The Phenomenology of SUSY models with a Gluino LSP

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Abstract

In this paper we study the experimental consequences of a theory which naturally has a heavy, stable (or almost stable) gluino. We define the boundary conditions at a messenger scale $M \sim 10^{14}$ GeV which lead to this alternative phenomenologically acceptable version of the minimal supersymmetric standard model.

In this theory, either the gluino or the gravitino is the lightest supersymmetric particle [LSP]. If the gravitino is the LSP, then the gluino is the next-to-LSP with a lifetime on the order of 100 years. In either case, the gluino is (for all practical purposes) a stable particle with respect to experiments at high energy accelerators. Thus the standard missing energy signature for SUSY fails. A stable gluino forms a color singlet hadron, the lightest of which is assumed to be an isoscalar gluino-gluon bound
state ($R_0$). The $R_0$ has strong interactions and will interact in a hadronic calorimeter; depositing some fraction of its kinetic energy. Finally, in the case the gravitino is the LSP, bounds from searches for stable heavy isotopes of hydrogen or oxygen do not apply to the metastable $R_0$. 

1 Introduction

Supersymmetry [SUSY] is a strongly motivated candidate for new physics beyond the Standard Model [SM]. The minimal supersymmetric particle content is well defined and the interactions of all the new superparticles [sparticles] are constrained by the observed SM interactions as long as the theory has a conserved R parity. In order to search for SUSY one necessarily focuses on specific signatures which may rise above the standard model backgrounds. These signatures depend on how supersymmetry is broken; i.e. on the particle spectrum. There are two classes of models for SUSY breaking which have been studied extensively. They are the minimal supergravity model [CMSSM][1] (or constrained minimal supersymmetric model [CMSSM][2]) defined by 5 parameters and low energy gauge-mediated SUSY breaking [GMSB][3, 4] defined by 3 parameters. The 5 parameters in the CMSSM model are a universal scalar mass $m_0$, a universal gaugino mass $M_1/2$, a trilinear scalar coupling proportional to $A$ and $\mu$, $\mu B = m_3^2$ in the Higgs sector. In this theory, the LSP is the lightest neutralino or a sneutrino. The dominant signature for SUSY is missing energy due to the escape of the LSP from the detector. In low energy GMSB, the effective scale of SUSY breaking is given by $\Lambda = F/M \approx M \sim 10^5$ GeV. The other two parameters are $\mu$, $\mu B = m_3^2$ as before. In this case, the gravitino with mass of order 10 - 1000 eV is the LSP. The next to lightest SUSY particle [NLSP], for example the lightest neutralino, can decay into a gravitino and a photon. This can lead to energetic photons plus missing energy, such as the single $e^+e^-\gamma\gamma$ + missing energy event seen at CDF.

In a recent paper[5] it was shown how GMSB with Higgs-messenger mixing in an SO(10) theory naturally leads to a gluino LSP. Since the gluino carries color, it will be confined. The stable LSP is then assumed to be a gluino-gluon bound state, a glueballino ($R_0$). $R_0$ is a hadron and thus interacts strongly in a hadronic calorimeter. As a result the standard missing energy signature for SUSY fails. This theory has a characteristically different spectrum of squark, slepton and gaugino masses. In this paper we analyze the low energy spectrum of the theory, starting with boundary conditions at the messenger scale $M \sim 10^{14}$ GeV, using two (one) loop renormalization group equations for dimensionless (dimensionful) parameters down to the weak scale, $m_Z$. We self-consistently require electroweak symmetry breaking using the one-loop improved Higgs potential. The model has six SUSY
breaking parameters — $\Lambda \sim 10^5$ GeV, $a \sim b \sim 10^{-1} - 10^{-2}$, $\mu$, $\mu B = m_3^2$ and $d$ (the size of the D term contribution in units of $M_2$). One of these parameters, $b$, allows us to arbitrarily vary the gluino mass.

Finally we note that the version of the model presented in this paper includes the SUSY breaking contribution of an anomalous $U(1)_X$, in addition to GMSB at the messenger scale $M[5]$. We show that the additional D term contribution is necessary in order to obtain a phenomenologically acceptable theory. In particular, without the D term the light Higgs boson $h$ would be unacceptably light. In the following, we assume that the contribution to scalar masses from GMSB and the $U(1)_X$ D term are comparable. In a recent paper[6], we have presented a model which dynamically breaks SUSY and leads to comparable SUSY breaking effects from gauge-mediated and D term sources. The messenger scale in this theory is determined by the Fayet-Iliopoulos D term[7] of $U(1)_X$. This theory suggests that the boundary conditions we consider may be “natural.”

2 Boundary conditions at the Messenger scale

The boundary conditions at the messenger scale are determined by two sources of SUSY breaking — gauge-mediated and D term.

2.1 Gauge mediated SUSY breaking

The messenger sector is defined by an SO(10) invariant superspace potential

$$W \supset \lambda_a \ 10_H \ A \ 10 + \lambda_s \ S \ 10^2$$

The magnitude of the Fayet-Iliopoulos D term $\xi$ is given by $\xi = \epsilon M_{Pl}^2$ where $\epsilon = \frac{g^2 X}{8 \pi^2} \left( \frac{\sqrt{2} M_{Pl}}{g X} \right)^2$ and the string scale $M_{st}$ is determined by the compactification scale. In weak coupling $M_{st} = \frac{1}{\sqrt{2}} M_{Pl}$. In the strong coupling limit of the $E_8 \times E_8$ heterotic string[8], on the other hand, the compactification scale is identified with the GUT scale. Note, however, in a recent paper[9] it has been argued that even in the strong coupling limit of the heterotic string, the scale of the D term is set by the Planck scale and not the compactification scale. This argument has some caveats as discussed in the paper. For example, it might not apply if the anomalous gauge symmetry emanates from D brane modes. Another possibility is that the anomalous $U(1)$ in the effective four dimensional theory is a linear combination of $U(1)$s on the two 10 dimensional boundaries of the $E_8 \times E_8$ heterotic string. In this case the axion which cancels the anomaly is partially in the $S$ and $T$ moduli of the theory. As a consequence the scale of the anomaly becomes arbitrary.
\( + \lambda_1 \mathbf{16}_1 \ A \ \mathbf{16}_1 + \lambda_2 \mathbf{16}_2 \ A \ \mathbf{16}_2 + \lambda \ S \ \mathbf{16}_1 \ \mathbf{16}_2 \) \quad (1)

The adjoint \( A \) is assumed to get a vacuum expectation value [vev]

\[ < A > = (B - L) \ M_G \] \quad (2)

where \( B - L \) (baryon number minus lepton number) is non-vanishing on color triplets and zero on weak doublets in the \( \mathbf{10} \) and the singlet \( S \) is assumed to get a vev

\[ < S(x, \theta) > = S + \theta^2 \ F_S \] \quad (3)

We take \( S \sim 10^{14} \) GeV with the effective SUSY breaking scale in the observable sector given by

\[ \Lambda = F_S / S \sim 10^5 \ \text{GeV}. \] \quad (4)

The field \( \mathbf{10}_H \) includes the weak doublet and color triplet Higgs, while \( \mathbf{10} \) contains the minimal set of messengers. The first two terms in \( W \) correspond to the natural doublet-triplet splitting mechanism in \( \text{SO}(10) \| 10 \). Note that the doublet messengers obtain mass at the scale \( S \), while the triplet messengers obtain mass at the GUT scale. This theory also includes a natural mechanism for suppressing proton decay when \( S \ll M_G \| 11 \). Finally, the last three terms in the superspace potential include four additional sets of messengers. They are introduced solely to give gluinos mass at one loop.\(^2\)

Note, the R symmetry which keeps gauginos massless is spontaneously broken at \( S \).

In this theory, gauginos obtain mass at one loop

\[
\begin{align*}
    m_{\tilde{g}} &= \frac{\alpha_3(M)}{\pi} b^2 \ \Lambda \\
    M_2 &= \frac{\alpha_2(M)}{4\pi} \Lambda(1 + 4b^2) \\
    M_1 &= \frac{3}{8} \frac{\alpha_1(M)}{4\pi} \Lambda(1 + \frac{20}{3}b^2)
\end{align*}
\]

where the parameter \( b = \lambda S / \sqrt{\lambda_1 \lambda_2} \ M_G \) is derived from the last 3 terms in \( W \). The parameter \( b \) is naturally of order \( 10^{-2} \lambda / \sqrt{\lambda_1 \lambda_2} \). For the purposes

\(^2\)The first two terms do not contribute to gaugino masses at one loop due to an accidental cancellation as shown by one of us (K.T.).
of our analysis, we take $b = 0.1$. However $b$ is clearly a free parameter which may be used to vary the gluino mass.

Scalars obtain mass at two loops

$$\tilde{m}^2 = 2\Lambda^2\left\{C_3\left(\frac{\alpha_3(M)}{4\pi}\right)^2(a^2 + 4b^2) + C_2\left(\frac{\alpha_2(M)}{4\pi}\right)^2(1 + 4b^2) + C_1\left(\frac{\alpha_1(M)}{4\pi}\right)^2\left(\frac{3}{5} + \frac{2}{5}a^2 + 4b^2\right)\right\}$$  (6)

where $C_3 = 4/3$ for color triplets and zero otherwise, $C_2 = 3/4$ for weak doublets and zero otherwise, $C_1 = \frac{3}{5}(Y/2)^2$ and $a = \lambda_S/\lambda_a M_G$. Note, the parameter $a$ derives from the first two terms in $W$. In our analysis, we take $a = 0.01$. The arbitrary parameters $a$ and $b$ are clearly independent.

Finally the gravitino mass is generically given by

$$m_{3/2} = F_S/\sqrt{3} M_{Pl}$$  (7)

with the reduced Planck scale, $M_{Pl} = 1/\sqrt{8\pi G_N} = 2.4 \times 10^{18}$ GeV. Note, supergravity corrections to squark and slepton masses are proportional to the gravitino mass.

In the particular model of SUSY breaking discussed in [6], the field which gets both a scalar and F component vev is the third component of an SU(2)$_F$ vector field, $S_3$. In this theory, the field $S$ is a composite superfield, with $S = S_3^2/M_{st}$ and $F_S = 2S_3 F_{S_3}/M_{st}$. The scale $M_{st}$ is associated with a string scale of order $M_G$. Hence with a maximum value of $S_3 \sim 10^{15}$ GeV we obtain $S \sim 10^{-2} M_G \sim 10^{14}$ GeV. The upper bound on $S_3 \sim 10^{15}$ GeV has been assumed in order to suppress flavor changing neutral current processes induced by supergravity corrections to squark and slepton masses. With the gravitino mass now given by$^3$

$$m_{3/2} = F_{S_3}/\sqrt{3} M_{Pl},$$  (8)

the ratio of the supergravity contribution to squark and slepton masses (fixed by the gravitino mass) to the GMSB contribution scales as

$$m_{3/2}/M_2 = \frac{2\pi}{\sqrt{3} a_2} (S_3/M_{Pl}) \sim 10^{3}(S_3/M_{Pl}) \sim 0.04.$$  (9)

$^3$The gravitino mass is set by the largest SUSY breaking vev in the theory. The D term contribution is of order the weak scale and thus is negligible. In the particular SUSY breaking sector of the theory given in [6], the fundamental SUSY breaking vev is given by $F_{S_3} = \frac{M_G}{S_3^2} F_S \sim 10 F_S$. This explains the change in going from eqn. 7 to 8. It is clear that the value of the gravitino mass is model dependent.
Hence with $S_3 \leq 10^{15}$ GeV, supergravity gives at most a 0.16% correction to squark and slepton masses squared. This may be sufficient to suppress large flavor changing neutral current processes.

As an example of the spectrum at $M$ we have (for $\Lambda = 10^5$ GeV)

\[
\begin{align*}
m_{\tilde{g}} &= 13 (b/1.1)^2 \text{ GeV} \\
M_2 &= 330 \text{ GeV} \\
m_{3/2} &= 12 \left( S_3/10^{15}\text{GeV} \right) \text{ GeV.}
\end{align*}
\]

(10)

In addition, scalar masses are fixed by the largest SUSY breaking scale in the problem, $M_2$; with right-handed squarks and sleptons being the lightest scalars obtaining mass only via weak hypercharge interactions.

Thus, in this model the gluino (or possibly the gravitino) is naturally the LSP. The gluino mass is suppressed by two powers of $S/M_G$. One power comes because the color triplet messengers have mass of order $M_G$ and the second comes because the R symmetry in this sector of the theory is only broken at the lower messenger scale $S$.

### 2.2 D term SUSY breaking

For phenomenological reasons we assume that SUSY is also broken by the D term of an anomalous $U(1)_X$ gauge symmetry. In a recent paper[6] we have shown that it is possible to obtain D term SUSY breaking and GMSB with contributions to scalar masses of the same order. In this paper we simply parametrize the D term contribution to scalar masses by the formula

\[
\delta_D \tilde{m}_a^2 = Q_a^X d M_2^2
\]

where $Q_a^X$ is the $U(1)_X$ charge of the field labelled by the index $a$ and $d$ is an arbitrary parameter of order 1 which measures the relative strength of D term vs. gauge mediated SUSY breaking.

Since we are working in the context of an SO(10) GUT, $Q_X$ necessarily commutes with the SO(10) generators. In order not to consider a random $U(1)$ gauge symmetry we shall assume that $U(1)_X$ is embedded into $E_6$ with

\[
Q_a^X = \begin{cases} 
1 & -2 & 4 \\
10 & 1 & \subset 27 \text{ of } E_6.
\end{cases}
\]

(12)
This U(1)$_X$ is naturally family independent and is thus safe from inducing large flavor changing neutral current processes. It is also well-motivated in the context of string theories[12].

2.3 Summary of boundary conditions at $M$

Thus the boundary conditions at the messenger scale $M$ from GMSB (for scalar and gaugino masses) and from D term SUSY breaking (for scalar masses only) are given in eqns. 5, 6, and 11. There are 3 additional parameters — the supersymmetric Higgs mass $\mu$, and the two soft SUSY breaking parameters, i.e. a scalar Higgs mass $m_3^2$ and the trilinear scalar coupling $A$. To leading order we have [13]

$$ A = 0 \quad . $$ (13)

On the other hand, for the purpose of this analysis, we leave $\mu$ and $m_3^2$ arbitrary. In principle, these parameters must be determined by new physics which we have not considered in this analysis. Using the 6 parameters $\Lambda, a, b, d, \mu, m_3^2$ defining the boundary conditions at $M = 10^{14}$ GeV (eqns. 5, 6, 11 and 13; $\mu$ and $m_3^2$ arbitrary) we renormalize the full set of dimensionful (dimensionless) parameters at one (two) loop down to $M_Z$. $\mu$ and $m_3^2$ are fixed by requiring electroweak symmetry breaking. For the bottom mass, we have included three loop QCD running of $m_b(\mu)$ from $\mu = m_b$ to $m_Z$ and one loop SUSY threshold corrections at $m_Z$, relevant for large $\tan \beta$[14].

3 Sparticle masses at the Z scale

3.1 Why introduce D term SUSY breaking?

Consider first the low energy spectrum in the absence of D term scalar masses, i.e. $d = 0$. In fig. 1, we show the squark, slepton, Higgs and gaugino masses for the case $M_2 = 110$ GeV, $\mu < 0$, $a = 0.01$, and $b = 0.1$ plotted as a function of $\tan \beta$. Note, the one-loop corrected light Higgs mass (the solid brown line) is below the present experimental lower bound, $m_h > 90$ GeV[15]. We find $m_h < 60$ GeV. This is due to the fact that at tree level the Higgs
Figure 1: Mass spectrum of squarks, sleptons, neutralinos, charginos, gluino, and Higgs for the case $M_2 = 110$ GeV, $\mu < 0$, $a = 0.01$, and $b = 0.1$ plotted as a function of $\tan \beta$. The right-handed up squarks (green), down squarks (red), sleptons (pink), neutralinos (blue), charginos (black), and gluino (dot-dashed green) masses are given by the respective colored lines. The first and second (third) generation squark and slepton masses are represented by solid (dashed) lines. The lightest Higgs mass is given by the solid (dashed) brown line at one loop level (tree level).
mass satisfies the approximate mass formula

\[ m_h \sim m_Z \cos 2\beta \]  \hspace{1cm} (14)

with \( \tan \beta \leq 2 \) in fig. 1. We could in principle increase the Higgs mass by increasing \( \tan \beta \), however, due to the relation

\[ \mu^2 = \frac{m_Z^2}{2} + \frac{m_{H_d}^2 - m_{H_u}^2 \tan^2 \beta}{\tan^2 \beta - 1} \]  \hspace{1cm} (15)

as \( \tan \beta \) increases, \( \mu \) decreases and as a consequence the chargino mass also decreases. The lower bound on the chargino mass \( m_{\tilde{\chi}^+} \geq 85 \text{ GeV}[16] \) is violated for \( \tan \beta > 1.83 \).\(^4\) Of course, we could in principle increase the chargino mass by increasing \( M_2 \), however, the right-handed SUSY breaking stop mass squared is negative at the weak scale and is driven further negative when \( M_2 \) increases since this is the dominant source of SUSY breaking. For \( M_2 \geq 150 \text{ GeV} \), the lightest stop mass becomes less than \( m_Z/2 \) (see fig. 2). Thus we cannot make \( M_2 \) very large. Hence since both \( \tan \beta \) and \( M_2 \) are severely constrained from above, the light Higgs mass is always below the experimental bound. We have checked that the problem of a light Higgs, chargino and stop is even worse for \( \mu > 0 \).

In order to solve this problem we have been forced to consider additional soft SUSY breaking mass contributions. In order not to lose predictability, we have considered the well motivated addition of an anomalous \( U(1)_X \) and its associated D term. We now consider how this one new contribution resolves the light Higgs problem.

### 3.2 The addition of D term scalar mass corrections

As discussed, we parametrize the D term contribution to soft SUSY breaking scalar masses at the messenger scale by

\[ \delta_D \tilde{m}_a^2 = Q_a^X d M_2^2 \]  \hspace{1cm} (16)

with

\[ Q^X = 1 \]  \hspace{1cm} (17)

\(^4\)This bound assumes the standard missing energy signature for SUSY. This limit is not applicable in our case. However, there will be a lower bound on the chargino mass coming from the visible Z width or \( e^+e^- \to \text{hadrons} \) which is comparable. We will discuss these limits on the chargino mass within our framework shortly.
Figure 2: Mass spectrum of squarks, sleptons, neutralinos, charginos, gluino, and Higgs for the case $\tan \beta = 1.9$, $\mu < 0$, and $b = 0.1$ as a function of the wino mass ($M_2$). The line notation is the same as in fig.1.
for squarks and sleptons and

$$Q^X = -2$$

(18)

for Higgs doublets. Thus squarks and sleptons obtain a universal positive mass squared correction. As a result we can now increase $M_2$ without the light stop mass becoming too small. Then increasing $M_2$ causes the chargino mass $m_{\tilde{\chi}^+}$ to increase. Finally $\tan \beta$ can now be increased which raises the light Higgs mass above the experimental lower bound.

In figs. 3(a,b) we plot the sparticle spectrum for $\tan \beta = 10$, $M_2/m_Z = 2$ as a function of $d$. We see in fig. 3a that $98 \lesssim m_h \lesssim 103$ GeV for $0.42 \lesssim d \lesssim 4.0$ (the lightest Higgs mass is given by the solid (dashed) brown line at one loop (tree level). Thus the light Higgs problem is solved with the addition of one new parameter, $d$. The Higgs mass is determined by electroweak symmetry breaking and is thus only weakly dependent on $d$. We also see that the mass of the other Higgs states, $A$, $H_0$, $H^+$, (evaluated at tree level) decrease as $d$ increases, with the pseudo-scalar $A$ remaining the lightest and for $d \lesssim 3.5$, we have $m_A \gtrsim 150$ GeV. Squark and slepton masses (fig. 3b), on the other hand, increase with $d$. The light first and second generation squarks and sleptons are right handed; the heavy ones are left handed. The third generation squarks and slepton eigenstates are mixtures of left and right handed states due to $A$ term mixing. The lightest squark is either a stop or sbottom. Note, for $d \leq 0.7$ the light stop is lighter than the top. This is because the right handed stop mass squared is driven negative by renormalization group running as a consequence of the large top yukawa coupling. This is the same effect which drives the Higgs mass squared negative. Finally chargino and neutralino masses implicitly depend on $d$ through their dependence on $\mu$. For small $d$ the lightest neutralino and chargino masses are naturally small. The lower bound on $d$ is obtained by requiring that the visible width of the $Z$ be consistent with the experimental bound (see next section). Recall, the value of the gluino mass depends on $b$. In the figures we have taken $b = 0.1$; the gluino mass scales as $b^2$.

In figs. 4, 5 we show the sparticle spectrum as a function of $\tan \beta$ and $M_2$ with the other parameters fixed. In figs. 4(a,b) we take $M_2/m_Z = 2$, $d = (0.8, 0.5)$, $b = 0.1$ and $\mu > 0$. We see (fig. 4a) that for $4 \lesssim \tan \beta \lesssim 39$ the light Higgs mass bound is satisfied. For low and moderate values of $\tan \beta$ the pseudo-scalar ($A$) and charged Higgs ($H^+$) are significantly heavier than
Figure 3 (a): Mass spectrum of the light neutral Higgs $h$ (solid (dashed) brown) at one loop level (tree level), pseudoscalar boson ($A$) (red), heavy neutral Higgs ($H_0$) (blue), and charged Higgs ($H^+$) (black) for the case $M_2/m_Z = 2$, $\tan\beta = 10$, $b = 0.1$ and $\mu > 0$ as a function of $d$. 

$M_2/m_Z = 2$, $\tan\beta = 10$, $b = 0.1$, $\mu > 0$
Figure 3 (b): Mass spectrum of up squarks (green), down squarks (red), sleptons (pink), sneutrinos (orange), neutralinos (blue), charginos (black), and gluino (dot-dashed green) for the case $M_2/m_Z = 2$, $\tan \beta = 10$, $b = 0.1$, and $\mu > 0$ as a function of $d$. The first and second (third) generation squark and slepton masses are given by solid (dashed) lines.
Figure 4 (a): Mass spectrum of squarks, sleptons, neutralinos, charginos, gluino, and Higgses for the case $d = 0.8, \frac{M_2}{m_Z} = 2, b = 0.1, \mu > 0$ as a function of $\tan \beta$. The line notation is the same as in fig.1.
Figure 4 (b): Same as fig.4 (a) except for $d = 0.5$.
Figure 4 (c): The mass spectrum of the heavy states (left-handed states) of the squarks and slepton (the line notation is same as in fig.3(b)), and pseudoscalar boson (solid brown line), heavy neutral Higgs (dotted brown), and charged Higgs (dashed brown) for the case $M_2/m_Z = 2$, $d = 0.8$, $b = 0.1$, and $\mu > 0$ as a function of $\tan \beta$. 
m_Z. Hence the phenomenology of the light Higgs is identical to that of the SM. However for large \( \tan \beta \sim 30 \) the bottom yukawa coupling is relevant. In this case, the pseudo-scalar mass \( m_A^2 = m_{H_u}^2 + m_{H_d}^2 \) becomes small (fig. 4c). The upper bound on \( \tan \beta \) is of order 38, since in this case both \( m_h \) and \( m_A \) are greater than \( \sim 75 \) GeV, the present experimental bound[15]. In figs. 5(a,b) we take \( \tan \beta = (10, 30), d = 0.8, b = 0.1 \) and \( \mu > 0 \). We see that the light Higgs mass increases slowly as \( M_2 \) increases with \( m_h \lesssim 115 \) GeV. Squark, slepton and gaugino masses also increase as \( M_2 \) increases. This is because all SUSY breaking masses scale with \( \Lambda \). For low values of \( M_2 \) of order 100 GeV, the lightest neutralino and chargino masses become small; the lower bound on \( M_2 \) is again determined by the visible width of the Z.

The gluino mass is arbitrary. Its value is bound from below by about 1 GeV since, in addition to the GMSB contribution to its mass (eqn. 5), it receives a dynamical QCD contribution to its mass and also radiative corrections at the weak scale dominated by the stop; both of order a GeV. The gluino mass also increases as \( M_2 \) increases and for very large \( M_2 \) of order a TeV (fig. 6) the gluino mass is above 100 GeV.

For a large range of the parameter \( b \), either the gluino or the gravitino is the LSP. In the case that the gravitino is the LSP, then the gluino is unstable with the decay rate for a gluino decaying into a goldstino (G) and a gluon given by

\[
\Gamma(\tilde{g} \to G + g) = \frac{m_{\tilde{g}}^3}{16\pi F_{\Sigma}} \left(1 - \frac{m_G^2}{m_{\tilde{g}}^2}\right)^3
\]

or

\[
\tau_{\tilde{g}} \sim 10^{20} \left(\frac{25 \text{GeV}}{m_{\tilde{g}}}\right)^5 \left(\frac{\sqrt{F_{\Sigma}}}{6 \times 10^9 \text{GeV}}\right)^4 \text{cm}
\]

where in eqn. 20 we have neglected the gravitino mass. Hence for all laboratory experiments the gluino may be considered stable, decaying outside the lab.
Figure 5 (a): Mass spectrum of squarks, sleptons, neutralinos, charginos, gluino, and Higgses for the case $\tan \beta = 10$, $d = 0.8$, $b = 0.1$, and $\mu > 0$ as a function of $M_2$. The line notation is the same as in fig.1.
Figure 5 (b): Same as fig.5 (a) except for $\tan \beta = 30$. 

$tan\beta = 30, d = 0.8, b = 0.1, \mu > 0$
Figure 6: Same as fig.5 (a)
detector. Thus the standard missing energy signature for SUSY is gone. The
 gluino will hadronize; forming a hadronic jet. For a heavy gluino, the jet
 may contain only the lightest stable hadron containing the gluino. In any
 event, some fraction of the gluino kinetic energy will be visible in a hadronic
 calorimeter.

We assume that the LSP is a gluino–gluon bound state, a glueballino
 \((R_0)\).\(^6\) It is (for all practical purposes) stable because of a conserved \(R\)
 parity. \(R_0\) will interact in a hadronic calorimeter with a strong interaction
 cross-section. The dominant process at small momentum transfer is governed
 by Regge exchange. In a recent paper by Baer et al.\([17]\) an estimate for the
 energy loss of \(R_0\) in a hadronic calorimeter was obtained using the triple
 pomeron amplitude for the single diffractive process \(R_0 \, N \rightarrow R_0 \, X\) where
 \(X\) denotes the inclusive sum over all final states. It was found that a 25
 GeV gluino with \(\beta \sim 1\) would deposit less than 20% of its kinetic energy
 in the detector assuming 8 hadronic collisions. An \(R_0\), however, could also
 have a significant charge exchange cross-section, given by a triple- Reggeon
 amplitude for the single diffractive process \(R_0 \, N \rightarrow \tilde{\rho}^+ \, X\). In this case
 there would be an additional ionization energy loss in the detector, assuming
 \(\tilde{\rho}^+\) is sufficiently long lived. Baer et al. have assumed that at each hadronic
 collision the light brown muck surrounding the gluino is stripped off and then
 the gluino re-hadronizes. They parametrize the probability for the resulting
 gluino bound state to be a charged \(\tilde{\rho}^+\) by \(P\). We shall discuss their results
 in the next section. It suffices now to say that their results are sensitive to
 this universal parameter \(P\). Clearly a better theoretical estimate of \(R_0 \, N\)
 scattering is needed. It is also clear that detailed detector simulations are
 necessary in order to confirm their results.

We can now test the theory with regards to LEP, CLEO and Tevatron
 data. Note that flavor changing neutral current processes are naturally sup-
 pressed within our framework of combined gauge-mediated and D term SUSY
 breaking.

\(^6\)The other likely candidate LSP is a gluino–\((u \bar{u} – d \bar{d})_8\) bound state, the neutral
 component of an isotriplet \(\tilde{\rho}\). We shall not consider this possibility further in this paper, see
 for eg. \([5]\).
3.3 Experimental tests

First let’s consider some tests from LEP. Since the LSP in this theory is a gluino, charginos and neutralinos can now decay into $q\bar{q}g$. Moreover since charginos and neutralinos are relatively light they can be produced at relatively low center of momentum energies at LEP. We have considered limits from two processes — the visible width of the Z ($\Delta\Gamma_Z$(visible)) and $e^+e^- \rightarrow$ hadrons. In figs. 3 - 6, the lower bound on $(d, \tan\beta, M_2)$, respectively is determined by the allowed contribution to $\Delta\Gamma_Z$(visible) $\leq 23.1 \times 10^{-3}$ GeV[18]. This provides a weak lower bound on neutralino and chargino masses. A better bound can be obtained by $e^+e^- \rightarrow$ hadrons. In figs. 7(a,b) we give the constraints on the chargino mass coming from the process $e^+e^- \rightarrow$ hadrons. The horizontal line is the OPAL $2\sigma$ bound on new physics at $\sqrt{s} = 172$ GeV[19]. In fig. 7(a) we show the contribution to $e^+e^- \rightarrow$ hadrons due to chargino and neutralino production times the branching fraction for quark decay as a function of the chargino mass (varying $d$) for fixed $\tan\beta = 10, M_2/m_Z = 2, \mu > 0, b = 0.1$. We find a lower bound $d \geq 0.44$. This translates into a lower bound on the lightest chargino mass $m_{\tilde{\chi}^+_1} \geq 66$ GeV (fig. 7(a)) and neutralino mass $m_{\tilde{\chi}^0_1} \geq 34$ GeV (fig. 3). In fig. 7(b) we show the same constraints now varying $M_2$ instead of $d$. For $d = 0.8$ and $M_2 > 122$ GeV, we find $m_{\tilde{\chi}^+_1} > 71$ GeV and $m_{\tilde{\chi}^0_1} > 27$ GeV.

An additional test from LEP is on the number of 4 jet events. The dominant decay mode for down squarks is $\tilde{d} \rightarrow d\tilde{g}$. In fig. 4(b) we see that the bottom squark can be as light as 75 GeV for large $\tan\beta$. Nevertheless the calculated cross section for $e^+e^- \rightarrow \tilde{b}\tilde{b}$ is significantly below the OPAL bound for excess four jet events at $\sqrt{s} = 184$ GeV[20]. Four jets are also possible via the direct process $e^+e^- \rightarrow q\bar{q}g\bar{g}$. Baer et al.[17] find strong limits using OPAL and L3 data[21] for $e^+e^- \rightarrow Z$ with $Z \rightarrow (2$ or $4)$ jets + missing energy. This data was used to place limits on the CMSSM process $e^+e^- \rightarrow \chi^0_1 \chi^0_2$ with the subsequent decay $\chi^0_2 \rightarrow q\bar{q} \chi^0_1$. They argue that a gluino mass in the range $5 \leq m_{\tilde{g}} \leq 25$ GeV is ruled out. This result seems to be relatively insensitive to the probability $P$. Results from LEP2 at higher center of mass energies do not improve upon this limit.

We have evaluated the rate for $b \rightarrow s\gamma$. The total contribution for $\mu < 0$

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Here we assume that the gluino produces a visible jet.
Figure 7 (a): The cross section of $e^+e^- \rightarrow \text{hadrons}$ at center of mass energy of 172 GeV from the chargino and neutralino pair production for the case $M_2/m_Z = 2$, $\tan \beta = 10$, $b = 0.1$, and $\mu > 0$ as a function of the chargino mass (varying $d$). We also show the OPAL $2\sigma$ bound on new physics (dot-dashed line).
Figure 7 (b): Same as fig. 7 (a) for the case $\tan \beta = 10$, $d = 0.8$, $b = 0.1$, and $\mu > 0$ as a function of the chargino mass (varying $M_2$).
Figure 8 (a): The ratio of $b \to s\gamma$ amplitude of SUSY contribution to Standard Model one. The colored lines represent the contribution of charged Higgs-top (red), chargino-stop (blue dashed), gluino-sbottom (pink long dashed), neutralino-sbottom (green dot-dashed), and total (all SUSY contribution plus SM contribution) (black solid line) for $M_2/m_Z = 2$, $\tan \beta = 10$, $b = 0.1$, and $\mu > 0$ as a function of $d$. 

(a)
Figure 8 (b): Same as (a) for $M_2/m_Z = 2, d = 0.8, b = 0.1, \mu > 0$ as a function of $\tan \beta$. 

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Figure 8 (c): Same as (a) for $\tan \beta = 10$, $d = 0.8$, $b = 0.1$, and $\mu > 0$ as a function of $M_2$. 
Figure 8 (d): Same as (a) for $\tan \beta = 30$, $d = 0.8$, $b = 0.1$, and $\mu > 0$ as a function of $M_2$. 
is always above the SM result, approaching it asymptotically for large $M_2$.\footnote{This is the reason figs. 3 - 6 are given only for $\mu > 0$.} The results for $\mu > 0$ are given in figs. 8(a,b,c,d). It is clear that there is a large range of parameters for which the total contribution is below the SM value. Our result however only includes one loop analysis and should be taken only as an indication that it is possible to obtain results consistent with the data\cite{22}.

For $d < 0.7$ (fig. 3) or $\tan \beta < 3.8$ (fig. 4a) or for any value of $\tan \beta$ and $d = 0.5$ (fig. 4b), we have $\tilde{m}_{t_1} < m_t$. Then for a sufficiently light gluino, the top can decay into a stop + gluino. This range of parameters is testable at the Tevatron. In the case that the stop is heavier than the lightest chargino, the dominant decay mode of the stop is $\tilde{t}_1 \rightarrow b \tilde{\chi}_1^+$. In this case stop decay can mimic a top signal. We are now studying this situation.

It has also been argued by Baer et al.\cite{17} that CDF data for jets + missing energy places stringent limits on the gluino mass. These results however are sensitive to the parameter $P$. For $P \leq 1/2$ they rule out a gluino with mass in the range $20 \lesssim m_{\tilde{g}} \lesssim 140$ GeV. However for $P = 3/4$ there is an allowed window for the gluino mass $25 \lesssim m_{\tilde{g}} \lesssim 35$ GeV. Clearly these results are more sensitive to the details of how an $R_0$ or $\tilde{\rho}^+$ interacts in the detector. For example, as discussed in \cite{17}, a $\tilde{\rho}^+$ has some probability of being mis-identified as a muon if it is seen in the muon tracking chambers. It is important that this analysis be redone by both CDF and D0. In addition, Baer et al. have explicitly assumed that squarks are very heavy and hence do not contribute to the jet + missing energy signal. Inclusion of squarks can, via coherent interference, decrease the gluino production cross section and hence the rate for jets + missing energy. At the same time, squark production and subsequent decay can increase the jets + missing energy rate. Hence the Baer et al. analysis must be redone for the model presented in this paper.

4 Conclusions

In this paper we have presented a phenomenologically acceptable model with a gluino or gravitino LSP. In case the gravitino is the LSP, then the gluino is the NLSP. It can decay but only on a cosmological time scale. Thus in either case the gluino is stable with respect to any accelerator experiment searching for SUSY. We have also discussed some experimental bounds and
possible tests of the theory at LEP, CLEO and the Tevatron. We believe that further experimental studies are now justified.

The theory includes a combination of gauge mediated and D term SUSY breaking. The latter contribution was shown to be necessary in order to obtain a Higgs with mass consistent with present experimental bounds.

The gluino is naturally light in this theory because the 10 of Higgs (in SO(10)) mixes with the 10 of messengers. It is light because of two effects. Firstly, the SU(3) triplet messengers are heavier than the electroweak doublet messengers; suppressing SUSY breaking mediated by color interactions. Secondly, R symmetry (which keeps gauginos massless) is broken at the electroweak doublet messenger scale; providing an additional suppression.

There are several other consequences of a theory with a gluino LSP which may have cosmological and/or astrophysical significance. It has been shown that electroweak baryogenesis is not possible in the SM. On the other hand, the necessary conditions for electroweak baryogenesis in the MSSM requires \( m_h \sim 100 \text{ GeV} \) and \( \tilde{m}_{t_R}^2 < 0 \)\(^{[23]}\). These conditions are naturally satisfied in our model. They guarantee that there is a first order phase transition which is sufficiently strong to shut off baryon violating interactions inside the bubbles of broken electroweak phase. If there is also sufficient CP violation; then baryogenesis is possible.

In addition, the \( R_0 \) provides a natural particle physics candidate for the source of ultra high energy cosmic rays, the UHECRON.\(^9\) It has recently been argued\(^{[24]}\) that an \( R_0 \) with mass in the range from 2 - 50 GeV could reproduce the highest energy cosmic ray shower observed by the Fly’s Eye Detector\(^{[25]}\).

Finally, the \( R_0 \) LSP is not a dark matter candidate since the annihilation cross section is too large and hence the relative abundance of \( R_0 \) to baryons is too small.\(^{10}\) There are nevertheless severe constraints on the abundance of \( R_0 \)s coming from searches for anomalous heavy isotopes of hydrogen and oxygen\(^{[26]}\). If \( R_0 \) binds to hydrogen, then any \( R_0 \) with mass greater than 2 GeV is ruled out. In addition, if the concentration of \( R_0 \) in oxygen is greater

\(^{9}\)This is even possible if the gluino is the NLSP and decays into a gravitino LSP. For example, a gluino with mass \( \sim 25 \text{ GeV} \) and energy characteristic of the highest energy cosmic rays \( \sim 10^8 \text{ TeV} \) can travel \( \sim 10^5 \text{ Mpc} \) before it decays.

\(^{10}\)The uncertainty in the annihilation cross section is quite large \( \sim 10^{\pm 1.5} \). Nevertheless the relative abundance of gluinos to baryons is of order \( (10^{-10} - 10^{-7})(m_\tilde{g}/1\text{GeV}) \). For a recent re-analysis, see Baer et al.\(^{[17]}\)
than one part in $10^{16} - 10^{19}$, it is also ruled out. Note however, if the gravitino is the LSP and $R_0$ is the NLSP, these limits are evaded.

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