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Low scale left-right symmetry and naturally small neutrino mass

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Abstract: We consider the low scale (10–100 TeV) left-right symmetric model with "naturally" small neutrino masses generated through the inverse seesaw mechanism. The Dirac neutrino mass terms are taken to be similar to the masses of charged leptons and quarks in order to satisfy the quark-lepton similarity condition. The inverse seesaw implies the existence of fermion singlets S with Majorana mass terms as well as the "left" and "right" Higgs doublets. These doublets provide the portal for S and break the left-right symmetry. The inverse seesaw allows to realize a scenario in which the large lepton mixing originates from the Majorana mass matrix of S fields which has certain symmetry. The model contains heavy pseudo-Dirac fermions, formed by S and the right-handed neutrinos, which have masses in the 1 GeV–100 TeV range and can be searched for at current and various future colliders such as LHC, FCC-ee and FCC-hh as well as in SHiP and DUNE experiments. Their contribution to neutrinoless double beta decay is unobservable. The radiative corrections to the mass of the Higgs boson and the possibility for generating the baryon asymmetry of the Universe are discussed. Modification of the model with two singlets (S_L and S_R) per generation can provide a viable keV-scale dark matter candidate.

Keywords: Beyond Standard Model, Neutrino Physics

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Contents

1 Introduction

The left-right symmetric models [1–5] based on the gauge symmetry group

$$\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}\,, \tag{1.1}$$

and parity P which ensures the equality between couplings in the left and right sectors, are still one of the most appealing and well-motivated extensions of the Standard Model (SM) [6–15]. Spontaneous breaking of $\mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L} \times P$ down to the SM symmetry group explains the observed low-energy asymmetry between the left and right as well as provides a natural framework for the generation of small neutrino masses via the seesaw mechanism [16–19]. The scale of left-right (L-R) symmetry breaking and the seesaw scale coincide. The key question is whether this scale can be at $\mathcal{O}(10-100)$ TeV energies that are accessible to LHC and the next generation of colliders.

The majority of the low scale L-R symmetric models constructed so far is at odds with the generation of "naturally" small neutrino masses. In what follows we will call the neutrino mass to be naturally small if mechanism of its generation employs the Dirac neutrino masses, m_ν^D, similar in size to the Dirac masses of charged leptons, m_l, and quarks, m_q, i.e.

$$m_\nu^D \approx m_q, m_l\,. \tag{1.2}$$

This relation facilitates the grand unification and we will refer to it as to the quark-lepton (q-l) similarity condition. The usual type-I seesaw mechanism realizes such a possibility provided that the scale of right-handed (RH) neutrino masses is about 10^{14} GeV. Lowering the scale of RH neutrinos down to e.g. 10^4 GeV requires $m_\nu^D \sim 1$ MeV, which is 5 orders of magnitude smaller than the top quark mass and thus is not in accord with eq. (1.2) for the third generation of neutrinos. If neutrinos acquire their masses dominantly via the type-II seesaw mechanism, the Dirac mass terms should be even more strongly suppressed.

In this paper, in order to reconcile the low scale L-R symmetry and the naturally small neutrino masses we assume that the latter are generated via the inverse seesaw mechanism [20, 21]. In such a framework, the small neutrino masses can be obtained for values of the Dirac mass terms that are in accord with eq. (1.2). The inverse seesaw mechanism requires the introduction of new fermionic singlets, S, which couple to the RH neutrinos and thus form the Dirac mass terms. We introduce three such singlets (one per generation), whose Majorana masses are much smaller than the electroweak scale. The generation of light neutrino masses via the inverse seesaw requires the right-handed $\mathrm{SU}(2)_R$ Higgs doublet and therefore, due to L-R symmetry, the left-handed $\mathrm{SU}(2)_L$ doublet. These doublets break the L-R symmetry, so that the Higgs triplets are not needed [22–24]. Hence, the problem of the absence of low-dimensional representations[1] in the model does not appear.

We do not assume any special smallness of the Yukawa couplings of the scalar doublets. These couplings are similar or even equal to the couplings of the bi-doublet. If equal, the so called screening of the Dirac structures is realized [25], and the large lepton mixing originates from a certain structure of the Majorana mass matrix of singlets S (the screening was previously studied for the double seesaw mechanism [26]). The mass matrix of S may have certain symmetry which leads, e.g., to the tribimaximal mixing.

In this paper we elaborate on such a scenario. We focus on generation of neutrino mass and mixing in the L-R model with Higgs doublets instead of triplets. Such models have been extensively explored before [4, 27] and the only new element here are singlet fermions S which allow to realize the inverse seesaw mechanism (as a dominant mechanism for the generation of neutrino mass) and certain selection of the Yukawa couplings. We explore here new features related to introduction of the fermion singlets, while the gauge and scalar boson sectors are the same as in several earlier publications. We confront the model with the existing experimental data from the beam-dump experiments, LHC, experiments searching for neutrinoless double beta decay, *etc.*, and also estimate the discovery potential of future colliders and neutrino oscillation facilities. We obtain bounds on relevant parameters and constrain the L-R symmetry breaking scale. We examine the possibility for the generation of the observed baryon asymmetry of the Universe and address the issue of the Higgs naturalness. We also consider the extensions of the aforementioned scenario, in particular the scenario with two S fields (left and right) per generation which contains a viable dark matter candidate.

[1] In the vast majority of studies, L-R symmetry is broken by introducing the Higgs triplets, while the Higgs doublets are absent. The existence of the higher representations (triplets) and the absence of low-dimensional representations (doublets here) should have a certain reason and a proper physical explanation.

The paper is organized as follows. In section 2, we describe the model and generation of the neutrino masses, and discuss the possibility to introduce flavor symmetries. The phenomenology of the model is presented in section 3. We elaborate on the various extensions of this scenario in section 4. Finally, in section 5 we conclude.

2 The model and neutrino masses

2.1 The model, linear and inverse seesaw

Leptons are organized in the following representations of the symmetry group (1.1)

$$L_L = \begin{pmatrix} \nu_l \\ l \end{pmatrix}_L \sim (2,1,-1)\,, \qquad L_R = \begin{pmatrix} \nu_l \\ l \end{pmatrix}_R \sim (1,2,-1)\,, \qquad S \sim (1,1,0)\,, \tag{2.1}$$

where $l = \{e, \mu, \tau\}$ and in brackets we indicate the corresponding quantum numbers. The Majorana leptons S are complete gauge singlets. We assume the existence of three such leptons — one per generation.

The scalar sector consists of the usual bi-doublet

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} \sim (2,2,0)\,, \tag{2.2}$$

and two doublets

$$\chi_L = \begin{pmatrix} \chi_L^+ \\ \chi_L^0 \end{pmatrix} \sim (2,1,1)\,, \qquad \chi_R = \begin{pmatrix} \chi_R^+ \\ \chi_R^0 \end{pmatrix} \sim (1,2,1)\,. \tag{2.3}$$

The latter are required for the realization of the inverse seesaw mechanism as well as for the L-R symmetry breaking. In the first step, the neutral component of the right-handed scalar doublet χ_R obtains a non-vanishing vacuum expectation value (VEV), which breaks the $\mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ down to $\mathrm{U}(1)_Y$, where Y is the hypercharge. In the second step, the neutral scalar fields from left-handed scalar doublet χ_L (χ_L^0) and bi-doublet Φ (ϕ_1^0, ϕ_2^0) break the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ symmetry down to $\mathrm{U}(1)_{\mathrm{EM}}$. Their VEVs should satisfy the relation

$$\sqrt{\langle\chi_L^0\rangle^2 + \langle\phi_1^0\rangle^2 + \langle\phi_2^0\rangle^2} \approx 246\,\mathrm{GeV}, \tag{2.4}$$

in order to reproduce the electroweak scale. Notice that for the electroweak symmetry breaking only one non-zero VEV among these three neutral fields is enough. The L-R symmetry breaking requires

$$\langle\chi_R^0\rangle \gg \langle\chi_L^0\rangle, \langle\phi_{1,2}^0\rangle\,. \tag{2.5}$$

The Higgs sector in eqs. (2.2) and (2.3) is identical to the one in earlier publications [4, 27]. Minimization of the potential had been done under certain simplifications. In particular, it was assumed *a priori* that the electic charge is conserved in minimum. The tri-linear terms in potential are absent due to additional symmetry. With such conditions, it was shown [4, 27] that the global minimum exists for a certain range of parameters of the potential, in which the inequality (2.5) is satisfied, i.e. the parity is spontaneosly

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broken. The requirements for such a scenario are inequalities of certain quartic couplings and positivity or small values of other couplings [4, 27]. The values of VEVs are controlled (at least in the case of mild hierarchy) by the mass parameters for the doublets and bi-doublet. For the most general form of the potential (and without above discussed assumptions) the minimization has not been done for the bi-doublet-doublet case. Such a study has been done recently [28] for the bi-doublet-triplet scenario and some information can be inferred from the obtained results. It has been found in [28] that there are significant regions in the parameter space where required minimum can be obtained. Note that bi-doublet-doublet and bi-doublet-triplet models have the same number of neutral bosons, and structure of the terms in the potential is similar. Interestingly, the $\beta-$ terms in the bi-doublet-triplet potential should be small in the minimum and the terms of similar type in the bidoublet-doublet potential are absent.

For the rest of our paper it is enough that hierarchy (2.5) is achieved at least for some choice of parameters. Since we are not discussing phenomenology of the Higgs sector (masses of scalars, decay rates, etc.) specific values of the parameters of the higgs potential are not important.

The lepton masses are generated by the following Lagrangian [22]

$$\mathcal{L} \supset -\bar{L}_R\, Y\, \Phi^\dagger L_L - \bar{L}_R\, \tilde{Y}\, \tilde{\Phi}^\dagger L_L - \bar{S}\, Y_L \tilde{\chi}_L^\dagger L_L - \bar{S}^c\, Y_R\, \tilde{\chi}_R^\dagger L_R - \frac{1}{2}\,\bar{S}^c\, \mu\, S + \text{h.c.}\,, \tag{2.6}$$

where Y, $\tilde{Y}$, Y_L, Y_R are 3×3 matrices of the Yukawa couplings and μ is the 3×3 Majorana mass matrix of S leptons. $\tilde{X} \equiv i\sigma_2 X^*$ $(X = \{\chi_L, \chi_R\})$, $S^c \equiv C\bar{S}^T$ and $\tilde{\Phi} \equiv \sigma_2 \Phi^* \sigma_2$ denote charge conjugated fields of scalars and fermions. The field transformations under parity

$$L_L \Longleftrightarrow L_R\,, \quad \chi_L \Longleftrightarrow \chi_R\,, \quad \Phi \Longleftrightarrow \Phi^\dagger\,, \quad S \Longleftrightarrow S^c\,, \tag{2.7}$$

impose the following relations

$$Y_L = Y_R\,, \qquad Y = Y^\dagger\,, \qquad \tilde{Y} = \tilde{Y}^\dagger\,, \qquad \mu = \mu^\dagger\,, \tag{2.8}$$

above the L-R symmetry breaking scale.

When the scalar fields acquire VEVs, the interactions (2.6) generate the mass matrix of neutral leptons

$$\mathcal{M} = \begin{pmatrix} 0 & m_D^T & m_D'^T \\ m_D & 0 & M_D^T \\ m_D' & M_D & \mu \end{pmatrix}, \tag{2.9}$$

given in the (ν_L, N_L, S^c) basis $(N_L \equiv \nu_R^c)$. Here,

$$m_D = \frac{1}{\sqrt{2}}\left(Y\langle\phi_1^0\rangle + \tilde{Y}\langle\phi_2^0\rangle\right), \qquad m_D' = \frac{1}{\sqrt{2}} Y_L\langle\chi_L^0\rangle\,, \qquad M_D = \frac{1}{\sqrt{2}} Y_R\langle\chi_R^0\rangle\,. \tag{2.10}$$

For simplicity we assume $\langle\phi_1^0\rangle \gg \langle\phi_2^0\rangle$ such that only the first term in m_D contributes.

The block diagonalization of $\mathcal{M}$ leads to the light neutrino mass matrix

$$m_\nu \simeq \frac{\langle\chi_L^0\rangle}{\langle\chi_R^0\rangle}\left(m_D + m_D^T\right) - m_D^T\, M_D^{-1}\, \mu\, \left(M_D^T\right)^{-1} m_D\,. \tag{2.11}$$

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Here, the first term is the linear seesaw contribution [20], whose existence is a generic consequence of the inverse seesaw realization in the L-R models. Notice that due to $Y_L = Y_R$, the Yukawa matrices cancel and this term is given by the Dirac mass matrix multiplied by the ratio of VEVs of the two doublets. For $\langle\chi_L^0\rangle$ of the order of electroweak scale, and in the absence of unnaturally small elements of m_D, this term yields too large neutrino masses. Furthermore, if $m_D \propto m_u$ (the subscript u denotes the up-type quarks), it has a wrong flavor structure with too strong mass hierarchy and small mixing. This is incompatible with the neutrino mass squared differences and large mixing angles observed in the oscillation experiments. Therefore, the linear seesaw contribution should be at most sub-dominant and the main contribution to the neutrino mass should arise from the inverse seesaw, given in the second term of eq. (2.11). Since $m_D \sim m_u$, this is achieved for

$$\frac{\langle\chi_L^0\rangle}{\langle\chi_R^0\rangle} < \frac{0.05\,\text{eV}}{2\,m_D^{\max}} \sim 10^{-12}, \tag{2.12}$$

where $m_D^{\max}$ denotes the largest entry of the Dirac neutrino mass matrix. It is worth noting that for $\langle\chi_L^0\rangle$, $\langle\phi_2^0\rangle \ll \langle\phi_1^0\rangle \approx 246\,\text{GeV}$, the SM Higgs boson is associated to the real part of ϕ_1^0 field.

In order to estimate $\langle\chi_L^0\rangle$ we consider the following terms of the scalar potential

$$V \supset h\chi_L^\dagger\tilde{\Phi}\chi_R - m_\chi^2\chi_L^\dagger\chi_L\,, \tag{2.13}$$

where h is the dimensionful coupling and the term $\lambda(\chi_L^\dagger\chi_L)^2$ can be neglected for small values of $\langle\chi_L^0\rangle$. The minimization condition, $\partial V/\partial\chi_L^\dagger = 0$, gives

$$\langle\chi_L^0\rangle = h\,\frac{\langle\phi_1^0\rangle}{\langle\chi_R^0\rangle}\,, \tag{2.14}$$

where we have taken into account that due to the L-R symmetry $m_{\chi_L}^2 = m_{\chi_R}^2 \sim \langle\chi_R^0\rangle^2$. Using the condition (2.12) we obtain

$$h \lesssim 40\,\text{keV}\left(\frac{\langle\chi_R^0\rangle}{10^5\ \text{GeV}}\right)^2, \tag{2.15}$$

which needs to be satisfied in order to generate neutrino masses mainly via the inverse seesaw mechanism. Notice that eq. (2.14) is a realization of the VEV seesaw [29]. In contrast to the triplet case, the electroweak scale VEV (of the bi-doublet) enters this relation linearly and it contains dimensionful coupling h. According to eq. (2.14), the VEV of χ_L^0 is controlled by free parameter h.

The coupling h in the potential (15) can be forbidden by symmetry with respect to transformation $\Phi \to e^{i\pi/2}\Phi$ [27]. This symmetry is explicitly broken by the Yukawa interactions of Φ, and therefore it does not prevent the appearance of the h term in higher orders of perturbation theory. Even if the coupling h vanishes at tree-level due to some symmetry, $\langle\chi_L^0\rangle \neq 0$ is generated radiatively via the one-loop diagram shown in figure 1. It can be estimated as

$$\langle\chi_L^0\rangle \simeq \frac{1}{16\pi^2}\langle\chi_R^0\rangle\frac{\langle\phi_1^0\rangle}{\langle\chi_R^0\rangle}\frac{\mu}{\langle\chi_R^0\rangle}. \tag{2.16}$$

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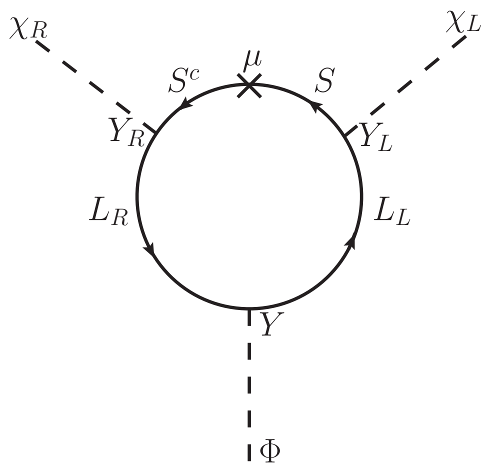

Figure 1. 1-loop diagram generating $\chi_L^\dagger \Phi \chi_R$ term in the potential.

For $\mu \simeq \mathcal{O}(10\!-\!100)\,\text{keV}$ (see below) and $\langle\chi_R^0\rangle = 10^5\,\text{GeV}$ eq. (2.16) gives $\langle\chi_L^0\rangle \simeq 10^{-14}\langle\chi_R^0\rangle$ which satisfies eq. (2.12).

Notice that for $\langle\chi_R^0\rangle \sim 10^5\,\text{GeV}$ (see section 3.1), h should be at most 100 keV, which (as we will see) is of the order of μ. This smallness can be associated to violation of the lepton number. In particular, one can reintroduce the global lepton number L_g and assign the charges $(1,-1,1)$ to (ν_L, N_L, S). In the limit $m'_D, \mu \to 0$, the conservation of $\tilde{L}$ is recovered and hence the small value of μ appears to be technically natural à la 't Hooft [30].

The inverse seesaw contribution in eq. (2.11) can be rewritten as

$$m_\nu \approx \frac{\langle\phi_1^0\rangle^2}{\langle\chi_R^0\rangle^2}\, Y^T\, Y_R^{-1}\, \mu\, \left(Y_R^T\right)^{-1} Y\,. \tag{2.17}$$

For $Y_3, Y_{R3} = 1$ we obtain $m_\nu \approx (\langle\phi_1^0\rangle/\langle\chi_R^0\rangle)^2\,\mu$ (hereafter we denote Yukawa matrices and the corresponding Dirac mass terms with only one subscript since these matrices are taken to be diagonal (see section 2.3)) which allows us to estimate $\langle\chi_R^0\rangle$ as

$$\langle\chi_R^0\rangle \simeq 3.5\cdot 10^5\ \text{GeV}\left(\frac{\mu}{100\text{keV}}\right)^{1/2}\left(\frac{0.05\,\text{eV}}{m_{\nu 3}}\right)^{1/2}, \tag{2.18}$$

where μ is the representative element of the corresponding matrix and $m_{\nu 3}$ is the largest active neutrino mass. Consequently,

$$\mu_{ij} \ll m_{Di}\,, M_{Di}\,. \tag{2.19}$$

The inverse seesaw not only explains the smallness of neutrino masses under the condition of q-l similarity, but also provides a rather appealing framework for the lepton mixing generation where the large mixing angles originate from the μ matrix of singlets S, i.e. from the hidden sector.

The leptonic mixing matrix (PMNS) can be written as

$$U_{\text{PMNS}} = U_l^\dagger U_\nu\,, \tag{2.20}$$

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where U_l follows from the diagonalization of the charged lepton mass matrix, whereas U_ν diagonalizes the neutrino mass matrix (2.17) generated by the inverse seesaw mechanism. A good agreement with experimental data can be achieved if

$$U_l \approx V_{\rm CKM}\,, \qquad U_\nu \sim U_{\rm TBM} \text{ or } U_{\rm BM}\,, \tag{2.21}$$

where $V_{\rm CKM}$ is the mixing matrix in the quark sector, and TBM and BM denote tribimaximal [31] and bimaximal [32, 33] mixing matrices. The first relation in eq. (2.21) could be a consequence of the grand unification or identical horizontal symmetry in the quark and lepton sectors. The second relation can follow from certain symmetry in the singlet sector.

2.2 Screening and *q*-*l* similarity

Let us consider the generation of neutrino mixing in the basis where Y, and therefore the Dirac mass matrix m_D, are diagonal

$$Y = Y^{\rm diag} \equiv {\rm diag}(Y_1,\ Y_2,\ Y_3)\,. \tag{2.22}$$

First, we assume that

$$Y_R = Y \tag{2.23}$$

(and also $Y_L = Y$), so that $m_D \propto M_D$. This equality can be a consequence (a remnant) of further unification when ν_L, N_L and S^c enter the same multiplet, e.g. 27-plet of the E_6 grand unification theory. This relation can also stem from certain horizontal symmetry [34]. The equality in eq. (2.23) leads to the screening of the Dirac structures: Y_R and Y cancel in the expression for the light neutrino mass matrix (see eq. (2.17)), so that

$$m_\nu \approx \xi^2 \mu\,, \tag{2.24}$$

where

$$\xi \equiv \frac{\langle \phi_1^0 \rangle}{\langle \chi_R^0 \rangle} = \frac{m_{Di}}{M_{Di}}\,. \tag{2.25}$$

According to eq. (2.24), the structure of the light neutrino mass matrix is given by the structure of the Majorana matrix μ. In particular, the neutrino contribution to the PMNS matrix is determined by μ:

$$U_\nu^T\, m_\nu\, U_\nu = \xi^2\, U_\nu^T\, \mu\, U_\nu = {\rm diag}\,(m_{\nu 1}, m_{\nu 2}, m_{\nu 3})\,. \tag{2.26}$$

Previously, such a cancellation of the couplings has been considered for the double seesaw model in refs. [25, 26, 34, 35].

Second, we assume the q-l similarity

$$Y \approx Y_u\,, \tag{2.27}$$

where Y_u is the up-type quark Yukawa coupling matrix. The screening and the q-l similarity conditions determine the phenomenology of this scenario.

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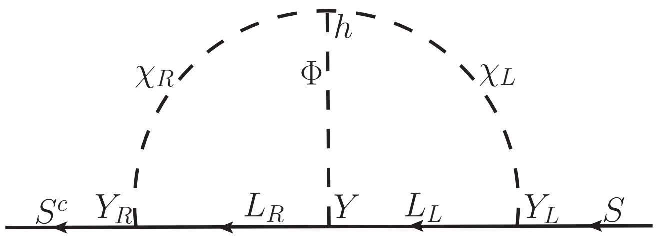

Figure 2. The leading radiative correction to the Majorana mass μ.

2.3 Flavor symmetries

The matrices Y and Y_R can be diagonal simultaneously due to the $G_{\rm basis} = Z_2 \times Z_2$ symmetry with $(-,-)$, $(+,-)$, $(-,+)$ charges for the three generations of fermions and uncharged scalar sector. We will call $G_{\rm basis}$ the basis fixing symmetry [34]. This symmetry is broken by the non-diagonal matrix μ and, in fact, the smallness of μ with respect to the other scales in the model can be related to this breaking. The μ term can arise from the interactions of S with the new gauge singlet bosons which carry non-trivial $Z_2 \times Z_2$ charge and develop non-zero VEVs. This Abelian symmetry, however, does not ensure the equality of the diagonal elements of Y and Y_R. The equality can be achieved by introducing, for instance, a permutation symmetry or by further unification mentioned above.

In the case of the exact screening, the matrix μ should have nearly tribimaximal form in the basis fixed by $G_{\rm basis} = Z_2 \times Z_2$. This can be achieved by introducing the non-Abelian (e.g., discrete) symmetry G_f which is broken down to $G_{\rm basis}$ in the visible sector and to another residual symmetry $G_{\rm hidden} = Z_2 \times Z_2$ in the $S-$sector [34]. In the visible sector, G_f can be broken explicitly. Similar construction has been realized for the double seesaw model [26] with masses of singlets at the Planck scale. In [26] the explicit symmetry breaking occurs at the lower (grand unification) scale and its impact on the neutrino mixing is suppressed by $\mathcal{O}(M_{\rm GUT}/M_{\rm Pl})$ factor. In the case of inverse seesaw, the μ scale (with certain symmetry) is much lower than the explicit symmetry breaking scale M_D. So, *a priori*, the corrections to μ can be large, thus destroying the structure of μ imposed by symmetry. To check this, let us assume that the required structure of μ, and consequently m_ν, is achieved at the tree-level and estimate the corresponding radiative corrections. The lowest order correction to μ is given by the two loop diagram shown in figure 2. It can be estimated as

$$\Delta\mu_{jj} \simeq \frac{1}{(16\pi^2)^2} Y_{Lj}^* \, Y_{Rj} \, Y_j \, h \,. \tag{2.28}$$

The corrections are diagonal due to screening. The hierarchical values of Yukawa couplings Y_{Lj}, Y_{Rj}, Y_j violate symmetry that sets the pattern of μ. The largest correction is the one to μ_{33} for which $Y_{L3} = Y_{R3} = Y_3 = 1$. Taking $h \sim 0.1\,$MeV (the value inferred from eq. (2.15) for $\langle \chi_R^0 \rangle \simeq 10^5\,$GeV) we obtain $\Delta\mu_{33} \sim 10\,$eV. The tree-level entries of the μ matrix are $\sim 0.1\,$MeV in order to reproduce $\sim 0.1\,$eV neutrino masses. Thus, the radiative corrections are much smaller than the tree-level contribution and the structure imposed by symmetries is preserved with high accuracy.

3 Phenomenology and naturalness

3.1 Heavy neutral lepton searches

Neglecting the linear seesaw contribution, we obtain from eq. (2.9) the mass matrix in (ν_L, N_L, S^c) basis

$$\mathcal{M} = \begin{pmatrix} 0 & m_D & 0 \\ m_D & 0 & M_D \\ 0 & M_D & \mu \end{pmatrix}, \tag{3.1}$$

where both m_D and M_D are simultaneously diagonal under the screening assumption, and furthermore $m_{Di} = \xi M_{Di}$. The diagonalization of $\mathcal{M}$ can be performed in several steps.

(i). We start with rotation in the ν-S plane

$$\mathcal{U}_S = \begin{pmatrix} c_\xi \mathbb{1} & 0 & s_\xi \mathbb{1} \\ 0 & \mathbb{1} & 0 \\ -s_\xi \mathbb{1} & 0 & c_\xi \mathbb{1} \end{pmatrix}, \tag{3.2}$$

where

$$s_\xi = \frac{\xi}{\sqrt{1+\xi^2}} \approx \xi\,, \tag{3.3}$$

and for brevity hereafter we denote $c \equiv \cos$, $s \equiv \sin$. After this rotation, in the new basis $(\nu',\ N_L,\ S')$, the 1-2 and 2-1 blocks vanish and the mass matrix becomes

$$\mathcal{M}_S = \begin{pmatrix} s_\xi^2\,\mu & 0 & -s_\xi\, c_\xi\, \mu \\ 0 & 0 & M_\xi \\ -s_\xi\, c_\xi\, \mu & M_\xi & c_\xi^2\, \mu \end{pmatrix}, \tag{3.4}$$

with

$$M_{\xi i} \equiv M_{Di}\sqrt{1+\xi^2}\,. \tag{3.5}$$

(ii). Next, we perform rotations in the $N_L - S'$ plane

$$\mathcal{U}_N = \begin{pmatrix} \mathbb{1} & 0 & 0 \\ 0 & \mathbf{c}_N \mathbb{1} & \mathbf{s}_N \mathbb{1} \\ 0 & -\mathbf{s}_N \mathbb{1} & \mathbf{c}_N \mathbb{1} \end{pmatrix}, \tag{3.6}$$

by the angles close to 45° which approximately diagonalize the N_L-S' block. Here $\mathbf{s}_N \equiv \mathrm{diag}\,(s_N^1, s_N^2, s_N^3)$ and

$$s_N^i \approx \frac{1}{\sqrt{2}}\left[1 - \frac{\mu_{ii}}{4M_{Di}}\right], \tag{3.7}$$

where μ_{ii} $(i = 1, 2, 3)$ are the diagonal elements of the matrix μ and we take $c_\xi \simeq 1$. As a result, the mass matrix in the new basis $(\nu',\ N^-,\ N^+)$ reads

$$\mathcal{M}_N \approx \begin{pmatrix} s_\xi^2 \mu & \frac{1}{\sqrt{2}}\, s_\xi\, \mu & -\frac{1}{\sqrt{2}}\, s_\xi\, \mu \\ \frac{1}{\sqrt{2}}\, s_\xi\, \mu & M^- & 0 \\ -\frac{1}{\sqrt{2}}\, s_\xi\, \mu & 0 & M^+ \end{pmatrix}. \tag{3.8}$$

Here,

$$M_i^- = -M_{\xi_i} + \frac{1}{2}\mu_{ii}\,, \qquad M_i^+ = M_{\xi_i} + \frac{1}{2}\mu_{ii}\,, \tag{3.9}$$

are the masses of N_i^- and N_i^+. Thus, the fields N_i^- and N_i^+ form a pair of quasi-degenerate heavy neutral leptons with $\sim \mu_{ii}$ mass splitting. For the phenomenology of heavy states, the 1-2 and 1-3 blocks in eq. (3.8) can be neglected.

(iii). Finally, we diagonalize the 1-1 block of the matrix (3.8), i.e. the light neutrino mass matrix, via

$$\mathcal{U}_\nu = \begin{pmatrix} U_\nu & 0 & 0 \\ 0 & \mathbb{1} & 0 \\ 0 & 0 & \mathbb{1} \end{pmatrix}. \tag{3.10}$$

The transition to the flavor basis requires an additional rotation which diagonalizes the mass matrix of the charged leptons

$$\mathcal{U}_l = \begin{pmatrix} U_l & 0 & 0 \\ 0 & \mathbb{1} & 0 \\ 0 & 0 & \mathbb{1} \end{pmatrix}. \tag{3.11}$$

Then, the total mixing matrix in the flavor basis is given by the product of rotations

$$\mathcal{U}_f = \mathcal{U}_l^\dagger\, \mathcal{U}_S\, \mathcal{U}_N\, \mathcal{U}_\nu = \begin{pmatrix} U_{\text{PMNS}} & -\mathbf{s}_N s_\xi\, U_l^\dagger & \mathbf{c}_N s_\xi\, U_l^\dagger \\ 0 & \frac{1}{\sqrt{2}}\mathbb{1} & \frac{1}{\sqrt{2}}\mathbb{1} \\ -s_\xi\, U_\nu & -\mathbf{s}_N\, \mathbb{1} & \mathbf{c}_N \mathbb{1} \end{pmatrix}. \tag{3.12}$$

Neglecting μ contribution in eq. (3.7) and taking $c_N = s_N = 1/\sqrt{2}$ yields

$$\nu_\alpha = U_{\text{PMNS}}\; \nu - \frac{1}{\sqrt{2}} s_\xi\, U_l^\dagger (N^- - N^+)\,, \tag{3.13}$$

where ν_α and ν are the flavor and light mass eigenstates, respectively. Since U_l is non-diagonal, each active neutrino state has admixtures of all pairs of heavy leptons. These admixtures can be constrained by various terrestrial experiments as well as by cosmology [36]. In figure 3 we show the bounds (adopted from ref. [37]) on the admixtures of $N_i^\pm$ $(i = 1, 2, 3)$ in $(\nu_e,\, \nu_\mu,\, \nu_\tau)$. Although these bounds have been derived for the mixing of a single heavy lepton, they are also applicable in our scenario where several heavy states are simultaneously present in the model.

According to eq. (3.13), the admixture of N_i^- and N_i^+ in ν_α equals

$$\left|U_{\alpha i}^N\right|^2 \equiv \left|U_{\alpha i}^{N^-}\right|^2 = \left|U_{\alpha i}^{N^+}\right|^2 = \frac{1}{2} s_\xi^2 \left|U_{l\,\alpha i}\right|^2 = \frac{1}{2}\left(\frac{m_{Di}}{M_i}\right)^2 \left|U_{l\,\alpha i}\right|^2, \tag{3.14}$$

where in the last equality we expressed s_ξ in terms of $N^\pm$ mass, and we define $M_i = (M_i^+ - M_i^-)/2$. For relevant cases of our model, the production coherence of the mass eigenstates N_i^+ and N_i^- is strongly broken, especially for the lightest leptons $(i = 1)$ that appear in low energy processes. That is, N_i^+ and N_i^- are produced (as components of N) and then decay incoherently without interference effects. Consequently, equal number of l^+ and l^- leptons will appear in the decays. For heavier leptons produced in very high energy processes (e.g. decays of W_R) the coherence can be maintained (see [14] and references therein). The experimental bounds, given in figure 3, are obtained for a production of a single heavy lepton. Since we deal here with two nearly degenerate states that are indistinguishable in experiments, the corresponding bounds on the individual mixing are two times stronger. In other words, we can treat the pair as a single particle and multiply the mixing by a factor of two so that the black lines in figure 3 correspond to

$$2|U^N_{\alpha i}|^2 = \frac{m^2_{Di}}{M^2_i}\,|U_{l\,\alpha i}|^2\,. \tag{3.15}$$

We use $m^2_{Di} = (m^2_u,\ m^2_c,\ m^2_t)$, $U_{l\,\alpha i} = (V_{\rm CKM})_{\alpha i}$ and do not impose here any relations between the heavy lepton masses M_i, in other words we are treating them as independent (the relaxation of the screening condition is discussed in section 4).

Notice that since $U_l \sim V_{\rm CKM} \sim \mathbb{1}$, the strongest bounds appear in the cases when the diagonal elements of U_l are involved. The most stringent bound on the mass of $N^\pm_1$ comes from its admixture in ν_e

$$\left|U^N_{e1}\right|^2 \approx \frac{1}{2}\left(\frac{m_{D1}}{M_1}\right)^2 \approx \frac{m^2_u}{2M^2_1}\,, \tag{3.16}$$

where $m_u \approx (1-2)\,$MeV is the mass of the up quark at the TeV scale. In the left panel of figure 3 we show with a black line (dashed in the excluded parameter space, solid elsewhere) the dependence of $2\left|U^N_{e1}\right|^2$ on M_1 for $m_u \simeq 2\,$MeV. From this figure we infer

$$M_1 \geq 2\,\text{GeV}\,, \tag{3.17}$$

which is set by the CHARM [38, 39] exclusion region. Varying m_u by a factor of 3 does not change the limit in eq. (3.17). However, for $m_u < 0.5\,$MeV the limit becomes weaker: $M_1 \geq 0.4\,$GeV.

Despite the involved CKM suppression, the admixture of $N^\pm_1$ in ν_μ yields practically identical (3.17) bound on M_1, which mainly stems from NuTeV [40]. From eq. (3.14) we have

$$\left|U^N_{\mu 1}\right|^2 \approx \frac{1}{2}\left(\frac{m_{D1}}{M^2_1}\right)^2 \left|U_{l\,\mu 1}\right|^2 \approx \frac{m^2_u}{2M^2_1}\sin^2\theta_{\rm c}\,, \tag{3.18}$$

where $\theta_{\rm c}$ is the Cabibbo angle. In turn, this gives the bound

$$\xi \leq 10^{-3}\,. \tag{3.19}$$

In our framework, the neutral lepton masses are related by screening and the q-l similarity. Employing the limit (3.17) we obtain $M_2 = M_1 m_c/m_u \geq 600\,$GeV, where $m_c \approx 0.5\,$GeV is the mass of the charm quark at the TeV scale [41]. From the screening relation for the mass of the third generation of heavy leptons, $M_3 = M_1 m_t/m_u$, we find the limit on the L-R symmetry breaking scale

$$\langle\chi^0_R\rangle \approx \sqrt{2}\,M_3 \geq 2\times 10^5\,\text{GeV}. \tag{3.20}$$

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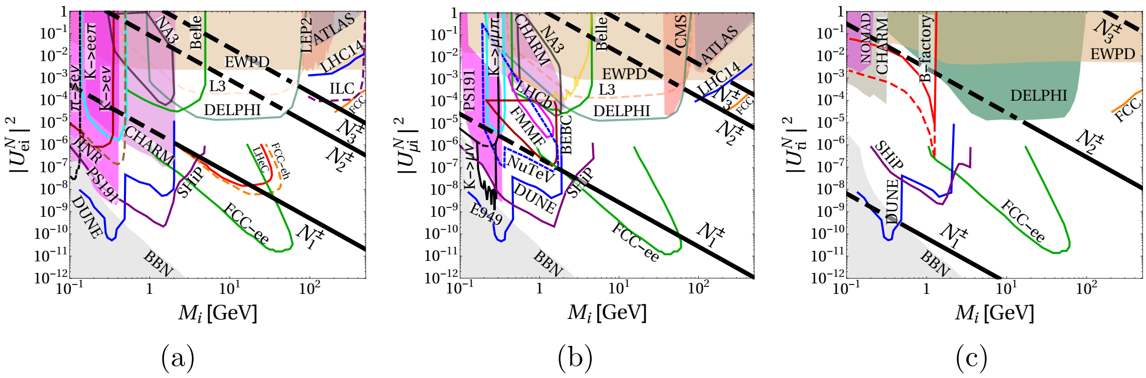

Figure 3. Experimental bounds on and future sensitivities to the heavy lepton mixing in ν_e (panel (a)), ν_μ (panel (b)) and ν_τ (panel (c)). The black lines show the predicted mixing of heavy leptons $N_i^\pm$ $(i = 1, 2, 3)$ in a given light neutrino flavor state as a function of M_i. We take $(Y_1, Y_2, Y_3) \simeq \left(2 \times 10^{-5}, 0.007, 1\right)$. Solid (dashed) parts of the black lines are expectations in the allowed (excluded) region.

The condition in eq. (3.20) makes discovery of the RH gauge bosons and scalars at present colliders unfeasible. The leptons $N_2^\pm$ and $N_3^\pm$ are beyond the reach of current and future experiments as well (see the lines in all panels of figure 3). However, as can be seen from the left and middle panel of figure 3, the lightest pair of neutral leptons is accessible to future colliders [42–44], beam-dump experiments [45] and neutrino oscillation facilities [46]. In particular, SHiP may improve the lower bound (3.17) on M_1 to approximately 5 GeV, whereas FCC-ee will be able to probe even larger masses (up to 60 GeV) and very tiny mixing angles. Hadron collider FCC-hh with the total center of mass energy around 100 TeV will be able to search for $N_2^\pm$ (see line denoted *FCC* in all panels of figure 3) and also $N_3^\pm$ in the absence of screening [47], as well as the RH gauge bosons and new scalar bosons from the Higgs doublets [48].

Without imposing the screening condition, the strongest *direct* bound on the mass of $N_2^\pm$ follows from its admixture in ν_μ

$$\left|U^N_{\mu 2}\right|^2 \approx \frac{1}{2}\left(\frac{m_{D2}}{M_2}\right)^2 \approx \frac{m_c^2}{2M_2^2}\,. \tag{3.21}$$

The regions excluded by DELPHI [49] and CMS [50] experiments give $M_2 > 70$ GeV (see $N_2^\pm$ line in the middle panel of figure 3). Decreasing m_{D2} by a factor of 3 relaxes the bound down to $M_2 > 40$ GeV.

The admixture of $N_2^\pm$ in ν_e is characterized by

$$\left|U^N_{e2}\right|^2 \approx \frac{1}{2}\left(\frac{m_{D2}}{M_2}\right)^2 \sin^2\theta_c \approx \frac{m_c^2}{2M_2^2}\sin^2\theta_c\,, \tag{3.22}$$

and according to figure 3 (left panel) this leads to $M_2 > 50$ GeV.

The strongest *direct* bound on the mass of $N_3^\pm$, $M_3 > 100$ GeV, is established by ATLAS. Note that the *direct* bounds on $N_2^\pm$ and $N_3^\pm$ are significantly weaker than those obtained from the bound on $N_1^\pm$ and the screening relation (see eqs. (3.17) and (3.20)).

Thus, in the absence of screening, the hierarchy of the heavy leptons can be much weaker. We will elaborate on such a scenario in section 4.

Let us note that the neutral leptons, $(N^+ + N^-)/\sqrt{2}$, are produced dominantly via the mixing with active neutrinos in the processes involving the left-handed gauge bosons [51]. The production of N^- and N^+ via the exchange of the off-shell right-handed bosons in e^+e^-, $p\bar{p}$ and pp collisions is subdominant as the corresponding cross-sections are suppressed by a factor $\xi^4 = \left(\langle\phi_1^0\rangle/\langle\chi_R^0\rangle\right)^4$.

3.2 $0\nu2\beta$ decay

The dominant contribution to the neutrinoless double beta $(0\nu2\beta)$ decay arises from the left-handed current since the right-handed current contribution scales as $\xi^4 \lesssim 10^{-12}$ [52–54] due to the bound (3.19). Then, the effective Majorana mass of the electron neutrino can be presented as

$$m_{ee} = m_{ee}^l + m_{ee}^h\,, \tag{3.23}$$

where the contributions from the light and heavy mass eigenstates read

$$\begin{aligned} m_{ee}^l &= \sum_{i=1,2,3} (U_{\mathrm{PMNS}\,ei})^2\, m_i \approx s_\xi^2\, \mu_{ee}\,, \\ m_{ee}^h &= r\, p^2 \sum_{i=1,2,3} \left((U_{ei}^{N^-})^2 \frac{M_i}{p^2 - M_i^2} + (U_{ei}^{N^+})^2 \frac{M_i'}{p^2 - M_i'^2} \right) \equiv \sum_{i=1,2,3} m_i^h\,. \end{aligned} \tag{3.24}$$

Here, $r \equiv M^l_{\beta\beta0\nu}/M^h_{\beta\beta0\nu}$ is the ratio of nuclear matrix elements for the exchange of heavy and light neutrinos, $-p^2 \sim (125\ \mathrm{MeV})^2$ is the neutrino momentum squared and μ_{ee} denotes the $(1,1)$ element of the μ matrix.

The contribution from the i-th pair of pseudo-Dirac neutrinos (N_i^- and N_i^+) can be estimated as

$$m_i^h = \xi^2 \mu_{ee} \frac{p^2}{M_i^2} \sim m_{ee}^l \frac{p^2}{M_i^2}\,, \tag{3.25}$$

where ξ^2 arises from the admixture of $N_i^\pm$ in ν_e and μ_{ee} stems from the sum $M_i^+ + M_i^-$ (see eq. (3.9)). Due to strong mass hierarchy, the contributions from the heavier pairs are negligible and only the lightest one (N_1^- and N_1^+) should be considered. Hence, $m_{ee}^h \approx m_1^h$.

To compute m_1^h one should retain the 1-2 and 1-3 blocks in the matrix (3.8) and perform further rotations to diagonalize it up to $\mathcal{O}(\mu^2)$. The terms proportional to the elements of μ matrix can not be neglected and, in particular, deviations of the $N_L - S$ mixing from 45° should be taken into account in eq. (3.12). To diagonalize the matrix in eq. (3.8) we perform two additional rotations:

$$\mathcal{U}_{14} = \begin{pmatrix} \mathbb{1} & U_{14} & 0 \\ -U_{14} & \mathbb{1} & 0 \\ 0 & 0 & \mathbb{1} \end{pmatrix}, \qquad \mathcal{U}_{17} = \begin{pmatrix} \mathbb{1} & 0 & U_{17} \\ 0 & \mathbb{1} & 0 \\ -U_{17} & 0 & \mathbb{1} \end{pmatrix}, \tag{3.26}$$

with

$$U_{14} \approx U_{17} = \begin{pmatrix} -\frac{s_\xi c_\xi}{\sqrt{2}} \frac{\mu_{ee}}{M_1} & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}. \tag{3.27}$$

Now the total mixing matrix equals

$$\mathcal{U}_f = \mathcal{U}_l^\dagger\, \mathcal{U}_S\, \mathcal{U}_N\, \mathcal{U}_\nu\, \mathcal{U}_{14}\, \mathcal{U}_{17}\,. \tag{3.28}$$

According to eq. (3.28), the flavor neutrino states can be expressed in terms of mass eigenstates

$$\nu_f = c_\xi\, U_{\rm PMNS}\, \nu + \left(c_\xi\, U_{\rm PMNS}\, U_{14} - s_\xi\, s_N\, U_l^\dagger\right)\, N_1^- + \left(c_\xi\, U_{\rm PMNS}\, U_{17} + s_\xi\, c_N\, U_l^\dagger\right)\, N_1^+\,, \tag{3.29}$$

where we neglected the small correction to the first term. Then, the admixtures of N_1^- and N_1^+ in ν_e are given explicitly by

$$\mp \frac{s_\xi}{\sqrt{2}}\left[1 \pm \frac{c_\xi^2 \mu_{ee}}{M_1}\left((U_{\rm PMNS})_{e1} - \frac{1}{4}(U_l^\dagger)_{e1}\right)\right]\,, \tag{3.30}$$

where the upper (lower) sign corresponds to $U_{e1}^{N^-}$ $(U_{e1}^{N^+})$. After inserting these expressions into eq. (3.24), we find

$$m_1^h = -2\,(U_{\rm PMNS})_{e1}\,(U_l^\dagger)_{e1}\,\frac{|p^2|}{M_1^2}\,s_\xi^2\,c_\xi^2\,\mu_{ee} = -2(U_{\rm PMNS})_{e1}\,(U_l^\dagger)_{e1}\,\frac{|p^2|}{M_1^2}\,m_{ee}^l\,. \tag{3.31}$$

Thus, the mass m_1^h is proportional to the contribution from light neutrinos, m_{ee}^l, with the additional suppression factor $|p^2|/M_1^2$, in agreement with the estimate in eq. (3.25). The ratio of the two contributions equals

$$\frac{m_1^h}{m_{ee}^l} = 1.67\,\frac{|p^2|}{M_1^2}\,. \tag{3.32}$$

For $M_1 = 0.4\,\mathrm{GeV}$ and $M_1 = 2\,\mathrm{GeV}$ this suppression is 0.1 and 0.004, respectively. Hence, the heavy lepton contribution is practically not observable in the $0\nu2\beta$ decay experiments.

3.3 Leptogenesis

The inverse seesaw mechanism with its pairs of quasi-degenerate heavy leptons appears, at a first sight, very suitable for realization of low scale leptogenesis mechanisms. Those include the resonant leptogenesis [55, 56], where the lepton asymmetry is produced from the decays of quasi-degenerate heavy leptons, and the ARS mechanism based on oscillations between heavy neutral leptons [57]. A vast majority of the previous studies on leptogenesis in the inverse seesaw framework have been performed in the SM gauge structure. In the case of L-R model with the q-l similarity we have large Yukawa couplings and additional processes which enhance the washout effect.

In its minimal setup, the inverse seesaw model is not compatible with successful thermal leptogenesis due to strong washout [58, 59]. The maximal achievable value of the baryon asymmetry is 8 orders of magnitude smaller than the observed one [60].

In our L-R symmetric setup, similar conclusions apply. Furthermore, what already prevents any possibility for viable leptogenesis in our model is the size of the Yukawa couplings chosen according to the q-l similarity. As a result, the out-of-equilibrium condition is not satisfied for all pairs of heavy leptons. In particular, for the second pair with $M_2 \sim$TeV there is no deviation from thermal equilibrium for couplings $Y_2 \gtrsim \mathcal{O}(10^{-7})$. For $Y_2 \lesssim \mathcal{O}(10^{-7})$ we obtain $(N^{\pm} - N^{\rm eq})/N_{\rm eq} \gtrsim 10^{-1}$ at $M/T \gtrsim \mathcal{O}(1)$. This should be compared with $Y_2 \sim 5 \cdot 10^{-3}$ in our model.

The observed baryon asymmetry can be produced in the non-minimal realization of the inverse seesaw [60, 61] which is motivated by explanation of the scale of Majorana masses μ. The extension includes an extra fermion singlet X with large Majorana mass $M_X \gg M_i$ and the Higgs singlet σ which develops a VEV. The authors [60] employ the additional global symmetry under which S is charged. This is, however, not possible in our scenario since S, σ and X are not charged under any global symmetry.

The leptogenesis mechanisms discussed so far either fail in producing the required amount of asymmetry or are not compatible with our model setup. For alternative and potentially successful low-scale leptogenesis scenarios (see below) it is necessary to avoid strong washout. At $T \sim M_N$, the lepton number violating process $L_L\Phi \to L_L^c\Phi^\dagger$ mediated by the Majorana fermions N^- and N^+ becomes operative. The strength of the washout is quantified by the parameter K_i which, for a given fermion generation i, reads

$$\frac{\Gamma_i}{H(T=M_i)}\left(\frac{\mu_i}{\Gamma_i}\right)^2, \qquad \text{for } \mu \ll \Gamma_i\,, \tag{3.33a}$$

$$\frac{\Gamma_i}{H(T=M_i)}, \qquad \text{for } \mu \geq \Gamma_i\,. \tag{3.33b}$$

Here, $\Gamma_i = Y_i^2 M_i/(8\pi)$ is the decay rate of $N_i^{\pm}$. The additional factor $(\mu_i/\Gamma_i)^2$ in eq. (3.33a) originates from the interference between diagrams containing quasi-degenerate particles (N^- and N^+) [62, 63].

Let us make general estimations by dropping the assumptions of the q-l similarity and screening. From the inverse seesaw formula we have $\mu = 2\,m_\nu M^2/(Y^2\langle\phi_1^0\rangle^2)$, and consequently

$$\frac{\mu}{\Gamma} = \frac{16\pi\, m_\nu}{\langle\phi_1^0\rangle^2}\frac{M}{Y^4} = 8.3\cdot 10^{-9}\,\frac{1}{Y^4}\left(\frac{m_\nu}{0.1\text{ eV}}\right)\left(\frac{M}{10^5\text{ GeV}}\right), \tag{3.34}$$

where for brevity we omitted flavor generation indices. For $m_\nu = 0.01\,$eV, $M = 200\,$GeV and $Y = 1.1\cdot 10^{-3}$, we obtain $\mu/\Gamma = 1$. For smaller Y, we should use expression (3.33b) which gives $K \gg 1$, implying a very strong washout unless $Y \lesssim \mathcal{O}(10^{-7})$. For $Y > 1.1{\cdot}10^{-3}$, the expression (3.33a) should be used. It can be rewritten as

$$K = 32\pi\frac{m_\nu^2 M}{Y^6 a\langle\phi_1^0\rangle^4} \simeq 1.95\cdot 10^{-5}\,\frac{1}{Y^6}\left(\frac{M}{10^5\text{ GeV}}\right)\left(\frac{m_\nu}{0.1\text{ eV}}\right)^2, \tag{3.35}$$

where $a = 1.66\sqrt{g_*}/M_{\rm Pl}$. The condition $K < 1$ gives even stronger bound on the coupling, namely $Y > 0.06$ for $M > 200\,$GeV. Together with the bound on the admixture of the heavy leptons in the flavor states $\xi = Y v_L/M < 10^{-2}$ (which is not applicable above TeV scale) this leads to a narrow range of the allowed parameters: $M = (10^3 - 10^5)\,$GeV and $Y > 0.1$, which is satisfied only for the third generation.

For the inverse seesaw, the ARS mechanism leads to successful leptogenesis if $Y \sim 10^{-7} - 10^{-6}$ [64]. For masses of N^- and N^+ at $\mathcal{O}(\text{TeV})$ the washout (3.33b) is suppressed. Note, however, that such parameter space is not in accord with $q - l$ similarity.

One can rely on the electroweak baryogenesis [65] which requires a strong first order electroweak phase transition. The rich scalar sector in our model may help to realize such a scenario. Still, the lepton number washout should be suppressed, since sphalerons partially carry the baryon asymmetry to the lepton section.

3.4 Corrections to the Higgs mass. Naturalness

In this section we will address a specific problem related to existence of heavy RH neutrinos (recall that for the third neutrino $M_3 \gg v_{EW}$) and their radiative corrections to the Higgs mass. It was noticed long time ago that the contribution to the Higgs mass from the loops formed by active neutrinos and heavy RH neutrinos in the type I seesaw mechanism equals [66] (see [67–69] for recent studies)

$$\delta m_H^2 \approx \frac{Y_i^2 M_i^2}{4\pi^2} = \frac{m_{\nu 3} M_3^3}{2\pi^2 \langle \phi_1^0 \rangle^2} \,. \tag{3.36}$$

Then, the condition that the correction δm_H^2 (3.36) does not exceed the Higgs mass itself ("naturalness"), $\delta m_H^2 \lesssim 10^4\,\text{GeV}^2$, leads to the upper bound on the mass [66]

$$M_i \lesssim \text{ few} \times 10^7\,\text{GeV}\,, \tag{3.37}$$

where $m_{\nu 3} \sim 0.1\,\text{eV}$ was used. This is smaller than the standard lower bound from thermal leptogenesis [70] (however, see [71]).

With $M_i \leq 10^5\,\text{GeV}$ (that we obtained in section 3.1), the bound (3.37) is satisfied. However, in the inverse seesaw, the correction to δm_H^2 becomes larger than in the type I seesaw model, being enhanced by the factor M_i/μ [72]

$$\delta m_H^2 \sim \frac{m_{\nu 3}\, M_i^4}{2\pi^2\, \mu_{\text{max}}\, \langle \phi_1^0 \rangle^2} \,. \tag{3.38}$$

Here, μ_{max} represents the largest entry of the μ matrix. The difference from expression (3.36) originates from different dependence of the neutrino mass on the mass of heavy leptons. Eq. (3.38) and the inequality $\delta m_H^2 < 10^4\,\text{GeV}^2$ give the bound $M_i < 10^4\,\text{GeV}$ for $\mu_{\text{max}} \sim \mathcal{O}(100)\,\text{keV}$. This bound is about 1 order of magnitude lower than M_3. At the same time, in our scenario (in contrast to, e.g., ref. [66]), there are other particles, gauge bosons and scalars, at the scale $10^5\,\text{GeV}$. Being bosons, these particles give corrections to the Higgs mass of the opposite sign and they can cancel contributions from the RH neutrinos. Let us consider such a possibility.

For the gauge boson contribution we consider the 1-loop corrections with heavy right-handed charged ($W_R^{\pm}$) and neutral (Z_R) vector bosons. To estimate δm_H^2 arising from the loops of $W_R^{\pm}$, Z_R and $N_3^{\pm}$ we use the 1-loop effective potential [73, 74]. Expanding the 1-loop potential and identifying the term containing two SM Higgs fields we obtain the condition required for a 1-loop cancellation between the fermion and gauge boson contributions to δm_H^2

$$\frac{Y_3^2}{g^2} \simeq \left(\frac{11}{8}\right)^{1/2} \approx 1.17\,. \tag{3.39}$$

In eq. (3.39), g is the SU(2) gauge coupling. In our framework, Yukawa couplings are fixed according to the q-l similarity condition and the gauge coupling for $\mathrm{SU}(2)_R$ equals the $\mathrm{SU}(2)_L$ one. The values of g and Y_3 at the top quark mass scale are 0.65 and 0.93, respectively, so that $Y_3^2/g^2 = 2.05$ which clearly fails to satisfy eq. (3.39). Even though the renormalization group effects may help a bit (we find 7% decrease of Y_3/g at $\langle \chi_R^0 \rangle \sim 10^5\,\mathrm{GeV}$) the cancellation can not be achieved without additional scalar contribution. Another possibility is that Y_3 differs from the top Yukawa coupling.

In the scalar sector a number of diagrams can contribute to the SM Higgs mass. If only one scalar with mass $m^2 \sim (c/2)\,\langle \chi_R^0 \rangle^2$ gives the dominant contribution to δm_H^2, the condition for complete cancellation of bosonic and fermionic 1-loop contributions to the SM Higgs mass reads

$$11\, g^4 - 8\, Y_3^4 + 2\, c^2 \simeq 0\,. \tag{3.40}$$

Note that for this cancellation one needs specific values of couplings in the scalar sector. There is no principle which would fix the couplings to such values implying that if the cancellation occurs it is accidental. If the full 1-loop cancellation is achieved, the dominant Higgs mass correction stems from 2-loop diagrams which can be estimated as [75]

$$\delta m_H^2 \sim \frac{g^2 Y_3^2}{(16\pi^2)^2} M_3^2 \approx \frac{g^2}{2(16\pi^2)^2} \langle \chi_R^0 \rangle^2 = 3.4 \cdot 10^5\,\mathrm{GeV}^2\,. \tag{3.41}$$

The value given in eq. (3.41) is only one order of magnitude larger than m_H^2, and therefore a rather acceptable level of fine-tuning is required. Note again that a small deviation from the q-l similarity can also solve the problem.

Let us underline that in our scenario new physics scale is just 3 orders of magnitude larger than the EW scale and hence we deal here with mild hierarchy in the scalar sector. We do not discuss the origin of this hierarchy, but we find that the hierarchy is supported by the price of moderate fine tuning. The hierarchy can be further weakened if we depart from the q-l similarity.

Of course, in the complete analysis one should take into account all corrections to the Higgs mass including corrections from the top quark. Note that in the dimensional regularization the leading corrections are proportional to the mass of particle propagating in the loop. Therefore, the correction from top and other EW scale particles are much smaller that those from heavy right-handed neutrinos.

4 Variation on the theme

4.1 Altering heavy fermion portal couplings

In the scenario described in section 3.1 the bound on the scale of L-R symmetry breaking has been obtained using the following three points: (i) the lower bound on the mass of the lightest heavy leptons $N_1^{\pm}$ (see eq. (3.17)), (*ii*) the screening (2.23), and (*iii*) the q-l similarity, eq. (2.27). The latter implies the equalities $M_3 = M_1 m_t/m_u$ and $Y_3 = 1$. Consequently,

$$\langle\chi_R^0\rangle \simeq \sqrt{2}\,M_3 = \sqrt{2}\,M_1\,\frac{m_t}{m_u}. \tag{4.1}$$

The scale of $\langle\chi_R^0\rangle$ can be reduced if the assumptions in eq. (2.23) and/or eq. (2.27) are relaxed. Let us discuss the two scenarios in which one of these assumptions is abandoned.

(i) Departing from the *q*-*l* similarity. We keep the screening which can be expressed as

$$Y_i = R_0\,Y_{Ri}, \tag{4.2}$$

where in general $R_0 \neq 1$. This implies

$$\xi_i \equiv \frac{m_{Di}}{M_{Di}} = \frac{\langle\phi_1^0\rangle}{\langle\chi_R^0\rangle}\frac{Y_i}{Y_{Ri}} = \frac{\langle\phi_1^0\rangle}{\langle\chi_R^0\rangle}R_0 \equiv \xi, \tag{4.3}$$

i.e. all ξ_i are equal and eq. (2.24) for m_ν is unchanged. Let us recall that ξ determines admixtures of heavy leptons in the flavor states. In the absence of the q-l similarity, the Dirac masses m_{Di} are free parameters.

The strongest bound on ξ is obtained from mixing of $N_1^{\pm}$ in ν_e since $\xi^2 = |U_{e1}^N|^2$. From figure 3 one finds $\xi \leq 0.006$ and $M_1 > 2\,\mathrm{GeV}$. This corresponds to $m_{D1} \leq 12\,\mathrm{MeV}$.

As the Dirac masses m_{Di} are not fixed, even the extreme scenario with all m_{Di} (and therefore M_i) being equal is possible. For $m_{Di} = 12\,\mathrm{MeV}$ all the heavy leptons are accessible to experiments. From the expression for the light neutrino masses (2.24), we find that for $\xi = 0.006$ and $m_{\nu 3} = 0.05\,\mathrm{eV}$ the entries of the Majorana mass μ should be $\mathcal{O}(1)\,\mathrm{keV}$.

The scale of L-R symmetry breaking can be substantially lowered in comparison to the value given in eq. (3.20). Namely, the bound can be as low as the bound from the direct RH gauge boson searches, which is roughly 3 TeV [50, 76].

(ii) Departing from the exact screening. Now we keep the q-l similarity only. In this case, $\xi_i \equiv m_{Di}/M_{Di}$ are different. Taking weaker hierarchy of M_i than that of m_{Di}, for instance $M_i = (2,\ 200,\ 5\cdot 10^3)\,\mathrm{GeV}$, we have $\xi_i = (1.15\cdot 10^{-3},\ \ 6.45\cdot 10^{-3},\ \ 3.45\cdot 10^{-2})$. The corresponding effective mixing parameters given by $|\xi_i|^2$ yield

$$2|U_{e1}^N|^2 \simeq 1.32\cdot 10^{-6}\,, \qquad 2|U_{\mu 2}^N|^2 \simeq 4.16\cdot 10^{-5}\,, \qquad 2|U_{\tau 3}^N|^2 \simeq 1.18\cdot 10^{-3}\,. \tag{4.4}$$

Now, $\langle\chi_R^0\rangle = \sqrt{2}\,M_3 \simeq 7\,\mathrm{TeV}$ for $Y_{R3} \simeq 1$, so that both additional neutral leptons and right-handed gauge bosons are accessible to LHC.

Due to departure from screening, the formula for light neutrino masses changes

$$m_\nu = \left(\frac{\langle\phi_1^0\rangle}{\langle\chi_R^0\rangle}\right)^2 R\,\mu\,R\,, \qquad\qquad R \equiv \mathrm{diag}\,(Y_1/Y_{R1},\ Y_2/Y_{R2},\ Y_3/Y_{R3})\,, \tag{4.5}$$

where R can be absorbed in redefinition of the matrix μ:

$$\mu_{ij} \to \mu'_{ij} = \mu_{ij}\,R_i\,R_j\,. \tag{4.6}$$

Numerically, for chosen M_i we obtain $R \simeq (0.033,\ 0.185,\ 1)$.

The appearance of R may, however, complicate the explanation of mixing pattern from symmetry arguments since now both Y and Y_R are non-trivially involved in the expression for the light neutrino mass matrix. Consequently, certain correlation between R and μ matrices should exist. Also, R would affect phenomenology of the heavy leptons (production, decay, etc.).

4.2 Left and right fermion singlets

So far we considered the scenario with single Majorana fermion S per generation, that is the common fermion S for the left and right sectors. This is consistent with L-R symmetry. Under P transformations we had $S_L \leftrightarrow (S^c)_R$. Minimal and logically straighforward extension of this scenario is a P-symmetric model with two independent singlets S_L and S_R for the left and the right sectors, respectively. This study allows to check wheather the L-R symmetry can be realized in the singlet sector.

Now the Yukawa interactions and mass terms read

$$\mathcal{L} \supset -\bar{L}_R Y \Phi^\dagger L_L - \bar{L}_R \tilde{Y} \tilde{\Phi}^\dagger L_L - \bar{S}^c_L Y_L \tilde{\chi}^\dagger_L L_L - \bar{S}^c_R Y_R \tilde{\chi}^\dagger_R L_R \\ - \frac{1}{2}\left[\bar{S}^c_L \mu_{LL} S_L + \bar{S}^c_R \mu_{RR} S_R + \bar{S}_L \mu_{LR} S_R\right] + \text{h.c.}\,. \tag{4.7}$$

Due to P-invariance, the relations (2.8) involving Y and $Y_{L(R)}$ still hold as before. Now, the Yukawa interactions in (4.7) are invariant with respect to global U(1) symmetry of the lepton number with charge prescription $L_g(L_L) = L_g(L_R) = 1$, $L_g(S_L) = L_g(S_R) = -1$, and zero charges for scalar fields. This symmetry forbids the Yukawa interactions $\bar{S}_{R(L)}\, \tilde{\chi}^\dagger_{L(R)} L_{L(R)}$. The symmetry is broken by the Majorana mass terms (second line in eq. (4.7)). The smallness of masses μ can be related to this breaking.

Invariance with respect to P transformation $S_L \leftrightarrow S_R$ would imply the following equalities: $\mu_{LL} = \mu_{RR}$, $\mu_{LR} = \mu^\dagger_{LR}$. However, as in the visible sector, masses can break parity, so that in general $\mu_{LL} \neq \mu_{RR}$. The breaking can be spontaneous if μ terms are generated by couplings of S with singlet scalars σ_L, σ_R σ_{LR}: $y_L \bar{S}^c_L S_L \sigma_L$, $y_R \bar{S}^c_R S_R \sigma_R$. Then, even if we impose P-symmetry which gives $y_R = y_L$, the scales of μ_{LL} and μ_{RR} can be different due to $\langle\sigma_L\rangle \neq \langle\sigma_R\rangle$, i.e. spontaneous violation of parity in the singlet sector. In this case $\mu_{LL} \propto \mu_{RR}$. With complicated singlet sector one can obtain also different structures of matrices μ_{LL} and μ_{RR}. Finally, the L-R symmetry may be explicitly broken in the S-sector. In what follows we will not specify origins of μ matrices, but consider *a priori* general structure of μ_{LL} and μ_{LR} assuming only that $\mu \ll v_{EW}$. As before, μ_{RR} is fixed via inverse seesaw by masses and mixing of light active neutrinos. Notice that in this extension the contribution to the $\chi^\dagger_L \tilde{\Phi} \chi_R$ coupling is generated by the loop diagram of figure 1 with μ substituted by μ_{LR}.

After the Higgs fields acquire VEVs, the mass matrix in the $(\nu_L,\ N_L,\ S_L,\ S^c_R)$ basis reads

$$\mathcal{M} = \begin{pmatrix} 0 & m_D & m'_D & 0 \\ m_D & 0 & 0 & M_D \\ m'_D & 0 & \mu_{LL} & \mu_{LR} \\ 0 & M_D & \mu^T_{LR} & \mu_{RR} \end{pmatrix}, \tag{4.8}$$

where all entries are 3×3 matrices and the expressions for m_D, m'_D and M_D are given in eq. (2.10). The procedure of diagonalization of this matrix is similar to the one outlined in section 3.1. First, we make a rotation in the $\nu_L - S^c_R$ plane by s_ξ. In the new basis $(\nu',\ N_L,\ S_L,\ S')$ we perform nearly maximal rotation in the $(N_L - S')$ plane. Then, in the rotated basis $(\nu',\ N^-,\ S_L,\ N^+)$, N^- and S_L are permuted and the pseudo-Dirac states (N^- and N^+) decouple. In the basis of light states (ν', S_L) the mass matrix reads

$$\begin{pmatrix} \mu_{RR}\, s_\xi^2 & c_\xi\, m'_D - s_\xi\, \mu_{LR} \\ c_\xi\, m'_D - s_\xi\, \mu_{LR} & \mu_{LL} \end{pmatrix} . \tag{4.9}$$

Interestingly, the decoupling of N^- and N^+ does not produce $\mathcal{O}(\mu_{ij}/M_D)$ corrections to this matrix.

Phenomenology of this extended version is largely identical to the one of the main scenario. In particular, properties of heavy pseudo-Dirac states formed now by ν_R and S_R are similar to ones presented in section 3.1. The light neutrino mass matrix is given by the $(1,1)$ element of (4.9), $\mu_{RR}\, s_\xi^2$, and the observed values of the neutrino masses are achieved for $\mu_{RR} \sim 10\,\text{keV}$ and $s_\xi^2 \lesssim 10^{-6}$.

The only substantial difference from the main scenario is the presence of three relatively light states S_L. Therefore, in what follows we will focus on new physics associated to S_L.

According to (4.7), the states S_L have Yukawa interactions with the left leptonic doublet L_L and heavy scalar doublet χ_L (4.7). The mass of χ_L is at the $\text{SU}(2)_R$ symmetry breaking scale and for estimations we will use $M_\chi = 3 \cdot 10^5\,\text{GeV}$. The Yukawa couplings Y_L are large, being $Y_3 \sim 1$ for the third generation. Furthermore, S_L mix with light (mostly active) neutrinos according to eq. (4.9). The mixing angles with light neutrinos are given by

$$\sin\theta_{Si} \approx \frac{1}{\mu_{LLi}} (c_\xi\, m'_{Di} - s_\xi\, \mu_{LRi}) , \tag{4.10}$$

where for simplicity we have taken the matrices m'_D, μ_{LL} and μ_{LR} to be diagonal. In the case of full cancellation in the above expression, S_L states and light neutrinos do not mix. The mixing gives additional contribution to the light neutrino masses

$$\delta m_\nu \approx \sin^2\theta_{Si}\, m_{LLi} . \tag{4.11}$$

Therefore, the condition that there is no significant contribution to the light neutrino masses from mixing with S_L gives the upper bound

$$\sin^2\theta_{Si} \ll \frac{m_\nu}{m_{LLi}} , \tag{4.12}$$

where $m_\nu \sim (0.01 - 0.02)\,\text{eV}$.

The states S_L are light sterile neutrinos and their properties (masses, mixing and interactions) are subject to strong cosmological bounds. At the same time, S_L can be a dark matter candidate or even a very light particle in meV–eV mass range with no observable contributions to the energy density of the Universe.

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The states S_L decay into three light neutrinos, $S_{iL} \to \nu\nu\bar{\nu}$, via mixing with light neutrinos and Z^0 exchange. In vacuum, the lifetime equals

$$\tau_i = 3.8 \cdot 10^{20}\,\text{sec} \left(\frac{10^{-6}}{\sin^2\theta_{Si}}\right)\left(\frac{10\text{ keV}}{\mu_{LLi}}\right)^5 . \tag{4.13}$$

Notice that, according to (4.12), for $\mu_{LLi} = 10\,\text{keV}$ and $\sin^2\theta_{Si} = 10^{-6}$ the contribution to the light masses equals $\sim 0.01\,\text{eV}$.

The heavier S_{iL} can also decay into lighter S_{jL} with the χ^0 exchange: $S_{jL} \to S_{iL}\nu\bar{\nu}$. Typical time for this mode is

$$\frac{1}{\Gamma_i} = 2.4 \cdot 10^{27}\text{sec}\left(\frac{10\text{ keV}}{\mu_{LLi}}\right)^5\left(\frac{M_\chi}{3 \cdot 10^5\text{ GeV}}\right)^4 \frac{1}{Y_j^2\, Y_i^2}\,, \tag{4.14}$$

which is much bigger than the decay time (4.13) via process where mixing in employed.

Thus, the lifetime of S_L with masses (10–100) keV is much larger than the age of the Universe. Therefore, these S_L can be candidates for the Dark matter particles if their appropriate number density is generated [77–83].

S_{iL} can be produced via mixing with active neutrinos and oscillations. In this case the conditions on parameters of S_{iL} to be a dark matter are similar to those in νMSM [81]. For masses in the ballpark of 10 keV, the mixing should to be $\sin^2\theta_S \lesssim 2.5 \cdot 10^{-11}$. The strongest limits on θ_S arise from X-ray searches [84, 85], Supernova 1987A [86, 87] and structure formation [88].

According to eq. (4.10), for $m'_D = 0$

$$\sin\theta_S \approx -s_\xi \frac{\mu_{LR}}{\mu}. \tag{4.15}$$

Then, for $s_\xi \simeq 10^{-3}$ eq. (4.15) yields $\mu_{LR} \lesssim 10^{-2}\mu$.

In addition, in our scenario S_{iL} can be produced in the process $\nu_i\bar{\nu}_j \to S_{kL}\bar{S}_{lL}$ with χ_L-exchange. Here, mixing is kept, so that S_{kL} and S_{lL} are the eigenstates of μ_{LL} matrix (with mixing due to μ_{LR} being neglected). Suppression of rate of this process, Γ_S, with respect to the rate of usual active neutrino reactions, Γ_ν, is

$$\frac{\Gamma_S}{\Gamma_\nu} = \frac{4}{g^4}\, Y_{iL}^2\, Y_{jL}^2\, |U_{ik}|^2 |U_{jl}|^2 \left(\frac{m_W}{M_\chi}\right)^4 \approx 1.14 \cdot 10^{-13} Y_{iL}^2 Y_{jL}^2\, |U_{ik}|^2 |U_{jl}|^2 \left(\frac{3 \cdot 10^5\text{ GeV}}{M_\chi}\right)^4 , \tag{4.16}$$

where U_{ik} is the mixing matrix elements of S_{kL} and g is the weak coupling constant.

Depending on values of masses and mixing of S_{kL} one can consider different possibilities. Let us consider two extreme cases.

1. If the L-R symmetry holds in the singlet sector, then mixing between S_L states is the same as the one for S_R, which is of the TBM type. Since the mixing is large, the ratio of rates (4.16) will be determined by the largest coupling $Y_{3L} = 1$ and the mixing matrix elements $U_{3k} = U_{\tau k}^{\text{TBM}}$. If S_{2L} is in the 10 keV range and other S_L states are lighter, then $|U_{32}|^2 = 1/6$ and from (4.16) we obtain the ratio of rates $\Gamma_S/\Gamma_\nu = 3 \cdot 10^{-15}$. Using this ratio we find that S_L exits the equilibrium with thermal bath at temperatures $T \sim 100\,\text{GeV}$.

Their density will be diluted due to decrease of number of degrees of freedom at lower temperatures. This dilution is, however, not enough and further suppression by factor of 30 is needed to match correct energy density of dark matter. This can be achieved if the reheating temperature, $T_{\rm reh}$, after inflation is below 100 GeV. Similar situation is for S_{3L} and S_{1L} which do not contribute substantially to the present energy density in the Universe due to smaller masses.

With such low $T_{\rm reh}$, however, it will be difficult to realize baryogenesis through leptogenesis, unless M_χ increases or the coupling Y_L decreases leading to higher S_L decoupling temperature.

2. If mixing in the S_L sector is absent, the DM component, taken to be S_{1L}, will have the smallest coupling $Y_{1L} = 10^{-5}$. (This scenario can be reconciled with spontaneous L-R symmetry introducing additional flavons σ). In this case, the ratio of rates (4.16) equals $\approx 10^{-13}\,|Y_{1L}|^2 = 10^{-23}$. Correspondingly, the decoupling (freeze-out) temperature will be $\sim 10^5$ GeV and the required number density of S_{1L} can be obtained at $T_{\rm reh} \sim 10^4$ GeV. This is much higher than in the previous case, and opens a possibility for low scale baryogenesis through leptogenesis.

Notice that in view of problems with generation of the number density via mixing in νMSM [81], the χ_L-exchange can be the main mechanism of the DM production, while the mixing is suppressed sufficiently to satisfy the bounds from $X-$ray observations.

5 Summary

The low scale left-right symmetric models accessible to the existing and planned colliders are at odds with generation of naturally small neutrino masses. The (1–100) TeV scale of the L-R symmetry breaking (and consequently, the scale of RH neutrino masses) requires very small Dirac masses of neutrinos which strongly break the natural condition of quark-lepton similarity $Y \sim Y_q$. This similarity can be retained using the inverse seesaw mechanism which requires introduction of three fermionic singlets as well as the left and the right handed Higgs doublets. These doublets break the L-R symmetry and provide the portal for interactions of singlets with the SM particles.

This setup allows to obtain the neutrino mixing of special form e.g. tribimaximal or bimaximal. Under the screening condition, $Y_R = Y$, the light neutrino mass matrix is proportional to the Majorana mass matrix μ of singlet fermions S. In turn, special form of μ can be governed by symmetry in the S-sector. This symmetry is generally broken by the other interactions in the model but we have shown that the corrections due to such breaking are small and do not destroy the structure of μ.

The contribution from the linear seesaw should be suppressed. If dominant, it would require unnaturally small Dirac neutrino mass terms with the structure that breaks the q-l similarity. Such suppression can be achieved by small VEV of the left-handed doublet.

The generic consequence of the inverse seesaw is the existence of three heavy pseudo-Dirac neutral leptons with a mass splitting of the order of μ. Under the conditions of q-l similarity and screening, these heavy leptons mix with active neutrinos with strength

$\xi \simeq \langle\phi_1^0\rangle/\langle\chi_R^0\rangle$ and have strongly hierarchical mass spectrum. Consequently, only the lightest states, $N_1^\pm$, are accessible to current and near-future experiments. From the present experimental searches, we obtain the bound on mass of $N_1^\pm$ to be $M_1 \geq (0.4-2)\,\text{GeV}$. This, in turn, leads to the lower bound on the scale of L-R symmetry breaking $\langle\chi_R^0\rangle >$ 200 TeV. Future experiments such as SHiP, DUNE and FCC-ee can strengthen these bounds significantly. The FCC-hh collider will be able to test the existence of other new particles: $N_2^\pm$, $N_3^\pm$, the RH gauge bosons and new scalars.

Contributions of the heavy leptons to the effective Majorana mass m_{ee} are suppressed, being proportional to the contributions of light states.

The leptogenesis scenarios do not yield the required amount of the baryon asymmetry in our model, primarily due to strong washout. However, with a certain extensions of the model (addition of new heavy fermion(s) X and another singlet S per generation) or departure from q-l similarity (ARS), the required lepton asymmetry can be produced. The electroweak baryogenesis is a viable option.

The corrections to the Higgs mass (δm_H^2), induced via loops of RH neutrinos are, for $\langle\chi_R^0\rangle = 200\,\text{TeV}$, about 4 orders of magnitude larger than the experimentally observed value ($m_H^2 \sim 10^4$ GeV^2). The contributions from gauge boson loops as well as additional scalars can lead to substantial reduction of this correction. We derived the conditions for the complete cancellation of the Higgs mass one-loop corrections. Admittedly, such cancellations would be accidental.

Finally, we considered a scenario with two singlets per generation: S_L in the left sector and S_R in the right sector. In such a case the L-R symmetry is explicit and furthermore the global lepton number can be introduced. This modification leads to the appearance of keV-scale leptons and the lightest of them can play the role of dark matter.

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