Gluino mass from dynamical supersymmetry breaking

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Abstract

We present a new mechanism for gluino mass generation in models of dynamical supersymmetry breaking. The mechanism requires two colored chiral superfields which feel a nonabelian gauge interaction such that a fermion condensate is formed at a scale of order 1 TeV. Renormalizable hidden sector models, which typically yield unacceptably light gauginos, become viable if a gauge singlet is coupled to these chiral superfields. Moreover, the interactions of the gauge singlet with the Higgs superfields substitute the $\mu$-term. Visible sector models can also incorporate this mechanism; however, the models of dynamical supersymmetry breaking previously analyzed cannot be fitted in a simple visible sector model because they lead to vacuum expectation values for charged scalars.

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1 Introduction

Supersymmetry (SUSY) ensures the stability of the electroweak scale against radiative corrections (i.e. it solves the naturalness problem) without explaining the small ratio of the electroweak scale to the Planck scale (the hierarchy problem). However, the hierarchy problem is solved if supersymmetry is broken dynamically, i.e. the vacuum is non-supersymmetric due to non-perturbative effects of some non-abelian gauge interaction, called supercolor.

There are known three classes of candidates for the supercolor sector (we include here only renormalizable models of calculable dynamical SUSY breaking): the $SU(3) \times SU(2)$ model [1] and its generalizations [2], the $SU(2n + 1)$ models [1, 3], and the $SU(2n) \times U(1)$ models [2, 4]. All these models have a spontaneously broken $R$-symmetry and, apparently, a massless $R$-axion. However, Bagger, Poppitz and Randall [5] showed that the cancellation of the cosmological constant implies the explicit breaking of the $R$-symmetry such that the $R$-axion is massive.

The outstanding phenomenological issue in this framework is to find a viable way of transmitting SUSY breaking from the supercolor sector to the minimal supersymmetric standard model (MSSM). In this approach it is difficult to produce a sufficiently large gluino mass, i.e. of the order of the weak scale (the possibility that light gluinos, with mass of order 1 GeV, are not experimentally ruled out is controversial [6]). The same problem arises in models with a continuous $R$-symmetry [7].

In the case of hidden sector models, SUSY breaking is transmitted by supergravity from the hidden sector (which includes supercolor interactions) to the visible sector (which includes the MSSM). Nonrenormalizable interactions suppressed by powers of the Planck scale, $M_P$, give masses to the scalars and gauginos of the visible sector. The scalar masses in the visible sector are of order $\Lambda_{SC}^2/M_P$, where $\Lambda_{SC}$ is the scale where the supercolor interaction becomes strong. Scalar masses of the order of the weak scale correspond to $\Lambda_{SC} \sim 10^{11}$ GeV. The gluino and photino masses, however, are produced by higher-dimensional operators, so they are suppressed by at least two powers of $M_P$ and are unacceptably small [1, 8].

The situation is improved if there are light color-octet chiral fields which can mix with the gluinos [7]. However, such mixing is associated with a large breaking of hypercharge unless hypercharge is unified in a non-abelian group broken at an intermediate scale [8]. Also, additional fields are required to enhance the photino mass. Thus, this alternative is
very complicated.

Another possibility is that the hidden sector contains a singlet whose $F$-term acquires a VEV [9]. In this case, the gluino mass is suppressed by only one power of $M_P$. However, the supercolor sector has chiral content, so that the $F$-term for the singlet can only be induced at one or two loops. As a result, the gluino mass is still smaller than the weak scale by at least one order of magnitude.

In visible sector models, SUSY breaking is communicated to the MSSM by gauge interactions. Affleck, Dine and Seiberg [1] constructed the simplest scheme of this type by gauging a global symmetry of the supercolor sector and identifying it with one of the gauge symmetries of the standard model. However, they found that in order to give a sufficiently large mass to the gluino the gauged symmetry should be $SU(3)_C$, such that color becomes strong above the SUSY breaking scale. Dine and Nelson [10] proposed a more sophisticated scheme which allows a large gluino mass without major phenomenological problems. This involves a new gauge interaction (the “messenger group”) which communicates SUSY breaking to some additional fields. These have interactions with a gauge singlet which obtains a VEV. Finally, some vector-like quarks and leptons have a non-supersymmetric spectrum due to the coupling to the gauge singlet. Subsequent versions of this approach were greatly simplified, firstly [11] by taking advantage of the observation that the $R$-axion is massive and secondly [2] by avoiding Fayet-Iliopoulos terms for the messenger $U(1)$. The result is a highly predictive model with all the soft breaking parameters of the MSSM determined in terms of only two unknown parameters. However, a large number of fields is still required and the problem of electroweak symmetry breaking is not settled.

The above approaches show that it is possible to construct realistic models of dynamical SUSY breaking but one has to introduce very complicated structures for avoiding light gluinos. Hence, it would be useful to find simpler methods for gluino mass generation.

In this paper we present a new mechanism for producing a gluino mass. The idea is to introduce a new gauge interaction and to arrange that a gluino mass arises dynamically when this interaction becomes strong. This has similar features to the fermion mass generation in technicolor models [12]. This mechanism for gluino mass generation is quite general and can be useful in hidden or visible sector models of dynamical SUSY breaking, or in models with a continuous $R$-symmetry.

In section 2 we discuss the mechanism in general, without considering the origin of SUSY breaking. Sections 3 and 4 specialize this mechanism to hidden sector models and
visible sector models, respectively. A summary of the results and some final remarks are given in section 5.

2 Gluino Mass

Consider the MSSM with massless gluinos at tree level. The radiative gluino masses are small, of order 1 GeV [13]. Let us introduce a new non-abelian gauge interaction ("G-color") which becomes strong at a scale $\Lambda_G$ larger than the weak scale. For simplicity we take this to be $SU(N_G)$. We also introduce two chiral superfields, $\psi_L$ and $\psi^c$, with $SU(N_G) \times SU(3)_C \times SU(2)_W \times U(1)_Y$ quantum numbers:

$$\psi_L : (N_G,3,1)_y \quad , \quad \psi^c : (\overline{N_G},\overline{3},1)_{-y}.$$  \hspace{1cm} (2.1)

The asymptotic freedom of QCD is preserved if $N_G = 2$; $N_G = 3$ or 4 is also acceptable, since the QCD coupling constant remains weak below the Planck scale ($\alpha_s(M_P) \lesssim 0.25$ for $N_G = 4$).

The notations we use for a superfield and for its upper-spin component are the same while the lower spin component is distinguished by a tilde. The only exception will be that we use the standard notation for Higgs fields (i.e. the tilde is on higgsinos).

We assume that the $G$-colored scalars have positive squared masses, $M_{\tilde{\psi}_L}^2$ and $M_{\tilde{\psi}_c}^2$, and that there is a dynamically generated mixing

$$M_{LR}^2 \tilde{\psi}_L \tilde{\psi}_c + h.c.$$ \hspace{1cm} (2.2)

Such soft SUSY breaking terms arise in models of dynamical SUSY breaking under certain conditions which are discussed in the following two sections.

If the scalar masses are larger than $\Lambda_G$, the low energy $G$-color theory consists of $G$-gluinos and $N_f = 3$ massless flavors of $G$-fermions in the fundamental representation. In the absence of scalar mixing, this theory has a non-anomalous $U(1)_R$ symmetry [14] in addition to the usual chiral symmetry $SU(3)_L \times SU(3)_R \times U(1)_Y$. Under this $R$-symmetry, the scalars $\tilde{\psi}_L$ and $\tilde{\psi}_c$ have charge $1 - N_G/N_f$ (in the normalization were the Grassmann variable $\theta$ has charge $-1$), so that their mass mixing violates $U(1)_R$. This is important in what follows because the gauginos have $R$-charge $+1$ and, thus, any gaugino mass violates $U(1)_R$ by $+2$ units.

The main effect of scalar mixing is given by a dimension-six operator involving two
$G$-gluinos, $\tilde{G}$, which arises due to the scalar exchange diagram shown in fig. 1:

$$\frac{2g_G^2}{M_{\tilde{G}}} \langle \tilde{G}^a \psi_L \rangle T^a T^b (\psi^c \tilde{G}^b) ,$$  
(2.3)

where $T^a$ are the generators of the fundamental representation, $g_G$ is the $G$-color gauge coupling and

$$M_{\tilde{G}}^2 = \frac{M_G^2 M_{\tilde{G}}^2 - M_{LR}^4}{M_{LR}^2} .$$  
(2.4)

When the $G$-interaction becomes strong, the fermions condense. It is possible that a $\langle \tilde{G}G \rangle$ or $\langle \psi_L \tilde{G}\psi^c \tilde{G} \rangle$ condensate forms, as suggested in ref. [15]. However, it is more likely that a $\langle \psi_L \psi^c \rangle$ condensate forms, as in non-SUSY QCD. This is due to the operator (2.3): if the fermion bilinear has a small VEV, a $G$-gluino mass is generated and the theory tends to become similar to non-SUSY QCD. Therefore, the condensate could increase up to the point where the $G$-gluino mass is larger than $\Lambda_G$. If this is the case, the $G$-gluino can be integrated out and the theory becomes a scaled up version of ordinary QCD with three flavors, so that chiral symmetry is broken by

$$\langle \psi_L \psi^c \rangle \approx \Lambda_G^3 .$$  
(2.5)

To check the consistency of this scenario, we compute the $G$-gluino mass by Fierz transforming the operator (2.3) and using eq. (2.5):

$$M_G \approx \frac{2\pi \alpha_G \Lambda_G^3}{N_G M_{\tilde{G}}^2} ,$$  
(2.6)

where $\alpha_G = g_G^2/(4\pi) \sim O(1)$. Note that the $R$-charge of the $G$-fermions is $-N_G/N_f$ such that the $R$-charge of the $G$-gluino mass term,

$$R(\tilde{G}G) = R(\langle \psi_L \psi^c \rangle) + R(\bar{\psi_L} \tilde{\psi}^c) ,$$  
(2.7)
is indeed +2. Eq. (2.6) shows that it is not possible without fine-tuning to have both the G-gluinos and the G-scalars much heavier than the G-color scale. We consider the case where \( M_{\psi L}, M_{\psi c}, M_{LR} \sim O(\Lambda_G) \) so that \( M_G \sim O(\Lambda_G) \). In this case, the low energy G-color dynamics might be influenced by the G-scalars and G-gluinos. However, this effect is probably not large enough to change the structure of the vacuum and we will neglect it in the rough estimate of the gluino mass.

The reversed argument, that a G-gluino condensation in connection to the operator (2.3) could lead to a large mass for the G-fermions, does not apply because the chirality-flip scalar mixing requires the G-fermion condensate. Still, it is possible that both the G-fermion condensate and the G-gluino condensate appear. This would induce a scalar mixing of order \( \alpha_G \Lambda_G^2 \), as shown in fig. 2.

![Diagram](image)

Fig. 2. Scalar mixing due to the \( \langle \bar{G}G \rangle \) and \( \langle \psi_L \psi^c \rangle \) condensates (represented by ●).

Whether or not this happens, the discussions that follow will not be affected since they rely only on the assumption that the G-fermions condense.

Now we have the tools for producing a gluino mass. The scalar exchange diagram of fig. 1 with the external G-gluinos replaced by ordinary gluinos yields an effective four-fermion interaction, which gives a gluino mass when the G-fermions condense:

\[
M_\tilde{g} \approx \frac{2\pi \alpha_s \Lambda_G^2}{3M_\psi^2},
\]

where \( \alpha_s \approx 0.1 \) is the strong coupling constant at the scale \( \Lambda_G \). A gluino mass of few hundred GeV can be easily obtained with \( \Lambda_G \sim O(1 \text{ TeV}) \). Note that the experimental lower bound for the gluino mass is model dependent and in the range \( 100 - 220 \text{ GeV} \) [6].

Similarly, the diagram of fig. 1, with photinos instead of G-gluinos, gives a Majorana mass for the photino:

\[
M_\tilde{\gamma} = \frac{2\pi \alpha_s \Lambda_G^2}{\cos^2 \theta_W M_\psi^2}.
\]

5
where $\alpha$ is the fine structure constant and $\theta_W$ is the weak angle. The condition of having the Landau pole of $U(1)_Y$ above $M_P$ imposes a bound on the hypercharge of $\psi_L$,

$$|y| \lesssim 1 \quad (2.10)$$

(we use the convention $Y/(T_3 - Q) \equiv 2$, where $Y$ is the hypercharge, $Q$ is the electric charge and $T_3$ is an $SU(2)_W$ generator). Thus, there is an upper bound on the ratio of the Majorana masses of the photino and gluino:

$$\frac{M_{\tilde{\gamma}}}{M_{\tilde{g}}} = \frac{3\alpha y^2}{2\alpha_s \cos^2 \theta_W} \lesssim 0.15 \quad (2.11)$$

3 Hidden Sector Models

If supersymmetry is broken dynamically in a hidden sector by supercolor interactions at a scale $\Lambda_{SC} \sim 10^{11}$ GeV, all the scalars get masses of order 1 TeV at this scale, while the gaugino masses are suppressed at least by an additional factor of $\Lambda_{SC}/M_P$.

In order to use the mechanism for gluino mass generation presented in the previous section, there is need for scalar mixing. Such a mixing could be produced by supergravitational interactions if the Kahler potential were non-minimal. Nevertheless, this situation is not likely to occur since the constraints from flavor-changing neutral currents suggest that the part of the Kahler potential responsible for squark and slepton masses is minimal.

In the rest of this section we present a more natural source of $G$-scalar mixing.

3.1 $G$-scalar mixing

We consider a gauge singlet, $S$, in the visible sector with interactions described by the following superpotential:

$$W = \lambda_1 S \psi_L \psi^c + \frac{\lambda_2}{3} S^3 \quad (3.1)$$

When the $G$-fermions condense, the Yukawa interaction of the $\tilde{S}$ scalar with the $G$-fermions gives rise to a tadpole term for the gauge singlet scalar. The potential for $\tilde{S}$ is given by:

$$V_{\tilde{S}} = \lambda_2^2 |\tilde{S}|^4 + M_{\tilde{S}}^2 |\tilde{S}|^2 + \lambda_1 A_{\Lambda_{SC}}(\tilde{S} + \tilde{S}^\dagger) \quad (3.2)$$

The scalar $\tilde{S}$ has positive mass squared at the scale $\Lambda_{SC}$, $M_{\tilde{S}}^2 \sim O((1 \text{ TeV})^2)$, but one loop corrections can drive $M_{\tilde{S}}^2$ negative at the scale $\Lambda_G$ if $\lambda_1$ is large enough. The minimum of
the potential is at
\[ \langle \tilde{S} \rangle = - \left( \frac{\lambda_1}{4\lambda_2} \right)^{1/3} \Lambda_G f(a), \]
where
\[ a = \frac{2^{1/3} M_{\tilde{S}}^2}{3(\lambda_1 \lambda_2)^{2/3} \Lambda_G^2}, \]
and
\[ f(a) = \begin{cases} 
[(1 + a^2)^{1/2} + 1]^{1/3} - a [(1 + a^2)^{1/2} + 1]^{-1/3}, & a \geq -1 \\
2 |a|^{1/2} \cos \left( \frac{1}{3} \arccos |a|^{-3/2} \right), & a < -1 \end{cases} \]

For Yukawa coupling constants \( \lambda_1, \lambda_2 \sim \mathcal{O}(1) \), eqs. (3.3)-(3.5) give a negative VEV for \( \tilde{S} \) of order \( \Lambda_G \), which is quite insensitive to the sign or value of \( M_{\tilde{S}}^2 \) as long as the ratio \( |M_{\tilde{S}}^2|/\Lambda_G^2 \) is not very large. Due to this VEV, the Yukawa interaction of the \( \tilde{S} \) scalar with \( S \) fermions induces a Majorana mass for \( S \):
\[ m_S = \lambda_2 |\langle \tilde{S} \rangle| \sim \mathcal{O}(\Lambda_G). \]

The \( S \) fermion exchange diagram shown in fig. 3 yields a dimension-five operator,
\[ \frac{\lambda_1^2}{m_S} \tilde{\psi}_L \psi^c \psi_L \psi^c, \]
which gives a scalar mixing when the fermions condense:
\[ M_{LR}^2 \approx \frac{\lambda_1^2 \Lambda_G^3}{m_S}. \]

\[ \begin{array}{c}
\psi^c \\
\tilde{\psi}_L \\
S \\
\psi_L \\
\tilde{\psi}^c
\end{array} \]

Fig. 3. Fermion exchange diagram responsible for scalar mixing.

Therefore, a Majorana gluino mass is generated as discussed in section 2 (see eqs. (2.4) and (2.8)). Note that this is the gluino mass at a scale of order 1 TeV. At a scale of 200 GeV, the renormalization group evolution gives the gluino mass larger by a factor of
about 1.2. If all the mass parameters are of order 1 TeV and the Yukawa couplings are of order one, then the gluino mass is of order 200 GeV, which is close to the experimental lower bound.

By contrast, if the Kahler potential is minimal, we expect the squark and slepton masses to be in the $0.5 - 1$ TeV range. Furthermore, since the gaugino masses are produced at the low scale $\Lambda_G$, the renormalization group equations are different than in the MSSM where it is assumed that the gluino and photino masses are produced at a high scale $\sim \Lambda_{SC}$. In the MSSM the squarks are typically heavier than the sleptons because there is a positive contribution at one-loop proportional to the gluino mass squared. In the present context, this contribution appears only below $\Lambda_G$ and is negligible. Therefore, the squark and slepton masses will be almost equal. Only the top scalars will be lighter due to negative one-loop contributions proportional to the square of the large top Yukawa coupling.

### 3.2 Linking the $G$-color scale with the $G$-scalar masses.

There is an artificial\(^2\) ingredient in the above discussion: the scale associated with the $G$-scalar masses and the scale of the $G$-color interaction have the same order of magnitude. Both these scales arise naturally due to nonperturbative dynamics, but the fact that they are close to each other requires an explanation. Note that a similar problem \cite{16} appears in supersymmetric technicolor models \cite{17}.

Since $M_{\tilde{\psi}_L}$ and $M_{\tilde{\psi}_c}$ have the same origin as the squark and slepton masses, we do not expect them to be larger than about 1 TeV. If $\Lambda_G$ were significantly smaller than $M_{\tilde{\psi}_L}$ (we assume $M_{\tilde{\psi}_L} = M_{\tilde{\psi}_c}$), then the gluino mass would be too small, as can be seen from eq. (2.8). If $\Lambda_G$ were much larger than $M_{\tilde{\psi}_L}$, $G$-color would probably be spontaneously broken by a $G$-scalar VEV when $N_G \geq 3$ or the chiral symmetry would be unbroken when $N_G = 2$ \cite{15}.

The equality of these two scales can be seen as a fine-tuning of the $G$-color coupling constant. Note that this is not a fine-tuning in the technical sense (it is natural in the sense of 't Hooft \cite{18} to adjust a gauge coupling), but rather in the colloquial sense, i.e. the range allowed for the $G$-color coupling constant is small. The renormalization group evolution of the $G$-color coupling constant between $\Lambda_G$, where $\alpha_G(\Lambda_G) \approx 1$, and $M_{\tilde{\psi}_L}$ gives the range allowed for $\alpha_G(M_{\tilde{\psi}_L})$. The one-loop coefficient of the $\beta$-function

\(^2\)I thank Lisa Randall for emphasizing this point.
is $b_0 = -3N_G + 2$ since the $G$-scalars decouple below $M_{\psi_L}$. For example, if we require $1 \leq M_{\psi_L}/\Lambda_G \leq 2$, then $0.56 \lesssim \alpha_G(M_{\psi_L}) \lesssim 1$ for $N_G = 3$. Although this does not seem to be a severe constraint, the range of $\alpha_G$ shrinks at higher scales. Using $b_0 = -3N_G + 3$ with $N_G = 3$ between $M_{\psi_L}$ and the supercolor scale gives a 4% fine-tuning in $\alpha_G(\Lambda_{SC})$ (we define the amount of fine-tuning by $\alpha_G(\Lambda_{SC})_{\text{max}}/\alpha_G(\Lambda_{SC})_{\text{min}} - 1$).

Such a fine-tuning might be explained if supercolor and $G$-color are unified at a high scale. This would provide a connection between the supercolor coupling constant, $\alpha_{SC}(\Lambda_{SC}) \approx 1$, and $\alpha_G(\Lambda_{SC})$.

A more convenient approach is to avoid excessive fine-tuning by slowing down the running of $\alpha_G$ above $M_{\psi_L}$. This requires additional $G$-colored chiral superfields. As an example, consider the case $N_G = 3$. We introduce a new chiral superfield in the adjoint representation of the $G$-color group. Its scalar component receives a mass of order $M_{\psi_L}$ from supergravitational interactions while the fermion component condenses when $G$-color becomes strong (we assume that this does not prevent the formation of the $(\psi_L \psi^c)$ condensate). The $\beta$-function above $M_{\psi_L}$ is given by [19]:

$$\beta(g_G) \approx \frac{g_G^3}{16\pi^2} \left(-3 + 2.7\alpha_G + \mathcal{O}(\alpha_G^2)\right).$$

We see that the asymptotic freedom is preserved while the rate of running for $\alpha_G \sim \mathcal{O}(1)$ is very small. However, previous studies of gauge theories with a slowly-running coupling constant in the context of walking technicolor [20] show that the convergence of the $\beta$-function is not reliable in this case. Therefore, it is not very useful to compute the amount of fine-tuning at the scale $\Lambda_{SC}$ using the two-loop $\beta$-function. We expect, though, that the range of $\alpha_G$ will not shrink dramatically above $M_{\psi_L}$, while the range of $\alpha_G(M_{\psi_L})$ will not be reduced significantly with the inclusion of the additional superfield.

In conclusion, if the $G$-color coupling constant evolves slowly, then the two scales, $M_{\psi_L}$ and $\Lambda_G$, are linked [16]: when the $G$-scalars decouple, $\alpha_G$ starts to increase faster, triggering chiral symmetry breaking.

### 3.3 Electroweak symmetry breaking

Since we try to construct a viable way of transmitting dynamical SUSY breaking, we have to specify the mechanism for electroweak symmetry breaking. At the scale $\Lambda_{SC}$ supergravitational interactions give positive squared masses to the Higgs scalars, $H_u$ and $H_d$, of order $M_{\psi_L}$. Due to the large Yukawa coupling of the top quark, radiative corrections
drive \( M_{H_u}^2 \) negative at the scale \( \mu_{\psi_L} \) and the electroweak symmetry breaks [21]. Though, additional features are necessary: \( H_d \) should also have a VEV and the Higgsinos should be massive. In the MSSM, these requirements are satisfied by a \( \mu H_u H_d \) term in the superpotential. In the present context it is not reasonable to introduce a \( \mu \)-term since the main goal of dynamical SUSY breaking is to explain the existence of any mass parameter other than \( M_P \).

However, we can introduce a term

\[
W' = \lambda_3 S H_u H_d
\]

(3.10)
in the superpotential. The potential for \( \tilde{S} \) is modified in this case and we have to minimize it simultaneously with respect to \( \tilde{S}, \tilde{H}_u \) and \( \tilde{H}_d \). However, for simplicity we consider the case were the Yukawa coupling constant is small, \( \lambda_3 \sim 0.2 - 0.3 \), so that eqs. (3.3)-(3.5) remain valid. The VEV of \( \tilde{S} \) then gives a higgsino mass term,

\[
-\lambda_3 \langle \tilde{S} \rangle H_u \tilde{H}_d,
\]

(3.11)
and the VEV of the \( F_S \) auxiliary field gives a scalar mixing \( B \)-term,

\[
-2\lambda_2 \lambda_3 \langle \tilde{S} \rangle^2 H_u H_d.
\]

(3.12)

If \( \langle \tilde{S} \rangle \sim \mathcal{O}(1\text{TeV}) \) and \( \lambda_2 \sim \mathcal{O}(1) \), then the higgsino mass is in the range of few hundred GeV, while the mass coefficient of the \( B \)-term is larger by a factor of 2 or 3. Therefore, the gauge singlet can perform the tasks of a \( \mu \)-term.

Since the Majorana wino mass is zero at tree level, the lightest chargino could have the mass above the experimental lower limit of \( \sim 45 \text{ GeV} \) only if \( \tan \beta = \frac{\mu_{\psi_L}}{y} \) is small (\( \tan \beta \lesssim 3 \)) [22]. The lightest neutralino has a mass of order \( M_{\tilde{\chi}} \), which is \( \sim 30 \text{ GeV} \) when \( M_{\tilde{g}} \sim 200 \text{ GeV} \) and \( y \approx 1 \) (see eq. (2.11)); this is also close to the experimental lower limit of 20 GeV [6]. It is possible, however, to give Majorana masses of order \( \sim 100 \text{ GeV} \) to the winos and zinos if there are two flavors of \( G \)-colored superfields which form \( SU(2)_W \) doublets. Note that \( \alpha_G \) runs slowly if we introduce four (six) new flavors when \( N_G = 3 \) (\( N_G = 4 \)). In this case, the masses of the lightest neutralino and chargino will increase.

We have to address a problem associated with gauge singlets. A quadratically divergent tadpole can appear at two loops in supergravity theories which contain a singlet under any gauge and global symmetry [23]. This tadpole leads to a large VEV for the singlet which destabilizes the mass hierarchy if the singlet has renormalizable couplings
to the visible sector. Thus, in the present context this problem is avoided only if $S$ is charged under a global symmetry. The term $S^3$ from the superpotential breaks any such continuous global symmetry (excepting an $R$-symmetry under which $S$ has charge $\frac{2}{3}$; this is broken by the $S\psi E\psi^c$ term). Still, the superpotential $W + W'$ has a discrete $Z_3$ symmetry. Since $S$ has non-zero $Z_3$ charge, the Higgs fields are also charged so that the quarks and leptons have to carry $Z_3$ charge. This symmetry can allow the most general couplings of the quarks and leptons to the Higgs fields, but also it can be used as a horizontal symmetry which restricts these couplings (this would make a connection between fermion masses and supergravity).

$Z_3$ is dynamically broken by the $G$-fermion condensate which implies that domain walls were produced when the Universe had a temperature of order $\Lambda_G$ [24]. As emphasized in refs. [2, 10], this may not be a phenomenological problem since there are several possibilities that lead to a sufficiently fast decay of the domain walls: for example the $Z_3$ could be broken by dimension-five operators or this discrete symmetry can be a remnant of a spontaneously broken continuous symmetry [24, 25].

4 Visible Sector Models

Visible sector models are appealing for several reasons: since SUSY breaking is transmitted by gauge interactions, squark and slepton degeneracy arises naturally; the coefficients of the soft SUSY breaking terms appearing in the MSSM can be computed as functions of few parameters [2]; Planck scale physics is employed only in the cancellation of the cosmological constant and in the generation of the $R$-axion mass [5]. In this type of models, the lower limit on the $R$-axion mass and the upper limit on the gravitino mass require the supercolor scale to be in the range $10^2 - 10^4$ TeV [11].

It is difficult to produce a large gluino mass in this type of models because the supercolored fields cannot carry ordinary color. In refs. [10, 11, 2] this problem is solved by an indirect transmission of SUSY breaking, where the last link before the MSSM is formed by vector-like quarks and leptons.

The mechanism for gluino mass generation presented in section 2 opens up the possibility of constructing simple visible sector models along the lines of ref. [1]. In the rest of this section we present some difficulties with this approach.

Consider a supercolor model with an unbroken global $U(1)$ symmetry which is gauged and identified with hypercharge. As a result, all the scalars of the MSSM and the $G$-
scalars get mass at two loops proportional to their hypercharge [1]. Thus, the right-handed charged sleptons are the heaviest scalars while the squark doublets are the lightest, the mass ratio between these being six. This might be a problem: depending on the experimental lower bound for squarks (which is in the range 90 – 220 GeV), the sleptons could be too heavy to insure the weak scale stability (a recent study of the fine-tuning as a function of the superpartner masses can be found in ref. [26]).

Since supercolored fields carry hypercharge, the photino Majorana mass is easily produced. The main problem encountered in ref. [1] is that gluinos get only a very small mass at three-loop level from the supercolor sector. Here, however, a large gluino mass is produced by the $G$-color sector supplemented by a gauge singlet superfield, as in section 3. Note that in addition to the $S$ fermion exchange diagram, there is a contribution to $G$-scalar mixing from a photino exchange. The supercolored fields do not contribute to the mass of the scalar singlet but this does not prevent $\tilde{S}$ from acquiring a VEV such that the electroweak breaking occurs as in the hidden sector models.

The upper bound on the hypercharge of the $G$-colored superfields (see eq (2.10)) implies that $M_{\tilde{\psi}_L}$ is at most half of the right-handed charged slepton mass, which might be too low for allowing a sufficiently large gluino mass. If this is the case, $M_{\tilde{\psi}_L}$ can be enhanced by gauging a second global symmetry of the supercolor sector and assigning its charges to the $G$-scalars. Moreover, if the $G$-scalars become sufficiently heavy, their two-loop contributions to the squark masses may reduce the slepton to squark mass ratios.

The above considerations are based on the assumption that the two-loop contributions of the supercolor sector give positive squared squark and slepton masses. The scalars (fermions) from the supercolor sector give negative (positive) contributions [11, 10] so that at least some of the fermions carrying hypercharge should be heavier than some of the charged scalars. This turns out to be a strong condition for supercolor models. For example, the simplest supercolor sector, the $SU(3) \times SU(2)$ model [1, 5], is ruled out from this scheme since the fermion component of the only light composite superfield charged under $U(1)_Y$ is massless.

Also, the supercolor model should not allow a large Fayet-Iliopoulos term for hypercharge at one loop since this would produce squark VEV’s [11, 27]. The $SU(6) \times U(1)$ model of ref. [2] is designed to cancel the Fayet-Iliopoulos term. Again, in this model the charged fermions are lighter than the scalars.

Yet another constraint is that all the charged fields should be massive. The $SU(5)$ model with two chiral superfields in the $\bar{5}$ representation and two in the 10 [3, 28], or the
SU(3) \times SU(2) model do not satisfy this condition.

However, it is quite possible that supercolor models with the desired properties exist. Only few of the known models of dynamical SUSY breaking were studied in detail and, probably, new models will be find soon. If a supercolor sector satisfying the above constraints is found, a visible sector model including the $G$-color sector would have certain advantages over the models of refs. [2, 11]: SUSY breaking is transmitted directly to the MSSM without need of additional fields beyond the $G$-color sector and the singlet; the supercolor scale is lower by an order of magnitude so that the gravitino is lighter, satisfying easier the cosmological bounds [29]; the role of the $\mu$-term is played successfully by the singlet. Alternately, if the mechanisms for generating a $\mu$-term discussed in refs. [2, 11] remain problematic, the models with indirect transmission of SUSY breaking may need a $G$-color sector to give a weak scale VEV to a scalar singlet.

5 Conclusions

The complications in constructing realistic models of dynamical SUSY breaking, mainly with producing a gluino mass and a $\mu$-term, suggest a need for additional dynamics, beyond the one required to break SUSY.

We have shown that a sufficiently large gluino mass can be dynamically generated if there are chiral superfields carrying both color and the charges of an additional non-abelian gauge interaction, $G$-color. When the $G$-colored scalars decouple, the $G$-color coupling constant starts running faster and the $G$-fermions condense at a scale of order 1 TeV. However, in order to generate a gluino mass without fine-tuning, this link between the $G$-color scale and the SUSY breaking scale should be very effective, which happens only if the $G$-color $\beta$-function nearly vanishes at scales above the scalar masses. A nice feature of this mechanism is that the absence of the $\mu$-term can be compensated naturally by the interactions of a gauge singlet.

Hidden sector models appear to become viable by including this mechanism. The only role of the supercolor sector in this case is to give masses to the scalars in the visible sector and, therefore, any model of dynamical SUSY breaking is adequate here.

By contrast, simple visible sector models could be constructed using this mechanism provided one finds a supercolor model which fulfills the following conditions: i) some of the supercolored fields carry hypercharge; ii) the non-supersymmetric vacuum of the supercolor sector preserves hypercharge; iii) in the effective low energy theory the charged
fermions are heavier than the charged scalars; \( iv \) there is no large Fayet-Iliopoulos term for hypercharge; \( v \) there is no massless charged field.

A more general issue which is relevant for this approach is the behaviour of SUSY gauge theories with a scale of the order of the soft SUSY breaking terms. This theories are in between the better understood gauge theories with small soft SUSY breaking terms and the non-SUSY gauge theories, and deserve more studies.

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