Gauge-mediated SUSY Breaking with a Gluino LSP

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Abstract

In gauge-mediated SUSY breaking models, messengers transmit SUSY breaking from a partially hidden sector to the standard model sector via common standard model gauge interactions. The minimal set of messengers has quantum numbers of a $5 + \bar{5}$ of SU(5); identical to the quantum numbers of the minimal Higgs sector of an SU(5) GUT. We show in a simple model with messenger masses of order the GUT scale that Higgs - messenger mixing quite naturally leads to a low energy MSSM with gluinos as the lightest supersymmetric particles [LSP]. We study the phenomenological consequences of such a model.
Gauge-mediated SUSY Breaking

Gauge-mediated SUSY breaking [GMSB] models \[1, 2\] solve the problem of flavor changing neutral currents inherent in the MSSM \[3, 4\]. Consider for the purposes of this short paper, flavor changing processes of charged leptons. Supersymmetric charged lepton mass terms are of the form

\[ \bar{e} m_e e \]

where \( e (\bar{e}) \) represents 3 families of left-handed (right-handed) fermions and their scalar partners and \( m_e \) is a complex 3 x 3 mass matrix. In addition, scalars necessarily have soft SUSY breaking mass terms given by

\[ \tilde{e}^* m^2_{\tilde{e}} \tilde{e} + \tilde{\bar{e}}^* m^2_{\tilde{\bar{e}}} \tilde{\bar{e}} \]

where \( \tilde{e} (\tilde{\bar{e}}) \) represents the left-handed (right-handed) sleptons and \( m^2_{\tilde{e}} (m^2_{\tilde{\bar{e}}} \) is an hermitian 3 x 3 mass squared matrix. One may always diagonalize the supersymmetric mass term \( m_e \) by a simultaneous rotation of the charged lepton and slepton fields. This rotation however will not, in general, diagonalize \( m^2_{\tilde{e}}, m^2_{\tilde{\bar{e}}} \), unless they are proportional to the identity matrix. Note, off diagonal slepton masses lead to flavor violating processes such as \( \mu \to e\gamma, \mu \to 3e, \mu \to e \) conversion, etc.

In GMSB, SUSY breaking occurs in an almost hidden sector of the theory due to the expectation value \( F_X \), the F component of a superfield \( X \). Moreover, standard model [SM] squarks, sleptons and gauginos do not couple directly to \( X \). Hence they do not obtain SUSY breaking masses at tree level. The states which couple directly to \( X \) are the messengers of SUSY breaking. They carry SM gauge interactions, but otherwise do not couple to squarks and sleptons directly. Thus SUSY breaking enters the SM sector at one loop to gauginos and at two loops to squarks and sleptons. These SUSY breaking effects are dimensionally of order \( \Lambda \equiv F_X / M \) where \( M \) is the messenger mass. Moreover, they are determined by gauge quantum numbers; thus, for example, the matrices \( m^2_{\tilde{e}}, m^2_{\tilde{\bar{e}}} \) are proportional to the identity matrix at \( M \). As a result individual lepton number is conserved. Hence processes such as \( \mu \to e\gamma \) are forbidden.\[1\]

\[1\]Our discussion ignored the possibility of new flavor violating interactions due to physics at the GUT scale, \( M_G \). These interactions can only enter through loops containing GUT mass states, hence they generate off diagonal mass squared terms suppressed by factors of \( (M/M_G)^2 \).
2 The Minimal Messenger Sector

The messenger states must carry both color and electroweak quantum numbers. In addition, the messengers should be in complete SU(5) representations, to preserve GUT predictions for gauge couplings. The minimal set of states satisfying these criteria transform as a $5 + \bar{5}$ with the color triplet (weak doublets) denoted as follows $t, \bar{t}$ $(d, \bar{d})$. In the minimal models, all messengers have a common mass $M$. The resulting soft breaking masses are as follows.

Gauginos obtain mass at one loop given by

$$m_{\lambda_i} = \frac{\alpha_i(M)}{4\pi} \Lambda \quad \text{(for } i = 1, 2, 3).$$

(1)

The scalar masses squared arise at two-loops

$$\tilde{m}^2 = 2\Lambda^2 \left[ \sum_{i=1}^{3} C_i \left( \frac{\alpha_i(M)}{4\pi} \right)^2 \right]$$

(2)

where $C_3 = \frac{4}{3}$ for color triplets and zero for singlets, $C_2 = \frac{2}{3}$ for weak doublets and zero for singlets, and $C_1 = \frac{3}{5} \left( \frac{Y}{2} \right)^2$, with the ordinary hypercharge $Y$ normalized as $Q = T_3 + \frac{1}{2} Y$ and $\alpha_1$, GUT normalized.

In the limit $M << M_G$ ($M_G$ is the GUT scale), we have $\alpha_3(M) >> \alpha_2(M) > \alpha_1(M)$. Thus right-handed sleptons are expected to be the lightest SUSY partners of SM fermions and binos are the lightest gauginos.

3 SUSY GUT and Higgs – Messenger Mixing

In the minimal SU(5) SUSY GUT, Higgs doublets are contained in a $5_H + \bar{5}_H$. In order to avoid large baryon number violating nucleon decay rates, the color triplet Higgs $t_H, \bar{t}_H$ must have mass of order $M_G$, while the Higgs doublets $d_H, \bar{d}_H$ remain massless at the GUT scale. The latter are responsible for electroweak symmetry breaking at $M_Z$.

Our main observation is that the Higgs in a SUSY GUT and the messengers of GMSB have identical quantum numbers. Thus, for messengers with mass at an intermediate scale, Higgs-Messenger mixing is natural. Moreover as a result of doublet-triplet splitting in the Higgs sector, the doublet and triplet messengers will also be split. This can have significant consequences for SUSY breaking masses.
As a simple example, consider the natural doublet-triplet splitting mechanism in SO(10) \cite{6}. The 10 of SO(10) decomposes into a $5 + \bar{5}$ of SU(5) and the adjoint 45 can be represented by an anti-symmetric $10 \times 10$ matrix. The Higgs sector superspace potential is given by

$$\begin{align*}
W_{\text{Higgs}} &= 10_H 45 10 + X 10^2
\end{align*}$$

where $10_H$ contains $5_H + \bar{5}_H$, 10 is an auxiliary $5 + \bar{5}$ introduced for doublet-triplet splitting and $X$ is a singlet. Assuming $<45> = M_G (B - L)$, i.e. 45 obtains an SO(10) breaking vacuum expectation value in the $B - L$ direction and $<X> = M$, we obtain the triplet (doublet) mass terms given by

$$\begin{align*}
W_{\text{Higgs}} &= t_H M_G \tilde{t} + t M_G \tilde{t}_H + M t \tilde{t} \\
&\quad + M d \bar{d}
\end{align*}$$

Note, the triplets naturally have mass of order the GUT scale, while the auxiliary doublets have mass $M$ and the Higgs doublets are massless.\footnote{The doublet mass is necessarily smaller than the triplets in order to suppress baryon number violating interactions \cite{7}. Specifically, if only $10_H$ couples to quarks and leptons, then the effective color triplet Higgs mass $\tilde{M}_t$ which enters baryon decay amplitudes is given by $\tilde{M}_t = M_G^2 / M$. Hence $\tilde{M}_t > M_G$ implies $M/M_G < 1$.} The theory we propose, with Higgs-messenger mixing, is quite simple. Assume the auxiliary 10 is the messenger of SUSY breaking, i.e. assume that $X$ gets both a SUSY conserving vev $M$ and SUSY breaking vev $F_X -$

$$\begin{align*}
<X> &= M + F_X \theta^2; \\
\Lambda &= \frac{F_X}{M} \sim 10^5 \text{ GeV}; \\
A &\equiv \frac{M}{M_G} \sim 0.1.
\end{align*}$$

Since the triplet messengers (mass $O(M_G)$) are heavier than the doublets (mass $O(M)$), SUSY breaking effects mediated by color triplets are suppressed. This has significant consequences for gluinos which only receive SUSY violating mass corrections through colored messengers.

\footnote{There are several different ways that a $\mu$ term for the Higgs doublets can be generated once SUSY is broken. We will not discuss this issue further here.}
Gauginos obtain mass at one loop given by
\[ m_{\lambda_i} = D_i \frac{\alpha_i(M)}{4\pi} \Lambda + \frac{\alpha_i(M)}{2\pi} \Lambda B^2 \quad \text{(for } i = 1, 2, 3) \] (6)

where \( D_1 = \frac{3}{5}, D_2 = 1, D_3 = 0. \)

- In order to generate SUSY violating gaugino masses, both SUSY and R
  symmetry must be broken. In this theory, \( F_X \) breaks SUSY and the scalar vev
  \( M \) breaks the R symmetry which survives GUT symmetry breaking. Thus
  both are necessary to generate the SUSY violating effective mass operator
  given by
  \[ \frac{1}{M} \int d^4 \theta \ X^\dagger X \ W^\dagger_\alpha W_{\alpha i} \quad \text{for } i = 1, 2, 3 \] (7)
  where \( M \) is determined by the heaviest messenger entering the loop.

- Note the terms proportional to \( B^2 \). Without them the gluino mass vanishes
  at one loop due to an accidental cancellation. In order to compensate for
  this one loop cancellation we include additional messengers with a common
  mass of order \( M_G \), and an R symmetry breaking mass \( M \). This sector thus
  contributes a common mass correction proportional to \( B \sim M/M_G \).

The scalar masses squared arise at two-loops. We obtain
\[ \tilde{m}^2 = 2\Lambda^2 \left[ C_3 \left( \frac{\alpha_3(M)}{4\pi} \right)^2 (A^2 + 2B^2) + C_2 \left( \frac{\alpha_2(M)}{4\pi} \right)^2 (1 + 2B^2) \right] + 2\Lambda^2 \left[ C_1 \left( \frac{\alpha_1(M)}{4\pi} \right)^2 \left( \frac{3}{5} + \frac{2}{5} A^2 + 2B^2 \right) \right] \] (8)

where \( C_i, \ i = 1, 2, 3 \) are defined after equation 2.

### 4 Low Energy Spectrum

The heaviest SUSY particles are electroweak doublet squarks and sleptons
and weak triplet winos, while right-handed squarks, sleptons and binos are
lighter. Finally gluinos are expected to be the lightest SUSY particle [LSP].
We have the approximate mass relations\(^4\) after renormalization group running to \( M_Z \),
\[ M_2 = \frac{\alpha_2(M_Z)}{4\pi} \Lambda \approx 3 \times 10^{-3} \Lambda \]
\[ M_3 = \frac{\alpha_3(M_Z)}{2\pi} B^2 \Lambda \approx 9 \times 10^{-5} (B/0.1)^2 \Lambda. \] (9)

\(^3\)I thank Kazuhiro Tobe for pointing this out to me. Note, eqns. (6, 8) are corrections
for similar equations in ref. [5].

\(^4\)neglecting terms of order \( A^2 \) or \( B^2 \), when possible
In addition, the gravitino mass (which sets the scale for supergravity mediated soft SUSY breaking effects) is given by

\[ m_{3/2} = \left( \frac{F_X}{\sqrt{3}M_{pl}} \right). \]  

Hence

\[ m_{3/2} = \left( \frac{M_G}{\sqrt{3}M_{pl}} \right) B \Lambda \approx 6 \times 10^{-4} \left( B/0.1 \right) \Lambda. \]  

The gluino mass \( M_3 \) depends on the arbitrary parameter \( B \), the ratio of the R symmetry breaking scale \( M \) to the typical messenger mass of order \( M_G \). The gravitino mass also depends parametrically on \( B \) when expressed in terms of the SUSY breaking scale \( \Lambda \).

With \( \Lambda = 10^5 \) GeV, we obtain

\[
\begin{align*}
M_2 & \approx 300 \text{ GeV} \\
m_{3/2} & \approx 60 \left( B/0.1 \right) \text{ GeV} \\
M_3 & \approx 9 \left( B/0.1 \right)^2 \text{ GeV}
\end{align*}
\]  

5 Signatures of SUSY with a Gluino LSP

Gluinos are stable.\(^5\) They form color singlet hadrons, with the lightest of them\(^6\) given by

\[
\begin{align*}
R_0 &= \tilde{g} \, g \\
\tilde{\rho} &= \tilde{g} \, q \, \bar{q} \quad \text{with} \quad q = u, \, d \\
S_0 &= \tilde{g} \, u \, d \, s
\end{align*}
\]  

where \( R_0 \) is an iso-scalar fermion [glueballino]; \( \tilde{\rho} \) is an iso-vector fermion and \( S_0 \) is an iso-scalar boson with baryon number 1.

It is unclear which one is the stable color singlet LSP. For this paper, we assume \( R_0 \) is lighter and that both \( \tilde{\rho}, \; S_0 \) are unstable, decaying via the processes \( \tilde{\rho} \to R_0 + \pi \) and \( S_0 \to R_0 + n \).

Consider the consequences of a gluino LSP. First, the missing energy signal for SUSY is seriously diluted. An energetic gluino, produced in a high

\(^5\)Assuming R parity is conserved.
energy collision, will fragment and form an hadronic jet containing an $R_0$. The $R_0$ will deposit energy in the hadronic calorimeter. Thus collider limits on squark and gluino masses must be re-evaluated. Gluinos with mass as large as 50 GeV may have escaped detection.\footnote{A lower bound on the gluino mass of 6.3 GeV has been obtained using LEP data on the running of $\alpha_s$ from $m_t$ to $M_Z$.[9] Thus glueballinos may be expected in the range from 6 – 50 GeV.}

At LEP a 4 jet signal is expected above the squark threshold, since squarks decay into a quark plus gluino.

Now consider possible constraints from exotic heavy isotope searches. Stringent limits exist on heavy isotopes of hydrogen. However, an $R_0$ must be in a bound state with a proton in order for these searches to be relevant. Such a bound state is unlikely due to the short range nature of the interaction of $R_0$ with hadrons; predominantly due to the exchange of a glueball (the lightest of which is 10 times heavier than a pion) or multiple pions. Strong limits on heavy isotopes of oxygen also exist. An $R_0$ can certainly be trapped in the potential well of a heavy nucleus. However, for these searches to be restrictive, the expected abundance of $R_0$-nucleus bound states must be above the experimental bounds. The dominant process for forming such bound states is for $R_0$s, produced by cosmic ray collisions in the earth’s atmosphere, to be captured into nuclei. A back of the envelope estimate gives an expected abundance bordering on the observable limit. A more detailed calculation is therefore needed to say more.

Finally, what about the cosmological abundance of $R_0$s. Since $R_0$s annihilate via strong processes $R_0 + R_0 \to 2\,\pi$, the cosmological abundance is quite suppressed. A rough estimate gives

$$n_{R_0} = 10^{-10} \left(\frac{m_{R_0}}{m_\pi}\right) n_B$$

where $n_B$ is the cosmological baryon density. As a result, $R_0$s are NOT dark matter candidates.

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