A More Minimal Messenger Model of
Gauge-Mediated Supersymmetry Breaking?

G. Dvali
Theory Division, CERN, CH-1211 Geneva 23, Switzerland

and

M. Shifman
Theoretical Physics Institute, University of Minnesota, Minneapolis, MN 55455, USA

Abstract

This Letter addresses a provocative question: “Can the standard electroweak Higgs
doublets and their color-triplet partners be the messengers of a low energy gauge-
mediated SUSY breaking?” Such a possibility does not seem to be immediately ruled
out. If so, it can lead to a very economical scheme with clear-cut predictions quite
distinct from those of the conventional gauge-mediated scenario. Namely, we get
(i) a single light Higgs below the original SUSY-breaking scale; (ii) tanβ = 1; (iii)
flavor non-universal, but automatically flavor-conserving soft scalar masses; (iv) a
light colored scalar with peculiar phenomenology. The familiar µ problem looses its
meaning in this approach.

December 1996
1 Introduction

The origin of supersymmetry (SUSY) breaking remains a key problem for modern models of fundamental interaction. For phenomenological implications the most important question, however, is not the precise nature of supersymmetry breaking per se but, rather, how this breaking is communicated to the low-energy observable sector: quarks, leptons and gauge fields. Recently interest has been revived \[1\] in the so-called “gauge-mediated” low-energy SUSY breaking scenarios \[2\]. This approach is motivated predominantly by its predictivity and a potential for solving the problem of the flavor-changing neutral currents \[3\]. The main ingredients of the approach are as follows. Supersymmetry breaking occurs through one of the known mechanisms (usually, through a dynamical mechanism \[4\]) in some (usually strongly coupled) hidden sector of the theory. The role of this sector is to provide a \( G_W = SU(3) \otimes SU(2) \otimes U(1) \)-singlet chiral superfield \( X \) with non-vanishing vacuum expectation values (VEV) of its auxiliary and scalar components \( \langle X \rangle \neq 0, \quad \langle F_X \rangle \neq 0 \).

The above VEVs ensure breaking of both supersymmetry and the \( R \) symmetry. A key role in transmitting this breaking from the hidden sector to the visible matter belongs to the so-called “messenger” sector composed of the superfields \( \phi, \bar{\phi} \) in the vector-like representations of the standard model (SM) gauge group. Sometimes this mechanism is considered in the context of grand unified theories (GUT). In this case, \( \phi, \bar{\phi} \) are assumed to belong to the vector-like representations of the GUT gauge group. The fields in \( \phi, \bar{\phi} \) experience a tree level Fermi-Bose mass splitting due to a direct coupling with the \( X \) superfield in the superpotential

\[
W = h X \bar{\phi} \phi .
\]

A standard “minimal” choice for the messengers in the GUT version is \( \phi \sim 5 \) and \( \bar{\phi} \sim \bar{5} \). Even if grand unification is not considered, the minimal choice remains essentially the same, the components of \( \phi, \bar{\phi} \) are assumed to transform under \( G_W \) as

\[
\{ \bar{3}, 1, (-2/3) \} + \{ 1, 2, 1 \} ,
\]

and their conjugates, respectively. This assignment ensures non-vanishing one-loop soft masses for the “observable” gauginos

\[
m_{\lambda_i} \sim K_i \frac{\alpha_i}{4\pi} \left( \frac{F_X}{X} \right) ,
\]

and two-loop soft masses for the scalars

\[
m_{\lambda_i} \sim K_i C_i \frac{\alpha_i}{4\pi} \left( \frac{F_X}{X} \right)^2 ,
\]

\[1\] Below we will denote the chiral superfields as well as their lower components by the same symbols. In each case it will be clear from the context to which component we refer to.
where $\alpha_i$ are the gauge couplings, and $K_i$ and $C_i$ are group-theoretical factors which only depend on the gauge quantum numbers. Thus, in the minimal case we can parameterize our ignorance of the messenger sector by a single ratio

$$\Lambda = \frac{F_X}{X}$$

which, for realistic soft masses, must be $\sim 100 \text{ TeV}$ or so.

Along with obvious attractive properties, the low-energy gauge-mediated scenario suffers from an aesthetically ugly feature: the messenger sector is composed of ad hoc fields whose sole raison d'etre is to connect, in a SUSY breaking way, the hidden and observable sectors. These superfields are not otherwise motivated, and this brings in a certain degree of arbitrariness in the theory. Moreover, they may lead to a serious cosmological difficulty [5], as the lightest of the messengers tends to be a stable particle. To avoid the problem, an instability must be ensured by postulating messenger couplings to ordinary matter. Such couplings may introduce back the flavor non-universality and unacceptably strong baryon number violation (unless they are strongly suppressed by hand).

Another serious difficulty of this approach is the $\mu$ problem [6]. One needs to generate both a supersymmetric Higgs mass term (Higgsino mass) in the superpotential

$$\Delta W = \mu \bar{H}H$$

and a soft scalar bilinear mass ($B\mu$ term) in the potential

$$\Delta V_{\text{soft}} = B\mu \bar{H}H + \text{h. c.}$$

of the right order of magnitude,

$$B\mu \sim \mu^2 \sim (100\text{GeV})^2.$$  \hspace{1cm} (8)

Usually this is difficult to achieve in the minimal schemes, since, once forbidden at the tree level, one tends to end up with a problematic relation $B\mu \gg \mu^2$ (for a more detailed discussion and possible ways out see [6] and references therein).

In view of the above, a natural question arises: can one exploit the fields which exist in the modern theory anyway, to assign to them the messenger role? In the present Letter we address this issue, and analyze whether the electroweak Higgs doublets $H, \bar{H}$ and their color-triplet GUT partners $T, \bar{T}$ can play the role of the messengers of the low-energy gauge-mediated SUSY breaking.

Thus, our task is the search for a minimal messenger model (MMM).

If $T, \bar{T}$ can play the role of the messengers, the aesthetically unpleasant feature of the approach is eliminated. On the practical side, our minimal messenger model exhibits no flavor problem. Its solution is automatic – the Yukawa coupling constants are diagonalized simultaneously with the fermion masses. The $\mu$ problem gets a different (essentially no) meaning, since, as shown below, in the low-energy sector there is a single light Higgs scalar

$$h = \frac{H + \bar{H}^+}{\sqrt{2}}.$$ \hspace{1cm} (9)
(The mass of this particle will be denoted below as $M_h$.) The orthogonal superposition and Higgsinos are heavy. Their masses are of order $\Lambda$, and they decouple. This does not lead to the usual naturalness problems, however, since the Higgs mass is only two-loop corrected.

Below we will argue that such a scenario is not ruled out and leads to clear-cut predictions, which can be tested in present and future experiments. Apart from a very different low-energy Higgs spectrum mentioned above, MMM implies the existence of light color-triplet scalar Higgs particles $T$, whose mass can be close to $M_h$. Usually the light color-triplet Higgs particles are not considered because of the menace of a fast proton decay. In supersymmetric theories the proton decay can be naturally suppressed, however, by the so-called Clebsch-factor mechanism, see Ref. [7]. Depending on the Clebsch factors emerging in the underlying GUT, there are three possible outcomes: the light Higgs triplet decays (i) only in the quark channels; (ii) only in the lepton channels; (iii) appears to be stable in the detector and must be observed in the form of stable charged or neutral hadrons (more exactly, it is not absolutely stable, but the lifetime is large).

## 2 Higgs weak doublet and colour-triplet as messengers

Since we want the ordinary Higgs doublets and their color-triplet partners to be the only messengers of the low-energy gauge-mediated SUSY breaking, we assume the source of their masses to be a coupling with the $X$ superfield in the superpotential, with a VEV of order $\Lambda$. At first sight, this sounds impossible, since such a light color triplet is believed to lead to unacceptably fast proton decay. However, generically this is not true [7]. The proton decay can be eliminated by Clebsch factors; these are certain dimensionless combinations of the GUT Higgs VEVs which control the strength of the effective Yukawa coupling constants after the GUT symmetry breaks down. These effective Clebsch factors are low-energy remnants of the heavy sector integrated out at the GUT scale. They can naturally decouple the Higgs triplet $T$ from some (or all) of the species of quarks and leptons, thus automatically suppressing the proton decay.

Let us briefly discuss the main idea of the Clebsch-suppression mechanism. Consider a grand unified group $G$ with quarks and leptons transforming in the irreducible (or reducible) representation $\Psi^\alpha$ with $\alpha = 1, 2, 3$ being a family index. Let $\Sigma$ denote the Higgs representation(s) that break(s) $G$ to $G_W$ (generically, there can be more than one such Higgs field). Let Higgs doublet and triplet be placed in the irreducible representation $R$. Then the masses of the ordinary fermions are generically induced from a set of effective $G$ invariant operators (a $G$ invariant contraction of the group indices is assumed)

$$R \left( \frac{\Sigma}{M_G} \right)^{n_{\alpha,\beta}} \Psi^\alpha \Psi^\beta.$$  \hspace{1cm} (10)
These operators are induced after integrating out all heavy fields at $M_G$ and, depending on the precise structure of the theory up there, may have different contractions of the indices and different flavor dependence. In view of the fact that none of the minimal SUSY GUTs (e.g. $SU(5)$ or $SO(10)$) with the minimal Yukawa interactions (corresponding to $n = 0$ in Eq. (10)) can account for the observed pattern of the fermion masses, such operators are very much motivated. An important consequence of such a construction is that, after breaking the GUT symmetry by the vacuum expectation value of $\Sigma$, the universality of the resulting doublet and triplet Yukawa coupling constants is generically 100% violated. The relation between the couplings is determined by the group theoretical (Clebsch) factors. It is perfectly natural that for a certain choice of the above operators the triplet turns out to be decoupled from some (or all) species of the quark and lepton superfields. In other words, the triplet is coupled to matter only in combination with certain components of $\Sigma$ and is automatically decoupled if the latter have vanishing (or small) expectation values. In short, the Clebsch mechanism \[7\] insures the decoupling of the triplet Higgs not by adjusting its mass to be huge, as in the standard scenario \[8\], but, rather, through suppressing the corresponding coupling constants. It exhibits certain advantages over the standard doublet-triplet mass-splitting solutions \[8\], since it kills simultaneously both dimension-5 \[9\] and dimension-6 proton decaying-mediating operators. For further details the reader is referred to Ref. \[7\]. Here we simply parameterize the Clebsch factors by independent parameters subject to the constraint of the proton stability, and then consider phenomenologically the most promising possibilities.

Consider first the Higgs spectrum. By assumption, the states $H, \bar{H}$ and $T, \bar{T}$ get supersymmetric and non-supersymmetric contributions to their masses from the couplings to the $X$ field in the superpotential

$$W = gX\{\bar{H}H + (1 + a)T\bar{T}\}.$$ \hspace{1cm} (11)

To allow for different relative wave function renormalizations from the GUT scale down to the scale of the messenger sector we have introduced above a factor $a$. Its value will be discussed shortly. After supersymmetry breaking takes place the above superpotential leads to the following tree-level masses

$$g^2|X|^2 \left(|H|^2 + |\bar{H}|^2 + (1 + a)^2(|T|^2 + |\bar{T}|^2)\right) + \left[gF_X \left(H\bar{H} + (1 + a)T\bar{T}\right) + \text{h.c.}\right].$$ \hspace{1cm} (12)

If we assume, for definiteness, that $F_X < 0$, the condition of existence of a light doublet is

$$g^2|X|^2 + gF_X = M_h^2 \ll \Lambda^2.$$ \hspace{1cm} (13)

Then the combination indicated in Eq. (14) is the lightest Higgs, with the mass $M_h$, while the orthogonal combination $\sim (H - \bar{H}^+)$ is heavy, with the mass $\sim g|X|$. The lightest scalar triplet is given by a similar combination,

$$T_h = \frac{T + \bar{T}^+}{\sqrt{2}}.$$ \hspace{1cm} (14)
In fact, the most obvious choice is $M^2_h = 0$ (at the scale $\Lambda$). Some ideas as to how this cancellation may naturally take place will be discussed shortly. Here we want to mention that, even being regarded as an explicit input fine-tuning, Eq. (13) still is a much less severe condition than the fine-tuning in the standard $SU(5)$ \[8\], since here we fine-tune two orders of magnitude versus 14 orders in the standard version.

Another advantage over the standard $SU(5)$ approach, where the fine-tuning (to zero) does exhibit the $\mu$ problem in the gauge-mediated supersymmetry breaking framework, is that our suggestion eliminates the $\mu$ problem altogether.

### 2.1 Minimal $SU(5)$

This is the most economic version. We have only $T$'s and $H$'s as messengers of SUSY breaking. The mass of Eq. (14), however, generically is substantially heavier than $M_h$, see Eq. (16) below. This is due to the fact that the breaking $SU(5) \to G_W$ occurs at a very high scale, and evolving down to $\Lambda$ brings in the difference in the wave function renormalizations. Generically, the constant $a$ is several units. Although this does not preclude the messengers from their mission of SUSY breaking, other potentially appealing features appearing due to $M_h \sim M_{T_h}$, discussed in Sect. 3, are lost. Note, however, that $T$ can have a large Yukawa coupling with some of the species of the third generation, of order unity. The wave function renormalization due to the gauge coupling and due to the Yukawa couplings have opposite signs, and tend to cancel each other. It may happen that, thanks to this cancellation, $a$ is numerically rather small, and $M_h \sim M_{T_h}$ is still valid.

### 2.2 Advanced GUT’s

Now let us briefly discuss how the cancellation (13) may happen due to symmetries of the theory. The same symmetries will ensure also that $a \ll 1$. In this paper we would not like to enter into a detailed discussion of specific models; an example we give must be rather regarded as an “existence proof”.

The cancellation may be ensured by a pseudo-Goldstone nature \[10\] of $T$ and $H$\[1\]. Imagine that the GUT symmetry is $SU(6)$; it is broken to $G_{\text{intermediate}} = SU(3) \otimes SU(3) \otimes U(1)$ at some high scale $M_G$ and then, at a lower scale $M$, not far from $\Lambda$, is further broken to $G_W$. The sector of the theory responsible for the breaking of $SU(6)$ down to $G_W$ contains a $35$-plet and 6- ($\bar{6}$)-plets. Assume that $H \ (\bar{H})$ and $T \ (\bar{T})$ states belong to the fundamental $6$-plet ($\bar{6}$-plet) of $SU(6)$. To this end they should be supplemented by an additional $G_W$-singlet field $S \ (\bar{S})$. The embedding is such that $H$ and $S$ compose a triplet under one of the $SU(3)$ subgroups, while $T$ is a triplet under another $SU(3)$ subgroup.

\[2\]The idea was used previously, in particular, for the interpretation of the solution in Ref. \[8\]. Moreover, it was pointed out, in a different context \[13\], that it can ensure the cancellation even for large $\mu$ and $B\mu$ parameters.
With the gauge interaction switched off, and with both 6- and 6-plets developing VEV’s, we would get twice more “phase” fields than the number of such fields that are actually eaten in the (super)Higgs mechanism. Combinations orthogonal to those that are eaten in the Higgs mechanism remain massless at the tree level. Their masses appear only after SUSY breaking and are small. If we additionally assume that the VEV of the \((T,H,S)\) 6-plet are significantly smaller than that of the “other” 6-plet, then the pseudo-Goldstone bosons discussed above will be almost pure \((T,H,S)\).

A relevant superpotential can be written as

\[
W = gX \left( (1 + a)T\bar{T} + H\bar{H} + S\bar{S} - M^2 \right).
\]

(15)

In the example at hand \(a = 0\) (above the scale \(M\)) if no additional states are introduced at \(M_G\). The factors of the relative wave function renormalization of doublet(s) and singlet(s) are equal because of the \(SU(3)\) symmetry. The equality of the relative wave function renormalization of \(T\) and \(H,S\) (i.e. \(a = 0\)) is due to an obvious symmetry of the theory under the interchange of two \(SU(3)\) groups. A very small value of \(a\) is only generated below the scale \(M\) of the second breaking \(3\).

Now, SUSY is spontaneously broken whenever dynamics induces a non-vanishing VEV of \(X\). Indeed the minimization with respect to all other fields automatically gives Eq. (13) provided \(X \neq 0\). This fact is not surprising. Indeed, the state \(\text{(9)}\) appears to be the pseudo-Goldstone of the broken \(SU(3)\) and, thus, must be massless at the tree level. At the same time, the mass of the lightest scalar triplet state \(\text{(14)}\) is given by

\[
m^2_T = g^2|X|^2a(1 + a).
\]

(16)

It vanishes in the limit \(a \to 0\), as it should. The precise value of \(a\) depends on details of the GUT scheme in question. The most interesting phenomenology results for \(a \ll 1\) which we briefly discuss below.

### 2.3 Soft mass terms of matter

Now, let us consider how the soft masses of the matter fields are generated through our gauge-mediated scenario. As usual, there are one-loop gaugino \(\text{(4)}\) and flavor-universal two-loop scalar \(\text{(5)}\) soft masses. A new crucial point is the generation of the soft scalar masses at one-loop level. These mass terms are generated due to direct Yukawa couplings of our Higgs messengers to the squarks and sleptons and are not flavor-universal. They have the form

\[
m^2_{\alpha\beta} = \frac{G^\alpha_{\gamma\gamma}G^\gamma_{\beta\beta}}{16\pi^2}O(\Lambda^2),
\]

(17)

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3This scheme above can exhibit a difficulty in accommodating the standard unification of the gauge couplings. This problem may be solved by a suitable extension of the field content. We thank I. Gogoladze for bringing such examples to our attention.
where $G_{\alpha\beta}$ are the Yukawa coupling constants, and we have taken into account the fact of SUSY breaking in the Higgs sector, with the scale $O(\Lambda)$. These masses are manifestly non-universal; by far the largest appears in the third generation. What is remarkable, the large and flavor-non-universal squark masses do not lead to flavor violation since they are diagonalized simultaneously with the fermion masses. More explicitly the one-loop soft scalar masses for the different species $Q^\alpha, d^\alpha, u^\beta, L^\alpha, e^\beta_c$ are proportional to:

\[
m^2_{Q\alpha\beta} \propto G_{\alpha\gamma}^{u*} G_{\gamma\beta}^{u},
\]

\[
m^2_{d\alpha\beta} \propto G_{\alpha\gamma}^{d*} G_{\gamma\beta}^{d},
\]

\[
m^2_{u\beta} \propto G_{\alpha\gamma}^{e*} G_{\gamma\beta}^{e},
\]

\[
m^2_{L\alpha\beta} \propto G_{\alpha\gamma}^{e*} G_{\gamma\beta}^{e},
\]  

where $G_{Q,u,d,L,e}^{u,d}$ are the respective Yukawa coupling constants. Unsurprisingly, the flavor violation is suppressed in this scheme and is controlled by the CKM mixing angles. Note that the parameter $\tan \beta$ is extremely close to 1, and the Higgs sector essentially looks as that of the standard model. Since the above soft masses are one-loop induced, for the third family our new contribution will dominate over the usual gauge-mediated two-loop mass. Thus, the above approach leads to: (i) the hierarchical, but flavor conserving (aligned), pattern of the soft masses; (ii) a single light Higgs doublet $h$; (iii) a light color-triplet Higgs with mass $\sim a\Lambda$. Since the corrections to $M_h$ appear only at two-loop level the two-loop hierarchy $\Lambda \gg M_h$ is natural.

3 **Light color-triplet Higgs: phenomenological implications**

As was mentioned, the most interesting phenomenological consequences will take place for $a \ll 1$. In this case the light color-triplet scalar can be the subject of experimental study at existing facilities or those planned for the near future. To ensure the proton stability we assume that the couplings of the Higgs triplet to the matter is suppressed by the Clebsch factor via the mechanism of Ref. [7]. For simplicity we parametrize the couplings of triplets with the matter superfields in terms of the flavor independent Clebsch factors (more complicated versions, with flavor-dependent structures, are also possible). Then the most general Yukawa couplings of the Higgs triplets in the superpotential are

\[
TG_{\alpha\beta}^u \left(C_{ue}^u e^\beta_c + C_{QQ}^u Q^\alpha Q^\beta\right) + \bar{T}G_{\alpha\beta}^d \left(C_{du}^d u^\beta_c + C_{QL}^d L^\alpha Q^\beta\right),
\]  

where $C_{ue}, ...$ are the Clebsch factors. To ensure the proton stability, some of these factors (or all of them) must be suppressed. For example, it is enough to have $C_{QL} = C_{ue} = 0$. In this case the light triplet, once produced, will tend to decay into the quark-quark pairs. Note that it is phenomenologically impossible to allow both the quark-quark and quark-lepton decay channels, since this would lead to unacceptable proton decay.
If so, these Higgs color triplets must be carefully considered as possible candidates for the ALEPH four-jet events. As well known, these events continue to accumulate, on the one hand, and continue to defy any reasonable explanation, on the other. Phenomenologically, the light triplet Higgses in this aspect will look similar to the down right-handed quarks in the models with the $R$ parity violation (for a review of such models see [12, 13]). Phenomenology of the two-quark decays is the same. If the $R$ parity violation explanation goes through (see e.g. [14]), the same should be valid for the light triplet Higgses. The opposite is also true.\footnote{This remark may not apply to the explanation of Ref. [13], as this work does not assume the pair production but, rather, assumes production of left-handed plus right-handed selectrons.} Theoretically, there are two important distinctions, however. First, since we do not violate the $R$ parity, the stability of LSP is preserved. Second, although [14] mimics the right-handed down squarks, their fermion partners are much heavier, with masses of order $\Lambda$, in sharp contradistinction to the situation with the right-handed down squarks, whose fermion partners are lighter than the squarks themselves.

Two further theoretical points deserve mentioning. If the triplet Higgs is light, it gives, through a loop, a noticeable contribution to $Z \to b\bar{b}$ yield, roughly at the level of that associated with the light stop and chargino (see e.g. Ref. [15]). This puts the theoretical prediction for $R_b$ right on top of the existing world average for $R_b$, which is slightly higher than the standard model prediction. Simultaneously, the genuine value of $\alpha_s(M_Z)$ goes down to 0.112, which is also welcomed [16].

On the other hand, the light triplet Higgs spoils unification of the gauge couplings within $SU(5)$ GUT. Given the experimental values of $\alpha_1$ and $\alpha_2$ we get a value of $\alpha_3$ too low to be compatible with data.

What if all Clebsch factors vanish in the supersymmetric limit? Their typical value after supersymmetry breaking, induced due to the shift of the heavy VEVs, is $\alpha\Lambda/4\pi M_G$ [7]. This is certainly not enough to mediate (an observable) proton decay or to make triplets decay in the detector. Therefore, experimentally such a decoupled scalar triplet should be observed in the form of stable possibly charged hadrons. (They are not truly stable, and can not accumulate in matter, but rather the lifetime is large.) Phenomenology of such states will be somewhat similar to that discussed in [17] in a different context.

### 3.1 Conclusions and Outlook

In this Letter we have suggested the possibility that the standard electroweak Higgs doublets and their color-triplet partners are the messengers of a gauge-mediated low energy supersymmetry breaking. While \textit{a priori} it is not obvious that such a possibility is free of inconsistencies, it does not seem to be ruled out so far. If so, a very economical, predictive and exciting scenario may emerge. The MMM approach, in its simplest form, is quite restrictive and leads to clear-cut predictions different from those of more conventional messenger scenarios:

1) a single light Higgs scalar below the scale $\Lambda$, in particular, implying $\tan\beta = 1$;
2) flavor non-universal, but automatically flavor-conserving soft scalar masses;
3) the possibility of the light scalar Higgs triplet (with quantum numbers of
dc) decaying only into quarks (only into leptons), or not decaying at all (in detector). The phenomenology of such colored scalar may be a subject of speculation in connection with the recent ALEPH four-jet events.

Potential difficulties of the above approach, which may require further assumptions about physics above the scale Λ, are related to the gauge coupling unification. Also, the electroweak symmetry breaking in this scheme deserves a careful study.

Acknowledgments:
We are grateful to Gian Giudice for very useful comments and discussions. One of the authors (M.S.) would like to thank the CERN Theory Division, where this work began, for its kind hospitality.

This work was supported in part by DOE under the grant number DE-FG02-94ER40823.

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