PHYSICAL REVIEW D **84,** 055015 (2011)

Low energy supersymmetry with baryon and lepton number gauged

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(Received 6 June 2011; published 15 September 2011)

We investigate the spontaneous breaking of the Baryon (B) and Lepton (L) number at the TeV scale in supersymmetric models. A simple extension of the minimal supersymmetric standard model where B and L are spontaneously broken local gauge symmetries is proposed. The B and L symmetry breaking scales are defined by the supersymmetry breaking scale. By gauging B and L, we understand the absence of the baryon and lepton number violating interactions of dimension four and five in the minimal supersymmetric standard model. Furthermore, we show that even though these symmetries are spontaneously broken there are no dangerous operators mediating proton decay. We discuss the main properties of the spectrum, the possible baryon number violating decays and the implications for the dark matter candidates. In this model, one can have lepton number violating signals from the decays of the right-handed neutrinos and baryon number violating signals from the decays of squarks and gauginos without conflict with the bounds coming from proton decay, $n - \bar{n}$ oscillations and dinucleon decays.

DOI: 10.1103/PhysRevD.84.055015 PACS numbers: 12.60.Jv

I. INTRODUCTION

Experimental data are consistent with baryon number (B) conservation and lepton number (L) conservation. In neutrino experiments, we have observed the violation of the individual lepton numbers $L_{e,\mu,\tau}$ but not of the total lepton number $L = L_e + L_\mu + L_\tau$. It is interesting to explore the possibility that the observed B and/or L conservation has its origin in the principle of gauge invariance and construct models where B and L are spontaneously broken gauge symmetries. To gauge B and/or L, additional fermions beyond those in the minimal standard model must be added to cancel anomalies. Solutions to the anomaly constraint equations were found in Ref. [1–3].

The authors in Ref. [2] explored models where baryon number is gauged with the anomalies canceled by adding a fourth generation of quarks and leptons. Since three generations have been observed, and we do not understand why there should be only three, we view this way of canceling anomalies as less arbitrary than the other possibilities for canceling anomalies that introduce fermions with quantum numbers unrelated to those of the observed standard model fermions. Recently, we constructed two explicit models where both B and L are spontaneously broken local gauge symmetries [3]. In these models, B and L are on the same footing and the anomalies are canceled by adding a single new fermionic generation. There is a natural suppression of flavor violation in the quark and leptonic sectors since the gauge symmetries and particle content forbid tree level flavor changing neutral currents involving the quarks or charged leptons. Also, there is a dark matter candidate that is automatically stable. In these models, the symmetry breaking scale for the $U(1)_B$ and $U(1)_L$ symmetries are not necessarily related to the weak scale, however, we explored some of their phenomenology with that assumption.

In the standard model, operators that violate baryon number (schematically $qqql$) do not occur until dimension six and experimental constraints on the nucleon decay rate imply that the mass scale that suppresses them, Λ must satisfy, $\Lambda > 10^{15}$ GeV. Hence the observed conservation of baryon number is explained if there is no new physics below this scale, i.e., a desert. However, in models where baryon number is gauged the observed conservation of baryon number can be understood, even if there is new physics at scales much lower than 10^{15} GeV, since without spontaneous symmetry breaking operators that violate B are forbidden and (depending on the charges of the fields that break baryon number) the spontaneous breaking of baryon number may not induce operators that cause observable proton decay.

Supersymmetry (SUSY), softly broken at the weak scale, solves the hierarchy problem. Today, the minimal supersymmetric extension of the standard model (MSSM) is considered one of the most appealing scenarios for physics beyond the standard model. For a review on supersymmetric models, see Ref. [4]. One of the open issues for these models is the presence of renormalizable and dimension five operators that violate baryon and lepton number. These can be forbidden by gauging a linear combination of B and L [5] and it is interesting to consider extending the work in Ref. [3] to a supersymmetric model since it can also achieve that goal.

In this letter, we investigate the simplest supersymmetric extension of one of the models in Ref. [3]. Unlike the nonsupersymmetric case, here (if there are no large Fayet

Illiopoulos D terms) the B and L symmetry breaking scales are necessarily of the order of the soft supersymmetry breaking scale. We discuss the main features of the model including the properties of the spectrum and dark matter candidates. We show that there are no dangerous operators that cause proton decay even after baryon and lepton number are spontaneously broken. This model should be interpreted as an effective theory below a scale that is at most a few orders of magnitude above the weak scale because beyond that point the Yukawa couplings of the fourth generation become strong [6]. Consequently, the evidence for a supersymmetric extension of the standard model based on the meeting of the gauge couplings is not applicable in models with a fourth generation.

Within the effective field theory approach, it is possible to consider gauge theories that are anomalous. With a cutoff that is only a few orders of magnitude above the weak scale, it is possible to do this in theories that gauge B and L [7]. However, we prefer not to take that approach and cancel the anomalies in B and L using a fourth generation.

This paper is organized as follows: In Sec. II, we discuss baryon number violation in models where B and L are spontaneously broken. The B and L violation in the MSSM is discussed in Sec. III. In Sec. IV, we propose the simplest supersymmetric extension of the model in Ref. [3], while in Sec. IV we summarize our main findings.

II. BARYON NUMBER VIOLATION IN MODELS WITH B AND L SPONTANEOUSLY BROKEN GAUGE SYMMETRIES

Recently, we proposed simple extensions of the standard model where B and L are local gauge symmetries [3]. These models are based on the gauge symmetry, $SU(3)_C \bigotimes SU(2)_L \bigotimes U(1)_Y \bigotimes U(1)_B \bigotimes U(1)_L$ and one introduces a new fermionic family to cancel all baryonic and leptonic anomalies. There are two ways to cancel all baryonic and leptonic anomalies, with a new family of fermions that has the following properties. In Model I, one adds Q'_L, u'_R, d'_R with $B = -1$, and l'_L, e'_R, ν'_R with $L = -3$, while in Model II the new generation has different chirality: Q'_R, u'_L, d'_L, with $B = 1$, and l'_R, e'_L and ν'_L with $L = 3$. Since the new fourth generation fermions have different B and L quantum numbers than the quarks and leptons in the first three generations, it was easy to determine that there are no flavor changing neutral currents at tree level. In order to avoid a stable fourth generation quark, we introduced a new scalar field which is a cold dark matter candidate that coupled the fourth generation fermions to first three generations. For a discussion of the cosmology of these models including the generation of the baryon excess, see [8]. For generic studies of models with fourth generations, see Ref. [9]. Since the fourth generation Yukawa couplings get strong at an energy scale not very far above the TeV scale, these models have a fairly low ultraviolet cutoff and hence it is important that nucleon decay is forbidden even including nonrenormalizable operators of very high dimension. In these models, lepton number and baryon number are broken by the vacuum expectation value of fields S_L and S_B with L and B charges n_L and n_B, respectively. In the calculation of S-matrix elements, lepton number and baryon number violation arises from insertions of the vacuum expectation values of these fields. Possible nucleon decay modes are: $p \to \pi^0 e^+$, $p \to \pi^0 e^+ \nu\nu$, $p \to \pi^0 e^+ \nu\bar{\nu}$, *etc*. All possible nucleon decay modes have $\Delta B = -1$ and $\Delta L = \pm$ an odd natural number. Hence, if $k|n_B| \neq 1$ and/or $k|n_L| \neq$ an odd natural number, for $k = 1, 2, \ldots,$ proton decay is forbidden even allowing nonrenormalizable operators of arbitrarily high dimension. Clearly, it is not difficult to determine that baryon number violating nucleon decay is forbidden in models where baryon number and lepton number are gauged by a suitable choice of the charges n_B and n_L even though these symmetries are spontaneously broken. Note we are assuming here that the gravitino mass is greater than the proton mass. If it is lighter, then final states without a lepton would be allowed. A two-body scattering process that violates baryon number can occur inside of the nucleus. For example, $p + n \to \pi^+\pi^0$, $p + p \to \pi^+\pi^+$, K^+K^+, etc. These along with $n - \bar{n}$ oscillations are forbidden if $k|n_B| \neq 2$. If they are not forbidden, by the value of n_B, the limits they impose on the scale of baryon number symmetry breaking are typically not extremely strong because in the low energy effective theory (below the scales of spontaneous baryon number and weak symmetry breaking) the lowest dimension operators that induce $\Delta B = 2$ transitions have six quark fields and are dimension nine. For a discussion of discrete symmetries that enforce baryon number and lepton number conservation in supersymmetric versions of the standard model, see [10]. Models I and II in Ref. [3] have several scalars with masses that are at or below the weak scale and this requires multiple fine-tunings (i.e., the hierarchy puzzle). Furthermore, even though we assumed the breaking of B and L occurred at the weak scale there was no reason for this to be the case. Motivated by these issues, we study in this letter a simple supersymmetric extension of Model I. The quantum numbers of the quark and lepton fields are the same as Model I in [3] but the scalar representations used to break the symmetry are different. Furthermore, no additional scalars are introduced to prevent the stability of the fourth generation quarks. In the supersymmetric version of Model I that we discuss below, they decay through nonrenormalizable interactions.

III. B AND L VIOLATION IN THE MSSM

The MSSM superpotential up to dimension five is given by

$$\mathcal{W}_{\mathrm{MSSM}} = \mathcal{W}_M + \mathcal{W}_L + \mathcal{W}_B + \mathcal{W}_5. \qquad (1)$$

The first term in the superpotential,

$$\mathcal{W}_M = g_u \hat{Q} \hat{u}^c \hat{H}_u + g_d \hat{Q} \hat{d}^c \hat{H}_d + g_e \hat{L} \hat{e}^c \hat{H}_d + \mu \hat{H}_u \hat{H}_d, \tag{2}$$

contains all the renormalizable terms conserving matter parity, $M = (-1)^{3(B-L)}$. The terms violating L at the renormalizable level appear in

$$\mathcal{W}_L = \epsilon \hat{L} \hat{H}_u + \lambda \hat{L} \hat{L} \hat{e}^c + \lambda' \hat{Q} \hat{L} \hat{d}^c. \tag{3}$$

There is only one term in the MSSM superpotential which violates B at the renormalizable level and it is given by

$$\mathcal{W}_B = \lambda'' \hat{u}^c \hat{d}^c \hat{d}^c. \tag{4}$$

Now, at the nonrenormalizable level one also has the following dimension five operators that violate B and/or L[1]:

$$\mathcal{W}_5 = \frac{\lambda_1}{\Lambda} \hat{Q} \hat{Q} \hat{Q} \hat{L} + \frac{\lambda_2}{\Lambda} \hat{u}^c \hat{d}^c \hat{u}^c \hat{e}^c + \frac{\lambda_3}{\Lambda} \hat{L} \hat{L} \hat{H}_u \hat{H}_u. \tag{5}$$

Using the interaction $\lambda' \hat{Q} \hat{L} \hat{d}^c$ and the term in $\mathcal{W}_B$, one gets the dimension four contributions to proton decay, which predict a lifetime of order $\tau_p \sim 10^{-15}$ years, if the couplings are order one and the squark masses are around a 1 TeV. With similar assumptions, the dimension five operators in $\mathcal{W}_5$ also give unacceptably fast contributions to the decay of the proton even if the Λ scale is close to the Planck scale. For a review on proton decay and a detailed discussion about these contributions, see Ref. [11].

In order to clarify our notation, we list the MSSM superfields:

$$\hat{Q} = \begin{pmatrix} \hat{u} \\ \hat{d} \end{pmatrix} \sim (3, 2, 1/6, 1/3, 0),$$
$$\hat{u}^c \sim (\bar{3}, 1, -2/3, -1/3, 0),$$
$$\hat{d}^c \sim (\bar{3}, 1, 1/3, -1/3, 0),$$
$$\hat{L} = \begin{pmatrix} \hat{\nu} \\ \hat{e} \end{pmatrix} \sim (1, 2, -1/2, 0, 1),$$

and $\hat{e}^c \sim (1, 1, 1, 0, -1)$. Notice that we have included their transformation properties under the gauge group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_B \otimes U(1)_L$, anticipating that we will eventually gauge B and L.

The two MSSM Higgses are given by

$$\hat{H}_u = \begin{pmatrix} \hat{H}_u^+ \\ \hat{H}_u^0 \end{pmatrix} \sim (1, 2, 1/2, 0, 0),$$
$$\hat{H}_d = \begin{pmatrix} \hat{H}_d^0 \\ \hat{H}_d^- \end{pmatrix} \sim (1, 2, -1/2, 0, 0).$$

Adding right-handed neutrinos, $\hat{\nu}^c \sim (1, 1, 0, 0, -1)$, we have the following extra terms in the superpotential

$$\mathcal{W}_\nu = g_\nu \hat{L} \hat{H}_u \hat{\nu}^c + M_\nu \hat{\nu}^c \hat{\nu}^c + \frac{\lambda_4}{\Lambda} \hat{L} \hat{L} \hat{e}^c \hat{\nu}^c + \frac{\lambda_5}{\Lambda} \hat{Q} \hat{L} \hat{d}^c \hat{\nu}^c + \frac{\lambda_6}{\Lambda} \hat{u}^c \hat{d}^c \hat{d}^c \hat{\nu}^c. \tag{6}$$

It is well-known that adding three copies of right-handed neutrinos one can gauge $B - L$ and the dimension four operators that violate baryon and/or lepton number in $\mathcal{W}_B$ and $\mathcal{W}_L$ are not allowed. However, even if we impose $B - L$ as a gauge symmetry, the dimension five contributions to proton decay that arise from couplings in $\mathcal{W}_5$ are allowed. Therefore, one does not resolve the issue of an unacceptably large proton decay rate in SUSY theories just by gauging $B - L$. For a study of the origin of B and L violating interactions in $B - L$ models, see Ref [12]. This is one of the main motivations to consider the SUSY version of the model proposed in Ref. [3].

In Ref. [13], the authors studied a supersymmetric extension of our model in Ref. [3]. However, their motivation was primarily a study of dark matter candidates in the model while our motivation is to construct the simplest possible SUSY extensions of our model that do not permit proton decay even including nonrenormalizable terms of high dimension. We use nonrenormalizable interactions to render the fourth generation quarks unstable instead of adding additional multiplets as was done in [3]. Note that stable color triplet heavy particles give rise to exotic nuclei that form atoms. Limits on the density of such atoms and constraints from big bang nucleosynthesis suggest that stable heavy quarks with masses of a few hundred GeV are not acceptable.

IV. THE MSSM WITH B AND L GAUGED

In order to write the simplest supersymmetric model based on the gauge symmetry $G_{BL} = SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_B \otimes U(1)_L$ and cancel anomalies, we need to introduce chiral superfields for a new fermionic generation. They are:

$$\hat{Q}_4 = \begin{pmatrix} \hat{u}_4 \\ \hat{d}_4 \end{pmatrix} \sim (3, 2, 1/6, -1, 0),$$
$$\hat{u}_4^c \sim (\bar{3}, 1, -2/3, 1, 0),$$
$$\hat{d}_4^c \sim (\bar{3}, 1, -1/3, 1, 0),$$
$$\hat{L}_4 = \begin{pmatrix} \hat{\nu}_4 \\ \hat{e}_4 \end{pmatrix} \sim (1, 2, -1/2, 0, -3),$$
$$\hat{e}_4^c \sim (1, 1, 1, 0, 3),$$
$$\hat{\nu}_4^c \sim (1, 1, 0, 0, 3).$$

We have shown explicitly how the new fermions transform under G_{BL}. We need additional chiral superfields that acquire vacuum expectation values that break B and L. The required new Higgses to break $U(1)_B$ are: $\hat{S}_B \sim (1, 1, 0, -1/3, 0)$ and $\hat{\bar{S}}_B \sim (1, 1, 0, 1/3, 0)$. For the chiral

[1]Note we have not yet gauged B and L.

superfields that break $U(1)_L$, there are two possibilities that we consider: either (i) $\hat{S}_L \sim (1,1,0,0,-6)$ and $\hat{\bar{S}}_L \sim (1,1,0,0,6)$ or (ii) $\hat{S}_L \sim (1,1,0,0,-2)$ and $\hat{\bar{S}}_L \sim (1,1,0,0,2)$.

The superpotential of the theory is given by

$$\mathcal{W}_{BL} = \mathcal{W}_{\text{Yukawa}} + \mathcal{W}_{\text{Higgs}} + \mathcal{W}^5_{BL}, \tag{7}$$

where in case (i),

$$\begin{aligned}\mathcal{W}^{(i)}_{\text{Yukawa}} = {}& g_u \hat{Q}\hat{u}^c \hat{H}_u + g_d \hat{Q}\hat{d}^c \hat{H}_d + g_e \hat{L}\hat{e}^c \hat{H}_d \\ &+ g_\nu \hat{L}\hat{H}_u \hat{\nu}^c + Y_u \hat{Q}_4 \hat{H}_u \hat{u}^c_4 + Y_d \hat{Q}_4 \hat{H}_d \hat{d}^c_4 \\ &+ Y_e \hat{L}_4 \hat{H}_d \hat{e}^c_4 + Y_\nu \hat{L}_4 \hat{H}_u \hat{\nu}^c_4 + \lambda_{\nu^c_4} \hat{\nu}^c_4 \hat{\nu}^c_4 \hat{S}_L.\end{aligned} \tag{8}$$

Here the ordinary three generation neutrinos have Dirac masses and the fourth generation neutrino has both Dirac and Majorana mass terms. The fourth generation neutrino mass must be greater than $M_Z/2$. On the other hand, for case (ii)

$$\begin{aligned}\mathcal{W}^{(ii)}_{\text{Yukawa}} = {}& g_u \hat{Q}\hat{u}^c \hat{H}_u + g_d \hat{Q}\hat{d}^c \hat{H}_d + g_e \hat{L}\hat{e}^c \hat{H}_d \\ &+ g_\nu \hat{L}\hat{H}_u \hat{\nu}^c + Y_u \hat{Q}_4 \hat{H}_u \hat{u}^c_4 + Y_d \hat{Q}_4 \hat{H}_d \hat{d}^c_4 \\ &+ Y_e \hat{L}_4 \hat{H}_d \hat{e}^c_4 + Y_\nu \hat{L}_4 \hat{H}_u \hat{\nu}^c_4 + \lambda_{\nu^c} \hat{\nu}^c \hat{\nu}^c \hat{\bar{S}}_L \\ &+ \lambda'_{\nu^c} \hat{\nu}^c \hat{\nu}^c_4 \hat{S}_L.\end{aligned} \tag{9}$$

The ordinary light three generations of neutrinos have both Majorana and Dirac mass terms and so extremely small Yukawa coupling constants can be avoided using the type I seesaw mechanism [14]. The fourth generation neutrino has a Dirac mass term and a Majorana mass term that mixes it with the first three generations of neutrinos.

The Higgs part of the superpotential is

$$\mathcal{W}_{\text{Higgs}} = \mu \hat{H}_u \hat{H}_d + \mu_B \hat{S}_B \hat{\bar{S}}_B + \mu_L \hat{S}_L \hat{\bar{S}}_L. \tag{10}$$

Finally, the dimension five terms that allow fourth generation particles to decay to the ordinary generations are

$$\mathcal{W}^5_{BL} = \frac{a_1}{\Lambda} \hat{u}^c_4 \hat{d}^c \hat{d}^c \hat{S}_B + \frac{a_2}{\Lambda} \hat{u}^c \hat{d}^c_4 \hat{d}^c \hat{S}_B + \frac{a_3}{\Lambda} \hat{\nu}^c \hat{\nu}^c \hat{\nu}^c \hat{\nu}^c_4. \tag{11}$$

The terms proportional to the a_1 and a_2 couplings are needed to avoid a stable quark from the 4th generation. In case (i), the term proportional to a_3 avoids the presence of a stable heavy Dirac neutrino. Notice that here we write only the relevant dimension five operators.

For simplicity, in our discussions we ignore kinetic mixing between the $U(1)$'s and the possible Fayet-Illiopoulos D-terms.

Symmetry breaking.—Here we investigate the symmetry breaking mechanism to show that $U(1)_B$ and $U(1)_L$ can be broken at the TeV scale. In the case of the $U(1)_B$ symmetry, it is broken by the vacuum expectation value (VEV) of the scalar fields, S_B and $\bar{S}_B$. These vacuum expectation values can be chosen real and positive. The relevant soft terms for our discussion are:

$$-\Delta \mathcal{L}_{\text{Soft}} = (-b_B S_B \bar{S}_B + \text{H.c.}) + m^2_{S_B} |S_B|^2 + m^2_{\bar{S}_B} |\bar{S}_B|^2. \tag{12}$$

For simplicity of notation, we take b_B to be real. Using $\langle S_B \rangle = v_B/\sqrt{2}$ and $\langle \bar{S}_B \rangle = \bar{v}_B/\sqrt{2}$ for the VEVs, one finds

$$\begin{aligned}V_B = {}& \frac{1}{2}|\mu_B|^2 (v_B^2 + \bar{v}_B^2) - b_B v_B \bar{v}_B + \frac{1}{2} m^2_{S_B} v_B^2 \\ &+ \frac{1}{2} m^2_{\bar{S}_B} \bar{v}_B^2 + \frac{g_B^2}{32} n_B^2 (v_B^2 - \bar{v}_B^2)^2.\end{aligned} \tag{13}$$

Now, assuming that the potential is bounded from below along the D-flat direction, we get:

$$2 b_B < 2|\mu_B|^2 + m^2_{S_B} + m^2_{\bar{S}_B}, \tag{14}$$

while

$$b_B^2 > (|\mu_B|^2 + m^2_{S_B})(|\mu_B|^2 + m^2_{\bar{S}_B}), \tag{15}$$

in order to have a nontrivial vacuum. Minimizing with respect to v_B and $\bar{v}_B$, one finds that

$$|\mu_B|^2 + m^2_{S_B} - \frac{1}{2} M^2_{Z_B} \cos 2\beta_B - b_B \cot\beta_B = 0, \tag{16}$$

$$|\mu_B|^2 + m^2_{\bar{S}_B} + \frac{1}{2} M^2_{Z_B} \cos 2\beta_B - b_B \tan\beta_B = 0, \tag{17}$$

with $\tan\beta_B = v_B/\bar{v}_B$ and $M^2_{Z_B} = (n_B g_B)^2 (v_B^2 + \bar{v}_B^2)/4$. Here $n_B = 1/3(-1/3)$ for $\bar{S}_B(S_B)$. The above equations can be written as

$$\frac{1}{2} m^2_{Z_B} = -|\mu_B|^2 - \left(\frac{m^2_{S_B} \tan^2\beta_B - m^2_{\bar{S}_B}}{\tan^2\beta_B - 1} \right), \tag{18}$$

$$b_B = \frac{\sin 2\beta_B}{2} (2|\mu_B|^2 + m^2_{S_B} + m^2_{\bar{S}_B}). \tag{19}$$

For symmetry breaking to occur, $m^2_{Z_B}$ must be positive and it is clear from the above equation that the gauge boson mass is set by the soft supersymmetry breaking terms since they must overpower the negative contribution from the μ_B piece. The $U(1)_B$ symmetry is broken at the SUSY scale.

The analysis of $U(1)_L$ breaking is similar to the breaking of $B-L$ studied in Ref. [15]. Several fields can get a VEV: $\langle S_L \rangle$, $\langle \bar{S}_L \rangle$, $\langle \tilde{\nu} \rangle$ and $\langle \tilde{\nu}^c \rangle$. There are two different cases: i) R-parity conservation, where only S_L and $\bar{S}_L$ can get a VEV, and ii) R-parity is spontaneously broken due to the VEV of sneutrinos. In the latter case, one needs a tachyonic mass term [15] for the "right-handed" sneutrinos. In this paper, we assume that the soft supersymmetry breaking mass terms for the sneutrinos are not tachyonic so that the only fields with lepton number that get a VEV are S_L and $\bar{S}_L$. In the case where the sneutrinos get a VEV, one has R-parity and L violating interactions, which together with the interactions coming from Eq. (20) give rise to proton decay. Since the cutoff in the theory is low due to the existence of the Landau poles for the fourth generation Yukawa couplings, one finds that these contributions give rise to unacceptably fast proton decay. For a study where the sneutrino VEV breaks the leptonic symmetry, see Ref. [12].

In this paper, we do not address the μ problem. The supersymmetric parameters μ, μ_B and μ_L are taken to be of the order of the supersymmetry breaking scale, even though there is no clear reason for this to be the case.

Baryon number violation.—One does not generate operators that mediate proton decay because S_L has an even lepton number charge (see Sec. III). In the MSSM, typically we define matter parity as $M = (-1)^{3(B-L)} = M_L \times M_B$, where $M_L = (-1)^{-3L}$ and $M_B = (-1)^{3B}$ can be called leptonic parity and baryonic parity, respectively. Notice that $M_L = -1$ for all leptons and $+1$ for $\hat{S}_L$ and $\hat{\bar{S}}_L$. All the fields with baryon number have $M_B = -1$. After symmetry breaking, M_L is conserved but M_B is broken. The fact that M_L is conserved tells us that one cannot generate any operator which induces proton decay, because one must break M_L to allow the proton to decay. Note that the absence of proton decay is true even if we include nonrenormalizable operators of arbitrarily high dimension. One can, however, generate $|\Delta B| = 2$ operators that mediate nucleus decay. For example, a dimension seven operator in the superpotential

$$\Delta \mathcal{W}_B^7 = \frac{\tilde{\lambda}''}{\Lambda^3} \hat{u}^c \hat{d}^c \hat{d}^c \hat{\bar{S}}_B^3 \tag{20}$$

generates a contribution to the reaction $^{16}O(pp) \to {}^{14}CK^+K^+$ after integrating out the squarks and the gluino. The relevant dimension nine operator is $C_9 \overline{u^c}\,\overline{d^c}\,\overline{s^c}\,\overline{u^c}\,\overline{d^c}\,\overline{s^c}$, with C_9:

$$C_9 = \left(\frac{\tilde{\lambda}''_{uds} v_B^3}{\Lambda^3}\right)^2 \times \frac{4\pi\alpha_s}{M_{\tilde{s}^c}^4 M_{\tilde{g}}}. \tag{21}$$

Assuming that $M_{\tilde{s}^c}, M_{\tilde{g}} \sim 1$ TeV, and using the experimental limit on this channel from the Super-Kamiokande collaboration, one finds that $\tilde{\lambda}''_{uds} v_B^3/\Lambda^3 < 10^{-8}$ [16,17]. Notice that C_9 can induce $n - \bar{n}$ oscillation at tree level if one assumes flavor violation in the squark sector. Here, for simplicity we do not consider this possibility. At one loop level, one has a contribution to $n - \bar{n}$ oscillations where inside the loops one has the charginos (winos) and the standard model quarks. However, constraints from $n - \bar{n}$ oscillations are weaker than the one from dinucleon decays discussed above. For a review on $n - \bar{n}$ oscillation, see Ref. [18].

The couplings above allow the squarks to decay to two quarks with a partial width of order $\Gamma(\tilde{q} \to q\bar{q}) \sim (\tilde{\lambda}'')^2 \times (v_B/\Lambda)^6/(64\pi)$. This of course means that (apart from the gravitino in models with a high enough scale of spontaneous supersymmetry breaking), the lightest neutralino is not a dark matter candidate.[2] It decays through a virtual squark to three light quarks. Baryon number violating neutralino decay was discussed in Ref. [19]. In this model, the fourth generation squarks decay to quark pairs, which also violates baryon number.

Gauge bosons.—Neglecting kinetic gauge boson mixing, in this theory we have a leptophobic Z_B and quark-phobic Z_L neutral gauge bosons associated to the new symmetries $U(1)_B$ and $U(1)_L$, respectively. For a review on Z' models, see Ref. [20]. The masses of the new neutral gauge bosons are given by

$$m_{Z_B} = \frac{g_B}{6}(v_B^2 + \bar{v}_B^2)^{1/2}, \tag{22}$$

$$m_{Z_L} = \frac{n_L}{2} g_L (v_L^2 + \bar{v}_L^2)^{1/2}, \tag{23}$$

where in case (i) $n_L = 6$ and in case (ii) $n_L = 2$. The collider constraints on a quark-phobic Z' are more severe than the case of Z_B. For the case of Z_L, one can use the LEP2 bounds [21], while for Z_B it is possible to use the UA2 bounds [22].

Neutralinos.—The neutralino sector now has B and L neutralinos in addition to the MSSM neutralinos. In total, one has the MSSM neutralinos $\tilde{\chi}_i^0$, the baryonic neutralinos, $\tilde{\chi}_B^0 = (\tilde{B}_B, \tilde{S}_B, \tilde{\bar{S}}_B)$. Here $\tilde{B}_B$ is the $U(1)_B$ gaugino, and the $\tilde{S}_B$ Higgsinos. Finally, one also has the leptonic neutralinos, $\tilde{\chi}_L^0 = (\tilde{B}_L, \tilde{S}_L, \tilde{\bar{S}}_L)$. Here $\tilde{B}_L$ is the $U(1)_L$ gaugino, and $\tilde{S}_L$ and $\tilde{\bar{S}}_L$ are the superpartners of the Higgses breaking the local leptonic symmetry. It is straightforward to work out the neutralino mass matrices. For example, the neutralino $\tilde{\chi}_B^0$ mass matrix is,

$$\mathcal{M}_{\tilde{\chi}_B^0} = \begin{pmatrix} m_B & -\frac{g_B v_B}{6} & \frac{g_B \bar{v}_B}{6} \\ -\frac{g_B v_B}{6} & 0 & -\mu_B \\ \frac{g_B \bar{v}_B}{6} & -\mu_B & 0 \end{pmatrix}, \tag{24}$$

where m_B is the bino mass, and μ_B is the mass term of the Higgsinos in the baryonic sector. Notice that only when the Higgsino term is small can one have a light neutralino in this sector.

Sfermions and new Higgs spectrum.—After symmetry breaking, the sfermion masses get an extra contribution due to the new D-terms for $U(1)_L$ and $U(1)_B$. See Ref. [23] for a similar study of the spectrum of sfermions of a $U(1)$ extension of the MSSM. Of course we have additional sfermions associated with the fourth generation. In order

[2]Another well motivated dark matter candidate is the axion since it is associated with the solution of the strong CP puzzle.

to illustrate this point, we show the charged MSSM slepton masses

$$M^2_{\tilde{e}_{L_i}} = m^2_{\tilde{L}_i} + m^2_{e_i} - \left(\frac{1}{2} - \sin^2\theta_W\right) M_Z^2 \cos 2\beta + D_L, \tag{25}$$

$$M^2_{\tilde{e}^c_i} = m^2_{\tilde{e}^c_i} + m^2_{e_i} - M_Z^2 \sin^2\theta_W \cos 2\beta - D_L, \tag{26}$$

with

$$D_L = \frac{1}{2n_L} m^2_{Z_L} \cos 2\beta_L. \tag{27}$$

Here $m_{\tilde{L}}$ and $m_{\tilde{e}^c}$ are the soft terms for left- and right-handed sleptons, respectively. The new angle β_L is defined as $\tan\beta_L = v_L/\bar{v}_L$.

There are three new neutral L-Higgses: two CP-even S_{L_1}, S_{L_2} and one CP-odd A_L, while in the $U(1)_B$-sector one has the neutral Higgses S_{B_1}, S_{B_2} and A_B. These two sectors are not coupled to the MSSM sector at tree level through renormalizable interactions. For a recent study of the Higgs decays in the MSSM with four generations, see Ref. [24].

The masses of the Higgses in the baryonic sector are

$$m^2_{S_{B_1}, S_{B_2}} = \frac{1}{2}(m^2_{A_B} + m^2_{Z_B} \mp \sqrt{D}), \tag{28}$$

with

$$D = (m^2_{A_B} - m^2_{Z_B})^2 + 4 m^2_{Z_B} m^2_{A_B} \sin^2(2\beta_B), \tag{29}$$

where

$$m^2_{A_B} = \frac{2b_B}{\sin 2\beta_B}. \tag{30}$$

Notice that the Higgses in this sector can light because the limit on the mass of Z_B is not very strong [22]. In this way, we conclude the discussion of the properties of the spectrum of our model.

V. SUMMARY AND OUTLOOK

In this paper, we have proposed a simple model with baryon and lepton number gauged and spontaneously broken at the supersymmetry breaking scale. After symmetry breaking, the leptonic matter parity is conserved and so proton decay is forbidden (provided the gravitino is heavier than the proton) even when nonrenormalizable operators of arbitrarily high dimension are included.

We have noted some of the important features associated with the spontaneous breaking of baryon number including the implications for dark matter candidates. We have pointed out some properties of the spectrum and possible baryon number violating decays. It is important to mention that in this model one can have lepton number violating signals from the decays of the right-handed neutrinos and baryon number violating signals from the decays of squarks and gauginos without conflict with the bounds coming from proton decay, $n - \bar{n}$ oscillations and dinucleon decays. It would be interesting to investigate the collider signals and cosmological aspects of this model including the possibility of weak scale baryogenesis.

ACKNOWLEDGMENTS

P. F. P. would like to thank the Institute for Advanced Study in Princeton for hospitality in the last part of this project. P. F. P. thanks S. Spinner for several discussions. The work of P. F. P. has been supported in part by the U.S. Department of Energy under Grant No. DE-FG02-95ER40896, and by the Wisconsin Alumni Research Foundation. The work of M. B. W. was supported in part by the U.S. Department of Energy under Contract No. DE-FG02-92ER40701.

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