***B* and *L* at the supersymmetry scale, dark matter, and *R*-parity violation**

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We present a simple theory where baryon and lepton numbers are spontaneously broken at the supersymmetry scale. In this context *R* parity must be spontaneously broken but the theory still contains a stable field which can play the role of the cold dark matter of the Universe. We discuss the spectrum of the theory, the properties of the dark matter candidate and the predictions for direct detection experiments. This theory provides a concrete example of exotic supersymmetric signatures associated with having the simultaneous presence of *R*-parity-violating and missing energy signals at the Large Hadron Collider.

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I. INTRODUCTION

The minimal supersymmetric Standard Model (MSSM) is considered one of the most appealing candidates for physics beyond the Standard Model. While the recent results from the Large Hadron Collider (LHC) have set serious constraints on the masses of the supersymmetric particles, if one suspects that new physics exists at an LHC accessible scale, an MSSM-like theory still highly recommends itself as a candidate theory.

Despite its various appealing properties, the MSSM poses a challenge for proton stability. This is because it introduces two separate sets of operators which induce proton decay: tree-level terms, which separately violate baryon and lepton number, and nonrenormalizable terms which individually violate both baryon and lepton number. The first of these are

$$\hat{L}\hat{H}_u, \quad \hat{L}\hat{L}\hat{e}^c, \quad \hat{Q}\hat{L}\hat{d}^c, \quad \text{and} \quad \hat{u}^c\hat{d}^c\hat{d}^c,$$

where the first three operators violate lepton number and the last baryon number. Any combination of the first three operators with the last one leads to rapid proton decay. Their absence is typically explained by invoking *R* parity, an *ad hoc* discrete symmetry defined as $R = (-1)^{3(B-L)+2S}$, which forbids all of these terms. However, the fate of such operators is most simply divined from models of gauged $B - L$. The most minimal of such models lead to lepton number violating *R*-parity violation (and therefore no tree-level proton decay) [1] but *R*-parity-conserving models are also possible [2–4]. Regardless of the type of $B - L$ model,

the second type of proton decay inducing operators exist. These are nonrenormalizable operators which conserve $B - L$ but violate B and L separately, e.g.,

$$\hat{Q}\hat{Q}\hat{Q}\hat{L}/\Lambda \quad \text{and} \quad \hat{u}^c\hat{u}^c\hat{d}^c\hat{e}^c/\Lambda.$$

Despite the suppression in these terms due to the scale of new physics, the bounds on proton decay are strong enough to motivate a mechanism for suppressing them. See [5] for a review of proton decay.

Recently, a simple theory for the spontaneous breaking of local baryon and lepton numbers has been proposed in Ref. [6]. In this context one can define an anomaly-free theory using fermionic leptoquarks which have both baryon and lepton number charges. Furthermore, even after symmetry breaking, the lightest leptoquark is stable due to a remnant Z_2 symmetry and can therefore be a dark matter candidate. See also Refs. [7–9] for similar studies. This idea can be applied in the context of supersymmetric theories to establish not only the origin of the *R*-parity-violating terms, as in the $B - L$ models, but also determine the fate of the nonrenormalizable terms which violate B and L separately.

In this paper we investigate an extension of the MSSM where the local baryonic and leptonic symmetries are spontaneously broken at the supersymmetry scale. We find that the minimal model predicts that *R* parity must be spontaneously broken in the MSSM sector (leading only to lepton number violation). Despite the breaking of *R* parity, the remnant Z_2 symmetry from the breaking of the baryonic and leptonic symmetries ensures that the lightest leptoquark is stable and may be a candidate for the cold

dark matter of the Universe. We investigate the spectrum of the theory and the predictions for dark matter direct detection. This article is organized as follows: In Sec. II we discuss the model with local B and L symmetries, and in Sec. III we discuss the possible dark matter candidates and the predictions for dark matter experiments. Finally, we summarize our results in Sec. IV.

II. SPONTANEOUS BREAKING OF B AND L

In order to define a theory for local baryon and lepton numbers we use the gauge group

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_B \otimes U(1)_L.$$

An anomaly-free theory can be achieved by adding the following new leptoquark fields with B and L numbers [6]:

$$\begin{aligned} \hat{\Psi} &\sim (1, 2, -1/2, B_1, L_1), & \hat{\Psi}^c &\sim (1, 2, 1/2, B_2, L_2), \\ \hat{\eta}^c &\sim (1, 1, 1, -B_1, -L_1), & \hat{\eta} &\sim (1, 1, -1, -B_2, -L_2), \\ \hat{X}^c &\sim (1, 1, 0, -B_1, -L_1), & \text{and } \hat{X} &\sim (1, 1, 0, -B_2, -L_2). \end{aligned}$$

Notice that these fields are vectorlike with respect to the SM transformations. The anomalies can be canceled for any values of B_i and L_i ($i = 1, 2$) which satisfy the conditions

$$B_1 + B_2 = -3 \quad \text{and} \quad L_1 + L_2 = -3. \quad (1)$$

In order to generate masses for the new fields and for symmetry breaking we need the chiral superfields

$$\hat{S}_1 \sim (1, 1, 0, 3, 3) \quad \text{and} \quad \hat{S}_2 \sim (1, 1, 0, -3, -3).$$

Therefore, the superpotential of this theory is given by

$$\mathcal{W}_{\text{BL}} = \mathcal{W}_{\text{RpC}} + \mathcal{W}_{\text{LB}}, \quad (2)$$

where

$$\begin{aligned} \mathcal{W}_{\text{RpC}} &= Y_u \hat{Q} \hat{H}_u \hat{u}^c + Y_d \hat{Q} \hat{H}_d \hat{d}^c + Y_e \hat{L} \hat{H}_d \hat{e}^c \\ &\quad + Y_\nu \hat{L} \hat{H}_u \hat{\nu}^c + \mu \hat{H}_u \hat{H}_d \end{aligned} \quad (3)$$

contains the R -parity-conserving terms present in the MSSM (plus a Yukawa coupling for the neutrinos, Y_ν) and

$$\begin{aligned} \mathcal{W}_{\text{LB}} &= Y_1 \hat{\Psi} \hat{H}_d \hat{\eta}^c + Y_2 \hat{\Psi} \hat{H}_u \hat{X}^c + Y_3 \hat{\Psi}^c \hat{H}_u \hat{\eta} \\ &\quad + Y_4 \hat{\Psi}^c \hat{H}_d \hat{X} + \lambda_1 \hat{\Psi} \hat{\Psi}^c \hat{S}_1 + \lambda_2 \hat{\eta} \hat{\eta}^c \hat{S}_2 \\ &\quad + \lambda_3 \hat{X} \hat{X}^c \hat{S}_2 + \mu_{\text{BL}} \hat{S}_1 \hat{S}_2 \end{aligned} \quad (4)$$

is the superpotential of the leptoquark sector needed for anomaly cancellation. Of course, because of the conservation of B and L , both the R -parity-violating terms and the nonrenormalizable terms leading to proton decay are forbidden. Notice that when $B_1 = B_2$ and $L_1 = L_2$ we can have Majorana masses for the $\hat{X}$ and $\hat{X}^c$, but we stick to the general case where $B_1 \neq B_2$ and $L_1 \neq L_2$.

An interesting consequence of the leptoquark sector is the presence of a Z_2 symmetry once S_1 and S_2 acquire a

vacuum expectation value (VEV). Under this symmetry, all leptoquarks are odd: $\Psi \rightarrow -\Psi$, $\Psi^c \rightarrow -\Psi^c$, $\eta \rightarrow -\eta$, $\eta^c \rightarrow -\eta^c$, $X \rightarrow -X$ and $X^c \rightarrow -X^c$. The consequence of this is that the lightest leptoquark is stable (must be neutral) and therefore a dark matter candidate.

A. Symmetry breaking and gauge boson masses

Symmetry breaking in the baryon and lepton number sector proceeds through the following scalar potential:

$$\begin{aligned} V &= (M_1^2 + |\mu_{\text{BL}}|^2)|S_1|^2 + (M_2^2 + |\mu_{\text{BL}}|^2)|S_2|^2 \\ &\quad + M_{\tilde{\nu}^c}^2 |\tilde{\nu}^c|^2 + \frac{9}{2} g_B^2 (|S_1|^2 - |S_2|^2)^2 \\ &\quad + \frac{1}{2} g_L^2 (3|S_1|^2 - 3|S_2|^2 - |\tilde{\nu}^c|^2)^2 - (b_{\text{BL}} S_1 S_2 + \text{H.c.}), \end{aligned} \quad (5)$$

where M_1 , M_2 and $M_{\tilde{\nu}^c}$ are the soft terms for the scalar fields S_1 , S_2 and $\tilde{\nu}^c$, respectively. Here b_{BL} is the bilinear term between S_1 and S_2 and we define the VEVs as

$$\sqrt{2} S_1 = v_1 + h_1 + i a_1, \quad (6)$$

$$\sqrt{2} S_2 = v_2 + h_2 + i a_2, \quad (7)$$

$$\sqrt{2} \tilde{\nu}^c = v_R + h_R + i a_R. \quad (8)$$

The squared mass matrix for the new gauge bosons can be written as

$$\mathcal{M}_{Z'}^2 = 9 \begin{pmatrix} g_B^2(v_1^2 + v_2^2) & g_B g_L(v_1^2 + v_2^2) \\ g_B g_L(v_1^2 + v_2^2) & g_L^2(v_1^2 + v_2^2) + \frac{1}{9} g_L^2 v_R^2 \end{pmatrix}, \quad (9)$$

which has a zero determinant if $v_R = 0$; note that this cannot be modified even in the case where $\langle X \rangle \neq 0$ when $B_i = L_i$. We stick to these scenarios for simplicity. This is a consequence of the fact that when S_1 and S_2 acquire VEVs the symmetry group $U(1)_B \otimes U(1)_L$ is broken to $U(1)_{B-L}$. The $B-L$ symmetry can only be broken by the VEV of the right-handed sneutrino as in Ref. [1]. Therefore, we conclude that *R parity must be spontaneously broken in this theory*. However, it is only lepton number violating R -parity violation and therefore the proton remains safe.

The minimization conditions read as

$$\begin{aligned} 0 &= (M_1^2 + |\mu_{\text{BL}}|^2) - b_{\text{BL}} \frac{v_2}{v_1} + \frac{9}{2} g_B^2 (v_1^2 - v_2^2) \\ &\quad + \frac{3}{2} g_L^2 (3v_1^2 - 3v_2^2 - v_R^2), \end{aligned} \quad (10)$$

$$\begin{aligned} 0 &= (M_2^2 + |\mu_{\text{BL}}|^2) - b_{\text{BL}} \frac{v_1}{v_2} - \frac{9}{2} g_B^2 (v_1^2 - v_2^2) \\ &\quad - \frac{3}{2} g_L^2 (3v_1^2 - 3v_2^2 - v_R^2), \end{aligned} \quad (11)$$

$$0 = M_{\tilde{\nu}^c}^2 - \frac{1}{2} g_L^2 (3v_1^2 - 3v_2^2 - v_R^2) \quad (12)$$

and can be reformulated as

$$v_R^2 = \frac{2}{g_L^2} \left[-M_{\tilde{\nu}^c}^2 + \frac{3}{2} g_L^2 (v_1^2 - v_2^2) \right], \quad (13)$$

$$\sin(2\gamma) = \frac{2b_{BL}}{M_1^2 + M_2^2 + 2|\mu_{BL}|^2}, \quad (14)$$

where we have defined

$$\tan \gamma = \frac{v_2}{v_1}. \quad (15)$$

One can easily prove that there is no symmetry breaking in the supersymmetry (SUSY) limit. Therefore, the B and L breaking scales are determined by the SUSY scale. In order to have a potential bounded from below we must satisfy the condition

$$2b_{BL} < M_1^2 + M_2^2 + 2|\mu_{BL}|^2, \quad (16)$$

and in order to break the symmetry we need the condition

$$b_{BL}^2 > \left(M_1^2 + |\mu_{BL}|^2 - \frac{3}{2} g_L^2 v_R^2 \right) \left(M_2^2 + |\mu_{BL}|^2 + \frac{3}{2} g_L^2 v_R^2 \right). \quad (17)$$

The mixing angle between Z_1 and Z_2 is defined by

$$\begin{pmatrix} Z_B \\ Z_L \end{pmatrix} = \begin{pmatrix} \cos \theta_{BL} & \sin \theta_{BL} \\ -\sin \theta_{BL} & \cos \theta_{BL} \end{pmatrix} \begin{pmatrix} Z_1 \\ Z_2 \end{pmatrix}, \quad (18)$$

where $M_{Z_1} < M_{Z_2}$. The eigenvalues for the new neutral gauge bosons are

$$M_{Z_{1,2}}^2 = \frac{1}{2} (M_{Z_L}^2 + M_{Z_B}^2 \pm \sqrt{\Delta_{BL}^2}), \quad (19)$$

where

$$\Delta_{BL}^2 = (M_{Z_L}^2 - M_{Z_B}^2)^2 + 4g_L^2 M_{Z_B}^4 / g_B^2, \quad (20)$$

$$M_{Z_B}^2 \equiv 9g_B^2 (v_1^2 + v_2^2), \quad (21)$$

$$M_{Z_L}^2 \equiv 9g_L^2 \left(v_1^2 + v_2^2 + \frac{1}{9} v_R^2 \right), \quad (22)$$

and the mixing angle is given by

$$\sin(2\theta_{BL}) = \frac{2g_B g_L (v_1^2 + v_2^2)}{M_{Z_2}^2 - M_{Z_1}^2}. \quad (23)$$

Note that this produces a Z_1 lighter than Z_2 only for $M_{Z_B} < M_{Z_L}$. For the opposite case we take $\theta_{BL} \rightarrow -\theta_{BL}$ and $Z_1 \leftrightarrow Z_2$. In the limit $v_R^2 \gg v_1^2 + v_2^2$ the eigenvalues are

$$M_{Z_1} \sim 9g_B^2 (v_1^2 + v_2^2) (1 - 9\epsilon), \quad (24)$$

$$M_{Z_2} \sim g_L^2 v_R^2 (1 + 9\epsilon), \quad (25)$$

where $\epsilon \equiv (v_1^2 + v_2^2)/v_R^2$ and the mass eigenstates are

$$Z_1 = \left(1 - \frac{81}{2} \frac{g_B^2}{g_L^2} \epsilon^2 \right) Z_B - 9 \frac{g_B}{g_L} \epsilon Z_L, \quad (26)$$

$$Z_2 = 9 \frac{g_B}{g_L} \epsilon Z_B + \left(1 - \frac{81}{2} \frac{g_B^2}{g_L^2} \epsilon^2 \right) Z_L. \quad (27)$$

This is an interesting limit since the lighter Z_1 eigenstate is predominately Z_B -like and therefore has lower collider bounds [10,11] compared to a Z' that significantly couples to leptons [12].

Finally, we note that when baryon and lepton numbers are broken at the SUSY scale, one expects operators mediating proton decay. However, in this theory, the proton is stable because baryon number is broken by three units. The least suppressed nonrenormalizable terms which generate baryon- and lepton-number-violating interactions occur at dimension 14, e.g.,

$$\begin{aligned} \mathcal{W}_{14} = \frac{1}{\Lambda^{10}} [& c_1 \hat{S}_1 (\hat{u}^c \hat{u}^c \hat{d}^c \hat{e}^c)^3 + c_2 \hat{S}_1 (\hat{u}^c \hat{d}^c \hat{d}^c \hat{e}^c)^3 \\ & + c_3 \hat{S}_2 (\hat{Q} \hat{Q} \hat{Q} \hat{L})^3]. \end{aligned} \quad (28)$$

Because of this large suppression, there is no need to assume a large scale to be in agreement with experiments.

B. Spontaneous R -parity violation

As we saw earlier, in order to avoid a long range $B - L$ force, the sneutrino must acquire a VEV. The consequences of this are very similar to those in the minimal supersymmetric $B - L$ model [1] and we briefly review them here.

The first and most obvious of these consequences is that R parity is spontaneously broken. This induces a mixing between SUSY and non-SUSY fields with the same quantum numbers: neutralinos with neutrinos, charginos with charged leptons, sneutrinos with neutral Higgs and charged sleptons with charged Higgs. Typically, the most important of these mixings proceeds through the neutrino Yukawa coupling in the superpotential once the right-handed sneutrino acquires a VEV, and one obtains

$$W \supset \frac{1}{\sqrt{2}} Y_\nu v_R \ell \tilde{H}_u, \quad (29)$$

which is the so-called bilinear R -parity-violating term usually referred to as μ' . This term also induces a VEV for the left-handed sneutrino which leads to various mixing terms of gauge coupling strength such as $\frac{1}{2} g_1 \tilde{B} \nu \nu_L$ and $g_L \tilde{B}_L \nu \nu_L$, where $\tilde{B}$ and $\tilde{B}_L$ are the hypercharge and lepton number gauginos, respectively. The size of R -parity violation is related to the neutrino sector and is therefore small. Phenomenologically, this means that SUSY processes proceed as if R parity is conserved except for the decay of the lightest supersymmetric particle (LSP), which can now

decay into SM states. More specifically, SUSY particles are still pair produced. For specific decay channels for a given LSP, see for example [13].

A further interesting consequence is that a sizable VEV can only be realized for one generation of right-handed sneutrinos. This means that lepton number is broken by one unit only in one generation and it is only the corresponding generation of right-handed neutrinos which attains a TeV scale mass; the other two right-handed, or sterile neutrinos, attain sub-eV masses [14–16]. This has important consequences for cosmology in the form of dark radiation in the early Universe and for neutrino oscillation anomalies.

III. DARK MATTER CANDIDATES

After symmetry breaking, the lightest leptoquark is stable due to the remnant Z_2 symmetry as discussed earlier. This particle must be neutral and could play the role of dark matter. Furthermore, unlike in R -parity conservation, the lightest leptoquark can be either a fermion or a scalar. The best candidates are the $\hat{X}$ and $\hat{X}^c$ superfields since they do not couple to the Z . In this study we assume the lightest leptoquark is the fermionic component of $\hat{X}$ and $\hat{X}^c$, whose Dirac spinor we refer to as $\tilde{X}$, and focus on its properties. It is also interesting to note that because the mass of $\tilde{X}$ is given by the VEV of S_2 , it is automatically at the SUSY scale and therefore weakly interacting massive particle like. This would not be true if its mass was simply a parameter in the superpotential whose magnitude would be arbitrary. Of course, there is a trade-off here with the μ -type problem associated with the μ_{BL} parameter.

The fermionic dark matter candidate can annihilate into two fermions through the neutral gauge bosons present in the theory:

$$\tilde{X}\tilde{X} \rightarrow Z_i \rightarrow \bar{f}f. \quad (30)$$

The relevant interactions in this case are

$$-\mathcal{L} = g_B \tilde{X} \gamma^\mu (-B_2 P_L + B_1 P_R) Z_{B\mu} \tilde{X} + g_L \tilde{X} \gamma^\mu (-L_2 P_L + L_1 P_R) Z_{L\mu} \tilde{X}, \quad (31)$$

which in the physical basis reads as

$$-\mathcal{L} = g_B \tilde{X} \gamma^\mu (C_{11} P_L + C_{12} P_R) Z_{1\mu} \tilde{X} + g_B \tilde{X} \gamma^\mu (C_{21} P_L + C_{22} P_R) Z_{2\mu} \tilde{X}, \quad (32)$$

where

$$C_{11} = -B_2 \cos \theta_{BL} + \frac{g_L}{g_B} L_2 \sin \theta_{BL}, \quad (33)$$

$$C_{12} = B_1 \cos \theta_{BL} - \frac{g_L}{g_B} L_1 \sin \theta_{BL}, \quad (34)$$

$$C_{21} = -B_2 \sin \theta_{BL} - \frac{g_L}{g_B} L_2 \cos \theta_{BL}, \quad (35)$$

$$C_{22} = B_1 \sin \theta_{BL} + \frac{g_L}{g_B} L_1 \cos \theta_{BL}. \quad (36)$$

Assuming the contribution from Z_1 dominates, we find an annihilation cross section

$$\sigma v = \sum_q \frac{g_B^4 \tilde{C}}{36\pi s} \frac{\sqrt{1 - 4m_q^2/s}}{(s - M_{Z_1}^2)^2 + \Gamma_{Z_1}^2 M_{Z_1}^2}, \quad (37)$$

where

$$\tilde{C} = [(C_{11}^2 + C_{12}^2)(s + 2m_q^2)(s - M_{\tilde{X}}^2) + 6C_{11}C_{12}M_{\tilde{X}}^2(s + 2m_q^2)] \cos^2 \theta_{BL}. \quad (38)$$

This cross section is given by

$$(\sigma v)_{NR} = \sum_q \frac{g_B^4}{24\pi} \frac{\sqrt{1 - m_q^2/M_{\tilde{X}}^2}}{(4M_{\tilde{X}}^2 - M_{Z_1}^2)^2 + \Gamma_{Z_1}^2 M_{Z_1}^2} C^2 (2M_{\tilde{X}}^2 + m_q^2), \quad (39)$$

in the nonrelativistic limit. Here we have defined

$$C = (C_{11} + C_{12}) \cos \theta_{BL}. \quad (40)$$

In the present epoch the energy density of the relic dark matter particles $\tilde{X}$ would be

$$\Omega_{\tilde{X}} h^2 \simeq \frac{x_f}{\sqrt{g_*} \sigma_0 M_P} \frac{(1.07 \times 10^9)}{\text{GeV}}. \quad (41)$$

We adopt the value $\Omega_{DM}^{obs} h^2 = 0.1199 \pm 0.0027$ [17].

The freeze-out temperature $x_f = M_{\tilde{X}}/T_f$ is then given by

$$x_f = \ln \left(\frac{0.038 g M_P M_{\tilde{X}} \sigma_0}{\sqrt{g_*}} \right) - \frac{1}{2} \ln \left[\ln \left(\frac{0.038 g M_P M_{\tilde{X}} \sigma_0}{\sqrt{g_*}} \right) \right], \quad (42)$$

where g is the number of internal degrees of freedom (in our case $g = 4$), g_* is the effective number of relativistic degrees of freedom evaluated around the freeze-out temperature, M_P is the Planck mass, and we use the expansion $\sigma v = \sigma_0 + \sigma_1 v^2$.

The direct detection also proceeds through the Z_i :

$$\tilde{X} N \rightarrow Z_i \rightarrow \tilde{X} N, \quad (43)$$

and the spin-independent nucleon-dark matter cross section is then given by

$$\sigma_{SI} = \frac{1}{4\pi} \frac{M_{\tilde{X}}^2 M_N^2}{(M_{\tilde{X}} + M_N)^2} \frac{g_B^4}{M_{Z_1}^4} C^2, \quad (44)$$

assuming that the dominant contribution comes from the Z_1 gauge boson. Because both the dark matter annihilation and

FIG. 1 (color online). Predictions for the elastic nucleon-dark matter cross section for different values of the dark matter mass when $0.11 < \Omega_{\tilde{\chi}} h^2 < 0.13$.

FIG. 2 (color online). Allowed values for the gauge coupling and the dark matter mass when $0.11 < \Omega_{\tilde{\chi}} h^2 < 0.13$ and $M_{Z_1} = 2, 3, 4$ TeV.

direct detection proceed through Z_1 , they are intimately related to each other. Specifically, once one determines the parameters that yield the correct relic density for a given dark matter mass, there are no free parameters left to hide it from direct detection. Keeping this in mind we present the predictions for the direct detection experiments.

In Fig. 1 we show the values for the spin-independent cross section versus the dark matter mass when $C = 1$, $0.1 \leq g_B \leq 0.3$, $2.5 \text{ TeV} \leq M_{Z_1} \leq 5 \text{ TeV}$, and assuming that the relic density is in the range $0.11 < \Omega_{\tilde{\chi}} h^2 < 0.13$. One can appreciate in Fig. 1 that the allowed solutions are below the XENON100 bounds [18] but could be tested in future dark matter experiments such as XENON1T or LUX.

In Fig. 2 we show some solutions when the mass of new lightest neutral gauge boson is 2, 3, or 4 TeV. One can see that there is no need to be very close to the resonance to achieve the required cross section for the relic density.

IV. SUMMARY AND DISCUSSIONS

In this article we have presented the simplest supersymmetric extension of the model proposed in Ref. [6] where baryon and lepton number are local symmetries. In this context the baryonic and leptonic gauge symmetries are broken at the SUSY scale and the proton is stable.

One of the main predictions of this theory is that R parity must be spontaneously broken in the MSSM sector because the right-handed sneutrino VEVs are needed to break the remnant local $U(1)_{B-L}$ that results from the VEVs of S_1 and S_2 . Even though R parity is broken, the lightest leptoquark is stable and can be a cold dark matter candidate. The dark matter candidate can be either the spin one-half or spin zero SM singlet leptoquark; we have focused on the former in this paper. It furthermore has baryon and lepton number and therefore couples to the two Z 's in the model.

There are many interesting predictions for the Large Hadron Collider searches in this theory. Since R parity is broken in the MSSM sector we will have lepton-number-violating signatures at the LHC. For example, one can have exotic channels with four leptons and four jets where three of the leptons have the same electric charge [13,19]. On the other hand there is a stable dark matter candidate in the theory which can be produced through the new neutral gauge bosons. Therefore, one can also expect signatures with missing energy at the LHC. This theory provides a simple example of very exotic supersymmetric signatures at colliders since one can have the simultaneous presence of R -parity-violating and missing energy signatures at the LHC.

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