Closing the Light Gluino Window in Supersymmetric Grand Unified Models

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Abstract

We study the light gluino scenario giving special attention to constraints from the masses of the light CP-even neutral Higgs $m_h$, the lightest chargino $m_{\chi^\pm_1}$, and the second lightest neutralino $m_{\chi^0_2}$, and from the $b \to s\gamma$ decay. We find that minimal $N = 1$ supergravity, with a radiatively broken electroweak symmetry group and universality of scalar and gaugino masses at the unification scale, is incompatible with the existence of a light gluino.

Revised Version
Discussions about the existence of a light gluino have been in the literature for a long time\footnote{1}. Recently, motivated by the discrepancy between the value of the strong coupling constant determined by low energy deep inelastic lepton-nucleon scattering, and the one determined by high energy $e^+e^-$ LEP experiments, there has been a renewed interest in this possibility\footnote{2-4}.

As it was explained by H.E. Haber\footnote{5}, the light gluino window is:

$$2.6 \lesssim m_{\tilde{g}} \lesssim 3 \text{ GeV},$$

where the lower limit comes from the non-observation of a pseudoscalar $\tilde{g}\tilde{g}$ bound state in quarkonium decays, and the upper limit follows from an analysis of CERN $p\bar{p}$ Collider data\footnote{6}.

In supersymmetric grand unified theories (SUSY-GUT)\footnote{7}, the three gaugino masses $M_s$, $M$, and $M'$ are different at the weak scale but equal to a common gaugino mass $M_{1/2}$ at the grand unification scale $M_X$. The difference at the weak scale is due to the fact that the evolution of the three masses is controlled by different renormalization group equations (RGE). The approximated solution of these RGE is:

$$M_s \approx M_{1/2} \left[ 1 + \frac{3g_s^2}{8\pi^2} \ln \frac{M_X}{m_Z} \right], \quad m_{\tilde{g}} = |M_s|,$$

$$M \approx M_{1/2} \left[ 1 - \frac{g^2}{8\pi^2} \ln \frac{M_X}{m_Z} \right],$$

$$M' \approx M_{1/2} \left[ 1 - \frac{11g'^2}{8\pi^2} \ln \frac{M_X}{m_Z} \right],$$

where we are neglecting the supersymmetric threshold effects. Taking $M_X = 10^{16}$ GeV, we find that $M \approx 0.30 m_{\tilde{g}}$ and $M' \approx 0.16 m_{\tilde{g}}$ at the weak scale.

Similarly, the scalar masses are also degenerate at the unification scale, and equal to $m_0$. The RGE make both the Higgs mass parameters $m_1$ and $m_2$, and the squark and slepton mass parameters, evolve differently. A third independent parameter at the unification scale is the mass parameter $B$. This mass defines
the value of the unified trilinear mass parameter $A$ at $M_X$ by $A = B + m_0$, a relation valid in models with canonical kinetic terms. Moreover, it also defines the third Higgs mass parameter $m_{12}^2 = -B\mu$, valid at every scale, where $\mu$ is the supersymmetric Higgs mass parameter. The set of independent parameters we choose to work with, given by $M_{1/2}$, $m_0$, and $B$ at the unification scale, is completed by the value of the top quark Yukawa coupling $h_t = g m_t/(\sqrt{2} m_W \tan \beta)$ at the weak scale. Here the angle $\beta$ is defined through $\tan \beta = v_2/v_1$, where $v_1$ and $v_2$ are the vacuum expectation values of the two Higgs doublets.

Knowing the parameters of the Higgs potential at the weak scale $m_1^2$, $m_2^2$, and $B$, we can calculate the more familiar parameters $m_t$, $t_\beta$, $m_A$, and $\mu$, for a given value of the top quark Yukawa coupling $h_t$, through the following formulas

\begin{align}
m_{1H}^2 &\equiv m_1^2 + \mu^2 = -\frac{1}{2} m_Z^2 c_2\beta + \frac{1}{2} m_A^2 (1 - c_2\beta), \\
m_{2H}^2 &\equiv m_2^2 + \mu^2 = \frac{1}{2} m_Z^2 c_2\beta + \frac{1}{2} m_A^2 (1 + c_2\beta), \\
m_{12}^2 &= -B\mu = \frac{1}{2} m_A^2 s_2\beta,
\end{align}

where $s_{2\beta}$ and $c_{2\beta}$ are sine and cosine functions of the angle $2\beta$, and it is understood that all the parameters are evaluated at the weak scale. We alert the reader that for a given set of values $M_{1/2}$, $m_0$, $B$, and $h_t$ there may exist more than one solution for the parameters at the weak scale $m_t$, $t_\beta$, $m_A$, and $\mu$. According to ref. [4], and we will confirm this, the relevant region of parameter space in the light gluino scenario is characterized by low values of the top quark mass and values of $\tan \beta$ close to unity. Considering the low values of the top quark mass relevant for our calculations, radiative corrections to the chargino and neutralino masses (recently calculated in ref. [8]) will have a minor effect.

The region $\tan \beta$ close to unity has been singled out by the grand unification condition $m_b = m_\tau$ at $M_X$ [9], and was analyzed in detail in ref. [10]. Here we stress the fact that if $\tan \beta = 1$, the lightest CP-even neutral Higgs is massless at tree level. Nevertheless, the supersymmetric Coleman-Weinberg mechanism [11] generates a mass $m_h$ different from zero via radiative corrections. The
fact that $m_t$ is also small will result in a radiatively generated $m_h$ close to the experimental lower limit $m_h \gtrsim 56$ GeV, valid for $m_A > 100$ GeV\textsuperscript{[12]}. Therefore, experimental lower limits on $m_h$ impose important restrictions on the light gluino window.

It has been pointed out that the branching ratio $B(b \rightarrow s\gamma)$ has a strong dependence on the supersymmetric parameters\textsuperscript{[13,14]}. The theoretical branching ratio must remain within the experimental bounds $0.65 \times 10^{-4} < B(b \rightarrow s\gamma) < 5.4 \times 10^{-4}$. We calculate this ratio, including loops involving $W^\pm / U$-quarks, $H^\pm / U$-quarks, $\chi^\pm / U$-squarks, and $\tilde{g}/ D$-squarks, neglecting only the contribution from the neutralinos, which were reported to be small\textsuperscript{[13]}. We also include QCD corrections to the branching ratio\textsuperscript{[15]} and one loop electroweak corrections to both the charged Higgs mass\textsuperscript{[16]} and the charged Higgs-fermion-fermion vertex\textsuperscript{[17]}.

Another important source of constraints comes from the chargino/neutralino sector. For $\tan \beta > 4$, a neutralino with mass lower than 27 GeV is excluded, but the lower bound decreases when $\tan \beta$ decreases, and no bound is obtained if $\tan \beta < 1.6$\textsuperscript{[18]}. The lower bound for the heavier neutralinos (collectively denoted by $\chi'$) is $m_{\chi'} > 45$ GeV for $\tan \beta > 3$, and this bound also decreases with $\tan \beta$ and eventually disappears\textsuperscript{[19]}. On the other hand, if the lightest neutralino has a mass $\lesssim 40$ GeV (as we will see, in the light gluino scenario, the lightest neutralino has a mass of the order of 1 GeV), the lower bound for the lightest chargino mass is 47 GeV\textsuperscript{[19]}. For notational convenience, this latest experimental bound will be denoted by $\bar{m}_{\chi^\pm_1} \equiv 47$ GeV.

In the following, we study the chargino/neutralino sector in more detail by analysing the mass matrices. The chargino mass matrix\textsuperscript{[20]} has eigenvalues denoted by $m_{\chi^\pm_i}$, $i = 1, 2$ and $m_{\chi^\pm_1} < m_{\chi^\pm_2}$. In the light gluino case we have $M \ll m_W$, and the chargino masses can be approximated by

$$m_{\chi^\pm_1,2}^2 = \frac{1}{2} \mu^2 + m_W^2 \pm \frac{1}{2} \sqrt{R} \pm \frac{2m_W^2 \mu M s_2 \sqrt{R}}{\sqrt{R}} + O(M_{1/2}),$$

where $R = \mu^4 + 4m_W^2 \mu^2 + 4m_W^4 c_2^2$. Since the lightest chargino mass is bounded
from below, we get the following constraint:

\[
m_W^4 c_{2\beta}^2 + \mu^2 \bar{m}_{\chi_1^+}^2 < \left( m_W^2 - \bar{m}_{\chi_1^+}^2 \right)^2 - \frac{4 m_W^2 \mu M s_{2\beta} (\frac{1}{2} \mu^2 + m_W^2 - \bar{m}_{\chi_1^+}^2)}{\sqrt{\mu^2 + 4 m_W^2 \mu^2 + 4 m_W^4 c_{2\beta}^2}},
\]

(5)

plus terms of \(O(M_{1/2}^2)\). This limits the values of \(\mu\) and \(\tan \beta\):

\[
\mu^2 < \bar{m}_{\chi_1^+}^2 \left( \frac{m_W^2}{\bar{m}_{\chi_1^+}^2} - 1 \right)^2 - \frac{4 m_W^2 \mu_0 M (\frac{1}{2} \mu_0^2 + m_W^2 - \bar{m}_{\chi_1^+}^2)}{\bar{m}_{\chi_1^+}^2 |\mu_0| \sqrt{\mu_0^2 + 4 m_W^2}} + O(M_{1/2}^2)
\]

\[
\implies |\mu| \lesssim (90 \mp 0.87 m_{\tilde{g}}) \text{ GeV, with } \pm = \text{sign}(\mu M),
\]

(6)

\[
|c_{2\beta}| < 1 - \frac{\bar{m}_{\chi_1^+}^2}{m_W^2} + O(M_{1/2}^2) \implies 0.46 < t_\beta < 2.2,
\]

where \(\mu_0^2 = \bar{m}_{\chi_1^+}^2 (m_W^2 / \bar{m}_{\chi_1^+}^2 - 1)^2 \approx 90\) GeV is the zeroth order solution (\(M = 0\)), and \(\mp 0.87 m_{\tilde{g}}\) correspond to the first order correction. The type of constraints given in eq. (6) were already found in ref. [4] at zeroth order, but as we will see, the neutralino sector will restrict the parameter space even more.

The neutralino mass matrix\(^{[21]}\) in the zero gluino mass limit \((M = M' = 0)\) has one eigenvalue equal to zero. Calculating the first order correction, we find for the lightest neutralino mass:

\[
m_{\chi_1^0} = M s_W^2 + M' c_W^2 + O(M_{1/2}^2) \approx 0.19 m_{\tilde{g}},
\]

(7)

where we used the relations between \(M, M',\) and \(m_{\tilde{g}}\) given below eq. (2). Considering eq. (1) we get

\[
0.49 \lesssim m_{\chi_1^0} \lesssim 0.57 \text{ GeV.}
\]

(8)

This light neutralino (the lightest supersymmetric particle, or LSP) is, up to terms of \(O(M_{1/2}^2/m_Z^2)\), almost a pure photino, and there is no bound on its mass from LEP collider data. Nevertheless, in the case of a stable LSP (R-parity conserving
models), ref. [4] pointed out some cosmological implications that make this scenario less attractive. On the other hand, the possibility of having a small amount of R-parity violation is not ruled out, in which case the LSP would not be stable[3].

The other three neutralino masses are, in first approximation, solutions of the cubic equation

\[ m_{\chi^0}^3 - (\mu^2 + m_Z^2)m_{\chi^0} - s_{2\beta}\mu m_Z^2 = 0. \]  

(9)

According to eq. (6), the value of \( \tan \beta \) will be close to unity, i.e., \( s_{2\beta} \approx 1 \). If we expand around this value we get for the other neutralino masses:

\[ m_{\chi^0_2} = -\mu - \frac{m_Z^2(1 - s_{2\beta})}{2\mu^2 - m_Z^2}, \]

\[ m_{\chi^0_{3,4}} = m_\pm + \frac{m_Z^2(\mu + m_\pm)(M_W^2 + M_W' s_W^2)}{3\mu m_Z^2 + 2(\mu^2 + m_Z^2)m_\pm} - \frac{\mu m_Z^2 m_\pm(1 - s_{2\beta})}{3\mu m_Z^2 + 2(\mu^2 + m_Z^2)m_\pm}, \]  

(10)

where \( m_\pm \equiv \frac{1}{2}\mu \pm \frac{1}{2}\sqrt{\mu^2 + 4m_Z^2} \) and we neglect terms of \( O(1 - s_{2\beta})^2 \) and \( O(M_{1/2}^2) \). It is understood that if an eigenvalue of the neutralino mass matrix is negative, a simple rotation of the fields will give us a positive mass. The approximation in eq. (10) breaks down when \( \mu^2 \approx \frac{1}{2}m_Z^2 \) except for \( t_\beta = 1 \).

Now we turn to the exact numerical calculation of the chargino and neutralino masses. In Fig. 1 we plot contours of constant masses in the \( \mu - t_\beta \) plane. The curve \( m_{\chi^\pm_1} = 47 \text{ GeV} \) corresponds to the constraint expressed in eq. (5). We also plot contours defined by \( m_{\chi^0_2} = 5 - 45 \text{ GeV} \), and the \( \tan \beta \) dependent experimental bound on \( m_{\chi^0_2} \) is represented by the solid line that joins the crosses. In this way, the “allowed” region (including chargino/neutralino searches only) corresponds to the region below the two solid lines. For \( \mu < 0 \) the allowed region is almost an exact reflection. The approximate bounds for \( \mu \) we got in eq. (6) are confirmed numerically: \( \mu < 87.4 \text{ GeV} \) for \( m_{\tilde{g}} = 3 \text{ GeV} \). Nevertheless, the bounds on \( \tan \beta \) come only from the experimental result \( m_{\chi^\pm_1} > 47 \text{ GeV} \), and we must include also the experimental results on \( m_{\chi^0_2} \). From Fig. 1 we see that this bound restricts the model to \( \tan \beta \lesssim 1.82 \), with the equality valid for \( \mu = 49.4 \text{ GeV} \). Since for
tan β ≲ 1 there is no solution for the radiatively broken electroweak symmetry group, the allowed values of tan β in the light gluino scenario and with μ > 0 are

\[ 1 ≲ \tan β ≲ 1.82 \]  \hspace{1cm} (11)

If μ < 0, the upper bound is tan β ≲ 1.85 with the equality valid for μ = −51.8 GeV. We go on to analyze the viability of the “allowed” region in Fig. 1. We will find that the region allowed by the χ± and χ0 analysis is in fact disallowed by the experimental bound on m_h and m_t.

In ref. [22] the RGE are solved for the special case in which only the top quark Yukawa coupling is different from zero. In the case of a light gluino (M_{1/2} ≈ 0), the value of μ at the weak scale can be approximated by

\[ \frac{1}{2} m_Z^2 + \mu^2 = -m_0^2 + \frac{z - 1}{z(1 - t_{\beta}^{-2})} \left[ \frac{3m_0^2}{2} + \frac{A^2}{2z} \right], \]  \hspace{1cm} (12)

with

\[ z^{-1} = 1 - (1 + t_{\beta}^{-2}) \left( \frac{m_t}{193 \text{GeV}} \right)^2. \]  \hspace{1cm} (13)

As it was reported in ref. [4], there is a fine-tuning situation in which we can have m_0 ≫ |μ| (producing larger radiative corrections to m_h) and it is obtained when the coefficient of m_0^2 in eq. (12) is zero. Ref. [4] concluded that constraints on m_h can be satisfied in a small window around tan β = 1.88 − 1.89 (they did not consider the constraint on the second lightest neutralino). We will see that if the relation A = B + m_0 holds we do not find this type of solution (m_0 ≫ |μ|) as opposed to the case in which A = 0. However, the later is obtained for a value of the top quark mass below the value of the experimental lower bound m_t ≥ 131 GeV[23].

We survey the parameter space m_0, B, M_{1/2}, and h_t, looking for the maximum value of tan β allowed by collider negative searches in the chargino/neutralino sector, using the SUSY-GUT model described earlier. We first consider models in
which the relation $A = B + m_0$ holds. We expect maximum $\tan \beta$ to maximize $m_h$. For example, for the value $h_t = 0.87$ and $M_{1/2} = 1$ GeV (essentially fixed by the light gluino mass hypothesis) we find that $m_0 = 132.9$ and $B = -225.5$ GeV (at the unification scale) gives us $\tan \beta = 1.82$ and $\mu = 49.4$ GeV, i.e., the critical point with maximum $\tan \beta$ in the upper corner of the allowed region in Fig. 1. The values of other important parameters at the weak scale are $m_{\chi_1^\pm} \approx 47.1$, $m_{\chi_2^0} \approx 36.8$, $m_t = 131.1$, $m_A = 152.1$, and $m_\tilde{g} = 2.75$ GeV. We find a value for $B(b \rightarrow s\gamma) = 5.35 \times 10^{-4}$ consistent with the CLEO bounds. However, the lightest CP-even neutral Higgs fails to meet the experimental requirement: we get $m_h = 47.7$ GeV, inconsistent with LEP data.

From the two fixed parameters, $h_t$ and $M_{1/2}$, the one that could affect the mass of the CP-even neutral Higgs is the first one; for a fixed value of $\tan \beta$, a larger value of the top quark Yukawa coupling will give us a larger $m_t$, and this will increase $m_h$. However, $h_t$ also enters the RGE for the Higgs mass parameters, and in order to get the correct electroweak symmetry breaking, a smaller value of $m_0$ is necessary. This implies smaller squark masses, which in turn reduce $m_h$ through radiative corrections. As an example with a larger $h_t$, we have found that for $h_t = 0.97$ and $M_{1/2} = 1$ GeV, the critical point is obtained at $m_0 = 103.8$ and $B = -132.5$ GeV. As expected, the value of the top quark mass is larger ($m_t = 146.2$ GeV), but we get smaller values for the squark masses. The net effect is that now $m_h$ is even smaller, 43.5 GeV, also in conflict with the experimental lower bound. (We caution the reader that at the small values of $m_t$ and $m_0$ used here, the contributions to $m_h$ coming from the Higgs/Gauge-boson/neutralino/chargino are also important\cite{11}; we include these in our analysis.)

We go back to $h_t = 0.87$ to analyze the case $\mu < 0$. In this case the critical point, given by $\tan \beta = 1.85$ and $\mu = -51.8$ GeV, is obtained for $m_0 = 71.1$ and $B = 111$ GeV. However the light CP-even Higgs is lighter than before: $m_h = 40.4$ GeV, incompatible with LEP data.

Models in which $A$ and $B$ are independent parameters have one extra degree
of freedom that may help to satisfy the experimental constraints. According to
eq. (12) the fine tuning \( m_0 \gg |\mu| \) is obtained for \( A = 0 \). Adopting that value and
considering \( \mu > 0 \), for \( h_t = 0.87 \) and \( M_{1/2} = 1 \) GeV we obtain the critical point
for \( m_0 = 151.6 \) and \( B = -256.8 \) GeV, which implies \( m_{\chi^\pm_1} = 47.1 \), \( m_{\chi^0_2} = 36.8 \),
\( m_t = 131.1 \), \( m_A = 173.4 \), and \( m_{\tilde{g}} = 2.75 \) GeV. However, we get \( B(b \rightarrow s\gamma) = 7.15 \times 10^{-4} \) and \( m_h = 48.8 \) GeV, both inconsistent with experimental bounds.

In order to illustrate the fine-tuning we consider \( h_t = 0.77 \) and \( M_{1/2} = 1 \) GeV.
The critical point is obtained for \( m_0 = 930 \) and \( B = -9576 \) GeV. The masses
of \( \chi^\pm_1 \), \( \chi^0_2 \) and \( \tilde{g} \) are the same as before, and we also get \( m_A = 1059 \), \( m_h = 64.8 \)
GeV consistent with LEP bound, and \( B(b \rightarrow s\gamma) = 4.14 \times 10^{-4} \) consistent with
the CLEO bound, but this time it is the top quark mass that does not meet the
experimental bound: we get \( m_t = 116.0 \) GeV, incompatible with the \( D0 \) lower
bound of 131 GeV.

If \( \mu < 0 \) no big changes are found. For \( h_t = 0.87 \) the critical point, defined
now by \( \tan \beta = 1.85 \) and \( \mu = -51.8 \) GeV, is obtained for \( m_0 = 159.4 \) and \( B = 268 \)
GeV and, as before, the two quantities inconsistent with experimental results are
\( B(b \rightarrow s\gamma) = 8.42 \times 10^{-4} \) and \( m_h = 50.5 \) GeV.

Our conclusion is that N=1 Supergravity with a radiatively broken electroweak
symmetry group is incompatible with a light gluino with a mass of a few GeV. This
is valid in models where the relation \( A = B + m_0 \) holds as well as in models where
\( A \) and \( B \) are independent parameters.

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FIGURE CAPTIONS

1) Contours of constant value of the lightest chargino and the second lightest neutralino masses, for a gluino mass $m_{\tilde{g}} = 3$ GeV. The contour corresponding to the chargino mass is defined by the experimental lower bound $m_{\chi^\pm_1} = 47$. For $\chi^0_2$ we plot contour of constant mass from 5 to 45 GeV (dashed lines). The solid line that joins the crosses represent the $tan\beta$ dependent bound on $m_{\chi^0_2}$. The “allowed” region lies below the two solid lines. We are considering in this graph experimental restrictions from the chargino/neutralino searches only.