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Matter-parity as a residual gauge symmetry: Probing a theory of cosmological dark matter

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ABSTRACT

We discuss a non-supersymmetric scenario which addresses the origin of the matter-parity symmetry, $P_M = (-1)^{3(B-L)+2s}$, leading to a viable Dirac fermion dark matter candidate. Implications to electroweak precision, muon anomalous magnetic moment, flavor changing interactions, lepton flavor violation, dark matter and collider physics are discussed in detail. We show that this non-supersymmetric model is capable of generating the matter-parity symmetry in agreement with existing data with gripping implications to particle physics and cosmology.

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1. Introduction

The nature of dark matter is one of the most challenging problems in science, requiring physics beyond the Standard Model as well as a new symmetry capable of making the corresponding particle stable on cosmological scales. R-parity is a symmetry imposed by hand in supersymmetry in order to avoid fast proton decay, leading also to the existence of a stable Weakly Interacting Massive Particle (WIMP), one of the most compelling dark matter candidates [1]. Even if imposed by hand, R-parity may still break through high dimension operators [2–4] or spontaneously [5,6]. While the second case leads to an attractive neutrino mass generation scheme [7], one loses the WIMP dark matter scenario [8]. Generally, some sort of symmetry should be invoked in order to stabilize the dark matter candidate, and it is desirable that this stability arises from gauge principles [9–13]. For example, an alternative to R-parity in non-supersymmetric schemes is to impose a discrete lepton number symmetry to stabilize the WIMP dark matter particle [14–16].

In this work, we discuss a non-supersymmetric model where dark matter stability results from the matter-parity symmetry $P_M = (-1)^{3(B-L)+2s}$, naturally arising as a consequence of the spontaneous breaking of the gauge symmetry inspired by the works done in [17,18]. In order to implement this idea we consider an extension of the standard model based upon an extended $\mathrm{SU(3)_c \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N}$ electroweak symmetry broken by Higgs triplets preserving $B-L$. Note that the $\mathrm{SU}(3)_L$ symmetry is well-motivated due to its ability to determine the number of generations to match that of colors by the anomaly cancellation requirement [19–21]. matter-parity symmetry $P_M = (-1)^{3(B-L)+2s}$ arises in our model as a result of spontaneous gauge symmetry breaking, and the stability of the lightest R_P-odd particle leads to a viable Dirac fermion WIMP dark matter candidate. We work out the expected rates for direct detection experiments, flavor changing neutral currents, lepton flavor violation processes such as $\mu \to e\gamma$, as well as high energy collider signatures. We also comment on possible connections to cosmological inflation and leptogenesis.

2. The model

Our non-supersymmetric model is based on the $\mathrm{SU(3)_c \otimes SU(3)_L \otimes U(1)_X \otimes U(1)_N}$ gauge group, in which the matter generations are arranged in the fundamental representation of $SU(3)_L$ as follows,

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Table 1
The X and N charges of the various multiplets. Gauge fields have $X = N = 0$ and are not listed.

Multiplet	l_{aL}	ν_{aR}	e_{aR}	N_{aR}	$q_{\alpha L}$	q_{3L}	u_{aR}	d_{aR}	U_R	$D_{\alpha R}$	η	ρ	χ	ϕ
X	$-1/3$	0	-1	0	0	$1/3$	$2/3$	$-1/3$	$2/3$	$-1/3$	$-1/3$	$2/3$	$-1/3$	0
N	$-2/3$	-1	-1	0	0	$2/3$	$1/3$	$1/3$	$4/3$	$-2/3$	$1/3$	$1/3$	$-2/3$	2

Leptons	1-2nd Generations	3th Generation
$l_{aL} = \begin{pmatrix} \nu_a \\ e_a \\ N_a \end{pmatrix}_L$	$q_{\alpha L} = \begin{pmatrix} d_\alpha \\ -u_\alpha \\ D_\alpha \end{pmatrix}_L$	$q_{3L} = \begin{pmatrix} u_3 \\ d_3 \\ U \end{pmatrix}_L$
ν_{aR}, e_{aR}, N_{aR}	$u_{\alpha R}, d_{\alpha R}, D_{\alpha R}$	u_{3R}, d_{3R}, U_R

Scalars

$$\eta = \begin{pmatrix} \eta_1^0 \\ \eta_2^- \\ \eta_3^0 \end{pmatrix} \qquad \rho = \begin{pmatrix} \rho_1^+ \\ \rho_2^0 \\ \rho_3^+ \end{pmatrix} \qquad \chi = \begin{pmatrix} \chi_1^0 \\ \chi_2^- \\ \chi_3^0 \end{pmatrix}, \qquad \phi$$

where we have adopted the generation indices $a = 1, 2, 3$ and $\alpha = 1, 2$. We emphasize that all leptons are placed in the fundamental representation of $SU(3)_L$, along with the third generation of quarks. However, the first and second generation of quarks are arranged in the anti-fundamental representation of $SU(3)_L$ to cancel the gauge anomalies. Models based on this gauge group can feature different realizations and fermion contents [22–40].

The generators of the Abelian $U(1)_X$ and $U(1)_N$ groups obey the following relations,

$$Q = T_3 - \frac{1}{\sqrt{3}} T_8 + X, \qquad B - L = -\frac{2}{\sqrt{3}} T_8 + N, \tag{1}$$

where T_i $(i = 1, 2, 3, ..., 8)$, X and N are the charges of $SU(3)_L$, $U(1)_X$ and $U(1)_N$, respectively [17,18,41–45]. The exotic quarks U and D have electric charge $2/3$ and $-1/3$ respectively. The quantum numbers associated to the $U(1)_X$ and $U(1)_N$ symmetries are collected in Table 1.

Gauge symmetry breaking by these SU(3) Higgs triplets and singlet addresses the origin of matter-parity conservation and dark matter stability by preserving $B - L$. Indeed, after the scalar ϕ develops a vacuum expectation value (VEV) at scale Λ, the continuous $U(1)_N$ symmetry is spontaneously broken down to the discrete matter-parity given as $P_M = (-1)^{3(B-L)+2s} = (-1)^{-2\sqrt{3}T_8+3N+2s}$.

We emphasize that this is the only plausible way to embed the $B - L$ symmetry in the model and naturally explain the origin of the matter-parity, since $SU(3)_L$ and $B - L$ symmetries neither commute nor close algebraically as shown in detail in [18]. We also note that the exotic fermions have the following $B - L$ quantum numbers, $[B - L](N_a, D_\alpha, U) = 0, -2/3, 4/3$, and hence are R_P-odd. The new Abelian gauge groups give rise to two new neutral gauge bosons with masses proportional to the B-L and $SU(3)_L$ symmetry breaking scales, respectively. Unless otherwise stated we will assume that the B-L symmetry is broken at very high energy scales, implying that only one new neutral gauge boson, Z', will be phenomenologically relevant. Concerning the exotic quarks they are sufficiently heavy since their masses are proportional to $w = \langle \chi_3^0 \rangle$, the VEV of the χ_3^0 field, taken to be larger than 10 TeV.

Note that in our model the N_{aR} are truly singlets under the gauge group, in contrast to the ν_{aR} which transform under $U(1)_N$, so they have Dirac masses $h^N_{ab} w$ proportional to $w = \langle \chi_3^0 \rangle$. For all cases, N_a can be made the lightest odd particle under matter-parity, and therefore it is a Dirac dark matter candidate (see *Supplemental Material* for alternative assumptions). In what follows, we will investigate the phenomenological consequences of our model. We start by addressing electroweak limits.

3. CKM unitarity

Quantum loop corrections to the quark mixing matrix resulting from additional neutral gauge bosons induce deviations from unitarity of the CKM matrix. Other new particles such as right-handed neutrinos can also give rise to violation of CKM unitarity but these are suppressed by the mass-mixing with active neutrinos and for this reason we neglect them [46–48]. These contributions appear as box-diagrams involving W-gauge bosons and the Z' gauge boson leading to hadronic β-decay, where the CKM matrix can be extracted from. Such contribution can be parametrized by $\Delta_{\rm CKM} = 1 - \sum_{q=d,s,b} |V_{u,q}|^2$ [49]. Applying this to the neutral current we find,

$$\Delta_{\rm CKM} = -0.0033 \frac{M_W^2}{M_{Z'}^2} \ln\left(\frac{M_W^2}{M_{Z'}^2}\right) \tag{2}$$

which implies into $M_{Z'} \gtrsim 200$ GeV.

4. Electroweak precision tests

New physics contributions to the ρ-parameter come from the mixing among the neutral gauge bosons, which is evaluated by [18],

$$\Delta\rho \equiv \frac{M_W^2}{c_W^2 M_{Z_1}^2} - 1 \simeq \frac{(c_{2W} u^2 - v^2)^2}{4 c_W^4 (u^2 + v^2) w^2} + \frac{t_W^4 (u^2 + v^2)}{36 \Lambda^2}, \tag{3}$$

where u, v, w, and Λ are the VEVs of η_1, ρ_2, χ_3, and ϕ, respectively. Here Z_1 is the lightest of the massive neutral gauge bosons, i.e. the Standard Model Z boson in the limit where the scale w is sufficiently high. By enforcing the experimental limit $\Delta\rho < 0.0006$ [50], we find the bounds summarized in Table 3, taking into account that $u^2 + v^2 = v_{SM}^2$, $v_{SM} = 246$ GeV, $s_W^2 = 0.231$, $\alpha = 1/128$, and $g^2 = 4\pi\alpha/s_W^2$. We have checked that one-loop new physics corrections to the ρ-parameter dominantly arise from the new non-Hermitian gauge bosons, but they are much smaller than the tree-level one.

A more robust bound, insensitive to the VEV hierarchy, stems from the amazing precision achieved by LEP. This still provides a good test for new neutral gauge bosons that couple to leptons via the $e^+e^- \to Z' \to f\bar{f}$ production channel with Z' being off-shell. The bound can be obtained using the parametrization [51],

$$\mathcal{L} = \frac{g^2}{c_W^2 M_{Z'}^2} [\bar{e}\gamma^\mu (a_L^e P_L + a_R^e P_R) e][\bar{f}\gamma_\mu (a_L^f P_L + a_R^f P_R f] \tag{4}$$

where $a_L^f = (g_V^f + g_A^f)/2$ and $a_R^f = (g_V^f - g_A^f)/2$.

Using the LEP-II results for final state dileptons we get [52],

$$\frac{g^2}{\cos^2(\theta_W)} \left[\frac{\cos(2\theta_W)}{2\sqrt{3 - 4\sin^2(\theta_W)}} \right] \frac{1}{M_{Z'}^2} < \frac{1}{(6\,\text{TeV})^2} \tag{5}$$

which translates into $M_{Z'} > 1.93$ TeV.

Table 2
The couplings of Z' with fermions.

f	g_V^f	g_A^f
ν_a	$\frac{c_{2W}}{2\sqrt{3-4s_W^2}}$	$\frac{c_{2W}}{2\sqrt{3-4s_W^2}}$
e_a	$\frac{1-4s_W^2}{2\sqrt{3-4s_W^2}}$	$\frac{1}{2\sqrt{3-4s_W^2}}$
N_a	$-\frac{c_W^2}{\sqrt{3-4s_W^2}}$	$-\frac{c_W^2}{\sqrt{3-4s_W^2}}$
u_α	$-\frac{3-8s_W^2}{6\sqrt{3-4s_W^2}}$	$-\frac{1}{2\sqrt{3-4s_W^2}}$
u_3	$\frac{3+2s_W^2}{6\sqrt{3-4s_W^2}}$	$\frac{c_{2W}}{2\sqrt{3-4s_W^2}}$
d_α	$-\frac{(3-2s_W^2)}{6\sqrt{3-4s_W^2}}$	$-\frac{c_{2W}}{2\sqrt{3-4s_W^2}}$
d_3	$\frac{\sqrt{3-4s_W^2}}{6}$	$\frac{1}{2\sqrt{3-4s_W^2}}$

Table 3
Bounds on the w symmetry breaking scale from electroweak measurements on the ρ-parameter. This bound can be translated into limits on the gauge boson masses $M_{Z'}^2 \simeq \frac{g^2 c_W^2 w^2}{3-4s_W^2}$ and $M_{W'}^2 = \frac{g^2(v_{SM}^2+w^2)}{4}$.

u/v	0	1	∞
w [TeV]	6.53	1.5	3.51
$M_{Z'}$ [TeV]	2.58	0.593	1.4
$M_{W'}$ [TeV]	2.12	0.494	1.14

5. Muon magnetic moment

Any fundamental charged particle has a magnetic dipole moment (g) which is parametrized in terms of $a = (g-2)/2$. In the case of the electron the SM prediction agrees quite well with the experimental observation, constituting a capital example of the success of quantum field theory. On the other hand, for the muon, there is a long standing discrepancy between theory and measurement of about 3.6σ [53]. This translates into $\Delta a_\mu = (287 \pm 80) \times 10^{-11}$ [54,55]. The ongoing $g-2$ experiment at FERMILAB will shed light into this problem and, should the central value remain intact, a 5σ evidence for new physics would result, with $\Delta a_\mu = (287 \pm 34) \times 10^{-11}$. The model presented here cannot account for $g-2$, since the required gauge boson masses would be too small to fulfill current experimental limits from high energy colliders (see below). Hence one can only require their contribution to lie within the error bars. Using Table 2, and current (projected) sensitivities on the muon magnetic moment we find, $M_{Z'} > 180\,\text{GeV}$ $(273\,\text{GeV})$, $M_{W'} > 100\,\text{GeV}$ $(145\,\text{GeV})$.

6. Flavor changing neutral current

Mesons are unstable systems, but if their lifetime is sufficiently long we can observe them at colliders. The K^0 meson, a bound state of $d\bar{s}$, is necessarily different from its antiparticle due to strangeness. As a result of CP violation in weak interactions, these mesons decay differently, and their mass difference has been used as a sensitive probe for flavor changing interactions [56]. Similar discussion holds for the $B_s^0 - \bar{B}_s^0$ meson system. Defining $V_{\text{CKM}} = U_L^\dagger V_L$, with $U_L(V_L)$ being the matrix relating the flavor states to the mass-eigenstates of positive (negative) isospin, one can find that the contribution to the mass difference for meson systems is [18],

$$K^0 - \bar{K}^0 : \frac{g^2}{(3-t_W^2)M_{Z'}^2}[(V_L^*)_{31}(V_L)_{32}]^2 < \frac{1}{(10^4\,\text{TeV})^2}, \tag{6}$$

$$B_s^0 - \bar{B}_s^0 : \frac{g^2}{(3-t_W^2)M_{Z'}^2}[(V_L^*)_{32}(V_L)_{33}]^2 < \frac{1}{(100\,\text{TeV})^2}. \tag{7}$$

The bound derived on the mass of the Z' gauge boson is rather sensitive to the parametrization used for the V matrix that diagonalizes the CKM matrix. In [57] two possible parametrizations were considered, that yield either optimistic or conservative limits, while keeping the CKM matrix in agreement with data. In the optimistic one, one finds $V_{31} = 0.43$, $V_{32} = 0.089$, $V_{33} = 0.995$, with the $K - \bar{K}^0$ system producing the strongest limit, $M_{Z'} >$ 14 TeV. Taking a conservative approach, one finds $V_{31} = 0.00037$, $V_{32} = 0.052$, $V_{33} = 0.998$, with the $B_s^0 - \bar{B}_s^0$ system offering a better probe, implying the lower bound $M_{Z'} > 1.95$ TeV. Thus it is clear that meson systems can be powerful tests for new physics effects, although suffer from sizeable uncertainties. In this work we will adopt the conservative bound, but bear in mind that more stringent limit may be applicable. See [58–70] for complementary studies on flavor changing interactions.

7. Dilepton resonance searches at the LHC

Dilepton resonance searches are the gold channel for heavy neutral gauge bosons with un-suppressed couplings to leptons [71]. Since signal events are peaked at the Z' mass, the use of cuts on the dilepton invariant mass is a powerful discriminating tool. The background comes mainly from Drell–Yann processes and is well understood [72,73]. Using 13 TeV center-of-energy with integrated luminosity of $\mathcal{L} = 36.1\,\text{fb}^{-1}$, and applying the cuts,

- $E_T(e_1) > 30\,\text{GeV}$, $E_T(e_2) > 30\,\text{GeV}$, $|\eta_e| < 2.5$,
- $p_T(\mu_1) > 30\,\text{GeV}$, $p_T(\mu_2) > 30\,\text{GeV}$, $|\eta_\mu| < 2.5$,
- $500\,\text{GeV} < M_{ll} < 6000\,\text{GeV}$,

with M_{ll} denoting the dilepton invariant mass, one can find a bound on the Z' mass [74].[1] We have generated events with MadGraph5 [78,79], adopting the CTEQ6L parton distribution function [80] and efficiencies/acceptances as described in [74]. The resulting limit was found to be $M_{Z'} > 4.25$ TeV, superseding previous studies [75–77,81–84]. Keeping a similar detector response we expect that upcoming LHC runs with $\mathcal{L} = 100(1000)\,\text{fb}^{-1}$ will probe $M_{Z'} = 4.9\,(6.1)$ TeV, respectively.

8. Charged lepton + MET at the LHC

The presence of a charged gauge boson (W') is a feature shared among all models based on the $SU(3)_L$ gauge group. In order to constrain the mass of this charged gauge boson one looks for high transverse mass signal events [85,86]. Here one can use the lepton plus missing energy data, via the $pp \to W' \to l\nu$ production channel at the LHC with $\mathcal{L} = 13.3\,fb^{-1}$ and 13 TeV center of mass energy. No significant excess above Standard Model predictions was seen, leading to $M_{W'} > 4.74$ TeV [86]. In this model, the charged current contains,

$$\mathcal{L} \supset -\frac{g}{\sqrt{2}}\left(\bar{e}_{aL}\gamma^\mu N_{aL}\right) W'_\mu. \tag{8}$$

Since this charged gauge boson will be assumed to be much heavier than the lightest N (i.e. odd fermion in our case), we expect that the signal events will have approximately the same cut efficiencies observed in the ATLAS study. Given that the interactions of the Lagrangian in Eq. (8) is similar to the one considered in W' searches, the bound above is expected to be applicable to

[1] See [57] and [75–77] for previous studies.

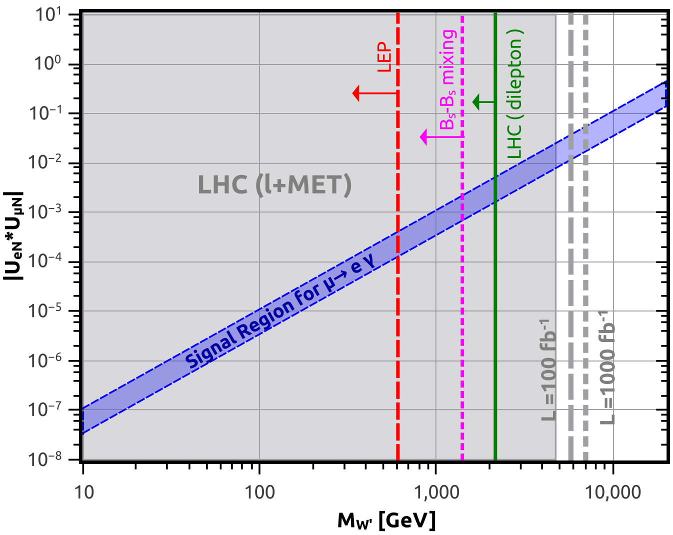

Fig. 1. Region of parameters yielding $4.2 \times 10^{-13} < \mathrm{Br}(\mu \to e\gamma) < 4 \times 10^{-14}$ in blue, overlaid with bounds from LEP (dashed red), $B_s^0 - \bar{B}_s^0$ mixing (dashed pink), dilepton data from LHC (solid green), and l+MET data from LHC in gray. The upper blue line in the region represents the current limit $\mathrm{Br}(\mu \to e\gamma) < 4.2 \times 10^{-13}$. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

our model. This limit is represented by the gray region in Fig. 1. We also looked at the prospects for future runs from the LHC at 13 TeV, with $\mathcal{L} = 100(1000)\ \mathrm{fb}^{-1}$ which turn out to be sensitive to $M_{W'} = 5.8\,(7)$ TeV.

9. Lepton flavor violation

In the Standard Model lepton flavor is conserved and neutrinos are massless. However, neutrinos experience flavor oscillations [87–89] which is a direct confirmation that leptonic flavor is violated. An observation of charged lepton flavor violation would have enormous impact on our understanding of the lepton sector and could have important implications for new physics. Indeed, the existence of lepton flavor violation in neutrino propagation suggests that it should also exist in the charged lepton sector, leading to decays such as $\mu \to e\gamma$. Unfortunately the connection is highly model-dependent [90]. In our model, the presence of right-handed neutrinos (i.e. odd fermions), with the lightest one constituting the dark matter, can mediate a fast decay $\mu \to e\gamma$ via W' exchange, with a branching ratio found to be [55],

$$\mathrm{Br}(\mu \to e\gamma) = 6.43 \times 10^{-6} \left(\frac{1\ \mathrm{TeV}}{M_{W'}}\right)^4 \sum_f (g^{fe*} g^{f\mu})^2, \tag{9}$$

with $g^{fe} = g\, U^{Ne*}/(2\sqrt{2})$ and $g^{f\mu} = g\, U^{N\mu*}/(2\sqrt{2})$.

Current (projected) sensitivity as reported by the MEG Collaboration [91] implies that $\mathrm{Br}(\mu \to e\gamma) < 4.2 \times 10^{-13}$ (4×10^{-14}). Thus one can translate this bound into a limit on the product $U^{Ne*}U^{N\mu}$ as function of the W' mass as shown in Fig. 1. There we have overlaid the aforementioned constraints altogether as indicated in the caption. There we have converted the limits on the Z' mass into bounds on the W' knowing that their mass are determined by a common energy scale w. We conclude that depending on the value for the product $U^{Ne*}U^{N\mu}$ the $\mu \to e\gamma$ search may outperform collider probes. We now investigate the feasibility of this model concerning dark matter searches.

10. WIMP Dirac dark matter

In our model, one can have either a Dirac or Majorana fermionic dark matter [30,92–99], though in this work we focus on the Dirac possibility, since the Majorana case is already excluded by combining the existing constraints (*see Appendix A*).

The Dirac dark matter candidate is the neutral fermion N. As we discussed earlier the dark matter mass can be regarded as a free parameter. The current dark matter relic density and scattering rate at underground detectors are dictated, respectively, by the s-channel and t-channel Z'-induced interactions that result from the neutral current. The vector and axial-vector couplings are exhibited in Table III of [18]. The W' boson also mediates t-channel interactions, which are nevertheless subdominant, and thus neglected in our computations. The s-channel Z' mediated process induces the dark matter annihilation into SM fermions, whereas the t-channel diagram accounts for the annihilation into Z' pairs. The relic density calculation in the context of vector mediators has been performed (see [100]). Our findings fully agree with their result and they read,

$$\begin{aligned}\sigma v\left(N\bar{N} \to f\bar{f}\right) \approx\ & \frac{n_c\sqrt{1-\frac{m_f^2}{M_N^2}}}{2\pi m_{Z'}^4\left(4M_N^2 - m_{Z'}^2\right)^2} \\ & \times \Bigg\{ g_{fa}^2 \Big[2 g_{\chi v}^2 m_{Z'}^4 \left(M_N^2 - m_f^2\right) \\ & + g_{\chi a}^2 m_f^2 \left(4M_N^2 - m_{Z'}^2\right)^2 \Big] \\ & + g_{\chi v}^2 g_{fv}^2 m_{Z'}^4 \left(2M_N^2 + m_f^2\right) \Bigg\}, \end{aligned} \tag{10}$$

where v is the relative velocity of the annihilating DM pair, n_c is the number of colors of the final state SM fermions. E.g. $n_c = 3$ for quarks. Sufficiently near resonance, the Z' must be included in Eq. (10). Anyhow, if $M_N > m_{Z'}$, then N may also self-annihilate into on-shell Z' bosons yielding,

$$\begin{aligned} & \sigma v\left(N\bar{N} \to Z'Z'\right) \\ & \approx \frac{1}{16\pi M_N^2 m_{Z'}^2}\left(1 - \frac{m_{Z'}^2}{M_N^2}\right)^{3/2}\left(1 - \frac{m_{Z'}^2}{2M_N^2}\right)^{-2} \\ & \quad \times \left[8 g_{\chi v}^2 g_{\chi a}^2 M_N^2 + \left(g_{\chi v}^4 + g_{\chi a}^4 - 6 g_{\chi v}^2 g_{\chi a}^2\right) m_{Z'}^2\right]. \end{aligned} \tag{11}$$

We have handled our numerical calculations within Micromegas where the Z' width is properly accounted for [101,102].

Since the WIMP might annihilate into SM fermions or Z′ pairs, which have different annihilation cross sections as shown above, the relic density curve represented in *red* in Fig. 2 features different dependences on the dark matter and mediator mass (M_N and $m_{Z'}$). We enforced the dark matter relic density to be 27% of the universe budget as indicated by PLANCK data [103].

The WIMP-nucleon scattering cross section which is mediated by a t-channel Z' exchange is spin-independent and thus suffers from stringent limits from XENON1T experiment [104]. It reads,

$$\sigma^{SI} \approx \frac{\mu_{Nn}^2}{\pi}\left[\frac{Zf_p + (A-Z)f_n}{A}\right]^2 \tag{12}$$

where,

$$f_p \equiv \frac{1}{m_{Z'}^2}\left(2g_{uV} + g_{dV}\right), \tag{13}$$

and

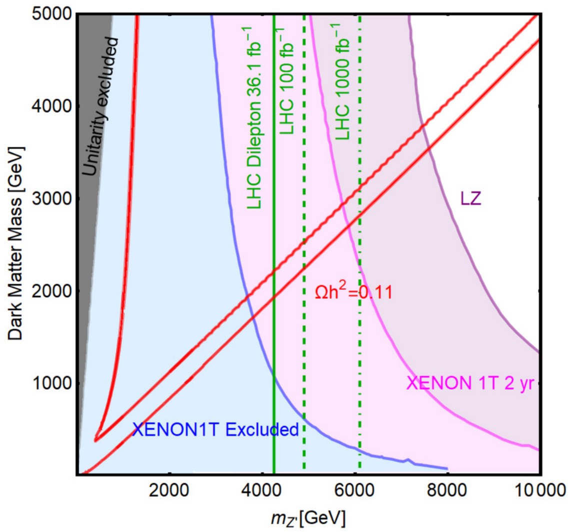

Fig. 2. Region of parameters that yields the right relic density curve in red. The existing limits from the non-observation of dark matter matter-nucleon scattering by the LUX Collaboration are indicated in light blue [105]. The prospects for the XENON1T experiment with 2-years exposure [105], as well as the projected sensitivity of the LZ dark matter experiment are also indicated [106]. Current limits as well as projected sensitivities from LHC searches of dilepton resonances for luminosities of 100 fb^{-1} and 1000 fb^{-1} are also shown. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

$$f_n \equiv \frac{1}{m_{Z'}^2}\,(g_{uV} + 2g_{dV})\,, \tag{14}$$

where g_{uV} and g_{dV} are the N couplings to up- and down-quarks which are simply the product of the $Z' - q - q$ with $N - N - Z'$ couplings, where $q = u, d$ given in Table III of [18]. $\mu_{\chi n}$ is the WIMP-nucleon reduced mass, and where Z and A are the atomic number and atomic mass of the target nucleus, respectively.

In Fig. 2 we delimit current exclusion limit of XENON1T [104] with 34 days-Ton exposure, in *blue* the projected exclusion bounds from XENON1T [105] with 2-years-ton exposure and LZ [106] experiments in *pink* and *purple* respectively. It is clear that current limits from the LHC and dark matter direct detection are rather complementary; in particular LHC can test the WIMP paradigm for higher values of the dark matter mass. It is nevertheless exciting to observe that next generation direct detection experiments, i.e. XENON1T and LZ, are expected to probe the model for Z' mases up to 10 TeV outperforming the LHC.

In conclusion we have shown that this UV complete dark matter model addresses the origin of matter-parity from first principles is amenable to a variety of constraints with gripping implication to cosmology, and additionally offers a viable dark matter candidate for $m_{Z'} > 4$ TeV and dark matter masses of 2–5 TeV.

11. Conclusion

Is summary, we have discussed an elegant solution to the origin of a matter-parity symmetry, $P_M = (-1)^{3(B-L)+2s}$, and consequently the stability of the dark matter particle, in a non-supersymmetric model. The fact that the $B - L$ symmetry is preserved at high scales plays a key role in accounting for the origin of matter-parity conservation. The lightest P_M-odd particle constitutes a viable dark matter candidate, whose stability follows naturally from the breaking of the gauge symmetry. We have shown that the scheme offers good prospects for dark matter detection in nuclear recoil experiments, as well as flavor changing neutral currents in the neutral meson systems $K^0 - \bar{K}^0$ and $B^0 - \bar{B}^0$, searches for lepton flavor violating processes such as $\mu \to e\gamma$, as well as dilepton and l + MET event searches at the LHC.

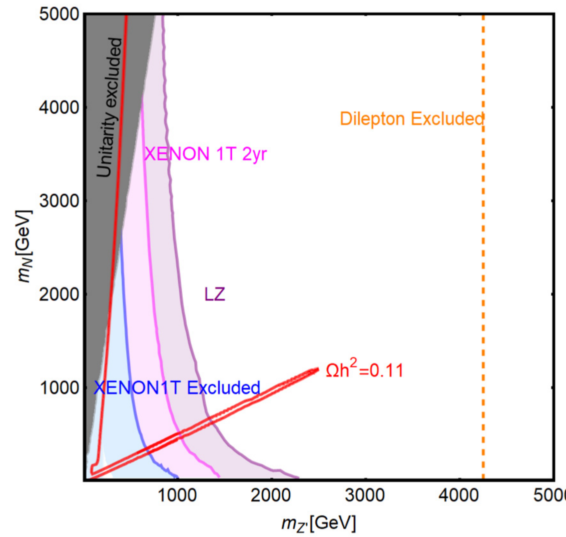

Fig. 3. Inviability of Majorana dark matter: the plot shows how existing dilepton event search limits preclude a viable Majorana dark matter candidate.

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Appendix A. Inviability of Majorana dark matter

In the model discussed thus far the neutral fermions have a Lagrangian term $h^N \bar{l}_L \chi N_R$ where h^N is the relevant Yukawa coupling. After χ develops a nonzero VEV $\langle\chi\rangle = (0, 0, w)/\sqrt{2}$ one obtains three heavy Dirac fermions N, with masses at the large symmetry breaking scale $\langle\chi\rangle$. Notice however, that one can also add a bare mass term $N_R N_R$ proportional to a mass parameter μ. For $\mu \to 0$ we have Dirac fermions, while for $\mu \ll w$ the global symmetry enforcing Diracness is only approximate, and the N_R become quasi-Dirac fermions [107]. On the other hand, for arbitrary $\mu \sim w$ they are generic Majorana fermions.

Such a model would be perfectly consistent except that the dark matter interpretation would no longer be viable. Indeed, if the lightest of the N_a is a Majorana fermion its vectorial coupling with the Z' gauge boson vanishes, affecting its relic density calculation as well its direct detection rates. In Fig. 3 we show the final result of having N_a as Majorana fermions after implementing collider and spin-dependent direct detection limits. One sees that a Majorana dark matter fermion is already excluded in view of the current limits.

This highlights the testability of the model, since the couplings are all fixed by gauge symmetry. Therefore, the bare mass term must be suppressed, making the N_a (mainly) Dirac fermions and restoring the results discussed the main text.

Appendix B. Neutrino seesaw mechanism, leptogenesis and cosmological inflation

Here we note that the neutrinos have Yukawa Lagrangian terms given by

$$\mathcal{L} \supset h^{\nu}_{ab}\bar{l}_{aL}\eta\nu_{bR} + \frac{1}{2} f^{\nu}_{ab}\bar{\nu}^{c}_{aR}\phi\nu_{bR} + H.c. \tag{15}$$

After the scalars develop nonzero vacuum expectation values, $\langle\eta\rangle = (u, 0, 0)/\sqrt{2}$ and $\langle\phi\rangle = \Lambda/\sqrt{2}$, this leads to $m_\nu \simeq -m_D m_R^{-1} m_D^T \sim u^2/\Lambda$, where $m_D = -h^\nu u/\sqrt{2}$. Since $\Lambda \gg u$ the light neutrino masses are naturally small thanks to the canonical type-I seesaw mechanism. On the other hand the heavy right-handed neutrinos ν_R have large Majorana masses $m_R = -f^\nu\Lambda/\sqrt{2}$, at the $U(1)_N$ breaking scale.

Here we note that the $U(1)_N$ breaking that defines R-parity can not only lead to neutrino masses, but also potentially lead to leptogenesis, and cosmological inflation, in certain correlations to dark matter.

The leptogenesis mechanism is governed by CP-asymmetric ν_R decays which, besides the usual mode $\nu_R \to e\eta_2$ include a new mode $\nu_R \to N\eta_3$ to R_P-odd states. This channel transforming into the dark sector is enhanced with respect to the former. We expect there is a link between the fermionic dark matter and the matter-antimatter asymmetry. Note also that the effective potential of ϕ which has ν_R, Higgs triplets, and $U(1)_N$ gauge field contributions, may easily satisfy slow-roll conditions required for cosmic inflation. Inflation would end seemingly due to an instability triggered by ϕ when it reaches a critical value defined by the largest Higgs triplet mass. The inflaton eventually decays into odd scalars, $\phi \to \eta_3\eta_3$, while the channel into fermionic dark matters is loop-induced. This would ensure that fermionic dark matter particles are thermally produced.

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