R-parity as a residual gauge symmetry: probing a theory of cosmological dark matter

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We present a non-supersymmetric scenario in which the R-parity symmetry \( R_P = (-1)^{3(B-L)+2s} \) arises as a result of spontaneous gauge symmetry breaking, leading to a viable Dirac fermion WIMP dark matter candidate. Direct detection in nuclear recoil experiments probes dark matter masses around 2 – 5 TeV for \( M_{\chi} \lesssim 3 – 4 \) TeV consistent with searches at the LHC, while lepton flavor violation rates and flavor changing neutral currents in neutral meson systems lie within reach of upcoming experiments.

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. Even if imposed by hand, R-parity may still break through high dimension operators or spontaneously. While the second case leads to an attractive neutrino mass generation scheme, one loses the WIMP dark matter scenario. Generally, some sort of R-parity symmetry should be invoked in order to stabilize the dark matter candidate. 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. In this work, we discuss a non-supersymmetric model where dark matter stability results from the R-parity symmetry \( R_P = (-1)^{3(B-L)+2s} \), naturally arising as a consequence of the spontaneous breaking of the gauge symmetry. In order to implement this idea we consider an extension of the standard model based upon an extended 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 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. R-parity symmetry \( R_P = (-1)^{3(B-L)+2s} \) 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 \( \mu \rightarrow e\gamma \), as well as high energy collider signatures. We also comment on possible connections to cosmological inflation and leptogenesis.

The model.— Our non-supersymmetric model is based on the 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:

\[
\begin{array}{ccc}
\text{Leptons} & 1\text{-2nd Generations} & 3\text{th Generation} \\
\begin{pmatrix}
\nu_a \\
e_a \\
N_a \\
\end{pmatrix}_L & 
\begin{pmatrix}
d_\alpha \\
-u_\alpha \\
D_\alpha \\
\end{pmatrix}_L & 
\begin{pmatrix}
u_3 \\
d_3 \\
U \\
\end{pmatrix}_L \\
\end{array}
\]

\[
\begin{array}{c}
v_{aR}, e_{aR}, N_{aR}, u_{aR}, d_{aR}, D_{aR}, u_{3R}, d_{3R}, U_R \\
\end{array}
\]

where we have adopted the generation indices \( a = 1, 2, 3 \) and \( \alpha = 1, 2 \).

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

\[
Q = T_3 - \frac{1}{\sqrt{3}} T_8 + X, \quad B - L = -\frac{2}{\sqrt{3}} T_8 + N,
\]

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. 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 I.

\[\text{ Scalars}\]

\[
\begin{pmatrix}
\eta_1^0 \\
\eta_2^0 \\
\eta_3^0 \\
\end{pmatrix} \quad \rho = \begin{pmatrix} \rho_1^+ \\ \rho_2^0 \\ \rho_3^0 \end{pmatrix} \quad \chi = \begin{pmatrix} \chi_1^0 \\ \chi_2^0 \\ \chi_3^0 \end{pmatrix} \quad \phi
\]

where we have adopted the generation indices \( a = 1, 2, 3 \) and \( \alpha = 1, 2 \).

\[\text{ These can however be forbidden with further symmetries.} [2][3].\]
TABLE I. 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_{aL} | q_{aL} | u_{aR} | d_{aR} | U_{aR} | D_{aR} | η | ρ | χ | φ |
|-----------|--------|--------|--------|--------|--------|--------|--------|--------|--------|--------|-----|-----|-----|-----|
| X         | −1/3   | 0      | −1     | 0      | 1/3    | 2/3    | −1/3   | 2/3    | −1/3   | 2/3    | −1/3| 0   |     |     |
| N         | −2/3   | −1     | −1     | 0      | 2/3    | 1/3    | 4/3    | −2/3   | 1/3    | 1/3    | −2/3| 2   |     |     |

Gauge symmetry breaking by these SU(3) Higgs triplets and singlet addresses the origin of R-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 R-parity given as 

\[ R_P = (−1)^{B−L+2s} = (−1)^{−2\sqrt{3}f_b+3N+2s} \]

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 R-parity, since SU(3)g and B − L symmetries neither commute nor close algebraically. We also note that the exotic fermions have the following B − L quantum numbers, [B − L](N_{aR}, D_{aR}) = 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 \( w = \langle \chi^0_3 \rangle \), the VEV of the \( \chi^0_3 \) 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 \( h_{ab}^\nu w \) proportional to \( w = \langle \chi^0_3 \rangle \). For all cases, N_{a} can be made the lightest odd particle under R-parity, and therefore it is a Dirac dark matter candidate (see Appendix A for alternative assumptions). In what follows, we will investigate the phenomenological consequences of our model. We start by addressing electroweak limits.

**CKM unitarity.** — Quantum loop corrections to the quark mixing matrix resulting from additional neutral gauge bosons induce deviations from unitarity of the CKM matrix \[2\]. 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 \( \Delta_{\text{CKM}} = 1 - \sum_{q=d,s,b} |V_{u,q}|^2 \). Applying this to the neutral current we find,

\[ \Delta_{\text{CKM}} = -0.0033 \frac{M_{W}^{2}}{M_{Z'}^{2}} \ln \left( \frac{M_{W}^{2}}{M_{Z'}^{2}} \right) \]

which implies into \( M_{Z'} > 200 \text{ GeV} \).

**TABLE II.** The couplings of Z’ with fermions.

| Multiplet | \( f \) | \( g_f \) | \( g_f' \) |
|-----------|-------|--------|--------|
| \( \nu_{aL} \) | \( \frac{c_{2W}}{2\sqrt{3}+4t_{W}} \) | \( c_{2W} \) | \( 2\sqrt{3}+4t_{W} \) |
| \( e_{aL} \) | \( \frac{c_{3W}}{2\sqrt{3}+4t_{W}} \) | \( 1 \) | \( 2\sqrt{3}+4t_{W} \) |
| \( N_{aR} \) | \( \frac{c_{3W}}{\sqrt{3}+4t_{W}} \) | \( c_{2W} \) | \( c_{2W} \) |
| \( u_{aR} \) | \( \frac{c_{3W}}{3+2\sqrt{3}+4t_{W}} \) | \( 1 \) | \( 3+2\sqrt{3}+4t_{W} \) |
| \( d_{aR} \) | \( \frac{c_{3W}}{\sqrt{3}+4t_{W}} \) | \( 1 \) | \( 2\sqrt{3}+4t_{W} \) |

**Electroweak precision tests.** — New physics contributions to the \( \rho \)-parameter come from the mixing among the neutral gauge bosons, which is evaluated by,

\[ \Delta \rho = \frac{M_{W}^{2}}{c_{W}^{2}M_{Z'}^{2}} - 1 \approx \left( \frac{c_{2W}w^{2} - v^{2}}{4c_{W}^{2}(u^{2} + v^{2})w^{2}} + \frac{t_{W}^{4}(u^{2} + v^{2})}{36\Lambda^{2}} \right) \]

where \( u, v, w, \) and \( \Lambda \) are the VEVs of \( \eta_{1}, \rho_{2}, \chi_{3}, \) and \( \phi \), 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 \( \Delta \rho < 0.0006 \) \[21\], we find the bounds summarized in Table II, taking into account that \( u^{2} + v^{2} = v_{SM}^{2}, v_{SM} = 246 \text{ GeV}, s_{W}^{2} = 0.231, \alpha = 1/128, \) and \( g^{2} = 4\pi\alpha/s_{W}^{2} \). As we shall see these limits are all surpassed by LHC probes. We have checked that one-loop new physics corrections to the \( \rho \)-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 \( e^{+}e^{-} \rightarrow Z' \rightarrow f \bar{f} \)

\footnotetext{2 Here we neglect unitarity violation from \( \nu_{aR} \) admixture \[17\] \[19\].}
production channel with $Z'$ being off-shell. The bound can be obtained using the parametrization [22],
\[
\mathcal{L} = \frac{g^2}{c_W M_{Z'}} [\bar{e} \gamma^\mu (a_L^f P_L + a_R^f P_R) e] [\bar{f} \gamma_\mu (a_L^f P_L + a_R^f P_R) f]
\]
where $a_L^f = (g_L^f + g_1^f)/2$ and $a_R^f = (g_L^f - g_1^f)/2$.

Using the LEP-II results for final state dileptons we get [23],
\[
\frac{g^2}{\cos^2(\theta_W)} \left[ \frac{\cos(2\theta_W)}{2 \sqrt{3 - 4 \sin^2(\theta_W)}} \right] \frac{1}{M_{Z'}^2} < \frac{1}{(6 \text{ TeV})^2}
\]
which translates into $M_{Z'} > 1.93$ TeV.

**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 [24]. This translates into $\Delta a_\mu = (287 \pm 80) \times 10^{-11}$ (see [25] for an extensive and recent review). 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 derive limits, by requiring their contribution to lie within the error bars. Using the current (projected) sensitivities we find, $M_{Z'} > 180$ GeV (273 GeV), $M_{W'} > 100$ GeV (145 GeV).

**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. Similar discussion holds for the $B_0^0 - \bar{B}_0^0$ meson system. Defining $V_{\text{CKM}} = U_L^T V_L$, 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,
\[
K^0 - \bar{K}^0 : \frac{g^2}{(3 - i W_Z) M_{Z'}} \left[ (V_{31}^*) (V_{32}) \right]^2 < \frac{1}{(10^4 \text{ TeV})^2},
\]
\[
B_0^0 - \bar{B}_0^0 : \frac{g^2}{(3 - i W_Z) M_{Z'}} \left[ (V_{32}^*) (V_{33}) \right]^2 < \frac{1}{(100 \text{ TeV})^2}.
\]

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 [26] 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 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 $B_0^0 - \bar{B}_0^0$ 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 important uncertainties. In this work we will adopt the conservative bound, but bear in mind that more stringent limit may be applicable.

**Dilepton Resonance Searches at the LHC.**— Dilepton resonance searches are the gold channel for heavy neutral gauge bosons with un-suppressed couplings to leptons [27]. 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 [28, 29]. Using 13 TeV center-of-energy and 3.2 fb$^{-1}$ of integrated luminosity the ATLAS collaboration has placed restrictive limits on the mass of gauge bosons arising in some new physics models [30]. Using the just released LHC data with integrated luminosity of $\mathcal{L} = 13.3$ fb$^{-1}$, and applying the cuts,

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

with $M_{ll}$ denoting the dilepton invariant mass, one can find a bound on the $Z'$ mass [31]. We have generated events with MadGraph5 [32, 33], adopting the CTEQ6L parton distribution function [34] and efficiencies/acceptances

\[\text{See [29] and [31, 33] for previous studies.}\]
as described in [30]. The resulting limit was found to be $M_{Z'} > 3.8$ TeV. Keeping a similar detector response we expect that upcoming LHC runs with $L = 100(1000)$ fb$^{-1}$ will probe $M_{Z'} = 4.9 (6.1)$ TeV, respectively.

**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 [37, 38]. Here one can use the lepton plus missing energy data, via the $pp \to W' \to l\nu$ production channel at the LHC with $L = 13.3 fb^{-1}$ and 13 TeV center of mass energy. No significant excess above Standard Model predictions was seen, leading to $M_{W'} > 4.74$ TeV [38]. In this model, the charged current contains,

$$\mathcal{L} \supset - \frac{g}{\sqrt{2}} (\epsilon_{\alpha L} \gamma^\mu N_{\alpha L}) W'_\mu. \quad (8)$$

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 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 $L = 100(1000)$ fb$^{-1}$ which turn out to be sensitive to $M_{W'} = 5.8 (7)$ TeV.

![Fig. 1](image_url)  
**Fig. 1.** Region of parameters yielding $4.2 \times 10^{-13} < \text{Br}(\mu \to e\gamma) < 4 \times 10^{-14}$ 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 $l+\text{MET}$ data from LHC in gray. The upper blue line in the region represents the current limit $\text{Br}(\mu \to e\gamma) < 4.2 \times 10^{-13}$.

**Lepton flavor violation.**— In the Standard Model lepton flavor is conserved and neutrinos are massless. However, neutrinos experience flavor oscillations [39, 41] 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 [42]. 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 [29].

$$\text{Br}(\mu \to e\gamma) = 6.43 \times 10^{-6} \left( \frac{1 \text{TeV}}{M_{W'}} \right)^4 \sum_f (g_{f e}^\ell g_{f \mu}^\ell)^2,$$

with $g_{f e}^\ell = g U_{N e}/(2\sqrt{2})$ and $g_{f \mu}^\ell = g U_{N \mu}/(2\sqrt{2})$.

Current (projected) sensitivity as reported by the MEG collaboration [43] implies that $\text{Br}(\mu \to e\gamma) < 4.2 \times 10^{-13}$ ($4 \times 10^{-14}$). Thus one can translate this bound into a limit on the product $U_{N e}U_{N \mu}$ as function of the $W'$ mass as shown in Fig. 1. Here 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 $U_{N e}U_{N \mu}$ the $\mu \to e\gamma$ search may outperform collider probes. We now investigate the feasibility of this model concerning dark matter searches.

**WIMP Dirac dark matter.**— In our model, one can have either a Dirac or Majorana fermionic dark matter [44, 50], 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). 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. The $W'$ boson also mediates t-channel interactions, which are nevertheless subdominant, and thus neglected in our computations. In Fig. 2 we exhibit the result of the dark matter phenomenology, encompassing the relic density using PLANCK data [61], dark matter-nucleon scattering limits from searches by the LUX collaboration [52], expected XENON1T [53] and LZ [54] sensitivities, as well as current and future prospects at the LHC for searches for neutral gauge bosons. 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 $Z'$ masses up to 10 TeV outperforming the LHC. In conclusion we have shown that this UV complete dark matter model addresses the origin of R-parity from first principles and offers a viable dark matter candidate for $M_{Z'} > 4$ TeV and dark matter masses of 2 – 5 TeV.

**FIG. 2.** Region of parameters that yields the right relic density curve in red. The existing limits from the non-observation of dark matter nuclear scattering by the LUX collaboration are indicated in light blue [53]. The prospects for the XENON1T experiment with 2-years exposure [53], as well as the projected sensitivity of the LZ dark matter experiment are also indicated [54]. 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.

**Conclusion.**— Is summary, we have presented a non-supersymmetric SU(3)$_c$ ⊗ SU(3)$_b$ ⊗ U(1)$_X$ ⊗ U(1)$_N$ model in which a conserved R-parity symmetry $R_P = (-1)^{3(B-L)+2s}$ arises as a residual unbroken discrete gauge symmetry. The fact that the $B-L$ symmetry is preserved at high scales plays a key role in accounting for the origin of R-parity conservation. The lightest $R_P$-odd particle constitutes a viable WIMP 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 $\mu \rightarrow e\gamma$, as well as dilepton event searches at the LHC.

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**Appendix A: Inviability of Majorana dark matter**

In the model discussed thus far the neutral fermions have a Lagrangian term $h N \bar{L} \chi N_R$ where $h^N$ is the relevant Yukawa coupling. After $\chi$ 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 $\mu$. For $\mu \rightarrow 0$ 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 [55]. 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.

**FIG. 3.** Inviability of Majorana dark matter: the plot shows how existing dilepton event search limits preclude a viable Majorana dark matter candidate.
Appendix B: Neutrino seesaw mechanism, leptogenesis and cosmological inflation

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

$$\mathcal{L} \supset h_{ab}^T \eta a_L \nu_b R + \frac{1}{2} f^{\nu} \nu a_R \phi \nu b_R + H.c.$$  \hspace{1cm} (10)

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 $m_{\nu} \approx -m_D^2 m_D^2 / a_2^2 \approx u^2 / \Lambda$. Here we note that the neutrinos have Yukawa Lagrangian terms given by [5] A. Masiero and J. W. F. Valle, Phys. Lett. B251, 273 (1990).

$$\mathcal{L} \supset h_{ab}^T \eta a_L \nu_b R + \frac{1}{2} f^{\nu} \nu a_R \phi \nu b_R + H.c.$$  \hspace{1cm} (10)

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 $m_{\nu} \approx -m_D^2 m_D^2 / a_2^2 \approx u^2 / \Lambda$. In contrast 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 $\phi$ has $\nu_R$, Higgs triplets, and U(1)$_N$ gauge field contributions, may easily satisfy slow-roll conditions required for cosmic inflation. Inflation would ensue accordingly due to an instability triggered by $\phi$ when it reaches a critical value defined by the largest Higgs triplet mass. The inflation eventually decays into odd scalars, $\phi \rightarrow \eta_R$, while the channel into fermionic dark matters is loop-induced. This would ensure that fermionic dark matter particles are thermally produced.

4 For the systematic seesaw expansion formalism see [50].

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