Observing left-right symmetry in the cosmic microwave background

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We consider the possibility of probing the left-right symmetric model (LRSM) via the cosmic microwave background (CMB). We adopt the minimal LRSM with Higgs doublets, also known as the doublet left-right model (DLRM), where all fermions including the neutrinos acquire masses only via their couplings to the Higgs bidoublet. Because of the Dirac nature of light neutrinos, there exist additional relativistic degrees of freedom which can thermalize in the early Universe by virtue of their gauge interactions corresponding to the right sector. We constrain the model from Planck 2018 bound on the effective relativistic degrees of freedom and also estimate the prospects for planned CMB stage IV experiments to constrain the model further. We find that the W_R boson mass below 4.06 TeV can be ruled out from the Planck 2018 bound at 2σ C.L. in the exact left-right symmetric limit which is equally competitive as the LHC bounds from dijet resonance searches. On the other hand, the Planck 2018 bound at 1σ C.L. can rule out a much larger parameter space out of reach of present direct search experiments, even in the presence of additional relativistic degrees of freedom around the TeV corner. We also study the consequence of these constraints on dark matter in the DLRM by considering a right-handed real fermion quintuplet to be the dominant dark matter component in the Universe.

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I. INTRODUCTION

Left-right symmetric models (LRSMs) [1–12] have been one of the most popular beyond the standard model (BSM) frameworks studied in the literature. Here the gauge symmetry of the SM is extended to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ so that the right-handed fermions (which are singlet in the SM) can form doublets under the new $SU(2)_R$. This not only makes the inclusion of a right-handed neutrino automatic, but also puts the left- and right-handed fermions on equal footing. Incorporating an additional discrete symmetry (or left-right parity) ensures that the theory is invariant under $SU(2)_L \leftrightarrow SU(2)_R$. Thus, the model not only explains the origin of parity violation in electroweak interactions through spontaneous breaking of a parity symmetric theory at the high-energy scale but also incorporates a right-handed neutrino, crucial to generate light neutrino masses and mixing observed at neutrino

oscillation experiments [13,14]. Apart from the possibility of realizing left-right symmetry as an intermediate symmetry in popular grand unified theories like $SO(10)$, a TeV-scale realization of a LRSM can also have very interesting consequences at collider experiments like the Large Hadron Collider (LHC) [15–21].

Conventionally, the very first proposals and studies of LRSMs [1–5] considered a scalar bidoublet for generating fermion masses and also for electroweak symmetry breaking, whereas a pair of scalar doublets were introduced for the purpose of left-right symmetry breaking at the high-energy scale. A very recent detailed study of this model can be found in Ref. [22]. On the other hand, the LRSM proposals put forward later [7–12] received much more attention due to the possibility of a seesaw origin of light neutrino masses through a combination of a type-I seesaw [23–26] and a type-II seesaw [8,27–30] or a type-III seesaw [31]. In the doublet left-right model (DLRM), in its minimal version, there is no such seesaw mechanism, as all fermions including neutrinos acquire Dirac masses by virtue of their couplings to the bidoublet scalar. While generating a sub-eV neutrino mass in this fashion requires relevant Yukawa couplings at the level of $<10^{-12}$, we adopt this minimal scenario to study some of the interesting phenomenological consequences. Radiative generation of light Dirac neutrinos in different left-right symmetric models has also been discussed over the past few decades

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[32–39] which may provide a UV completion of the minimal DLRM we discuss here. Since such UV completions do not drastically change the conclusions we reach in the present work, we stick to the DLRM for the sake of simplicity.

The Dirac nature of light neutrinos in the DLRM gives rise to additional relativistic degrees of freedom which can be thermalized in the early Universe due to their gauge interactions mediated by right sector gauge bosons. Such additional light degrees of freedom can be probed by precise measurements of the cosmic microwave background (CMB) anisotropies. Recent 2018 data from the CMB measurement by the *Planck* satellite [40] suggests the effective degrees of freedom for neutrinos as

$$N_{\text{eff}} = 2.99_{-0.33}^{+0.34} \quad (1)$$

at 2σ or 95% C.L. including baryon acoustic oscillation data. At 1σ C.L., it becomes more stringent to $N_{\text{eff}} = 2.99 \pm 0.17$. Both these bounds are consistent with the SM prediction $N_{\text{eff}}^{\text{SM}} = 3.045$ [41–43]. Upcoming CMB stage IV (CMB-S4) experiments are expected to put much more stringent bounds than *Planck* due to their potential of probing all the way down to $\Delta N_{\text{eff}} = N_{\text{eff}} - N_{\text{eff}}^{\text{SM}} = 0.06$ [44]. We use the existing constraints and put strong limits on the scale of left-right symmetry or, equivalently, the right sector gauge bosons W_R and Z_R . For comparison, we also check the corresponding bounds for the left-right asymmetric scenario by considering different $SU(2)_R$ gauge couplings. Interestingly, we find that the bounds on W_R and Z_R mass from the Planck 2018 bound on ΔN_{eff} at 2σ C.L. are equally competitive as the latest LHC bounds [16,17,20] and much stronger than the corresponding bounds from flavor physics [45]. On the other hand, the same Planck 2018 bound at 1σ C.L. can rule out a much larger mass window for W_R and Z_R out of reach of present collider experiments. In fact, CMB-S4 will be able to probe a much larger region of W_R and Z_R masses out of existing collider reach and, hence, can probe or rule out the minimal model. Since there have been a few recent studies on the gauged $B-L$ model with light Dirac neutrinos [46–49] and corresponding constraints due to the Planck 2018 bound on ΔN_{eff} , we also reproduce the corresponding

parameter space in the gauged $B-L$ model and compare with the one obtained in the DLRM. We point out the important difference due to the restricted range of DLRM gauge couplings g_R and g_{BL} unlike that in the gauged $B-L$ model. We also show the impact of these constraints on dark matter (DM) parameter space in the DLRM by considering a right-handed fermion quintuplet to be the dominant component of DM which can thermalize by virtue of its interactions with the SM mediated by right sector gauge bosons. We calculate the parameter space allowed from the observed DM relic and find the leftover parameter space after applying the ΔN_{eff} bound. Finally, we comment on the more stringent Planck 2018 1σ bound which can be satisfied if more light fields below the scale of left-right symmetry breaking in addition to the SM plus three right-handed neutrinos exist. In fact, we show that the DLRM with right-handed fermion quintuplet DM can give rise to a sufficient number of additional degrees of freedom to serve this purpose.

This paper is organized as follows. In Sec. II, we discuss the doublet left-right symmetric model followed by a discussion of additional relativistic degrees of freedom due to light Dirac neutrinos in Sec. III. In Sec. IV, we briefly discuss dark matter in the DLRM, particularly focusing on fermion quintuplet DM, followed by results and discussion in Sec. V. We finally conclude in Sec. VI.

II. THE DLRM

We briefly discuss the doublet left-right symmetric model in this section. The fermion and scalar content of the model are given in Tables I and II, respectively. The relevant Yukawa Lagrangian giving masses to the three generations of leptons is given by

$$\mathcal{L} = h_{ij} \bar{l}_{L,i} \Phi l_{R,j} + \widetilde{h}_{ij} \bar{l}_{L,i} \widetilde{\Phi} l_{R,j} + \text{H.c.}, \quad (2)$$

where the indices $i, j = 1, 2, 3$ represent the family indices for the three generations of fermions, $\widetilde{\Phi} = \tau_2 \Phi^* \tau_2$, and τ_2 is the Pauli matrix. The gauge structure of the model prevents any renormalizable Yukawa couplings involving the scalar doublets $\chi_{L,R}$. The scalar potential V_{scalar} is given by [22]

$$\begin{aligned} V_{\text{scalar}} = & -\mu_1^2 \text{Tr}[\Phi^\dagger \Phi] - \mu_2^2 \text{Tr}[\Phi^\dagger \widetilde{\Phi} + \widetilde{\Phi}^\dagger \Phi] - \mu_3^2 (\chi_L^\dagger \chi_L + \chi_R^\dagger \chi_R) + \lambda_1 (\text{Tr}[\Phi^\dagger \Phi])^2 + \lambda_2 \{ (\text{Tr}[\Phi^\dagger \widetilde{\Phi}])^2 + (\text{Tr}[\widetilde{\Phi}^\dagger \Phi])^2 \} \\ & + \lambda_3 \text{Tr}[\Phi^\dagger \widetilde{\Phi}] \text{Tr}[\widetilde{\Phi}^\dagger \Phi] + \lambda_4 \text{Tr}[\Phi^\dagger \Phi] \text{Tr}[\Phi^\dagger \widetilde{\Phi} + \widetilde{\Phi}^\dagger \Phi] + \mu'_1 (\chi_L^\dagger \Phi \chi_R + \chi_R^\dagger \Phi^\dagger \chi_L) + \mu'_2 (\chi_L^\dagger \widetilde{\Phi} \chi_R + \chi_R^\dagger \widetilde{\Phi}^\dagger \chi_L) \\ & + \rho_1 [(\chi_L^\dagger \chi_L)^2 + (\chi_R^\dagger \chi_R)^2] + \alpha_1 \text{Tr}[\Phi^\dagger \Phi] [(\chi_L^\dagger \chi_L) + (\chi_R^\dagger \chi_R)] + \alpha_2 e^{i\delta} [\text{Tr}[\widetilde{\Phi} \Phi^\dagger] (\chi_L^\dagger \chi_L) + \text{Tr}[\Phi \widetilde{\Phi}^\dagger] (\chi_R^\dagger \chi_R)] \\ & + \alpha_2 e^{-i\delta} [\text{Tr}[\widetilde{\Phi}^\dagger \Phi] (\chi_L^\dagger \chi_L) + \text{Tr}[\Phi^\dagger \widetilde{\Phi}] (\chi_R^\dagger \chi_R)] + \alpha_3 (\chi_L^\dagger \Phi \Phi^\dagger \chi_L + \chi_R^\dagger \Phi^\dagger \Phi \chi_R) + \alpha_4 (\chi_L^\dagger \widetilde{\Phi} \widetilde{\Phi}^\dagger \chi_L + \chi_R^\dagger \widetilde{\Phi}^\dagger \widetilde{\Phi} \chi_R). \end{aligned} \quad (3)$$

For details of the minimization of the scalar potential and resulting symmetry breaking, please refer to Ref. [22]. In the symmetry-breaking pattern, the neutral component of the Higgs doublet χ_R acquires a vacuum expectation

value (VEV) to break the gauge symmetry of the DLRM into that of the SM and then to the $U(1)$ of electromagnetism by the VEV of the neutral components of Higgs bidoublet Φ :

TABLE I. Fermionic fields of the present model including the SM fermions.

Particles	$SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$
$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$(3, 2, 1, \frac{1}{6})$
$Q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix}$	$(3, 1, 2, \frac{1}{3})$
$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	$(1, 2, 1, -\frac{1}{2})$
$\ell_R = \begin{pmatrix} \nu_R \\ e_R \end{pmatrix}$	$(1, 1, 2, -1)$

TABLE II. Scalar fields and their corresponding charges under all the symmetry groups.

Particles	$SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$
$\Phi = \begin{pmatrix} \phi^0 & \phi'^+ \\ \phi^- & \phi^0 \end{pmatrix}$	$(1, 2, 2, 0)$
χ_L	$(1, 2, 1, 1)$
χ_R	$(1, 1, 2, 1)$

$$SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow{\langle \chi_R \rangle} SU(2)_L \times U(1)_Y \xrightarrow{\langle \Phi \rangle} U(1)_{\text{em}}.$$

The VEVs of the neutral components of the Higgs fields can be denoted as

$$\langle \Phi \rangle = \begin{pmatrix} \frac{k_1}{\sqrt{2}} & 0 \\ 0 & \frac{k_2}{\sqrt{2}} \end{pmatrix}, \quad \langle \chi_L \rangle = \begin{pmatrix} 0 \\ \frac{v_L}{\sqrt{2}} \end{pmatrix}, \quad \langle \chi_R \rangle = \begin{pmatrix} 0 \\ \frac{v_R}{\sqrt{2}} \end{pmatrix},$$

where the VEV's k_1 and k_2 satisfy the VEV of the SM; namely, $v_{\text{SM}} = \sqrt{k_1^2 + k_2^2} \approx 246$ GeV. The spontaneous breaking of DLRM gauge symmetry down to $U(1)_{\text{em}}$ results in two charged massive vector bosons W_L and W_R , two neutral massive bosons Z_L and Z_R , and a massless photon as expected. The details of the mass spectrum of gauge bosons are shown in Appendix A.

The light Dirac neutrino mass and charged lepton mass are given, respectively, by

$$M_\nu = \frac{1}{\sqrt{2}}(k_1 h + k_2 \tilde{h}), \quad M_l = \frac{1}{\sqrt{2}}(k_2 h + k_1 \tilde{h}), \quad (4)$$

where the family indices are suppressed. Without any loss of generality, we make use of rotation in the $SU(2)_L \times SU(2)_R$ space so that only one of the neutral components of the Higgs bidoublet acquires a large vacuum expectation value, $k_1 \approx v_{\text{SM}}$ and $k_2 \approx 0$. Under these assumptions, the Dirac neutrino mass matrix is

$$M_\nu = \frac{1}{\sqrt{2}}(k_1 h), \quad (5)$$

while the charged lepton mass matrix is

$$M_l = \frac{1}{\sqrt{2}}(k_1 \tilde{h}). \quad (6)$$

Therefore, tiny sub-eV Dirac neutrino mass arises due to the smallness of Yukawa coupling h , while charged lepton masses are generated by corresponding Yukawa coupling $\tilde{h}$. The details of fermion-gauge boson couplings are shown in Appendix B. The details of the scalar mass spectrum are not derived here, as we do not need them for our analysis, and we refer to Ref. [22] for details of the same.

III. ΔN_{eff} IN DLRM

The effective number of relativistic degrees of freedom is defined as

$$N_{\text{eff}} \equiv \frac{8}{7} \left(\frac{11}{4} \right)^{4/3} \left(\frac{\rho_{\text{rad}} - \rho_\gamma}{\rho_\gamma} \right),$$

where $\rho_{\text{rad}} = \rho_\gamma + \rho_\nu$ is the net radiation content of the Universe. As mentioned earlier, the SM prediction is $N_{\text{eff}}^{\text{SM}} = 3.045$ [41–43], which is also consistent with the constraint from precision measurement of Z boson decay width at large electron positron (LEP) collider $N_\nu = 2.984 \pm 0.008$ [14]. Any deviation of N_{eff} from $N_{\text{eff}}^{\text{SM}}$ will therefore indicate the presence of additional relativistic species thermalized in the early Universe. While these additional relativistic degrees of freedom cannot fully thermalize with the SM bath through interactions mediated by a Z boson due to the strong LEP bound, they can thermalize via additional interactions or mediating particles not yet observed in direct search experiments. The right-handed neutrinos in the DLRM provide such an example. They can thermalize with the SM bath in the early Universe due to the interactions mediated by right sector gauge bosons, as depicted by the Feynman diagrams shown in Fig. 1. We consider negligible mixing between left and right sector gauge bosons and, hence, ignore the contributions coming from processes like $\bar{\nu}_R \nu_L \rightarrow f \bar{f}$ and $\bar{\nu}_L \nu_R \rightarrow f \bar{f}$. Additionally, the scalar-mediated interactions are negligible due to tiny Dirac Yukawa couplings.

To estimate the contribution in ΔN_{eff} , we need to check the decoupling temperature of the right-handed neutrinos. The decoupling occurs when the expansion rate of the Universe becomes more than the interaction rate. Hence, the decoupling temperature can be calculated from the following equality:

$$\Gamma(T_{\nu_R}^{\text{d}}) = H(T_{\nu_R}^{\text{d}}), \quad (7)$$

FIG. 1. Feynman diagrams of ν_R annihilation. Here, $l_R \equiv e_R, \mu_R, \tau_R$ and $f \equiv e_R, \mu_R, \tau_R, q_R$ with q being a quark.

where $\Gamma(T)$ is the interaction rate and $H(T)$ is the expansion rate of the Universe. The interaction rate can be written as

$$\Gamma(T) = n_{\nu_R}(T) \langle \sigma_{\text{Tot}} v \rangle, \quad (8)$$

where the number density n_{ν_R} for a relativistic neutrino can be written as

$$n_{\nu_R}(T) = \frac{3g_{\nu_R}}{4\pi^2} \zeta(3) T^3 \quad (9)$$

and the annihilation cross sections of right-handed neutrinos are given in Appendix C.

The expansion rate of the Universe can be written as

$$H(T) = \sqrt{\frac{8\pi G_N \rho(T)}{3}} = \sqrt{\frac{4\pi^3 G_N}{45} \left(g_*(T) + 3 \frac{7}{8} g_{\nu_R} \right) T^2}, \quad (10)$$

where g_{ν_R} is the internal degrees of freedom for right-handed neutrinos. Thus, the contribution of ν_R to effective relativistic degrees of freedom can be estimated as

$$\Delta N_{\text{eff}} = N_{\text{eff}} - N_{\text{eff}}^{\text{SM}} = N_{\nu_R} \left(\frac{T_{\nu_R}}{T_{\nu_L}} \right)^4 = N_{\nu_R} \left(\frac{g_{*S}(T_{\nu_L}^d)}{g_{*S}(T_{\nu_R}^d)} \right)^{4/3}, \quad (11)$$

where N_{ν_R} represents the number of relativistic right-handed neutrinos, $g_*(T)$ corresponds to the relativistic degrees of freedom at temperature T , $g_{*S}(T)$ corresponds to the relativistic entropy degrees of freedom at temperature T ,¹ and $T_{\nu_R}^d$ and $T_{\nu_L}^d$ are the decoupling temperatures for ν_R and ν_L , respectively. Thus, depending upon the decoupling temperature of ν_R and, hence, $g_*(T_{\nu_R}^d)$, the additional contribution to ΔN_{eff} can be kept within experimental upper limits. The lower the strength of ν_R interaction with the SM bath or the higher the mediator mass of ν_R -SM interactions, the larger will be $g_*(T_{\nu_R}^d)$ and, hence, the smaller will be ΔN_{eff} . A similar analysis for $U(1)_{B-L}$ extension of the SM can be found in Refs. [46–49], whereas some estimates in the context of radiative Dirac neutrino mass in LRSM were made in Refs. [35,37].

IV. DARK MATTER IN DLRM

The data from the Planck experiment which restrict the effective relativistic degrees of freedom in our Universe also reveal that more than 26% of the present Universe's energy density is composed of a nonluminous and nonbaryonic form of matter, known as dark matter. Apart from recent cosmology-based experiments like Planck, there have been

several astrophysical evidences for many decades suggesting the presence of DM [50–52]. In terms of density parameter Ω_{DM} and h = Hubble parameter/(100 km s⁻¹ Mpc⁻¹), the present DM abundance is conventionally reported as [40] $\Omega_{\text{DM}} h^2 = 0.120 \pm 0.001$ at 68% C.L. Given that none of the SM particles can be a viable DM candidate, several BSM proposals have been put forward, among which the weakly interacting massive particle (WIMP) paradigm is the most popular one. In this framework, a DM particle having masses and interactions similar to those around the electroweak scale gives rise to the observed relic after thermal freeze-out, a remarkable coincidence often referred to as the *WIMP miracle* [53].

The minimal DLRM discussed above does not have a stable DM candidate. One can, however, minimally extend the model by including additional scalar or fermionic multiplets in the spirit of the minimal dark matter scenario [54–56]. In these models, the dark matter candidate is stabilized either by a $Z_2 = (-1)^{B-L}$ subgroup of the $U(1)_{B-L}$ gauge symmetry or due to an accidental symmetry at the renormalizable level due to the absence of any renormalizable operator leading to dark matter decay. Such a minimal dark matter scenario in the LRSM has been studied recently by the authors of Refs. [57,58]. Some more recent works on DM in the LRSM can be found in Refs. [35–37,59–66]. Unlike in the triplet LRSM, where $SU(2)_R \times U(1)_{B-L}$ gauge symmetry is spontaneously broken by a scalar triplet with even $(B-L)$ charge (± 2), in the DLRM the same happens due to a scalar doublet with odd $(B-L)$ charge (± 1). Thus, there is no stabilizing symmetry like $Z_2 = (-1)^{B-L}$ in the DLRM to stabilize DM. Therefore, some DM candidates like the fermion triplet, fermion bidoublet, and scalar doublet discussed in the context of the triplet LRSM [57,58] are no longer stable in the DLRM due to the presence of renormalizable interactions with lighter fields. We therefore consider the option of a larger fermion multiplet as DM, and the minimal scenario is to consider a real fermion quintuplet of $B-L$ charge 0. Since we want to constrain the right sector gauge bosons from the cosmology bound on ΔN_{eff} , we particularly focus on right-handed fermion quintuplet DM whose relic abundance depends upon the strength of its annihilation through right sector gauge bosons.

In the pure left-right symmetric setup, one has to introduce a pair of left- and right-handed fermion quintuplets (having the same mass) which can be written in component form as

$$\Omega_L = \begin{pmatrix} \Omega_L^{++} \\ \Omega_L^+ \\ \Omega_L^0 \\ \Omega_L^- \\ \Omega_L^{--} \end{pmatrix}, \quad \Omega_R = \begin{pmatrix} \Omega_R^{++} \\ \Omega_R^+ \\ \Omega_R^0 \\ \Omega_R^- \\ \Omega_R^{--} \end{pmatrix}. \quad (12)$$

¹We use g_* and g_{*S} interchangeably, which is true in the SM at high temperatures.

Since we are discussing a general scenario with $g_L \neq g_R$, we consider the left fermion quintuplet to be very heavy and decoupled from the low-energy phenomenology. Even in the pure left-right symmetric limit $g_L = g_R$, one can make the left quintuplet decouple from the low-energy phenomenology by introducing a parity-odd scalar singlet whose nonzero VEV at a very high scale splits the right and left fermion masses. Such proposals where the left-right discrete symmetry or parity gets broken spontaneously before $SU(2)_R \times U(1)_{B-L}$ gauge symmetry were put forward long ago in Refs. [67–69]. While all the components of the fermion multiplet have the same tree-level masses, at the radiative level, there arises a mass splitting between charged (with Q) and neutral components given by [57,58]

$$M_{\Omega_R^Q} - M_{\Omega_R^0} \simeq \frac{\alpha_2}{4\pi} \frac{g_R^2}{g_L^2} M Q^2 [f(r_{W_R}) - c_M^2 f(r_{Z_R}) - s_W^2 s_M^2 f(r_{Z_L}) - c_W^2 s_M^2 f(r_\gamma)], \quad (13)$$

where $s_M = \sin \theta_M \equiv \tan \theta_W \frac{g_L}{g_R}$, $s_W = \sin \theta_W$, $r_X = M_X/M$, and

$$f(r) \equiv 2 \int_0^1 dx (1+x) \log [x^2 + (1-x)x^2].$$

The bare mass of the multiplet is denoted by M . Here, the one-loop self-energy corrections through mediations of gauge bosons are presented within the square bracket of the second expression. Because of such a tiny one-loop mass splitting, the next-to-lightest component of each DM multiplet can be thermally accessible during the dark matter freeze-out, and, hence, the coannihilation effects play a crucial role [70]. While in the triplet LRSM there exists the possibility that, for $M \gg M_{W_R}$, the neutral component of the multiplet can become heavier compared to the charged components, such possibilities do not arise in the DLRM. This was noted in Ref. [66], which we also confirm. Apart from its role in enhancing DM coannihilations, such a mass splitting, if small enough, may also induce inelastic DM-nucleon scattering mediated by W_R bosons. However, for our region of interest, such a tiny mass splitting does not arise. This ensures that W_R -mediated DM-nucleon scattering occurs only at one-loop level and, hence, remains suppressed. Spin-independent elastic DM nucleon scattering mediated by Z_R at tree level remains absent due to vanishing $B-L$ charge of the fermion quintuplet discussed in this work. Thus, the DM phenomenology of a right-handed fermion quintuplet is mainly governed by its gauge interactions, which are given by [66]

$$\begin{aligned} \mathcal{L}_{\Omega_R} \supset & -s_W s_M g_R \overline{Q} \Omega_R^Q Z_L^\mu \gamma_\mu \Omega_R^Q + c_M g_R \overline{Q} \Omega_R^Q Z_R^\mu \gamma_\mu \Omega_R^Q \\ & + c_W s_M g_R \overline{Q} \Omega_R^Q A^\mu \gamma_\mu \Omega_R^Q \\ & + \frac{g_R}{\sqrt{2}} \left(c_Q \overline{\Omega_R^{Q+1}} W_R^\mu \gamma_\mu \Omega_R^Q + \text{H.c.} \right), \end{aligned} \quad (14)$$

where $c_Q = \sqrt{(2+Q+1)(2-Q)}$ and Q is the electromagnetic charge of the quintuplet component.

V. RESULTS AND DISCUSSION

Using the recipe discussed in the previous section, we first calculate the decoupling temperature of right-handed neutrinos from the thermal bath for different values of W_R and Z_R mass and gauge coupling g_R . The variation of decoupling temperature with W_R mass for different values of g_R is shown in Fig. 2. Although both W_R and Z_R masses play a role in right-handed neutrino interactions with the thermal bath, we show the variation of the decoupling temperature as well as other physical quantities only in terms of W_R mass. This is due to the fact that Z_R mass typically depends upon W_R mass and is heavier than it, similar to Z and W masses of the SM. Also, we are not restricting ourselves to pure left-right symmetric limit $g_R = g_L$ and considering different values of g_R as well. The decoupling temperature rises for lower values of gauge coupling as well as higher values of W_R mass as seen from Fig. 2, which is expected as the corresponding rate of interactions decreases. Typically, the interaction cross section $\langle \sigma v \rangle$ of right-handed neutrinos depends upon the temperature as $g_R^4 T^2 / M_{W_R}^4$, and, hence, the rate of interaction is $\Gamma(T) \propto g_R^4 T^5 / M_{W_R}^4$. When this is compared with the Hubble parameter $H(T) \propto T^2$ as in Eq. (7) to find the decoupling temperature $T_{\nu_R}^d$, it is clear that, with an increase in W_R mass, the decoupling temperature will also rise. Similarly, a decrease in g_R will lead to an increase in the decoupling temperature for same value of M_{W_R} . This

FIG. 2. Decoupling temperature of right-handed neutrinos as a function of the W_R mass for different gauge couplings g_R .

FIG. 3. ΔN_{eff} as a function of the decoupling temperature (left panel) and W_R mass (right panel).

conclusion reached from approximate analytical formulas for the cross section and Hubble parameter also agrees with our numerical results shown in Fig. 2.

We then show the contribution to ΔN_{eff} in Fig. 3 as functions of the decoupling temperature as well as the W_R mass. The dependence of ΔN_{eff} on $T_{\nu_R}^d$ in the left panel in Fig. 3 can be understood from Eq. (11). The total number of relativistic degrees of freedom at decoupling temperature $g_{*s}(T_{\nu_R}^d)$ increases with the increasing $T_{\nu_R}^d$, which thereby decreases the value of ΔN_{eff} . However, at some point g_{*s} it will reach its maximum value and ΔN_{eff} becomes almost constant, as seen from the plateau region on the bottom left corner of the left panel in Fig. 3. Since $T_{\nu_R}^d$ is being varied independently and the corresponding $g_{*s}(T_{\nu_R}^d)$ is used to find ΔN_{eff} using Eq. (11), the behavior of this plot does not depend upon g_R and M_{W_R} . To show the dependence of ΔN_{eff} on such model parameters, we have made the plot shown in right panel in Fig. 3. The right panel in Fig. 3 represents the dependence of ΔN_{eff} on M_{W_R} for three benchmark values of g_R . One can see that the contribution to the ΔN_{eff} decreases with increasing M_{W_R} . This is because, for higher values of M_{W_R} , ν_R will decouple at some higher temperature and the contribution to the ΔN_{eff} will become smaller. Along with the Planck 2018 bound mentioned earlier, we also show the CMB-S4 sensitivity [71] as well as the Planck 2018 1σ limit, while the latter is the same as SPT-3G sensitivity [72]. Clearly, the Planck 2018 bound at 2σ C.L. itself rules out a W_R mass below 4.06 TeV with gauge coupling $g_R = g_L$. On the other hand, a future probe will be able to either confirm or rule out the model, even for very high W_R masses, out of reach of direct search experiments. Finally, we show the final parameter space in the g_R - M_{W_R} plane after applying Planck 2018 2σ constraints in Fig. 4.

Unlike in Refs. [46–49], where similar constraints on the $U(1)_{B-L}$ gauge boson were obtained, the crucial difference in the DLRM is that here one cannot tune the gauge couplings for a particular value of gauge boson mass in order to suppress the contribution to ΔN_{eff} . This is because the gauge couplings of $SU(2)_R$ and $U(1)_{B-L}$ are not arbitrary but related to the gauge coupling of $U(1)_Y$ (at the scale of left-right symmetry breaking) as

$$\frac{1}{g_Y^2} = \frac{1}{g_R^2} + \frac{1}{g_{BL}^2}. \quad (15)$$

FIG. 4. Parameter space in the g_R - M_{W_R} plane from Planck 2018 2σ constraints on ΔN_{eff} .

FIG. 5. Allowed values of g_{BL} and g_R which will reproduce the value g_Y after the spontaneous breaking of $SU(2)_R \times U(1)_{B-L}$ to the remaining $U(1)_Y$.

Since g_Y is known, one cannot change g_R and g_{BL} arbitrarily within their perturbative limits.² We show the allowed region of these two gauge couplings in Fig. 5. While we still have a large region within perturbative limits, we have chosen g_R to be either equal to g_L or smaller while keeping g_{BL} also below the order of one for our benchmark analysis. For $g_R > g_L$, the Planck bound becomes even more stringent, as we found in the scan plot shown in the g_R - M_{W_R} plane in Fig. 4. Thus, compared to $U(1)_{B-L}$ or other Abelian gauge models of Dirac neutrinos, the DLRM is much more constrained. For a comparison, we show the parameter space for the gauged $U(1)_{B-L}$ model with Dirac neutrinos in Fig. 6. This is a minimal gauged $B-L$ model, where there are three right-handed neutrinos having $B-L$ charge -1 each apart from the SM fermion content and neutrinos get a sub-eV Dirac mass by virtue of their tiny couplings with the SM Higgs boson. Since the $B-L$ gauge coupling is a free parameter and not related to SM gauge couplings in this model, one can tune the gauge coupling arbitrarily to evade the stringent Planck bound on ΔN_{eff} as can be seen from Fig. 6.

To check the impact of these constraints on DM parameter space, we then calculate the relic of right fermion quintuplet DM. For DM relic calculation, we first implement the model in SARAH [74] and then feed the model files into micrOMEGAs [75] for relic calculations. We then consider three benchmark combinations of (g_R, W_R) while

²For a discussion related to perturbativity constraints on similar models, please refer to Ref. [73] and references therein.

FIG. 6. Allowed $g_{BL} - M_{Z_{BL}}$ parameter space from the Planck 2018 bound on ΔN_{eff} at 2σ in the minimal $U(1)_{B-L}$ gauge model with light Dirac neutrinos.

keeping the scale of left-right gauge symmetry breaking v_R fixed. The resulting variation of DM relic as a function of the DM mass is shown in left panel in Fig. 7. The resonances corresponding to W_R and Z_R masses are clearly visible in this plot. Unlike in the quintuplet DM scenario in the triplet LRSB [58], here the two resonances are quite close to each other due to a smaller ratio of Z_R to W_R mass in the DLRM. The results shown in left panel in Fig. 7 also agree with that shown in Ref. [66]. We then scan the parameter space of W_R and Ω_R^0 masses and show the region satisfying the correct DM relic in the right panel in Fig. 7. Multiple allowed values of DM mass for a fixed W_R mass are arising due to annihilation and coannihilations of Ω_R^0 , $\Omega_R^\pm$, and $\Omega_R^{\pm\pm}$ mediated by W_R and Z_R bosons where Z_R is slightly heavier than W_R ($M_{Z_R} \approx 1.2M_{W_R}$ in pure left-right symmetric limit $g_R = g_L$). Furthermore, the red and green solid lines in the right panel in Fig. 7 show the possibility of having three different values of DM mass with the correct relic abundance for a fixed W_R mass. The same behavior is also seen in the left panel in Fig. 7 with a fixed W_R mass, where the red and green lines satisfy the relic at three different values of DM mass. On the other hand, the blue dashed line in the left panel in Fig. 7 satisfies the correct relic abundance for five different values of DM mass, a feature which is also depicted by the blue solid line in the scan plot (right panel in Fig. 7). We also apply the corresponding bounds on the W_R mass from Planck constraints on ΔN_{eff} at 2σ C.L. as horizontal shaded lines so that the region below the respective lines is disallowed. Clearly, some part of the DM parameter space on the right

FIG. 7. Left panel: DM relic as a function of the DM mass for different benchmark combinations of g_R and W_R . Right panel: Parameter space satisfying relic abundance criteria of right-handed fermion quintuplet dark matter in the DLRM. The shaded regions are ruled out by the Planck 2018 bound on ΔN_{eff} at 2σ for respective values of g_R .

side of the scan plot (right panel in Fig. 7) gets disfavored for all three values of g_R by ΔN_{eff} bounds, as the lower region of the dashed lines is excluded from Planck constraints on ΔN_{eff} at 2σ C.L. In the parabolic part of the parameter space in the scan plot, a small part of the parameter space for the $g_R = g_L$ case is disfavored by ΔN_{eff} bounds, while the parabolic lines corresponding to $g_R \neq g_L$ still remain allowed from respective ΔN_{eff} . We do not show other existing bounds on the W_R mass from flavor or LHC data, as they are either equally or less strong compared to the bounds derived here.

VI. CONCLUSION

We have studied the minimal left-right symmetric model with Higgs doublets, known as the doublet left-right model, where the left-right symmetry is broken spontaneously by a Higgs doublet instead of Higgs triplets in the LRSM with type-I plus type-II seesaw for light neutrino masses. In the minimal DLRM, a light neutrino mass can be realized only through the Higgs bidoublet with tiny coupling to neutrinos leading to a sub-eV Dirac neutrino mass. Because of $SU(2)_R \times U(1)_{B-L}$ gauge interactions, the right-handed neutrinos can thermalize in the early Universe, thereby contributing to the effective relativistic degrees of freedom N_{eff} which are tightly constrained by CMB measurements. We constrain the scale of left-right symmetry from the requirement of satisfying the Planck 2018 bound on ΔN_{eff} not only for pure left-right symmetric limit $g_R = g_L$ but also for other values of g_R . While g_R and g_{BL} cannot be chosen arbitrarily in the LRSM due to their relation with $U(1)_Y$ coupling of the SM, the parameter space gets tightly

constrained. For $g_R = g_L$, the Planck 2018 bound at 2σ C.L. can rule out W_R mass all the way up to 4.06 TeV, which is as competitive as the existing collider bounds on W_R from dijet resonance searches. For example, the ATLAS dijet resonance search rules out such an additional charged vector boson mass up to 4 TeV [20], while a similar analysis by CMS rules out up to 3.6 TeV for SM-like gauge couplings [21]. Our conclusions also agree with the recent model-independent calculations [76] where, considering four fermion interactions of right-handed neutrinos and their contribution to N_{eff} , the authors constrained the interaction strength to be 10^{-5} – 10^{-3} times the usual Fermi coupling constant. This bound we derive here, that is, $M_{W_R} > 4.06$ TeV, can be made weaker for smaller values of g_R . For example, in the case of $g_R = 0.5$, we get $M_{W_R} > 3.17$ TeV and $g_R = 0.4$ leads to $M_{W_R} > 2.77$ TeV. We also make a comparison with similar constraints derived in the $U(1)_{B-L}$ gauge model with light Dirac neutrinos, where, due to the freedom in choosing the gauge coupling arbitrarily, one can have a much lighter $B-L$ gauge boson as well. However, the choices of g_R as well as g_{BL} are not arbitrary in the DLRM due to their nontrivial connection to g_Y of the standard model. While we perform our analysis without explicitly solving the Boltzmann equations, the conclusions do not change much as long as the decoupling temperature of right-handed neutrinos remains much higher compared to that of active neutrinos. This was also noted by the authors of Ref. [76], and our parameter space is precisely confined to this regime.

We also show the impact of ΔN_{eff} constraints on dark matter parameter space in the DLRM. While the DLRM does not have a dark matter candidate on its own, we

FIG. 8. Left panel: Additional relativistic degrees of freedom required to bring N_{eff} within Planck 2018 1σ bound $N_{\text{eff}} = 2.99 \pm 0.17$. Right panel: Allowed parameter space in the g_R - M_{W_R} plane from Planck 2018 1σ bound $N_{\text{eff}} = 2.99 \pm 0.17$ after considering additional relativistic degrees of freedom.

incorporate the presence of an additional fermion quintuplet DM in the minimal DM spirit. Since such a real fermion quintuplet does not have any renormalizable coupling with other fermions or scalars of the DLRM, the relic abundance of DM, the neutral component of the right-handed fermion quintuplet, depends crucially on its annihilation and coannihilation mediated by W_R and Z_R gauge bosons. We constrain the parameter space satisfying the correct DM relic by using the respective ΔN_{eff} bounds for different g_R . We find available parameter space satisfying the correct DM relic even after applying the Planck 2018 bound on ΔN_{eff} at 2σ C.L.

We also compare our results in view of the more stringent Planck 2018 1σ bound $N_{\text{eff}} = 2.99 \pm 0.17$, which rules out all the parameter space if we assume only the SM plus three right-handed neutrinos to be contributing to the relativistic degrees of freedom (DOF) below the scale of left-right symmetry breaking. However, if there are more relativistic degrees of freedom due to the presence of light physical fields resulting from the scalar and fermion multiplets, one can satisfy the Planck 2018 1σ bound as well. We show the required additional DOF in the left panel in Fig. 8 as a function of the right-handed neutrino decoupling temperature. All the points on the dashed line in the left panel in Fig. 8 give rise to $\Delta N_{\text{eff}} = 0.12$, so that the points below this line are ruled out. These DOF can arise from the DLRM with right-handed fermion quintuplet DM. For example, the scalar bidoublet has four physical DOF apart from the SM Higgs boson, while the pair of Higgs doublets can give rise to five more physical DOF. Similarly, right-handed fermion quintuplet DM has approximately ten DOF. Thus, one can have 19 additional DOF in the DLRM

with right-handed fermion quintuplet DM. Similar ways of avoiding such strict cosmological bounds on ΔN_{eff} have also been discussed in the recent work [76]. In the right panel in Fig. 8, we show the allowed parameter space on the g_R - M_{W_R} plane from the Planck 2018 1σ bound after considering the required additional relativistic DOF at decoupling temperatures of 200 and 650 GeV, respectively. Since a decoupling temperature of 200 GeV can be achieved with a lighter W_R gauge boson, we get more allowed parameter space compared to the case with a higher decoupling temperature of right-handed neutrinos. However, the presence of such additional light degrees of freedom will face stringent tests from collider as well as flavor physics constraints. Additionally, possible UV completion of this minimal model in order to explain the tiny origin of Dirac neutrino mass naturally [32–39] may involve additional interaction portals of right-handed neutrinos to thermalize with the standard bath. While the strength of these portals can be tuned with more freedom compared to the gauge portals discussed in this work and, hence, are unlikely to put stronger constraints than what we obtain here, it will be interesting to study the details of such scenarios, especially in the context of dark matter and collider phenomenology. A more detailed investigation of such scenarios is left for future studies.

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APPENDIX A: PHYSICAL MASSES OF GAUGE BOSONS

Covariant derivatives of the scalar fields in the DLRM can be written as

$$\begin{aligned} D_\mu \Phi &= \partial_\mu \Phi - i \frac{g_L}{2} (\sigma \cdot W_{L\mu}) \Phi + i \frac{g_R}{2} \Phi (\sigma \cdot W_{R\mu}), \\ D_\mu \chi_L &= \partial_\mu \chi_L - i \frac{g_L}{2} (\sigma \cdot W_{L\mu}) \chi_L - i g_{BL} Q_{BL} B'_\mu \chi_L, \\ D_\mu \chi_R &= \partial_\mu \chi_R - i \frac{g_R}{2} (\sigma \cdot W_{R\mu}) \chi_R - i g_{BL} Q_{BL} B'_\mu \chi_R, \end{aligned} \quad (A1)$$

where

$$\sigma \cdot W_{L/R\mu} = \begin{pmatrix} W_{L/R\mu}^3 & \sqrt{2} W_{L/R\mu}^+ \\ \sqrt{2} W_{L/R\mu}^- & -W_{L/R\mu}^3 \end{pmatrix}. \quad (A2)$$

The corresponding kinetic Lagrangian of scalar fields is

$$M_{V^0}^2 = \begin{pmatrix} \frac{1}{8} g_L^2 (k_1^2 + k_2^2 + v_L^2) & -\frac{1}{4} g_L g_R (k_1^2 + k_2^2) & -\frac{1}{4} g_{BL} g_L v_L^2 \\ -\frac{1}{4} g_L g_R (k_1^2 + k_2^2) & \frac{1}{8} g_L^2 (k_1^2 + k_2^2 + v_R^2) & -\frac{1}{4} g_{BL} g_R v_R^2 \\ -\frac{1}{4} g_{BL} g_L v_L^2 & -\frac{1}{4} g_{BL} g_R v_R^2 & \frac{1}{8} g_{BL}^2 (v_L^2 + v_R^2) \end{pmatrix}. \quad (A5)$$

As expected, the neutral gauge boson mass matrix has one vanishing eigenvalue, corresponding to the massless photon. After diagonalization of the mass matrices, we can represent the gauge fields in terms of physical gauge boson states as

$$\begin{aligned} W_{L\mu}^3 &= \frac{e}{g_Y} Z_{L\mu} + \frac{e}{g_L} A_\mu + 0 Z_{R\mu}, \\ W_{R\mu}^3 &= -\frac{e g_Y}{g_L g_R} Z_{L\mu} + \frac{e}{g_R} A_\mu + \frac{g_Y}{g_{BL}} Z_{R\mu}, \\ B'_\mu &= -\frac{e g_Y}{g_L g_{BL}} Z_{L\mu} + \frac{e}{g_{BL}} A_\mu - \frac{g_Y}{g_R} Z_{R\mu}. \end{aligned} \quad (A6)$$

Also, we can express these couplings as $\sin \theta_W = \frac{e}{g_L}$ and $\cos \theta_W = \frac{e}{g_Y}$ with θ_W being the Weinberg angle. In the DLRM, $Z_{L\mu}$ and $Z_{R\mu}$ will also mix as the bidoublet Φ transforms nontrivially under both $SU(2)_L$ and $SU(2)_R$ gauge groups. The mixing can be represented as

$$\begin{aligned} Z_{L\mu} &= \cos \delta Z_\mu - \sin \delta Z'_\mu, \\ Z_{R\mu} &= \sin \delta Z_\mu + \cos \delta Z'_\mu, \end{aligned} \quad (A7)$$

where the mixing angle can be written as

$$\tan 2\delta = \frac{2(M_{L,R}^0)^2}{(M_L^0)^2 - (M_R^0)^2} \quad (A8)$$

$$\begin{aligned} \mathcal{L}_{\text{scalar}} &= \text{Tr}[(D_\mu \Phi)^\dagger (D_\mu \Phi)] + (D_\mu \chi_L)^\dagger (D_\mu \chi_L) \\ &\quad + (D_\mu \chi_R)^\dagger (D_\mu \chi_R). \end{aligned} \quad (A3)$$

Considering the scalar VEVs as

$$\langle \Phi \rangle = \begin{pmatrix} \frac{k_1}{\sqrt{2}} & 0 \\ 0 & \frac{k_2}{\sqrt{2}} \end{pmatrix}, \quad \langle \chi_L \rangle = \begin{pmatrix} 0 \\ \frac{v_L}{\sqrt{2}} \end{pmatrix}, \quad \langle \chi_R \rangle = \begin{pmatrix} 0 \\ \frac{v_R}{\sqrt{2}} \end{pmatrix}$$

the charged vector boson mass matrix can be written as

$$M_{V^\pm}^2 = \begin{pmatrix} \frac{1}{4} v_L^2 g_L^2 + \frac{1}{4} (k_1^2 + k_2^2) g_L^2 & -\frac{1}{2} g_L g_R k_1 k_2 \\ -\frac{1}{2} g_L g_R k_1 k_2 & \frac{1}{4} v_R^2 g_R^2 + \frac{1}{4} (k_1^2 + k_2^2) g_R^2 \end{pmatrix}, \quad (A4)$$

whereas the neutral vector boson mass matrix is

with

$$(M_L^0)^2 = \frac{e^2 (g_L^2 + g_Y^2) (k_1^2 + k_2^2 + v_L^2)}{8 g_L^2 g_Y^2}, \quad (A9)$$

$$(M_{L,R}^0)^2 = \frac{e (g_L^2 + g_Y^2) (g_R^2 (k_1^2 + k_2^2) - g_{BL}^2 v_L^2)}{4 g_{BL} g_L g_R}, \quad (A10)$$

$$(M_R^0)^2 = \frac{g_Y^2 (2 g_{BL}^2 g_R^2 v_R^2 + g_R^4 (k_1^2 + k_2^2 + v_R^2) + g_{BL}^4 (v_L^2 + v_R^2))}{8 g_{BL}^2 g_R^2}. \quad (A11)$$

The charged vector boson states are

$$\begin{aligned} W_{L\mu}^\pm &= \cos \zeta W_{1\mu}^\pm - \sin \zeta W_{2\mu}^\pm, \\ W_{R\mu}^\pm &= \sin \zeta W_{1\mu}^\pm + \cos \zeta W_{2\mu}^\pm \end{aligned} \quad (A12)$$

with

$$\tan 2\zeta = \frac{2M_{LR}^2}{M_L^2 - M_R^2}, \quad (A13)$$

where

$$M_L^2 = \frac{1}{4} g_L^2 (k_1^2 + k_2^2 + v_L^2), \quad (A14)$$

$$M_R^2 = \frac{1}{4} g_R^2 (k_1^2 + k_2^2 + v_R^2), \quad (\text{A15})$$

$$M_{LR}^2 = -\frac{1}{2} g_L g_R k_1 k_2. \quad (\text{A16})$$

After diagonalization of $\{W_L^\pm, W_R^\pm\}$ to $\{W_1^\pm, W_2^\pm\}$, we can have the corresponding mass-squared terms for the charged physical gauge bosons as

$$M_1^2 = \frac{1}{4} (g_L^2 (k_1^2 + k_2^2 + v_L^2) \cos^2 \zeta - 2g_L g_R k_1 k_2 \sin 2\zeta + g_R^2 (k_1^2 + k_2^2 + v_R^2) \sin^2 \zeta), \quad (\text{A17})$$

$$M_2^2 = \frac{1}{4} (g_R^2 (k_1^2 + k_2^2 + v_R^2) \cos^2 \zeta + 2g_L g_R k_1 k_2 \sin 2\zeta + g_L^2 (k_1^2 + k_2^2 + v_L^2) \sin^2 \zeta). \quad (\text{A18})$$

Note that we have taken $k_2 = 0$, which is equivalent to vanishing tree-level mixing angle ζ . One can, however, generate radiative mixing between charged vector bosons, but that is typically very small $< 10^{-7}$ [37]. Note that, although we write W_1 , W_2 , Z , and Z' as physical massive gauge boson states here to show the details, in the main text we continue to use W_L , W_R , Z_L , and Z_R for better clarity.

APPENDIX B: FERMION-GAUGE BOSON INTERACTIONS IN DLRM

In this section, we note down the fermion interactions with massive vector bosons. The kinetic term of leptons in the DLRM is given by

$$\mathcal{L}_\ell = i\bar{\ell}_L \not{D}^L \ell_L + i\bar{\ell}_R \not{D}^R \ell_R, \quad (\text{B1})$$

where

TABLE IV. Fermion interaction with the Z' boson.

Fermions	Z' in LRSM
$\bar{e}_L \gamma_\mu e_L$	$i \left(\frac{g_{BL} g_L \tan \theta_W \cos \delta}{2g_R} + \frac{g_L \cos 2\theta_W \sin \delta}{\cos \theta_W} \right)$
$\bar{e}_R \gamma_\mu e_R$	$i \left(\frac{1}{2} g_L \tan \theta_W \left(\frac{g_{BL}}{g_R} - \frac{g_R}{g_{BL}} \right) \cos \delta - 2 \sin \theta_W \sin \delta \right)$
$\bar{\nu}_L \gamma_\mu \nu_L$	$i \left(\frac{g_{BL} g_L \tan \theta_W \cos \delta}{2g_R} - \frac{g_L (1 - 2 \sin^2 \theta_W \cos^2 \theta_W) \sin \delta}{2 \cos^2 \theta_W \sin \theta_W} \right)$
$\bar{\nu}_R \gamma_\mu \nu_R$	$i \left(\frac{(g_{BL}^2 + g_R^2) g_L \tan \theta_W \cos \delta}{2g_{BL} g_R} \right)$
$\bar{u}_L \gamma_\mu u_L$	$i \left(-\frac{g_{BL} g_L \tan \theta_W \cos \delta}{6g_R} + g_L \sin \theta_W \left(-\frac{\cot \theta_W}{2} + \frac{\tan \theta_W}{6} \right) \sin \delta \right)$
$\bar{u}_R \gamma_\mu u_R$	$i \left(\frac{1}{6} g_L \tan \theta_W \left(-\frac{g_{BL}}{g_R} + \frac{3g_R}{g_{BL}} \right) \cos \delta + 4 \sin \theta_W \sin \delta \right)$
$\bar{d}_L \gamma_\mu d_L$	$i \left(-\frac{g_{BL} g_L \tan \theta_W \cos \delta}{6g_R} + g_L \sin \theta_W \left(\frac{\cot \theta_W}{2} + \frac{\tan \theta_W}{6} \right) \sin \delta \right)$
$\bar{d}_R \gamma_\mu d_R$	$i \left(-g_L \tan \theta_W \left(\frac{(g_{BL}^2 + 3g_R^2) \cos \delta}{6g_{BL} g_R} + \frac{\sin \theta_W \sin \delta}{3} \right) \right)$

$$D_\mu^L \ell_L = \partial_\mu \ell_L - i \frac{g_L}{2} (\sigma \cdot W_{L\mu}) \ell_L + i \frac{g_{BL}}{2} B_\mu \ell_L, \quad (\text{B2})$$

$$D_\mu^R \ell_R = \partial_\mu \ell_R - i \frac{g_R}{2} (\sigma \cdot W_{R\mu}) \ell_R + i \frac{g_{BL}}{2} B_\mu \ell_R. \quad (\text{B3})$$

The same kinetic Lagrangian is, in fact, applicable to quarks, too, if we include gluons in the covariant derivative. We show the interactions with neutral massive vector bosons in Tables III and IV, respectively.

The interaction of neutrino ν and charged leptons ℓ with W_2 (or W_R) is similar to the ones with W_L except that g_L is replaced by g_R :

$$-\frac{ig_R}{\sqrt{2}} \gamma_\mu P_R,$$

where we have ignored the details of the right-handed lepton mixing matrix, taking it to be a unit matrix.

TABLE III. Fermion interactions with the Z boson.

Fermions	Z bosons	
	Z in SM	Z in DLRM
$\bar{e}_L \gamma_\mu e_L$	$-i \frac{g_L}{2 \cos \theta_W} \cos 2\theta_W$	$i \left(-\frac{g_L \cos 2\theta_W}{2 \cos \theta_W} \cos \delta + \frac{g_{BL} g_L \tan \theta_W \sin \delta}{2g_R} \right)$
$\bar{e}_R \gamma_\mu e_R$	$i \frac{g_L}{\cos \theta_W} \sin^2 \theta_W$	$i \left(\frac{g_L \sin^2 \theta_W \cos \delta}{\cos \theta_W} + \frac{1}{2} \left(\frac{g_{BL}}{g_R} - \frac{g_R}{g_{BL}} \right) g_L \tan \theta_W \sin \delta \right)$
$\bar{\nu}_L \gamma_\mu \nu_L$	$i \frac{g_L}{2 \cos \theta_W}$	$i \left(\frac{g_L (1 - 2 \sin^2 \theta_W \cos^2 \theta_W) \cos \delta}{2 \cos^2 \theta_W \sin \theta_W} + \frac{g_{BL} g_L \tan \theta_W \sin \delta}{2g_R} \right)$
$\bar{u}_L \gamma_\mu u_L$	$-i \frac{g_L}{2 \cos \theta_W} \left(\cos^2 \theta_W - \frac{\sin^2 \theta_W}{3} \right)$	$i \left(-g_L \sin \theta_W \left(\frac{\cot \theta_W}{2} - \frac{\tan \theta_W}{6} \right) \cos \delta + \frac{g_{BL} g_L \tan \theta_W \sin \delta}{6g_R} \right)$
$\bar{u}_R \gamma_\mu u_R$	$-i \frac{2}{3} g_L \frac{\sin^2 \theta_W}{\cos \theta_W}$	$i \left(\frac{1}{6} g_L \tan \theta_W \left(-4 \sin \theta_W \cos \delta + \left(-\frac{g_{BL}}{g_R} + \frac{3g_R}{g_{BL}} \right) \sin \delta \right) \right)$
$\bar{d}_L \gamma_\mu d_L$	$-i \frac{g_L}{2 \cos \theta_W}$	$i \left(-\left(\frac{g_L \cos \theta_W}{2} + \frac{g_L \sin^2 \theta_W}{6 \cos \theta_W} \right) \cos \delta - \frac{g_{BL} g_L \tan \theta_W \sin \delta}{6g_R} \right)$
$\bar{d}_R \gamma_\mu d_R$	$i \frac{1}{3} g_L \frac{\sin^2 \theta_W}{\cos \theta_W}$	$i \left(\frac{1}{6} g_L \tan \theta_W \left(2 \sin \theta_W \cos \delta + \left(-\frac{g_{BL}}{g_R} - \frac{3g_R}{g_{BL}} \right) \sin \delta \right) \right)$
$\bar{\nu}_R \gamma_\mu \nu_R$	0	$i \left(\frac{(g_{BL}^2 + g_R^2) g_L \tan \theta_W \sin \delta}{2g_{BL} g_R} \right)$

APPENDIX C: ANNIHILATION CROSS SECTIONS OF RIGHT-HANDED NEUTRINOS

The annihilation cross sections of ν_R mediated by right sector gauge bosons are

$$\sigma_{\nu_R \bar{\nu}_R \rightarrow q \bar{q}} = \frac{(a^2 + b^2) \sqrt{1 - \frac{4m_q^2}{s}} (c^2(2m_q^2 + s) + d^2(s - 4m_q^2))}{192\pi M_{Z_R}^4}, \quad (C1)$$

$$\begin{aligned} \sigma_{\nu_R \bar{\nu}_R \rightarrow \ell_R \bar{\ell}_R} = & \frac{\sqrt{1 - \frac{4m_\ell^2}{s}}}{192\pi M_{W_R}^4 M_{Z_R}^4} (2m_\ell^2 (M_{W_R}^4 (a^2 + b^2) (c^2 - 2d^2) - M_{W_R}^2 M_{Z_R}^2 (a + b) (c - 2d) - M_{Z_R}^4 (a + b)) \\ & + s (M_{W_R}^4 (a^2 + b^2) (c^2 + d^2) - M_{W_R}^2 M_{Z_R}^2 (a + b) (c + d) + 2M_{Z_R}^4 (a + b))), \end{aligned} \quad (C2)$$

where

$$a = b = \frac{g_L \tan \theta_W}{2} \frac{g_{BL}^2 + g_R^2}{2g_{BL}g_R}, \quad (C3)$$

$$c = \begin{cases} \frac{g_L \tan \theta_W}{2} \frac{2g_{BL}^2 - g_R^2}{2g_{BL}g_R} & \text{(for charged leptons),} \\ \frac{g_L \tan \theta_W}{2} \frac{3g_R^2 - 2g_{BL}^2}{6g_{BL}g_R} & \text{(for up-type quark),} \\ -\frac{g_L \tan \theta_W}{2} \frac{2g_{BL}^2 + g_R^2}{2g_{BL}g_R} & \text{(for down-type quark),} \end{cases} \quad (C4)$$

$$d = \begin{cases} -\frac{g_L \tan \theta_W}{4g_{BL}} g_R & \text{(for charged leptons),} \\ \frac{g_L \tan \theta_W}{4g_{BL}} g_R & \text{(for up-type quark),} \\ -\frac{g_L \tan \theta_W}{4g_{BL}} g_R & \text{(for down-type quark).} \end{cases} \quad (C5)$$

We consider the mixing between left and right sector gauge bosons to be negligible and, hence, do not take it into account in our analysis.

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