

Right-handed neutrino dark matter in left-right symmetric models

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Abstract. We show that in a class of non-supersymmetric left-right extensions of the Standard Model (SM), the lightest right-handed neutrino (RHN) is naturally stable and can therefore play the role of thermal Dark Matter (DM) in the Universe for a wide mass range from TeV to PeV. Our model is based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{Y_L} \times U(1)_{Y_R}$ in which a heavy copy of the SM fermions are introduced and the stability of the RHN DM is guaranteed by an automatic Z_2 symmetry present in the leptonic sector. The active neutrino masses in the model arise from the type-II seesaw mechanism. We find a lower bound on the RHN DM mass of order TeV from relic density constraints, as well as an unitarity upper bound in the multi-TeV to PeV scale, depending on the entropy dilution factor. The RHN DM could be made long-lived by soft-breaking of the Z_2 symmetry and provides a concrete example of decaying DM interpretation of the PeV neutrinos observed at IceCube.

1. Introduction

The existence of dark matter (DM) in our Universe has been well established by astrophysical and cosmological observations. To explain this, one needs some new particle(s) and symmetries beyond the Standard Model (SM). Being colorless and electrically neutral, SM-singlet right-handed neutrino (RHN), is a natural DM candidate. Although the right handed neutrino is introduced to implement the seesaw mechanism to understand small neutrino masses, in the conventional seesaw models [1, 2, 3, 4, 5], it has Yukawa couplings to SM lepton and Higgs doublets, making it highly unstable, and therefore, unsuitable for DM. The only exception is when its mass is in the keV range and it has appropriately small Yukawa couplings [6]. In ultraviolet (UV) complete scenarios of seesaw, e.g. those based on the Left-Right (LR) gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [7, 8, 9], even having tiny Yukawa couplings do not help since the gauge interactions of RHNs induce further decay modes, which have to be suppressed/forbidden by imposing additional discrete symmetries to keep the lightest RHN cosmologically stable [10].

In this paper we discuss an alternative class of LR models based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{Y_L} \times U(1)_{Y_R}$ [11] which has an accidental Z_2 symmetry, at the renormalizable level that keeps the lightest RHN stable making it a natural DM candidate.



In this model there exists a right-handed copy of the SM electroweak sector at the TeV scale (or higher), charged under $SU(2)_R \times U(1)_{Y_R}$. The color $SU(3)$ is common to both sectors. The breaking of the weak gauge symmetry down to $U(1)_{\text{em}}$ is similar to the LR models with universal seesaw [13, 14, 15]. But one of the key differences from the universal seesaw models is that the right-handed lepton sector in our model is prevented by gauge symmetry from having any masses connecting the heavy sector to the light ones. Because of this, there appears a remnant automatic Z_2 symmetry in the leptonic sector at the renormalizable level that keeps the lightest RHN stable. If the stabilizing Z_2 symmetry is softly broken by a mass term connecting the heavy and light sector, the RHN DM can be made very long-lived but allowing for its decay into SM leptons. This raises the possibility that for a PeV-scale DM, one can explain the ultra-high energy neutrino events observed at IceCube [16]. We also point out a dilution mechanism which allows us to relax the unitarity bound on the dark matter to the PeV mass level so that it can be applied to the discussion of PeV neutrinos in the context of a thermal annihilating DM.

2. The Model

The model is based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{Y_L} \times U(1)_{Y_R}$ [17, 18]. For each generation of SM quark and lepton doublets Q_L, ψ_L and singlets u_R, d_R, e_R , we have the corresponding heavy fermion partners $\mathcal{Q}_R, \Psi_R$ and $\mathcal{U}_L, \mathcal{D}_L, \mathcal{E}_L$, all of which are collected in the upper rows of Table 1. Note that the heavy doublets transform nontrivially under the $SU(2)_R$ gauge symmetry, required by the parity $L \leftrightarrow R$. Analogous to other LR models, the electric charge formula is generalized to

$$Q = I_{3L} + I_{3R} + \frac{1}{2}(Y_L + Y_R). \quad (1)$$

The Higgs sector of the model consists of $SU(2)_{L,R}$ doublets $\chi_{L,R}$ and triplets $\Delta_{L,R}$, which are also presented in Table 1. In the limit of vanishing mixing of the scalars, the doublet χ_L plays the role of the SM Higgs doublet. The gauge symmetries $SU(2)_L$ and $SU(2)_R$ are broken by the vacuum expectation values (VEVs) of the neutral components of these Higgs fields:

$$\langle \chi_L^0 \rangle = v_L \equiv v_{\text{EW}}, \quad \langle \chi_R^0 \rangle = v_R, \quad \langle \Delta_L^0 \rangle = w_L, \quad \langle \Delta_R^0 \rangle = w_R, \quad (2)$$

which gives also masses to the heavy light gauge bosons. The SM and heavy charged fermion masses are obtained through their Yukawa couplings to the Higgs doublets:

$$-\mathcal{L}_Y \supset y_u \bar{Q}_L \tilde{\chi}_L u_R + y_d \bar{Q}_L \chi_L d_R + y_e \bar{\psi}_L \chi_L e_R \\ + y'_u \bar{\mathcal{Q}}_R \tilde{\chi}_R \mathcal{U}_L + y'_d \bar{\mathcal{Q}}_R \chi_R \mathcal{D}_L + y'_e \bar{\Psi}_R \chi_R \mathcal{E}_L + \text{H.c.}, \quad (3)$$

where $\tilde{\chi}_L = i\sigma_2 \chi_L^*$ and similarly for χ_R (σ_2 being the second Pauli matrix). Exact parity symmetry implies that the Yukawa couplings $y_f = y'_f$ (and also $f = f'$ in Eq. (4) below). If the $SU(2)_R$ symmetry breaking is at the TeV scale, there would be new quarks in the few GeV range, which is inconsistent with the current LHC bounds on vector-like fermion masses of order of TeV. To have heavy quarks and leptons at the TeV scale, we adopt the parity violating scenario in which the couplings y_f and y'_f are independent and choose $y'_f \sim \mathcal{O}(1)$.¹ It is worth noting that at this stage the lower bound on M_{W_R} from low-energy flavor changing effects such as $K^0 - \bar{K}^0$ mixing does not apply to our model.

¹ The parity symmetric version could be realized at the TeV scale by doubling the doublet scalars in both the SM and heavy sectors, such that one set of the Yukawa couplings $y_{f,1}$ is responsible for the SM fermion masses and the other one $y_{f,2}$ dominates the heavy fermion masses [11].

Table 1. Fermion and scalar contents of the LR model in this paper and their representations under the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{Y_L} \times U(1)_{Y_R}$ [11].

field	representation	field	representation
$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$(\mathbf{3}, \mathbf{2}, \mathbf{1}, \frac{1}{3}, 0)$	$Q_R = \begin{pmatrix} \mathcal{U}_R \\ \mathcal{D}_R \end{pmatrix}$	$(\mathbf{3}, \mathbf{1}, \mathbf{2}, 0, \frac{1}{3})$
u_R	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, \frac{4}{3}, 0)$	$\mathcal{U}_L$	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, 0, \frac{4}{3})$
d_R	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, -\frac{2}{3}, 0)$	$\mathcal{D}_L$	$(\mathbf{3}, \mathbf{1}, \mathbf{1}, 0, -\frac{2}{3})$
$\psi_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	$(\mathbf{1}, \mathbf{2}, \mathbf{1}, -1, 0)$	$\Psi_R = \begin{pmatrix} N \\ \mathcal{E}_R \end{pmatrix}$	$(\mathbf{1}, \mathbf{1}, \mathbf{2}, 0, -1)$
e_R	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, -2, 0)$	$\mathcal{E}_L$	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, 0, -2)$
$\chi_L = \begin{pmatrix} \chi_L^+ \\ \chi_L^0 \end{pmatrix}$	$(\mathbf{1}, \mathbf{2}, \mathbf{1}, 1, 0)$	$\chi_R = \begin{pmatrix} \chi_R^+ \\ \chi_R^0 \end{pmatrix}$	$(\mathbf{1}, \mathbf{1}, \mathbf{2}, 0, 1)$
$\Delta_L = \begin{pmatrix} \Delta_L^+/\sqrt{2} & \Delta_L^{++} \\ \Delta_L^0 & -\Delta_L^+/\sqrt{2} \end{pmatrix}$	$(\mathbf{1}, \mathbf{3}, \mathbf{1}, 2, 0)$	$\Delta_R = \begin{pmatrix} \Delta_R^+/\sqrt{2} & \Delta_R^{++} \\ \Delta_R^0 & -\Delta_R^+/\sqrt{2} \end{pmatrix}$	$(\mathbf{1}, \mathbf{1}, \mathbf{3}, 0, 2)$
Σ_1	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, -\frac{4}{3}, \frac{4}{3})$	Σ_2	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, \frac{2}{3}, -\frac{2}{3})$

The masses of light and heavy neutrinos are generated from couplings to the triplet scalars $\Delta_{L,R}$ via the type-II seesaw mechanism [19, 20, 21, 22]:

$$-\mathcal{L}_Y \supset f \bar{\psi}_L^C i \sigma_2 \Delta_L \psi_L + f' \bar{\Psi}_R^C i \sigma_2 \Delta_R \Psi_R + \text{H.c.}, \quad (4)$$

which implies that the VEV w_L should be at the eV scale and w_R is at the TeV scale or above. This can be easily obtained by choosing appropriate values of the parameters in the scalar potential.

3. Depleting the heavy quarks

After symmetry breaking at the right-handed scale and the electroweak scale, the model has a large global symmetry: $U(1)_{B,L} \times U(1)_{B,R} \times Z_{2\ell,L} \times Z_{2\ell,R}$, with the two $U(1)_B$ the baryon numbers in the SM and heavy sectors, and the two Z_2 's the corresponding lepton numbers.

Protected by the baryon number symmetry, the lightest hadron in the $SU(2)_R$ sector is absolutely stable, which can be one of the heavy baryons of form QQQ , QQq , Qqq , or heavy mesons $\bar{Q}Q$ and $\bar{Q}q$. For the phenomenological purpose of avoiding the heavy lightest baryon also becoming DM and affecting the evolution of the Universe, we add two more Higgs singlets $\Sigma_1(\mathbf{1}, \mathbf{1}, -\frac{4}{3}, +\frac{4}{3})$ and $\Sigma_2(\mathbf{1}, \mathbf{1}, +\frac{2}{3}, -\frac{2}{3})$ which connect the heavy and light singlet quarks via the terms

$$\mathcal{L}_{\text{mix}} = \lambda_U \bar{\mathcal{U}}_L u_R \Sigma_1^* + \lambda_D \bar{\mathcal{D}}_L d_R \Sigma_2^* + \text{H.c.}. \quad (5)$$

Once the Σ fields acquire VEVs, the heavy-light quark mass mixing terms $\delta_U = \lambda_U \langle \Sigma_1 \rangle$ and $\delta_D = \lambda_D \langle \Sigma_2 \rangle$ appear in the Lagrangian, and the heavy quark can then decay into the SM fermions, leaving only the lightest RHN stable due to the leptonic Z_2 symmetry. The singlets $\Sigma_{1,2}$ are charged under $U(1)_{Y_L} \times U(1)_{Y_R}$ and break this symmetry down to $U(1)_Y$. At this stage this model is similar to the usual LR models with $U(1)_{B-L}$, however with the quantum number $\mathcal{Y}$ different from $B-L$ and also from the SM hypercharge Y_L (which is a linear combination of $\mathcal{Y}$ and the $U(1)$ subgroup of $SU(2)_R$ after symmetry breaking).

The decay rate of heavy quarks $Q \rightarrow q + Z$ can be estimated as

$$\Gamma_Q \simeq \frac{g_L^2}{64\pi \cos^2 \theta_w} \frac{\delta_Q^2}{v_R^2} \frac{M_Q^3}{M_Z^2}, \quad (6)$$

where θ_w is the SM weak mixing angle and g_L is the gauge coupling of $SU(2)_L$. Requiring that the heavy quarks are depleted before the QCD phase transition epoch $T_{\text{QCD}} \approx 200$ MeV imposes a lower bound on the magnitude of $\delta_Q \gtrsim 10^{-6}$ GeV [11]. When this limit is close to being saturated, we expect the heavy quark decay length to be several cm, leading to displaced vertex signatures at hadron colliders.

4. Connecting the SM to DM

The heavy gauge boson Z_R is a linear combination of the $U(1)$ subgroup of $SU(2)_R$, $U(1)_{Y_L}$ and $U(1)_{Y_R}$, thus both the SM and heavy fermions are charged under Z_R . Even in the absence of any scalar and fermion mixings connecting the heavy and light sectors, the RHN N could interact pairwise with the SM particles through the heavy Z_R boson. Up to the leading order of mixings of $h - \Delta_R^0$ and $Z - Z_R$, two N 's could also annihilate to the SM fields via the scalar portal (h and Δ_R^0) and gauge portal (Z boson). The SM Higgs h is assumed to be predominantly from the doublet χ_L , and thus the scalar mixing is expected to be of order $\zeta_S \simeq \lambda v_{\text{EW}}/w_R$ from the quartic term of form $\lambda(\chi_L^\dagger \chi_L)\text{Tr}(\Delta_R^\dagger \Delta_R)$, where for simplicity we have neglected the mixing with other scalars. The neutral gauge boson mixing is of order $\zeta_Z \simeq \xi_Z(M_Z/M_{Z_R})^2$, where $\zeta_Z = \sin\theta_w(\cot^2\theta_w g_R^2/g_L^2 - 1)^{-1/2}$. Note that when the gauge coupling g_R approaches the theoretical lower limit of $g_L \tan\theta_w$ [23], which is independent of the symmetry breaking pattern [24], the $Z - Z_R$ mixing could be significantly enhanced.

In the minimal version of our model, we do not have the bi-multiplet scalars which transform non-trivially under both $SU(2)_L$ and $SU(2)_R$, thus the charged gauge bosons W and W_R can not mix at the tree level. A tiny mixing arises at one-loop level, induced by the (3rd generation) heavy-light quark mixings, which is estimated to be very small, of order $\zeta_W \sim 10^{-13}$, given the heavy quarks are of TeV mass, the W_R boson with mass of 5 TeV and the parameters $\delta_{U,D} \simeq \text{GeV}$ [11].

5. DM annihilation and relic density

In the early universe, all the particles are in equilibrium due to the common $SU(3)_c$ color and $U(1)_Y$ interactions. As the universe cools down below the DM mass M_N , the DM density goes down and freezes out, as in case of a generic thermal DM candidate. The primary annihilation channels to the SM particles proceed via the scalar and the neutral gauge boson portals:

$$NN \rightarrow \begin{cases} h^{(*)}/\Delta_R^{0(*)} \\ Z^{(*)}/Z_R^{(*)} \end{cases} \rightarrow \text{SM particles}. \quad (7)$$

The h , Δ_R^0 and Z portals are suppressed respectively by the small mixing angles ζ_S and ζ_Z . On the other hand, though the couplings of Z_R to N and the SM fermions are of order one, the Z_R portal is suppressed by the large Z_R mass, except for the Breit-Wigner resonance when $2M_N \simeq M_{Z_R}$. Note that the Δ_R^0 portal is further suppressed by the DM mass when DM is light, as the Yukawa coupling f' in Eq. (4) is proportional to the DM mass via $M_N = 2f'w_R$. Combining all the available channels, we calculate the thermally averaged DM annihilation cross section times velocity $\langle\sigma v\rangle = a + b\langle v^2\rangle + \mathcal{O}(v^4)$, from which we can obtain the current DM relic density as for a standard thermal WIMP particle.

An example is given in Fig. 1, where we set explicitly the quartic coupling $\lambda = 1$ and the heavy scalar mass $M_{\Delta_R^0} = 2$ TeV with width 30 GeV. Three different value of the gauge coupling g_R are chosen with $g_R/g_L = 0.6, 1$ and 1.5 . The corresponding Z_R mass and the RH scale $v_R = w_R$ are set to their current experimental lower bounds [25, 26], viz. $M_{Z_R} = 4.6, 3.9, 3.5$ TeV and $w_R = 3.0, 3.2, 2.1$ TeV for $g_R/g_L = 0.6, 1, 1.5$ respectively. The horizontal dashed line shows the observed relic density, as measured by Planck [27]. The various peaks in Fig. 1 are respectively

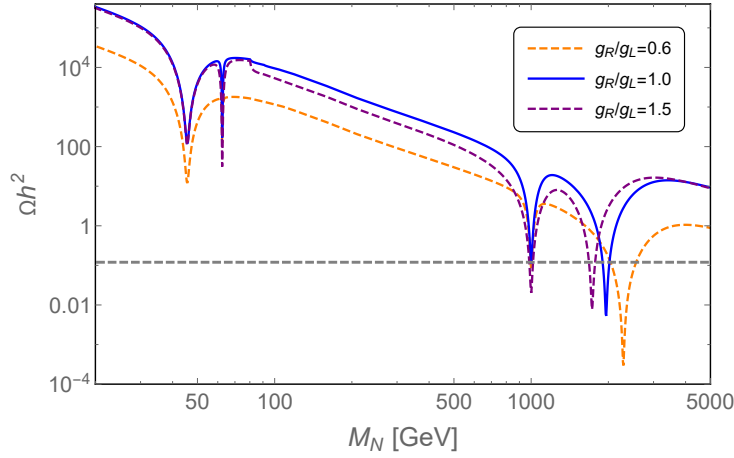


Figure 1. An illustration of the relic density of RHN DM as a function of its mass M_N for different values of g_R/g_L . The horizontal line gives the observed value from Planck data [27].

(from left to right) due to the SM Z and Higgs bosons, Δ_R^0 and Z_R . We find that a TeV scale RHN could accommodate the observed DM relic density, with the annihilation dominated by the heavy scalar and/or Z_R bosons, depending largely on the quartic coupling λ , the gauge coupling g_R and the VEV v_R .

6. Going beyond the unitarity bound

In generic annihilating thermal DM models, the DM mass cannot be arbitrarily large due to the well-known partial wave unitarity limit [28]. For the resonance $R(= \Delta_R^0, Z_R)$ just above the threshold, the annihilation cross section times velocity can be approximated by

$$\langle\sigma v\rangle \simeq 16\pi \frac{\Gamma_{NN}\Gamma_{\text{SM}}}{(4M_N^2 - M_R^2)^2 + \Gamma_R^2 M_R^2} \sim \frac{4\pi}{M_N^2} (B_{NN}B_{\text{SM}}), \quad (8)$$

where Γ_{NN} and Γ_{SM} are the partial decay widths of R to the DM pair and SM particles respectively, and B_{NN} and B_{SM} are the corresponding branching ratios. From Eq. (8), we find that the annihilation rate decreases with increasing DM mass, which leads to the unitarity bound of ~ 20 TeV in this model.

It turns out that there is a built-in mechanism for entropy dilution in our model, which makes it possible to go beyond the unitarity limit. The basic idea is that the next-to-lightest heavy neutrino (denoted here by N_2) does not need to be stable and can decay into SM leptons, e.g. $N_2 \rightarrow \ell\ell\nu$ via the charged lepton mixing terms in Eq. (10) involving only the heavier-than- N leptons in the $SU(2)_R$ sector. Note that the hadronic channels $N_2 \rightarrow \ell q\bar{q}$ are comparatively suppressed by the small quark mixing $\delta_{U,D}$, and the two-body decay $N_2 \rightarrow W^\pm \ell^\mp$ is heavily suppressed by the loop-induced $W - W_R$ mixing. Because of higher mass, N_2 freezes out before the DM N but could decay after N freezes out, i.e. $T_{N_2} < T_F$, with appropriate values of leptonic mixing. The decay of N_2 to relativistic SM species will generate entropy which dilutes the relic density of N , and therefore, the initial relic density of N can be significantly larger than the current value, thus allowing for larger M_N values. The dilution factor depends largely on the mixing parameters in the scalar and lepton sector and some of the Yukawa couplings, which could reach up to the factor of 10^6 to accommodate a PeV DM N to explain the IceCube data. More details can be found in [11, 12].

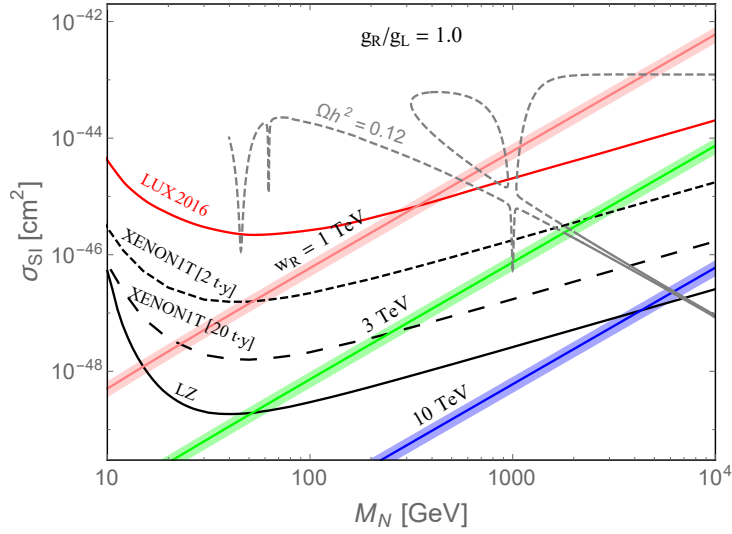


Figure 2. Predictions for the SI DM-nucleon scattering cross sections as functions of DM mass M_N with the gauge coupling $g_R = g_L$ and the RH scale $w_R = 1, 3$ and 10 TeV. The colorful bands are due to the sizable uncertainties of effective DM-nucleon couplings $f_{p,n}$. We show also the current upper limits from LUX [30], as well as the future reaches of XENON1T [32] and LZ [33]. The gray curves correspond to the parameter space producing the observed relic density $\Omega h^2 = 0.12$ (with the dashed part excluded by direct collider searches of Z_R in the dilepton channel).

7. Direct and indirect detections and collider searches

As a Majorana DM candidate, the RHN has both spin-independent (SI) and spin-dependent (SD) interactions with nuclei, which are mediated by the scalars and neutral gauge bosons, respectively.

The SI scattering cross section is dominated by the SM Higgs contribution, with the heavy scalar portal suppressed by its large mass,

$$\sigma_{\text{SI}} = \frac{\lambda^2 \mu^2 M_N^2}{2\pi m_h^4 w_R^4} [Z f_p + (A - Z) f_n]^2 \quad (9)$$

where μ is the reduced DM-nucleus mass, and $f_{p,n}$ the effective Higgs-proton/neutron couplings. A characteristic feature of our model is that $\sigma_{\text{SI}} \propto M_N^2$, which originates from the dependence of Yukawa coupling on the DM mass $f' \propto M_N$. As shown in Fig. 2, when the right-handed scale $v_R = w_R$ goes higher, the scalar mixing decreases and the scattering cross section σ_{SI} also gets smaller. In this plot we adopt the benchmark values of $v_R = w_R = 1, 3$ and 10 TeV, the gauge coupling $g_R = g_L$ and the quartic coupling $\lambda = 1$. The colorful bands are due to the sizable uncertainties of effective couplings $f_{p,n}$. The most stringent current direct detection limits come from LUX [30] (with comparable limits from PandaX-II [31]), which rule out the lighter RHN DM mass range, if the DM relic density is set to be the observed value (the gray contours). When the collider limits on Z_R mass are also taken into consideration (dashed gray lines), we can set more stringent lower limits of a TeV or so on the DM mass. We show also the projected limits from XENON1T with exposure values of 2 and 20 ton-year [32] and the LZ experiment [33] with orders of magnitude improvement in the SI scattering cross section limits, which can constrain the DM mass M_N and right-handed scale w_R up to the 10 TeV scale in near future experiments.

As for the SD scattering, it is mediated at leading order by the Z and Z_R bosons, both of which are suppressed by the small $Z - Z_R$ mixing angle ζ_Z . When the collider constraints on M_{Z_R} are taken into consideration, the $Z - Z_R$ mixing is so small that even with the future LZ limits it is rather challenging to test the model in the SD channels.

In the early universe, the RHN DM annihilates into the SM particles, e.g. W^+W^- , $\ell^+\ell^-$ and $q\bar{q}$, which could emit energetic gamma-rays. Thus the RHN DM can also be probed by the Fermi-LAT gamma-ray observations. Also DM annihilation injects energy into the thermal bath in the early universe, thus the annihilation cross section could also be constrained by the Planck data of the cosmic microwave background (CMB). However, limited by the small scalar and gauge mixing angles (ζ_S and ζ_Z), these indirect detections do not put additional constraints on the currently allowed model parameter space with the observed relic density.

As in usual thermal DM models, the RHN DM can be pair-produced at high energy hadron colliders, with monojet, mono- h , mono- W/Z and mono-photon signatures. However, suppressed again by the small mixing parameters, none of these channels could effectively probe our model at the LHC. It might be helpful to check some specific parameter regions at future higher energy colliders, such as FCC-hh and SPPC.

8. Decaying RHN DM and IceCube PeV neutrinos

In the model discussed in Section 2, the SM charged leptons e_R do not mix with their heavy partners $\mathcal{E}_L$, thus the lightest RHN is stabilized by the automatic Z_2 symmetry. If we add the following terms to the Lagrangian

$$\mathcal{L}_{\text{mix}} = \delta_\ell \bar{\mathcal{E}}_L e_R + \text{H.c.}, \quad (10)$$

to explicitly break the Z_2 symmetry (albeit softly), then the lightest RHN can decay into SM lepton and W boson through the loop mediated $W - W_R$ mixing, i.e. $N \rightarrow W^\pm \ell^\mp$. This raises the interesting possibility of explaining the IceCube PeV neutrinos [34] when we set $M_N \sim \text{few PeV}$ [16]. However, the two-body decays are disfavored by the IceCube data, as the neutrino spectrum is flatter than required [16].

To fit the IceCube neutrino data, we promote the scalars to be lepton-specific two Higgs doublet model (2HDM) in both the SM and heavy sectors, with the Yukawa couplings given by

$$\begin{aligned} -\mathcal{L}_Y \supset & y_u \bar{Q}_L \tilde{\chi}_{L,q} u_R + y_d \bar{Q}_L \chi_{L,q} d_R + y_\ell \bar{\psi}_L \chi_{L,\ell} e_R \\ & + y'_u \bar{Q}_R \tilde{\chi}_{R,q} \mathcal{U}_L + y'_d \bar{Q}_R \chi_{R,q} \mathcal{D}_L + y'_\ell \bar{\Psi}_R \chi_{R,\ell} \mathcal{E}_L + \text{H.c.}, \end{aligned} \quad (11)$$

where the electroweak VEV $v_{\text{EW}} = \sqrt{v_{L,\ell}^2 + v_{L,q}^2}$ with $v_{L,\ell} = \langle \chi_{L,\ell}^0 \rangle$ and $v_{L,q} = \langle \chi_{L,q}^0 \rangle$. When the VEV ratio $\tan \beta = v_{L,q}/v_{L,\ell} \gg 1$, the singly-charged scalar is predominantly from the lepton doublet, i.e. $\chi_L^\pm \simeq (\chi_{L,\ell}^\pm - \cot \beta \chi_{L,q}^\pm)$, and its hadronic decays are highly suppressed by $\tan \beta$. Mediated by the mixing of $\chi_L^\pm$ with its heavy partner $\chi_R^\pm$ which couples directly to the DM, we can have three-body decay of DM, i.e. $N \rightarrow \ell^+ \ell^- \nu$, which can successfully fit the IceCube data [16].

9. Conclusion

In summary, we have discussed the phenomenology and cosmology of a new TeV-scale RHN DM candidate in the context of an $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{Y_L} \times U(1)_{Y_R}$ model, where we introduced an $SU(2)$ analog of the electroweak interaction and matter contents but with a common QCD. After symmetry breaking, there is an automatic Z_2 symmetry in the leptonic sector that guarantees the stability of the lightest RHN, which plays the role of cold DM in the universe. The DM relic density and direct detection constraints imply a lower limit of order 1 TeV for the RHN DM. We have also discussed a dilution mechanism which can push the RHN

dark matter mass up to the multi-PeV range (beyond the conventional unitarity limit) so that together with a weak breaking of the Z_2 symmetry, the RHN DM can in principle explain the PeV neutrinos observed at IceCube.

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