

Decaying scalar dark matter in the minimal left-right symmetric model

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In the minimal left-right symmetric theory, the dark matter candidate is usually ascribed to the lightest right-handed neutrino. Here we present an alternative decaying dark matter candidate in this model in terms of the lightest neutral scalar from the $SU(2)_R$ -triplet field. This setup requires a vast hierarchy between the scalar mass and the left-right symmetry breaking scale, which renders the scalar dark matter sufficiently stable on cosmological time scales. The stability of the dark matter imposes constraints on the right-handed neutrino mass, which has consequences for the neutrino mass generation, as well as for leptogenesis. Although somewhat fine-tuned, it provides a very economical scenario wherein the minimal left-right model can simultaneously explain dark matter, neutrino masses, and the matter-antimatter asymmetry of the Universe.

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I. INTRODUCTION

The observation of nonzero neutrino masses through neutrino oscillations is the strongest proof that the otherwise wildly successful Standard Model (SM) of particle physics is incomplete. Ad-hoc SM extensions to accommodate neutrino masses provide interesting toy models, but ultimately a more complete SM successor that naturally generates neutrino masses is desirable. One popular model of this kind is the left-right (LR) symmetric model [1–5], which extends the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ to the manifestly parity-symmetric $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and *requires* right-handed (RH) neutrinos for anomaly cancellation. Neutrino masses are thus theoretically inevitable in LR models. If the LR gauge group is broken by scalar triplets $\Delta_{L,R}$ rather than doublets, lepton number is spontaneously broken and one obtains Majorana neutrino masses via both the type-I [6–9] and type-II [10–13] seesaw mechanisms, shedding light not just on the existence, but also on the *smallness* of neutrino masses.

LR models also have useful ingredients to explain the second biggest shortcoming of the SM, namely, the existence of dark matter (DM). For one, the LR breaking via triplets allows one to elegantly incorporate weakly interacting massive particles, stabilized by the leftover $U(1)_{B-L}$ subgroup $\mathbb{Z}_2$ [14–16]. Even without introducing additional particles, LR models can accommodate *unstable* DM candidates, the almost exclusively discussed example being the lightest RH neutrino [17–20]. Aside from the additional gauge interactions, this case is similar to the popular “keV sterile neutrino” DM candidate [21,22], which eventually decays into an active neutrino and a photon [23], and thus can be constrained by nonobservation of nearly monochromatic X-ray signatures [24–26].

As we discuss here, the simplest LR model actually features another potential DM candidate besides the sterile neutrino, namely, the Higgs boson associated with the $SU(2)_R \times U(1)_{B-L}$ breaking, i.e., the real part of the neutral triplet component, $\Re(\Delta_R^0)$. It is massive, electrically neutral, and can be sufficiently decoupled from the SM to be long lived on cosmological time scales. The possibility of long-lived $\Re(\Delta_R^0)$ being a DM candidate was briefly mentioned before in Refs. [18,20], but without a detailed discussion of its phenomenology. We fill this gap in the literature and provide a first detailed study of this kind of DM and the restrictions it imposes on the LR parameter space. We establish $\Re(\Delta_R^0)$ as a viable decaying DM candidate in the keV to multi-MeV mass range, eventually decaying to two photons. While severe fine-tuning is required to achieve the large hierarchy between the DM mass and the LR breaking scale, we show that it restricts the

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RH neutrino mass scale, thereby influencing the neutrino mass generation via seesaw mechanism, as well as the generation of the observed baryon asymmetry of the Universe via leptogenesis [27]. In this way, this economical scenario of the minimal LR model can simultaneously explain all three major pieces of empirical evidence for beyond SM physics, i.e., DM, neutrino masses, and the baryon asymmetry of our Universe. Moreover, we conjecture that the same DM candidate can be realized in LR extensions such as the Pati-Salam model [1] or the $SO(10)$ grand unified theory [28].

II. REVIEW OF THE MINIMAL LR MODEL

The LR model is based on the gauge group $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [1–5]. The fermion fields transform as either left- or right-handed $SU(2)$ doublets,

$$Q_{L,R} = \begin{pmatrix} u \\ d \end{pmatrix}_{L,R}, \quad \psi_{L,R} = \begin{pmatrix} \nu \\ e \end{pmatrix}_{L,R}. \quad (1)$$

It naturally incorporates the RH neutrino ν_R within an RH lepton doublet, thus enabling neutrino mass generation through the seesaw mechanism [7,12].

The Higgs sector comprises the triplets $\Delta_L(3, 1, 2)$, $\Delta_R(1, 3, 2)$ and the bidoublet $\Phi(2, 2, 0)$ under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$,

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}, \quad \Delta_{L,R} = \begin{pmatrix} \delta^+/\sqrt{2} & \delta^{++} \\ \delta^0 & -\delta^+/\sqrt{2} \end{pmatrix}_{L,R}. \quad (2)$$

After the neutral component of Δ_R develops a vacuum expectation value (VEV) $\langle \delta_R^0 \rangle = v_R/\sqrt{2}$, the $SU(2)_R \times U(1)_{B-L}$ symmetry is broken to $U(1)_Y$, giving masses to the $W_R^\pm$ and Z_R gauge bosons. The VEVs $\langle \phi_{1,2}^0 \rangle = \kappa_{1,2}/\sqrt{2}$ break the remaining $SU(2)_L \times U(1)_Y$ down to the usual $U(1)_{em}$, thereby setting the mass scale for $SU(2)_L$ gauge bosons $W_L^\pm$ and Z_L . The VEVs must obey the hierarchy $v_R \gg \kappa_{1,2} \gg v_L$, with

$$\kappa_+^2 \equiv \kappa_1^2 + \kappa_2^2 \simeq (246 \text{ GeV})^2 \quad (3)$$

defining the electroweak scale, to satisfy constraints from low-energy weak interactions. Here, $\langle \delta_L^0 \rangle = v_L/\sqrt{2}$ is the triplet VEV that cannot exceed $\mathcal{O}(\text{GeV})$ due to electroweak precision measurements, and in particular, from the ρ -parameter constraint [29–31].¹ Similarly, $v_R \gtrsim 20 \text{ TeV}$ in order to satisfy the current LHC constraints on the heavy gauge boson masses [34,35].

¹Here we do not consider the collider detector at fermilab anomalous W -mass measurement [32], which is in disagreement with the recent compact muon solenoid measurement [33].

For the theory to be LR symmetric, the Lagrangian must be invariant under parity (discrete LR symmetry)²:

$$Q_L \leftrightarrow Q_R, \quad \Psi_L \leftrightarrow \Psi_R, \quad \Delta_L \leftrightarrow \Delta_R, \quad \Phi \leftrightarrow \Phi^\dagger. \quad (4)$$

The relevant terms in the Lagrangian are

$$\mathcal{L}_Y = \bar{Q}_L(h_q\Phi + \tilde{h}_q\tilde{\Phi})Q_R + \bar{\psi}_L(h_\ell\Phi + \tilde{h}_\ell\tilde{\Phi})\psi_R + f[\psi_R^T C i\sigma_2 \Delta_R \psi_R + \psi_L^T C i\sigma_2 \Delta_L \psi_L] + \text{H.c.}, \quad (5)$$

where $\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2$ with the second Pauli matrix σ_2 . The mass matrices for charged leptons (M_ℓ) and Dirac neutrinos (M_{ν_D}), as well as the 6×6 mass matrix for the $\nu - N$ sector (with $\nu \equiv \nu_L$ and $N \equiv C(\bar{\nu}_R)^T$), are

$$M_{\ell,\nu_D} = \frac{h_\ell \kappa_{2,1} + \tilde{h}_\ell \kappa_{1,2}}{\sqrt{2}}, \quad M_\nu = \begin{pmatrix} \sqrt{2}f v_L & M_{\nu_D} \\ M_{\nu_D}^T & \sqrt{2}f v_R \end{pmatrix}. \quad (6)$$

In the limit $M_{\nu_D} \ll f v_R$, the states ν and N will approximately be mass eigenstates with mass matrices,

$$m_N \simeq \sqrt{2}f v_R, \quad m_\nu \simeq \sqrt{2}f v_L - M_{\nu_D} m_N^{-1} M_{\nu_D}^T. \quad (7)$$

Here, m_ν is the Majorana mass matrix for the light, mostly left-handed neutrinos, containing contributions from both type-II (the v_L term) and type-I (the m_N^{-1} term) seesaw.

The Higgs sector of the model contains 20 degrees of freedom: eight within neutral scalars, eight within singly charged scalars, and four within doubly-charged scalars. After LR symmetry is spontaneously broken down to $SU(3)_c \times U(1)_{em}$, six Goldstone modes are eaten up by the $W_{L,R}^\pm$ and $Z_{L,R}$ gauge bosons. We analyze the full scalar potential in Appendix A to identify all the Goldstone modes and give the masses of the physical scalar fields. The upshot is that one of the neutral scalars can be sufficiently light and long lived, thus opening the possibility to make it a decaying DM candidate.

III. THE $SU(2)_R$ -TRIPLT HIGGS AS DARK MATTER

We wish to identify the neutral scalar component of the RH triplet Higgs $\delta_R^0 \equiv \chi$ as a prospective DM candidate.

²It is also possible to define a discrete charge conjugation symmetry C within the model under which the scalar and fermion fields transform into their conjugate fields: $Q_{L,R} \leftrightarrow (Q^c)_{R,L}$, $\Psi_{L,R} \leftrightarrow (\Psi^c)_{R,L}$, $\Phi \leftrightarrow \Phi^*$, $\Delta_{L,R} \rightarrow (\Delta^c)_{R,L}$. If this symmetry is adopted, the Yukawa couplings h_q , $\tilde{h}_q$, h_ℓ , and $\tilde{h}_\ell$ would be complex symmetric matrices, rather than Hermitian matrices, with little impact on our discussion.

Given its involvement in spontaneous symmetry breaking and the Yukawa couplings in Eq. (5), χ can decay into SM particles. Requiring the DM lifetime to be longer than the age of the Universe is easiest to achieve if χ is very light. For now, we stay agnostic about the DM production mechanism, some choices being discussed below. For simplicity we do, however, assume a (semi) thermal mechanism rather than condensation or misalignment [36]. As a result, the DM mass should exceed $\mathcal{O}(100 \text{ eV})$, likely even several keV, to satisfy structure-formation constraints.

We define the mixing angles of $\delta_R^0 \equiv \chi$ with h_1^{0r} and h_2^{0r} as θ_1 and θ_2 , respectively [cf. Eqs. (A10) and (A11)]. The fields (h_1^{0r}, h_2^{0r}) in the rotated basis are defined in Eq. (A8) such that h_1^{0r} is approximately identified as the SM-like Higgs boson h and h_2^{0r} as the heavy CP even scalar H . In the limit $2\rho_1 v_R^2 \ll m_h^2, m_H^2$, the mass of χ can be written as

$$m_\chi^2 \simeq 2\rho_1 v_R^2 - \theta_1^2 m_h^2 - \theta_2^2 m_H^2, \quad (8)$$

where ρ_1 is one of χ 's quartic couplings [cf. Eq. (A1)]. Generically, χ will thus have a mass around v_R , so bringing it down to the electroweak scale requires considerable fine-tuning in the form of tiny parameters or cancellations. This is the price we pay for this DM candidate. Rather than canceling these terms against each other, we assume all of them to be minuscule, so $\rho_1, \theta_{1,2} \ll 1$. For concreteness, we effectively set $\theta_{1,2} = 0$ in the following.

Even for vanishing mixing angles of χ with the other two neutral scalars, it can decay into a pair of photons at the one-loop level via virtual charged scalars H^+ , δ_L^+ , and $\delta_{L,R}^{++}$, as well as the gauge boson W_R^+ , as shown in Fig. 1. In the low-mass limit $m_\chi \ll v_R$ of interest to us, the decay width reads (see Ref. [37] and Appendix B)

$$\Gamma_{\chi \rightarrow \gamma\gamma} \simeq \frac{\alpha^2 m_\chi^3}{1028\pi^3 v_R^2} \left| \frac{1}{3} + \frac{1}{3} + \frac{4}{3} + \frac{4}{3} - 7 \right|^2, \quad (9)$$

where $\alpha \equiv e^2/4\pi$ is the electromagnetic fine-structure constant, and the factors in brackets represent the contributions from the H^+ , δ_L^+ , δ_R^{++} , δ_L^{++} , and W_R^+ loops, respectively.³ We refrain from discussing other possible decay channels such as $\chi \rightarrow \nu\nu$ or $\chi \rightarrow e^+e^-$ because they are less important in the mass range of interest.

If χ makes up all DM in our Universe, the diphoton decay has to be enormously suppressed to satisfy the stringent experimental limits [38].⁴ This can be achieved by an extreme hierarchy $m_\chi \ll v_R$. The lower limit

³Had we kept the mixing angles $\theta_{1,2}$ there would be additional contributions and even the possibility to fine-tune $\Gamma_{\chi \rightarrow \gamma\gamma} \simeq 0$, which could make possible a much lower v_R scale.

⁴This turns out to be a much stronger requirement than simply the cosmological stability of the DM.

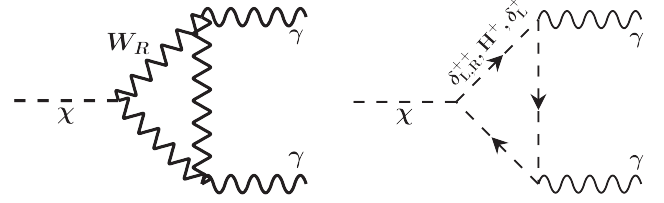


FIG. 1. Loop-induced decay of scalar DM into photons through gauge bosons $W_R^\pm$ and charged scalars $\delta_{L,R}^{++}$, δ_L^+ , and H^+ in the LR model.

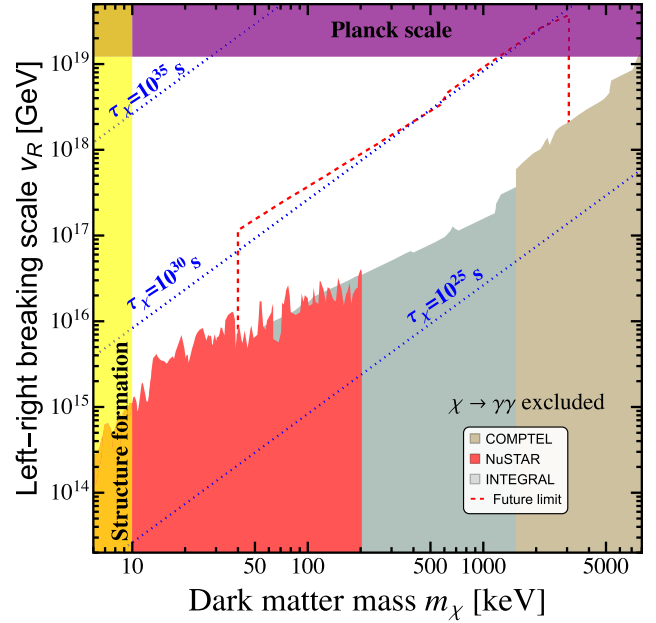


FIG. 2. Allowed parameter space (unshaded) for the DM mass m_χ as a function of the LR breaking scale v_R . The shaded regions are excluded. X-ray constraints are obtained from NuSTAR [41,43] (red), INTEGRAL [42,44] (gray), and COMPTEL [40] (brown). The red dashed line represents future X-ray reach [38,45]. The yellow region corresponds roughly to the excluded region due to Lyman-alpha data, assuming a (semi)thermal DM production mechanism [39]. The purple shaded region on top corresponds to $v_R > M_{\text{Pl}}$, where quantum gravity effects might invalidate our perturbative calculation. Shown in dotted blue are contours of the DM lifetime $\chi \rightarrow \gamma\gamma$.

$m_\chi \gtrsim 10 \text{ keV}$ from structure formation [39] forces v_R above 10^{15} GeV to satisfy X-ray line search limits [40–44], illustrated in Fig. 2. While such a high LR breaking scale is not unfamiliar from grand unified theories, it exacerbates the fine-tuning needed to make χ light. On the other hand, for $v_R < M_{\text{Pl}}$, we must have $m_\chi \lesssim 10 \text{ MeV}$ to satisfy the gamma-ray constraints.

We discuss the scalar potential and the conditions for stable parameters $\rho_1, \theta_{1,2} \ll 1$ in Appendixes A and C, respectively. ρ_1 , one of the quartic couplings of Δ_R , is particularly important in this setup,

$$\rho_1 \simeq \frac{m_\chi^2}{2v_R^2} \simeq 5 \times 10^{-43} \left(\frac{m_\chi}{100 \text{ keV}} \right)^2 \left(\frac{10^{17} \text{ GeV}}{v_R} \right)^2. \quad (10)$$

It is chosen to be near zero at the scale m_χ , and is thus highly sensitive to renormalization-group running effects. Loop corrections coming from bosonic loops will increase ρ_1 , but the contributions from the RH neutrinos N drive ρ_1 to *negative* values, which would destabilize the vacuum. This behavior is of course completely analogous to the hierarchy problem and vacuum stability issue in the SM [46,47]. If we assume for now that all non-SM bosons except χ have masses near the v_R scale, ρ_1 remains constant up to the mass scale m_N corresponding to the lightest N , and then picks up a negative contribution

$$\Delta\rho_1 \simeq -\frac{16|f|^4}{(4\pi)^2} \log\left(\frac{v_R}{m_N}\right) \quad (11)$$

on the way to v_R , coming from the loop contribution to the quartic coupling shown in Fig. 3. If $\rho_1 + \Delta\rho_1 > 0$ up to v_R , it will remain positive even at larger scales due to the large positive contributions from the heavy bosons. But to ensure this, N should either sit near, or ideally above, v_R as well (to avoid ρ_1 running altogether), or have a tiny Yukawa coupling f , roughly $|f| < \sqrt{\pi}\rho_1^{1/4}$, which translates into $m_N \lesssim \sqrt{m_\chi v_R}$. Here we assumed only one N ; if instead we have three degenerate N with same coupling f , the bound is marginally modified to $m_N \lesssim \sqrt{m_\chi v_R}/\sqrt{3}$, which makes little difference to our analysis.

Vacuum stability of this DM setup thus forbids the RH neutrino mass region

$$\sqrt{m_\chi v_R} \lesssim m_N \lesssim v_R. \quad (12)$$

Even lowering the masses of some of the charged scalars below v_R to utilize their positive contributions to ρ_1 does not change this conclusion, or at least our numerical scans did not reveal any viable scenarios. The new scalars can, however, cure the *electroweak* vacuum stability problem, as shown in Appendix D. Since this has little to no bearing on our DM phenomenology, we will not discuss it further. Below, we discuss the two viable mass windows for N and

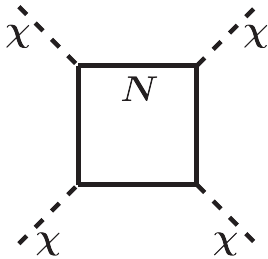


FIG. 3. One-loop contribution to the quartic-coupling constant ρ_1 from the Yukawa coupling f of Eq. (5).

possible ways to generate the correct DM abundance, neutrino masses, and the baryon asymmetry of our Universe.

IV. LIGHT RIGHT-HANDED NEUTRINOS

In the parameter region $m_N < \sqrt{m_\chi v_R}$, the RH neutrinos are far below the LR scale, but can still be far above the electroweak scale,

$$m_N \lesssim 3 \times 10^6 \text{ GeV} \left(\frac{m_\chi}{100 \text{ keV}} \right)^{\frac{1}{2}} \left(\frac{v_R}{10^{17} \text{ GeV}} \right)^{\frac{1}{2}}. \quad (13)$$

In this region of parameter space we can explain the active neutrino masses via type-I seesaw by adjusting the Dirac-mass matrix to be of order

$$M_{\nu_D} \simeq \sqrt{m_\nu m_N} \simeq 20 \text{ MeV} \left(\frac{m_\nu}{0.1 \text{ eV}} \right)^{\frac{1}{2}} \left(\frac{m_N}{3 \times 10^6 \text{ GeV}} \right)^{\frac{1}{2}}. \quad (14)$$

In addition, the type-II seesaw contribution could be sizeable despite the small $|f| < \sqrt{\pi}\rho_1^{1/4}$,

$$m_\nu^{II} \simeq \sqrt{2}f v_L \simeq 0.07 \text{ eV} \left(\frac{m_\chi}{100 \text{ keV}} \right)^{\frac{1}{2}} \left(\frac{10^{17} \text{ GeV}}{v_R} \right)^{\frac{1}{2}} \left(\frac{v_L}{\text{GeV}} \right). \quad (15)$$

Given the upper limit on v_L of order GeV from electroweak precision measurements [30,31], there is a sizable region of parameter space in Fig. 2 where most or even all of m_ν comes from the type-II seesaw mechanism. For very large v_R and small m_χ , type-I seesaw is going to be dominant though.

If the Universe reached temperatures around v_R , the scalar DM would have been thermalized through the RH gauge interactions. Assuming all triplet partners to have masses around v_R , the real scalar χ would have frozen out at a temperature T_f just below v_R while being relativistic, giving an abundance today [48]

$$\Omega_\chi h^2 \simeq 0.12 \left(\frac{106.75}{g_*(T_f)} \right) \left(\frac{m_\chi}{170 \text{ eV}} \right). \quad (16)$$

For $m_\chi \sim 170 \text{ eV}$, this would match the observed DM abundance $\Omega h^2 \simeq 0.12$ [49] but be *hot* DM, completely excluded. For $m_\chi > 170 \text{ eV}$, it would overclose the Universe, also excluded. The same problem occurs for the LR keV sterile-neutrino DM case—only the abundance is increased by $2 \times 3/4$ due to the fermionic degrees of freedom. Let us discuss a few possible solutions:

- (1) Entropy dilution through g_* : The DM abundance in Eq. (16) can equal the observed one even for $m_\chi \gg 170 \text{ eV}$ if we assume additional degrees of freedom

at the time of DM freeze-out, i.e., an increased $g_*(T_f)$. This solution is not particularly appealing since it requires the addition of many more particles beyond those in our LR model [17].

- (2) Entropy dilution from late decays: The relatively light RH neutrinos will be in thermal equilibrium around v_R and freeze out relativistically around a similar temperature T_f as χ since both couple to the $SU(2)_R$ gauge bosons. If the lifetime of one of the N is sufficiently long, it can dominate the Universe for a while once $T \lesssim m_N$, and then decay, thus reheating the Universe to a temperature [50]

$$T_{\text{reh}} \simeq 0.8(g_*(T_{\text{reh}}))^{-\frac{1}{4}} \sqrt{\Gamma_N M_{\text{Pl}}} \quad (17)$$

$$\simeq 700 \text{ GeV} \left(\frac{g_*(T_{\text{reh}})}{106.75} \right)^{-\frac{1}{4}} \left(\frac{10^{-12} \text{ s}}{\tau_N} \right)^{\frac{1}{2}}, \quad (18)$$

which effectively dilutes the DM temperature because N does not decay into DM—unlike in the keV sterile neutrino case [19,20]. This can give cold DM with the right abundance as long as m_N and τ_N are picked correctly so that enough entropy is produced (long N lifetime) while keeping the reheating temperature above MeV for successful big bang nucleosynthesis (short N lifetime). Concretely, Ω_χ is reduced by a factor [17,18]

$$S \simeq 80 \left(\frac{106.75}{g_*(T_{\text{reh}})} \right)^{\frac{3}{4}} \left(\frac{\tau_N}{10^{-12} \text{ s}} \right)^{\frac{1}{2}} \left(\frac{m_N}{10^7 \text{ GeV}} \right), \quad (19)$$

thus allowing for $m_\chi > 170 \text{ eV}$, and its temperature is reduced by a factor $S^{1/3}$, making it much colder. As a result, the Lyman- α limits in this case are those of a thermal relic [17], excluding masses below 4 keV [51]. We then need at least $S \gtrsim 25$. The required N lifetime is far longer than expected from the seesaw relation in Eq. (14), so we need to assume that one of the three N has much smaller Yukawa couplings, leading to a longer lifetime and one effectively massless active neutrino compatible with current data.⁵ Despite the upper bound on m_N from Eq. (13), there is even a region of parameter space in which all N decay before sphaleron decoupling, allowing for leptogenesis. The small masses compared to the Davidson-Ibarra bound [52] imply generically too small a CP asymmetry, but two or more degenerate N could lead to resonant leptogenesis to circumvent this problem [53,54]. Even thermal leptogenesis could work all the way down to 10^6 GeV if we consider flavor effects [55].

⁵The type-II seesaw contribution can and will lift the lightest neutrino mass.

- (3) Reheating temperature T_{reh} below v_R : If we obtain too much DM for $T_{\text{reh}} > v_R$ and too little DM for $T_{\text{reh}} \ll v_R$ —seeing as χ only interacts weakly with the SM through v_R -suppressed operators—there must be a sweet spot $T_{\text{reh}} < v_R$ where the correct DM abundance is produced for any m_χ . In this UV freeze-in [56] case, DM is produced with momenta $\langle p \rangle \sim 3T$, so masses below $\sim 10 \text{ keV}$ are excluded by Lyman α [51]. Since the RH neutrinos have by and large the same interactions as χ , they should generically be produced by the same amount at T_{reh} . They stay relativistic until $T \sim m_N$ and then decay. For seesaw-inspired M_{ν_D} values, the decay will be too fast to lead to a period of matter domination. The lepton asymmetry produced in their decay (subsequently converted to a baryon asymmetry by electroweak sphalerons) is generically too small to explain the matter asymmetry since the small m_N suppresses the CP asymmetry ϵ . However, *resonant* leptogenesis [53], i.e., two or more almost degenerate RH neutrinos, could work in this scenario [54], at the price of additional fine-tuning.

Overall, the light N window of the scalar-DM parameter space offers numerous ways to simultaneously explain neutrino masses, the DM abundance, and even leptogenesis. Let us now discuss the heavy- N scenario.

V. HEAVY RIGHT-HANDED NEUTRINOS

With $m_N > v_R$ and $v_R \gtrsim 10^{15} \text{ GeV}$ from Fig. 1, the type-I seesaw contribution to active neutrino masses is suppressed,

$$m_\nu^I \lesssim 0.04 \text{ eV} \left(\frac{M_{\nu_D}}{200 \text{ GeV}} \right)^2 \left(\frac{10^{15} \text{ GeV}}{m_N} \right), \quad (20)$$

and could at most account for the two lightest neutrino masses unless we push the Yukawa couplings near the nonperturbative regime. In this region of parameter space, neutrino masses have to come from type-II seesaw, and since f is of order one in this section we require a tiny $v_L \sim m_\nu$.

For $m_N > v_R$, we can still obtain the correct DM abundance through UV freeze-in or by severely increasing $g_*(T_f)$, but the late-entropy solution is no longer viable because the N will freeze-out nonrelativistically and thus cannot dump enough entropy into the SM bath to solve the overproduction problem. The minimal LR model does not contain any other particles that could lead to late entropy production. In that case, the heavy N were never produced and thus cannot be used for leptogenesis. Upon introduction of additional particles, one could arrange to produce N despite the low temperature, e.g., via inflaton decays [57] or bubble collisions [58], which could then lead to leptogenesis. In this case it is important to note that the large

coupling f of χ to N means that every N decay/production comes with an $\mathcal{O}(1)$ number of χ particles, e.g.,

$$\frac{\Gamma(N \rightarrow LH\chi)}{\Gamma(N \rightarrow LH)} \simeq \frac{|f|^2}{2\pi^2} \log \left[\frac{M_N}{m_\chi} \right], \quad (21)$$

enhanced by an infrared divergence since the virtual N goes on-shell for $m_\chi \rightarrow 0$ (and $p_\chi \rightarrow 0$, so the χ will be produced with very small momentum [39]). We will omit a discussion of such scenarios because they go beyond the minimal LR setup.

VI. CONCLUSION

The left-right model has been a popular extension of the SM for half a century. While it naturally explains neutrino masses and can accommodate the baryon asymmetry of the Universe, dark matter has proven elusive for a long time. A decade or so ago, it was noticed [18] that one of the right-handed neutrinos within the LR model can be tuned light enough to be stable on cosmological timescales, allowing for it to be viable DM. Here, we showed that the LR model contains yet another viable DM candidate in the form of the LR-breaking Higgs. With admittedly severe fine-tuning, it can be made light and stable enough to give DM. Vacuum stability imposes nontrivial constraints on the RH neutrino masses in this scenario, which however still allow for successful neutrino masses and leptogenesis. This DM candidate likely remains viable when we embed the LR model into Pati-Salam or even $SO(10)$, providing a completely different way to get DM there, as opposed to the

weakly interacting massive particle-like stable DM candidates proposed before [14–16]. Observing an X-ray or gamma-ray line signal from DM-rich regions would be a possible indication of this DM candidate. Finding a W_R boson at the LHC or future collider would exclude this DM model.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: SCALAR SECTOR

We analyze the Higgs potential of the model to see if a light $\Re(\Delta_R^0)$ is viable. The most general renormalizable Higgs potential involving Φ , Δ_L , and Δ_R fields is given by [59]

$$\begin{aligned} V(\Phi, \Delta_{L,R}) = & -\mu_1^2 \text{Tr}(\Phi^\dagger \Phi) - \mu_2^2 [\text{Tr}(\Phi^\dagger \tilde{\Phi} + \text{Tr}(\Phi \tilde{\Phi}^\dagger))] - \mu_3^2 [\text{Tr}(\Delta_L^\dagger \Delta_L) + \text{Tr}(\Delta_R^\dagger \Delta_R)] + \lambda_1 \text{Tr}(\Phi^\dagger \Phi)^2 \\ & + \lambda_2 [\text{Tr}(\Phi^\dagger \tilde{\Phi})^2 + \text{Tr}(\Phi \tilde{\Phi}^\dagger)^2] + \lambda_3 \text{Tr}(\Phi^\dagger \tilde{\Phi}) \text{Tr}(\Phi \tilde{\Phi}^\dagger) + \lambda_4 \text{Tr}(\Phi^\dagger \Phi) [\text{Tr}(\Phi \tilde{\Phi}^\dagger) + \text{Tr}(\Phi^\dagger \tilde{\Phi})] \\ & + \rho_1 [\text{Tr}(\Delta_L^\dagger \Delta_L)^2 + \text{Tr}(\Delta_R^\dagger \Delta_R)^2] + \rho_2 [\text{Tr}(\Delta_L \Delta_L) \text{Tr}(\Delta_L^\dagger \Delta_L) + \text{Tr}(\Delta_R \Delta_R) \text{Tr}(\Delta_R^\dagger \Delta_R)] \\ & + \rho_3 \text{Tr}(\Delta_L \Delta_L^\dagger) \text{Tr}(\Delta_R \Delta_R^\dagger) + \rho_4 [\text{Tr}(\Delta_L \Delta_L) \text{Tr}(\Delta_R^\dagger \Delta_R) + \text{Tr}(\Delta_L^\dagger \Delta_L) \text{Tr}(\Delta_R \Delta_R)] \\ & + \alpha_1 \text{Tr}(\Phi^\dagger \Phi) [\text{Tr}(\Delta_L^\dagger \Delta_L) + \text{Tr}(\Delta_R^\dagger \Delta_R)] + [\alpha_2 \{ \text{Tr}(\tilde{\Phi} \Phi^\dagger) \text{Tr}(\Delta_L^\dagger \Delta_L) + \text{Tr}(\tilde{\Phi}^\dagger \Phi) \text{Tr}(\Delta_R^\dagger \Delta_R) \} + \text{H.c.}] \\ & + \alpha_3 [\text{Tr}(\Phi \Phi^\dagger \Delta_L \Delta_L^\dagger) + \text{Tr}(\Phi^\dagger \Phi \Delta_R \Delta_R^\dagger)] + \beta_1 [\text{Tr}(\Phi \Delta_R \Phi^\dagger \Delta_L^\dagger) + \text{Tr}(\Phi^\dagger \Delta_L \Phi \Delta_R^\dagger)] \\ & + \beta_2 [\text{Tr}(\tilde{\Phi} \Delta_R \Phi^\dagger \Delta_L^\dagger) + \text{Tr}(\tilde{\Phi}^\dagger \Delta_L \Phi \Delta_R^\dagger)] + \beta_3 [\text{Tr}(\Phi \Delta_R \tilde{\Phi}^\dagger \Delta_L^\dagger) + \text{Tr}(\Phi^\dagger \Delta_L \tilde{\Phi} \Delta_R^\dagger)]. \end{aligned} \quad (A1)$$

Here all the couplings except for α_2 are made real by field redefinitions [59]. Furthermore, for simplicity, we take the phase of α_2 to be zero such that there is no CP violation in the scalar sector and no mixing between scalars and pseudoscalars. Moreover, we take all the VEVs to be real. After inserting the VEVs, we obtain the following conditions for the potential to be an extremum:

$$-\mu_1^2 + k_+^2 \lambda_1 + 2\kappa_1 \kappa_2 \lambda_4 + \frac{1}{2k_-^2} [2v_L v_R (\beta_2 \kappa_1^2 - \beta_3 \kappa_2^2) + (v_L^2 + v_R^2)(\alpha_1 k_-^2 - \alpha_3 \kappa_2^2)] = 0, \quad (A2)$$

$$-\mu_2^2 + \kappa_1 \kappa_2 (2\lambda_2 + \lambda_3) + \frac{1}{2} \lambda_4 k_+^2 + \frac{1}{4k_-^2} [v_L v_R (\beta_1 k_-^2 - 2\kappa_1 \kappa_2 (\beta_2 - \beta_3)) + (v_L^2 + v_R^2)(2\alpha_2 k_-^2 + \alpha_3 \kappa_1 \kappa_2)] = 0, \quad (A3)$$

$$-\mu_3^2 + \frac{1}{2} [\alpha_1 k_+^2 + 4\alpha_2 \kappa_1 \kappa_2 + \alpha_3 \kappa_2^2 + 2\rho_1 v_R^2 + 2\rho_1 (v_L^2 + v_R^2)] = 0, \quad (A4)$$

$$\beta_2 \kappa_1^2 + \beta_1 \kappa_1 \kappa_2 + \beta_3 \kappa_2^2 - v_L v_R (2\rho_1 - \rho_3) = 0. \quad (A5)$$

The last equation, Eq. (A5), is the VEV seesaw relation. Here and in what follows we take $\kappa_1 \geq \kappa_2$ without loss of generality and define

$$\kappa_{\pm}^2 \equiv \kappa_1^2 \pm \kappa_2^2. \quad (\text{A6})$$

Using the minimization conditions to eliminate μ_i^2 , we construct the mass matrix for the neutral Higgs bosons in the basis $\{\phi_1^{0r}, \phi_2^{0r}, \delta_R^{0r}, \delta_L^{0r}\}$ by expanding the potential about the minimum in VEVs to quadratic order. Here the basis $\{\phi_1^0, \phi_2^0, \delta_R^0, \delta_L^0\}$ consisting of complex fields is broken up into real and imaginary components; for instance $\phi_1^0 = (\phi_1^{0r} + i\phi_1^{0i})/\sqrt{2}$. To identify the SM-like Higgs, we use the following orthogonal matrix to transform the original basis to an intermediate flavor-diagonal basis:

$$O_r = \begin{pmatrix} \kappa_1/\kappa_+ & \kappa_2/\kappa_+ & 0 & 0 \\ -\kappa_2/\kappa_+ & \kappa_1/\kappa_+ & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}. \quad (\text{A8})$$

The symmetric mass matrix in the rotated basis $\{h_1^{0r}, h_2^{0r}, \delta_R^{0r}, \delta_L^{0r}\}$ with $v_L \simeq 0$ (for simplicity) then reads

$$\tilde{M}_r^2 = \begin{pmatrix} 2\lambda_1\kappa_+^2 + 8\lambda_4\kappa_1\kappa_2 + 8(2\lambda_2 + \lambda_3)\frac{\kappa_1^2\kappa_2^2}{\kappa_+^2} & \bullet & \bullet & \bullet \\ 2\lambda_4\kappa_-^2 + 4(2\lambda_2 + \lambda_3)\kappa_1\kappa_2\frac{\kappa_-^2}{\kappa_+^2} & 2(2\lambda_2 + \lambda_3)\frac{\kappa_-^4}{\kappa_+^2} + \alpha_3\frac{\kappa_+^2}{2\kappa_-^2}v_R^2 & \bullet & \bullet \\ \{\alpha_1\kappa_+^2 + (4\alpha_2\kappa_1 + \alpha_3\kappa_2)\kappa_2\}\frac{v_R}{\kappa_+} & (2\alpha_2\kappa_-^2 + \alpha_3\kappa_1\kappa_2)\frac{v_R}{\kappa_+} & 2\rho_1v_R^2 & \bullet \\ 0 & \frac{1}{2\kappa_1}(\beta_1\kappa_1 + 2\beta_3\kappa_2)\kappa_+v_R & 0 & \frac{1}{2}(\rho_3 - 2\rho_1)v_R^2 \end{pmatrix}. \quad (\text{A9})$$

The δ_L^{0r} field decouples in the limit $\beta_i \simeq 0$ or $\beta_1 = 2\beta_3\kappa_2/\kappa_1$. Note that such a choice is fully consistent with $v_L \simeq 0$ as evident by the seesaw relation of Eq. (A5). The field h_1^{0r} is the SM-like Higgs boson in the limit of small mixing with other states. The remaining 3×3 matrix can be further diagonalized to obtain the mass eigenvalues of the real scalar fields. Here we wish to identify an additional state, i.e., δ_R^{0r} light such that it is a DM candidate. We use the following transformation matrix to relate the flavor-diagonal basis $\{h_1^{0r}, h_2^{0r}, \delta_R^{0r}\}$ and mass eigenstates of the 3×3 real scalar matrix.

$$\begin{pmatrix} h \\ H \\ \chi \end{pmatrix} = \begin{pmatrix} \cos\theta_1 & 0 & \sin\theta_1 \\ 0 & 1 & 0 \\ -\sin\theta_1 & 0 & \cos\theta_1 \end{pmatrix} \times \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos\theta_2 & \sin\theta_2 \\ 0 & -\sin\theta_2 & \cos\theta_2 \end{pmatrix} \begin{pmatrix} h_1^{0r} \\ h_2^{0r} \\ \delta_R^{0r} \end{pmatrix}. \quad (\text{A10})$$

Here, $\sin\theta_{1,2}$ denotes the mixing between χ with h and H . We ignored the mixing between h and H since they are of order $\mathcal{O}(1/v_R^2)$, thus heavily suppressed. In the limit of small mixing angles, the DM mass is given by Eq. (8), where the mixing angles are given by

$$\theta_1 = \frac{C_1 \frac{v_R}{\kappa_+}}{m_h^2 - m_\chi^2}, \quad \theta_2 = \frac{(2\alpha_2\kappa_-^2 + \alpha_3\kappa_1\kappa_2)\frac{v_R}{\kappa_+}}{m_H^2 - m_\chi^2}, \quad (\text{A11})$$

with $C_1 = \alpha_1\kappa_+^2 + (4\alpha_2\kappa_1 + \alpha_3\kappa_2)\kappa_2$. In order for the mass of χ to be at or below the electroweak scale and its mixing with other Higgs fields negligible, one needs $\rho_1 \sim \mathcal{O}(1/v_R^2)$ and

$$\alpha_1 \simeq \frac{\alpha_3\kappa_2^2}{\kappa_-^2}, \quad \alpha_2 \simeq -\frac{\alpha_3\kappa_1\kappa_2}{2\kappa_-^2}. \quad (\text{A12})$$

In the pseudoscalar sector, we can make a similar rotation from the original basis $\{\phi_1^{0i}, \phi_2^{0i}, \delta_R^{0i}, \delta_L^{0i}\}$ to the mass basis $\{A, G_L^0, G_R^0, \delta_L^{0i}\}$, where G_L^0 and $G_R^0 (\equiv \delta_R^{0i})$ are the Goldstone modes associated with neutral gauge bosons $Z_{L,R}$ with $G_L^0 \equiv (-\kappa_1\phi_1^{0i} + \kappa_2\phi_2^{0i})/\kappa_+$. The mass of the physical pseudoscalars (in the limit $\beta_i \simeq 0$) then read as

$$m_A^2 = -2(2\lambda_2 - \lambda_3)\kappa_+^2 + \frac{1}{2}\alpha_3\frac{\kappa_+^2}{\kappa_-^2}v_R^2, \quad m_{\delta_L^{0i}}^2 = \frac{1}{2}(\rho_3 - 2\rho_1)v_R^2. \quad (\text{A13})$$

The singly charged scalars $\{\phi_1^+, \phi_2^+, \delta_R^+, \delta_L^+\}$ can be also rotated to the mass basis $\{H^+, G_L^+, G_R^+, \delta_L^+\}$. Here $G_{L,R}^+$ are the Goldstone modes associated with neutral gauge bosons $W_{L,R}^+$ and given by

$$G_L^+ = -\frac{-\kappa_1\phi_1^+ + \kappa_2\phi_2^+}{\kappa_+}, \quad G_R^+ = \frac{-\kappa_-^2\phi_1^+ + \sqrt{2}\kappa_2v_R\delta_R^+}{\sqrt{\kappa_-^4 + 2\kappa_2^2v_R^2}}. \quad (\text{A14})$$

TABLE I. Couplings of scalar DM χ in the limit of $\beta_i = v_L = 0$ and $v_R \gg \kappa_1, \kappa_2$ in the minimal LR model. Only the terms up to order $\mathcal{O}(\theta_1)$ and $\mathcal{O}(\theta_2)$ are kept. Here, we define $C_2 = (\alpha_1 + \alpha_3)\kappa_1^2 - 4\alpha_2\kappa_2\kappa_1 + \alpha_1\kappa_2^2$. C_1 is defined through θ_1 below Eq. (A11) and vanishes upon substitution of Eq. (A12).

Interaction	Couplings
χ^4	$\rho_1/4$
χ^3	$v_R\rho_1 - \theta_1 C_1/(2\kappa_+)$
$\chi\chi hh$	$C_1/(4\kappa_+^2)$
$\chi\chi h$	$C_1/(2\kappa_+) + \theta_1 v_R(3\rho_1 - \alpha_1 - \kappa_2(4\alpha_2\kappa_1 + \alpha_3\kappa_2)/\kappa_+^2)$
χhh	$v_R C_1/(2\kappa_+^2)$
$\chi\chi HH$	$C_2/(4\kappa_+^2)$
$\chi\chi H$	$-(\alpha_3\kappa_1\kappa_2 + 2\alpha_2\kappa_-^2)(\kappa_+^2 - 2\theta_1 v_R\kappa_+)/ (2\kappa_+^3)$
χHH	$v_R C_2/(2\kappa_+^2)$
$\chi\chi\delta_L^{0r}\delta_L^{0r}$	$\rho_3/4$
$\chi\delta_L^{0r}\delta_L^{0r}$	$v_R\rho_3/2 - \theta_1 C_1/(2\kappa_+)$
$\chi\chi AA$	$C_2/(4\kappa_+^2)$
χAA	$v_R C_2/(2\kappa_+^2)$
$\chi\chi\delta_L^{0i}\delta_L^{0i}$	$\rho_3/4$
$\chi\delta_L^{0i}\delta_L^{0i}$	$v_R\rho_3/2 - \theta_1 C_1/(2\kappa_+)$
$\chi\chi H^+ H^-$	$(C_2 v_R^2 + \kappa_-^4 \rho_1)/(2v_R^2 \kappa_+^2)$
$\chi H^+ H^-$	$(2C_2 v_R^2 + \alpha_3(\kappa_1^4 - \kappa_-^4) + 2\kappa_-^4 \rho_1)/(2v_R^2 \kappa_+^2)$
$\chi\chi\delta_L^+ \delta_L^-$	$\rho_3/2$
$\chi\delta_L^+ \delta_L^-$	$-\theta_1\kappa_+(2\alpha_1 + \alpha_3)/2 + v_R\rho_3 - 4\theta_1\alpha_2\kappa_1\kappa_2/\kappa_+$
$\chi\chi\delta_R^{++}\delta_R^{--}$	$\rho_1 + 2\rho_2$
$\chi\delta_R^{++}\delta_R^{--}$	$2v_R(\rho_1 + 2\rho_2) - \theta_1((\alpha_1 + \alpha_3)\kappa_1^2 + 4\alpha_2\kappa_2\kappa_1 + \alpha_1\kappa_2^2)/\kappa_+$
$\chi\chi\delta_L^{++}\delta_L^{--}$	$\rho_3/2$
$\chi\delta_L^{++}\delta_L^{--}$	$v_R\rho_3 - \theta_1((\alpha_1 + \alpha_3)\kappa_1^2 + 4\alpha_2\kappa_2\kappa_1 + \alpha_1\kappa_2^2)/\kappa_+.$

The orthogonal state to $\{G_L^+, G_R^+\}$ corresponds to the physical field H^+ . The masses of H^+ and δ_L^+ fields are

$$m_{H^+}^2 = \frac{1}{4}\alpha_3\left\{\kappa_-^2 + 2\frac{\kappa_+^2}{\kappa_-^2}v_R^2\right\},$$

$$m_{\delta_L^+}^2 = \frac{1}{2}\left[\alpha_3\kappa_-^2 - v_R^2(2\rho_1 - \rho_3)\right]. \quad (\text{A15})$$

The masses of the doubly charged scalars ($\delta_R^{++}, \delta_L^{++}$) in the limit $\beta_i \simeq 0$ read as

$$\begin{aligned} \lambda_1 < 4\pi/3, \quad (\lambda_1 + 4\lambda_2 + 2\lambda_3) < 4\pi, \quad (\lambda_1 - 4\lambda_2 + 2\lambda_3) < 4\pi, \quad \lambda_4 < 4\pi/3, \quad \alpha_1 < 8\pi, \quad \alpha_2 < 4\pi, \\ (\alpha_1 + \alpha_3) < 8\pi, \quad \rho_1 < 4\pi/3, \quad (\rho_1 + \rho_2) < 2\pi, \quad \rho_2 < 2\sqrt{2}\pi, \quad \rho_3 < 8\pi, \quad \rho_4 < 2\sqrt{2}\pi. \end{aligned} \quad (\text{C1})$$

The quartic terms $V^{(4)}(\Phi, \Delta_R)$ form a vector space spanned by the real-valued vectors $x^T = \{\kappa_1^2, \kappa_2^2, v_R^2\}$, which can be written as $V^{(4)} = x^T \hat{\lambda}_x$ with

$$m_{\delta_R^{++}}^2 = \frac{1}{2}\alpha_3\kappa_-^2 + 2\rho_2 v_R^2,$$

$$m_{\delta_L^{++}}^2 = \frac{1}{2}\left[\alpha_3\kappa_-^2 - v_R^2(2\rho_1 - \rho_3)\right]. \quad (\text{A16})$$

For completeness we also show the relevant couplings of $\delta_R^0 \equiv \chi$ in the mass basis in Table I in the limit of $\beta_i = v_L = 0$ and $v_R \gg \kappa_1, \kappa_2$.

APPENDIX B: DECAY WIDTH OF SCALAR DM TO PHOTONS

The decay of scalar dark matter χ to a pair of photons is induced at one loop order through the gauge bosons $W_R^\pm$ and the charged scalars $\delta_{L,R}^{\pm\pm}, \delta_L^\pm$ and $H^\pm$ as shown in Fig. 1. The decay width reads as [37]

$$\Gamma(\chi \rightarrow \gamma\gamma) \simeq \frac{\alpha^2 m_\chi^3}{1028\pi^3} \frac{4}{v_R^2} \left| A_0(\tau_{H^\pm}) + A_0(\tau_{\delta_L^\pm}) + 4A_0(\tau_{\delta_R^{\pm\pm}}) + 4A_0(\tau_{\delta_L^{\pm\pm}}) + A_1(\tau_{W_R^\pm}) \right|^2, \quad (\text{B1})$$

and the loop functions are given by

$$A_0(\tau) \equiv -[\tau - f(\tau)]\tau^{-2}, \quad (\text{B2})$$

$$A_1(\tau) \equiv -[2\tau^2 + 3\tau + 3(2\tau - 1)f(\tau)]\tau^{-2}, \quad (\text{B3})$$

with $\tau_i = m_\chi^2/4m_i^2$ and

$$f(\tau) \equiv \begin{cases} \arcsin^2 \sqrt{\tau} & (\text{for } \tau \leq 1), \\ -\frac{1}{4} \left[\log \left(\frac{1+\sqrt{1-1/\tau}}{1-\sqrt{1-1/\tau}} \right) - i\pi \right]^2 & (\text{for } \tau > 1). \end{cases} \quad (\text{B4})$$

For heavy particles running inside the loop ($m_i \gg m_\chi$), these loop functions can be approximated to $A_0(0) \approx 1/3$ and $A_1(0) \approx -7$.

APPENDIX C: THEORETICAL CONSTRAINTS

The two-body scalar scattering through contact interaction at tree level gives the following unitarity constraints on the quartic couplings [60],

$$\hat{\lambda} = \frac{1}{4} \begin{pmatrix} \lambda_1 & \lambda' & \alpha_1/2 \\ \lambda' & \lambda_1 & (\alpha_1 + \alpha_3)/2 \\ \alpha_1/2 & (\alpha_1 + \alpha_3)/2 & \rho_1 \end{pmatrix}. \quad (\text{C2})$$

Here, $\lambda' = \lambda_1 + 4\lambda_2 + 2\lambda_3$. The necessary conditions for the boundedness can be derived from the copositivity of the real symmetric matrix $\hat{\lambda}$ of Eq. (C2), which are [61–63],

$$\begin{aligned} \lambda_1 \geq 0, \quad \rho_1 \geq 0, \quad 2\lambda_2 + \lambda_3 \geq -\lambda_1, \quad \alpha_1 \geq -2\sqrt{\lambda_1\rho_1}, \quad \alpha_1 + \alpha_3 \geq -2\sqrt{\lambda_1\rho_1}, \\ \lambda'\sqrt{\rho_1} + \frac{1}{2}(\alpha_1 + \alpha_3)\sqrt{\lambda_1} + \frac{1}{2}\alpha_1\sqrt{\lambda_1} + \sqrt{\lambda_1^2\rho_1} \geq 0 \quad \text{or} \quad \det \hat{\lambda} \geq 0. \end{aligned} \quad (\text{C3})$$

Note that the inclusion of v_L leads to a symmetric matrix of order four. To determine copositivity of such a matrix one would need to adjudge eight different cases depending on the sign distributions of the off-diagonal elements [63]. The full set of analytical expressions for vacuum stability that yield the desirable vacuum expectation values ensuring correct symmetry breaking to SM is analyzed in Ref. [64,65].

APPENDIX D: RENORMALIZATION GROUP EQUATIONS

Here we discuss the difficulty in maintaining a stable light- χ solution due to renormalization group equations (RGEs). For the potential to remain bounded from below, the quartic coupling must satisfy the conditions outlined in Eq. (C3). Specifically, the quartic coupling $\rho_1\chi^4$ needs to be exceedingly small, making it highly sensitive to any new interaction. Lowering the masses of one or more RH neutrinos below v_R can potentially destabilize the vacuum, as the Yukawa coupling f could drive the quartic coupling ρ_1 to negative values. The beta function for ρ_1 is given by $\beta(\rho_1) \propto -16f^4$ and picks up the negative contribution as illustrated in Eq. (11). A possible mitigation strategy is to keep scalar particles at or below m_N to offset the negative $\Delta\rho_1$ from N with a positive one from a scalar. However, we did not find any consistent solution up to two-loop that satisfies both the boundedness as well the vacuum stability conditions. The two-loop beta functions were generated with the `PyR@TE` package [66] and were numerically solved in *Mathematica* by introducing various scalar particles of the LR theory below m_N .

Alternatively, by choosing a tiny Yukawa coupling f , roughly $|f| < \sqrt{\pi}\rho_1^{1/4}$ or equivalently $m_N \lesssim \sqrt{m_\chi v_R}$, we can ensure that ρ_1 remains positive up to the v_R scale. For instance, a benchmark point with $m_\chi = 1$ MeV, $v_R = 10^{17}$ GeV, and $f \lesssim 6 \times 10^{-11}$ yields a consistent solution. For this benchmark, the corresponding mass of the RH neutrino is $m_N \lesssim 8.5 \times 10^6$ GeV and the value $v_L \sim 0.59$ GeV ensures a neutrino mass of 0.1 eV via the type-II

seesaw mechanism. The phenomenology of such a parameter region ($f \ll 1$) is discussed in detail in Sec. IV.

Above the scale v_R , the gauge couplings of $SU(2)_R$ and $U(1)_{B-L}$ as well as all other scalar fields can contribute to the RGE running of ρ_1 , with the one-loop beta function given by

$$\begin{aligned} \beta^{(1)}(\rho_1) = & +28\rho_1^2 + 16\rho_1\rho_2 + 16\rho_2^2 + 3\rho_3^2 + 4\alpha_1^2 \\ & + 4\alpha_1\alpha_3 + 16|\alpha_2|^2 + 2\alpha_3^2 - 12g_{BL}^2\rho_1 - 24g_R^2\rho_1 \\ & + 6g_{BL}^4 + 12g_{BL}^2g_R^2 + 9g_R^4 + 8\rho_1\text{Tr}(f^\dagger f) \\ & - 16\text{Tr}(f^\dagger f f^\dagger f). \end{aligned} \quad (\text{D1})$$

The gauge couplings (g_R, g_{BL}) originate from the $SU(2)_R$ and $U(1)_{B-L}$ gauge bosons, while the quartic couplings $\rho_{1,2,3}$ ($\alpha_{1,2,3}$) arise from the fields $\Delta_{L,R}$ (ϕ and $\Delta_{L,R}$) at the LR breaking scale. Notably, the negative contribution from the Yukawa coupling f , arising from fermion field N , can be made to cancel by the bosonic contribution, allowing ρ_1 to remain small. Consequently, the DM mass can be kept at or below MeV scale across all scales.

Next we show that keeping the scalar fields below v_R can cure the *electroweak* vacuum stability problem. For simplicity we keep $m_H < v_R$. The interaction of the two doublet scalar fields $\{h, H\}$ and the singlet field χ read as

$$\begin{aligned} V = & -m_h^2|h|^2 + m_\chi^2\chi^2 + m_H^2|H|^2 + \{m_{hH}^2 h^\dagger H + \text{H.c.}\} \\ & + \mu_\chi\chi^3 + \mu_{H\chi}|H|^2 + \mu_{h\chi}|h|^2 + \tilde{\lambda}_1|h|^4 + \tilde{\lambda}_2|H|^4 \\ & + \tilde{\lambda}_3|h|^2|H|^2 + \{\tilde{\lambda}_4|h|^2 H^\dagger h + \tilde{\lambda}_5|H|^2 h^\dagger H \\ & + \tilde{\lambda}_6 h^\dagger H h^\dagger H + \text{H.c.}\} + \tilde{\lambda}_7 h^\dagger H H^\dagger h + \tilde{\rho}_1\chi^4 \\ & + \tilde{\alpha}_1\chi^2|h|^2 + \tilde{\alpha}_2\chi^2|H|^2. \end{aligned} \quad (\text{D2})$$

The RGE of the quartic couplings can be written as

$$\mu \frac{dX}{d\mu} \equiv \frac{1}{(4\pi)^2} \beta^{(1)}(X), \quad (\text{D3})$$

where $\beta^{(1)}(X)$ are the one-loop functions of the quartic couplings which, in the momentum range $m_H \leq \mu \leq v_R$, are given by

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_1) = & 24\tilde{\lambda}_1^2 + 2\tilde{\lambda}_3^2 + 2\tilde{\lambda}_3\tilde{\lambda}_7 + \tilde{\lambda}_7^2 + 2\tilde{\alpha}_1^2 - \frac{9}{5}\tilde{\lambda}_1g_1^2 \\ & - 9\tilde{\lambda}_1g_2^2 + \frac{27}{200}g_1^4 + \frac{9}{20}g_1^2g_2^2 + \frac{9}{8}g_2^4 \\ & + 12\tilde{\lambda}_1|y_t|^2 - 6|y_t|^4, \end{aligned} \quad (\text{D4})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_2) = & 24\tilde{\lambda}_2^2 + 2\tilde{\lambda}_3^2 + 2\tilde{\lambda}_3\tilde{\lambda}_7 + \tilde{\lambda}_7^2 + 2\tilde{\alpha}_2^2 - \frac{9}{5}\tilde{\lambda}_2g_1^2 \\ & - 9\tilde{\lambda}_2g_2^2 + \frac{27}{200}g_1^4 + \frac{9}{20}g_1^2g_2^2 + \frac{9}{8}g_2^4, \end{aligned} \quad (\text{D5})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_3) = & 12\tilde{\lambda}_1\tilde{\lambda}_3 + 4\tilde{\lambda}_1\tilde{\lambda}_7 + 12\tilde{\lambda}_2\tilde{\lambda}_3 + 4\tilde{\lambda}_2\tilde{\lambda}_7 + 4\tilde{\lambda}_3^2 + 2\tilde{\lambda}_7^2 \\ & + 4\tilde{\alpha}_1\tilde{\alpha}_2 - \frac{9}{5}\tilde{\lambda}_3g_1^2 - 9\tilde{\lambda}_3g_2^2 + \frac{27}{100}g_1^4 - \frac{9}{10}g_1^2g_2^2 \\ & + \frac{9}{4}g_2^4 + 6\tilde{\lambda}_3|y_t|^2, \end{aligned} \quad (\text{D6})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_4) = & 24\tilde{\lambda}_1\tilde{\lambda}_4 + 6\tilde{\lambda}_3\tilde{\lambda}_4 + 6\tilde{\lambda}_3\tilde{\lambda}_5 + 20\tilde{\lambda}_4\tilde{\lambda}_6 + 8\tilde{\lambda}_4\tilde{\lambda}_7 \\ & + 4\tilde{\lambda}_5\tilde{\lambda}_6 + 4\tilde{\lambda}_5\tilde{\lambda}_7 - \frac{9}{5}\tilde{\lambda}_4g_1^2 - 9\tilde{\lambda}_4g_2^2 + 9\tilde{\lambda}_4|y_t|^2, \end{aligned} \quad (\text{D7})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_5) = & 24\tilde{\lambda}_2\tilde{\lambda}_5 + 6\tilde{\lambda}_3\tilde{\lambda}_4 + 6\tilde{\lambda}_3\tilde{\lambda}_5 + 4\tilde{\lambda}_4\tilde{\lambda}_6 + 4\tilde{\lambda}_4\tilde{\lambda}_7 \\ & + 20\tilde{\lambda}_5\tilde{\lambda}_6 + 8\tilde{\lambda}_5\tilde{\lambda}_7 - \frac{9}{5}\tilde{\lambda}_5g_1^2 - 9\tilde{\lambda}_5g_2^2 + 3\tilde{\lambda}_5|y_t|^2, \end{aligned} \quad (\text{D8})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_6) = & 4\tilde{\lambda}_1\tilde{\lambda}_6 + 4\tilde{\lambda}_2\tilde{\lambda}_6 + 8\tilde{\lambda}_3\tilde{\lambda}_6 + 5\tilde{\lambda}_4^2 + 2\tilde{\lambda}_4\tilde{\lambda}_5 + 5\tilde{\lambda}_5^2 \\ & + 12\tilde{\lambda}_6\tilde{\lambda}_7 - \frac{9}{5}\tilde{\lambda}_6g_1^2 - 9\tilde{\lambda}_6g_2^2 + 6\tilde{\lambda}_6|y_t|^2, \end{aligned} \quad (\text{D9})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\lambda}_7) = & 4\tilde{\lambda}_1\tilde{\lambda}_7 + 4\tilde{\lambda}_2\tilde{\lambda}_7 + 8\tilde{\lambda}_3\tilde{\lambda}_7 + 4\tilde{\lambda}_7^2 - \frac{9}{5}\tilde{\lambda}_7g_1^2 \\ & - 9\tilde{\lambda}_7g_2^2 + \frac{9}{5}g_1^2g_2^2 + 6\tilde{\lambda}_7|y_t|^2, \end{aligned} \quad (\text{D10})$$

$$\beta^{(1)}(\tilde{\rho}_1) = 72\tilde{\rho}_1^2 + 2\tilde{\alpha}_1^2 + 2\tilde{\alpha}_2^2, \quad (\text{D11})$$

$$\begin{aligned} \beta^{(1)}(\tilde{\alpha}_1) = & 12\tilde{\lambda}_1\tilde{\alpha}_1 + 4\tilde{\lambda}_3\tilde{\alpha}_2 + 2\tilde{\lambda}_7\tilde{\alpha}_2 + 24\tilde{\alpha}_1\tilde{\rho}_1 + 8\tilde{\alpha}_1^2 \\ & - \frac{9}{10}\tilde{\alpha}_1g_1^2 - \frac{9}{2}\tilde{\alpha}_1g_2^2 + 6\tilde{\alpha}_1|y_t|^2, \end{aligned} \quad (\text{D12})$$

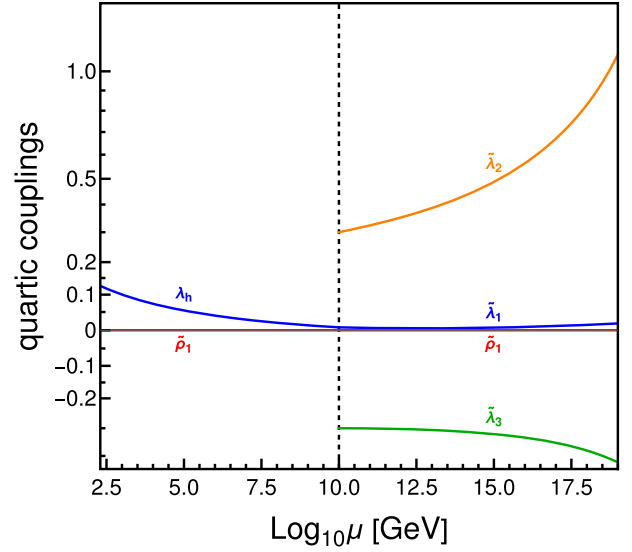


FIG. 4. Evolution of the quartic couplings to show that the electroweak vacuum is stable all the way up to the Planck scale. We have included full two-loop RGE effects in obtaining the results. The vertical dashed line corresponds to the mass of H , i.e., $m_H = 10^{10}$ GeV, and the LR scale $v_R = 10^{18}$ GeV here.

$$\begin{aligned} \beta^{(1)}(\tilde{\alpha}_2) = & 12\tilde{\lambda}_2\tilde{\alpha}_2 + 4\tilde{\lambda}_3\tilde{\alpha}_1 + 2\tilde{\lambda}_7\tilde{\alpha}_1 + 24\tilde{\alpha}_2\tilde{\rho}_1 + 8\tilde{\alpha}_2^2 \\ & - \frac{9}{10}\tilde{\alpha}_2g_1^2 - \frac{9}{2}\tilde{\alpha}_2g_2^2. \end{aligned} \quad (\text{D13})$$

Here we kept only the top quark contribution arising from the SM-like Higgs h . Below the scale m_H , only the quartic couplings $\{\tilde{\lambda}_1 \equiv \lambda_h, \tilde{\rho}_1, \tilde{\alpha}_1\}$ contribute. Numerically, we solve the RGEs up to two-loop beta functions and show that the inclusion of the scalar field H below the LR scale stabilizes the electroweak vacuum all the way up to the Planck scale, as illustrated in Fig. 4. This also preserves the necessary conditions for the boundedness of the potential. The two-loop beta functions were generated with the `PyR@TE` package [66] and were cross-checked with the known results [67,68]. The following input values at $m_t(m_t) = 162.8$ GeV were used for the gauge couplings [31],

$$\begin{aligned} \alpha_1(m_t) = 0.01024, \quad \alpha_2(m_t) = 0.0340, \\ \alpha_3(m_t) = 0.1094. \end{aligned} \quad (\text{D14})$$

Here $\alpha_i = g_i^2/4\pi$. The input values for the quartic couplings at the scale $\mu = m_H$ were chosen as

$$\begin{aligned} \tilde{\lambda}_1 = 7.8 \times 10^{-3}, \quad \tilde{\lambda}_2 = 3.0 \times 10^{-1}, \quad \tilde{\lambda}_3 = -3.0 \times 10^{-1}, \quad \tilde{\lambda}_4 = -1.8 \times 10^{-6}, \\ \tilde{\lambda}_5 = 1.8 \times 10^{-6}, \quad \tilde{\lambda}_6 = \tilde{\lambda}_7 = 0, \quad \tilde{\rho}_1 = 1 \times 10^{-46}, \quad \tilde{\alpha}_1 = 0, \quad \tilde{\alpha}_2 = 2 \times 10^{-18}. \end{aligned} \quad (\text{D15})$$

The boundedness conditions are similar to those given in Eq. (C2). Note that in Fig. 4, although $\tilde{\lambda}_3$ is negative above m_H scale, $\tilde{\lambda}_1 + 2\tilde{\lambda}_2 + \tilde{\lambda}_3$ remains positive all the way to the v_R scale, thus preserving the vacuum stability. The quartic couplings that are not shown in Fig. 4 barely evolve and play an insignificant role in the running of beta function $\beta(\tilde{\lambda}_1)$.

We note that above the scale v_R , the relation $\rho_1 = 4\tilde{\rho}_1$ holds, while the remaining quartic couplings $\rho_{2,3,4}$ are unconstrained and can take any values within the theoretical limits derived in Sec. C. For definiteness, we choose them to be of $\mathcal{O}(1)$, ensuring that the scalar fields δ_L^{0r} , δ_L^{0i} , δ_L^+ , δ_L^{++} , and δ_R^{++} acquire masses at or above the v_R scale. See Eqs. (A9)–(A16) for the corresponding analytical expressions.

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- [1] J. C. Pati and A. Salam, Lepton number as the fourth color, *Phys. Rev. D* **10**, 275 (1974); **11**, 703(E) (1975).
 - [2] R. N. Mohapatra and J. C. Pati, A natural left-right symmetry, *Phys. Rev. D* **11**, 2558 (1975).
 - [3] R. N. Mohapatra and J. C. Pati, Left-right gauge symmetry and an isoconjugate model of CP violation, *Phys. Rev. D* **11**, 566 (1975).
 - [4] G. Senjanovic and R. N. Mohapatra, Exact left-right symmetry and spontaneous violation of parity, *Phys. Rev. D* **12**, 1502 (1975).
 - [5] G. Senjanovic, Spontaneous breakdown of parity in a class of gauge theories, *Nucl. Phys. B* **153**, 334 (1979).
 - [6] P. Minkowski, $\mu \rightarrow e\gamma$ at a rate of one out of 10^9 muon decays?, *Phys. Lett.* **67B**, 421 (1977).
 - [7] R. N. Mohapatra and G. Senjanovic, Neutrino mass and spontaneous parity nonconservation, *Phys. Rev. Lett.* **44**, 912 (1980).
 - [8] T. Yanagida, Horizontal gauge symmetry and masses of neutrinos, *Conf. Proc. C* **7902131**, 95 (1979), <https://inspirehep.net/literature/143150>.
 - [9] M. Gell-Mann, P. Ramond, and R. Slansky, Complex spinors and unified theories, *Conf. Proc. C* **790927**, 315 (1979), <https://inspirehep.net/literature/9686>.
 - [10] M. Magg and C. Wetterich, Neutrino mass problem and gauge hierarchy, *Phys. Lett.* **94B**, 61 (1980).
 - [11] J. Schechter and J. W. F. Valle, Neutrino masses in $SU(2) \times U(1)$ theories, *Phys. Rev. D* **22**, 2227 (1980).
 - [12] R. N. Mohapatra and G. Senjanovic, Neutrino masses and mixings in gauge models with spontaneous parity violation, *Phys. Rev. D* **23**, 165 (1981).
 - [13] G. Lazarides, Q. Shafi, and C. Wetterich, Proton lifetime and fermion masses in an $SO(10)$ model, *Nucl. Phys. B* **181**, 287 (1981).
 - [14] J. Heeck and S. Patra, Minimal left-right symmetric dark matter, *Phys. Rev. Lett.* **115**, 121804 (2015).
 - [15] C. Garcia-Cely and J. Heeck, Phenomenology of left-right symmetric dark matter, *J. Cosmol. Astropart. Phys.* **03** (2016) 021.
 - [16] S. Ferrari, T. Hambye, J. Heeck, and M. H. G. Tytgat, $SO(10)$ paths to dark matter, *Phys. Rev. D* **99**, 055032 (2019).
 - [17] F. Bezrukov, H. Hettmansperger, and M. Lindner, keV sterile neutrino dark matter in gauge extensions of the Standard Model, *Phys. Rev. D* **81**, 085032 (2010).
 - [18] M. Nemevsek, G. Senjanovic, and Y. Zhang, Warm dark matter in low scale left-right theory, *J. Cosmol. Astropart. Phys.* **07** (2012) 006.
 - [19] M. Nemevsek and Y. Zhang, Dark matter dilution mechanism through the lens of large-scale structure, *Phys. Rev. Lett.* **130**, 121002 (2023).
 - [20] M. Nemevsek and Y. Zhang, Anatomy of diluted dark matter in the minimal left-right symmetric model, *Phys. Rev. D* **109**, 056021 (2024).
 - [21] M. Drewes *et al.*, A white paper on keV sterile neutrino dark matter, *J. Cosmol. Astropart. Phys.* **01** (2017) 025.
 - [22] A. Boyarsky, M. Drewes, T. Lasserre, S. Mertens, and O. Ruchayskiy, Sterile neutrino dark matter, *Prog. Part. Nucl. Phys.* **104**, 1 (2019).
 - [23] P. B. Pal and L. Wolfenstein, Radiative decays of massive neutrinos, *Phys. Rev. D* **25**, 766 (1982).
 - [24] C. Dessert, N. L. Rodd, and B. R. Safdi, The dark matter interpretation of the 3.5-keV line is inconsistent with blank-sky observations, *Science* **367**, 1465 (2020).
 - [25] J. W. Foster, M. Kongsore, C. Dessert, Y. Park, N. L. Rodd, K. Cranmer, and B. R. Safdi, Deep search for decaying dark matter with XMM-Newton Blank-Sky observations, *Phys. Rev. Lett.* **127**, 051101 (2021).
 - [26] R. A. Krivonos, V. V. Barinov, A. A. Mukhin, and D. S. Gorbunov, Strong limits on keV-scale galactic sterile neutrino dark matter with stray light from NuSTAR after 11 years of operation, *Phys. Rev. Lett.* **133**, 261002 (2024).
 - [27] M. Fukugita and T. Yanagida, Baryogenesis without grand unification, *Phys. Lett. B* **174**, 45 (1986).
 - [28] H. Fritzsch and P. Minkowski, Unified interactions of leptons and hadrons, *Ann. Phys. (N.Y.)* **93**, 193 (1975).
 - [29] W. Hollik, Nonstandard Higgs bosons in $SU(2) \times U(1)$ radiative corrections, *Z. Phys. C* **32**, 291 (1986).
 - [30] J. Heeck, W-boson mass in the triplet seesaw model, *Phys. Rev. D* **106**, 015004 (2022).
 - [31] S. Navas *et al.* (Particle Data Group), Review of particle physics, *Phys. Rev. D* **110**, 030001 (2024).
 - [32] T. Aaltonen *et al.* (CDF Collaboration), High-precision measurement of the W boson mass with the CDF II detector, *Science* **376**, 170 (2022).
 - [33] V. Chekhovsky *et al.* (CMS Collaboration), High-precision measurement of the W boson mass with the CMS experiment at the LHC, *arXiv:2412.13872*.
 - [34] A. Tumasyan *et al.* (CMS Collaboration), Search for a right-handed W boson and a heavy neutrino in proton-proton collisions at $\sqrt{s} = 13$ TeV, *J. High Energy Phys.* **04** (2022) 047.

- [35] G. Aad *et al.* (ATLAS Collaboration), Search for heavy Majorana or Dirac neutrinos and right-handed W gauge bosons in final states with charged leptons and jets in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Eur. Phys. J. C* **83**, 1164 (2023).
- [36] H. Baer, K.-Y. Choi, J. E. Kim, and L. Roszkowski, Dark matter production in the early Universe: beyond the thermal WIMP paradigm, *Phys. Rep.* **555**, 1 (2015).
- [37] P. S. B. Dev, R. N. Mohapatra, and Y. Zhang, Long lived light scalars as probe of low scale seesaw models, *Nucl. Phys. B* **923**, 179 (2017).
- [38] K. E. O'Donnell and T. R. Slatyer, Constraints on dark matter with future MeV gamma-ray telescopes, *Phys. Rev. D* **111**, 083037 (2025).
- [39] J. Heeck and D. Teresi, Cold keV dark matter from decays and scatterings, *Phys. Rev. D* **96**, 035018 (2017).
- [40] R. Essig, E. Kuflik, S. D. McDermott, T. Volansky, and K. M. Zurek, Constraining light dark matter with diffuse x-ray and gamma-ray observations, *J. High Energy Phys.* **11** (2013) 193.
- [41] K. C. Y. Ng, B. M. Roach, K. Perez, J. F. Beacom, S. Horiuchi, R. Krivonos, and D. R. Wik, New constraints on sterile neutrino dark matter from *NuSTAR* M31 observations, *Phys. Rev. D* **99**, 083005 (2019).
- [42] R. Laha, J. B. Muñoz, and T. R. Slatyer, INTEGRAL constraints on primordial black holes and particle dark matter, *Phys. Rev. D* **101**, 123514 (2020).
- [43] B. M. Roach, S. Rossland, K. C. Y. Ng, K. Perez, J. F. Beacom, B. W. Grefenstette, S. Horiuchi, R. Krivonos, and D. R. Wik, Long-exposure *NuSTAR* constraints on decaying dark matter in the Galactic halo, *Phys. Rev. D* **107**, 023009 (2023).
- [44] F. Calore, A. Dekker, P. D. Serpico, and T. Siebert, Constraints on light decaying dark matter candidates from 16 yr of INTEGRAL/SPI observations, *Mon. Not. R. Astron. Soc.* **520**, 4167 (2023).
- [45] J. Cooley *et al.*, Report of the topical group on particle dark matter for snowmass 2021, [arXiv:2209.07426](https://arxiv.org/abs/2209.07426).
- [46] D. Buttazzo, G. Degrandi, P. P. Giardino, G. F. Giudice, F. Sala, A. Salvio, and A. Strumia, Investigating the near-criticality of the Higgs boson, *J. High Energy Phys.* **12** (2013) 089.
- [47] G. Hiller, T. Höhne, D. F. Litim, and T. Staudner, Vacuum stability in the Standard Model and beyond, *Phys. Rev. D* **110**, 115017 (2024).
- [48] E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, Redwood City, USA, 1990), Vol. 69.
- [49] N. Aghanim *et al.* (Planck Collaboration), Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020); **652**, C4(E) (2021).
- [50] R. J. Scherrer and M. S. Turner, Decaying particles do not heat up the universe, *Phys. Rev. D* **31**, 681 (1985).
- [51] S. Boulebnane, J. Heeck, A. Nguyen, and D. Teresi, Cold light dark matter in extended seesaw models, *J. Cosmol. Astropart. Phys.* **04** (2018) 006.
- [52] S. Davidson and A. Ibarra, A lower bound on the right-handed neutrino mass from leptogenesis, *Phys. Lett. B* **535**, 25 (2002).
- [53] A. Pilaftsis and T. E. J. Underwood, Resonant leptogenesis, *Nucl. Phys. B* **692**, 303 (2004).
- [54] F. Bezrukov, A. Kartavtsev, and M. Lindner, Leptogenesis in models with keV sterile neutrino dark matter, *J. Phys. G* **40**, 095202 (2013).
- [55] K. Moffat, S. Pascoli, S. T. Petcov, H. Schulz, and J. Turner, Three-flavored nonresonant leptogenesis at intermediate scales, *Phys. Rev. D* **98**, 015036 (2018).
- [56] F. Elahi, C. Kolda, and J. Unwin, UltraViolet freeze-in, *J. High Energy Phys.* **03** (2015) 048.
- [57] B. Barman, D. Borah, and R. Roshan, Nonthermal leptogenesis and UV freeze-in of dark matter: Impact of inflationary reheating, *Phys. Rev. D* **104**, 035022 (2021).
- [58] M. Cataldi and B. Shakya, Leptogenesis via bubble collisions, *J. Cosmol. Astropart. Phys.* **11** (2024) 047.
- [59] N. G. Deshpande, J. F. Gunion, B. Kayser, and F. I. Olness, Left-right symmetric electroweak models with triplet Higgs, *Phys. Rev. D* **44**, 837 (1991).
- [60] J. Chakraborty, J. Gluza, T. Jelinski, and T. Srivastava, Theoretical constraints on masses of heavy particles in left-right symmetric models, *Phys. Lett. B* **759**, 361 (2016).
- [61] K. Hadeler, On copositive matrices, *Linear Algebra Appl.* **49**, 79 (1983).
- [62] K. G. Klimenko, On necessary and sufficient conditions for some Higgs potentials to be bounded from below, *Theor. Math. Phys.* **62**, 58 (1985).
- [63] J. Chakraborty, P. Konar, and T. Mondal, Copositive criteria and boundedness of the scalar potential, *Phys. Rev. D* **89**, 095008 (2014).
- [64] P. S. B. Dev, R. N. Mohapatra, W. Rodejohann, and X.-J. Xu, Vacuum structure of the left-right symmetric model, *J. High Energy Phys.* **02** (2019) 154.
- [65] G. Chauhan, Vacuum stability and symmetry breaking in left-right symmetric model, *J. High Energy Phys.* **12** (2019) 137.
- [66] L. Sartore and I. Schienbein, PyR@TE 3, *Comput. Phys. Commun.* **261**, 107819 (2021).
- [67] I. Z. Rothstein, Renormalization group analysis of the minimal left-right symmetric model, *Nucl. Phys. B* **358**, 181 (1991).
- [68] A. Maiezza, G. Senjanović, and J. C. Vazquez, Higgs sector of the minimal left-right symmetric theory, *Phys. Rev. D* **95**, 095004 (2017).