I present a realistic model of dynamical supersymmetry breaking, in which a $U(1)_{B-L}$ gauge interaction communicates supersymmetry breaking to the standard fields. A distinctive superpartner spectrum is predicted in this model.

1. Motivation

The large hierarchy between the electroweak scale, $v$, and the Planck scale indicates the existence of a symmetry which protects $v$ against quadratic divergences, and is dynamically broken. If there are light fundamental scalars, such as the Higgs doublet, then the “protective” symmetry must be supersymmetry.

If the superpartners of the standard model fields find out about supersymmetry breaking from the standard gauge interactions, then the superpartner spectrum can be computed in terms of few parameters, and unwanted flavor-changing neutral currents are suppressed. This standard gauge mediated supersymmetry breaking (SGMSB) scenario requires “messenger” superfields charged under the standard gauge group, with a nonsupersymmetric spectrum. In the usual SGMSB models [1], where the vacuum expectation values (VEV’s) of the scalar- and $F$- components of a gauge singlet superfield give the masses of the messenger fields, only false vacua have experimentally viable properties [2] [3]. The true vacuum may become viable if a second gauge singlet is included [4] [5], but in this case there are many independent parameters. If mass terms for the messenger fields are included in the superpotential, then simple SGMSB models can be constructed [6]. It remains, though, to be shown that this simplicity is preserved once a dynamical origin for these masses is specified. Other recent SGMSB models involve non-generic and nonrenormalizable interactions [7].

A different possibility is that the messenger of supersymmetry breaking is a spontaneously broken gauge interaction. The simplest choice is a $U(1)_{B-L}$ that couples to the $B-L$ number. The possibility of using $U(1)_{B-L}$ as messenger was first suggested in ref. [6], but model building efforts in this direction have been hampered by several phenomenological problems: i) it is difficult to give rise to positive squared-masses for squarks and sleptons; ii) the usual gauginos do not couple to $B-L$ so that they remain too light; iii) a natural mechanism of breaking $U(1)_{B-L}$ spontaneously should be found. A model which uses a combination of $U(1)_{B-L}$ and hypercharge as messenger is presented in ref. [7]. However, this is not a model of dynamical supersymmetry breaking (DSB) because supersymmetry breaking is introduced through Fayet-Iliopoulos terms.

Here I construct a renormalizable DSB model with $U(1)_{B-L}$ as messenger, which solves the phenomenological problems listed above, and as a consequence predicts a peculiar superpartner spectrum.

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2. The Model

In addition to the minimal supersymmetric standard model (MSSM) fields, the model I propose contains an $SU(5) \times SU(2) \times U(1)_{B-L}$ gauge group and the chiral superfields shown in table 1, which are singlets under the standard model gauge group.

|      | $SU(5)$ | $SU(2)$ | $U(1)_{B-L}$ |
|------|---------|---------|-------------|
| $A_1, A_2$ | [ ]    | 1       | 0           |
| $B$    | [ ]    | [ ]    | 0           |
| $\phi$ | 1      | [ ]    | 0           |
| $\chi_{0,\pm}$ | 1      | [ ]    | 0, $\pm y$ |

Table 1: Field content of the model.

The most general dimension-3 terms in the superpotential are given by

$$ W = \lambda_0 A_1 B^2 + \lambda_0^2 \phi + \kappa \chi + \chi_\phi \phi . \quad (2.1) $$

To avoid a hierarchy problem, the mass terms $\phi^2$ and $\chi_\phi \chi_\phi$ are excluded by invoking a discrete symmetry.

The $SU(2)$ group is infrared free, so that it must be in the weak coupling regime. The $SU(5)$ instantons generate the following effective superpotential $[8]$:

$$ W_{\text{np}} = \frac{\Lambda_5^{11}}{(A_1 A_2)^4 B^2} , \quad (2.2) $$

where $\Lambda_5$ is the scale of $SU(5)$.

The scalar potential can be written as

$$ V = V_5 + V_2 , \quad (2.3) $$

where

$$ V_5 = \frac{g_0^2}{2} D_5^a D_5^\dagger + |F_A|^2 + |F_B|^2 \quad (2.4) $$

is the potential of the $SU(5)$ DSB model with two generations $[8]$, with $D_5$ the $D$-term of $SU(5)$, and $F_A, F_B$ the usual $F$-terms of the $A$ and $B$ fields, while

$$ V_2 = \frac{g^2}{2} D_2^a D_2^\dagger + |F_\phi|^2 + |F_\chi|^2 \quad (2.5) $$

includes the $SU(2)$ $D$-potential and the remaining $|F|^2$ terms.

The most general parametrization of the $B$ scalar field, up to an $SU(5) \times SU(2)$ transformation, is given by

$$ B = \begin{pmatrix} b & 0 & 0 & 0 & 0 \\ 0 & b' & 0 & 0 & 0 \end{pmatrix} . \quad (2.6) $$

At the minimum at least one of $b$ and $b'$ is nonzero (otherwise $F_A \to \infty$), so that $SU(2)$ is completely broken.

At scales where the $SU(5)$ gauge coupling $g_0$ is much larger than $\lambda_0$, the global minimum of $V_5$ lies along the $D_5$ flat directions. In this case, i.e. to leading order in $\lambda_0^2/g_0^2$, it has been shown numerically $[9]$ that a flavor-$SU(2)$ symmetry of the $SU(5)$ DSB model is preserved, which requires $|b| = |b'| > 0$ at the minimum. Thus, the VEV of $B$ does not contribute to the $SU(2)$ $D$-term to leading order in $\lambda_0^2/g_0^2$. As a result, the deepest minimum of $V$ corresponds to the deepest minima of both $V_5$ and $V_2$.

Because $V_2$ has many classically flat directions, the vacuum energy is given entirely by $V_5$. At the minimum $b \sim \lambda_0^{-1/11} \Lambda_5$ and the vacuum energy is of order $\lambda_0^{1/2} b$, so that supersymmetry is spontaneously broken. Therefore, the flat directions of $V_2$ are lifted by radiative corrections.

The $\phi$ and $\chi$ scalars get masses at one-loop. In the Landau gauge, the only relevant diagrams are the ones shown in fig.1. The computation of these diagrams is difficult because there is mixing between the $A_1, A_2$ and $B$ states. It is possible, however, to estimate the one-loop masses of the $\phi$ and $\chi$ scalars by adapting to the present $SU(2)$ sector the Feynman rules given in ref. $[10]$ for the Higgs sector of the MSSM. This amounts to compute only the contributions from the $SU(2)$ doublets which acquire VEV’s, i.e. $B_1$ and $B_2$, where 1 and 2 are $SU(5)$
indices [see eq. (2.6)]. The result is finite and negative: the one-loop squared-mass of a scalar in the $R$ representation of $SU(2)$ is given by

$$M_R^2 = - C_R M^2 , \quad (2.7)$$

where $M^2 > 0$ and $C_R$ is the eigenvalue of the Casimir operator (3/4 for the doublet). $M^2$ has a simple expression if expanded in powers of $\lambda_0^2 / g^2$. The leading term, proportional to $g^4$, cancels because supersymmetry is exact in the $\lambda_0 \to 0$ limit. The next term in the expansion is positive:

$$M^2 = \frac{g^2}{8 \pi^2} \left( \mu_0^2 + m_0^2 \right) + \mathcal{O} (\lambda_0^4) . \quad (2.8)$$

Here $\mu_0 > 0$ is the analogue of the $\mu$-term from the Higgs sector, i.e. the coefficient of the $B_1 B_2$ fermion mass term ($\mu_0 \sim \lambda_0 b$), and $m_0^2 > 0$ is the coefficient of the $B_1 B_2$ scalar mass term ($m_0 \sim \lambda_0 b$). The mixing between the $A_1$, $A_2$ and $B$ states will change the numerical coefficient in eq. (2.8), but it appears reasonable to assume that the sign of $M^2$ will not change.

With the $\phi$ and $\chi$ negative mass terms included, the scalar potential

$$V' = V_2 - M^2 \left[ 2 |\phi|^2 + \frac{3}{4} \left( |\chi_0|^2 + |\chi_\pm|^2 \right) \right] \quad (2.9)$$

has a runaway direction: $V' \to - \infty$ for $\chi_{0,\pm} = 0$ and $|\phi| \to \infty$. However, the full scalar potential $V$ is positive definite, which implies that the runaway direction is lifted by higher dimensional terms. For example, one-loop diagrams similar to the ones in fig. 1 but with four external legs induce a $|\phi|^4$ term in the effective potential, with a coefficient of order $g^2 \lambda_0^2 / (4 \pi)^2$. There is also an infrared divergent contribution to the $|\phi|^4$ term which should be eliminated by summing up the complete one-loop effective potential, leading to a $\phi^2 \log \phi$ term. The balance between the $|\phi|^2$ term and the higher dimensional terms gives the global minimum at

$$\chi_{0,\pm} = 0 , \quad |\phi| \sim b . \quad (2.10)$$

The soft supersymmetry breaking terms generated in the MSSM (see the following sections) are only logarithmically sensitive to the value of $|\phi|$ because this VEV gives supersymmetric contributions to the $\chi_\pm$ masses. Note that the VEV of $\phi$ breaks a global $U(1)$ and the resulting Goldstone boson, which has anomalous couplings to the $SU(2) \times U(1)_{B-L}$ gauge bosons, is likely to get a Planck scale suppressed mass, of order $|\phi|^{3/2} M_P^{-1/2}$.

3. Squark and slepton spectrum

In the vacuum (2.10) all the $\chi$ and $\phi$ fields are massive. The four scalar components of the $\chi_\pm$ superfields are degenerate, with mass

$$M_{\chi_\pm}^2 = \frac{\kappa^2}{2} |\phi|^2 - \frac{3}{4} M^2 . \quad (3.1)$$

The fermion components of $\chi_\pm$ form two degenerate Dirac fermions of mass

$$m_{\chi} = \frac{\kappa}{\sqrt{2}} |\phi| . \quad (3.2)$$

Because the spectrum of $\chi_\pm$ is nonsupersymmetric, the squarks and sleptons, as well as any other scalar charged under $U(1)_{B-L}$, get masses at two loops (the one-loop contributions from the $\chi_+$ and $\chi_-$ scalars
cancel each other). The leading logarithmic term is given by

$$M_{B-L}^2 = - \left[ (B-L) y \frac{\alpha_{B-L}}{2\pi} \right]^2 \mathrm{Str}(M^2_{\chi}) \ln \left( \frac{\Lambda}{m_{\chi}} \right),$$

(3.3)

where the cut-off is of the order of the $SU(5)$ gauge boson mass, $\Lambda \sim g_0 b$, and the supertrace over the eight $\chi_{\pm}$ states,

$$\mathrm{Str}(M^2_{\chi}) = -3M^2,$$

(3.4)

is negative because of the $SU(2)$-induced one-loop mass [see eq. (2.7)] of the $\chi_{\pm}$ scalars.

Thus, the squark and slepton squared-masses are indeed positive, which is the primary motivation for the choice of the field content in table 1. Given that the squarks have baryon number $\pm 1/3$ and the sleptons have lepton number $\mp 1$, from eq. (3.3) follows an interesting prediction for the relation between the slepton mass, $m_{\tilde{L}}$, and the squark mass, $m_{\tilde{Q}}$:

$$m_{\tilde{L}} = 3m_{\tilde{Q}}.$$

(3.5)

This prediction is in contrast with the ones from SGMSB models and from supergravity scenarios [12], where the squarks are heavier than the sleptons. Hence, if the squarks and sleptons will be observed and their masses measured, then the prediction (3.5) will be an important test of the model presented here. Another feature is that all the sleptons are degenerate (the electroweak corrections are negligible), whereas in SGMSB models the left-handed sleptons are few times heavier than the right-handed ones. As discussed in the next section, additional fields are necessary for producing gaugino masses. These fields will contribute negatively to the squark squared-masses, so that the slepton-to-squark mass ratio increases further. Note that the squark degeneracy is slightly lifted by electroweak corrections

from the additional fields, and by the stop mixing due to the top Yukawa coupling.

Another difference from the SGMSB models is that the Higgs scalars, $H_u$ and $H_d$, do not get masses at two loops because they do not carry $B-L$ charge. However, they get masses at three loops from the interactions with the squarks and sleptons. These masses can be estimated by integrating out the $\chi$ fields and computing one-loop radiative corrections in the effective theory of heavy squarks and sleptons. Because of the large top Yukawa coupling, the stop-loop gives a negative squared-mass to $H_u$ which drives electroweak symmetry breaking, like in the SGMSB case. In order to set $v \sim 246$ GeV without fine-tuning, the squark mass should be of order few hundred GeV. The slepton mass, in turn, is of order 1 TeV.

The squark mass $m_{\tilde{Q}}$ being roughly known, the value of $b$ can be computed from eqs. (3.3) and (2.8), and then one can find out the vacuum energy. Some typical values of the parameters, $y, g \sim 1$, $\alpha_{B-L} \sim 10^{-2}$, $g_0/\kappa \sim 4\pi$, yield a vacuum energy of order $10^3$ TeV, which corresponds to a gravitino mass of few hundred eV. In this case, the lightest standard model superpartner could decay within the detector [13], and the gravitinos may be a dark-matter component [14].

4. Gaugino masses and $U(1)_{B-L}$ breaking

The gauginos of the standard gauge group do not couple to $U(1)_{B-L}$, so that their masses arise only at two loops and are of order 1 GeV. Although such light gauginos are not conclusively ruled out [7, 15], it appears more plausible that the gaugino masses are of the order of the electroweak scale. Another problem of the model presented in section 2 is that $U(1)_{B-L}$ is unbroken. These two phenomenological
problems can be solved by extending the chiral content of the model.

Consider two chiral superfields, \( q \) and \( \tilde{q} \), belonging to some vector-like representation of the standard gauge group, and three chiral superfields, \( S_{0,\pm} \), which carry \( B - L \) charges \( 0, \pm y_S \), and are singlets under the standard gauge group. The superpotential is given by

\[
W = \eta \tilde{q} q S_0 + \frac{\xi}{3} S_0^3 + \zeta S_0 S_+ S_- \quad \text{(4.1)}
\]

The other dimension-3 terms, \( S\phi^2 \) and \( S\chi^2 \), can be eliminated by discrete symmetries. The simplest choice of vector-like representation that preserves the gauge coupling unification is \( q + \tilde{q} \in n \times (5 + \overline{5}) \) of the grand unified \( SU(5)_{SM} \) group. Gauge coupling perturbativity up to the unification scale requires \( n \leq 4 \). If these superfields carry \( B - L \) charges \( \pm y_q \), then their scalar components get real masses equal to \( 3y_q m_{\tilde{q}} \). Consequently, the scalar \( S_0 \) receives a negative squared-mass via the one-loop diagram shown in fig. 2,

\[
m_{S_0}^2 = -5n \left( \frac{3y_q m_{\tilde{q}} \eta}{2\pi} \right)^2 \ln \left( \frac{\Lambda'}{3y_q m_{\tilde{q}}} \right) \quad \text{(4.2)}
\]

where \( \Lambda' \) is a cut-off of order \( m_\chi \). As a result, \( S_0 \) acquires a VEV of order \( m_{S_0} \), while an \( F_{S_0} \)-term of order \( m_{S_0}^2 \) is induced. Therefore, the \( q \) and \( \tilde{q} \) scalars end up with diagonal squared-masses equal to \( (3y_q m_{\tilde{q}})^2 + \eta^2 S_0^2 \), and off-diagonal squared-masses equal to \( \eta F_{S_0} \), while the \( q \) and \( \tilde{q} \) fermion pairs get Dirac masses equal to \( \eta S_0 \), so that the usual gauginos receive masses at one-loop [11]:

\[
m_{\tilde{g}_i} = m_S \alpha_i \frac{\eta S_0}{4\pi} \frac{r_1 \ln r_1 - r_2 \ln r_2}{1 - r_1} \quad \text{(4.3)}
\]

where \( \alpha_i \ (i = 1, 2, 3) \) are the \( SU(3)_C \times SU(2)_W \times U(1)_Y \) coupling constants, with the usual \( SU(5)_{SM} \) normalization of the hypercharge coupling constant

\[
m_{\tilde{g}_1} = m_S \alpha \frac{\eta S_0}{4\pi} \frac{r_1 \ln r_1 - r_2 \ln r_2}{1 - r_1} \quad \text{(4.4)}
\]

Acceptable gaugino masses can be produced for a range of parameters. For example, the values

\[
m_\tilde{Q} \sim 500 \text{ GeV}, \ y_q \sim 3, \ \eta \sim 1.5, \ n = 3 \quad \text{(4.5)}
\]

which correspond to \( |m_{S_0}| \sim 10 \text{ TeV} \), yield a gluino mass of about 200 GeV. The ratios between the three Majorana gaugino masses differ from the ones given by eq. (4.3) if \( q \) and \( \tilde{q} \) belong to other representations of the standard model gauge group [16].

The \( q \) and \( \tilde{q} \) fields have also contributions at two loops to the squark and slepton squared-masses, which are given by eq. (4.3) with the coupling constants, charges and masses appropriately replaced [11]. These contributions are negative because the \( q \) and \( \tilde{q} \) scalars are heavier than their fermion partners. For example, the values (4.5) give a decrease in \( m_{\tilde{Q}} \) of about 50%, while the effect on the slepton mass is negligible.

For \( R \equiv \zeta / \xi < 1 \) and \( R(1 - R)m_S^2 > 2(3y_S m_{\tilde{Q}})^2 \), the \( S_\pm \) scalars acquire VEV’s too, breaking \( U(1)_{B-L} \) as required, at a scale of order 10 TeV. Note that the \( D \)-term for \( U(1)_{B-L} \) cancels (because \( S_+ = S_- \) at the minimum), so that there are no problems with kinetic mixing between the \( U(1)_{B-L} \) gauge boson and the hypercharge gauge boson [17].

A mechanism similar to the one described above can be used to generate the \( \mu \)- and \( B \)-terms. For this purpose there is need for three new standard model
singlets $S^0_{0, \pm}$, which carry $B-L$ charges $0, \pm y'_S$. For $y'_S < y_S$, a VEV for $S^0_0$ in the few hundred GeV range can be produced, which then gives the higgsino mass and the $B$-term via an $S^0_0 H_u H_d$ term in the superpotential. Other mechanisms for generating the $\mu$- and $B$-terms may also be used [18].

It should also be mentioned that the $U(1)_{B-L}$ anomaly cancellation requires right-handed neutrinos. The VEV of $S_\pm$ can be used to generate a Majorana mass for the right-handed neutrinos, such that small neutrino masses arise by the see-saw mechanism. One has to worry though that the position of the vacuum may be changed if an $S_\pm \nu^c \nu^c$ term is included in the superpotential. Alternately, neutrino masses may be prevented by discrete symmetries (in this case the lower bound on the $U(1)_{B-L}$ gauge boson’s mass set by primordial nucleosynthesis is approximately 2 TeV [19]).

5. Outlook

The model described here contains a rather large number of parameters: 6 Yukawa couplings and 3 gauge couplings in addition to the standard model. It is also unsatisfactory that the fields which induce squark and slepton masses cannot be used to produce directly the gaugino masses and to break $U(1)_{B-L}$. Furthermore, a separate sector should be introduced for producing the $\mu$-term. However, this model is more economical than the other known viable DSB models. For example, the simplest complete model of SGMSB with a viable true vacuum [2], contains 13 parameters in the superpotential, 3 gauge couplings, and the sector that produces the $\mu$-term, in addition to the standard model parameters. Nevertheless, it would be desirable to find a common origin for the sectors that are responsible for gaugino and scalar masses, such as a grand unified theory in which the fields from different sectors belong to the same representation.

Another unpleasant feature of the SGMSB models shared by the model proposed here is that some renormalizable terms should be eliminated from the superpotential by discrete symmetries. It is unclear whether such symmetries are not badly violated by Planck scale effects. And in case they are preserved, one has to ensure that the domain walls can decay, or that a period of late inflation is possible. Inflation may be also needed for diluting the lightest $q-\bar{q}$ state [20].

Despite these drawbacks, the model proves the possibility of $B-L$ mediated supersymmetry breaking. Its importance lies in the distinctive predictions for the superpartner spectrum, most notably being the large ratio ($\geq 3$) between the slepton and squark masses. Finally, it is worth pointing out that there are other light neutral states besides the gravitino and the Goldstone boson discussed in section 2: the DSB sector contains a massless fermion [9], and an $R$-axion [21] with a mass of order 100 MeV given by supergravity effects.

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