, which is of relevance in Sec. V E below.

C. Supersymmetry breaks at the scale $\Lambda_{\rm conf}$

Can we be sure that the SU(3) coupling blows up close to the scale $\Lambda_{\rm conf}$ and not well below that scale? Consider the case $\Lambda_4\gg\Lambda_3$, with small λ_0, y_B, y_C. In this case, the Sp(4) theory confines at the scale Λ_4, leaving the theory with the fields in Table I. The SU(3) beta function coefficient b_0 changes from 2 to 3. The fields $B,\bar{B},C,\bar{C}$ have masses of order $y_B\Lambda_4, y_C\Lambda_4$; below these scales, $b_0=7$. The one-loop analysis for the SU(3) coupling is reasonably good in this case, and we find the coupling blows up at

$$\Lambda_{3L}=\Lambda_4[y_B^5 y_C]^{1/7}[\Lambda_4/\Lambda_3]^{2/7}. \tag{13}$$

The small fractional powers indicate that the Sp(4) dynamics controls the divergence of the SU(3) coupling even when g_3 is small at Λ_4. We expect this will be all the more true when y_B, y_C are large and $\Lambda_3>\Lambda_4$, in which case the SU(3) coupling will be substantial, and slow running, down to the scale $\Lambda_{\rm conf}$.

Since the strong SU(3) physics drives supersymmetry breaking, and since no couplings (except possibly y_A, see Sec. V E) are small, we expect supersymmetry breaking to occur at a mass scale $\sqrt{F}$ within an order of magnitude of $\Lambda_{\rm conf}$. We will show in Sec. V D that for sufficiently small g_2 there is a supersymmetry-breaking vacuum with $F \sim g_2^{3/7}\Lambda_{\rm conf}^2$ that preserves the standard model gauge group, and we will assume the theory lies in this vacuum even for $g_2\sim 1$.

D. An acceptable supersymmetry-breaking vacuum

Although we have shown supersymmetry is broken, this is far from showing that the vacuum of the theory is phenomenologically acceptable. In particular, the standard model gauge groups must not be broken.

For Λ_2 sufficiently small and $\Lambda_4>\Lambda_3$, there is a supersymmetry-breaking vacuum in which no field with standard model charges gets an expectation value. First, recall that when $\Lambda_4\to\infty$ the low-energy theory is the 3-2 model with massive messengers, and so when $\Lambda_2\to 0$, supersymmetry is restored, even for finite Λ_4. When $\Lambda_2=0$ there are flat directions labeled by the operators $q\bar{u}, q\bar{d}, l$ which carry only SU(2) quantum numbers. Classically, the superpotential sets $q\bar{u}$ to zero, but $q\bar{d}$ and l may still have expectation values, both of which may be nonzero as long as $ql \equiv q^i l^j \epsilon_{ij}=0$. This allows the expectation values

$$q=\begin{bmatrix} a & 0 & 0 \\ 0 & 0 & 0\end{bmatrix}; \quad \bar{d}=[\bar{a}\;0\;0]; \quad l=\begin{bmatrix} b \\ 0\end{bmatrix}, \tag{14}$$

where $|a|=|\bar{a}|$. Along this flat direction the Sp(4) gauge group is unbroken while the SU(3) gauge group breaks to $SU(2)_I$ with five massless flavors (the fields $q^1,\bar{d}$ are eaten and have mass $|g_3 a|$, while $\bar{u},l$ have mass $\lambda_0 b$; $B,\bar{T},C$ are massless.) The strong coupling scale of this group is $\Lambda_I = \Lambda_3^2\lambda_0\langle l/q\bar{d}\rangle$. Take $a,b\Lambda_4>\Lambda_3>\Lambda_I$; then the confining dynamics of Sp(4) gives mass to the messengers leaving the $SU(2)_I$ theory with no massless flavors and a strong coupling scale $\Lambda_L^6=y_B^5 y_C\lambda_0\Lambda_3^2\Lambda_4^5\langle l/q\bar{d}\rangle$. The low-energy superpotential is just given by the gaugino condensation in these variables:

$$W_L=\left[\frac{y_B^5 y_C\lambda_0\Lambda_3^2\Lambda_4^5\langle l\rangle}{\langle q\bar{d}\rangle}\right]^{1/2}. \tag{15}$$

Since $a,b\gg\Lambda_I$, the Kähler potential for these fields is trivial, and so the potential energy along this direction is

$$V_I(a,b)=|y_B^5 y_C\lambda_0\Lambda_3^2\Lambda_4^5|\left[\frac{1}{4|ba^2|}+\frac{|b|}{2|a|^4}\right]. \tag{16}$$

This is minimized at $a,b\to\infty$ with $|b/a^4|\to 0$.

We have assumed to this point that $\Lambda_2=0$. For $\Lambda_2 \ll\Lambda_L$, the only effect of the gauged SU(2) will be the potential energy from the D terms

$$V_2(a,b)=g_2^2(|a^2|+|b^2|)^2, \tag{17}$$

where g_2 is the SU(2) coupling at energy scales of order a,b. The minimum of the potential V_1+V_2 can be seen to be at $a\sim b\sim[g_2^{-2}y_B^5 y_C\lambda_0\Lambda_3^2\Lambda_4^5]^{1/7}\gg\Lambda_3,\Lambda_4$, where the vacuum energy density is of order

$$F^2=g_2^{6/7}[y_B^5 y_c\lambda_0\Lambda_3^2\Lambda_4^5]^{4/7}. \tag{18}$$

Notice that a,b go to infinity and F goes to zero as $g_2\to 0$, as expected; thus our assumption that $a,b\gg\Lambda_3$ is consistent for small g_2. The standard model gauge group is not broken in this vacuum, which, for sufficiently small g_2, is certainly the true minimum of the potential.

We have therefore found that our model's supersymmetry-breaking vacuum preserves the standard model gauge group when $\Lambda_4>\Lambda_3\gg\Lambda_2$. However, the regime of interest, $\Lambda_3>\Lambda_4\gg\Lambda_2$, is not calculable. While we cannot prove that the above vacuum continues to be stable (or sufficiently metastable) into this regime, it is reasonable to assume that it does so for some range of parameters, and that for $g_2\sim 1$ it leads to $F\sim\Lambda_{\rm conf}^2$. Nonetheless, since supersymmetry is restored for $y_A=0$, and the dynamics of the theory tend to drive y_A small at low energy, the effects of this parameter deserve further discussion.

E. Why y_A is small

We claimed that y_A is likely to be small, leading $\bar{A},A$ to be light. In this section, taking an idealized limit, we clarify why y_A is driven small, and explain why we cannot estimate its size. We also explain why we cannot determine a lower bound on y_A from the requirement that the global vacuum be phenomenologically acceptable.

In particular, suppose that the couplings g_3,λ_0,y_B,y_C are all large at the Planck scale m_P and chosen so that the theory reaches the (conjectured) infrared fixed point discussed in Sec. V A at energies near m_p. Suppose also that $g_4=g_i^{\rm SM}=0$. We do not know the anomalous dimensions $\gamma_{\bar{V}}$, γ_A, but the contribution from the couplings y_A, y_B are inevitably positive by unitarity (Sec. IV). In particular, $\gamma_{\bar{V}}$ need not be small in magnitude even when y_A is very small. Consequently y_A is irrelevant [see Eq. (8)], and is driven small as a power (most likely less than 1) of the energy:

$$y_A(\mu) \lesssim y_A(m_P)\left(\frac{\mu}{m_P}\right)^{\gamma_V^*}, \tag{19}$$

where γ_V^* is $\gamma_{\bar{V}}$ at the SU(3) fixed point (at which $y_A=0$).

Accounting for g_4, $g_i^{\rm SM}\neq 0$ makes very little difference until g_4 becomes substantial near Λ_4, making the anomalous dimension of $\bar{V}$ smaller and probably negative. The Sp(4) physics will cause y_A to change by some unknown factor κ_4, probably greater than one. As noted in Sec. V B, the regime in which g_4 is large may be extended by a small beta function, leading κ_4 to be larger than naively expected. But because the SU(3) physics is nearly scale invariant, and y_A is negligibly small at low energies, it follows that for Λ_4 and $\Lambda_{\rm conf}$ much smaller than m_P, the confining physics of the Sp(4) theory is independent of Λ_4/m_P. In turn this implies that κ_4 is independent of Λ_4 for small Λ_4. The low-energy value of y_A is thus a monotonic decreasing function of Λ_4:

$$y_A(\Lambda_{\rm conf}) \lesssim y_A(m_P)\left(\frac{\Lambda_4}{m_P}\right)^{\gamma_V^*}\kappa_4. \tag{20}$$

By varying the couplings at m_P, maintaining only $\Lambda_3 > \Lambda_4$, we may easily obtain any low-energy value of y_A that we want, as long as it is not very large. We therefore have no prediction for y_A (especially as κ_4 cannot be calculated,) and so we will always consider it a free parameter, which we expect to be smaller than unity.

The one additional concern one might have is that since supersymmetry is restored for $y_A=0$, due to the flat direction $A\bar{V}\bar{V}$ which breaks the standard model gauge group, perhaps the global minimum of the potential will break the standard model gauge group if y_A is driven small. However, as we will now see, it cannot be shown that this occurs, and so no conclusions may be drawn.

Specifically, in the limit $y_A=0$, the model develops a classical flat direction $A\bar{V}\bar{V}$. This flat direction is not lifted quantum mechanically, as can be seen through the following argument. The expectation value for $\bar{V}$ gives mass to four triplets in B and four antitriplets in $\bar{T}$, but leaves one triplet B^5 behind. The SU(3) gauge theory thus has three massless flavors q, B^5, $\bar{u}$, $\bar{d}$, C. This gauge group has a quantum modified moduli space, and so at large expectation values the classical moduli space is not modified. The theory therefore has vacua at infinite expectation values for A, $\bar{V}$ and the massless SU(3) fields, with a potential energy which is essentially flat at large vacuum expectation values (VEVs).

If y_A is nonzero and small, however, a quartic potential for A and $\bar{V}$ lifts these nearly flat directions. With the assumption of large $\langle A\rangle$ and $\langle\bar{V}\rangle$, there is no vacuum, even for arbitrarily small y_A. Therefore, $\langle A\rangle$ and $\langle\bar{V}\rangle$ must be small; but will they be zero in the vacuum, as required phenomenologically? We can look for a minimum at very small expectation values; in this case the confining of Sp(4) and the breaking of supersymmetry occurs as in Sec. V D, with the only effect of small y_A to make A, $\bar{A}$ massive. Locally, then, there is a minimum with $\langle A\rangle=\langle\bar{V}\rangle=0$. Unfortunately, we cannot be sure this is a global minimum; there could be a minimum for $\Lambda_{\rm conf}<\langle\bar{V}\rangle<\Lambda_3$, where strong coupling effects make a reliable computation impossible.

Note that our model involves a slightly different situation. We are not assuming that y_A is small at high scales, only that it may be driven small at low energy. The flat direction at large $\langle\bar{V}\rangle$ is therefore strongly lifted. Still, it is likely that there is a lower bound on the low-energy value of y_A below which the standard model may be broken in the global minimum of the potential. As this bound cannot be calculated, and depends on g_2, we leave the low-energy value of y_A as a free parameter without a lower bound.

Is it possible, even when the standard model gauge group is conserved, that the supersymmetry-breaking scale F depends on y_A, and perhaps becomes much less than $\Lambda_{\rm conf}$ when $y_A \ll 1$? This seems unlikely to us. We have shown there is a acceptable supersymmetry-breaking vacuum whose value for F is controlled by g_2 [Eq. (18)]; in the limit of small g_2, finite y_A, this vacuum is the true minimum. In the limit of small y_A, finite g_2, the conjectured vacuum discussed in this section, which breaks the standard model, might be the true minimum; its energy density would be proportional to a power of y_A. We would expect there to be a first-order transition between these two vacua, and a corresponding discontinuous transition in the functional dependence of F on g_2 and y_A. In the allowed minimum of Sec. V D, the y_A dependence of F is subleading; conversely, if F depends strongly on y_A, then the theory is probably in the unacceptable vacuum. Thus, although we cannot determine the range of g_2 and y_A for which the theory prefers the acceptable vacuum, we can assume that when it does so, the supersymmetry-breaking scale $\sqrt{F}$ depends very little on y_A and remains of order $g_2^{3/14}\Lambda_{\rm conf}$.

F. The SU(5) global symmetry and coupling constant unification

We have added extra matter to the standard model, and thus run the risk that we will drive the standard model couplings to a Landau pole below the string or Planck scale. Furthermore, we have added additional interactions which do not exactly satisfy SU(5) relations and which can become strong. In this section we confirm that unification at finite coupling is possible, and that SU(5) can be preserved as an approximate global symmetry, provided that there are no strong violations of SU(5) at the Planck scale.

First, we consider the limit in which SU(5) is preserved by the superpotential Eq. (1), and we verify that the standard model couplings can unify at finite coupling. We have added five triplets and antitriplets to $\mathrm{SU}(3)^{\rm SM}$ in Table I, which change the perturbative QCD beta function coefficient b_0 from $+3$ to -2. These fields disappear from the theory at the scale $\Lambda_{\rm conf}\sim 100$ TeV, except for the fields ψ_A, $\psi_{\bar{A}}$, A. A one-loop analysis indicates this is a borderline case, while a two-loop analysis ignoring the Yukawa couplings in Eq. (1) would suggest that the standard model gauge couplings hit a Landau pole below the Planck scale. However, the situation is modified by the Yukawa couplings.

In particular, the QCD beta function above $\Lambda_{\rm conf}$ reads

$$\beta_{g_3^{\rm SM}}=-\frac{(g_3^{\rm SM})^3-2+(3/2)\gamma_A+(3/2)\gamma_B+2\gamma_{\bar V}}{1-3(g_3^{\rm SM})^2/8\pi^2}. \quad (21)$$

(For simplicity of discussion, we ignore in this expression the anomalous dimensions of standard model fields, as generated by standard model gauge and Yukawa couplings.) At energy scales well above $\Lambda_{\rm conf}$ but below Λ_3, where g_3 is large, the expectation is that γ_A is positive, $\gamma_{\bar V}$ is positive, and γ_B is negative. If y_B were small, $\gamma_{\bar V}$ and γ_A would be very small and $g_3^{\rm SM}$ would run faster than the one-loop analysis would suggest. However, because y_B is large, and because y_A may be large at high scales, γ_A and $\gamma_{\bar V}$ need not be small, and likely make $g_s^{\rm SM}$ run more *slowly* than the one-loop analysis. Once we approach $\Lambda_{\rm conf}$ and g_4 becomes large, the analysis is even less under control; no argument can be constructed indicating either that $g_3^{\rm SM}$ must run faster or slower than perturbatively indicated. Finally, below $\Lambda_{\rm conf}$, the QCD beta function becomes negative, although the scalar A and the fermions ψ_A, $\psi_{\bar A}$ will reduce the QCD beta function slightly below its MSSM value. Altogether the uncertainties in the anomalous dimensions prevent us from demonstrating unambiguously that $g_3^{\rm SM}$ remains finite, but since the theory at one loop is marginally acceptable, there likely exists a region of parameter space in which this is the case.

Suppose that $g_3^{\rm SM}$ does not reach a Landau pole; what about the other standard model couplings? It is well-known that addition of complete SU(5) multiplets to the standard model does not ruin coupling constant unification. This is true at one loop in the standard model couplings and to all orders in other couplings. The proof in a supersymmetric theory is direct. As seen in Eq. (9), to leading order in a weak coupling constant the beta function is proportional simply to g^3 times $3C_2(G_k)-\Sigma_p T(\phi_p)[1-\gamma_{\phi_p}]$. The usual statement of coupling constant unification is that a complete SU(5) multiplet $\{\phi_j\}$ preserves unification because $\Sigma_j T(\phi_j)$ is the same for each standard model group factor, leading to equal shifts in $b_0=3C_2(G_k)-\Sigma_p T(\phi_p)$ for the three groups and preserving both unification and the unification scale. In this case, the SU(5) multiplets have large anomalous dimensions due to strong interactions involving the 4-3-2 sector. However, since the fields $\{\phi_j\}$ all have the same anomalous dimension $\gamma_{\{\phi_j\}}$ [by approximate SU(5) flavor symmetry] the sum $\Sigma_j T(\phi_j)[1-\gamma_{\phi_j}]=[1-\gamma_{\{\phi_j\}}]\Sigma_j T(\phi_j)$ is essentially the same in each standard model group factor. Again, unification is preserved.

Thus, if $g_3^{\rm SM}$ does not hit a Landau pole, neither will $g_2^{\rm SM}$ or $g_1^{\rm SM}$. Furthermore, coupling constant unification will be preserved despite the strong coupling effects.

Now, we consider the possibility that SU(5) is broken in the superpotential Eq. (1). In this case the multiplets $\bar V$, A, B must be broken up into their component multiplets under the standard model gauge group, each with its own anomalous dimension. As an simplified example, suppose that y_A is very small and that we can ignore it and the field A. Consider the fields B, $\bar V$, which contains as submultiplets a color triplet and antitriplet B_3, $\bar V_3$ and weak isodoublets B_2, $\bar V_2$. Let us consider the effect on the coupling $y_B B\overline{TV}$, which now becomes two couplings $y_3 B_3 \overline{TV}_3+y_2 B_2\overline{TV}_2$. The anomalous dimensions of the relevant fields are given at one loop by

$$16\pi^2\gamma_{B_2}\approx 2y_2^2-\frac{16}{3}g_3^2,\quad 16\pi^2\gamma_{B_3}\approx 2y_3^2-\frac{16}{3}g_3^2,$$

$$16\pi^2\gamma_{\bar V_2}\approx\frac{3}{2}y_2^2-\frac{15}{2}g_4^2,\quad 16\pi^2\gamma_{\bar V_3}\approx\frac{3}{2}y_3^2-\frac{15}{2}g_4^2, \quad (22)$$

$$16\pi^2\gamma_T\approx\frac{1}{2}(3y_3^2+2y_2^2)-\frac{16}{3}g_3^2-\frac{15}{2}g_4^2.$$

The beta functions for y_2 and y_3 are

$$\beta_{y_2}=\frac{1}{2}y_2(\gamma_{B_2}+\gamma_{\bar V_2}+\gamma_{\bar T});\quad \beta_{y_3}=\frac{1}{2}y_3(\gamma_{B_3}+\gamma_{\bar V_3}+\gamma_{\bar T}). \quad (23)$$

Consider the ratio $r=y_2/y_3$; its beta function can be written

$$\frac{1}{r}\frac{\partial r}{\partial\mu}=(\gamma_{B_2}-\gamma_{B_3}+\gamma_{\bar V_2}-\gamma_{\bar V_3})\approx\frac{7}{64\pi^2}y_2y_3\left(r-\frac{1}{r}\right). \quad (24)$$

Thus, if the product of the Yukawa couplings y_2, y_3 is small, both couplings will grow with the ratio r remaining fixed. However, the effect of the Yukawa couplings on the r beta function will cause r eventually to relax toward one. We see then that when the y_B couplings become large, as we expect them to be at low energy, SU(5) violation tends to be driven small.

A similar analysis shows that the couplings denoted y_A are also driven toward SU(5) universality if they are large, though not if they are small. Either way, the effects of SU(5) violation are not large and will not prevent unification of standard model gauge couplings.

In conclusion, this model probably allows the unification of standard model couplings. All strong couplings will be nearly SU(5) preserving as a result of strong dynamics; weak couplings may violate this global symmetry. We will see some physical consequences of this symmetry below.

VI. BELOW THE SCALE OF SUPERSYMMETRY BREAKING

In this section we discuss various predictions and interesting features of the model which are relevant for energies in the TeV region and below, including mass relations between sfermions and the A, ψ_A, $\psi_{\bar A}$ fields, large superoblique corrections, and possible decay modes for the A, ψ_A, $\psi_{\bar A}$ particles.

A. Spectrum of the light non-standard-model fields

It is helpful first to consider the limit where standard model gauge couplings and y_A are taken to be arbitrarily small (ignoring the appearance of a supersymmetry-preserving vacuum at $y_A=0$.) All the standard model fields and A, ψ_A are decoupled. The only interacting light fields are $\psi_{\bar{A}}$ and the light fields in the supersymmetry-breaking 3-2 sector. Here we discuss their interactions and properties.

For simplicity, we will assume that the 3-2 sector, which breaks supersymmetry, does not break its U(1)-''hypercharge'' flavor symmetry. This is true for small λ_0 [58] but may not be true for λ_0 large. With unbroken hypercharge, the 3-2 sector has two light particles after supersymmetry breaking: the Goldstino G, which is eaten by the gravitino and obtains a mass $F/(\sqrt{3}m_P)$, and a fermion η required to saturate the ''hypercharge'' anomalies. We will refer to the superfield containing η as S_η, and the superfield containing the Goldstino as X.[3] If instead the hypercharge symmetry is broken, there will be a massless Goldstone boson to replace η. However, the difference is irrelevant for present purposes, as its effect on phenomenology of standard-model-charged particles is limited to the decays of the heavy fields B, $\bar{B}$, which have mass of order $\Lambda_{\rm conf}$.

Before getting into the discussion of the specific scalar masses in our model, we note that supersymmetry-breaking scalar mass-squared terms are of two types, ''holomorphic'' and ''nonholomorphic,'' that are of differing size in the limit of small supersymmetry breaking. Holomorphic scalar mass terms are defined to be those which couple scalar fields of the same chirality. In the limit $F\ll\Lambda_{\rm conf}^2$, supersymmetry-breaking nonholomorphic scalar mass terms for a generic strongly coupled superfield Φ come from terms in the Kähler potential of the form

$$\frac{16\pi^2}{\Lambda_{\rm conf}^2}\int d^4\theta\Phi^\dagger\Phi X^\dagger X, \tag{25}$$

which will give Φ a nonholomorphic (of type $\Phi\Phi^\dagger$) supersymmetry-breaking mass squared of order $(4\pi F/\Lambda_{\rm conf})^2$. Holomorphic supersymmetry-breaking scalar mass terms (of type Φ^2) can also be induced. To see this, one can minimize the scalar potential and solve for the Φ auxiliary field in the presence of Kähler terms in the effective theory such as

$$\frac{4\pi}{\Lambda_{\rm conf}}\int d^4\theta\Phi^\dagger\Phi(X^\dagger+X), \tag{26}$$

and effective superpotential terms such as

$$\int d^2\theta m\Phi^2+\text{H.c.} \tag{27}$$

One then finds a scalar mass-squared term in the potential

$$\approx\frac{4\pi m}{\Lambda_{\rm conf}}F\Phi^2+\text{H.c.} \tag{28}$$

The superfield S_η is a participant in the strong coupling dynamics of the 3-2 sector. Naively one might guess that its scalar component gets a supersymmetry-breaking mass-squared of order F. In fact, a more careful analysis, using a supersymmetric effective Lagrangian, shows that for $\sqrt{F}\ll\Lambda_{\rm conf}$ its mass is much smaller than this, because the unbroken U(1) symmetry prevents the scalar from getting a holomorphic mass term. The mass squared of the S_η scalar thus gets only a nonholomorphic contribution, of order $(4\pi F/\Lambda_{\rm conf})^2$.

The field $\bar{A}$ is composite at the scale $\Lambda_{\rm conf}$, and so its scalar component also gets a nonholomorphic mass squared $m_{\bar{A}}^2$ of order $(4\pi F/\Lambda_{\rm conf})^2$. We will *assume* this mass squared is positive; since it is induced through strong coupling it cannot be computed.

Now let us consider turning on y_A. This gives the $A,\bar{A}$ multiplets a supersymmetry preserving mass, of size $m_\Psi\sim|y_A|\Lambda_{\rm conf}/(4\pi)$. In particular, ψ_A and $\psi_{\bar{A}}$ become a Dirac fermion Ψ_A. Also induced are a holomorphic supersymmetry-breaking mass squared m_{AH}^2 and a nonholomorphic supersymmetry-breaking mass squared m_A^2 for the scalar A. Their sizes are set by the following consideration: all supersymmetry-breaking interactions involving A are mediated through its coupling to $\bar{A}$, and are therefore suppressed by one power of $y_A/(4\pi)$ for each A on an external leg. We find therefore that $m_{AH}^2\sim y_AF$ and $m_A^2\sim(y_AF/\Lambda_{\rm conf})^2\sim(y_A/4\pi)^2m_{\bar{A}}^2$.

Next, when the standard model gauge couplings are made nonzero, the gauge bosons lead to a conventional positive gauge-mediated contribution $m_{\rm GM}^2$ to the $A,\bar{A}$ scalar masses squared, of order $(\alpha_k^{\rm SM}F/\Lambda_{\rm conf})^2$. The $A,\bar{A}$ scalar mass-squared matrix thus has the form

$$\begin{array}{cc} & \begin{array}{cc} A & \bar{A}^* \end{array} \\ \begin{array}{c} A^* \\ \bar{A} \end{array} & \begin{pmatrix} m_\Psi^2+m_A^2+m_{\rm GM}^2 & m_{AH}^2 \\ m_{AH}^2 & m_\Psi^2+m_{\bar{A}}^2+m_{\rm GM}^2 \end{pmatrix} \end{array}. \tag{29}$$

For small y_A, one linear combination of the scalars, which is mostly A, is relatively light.

The experimental signatures of the Ψ_A and A scalar depend on which one is lighter. When y_A is sufficiently small, i.e., for $y_A\Lambda_{\rm conf}/(4\pi)\ll\alpha_k^{\rm SM}F/\Lambda_{\rm conf}$, the fermion Ψ_A is lightest member of the $A,\bar{A}$ multiplet. Otherwise, it is not possible to say which is lighter. Note that the masses of the $B,\bar{B}$ multiplet have a similar form to those of $A,\bar{A}$, except that we expect y_B to be large, and so none of these particles will be light.

B. The message of supersymmetry breaking

In conventional models of GMSB, violations of supersymmetry in the MSSM sector can be reliably computed from diagrams containing loops of particles carrying ordi-

[3]Note that the F term of X is just F, the Goldstino decay constant.

nary gauge charges. In our model, one might be tempted to compute superpartner masses by considering loops containing the $A,\bar{A}$ and $B,\bar{B}$ fields. However, unless a theory has messengers with a canonical Kähler potential and a mass-squared matrix with vanishing supertrace, the squark and slepton masses come from ultraviolet divergent diagrams and are sensitive to the high-energy, strongly coupled physics. The squark, slepton and gaugino masses are computable in the limit $\Lambda_4 \gg \Lambda_3 > \Lambda_2$ and $y_A, y_B, y_C, \lambda_0 \ll 1$. In this limit, all the fields carrying standard model gauge charges are effectively weakly coupled below the scale Λ_4. In contrast, our model has some large Yukawa couplings and $\Lambda_3 > \Lambda_4$. Here, the situation is not so straightforward, even if y_A and y_B are small and the $A,\bar{A},B,\bar{B}$ messengers are much lighter than $\Lambda_{\rm conf}$.

To see the limitations of perturbative computation, consider the limit y_A, $y_B \ll 1$ and $\sqrt{F} \ll |y_A|\Lambda_{\rm conf}$, $|y_B|\Lambda_{\rm conf} \ll \Lambda_{\rm conf}$. In this case the nonholomorphic mass terms for the fields $A,\bar{A},B,\bar{B}$ are suppressed, as is the supertrace of their mass matrices, and all members of these supermultiplets appear as effectively weakly coupled fields below the scale $\Lambda_{\rm conf}$, with masses of order $y_A\Lambda_{\rm conf}/(4\pi)$, $y_B\Lambda_{\rm conf}/(4\pi)$. As we discussed in the previous section, holomorphic supersymmetry-breaking scalar mass squared terms m_{AH}^2, m_{BH}^2, of type $A\bar{A}$, $B\bar{B}$ respectively, appear in the low-energy effective theory, with $m_{AH}^2 \sim y_A F$ and $m_{BH}^2 \sim y_B F$. The nonholomorphic supersymmetry-breaking scalar mass terms are of order $(4\pi F/\Lambda_{\rm conf})^2$ and are relatively suppressed. In this limit, which we will call the ''holomorphic'' limit, the contributions of the A and B messengers to the ordinary superpartner masses are positive and of order $\alpha_k^{\rm SM} F/\Lambda_{\rm conf}$, and can be perturbatively computed once their masses are known. Note that these contributions are independent of y_A and y_B. But this is not the whole story. Near the scale $\Lambda_{\rm conf}$, there could also be a host of broad resonances, with standard model quantum numbers and supersymmetry-breaking holomorphic mass-squared terms of order $4\pi F$, which give an equally large contribution to the superpartner masses. Similar conclusions can be reached by applying naive dimensional analysis to graphs involving all the fundamental strongly coupled fields. It is therefore not appropriate to regard the $A,\bar{A}$ and $B,\bar{B}$ superfields as the only messengers—the message is carried by the supersymmetry breaking sector as a whole.

We do not expect the holomorphic limit to apply, however. We expect that y_A is small, so the holomorphic mass-squared term m_{AH}^2, which is proportional to y_A, does not dominate over the nonholomorphic term $m_{\bar{A}}^2$, which is proportional to $(4\pi F/\Lambda_{\rm conf})^2$. When the supertrace of the messenger masses squared does not vanish, the contribution of the messengers to the squark and slepton masses squared is logarithmically divergent in the low-energy effective theory. This divergence is cut off in the full theory, but logarithms of the ratios of messenger masses to $\Lambda_{\rm conf}$ can appear in the squark and slepton masses.

By considering the contribution of loops containing the $A,\bar{A}$ multiplet to ordinary sfermion masses squared [62,35], we can show that the regime with $(y_A/4\pi)\Lambda_{\rm conf}^2 \ll 4\pi F \ll \Lambda_{\rm conf}^2$ is ruled out because of a large negative logarithm. We will refer to this region of parameter space as ''log dominated.'' A perturbative calculation of the loop contribution from the $A,\bar{A}$ multiplet to squark and slepton masses squared [62,35], for $(y_A/4\pi)\Lambda_{\rm conf}^2 \ll 4\pi F$, is of order $\alpha_k^{\rm SM^2}(F^2\Lambda_{\rm conf}^2)\log(4\pi F/\Lambda_{\rm conf}^2)$ and is negative. The uncertainty in this calculation is of the same size as the effects from the rest of the strongly interacting sector and is of order $\alpha_k^{\rm SM^2}(F_2/\Lambda_{\rm conf}^2)$. For $4\pi F \ll \Lambda_{\rm conf}^2$ the log enhances the negative contribution of the $A,\bar{A}$ multiplet to squark and slepton masses squared. Since there is no other logarithmically enhanced contribution, we conclude that the log-dominated regime is excluded.

For $4\pi F \sim \Lambda_{\rm conf}^2$, we are not in the log-dominated limit, and there are equally large contributions to squark and slepton masses squared of unknown sign. We refer to the parameter region with $4\pi F \sim \Lambda_{\rm conf}^2$, $y_A \ll 4\pi$, a natural one for our model, as the ''light messenger'' regime. We will assume that the model has an acceptable region of parameter space in the light messenger regime with positive squark and slepton masses squared. Note that all contributions which are not suppressed by weak couplings are approximately SU(5) symmetric, so the two-loop contribution to all standard model sfermion masses squared has the same sign.

In such an acceptable regime, estimates for gauge-mediated supersymmetry-breaking masses can be made following the usual arguments. For a sfermion $\tilde{f}_r$ in a representation r under the standard model, we find

$$\tilde{m}_{\tilde{f}_r}^2 = c_m\left(\sum_{k=1}^{3} C_{rk}\alpha_k^{\rm SM^2}\right)\frac{F^2}{\Lambda_{\rm conf}^2}. \tag{30}$$

Here the C_{rk} are the Casimir indices for the representation r of the standard model gauge group [S]U(k), while c_m is an unknown constant, assumed positive. Because sfermion masses squared scale with the effective number of messengers, and perturbatively we have the equivalent of five messenger $5+5$'s of global SU(5), we estimate c_m is of order 5, and, due to approximate SU(5) symmetry, is approximately independent of the representation r. For the standard model gauginos M_i, we expect

$$M_i = d_m \alpha_k^{\rm SM}\frac{F}{\Lambda_{\rm conf}}. \tag{31}$$

Here d_m is another unknown constant, also of order 5.

The fields A and $\bar{A}$ have quantum numbers of a **10** and $\overline{\mathbf{10}}$ under SU(5). However because the global SU(5) is broken by weak couplings such as the y_A, the superfields $A,\bar{A}$ are each three multiplets with different masses, A_r, $\bar{A}_{\bar{r}}$ where $r=(3,2)_{1/6}$, $(\bar{3},1)_{-2/3}$, $(1,1)_1$ labels the representation of A_r under the standard model. We assume that the three different y_A couplings are all of the same order.

If the y_{A_r} are of order one, none of these particles will be observable in the near term. But if the y_{A_r} are small, then the fields Ψ_A, A might be light enough to be found soon. The color-neutral fields are likely to be the lightest since the scalar masses do not receive a large gauge-mediated contribution from color, and because the renormalization group predicts enhancement of the Yukawa couplings for colored particles. Since Ψ_A and A make up $\frac{3}{2}$ of an SU(5) supermultiplet, nonobservation of $\bar{A}$ in the near vicinity would suggest that $\bar{A}$, and the $\psi_{\bar{A}}$ component of Ψ_A, are participants in the supersymmetry-breaking dynamics. One might then probe this dynamics by studying the interactions of the Ψ_A particle.

The SU(5) global symmetry could be broken by the superpotential couplings. We have assumed (see Sec. V F) that all such couplings either are weak or, if strong, are drawn towards an SU(5) invariant fixed point. This assumption may be tested via the relations of Eqs. (30) and (31). Another interesting test is possible if the Ψ_A and A masses are measured. Global SU(5) symmetry gives the sum rule

$$m_{A_r}^2 - \tilde{m}_{\tilde{f}_r}^2 = x m_{\Psi_r}^2, \tag{32}$$

where x is a constant independent of r.

In addition, our model (in contrast to most GMSB models) may generate, through strong dynamics, relatively large ''superoblique'' corrections [59,60]—supersymmetry-violating contributions to the gaugino couplings. The typical expected size of such corrections can be estimated from naive dimensional analysis to be

$$\frac{g_i^{\rm SM} - \tilde{g}_i^{\rm SM}}{g_i^{\rm SM}} \sim \frac{\alpha_i^{\rm SM}}{4\pi}, \tag{33}$$

where the $\tilde{g}_i$ are the gaugino Yukawa couplings. In the light messenger limit, a logarithmic enhancement of this contribution can be reliably computed to be [59,60]

$$\frac{g_i^{\rm SM} - \tilde{g}_i^{\rm SM}}{g_i^{\rm SM}} = \frac{\alpha_i^{\rm SM}\xi_i}{4\pi}\log\left(\frac{4\pi F}{\Lambda_{\rm conf} m_\Psi}\right), \tag{34}$$

where $\xi_i = (\frac{5}{3}),1,1$ for $i=1,2,3$ respectively, and we have neglected differences between the different Ψ_r masses.

C. Trilinear scalar terms and the Higgs sector

As in most gauge-mediated supersymmetry models, the supersymmetry breaking parameter B which governs mixing of the H_u and H_d scalars, and the trilinear terms among scalar fields of the MSSM, are generated at two loops, and are relatively small, of order $(\alpha^{\rm SM})^2 F/(4\pi\Lambda_{\rm conf})$. The only way to generate such terms at one loop is to have heavy gauginos which couple to MSSM particles, such as when the standard model gauge group is embedded in a larger group.

Our model does not address the generation of the μ term, i.e., the supersymmetric Higgsino mass parameter in the MSSM superpotential. A long standing problem for supersymmetric theories is to explain why this parameter should be of order the weak scale, as is phenomenologically required [63]. Several solutions to this problem, which were proposed in the original models of GMSB with DSB [7–9], would also work for the present model. Basically these solutions rely on generating the μ parameter from the VEV of a fundamental singlet which is coupled to the MSSM sector. Since these original models, several newer solutions have been proposed. However none of these newer solutions [63–66] will work for the present model, either because they need $\sqrt{F}$ to be substantially larger than $\sim 10^5$ GeV, or because they require fundamental gauge singlet superfields with renormalizable couplings to the messenger sector. Witten's sliding singlet mechanism [67] has been claimed to generate an acceptable μ term in some models with $\sqrt{F}\sim 10^5$ GeV [68]; however, for this mechanism to generate an acceptable μ parameter, the B parameter must also be large, which is not straightforward to arrange in this model.

D. Decay of the messengers

So far we have assumed that there are no superpotential couplings between the MSSM and the supersymmetry-breaking sectors. This assumption could lead to stable charged particles in the messenger sector. Such particles are easily ruled out via, e.g., searches for heavy hydrogen.

The heavy $B\bar{B}$ messengers can decay via an Sp(4) instanton into two A messenger particles and a neutral C messenger. The C messengers are allowed to decay via the couplings h_1, h_2 into the light particles of the 3-2 sector, i.e., the Goldstino and the massless fermion mentioned in Sec. VI A. However, the lightest A messengers in each charge sector are stable unless new couplings are introduced.

The easiest way to eliminate the cosmological problems of stable charged messengers is to assume that dimension-five couplings between the supersymmetry breaking and MSSM sectors are allowed, e.g., $A\bar{U}\bar{D}\bar{E}$ (where $\bar{U},\bar{E}$ are, respectively, the MSSM up antiquarks and charged antileptons). Even if suppressed by the reduced Planck scale, dimension 5 couplings lead to lifetimes for messengers with TeV scale masses (such as most of the A multiplet might have) of around 10^{-2} s. In this case the lightest messengers would be irrelevant for cosmology but might be detectable in a collider experiment as heavy, long-lived charged particles.

Alternatively, (if A is in the **10**, not the $\overline{\mathbf{10}}$ representation) one could allow the following renormalizable couplings which are consistent with baryon and lepton number conservation and all the gauge symmetries.[4] Such couplings, if larger than $\sim 10^{-6}$ or so, will allow prompt decays of all new heavy particles.

$$\lambda_{LL}^{ij} L_i L_j A_{(1,1)_1} + \lambda_{LD}^{ij} L_i \bar{D}_j A_{(3,2)_{1/6}} + \lambda_{DD}^{ij} \bar{D}_i \bar{D}_j A_{(\bar{3},1)_{-2/3}}. \tag{35}$$

Here i, $j=1$, 2, 3, and $\bar{D}_i$, L_i are the left chiral superfields

[4]Note that the simple alternative of allowing the A particles to mix with those ordinary quarks and leptons with the same quantum numbers leads to rapid proton decay. We can impose a discrete symmetry to forbid such couplings and still allow those of Eq. (35).

for respectively the down antiquarks and the lepton doublets, and the coupling to the A field is to the appropriate component such that the coupling is gauge invariant. Hence when λ_{DD}, λ_{DL}, λ_{LL} are nonzero, A contains superfields with the quantum numbers of a diquark, a leptoquark, and a dilepton.

The addition of the λ_{LD}, λ_{LL}, λ_{DD} terms allows for a new classically flat direction parametrized by the superfields $\overline{V}^2$, $\overline{D}$, $\overline{L}$, along which the Sp(4) gauge symmetry is completely Higgsed. Along this flat direction the Sp(4) dynamics no longer lifts the classically flat directions parametrized by q^2B, B^3, and the SU(3)$\times$SU(2) factor is completely Higgsed as well. Thus there is a classically flat direction along which nonperturbative gauge effects are small, and so the couplings of Eq. (35) lead to a new, supersymmetric vacuum at infinity in field space. Still, provided the couplings are sufficiently small, there will still be a local supersymmetry-breaking minimum in the vicinity of the minimum which is there in the absence of such a coupling. Furthermore, when λ_{LD}, λ_{LL}, λ_{DD} are small, the barrier between the desired and the new minimum becomes very large and the lifetime of the desired vacuum becomes much longer than the lifetime of the observed universe. There is no known reason why we should not be living in such a false vacuum.

Another reason for requiring any couplings between A and the MSSM fields to be small is that such couplings can lead to flavor-changing neutral currents at tree level. The strongest constraint is on certain combinations of the couplings λ_{DL}^{11}, λ_{DL}^{22}, λ_{DL}^{12} and λ_{DL}^{21}, since the scalar leptoquark exchange can give a tree level contribution to rare K decays such as $K_L \rightarrow \mu e$ and $K \rightarrow \mu e \pi$. Furthermore, as argued in Sec. V E, the colored A scalars may be as light or lighter than the squarks. The leptoquark contribution to rare K decays will be compatible with all current bounds provided the λ_{DL} couplings are all smaller than $\sim 10^{-3}$. Fortunately it is natural for these couplings to be small since they are irrelevant in the energy regime between $\Lambda_{\rm conf}$ and Λ_3.

We conclude that another acceptable scenario is that the couplings λ^{LD}, λ^{LL}, λ^{DD} are all present but very small. In this case the only nonstandard stable or long-lived particles in the model are the gravitino $\tilde{G}$ and the massless fermion discussed in Sec. VI A, and neither of these lead to any phenomenological problems or cosmological difficulties. In this scenario, the lightest messengers (the Dirac fermion Ψ_A and the complex scalar A) might be pair produced and their decays observed in a hadron collider. The scalars are R parity even. The scalar dilepton, which decays into a charged lepton and a neutrino, could be the lightest of the even R parity messengers, potentially as light as the right handed charged sleptons. The leptoquark scalars are a nearly degenerate weak doublet, with the charge $\frac{2}{3}$ member decaying promptly into charged lepton and a down quark jet, and the charge $-\frac{1}{3}$ member into a down quark jet and a neutrino. These could be as light or lighter than the ordinary squarks. Such leptoquarks have already been searched for at Fermilab [69, 70] but will escape detection if heavier than 225 GeV. The Ψ_A fermions could be lighter than their scalar superpartners. If heavy enough, the messenger fermions decay into an ordinary quark and lepton and a squark or a slepton. If such decays are not kinematically allowed they must decay via virtual squarks and sleptons. In either case, the decay chain for a messenger fermion always leads to the NLSP and its typical decay signature.

VII. CONCLUSIONS

We have presented an explicit example of gauge-mediated dynamical supersymmetry breaking with a single supersymmetry breaking and messenger scale in the 10–30 TeV region. Our model has no fine tuning, explicit mass parameters, or *ad hoc* small numbers. The model is relatively simple, but exhibits a rich variety of phenomena. We have discussed the dynamics and properties of this model in detail. This requires a careful analysis of the effects of strong coupling that goes well beyond the use of holomorphy to constrain the superpotential.

While this model preserves the usual successes (no flavor changing neutral currents, a predictable spectrum) and difficulties (the μ and $B\mu$ problems) of gauge-mediated supersymmetry breaking, it has some unusual features and predictions as well. Some of these are central to the model and would be typical of any theory of this type.

(1) All couplings in the theory are of naturally of order one at high scales, although some are forced to be very large or very small at low energy by the effects of strong dynamics.

(2) Despite the large contribution of new strong interactions to the beta functions of the standard model gauge couplings, the usual supersymmetric GUT relations are protected by an approximate SU(5) global symmetry.

(3) Although the theory has two strongly coupled gauge groups, a single dynamical scale controls the physics of the supersymmetry-breaking and messenger sectors. This is due to a natural approximate fixed point, whose dynamics washes out the effects of the other energy scales.

(4) As a direct result of the lack of segregation between supersymmetry breaking and messenger dynamics, the supersymmetry-breaking scale is low, resulting in an NLSP which can decay promptly into an ordinary particle and a gravitino.

(5) The messenger sector is completely chiral with respect to the underlying gauge symmetries; it only becomes vector-like after one of the gauge groups undergoes a confining transition. The confining transition sets both the supersymmetry-breaking scale and the mass scale of the messenger sector.

Somewhat more model-dependent but still reasonably generic properties depend on the irrelevance, as a result of strong dynamics, of a certain coupling in the superpotential. This coupling may be driven much less than one, and if it is small (between 10^{-3} and 10^{-1}) it leads to a number of interesting effects.

(1) Certain messenger chiral superfields only couple weakly to the supersymmetry-breaking dynamics, while most of them are strongly coupled. This leads some of the particles in the messenger-DSB sector to have supersymmetry-preserving masses much smaller than the supersymmetry-breaking scale. These ''light messengers''

might be discovered soon by hadron colliders.

(2) The chiral structure of the model leads the light messenger supermultiplets to have supersymmetry-breaking mass splittings which differ substantially from those of their complex conjugates. Consequently, the light messenger particles do not form a set of complete supermultiplets.

(3) The large mass splittings in the light messenger supermultiplets cause ‘‘superoblique’’ radiative corrections to the gaugino couplings to be logarithmically enhanced, and thus much larger than in weakly coupled gauge mediated models.

(4) The usual constraints on light messengers, due to their tendency to give negative mass squared to standard model fermions, are evaded as a result of the strong-coupling dynamics in the supersymmetry-breaking sector.

(5) The SU(5) global symmetry gives a sum rule for the light messenger masses.

If indeed there are light messengers (which is the most natural regime for the theory), these effects make the model experimentally distinguishable from both gravity-mediated models and other proposed gauge-mediated models. We hope that the novel features of this model will be thought provoking, and will stimulate further research into the role that strong gauge dynamics may play in the process of supersymmetry breaking.

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