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Two loop neutrino model and dark matter particles with global $B - L$ symmetry

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Abstract. We study a two loop induced seesaw model with global $\mathrm{U}(1)_{B-L}$ symmetry, in which we consider two component dark matter particles. The dark matter properties are investigated together with some phenomenological constraints such as electroweak precision test, neutrino masses and mixing and lepton flavor violation. In particular, the mixing angle between the Standard Model like Higgs and an extra Higgs is extremely restricted by the direct detection experiment of dark matter. We also discuss the contribution of Goldstone boson to the effective number of neutrino species $\Delta N_{\mathrm{eff}} \approx 0.39$ which has been reported by several experiments.

Keywords: dark matter theory, neutrino theory

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Contents

1 Introduction 1

2 The two-loop radiative seesaw model 2
- 2.1 Model setup 2
- 2.2 Higgs potential 3
- 2.3 Constraints 3

3 Dark matter relics 5

4 Differences from the original model 9

5 Conclusions 10

1 Introduction

Even after the discovery of the standard model (SM) Higgs boson, there still exist some unsolved issues: the origin of neutrino masses and mixings, the nature of dark matter (DM), whether Higgs boson is the only elementary scalar particle or not, and so on. As for the neutrinos, the tiny mass scale is apparently different from the other sectors, i.e. the charged leptons and quarks. Hence many physicists believe there exist some mechanisms for neutrino mass generation which is different from the other fermion mass generation. One of elegant solutions is to generate the neutrino masses with radiative correction, which provides more natural explanation of its smallness. Moreover neutrinos often interact with some new mediating particles that can be frequently identified to be DM. Such kind of models have been proposed by many authors in refs. [1–35].

As for DM, its properties are being explored by various experiments such as direct detection and indirect detection experiments as well as Large Hadron Collider (LHC). For example, the current direct detection experiment LUX [36] tells us that the upper bound for the spin independent cross section is highly constrained to be $\mathcal{O}(10^{-46})$ cm^2 at around 50 GeV of DM mass. For indirect detection, the recent analysis of Fermi-LAT gamma-ray data has shown that there may be gamma-ray line peak near 130 GeV, which could be interpreted as annihilation or decay of DM [37–40]. AMS-02 experiment also has shown the anomaly in the positron fraction up to energy about 350 GeV, and its result is in good agreement with the previous PAMELA experiment [41, 42]. They also suggest that leptophilic DM [43–48] is preferable since PAMELA has reported the anti-proton-to-proton ratio which is consistent with the predicted background [49].

In this paper, we construct a two-loop radiative seesaw model with global $B-L$ symmetry at the TeV scale based on the paper [50].[1] We also analyze the multi-component DM properties, and we discuss their detectability in addition to the observed relic density [56, 57]. At the end, we discuss the discrepancy of the effective number of neutrino species $\Delta N_{\rm eff} \approx 0.39$ between theory and experiments which is recently suggested by ref. [58].

[1] See for example the recent works on local $B-L$ symmetries in non-supersymmetric theory [51–55].

JCAP06(2014)027

Particle	Q	u^c	d^c	L	e^c	N^c	S
$(\mathrm{SU}(2)_L, \mathrm{U}(1)_Y)$	$(\mathbf{2}, 1/6)$	$(\mathbf{1}, -2/3)$	$(\mathbf{1}, 1/3)$	$(\mathbf{2}, -1/2)$	$(\mathbf{1}, 1)$	$(\mathbf{1}, 0)$	$(\mathbf{1}, 0)$
$\mathrm{U}(1)_{B-L}$	$1/3$	$-1/3$	$-1/3$	-1	1	1	$-1/2$
$\mathbb{Z}_2$	$+$	$+$	$+$	$+$	$+$	$-$	$-$

Table 1. The particle contents and the charges for fermions.

Particle	Φ	η	χ	Σ
$(\mathrm{SU}(2)_L, \mathrm{U}(1)_Y)$	$(\mathbf{2}, 1/2)$	$(\mathbf{2}, 1/2)$	$(\mathbf{1}, 0)$	$(\mathbf{1}, 0)$
Y_{B-L}	0	0	$-1/2$	1
$\mathbb{Z}_2$	$+$	$-$	$+$	$+$

Table 2. The particle contents and the charges for bosons.

This paper is organized as follows. In section 2, we show our model and discuss the Higgs sector including the Higgs potential, its stability condition, S-T parameters and neutrino mass in the lepton sector. We analyze DM phenomenology in section 3 and the differences from our previous paper are summarized in section 4. Then finally we conclude in section 5.

2 The two-loop radiative seesaw model

2.1 Model setup

We revisit a two-loop radiative seesaw model [50] but with global $B-L$ symmetry.[2] We add three right-handed neutrinos N^c, three SM gauge singlet fermion S, a $\mathrm{SU}(2)_L$ doublet scalar η and $B-L$ charged scalars χ and Σ to the SM particles.[3] We do not need to add any $\overline{S}$ to avoid the gauge anomaly problem as in ref. [50] because of the global $B-L$ symmetry. Thus the particle contents are more economical as shown in table 1 and 2. The $\mathbb{Z}_2$ parity is also imposed as table 1 and 2 so as to forbid the type-I seesaw mechanism. As a consequence, the parity-odd particles N^c, S, and η can be DM candidates. The right handed neutrino N^c is naturally lighter than S and η because its mass is generated at one-loop level. The lightest one is stabilized by $\mathbb{Z}_2$ parity. The $\mathbb{Z}_6$ symmetry remains after the $B-L$ spontaneous breaking, and the $\mathbb{Z}_6$ charge of each particle is mathematically defined as $6(B-L)$ mod 6. The $\mathbb{Z}_6$ charge of S and χ is 3, and they can be called as odd particle under $\mathbb{Z}_6$ symmetry since their transformation is same with $\mathbb{Z}_2$ symmetry. Thus the stability of χ and S is assured by a remnant $\mathbb{Z}_6$ parity after the $B-L$ spontaneous breaking [50]. Although the remnant symmetry would be regarded as $\mathbb{Z}_2$ symmetry in a narrow meaning of the renormalizable model, the larger $\mathbb{Z}_6$ symmetry should be taken into account if higher dimensional operators such as $QQQL$ are considered.

[2] See another example in ref. [22].

[3] Note that several right-handed neutrinos N_i^c and SM gauge singlet fermions S_i are needed to induce the neutrino masses and mixing, but we do not have gauge anomaly problem. Adding multi-right-handed neutrinos are one of the minimal requirements to obtain the observed neutrino masses and mixing. Another way is to introduce two SU(2) doublet inert bosons, see e.g. ref. [16].

JCAP06(2014)027

The gauge invariant and renormalizable Lagrangian for Yukawa sector and Higgs potential are given by

$$\mathcal{L}_{\rm Y} = (y_\ell)_{\alpha\beta}\Phi^\dagger L_\alpha e^c_\beta + (y_\nu)_{\alpha i}L_\alpha\eta N^c_i + (y_N)_{ij}N^c_i\chi S_j + (y_S)_{ij}\Sigma S_i S_j + \text{h.c.}, \tag{2.1}$$

$$\begin{aligned}\mathcal{V} = {} & m_1^2\Phi^\dagger\Phi + m_2^2\eta^\dagger\eta + m_3^2\Sigma^\dagger\Sigma + m_4^2\chi^\dagger\chi + m_5[\chi^2\Sigma + \text{h.c.}] \\ & +\lambda_1(\Phi^\dagger\Phi)^2 + \lambda_2(\eta^\dagger\eta)^2 + \lambda_3(\Phi^\dagger\Phi)(\eta^\dagger\eta) + \lambda_4(\Phi^\dagger\eta)(\eta^\dagger\Phi) + \lambda_5[(\Phi^\dagger\eta)^2 + \text{h.c.}] \\ & +\lambda_6(\Sigma^\dagger\Sigma)^2 + \lambda_7(\Sigma^\dagger\Sigma)(\Phi^\dagger\Phi) + \lambda_8(\Sigma^\dagger\Sigma)(\eta^\dagger\eta) + \lambda_9(\chi^\dagger\chi)^2 \\ & +\lambda_{10}(\chi^\dagger\chi)(\Phi^\dagger\Phi) + \lambda_{11}(\chi^\dagger\chi)(\eta^\dagger\eta) + \lambda_{12}(\chi^\dagger\chi)(\Sigma^\dagger\Sigma),\end{aligned} \tag{2.2}$$

where the indices α, β, $i, j = 1-3$. We assume all the parameters are real.[4] The quartic couplings λ_1, λ_2, λ_6 and λ_9 have to be positive to stabilize the Higgs potential. While the scalars η and χ are assumed not to have a vacuum expectation value (VEV), the $B-L$ charged scalar Σ has the VEV $\langle\Sigma\rangle = v'/\sqrt{2}$ and is the source of the spontaneous global $B-L$ breaking. The VEV of Σ gives the masses to the singlet S. The active neutrino masses are obtained through two-loop level [50]. In general, we can choose a diagonal base of y_S and mass matrix of the right-handed neutrinos after the symmetry breaking.

2.2 Higgs potential

After the global $B-L$ and electroweak symmetry breaking, the scalar particles in the model mix each other and we need to rewrite them by mass eigenstates. Since the particle content for scalar in the model is same with that of ref. [50], the discussion of the potential is exactly same with the reference except the existence of the Goldstone boson G. The Φ^0 and Σ are given by

$$\Phi^0 = \frac{v+\phi^0(x)}{\sqrt{2}}, \qquad \Sigma = \frac{v'+\sigma(x)}{\sqrt{2}}e^{iG(x)/v'}. \tag{2.3}$$

and they are rewritten by the mass eigenstates h and H as

$$\begin{pmatrix}\phi^0\\ \sigma\end{pmatrix} = \begin{pmatrix}\cos\alpha & \sin\alpha\\ -\sin\alpha & \cos\alpha\end{pmatrix}\begin{pmatrix}h\\ H\end{pmatrix}. \tag{2.4}$$

The mass eigenstates of the other scalar particles are η^+, η_R, η_I, χ_R and χ_I where η_R and η_I are the real and imaginary parts of η^0, and χ_R and χ_I are the real and imaginary parts of χ. Their masses are expressed by m_η, m_{η_R}, m_{η_I}, m_{χ_R} and m_{χ_I} respectively. More detail is referred the ref. [50]. The requirements to obtain the proper vacuum $\langle\phi\rangle \neq 0$, $\langle\eta\rangle = \langle\chi\rangle = 0$, $\langle\Sigma\rangle \neq 0$ also have been discussed in the reference.

2.3 Constraints

There are some constraints we have to take into account. First, the radiative correction to gauge boson masses in the SM is constrained by the electroweak precision tests [60–62]. The constraint is expressed by S and T parameters, and can be rewritten in terms of a mass relation between neutral and charged component of η in the model. It is approximately given by

$$\sqrt{(m_\eta - m_{\eta_R})(m_\eta - m_{\eta_I})} \lesssim 133\ \text{GeV}. \tag{2.5}$$

as have discussed in ref. [50].

[4] If the parameters are allowed to be complex in general case, we may have darkogenesis similar to [59].

JCAP06(2014)027

Second, the constraint from neutrino masses and mixing is taken into account. The neutrino mass matrix is derived at two loop level and written as

$$(m_\nu)_{\alpha\beta} = \sum_{i=1}^{3} \left(y_\nu^T y_N^*\right)_{\alpha i} \Lambda_i \left(y_\nu^T y_N^*\right)_{\beta i}, \tag{2.6}$$

where the loop function Λ_i is defined as

$$\Lambda_i = \frac{m_{Si}}{4(4\pi)^4} \int_0^1 dx \int_0^{1-x} dy \frac{1}{x(1-x)} \Big[I\left(m_{Si}^2, m_{RR}^2, m_{RI}^2\right) - I\left(m_{Si}^2, m_{IR}^2, m_{II}^2\right) \Big], \tag{2.7}$$

with

$$I(m_1^2, m_2^2, m_3^2) = \frac{m_1^2 m_2^2 \log\left(\frac{m_2^2}{m_1^2}\right) + m_2^2 m_3^2 \log\left(\frac{m_3^2}{m_2^2}\right) + m_3^2 m_1^2 \log\left(\frac{m_1^2}{m_3^2}\right)}{(m_1^2 - m_2^2)(m_1^2 - m_3^2)}, \tag{2.8}$$

$$m_{ab}^2 = \frac{y m_{\eta_a}^2 + x m_{\chi_b}^2}{x(1-x)} \quad (a, b = R \,\mathrm{or}\, I). \tag{2.9}$$

Here the mass of S_i is given by m_{Si}. The particle S_i can obtain a mass after the $B-L$ symmetry breaking as one can see from the interaction in eq. (2.1). We use the Casas-Ibarra parametrization to express the Yukawa matrix with the constraint of neutrino masses and mixing [63]. Then the product of Yukawa matrix is written as

$$y_N^\dagger y_\nu = \sqrt{\Lambda}^{-1} C \sqrt{\hat{m}_\nu} U_{\mathrm{PMNS}}^\dagger, \tag{2.10}$$

where the matrix Λ is defined as $(\Lambda)_{ij} = \Lambda_i \delta_{ij}$, C is a complex orthogonal matrix which satisfies $C^T C = 1$, $\hat{m_\nu}$ is the diagonalized active neutrino mass matrix and U_{PMNS} is the Pontecorvo-Maki-Nakagawa-Sakata matrix. We need $\mathcal{O}(1)$ Yukawa couplings in order to produce the proper DM relic density as will be discussed later. This corresponds to $\Lambda \sim m_\nu \sim 10^{-10}$ GeV. In addition, for sum of active neutrino masses, the limit of $\sum m_\nu < 0.933$ eV at 95% confidence level is imposed from the cosmological observation [57].

Third, Lepton Flavor Violation (LFV) should be taken into account. The most stringent constraint comes from the LFV process $\mu \to e\gamma$. Note that the LFV process $\mu \to 3e$ would give a stronger constraint when the mass difference between the right-handed neutrino and charged scalar η^+ is sufficiently large [64]. The Branching Ratio (Br) of the process $\ell_\alpha \to \ell_\beta \gamma$ $(\alpha, \beta = e, \mu, \tau)$ is given by

$$\mathrm{Br}\left(\ell_\alpha \to \ell_\beta \gamma\right) = \frac{\alpha_{\mathrm{em}} \left| \left(y_\nu y_\nu^\dagger\right)_{\alpha\beta} \right|^2}{768\pi G_F^2 m_\eta^4} \mathrm{Br}\left(\ell_\alpha \to \ell_\beta \nu_\alpha \overline{\nu_\beta}\right), \tag{2.11}$$

where the right-handed neutrino masses are neglected. The latest limit for $\mu \to e\gamma$ is given by MEG experiment [65] as

$$\mathrm{Br}(\mu \to e\gamma) < 5.7 \times 10^{-13}, \tag{2.12}$$

at 90% confidence level. For example, if the matrix y_N is diagonal, the constraint of $\mu \to e\gamma$ imposes that the orthogonal matrix C should be almost unit matrix $C \sim 1$. In other words, we can see from the Casas-Ibarra parametrization that the product of the Yukawa matrix $y_\nu y_\nu^\dagger$ becomes almost diagonal since the PMNS matrix is cancelled. Thus it does not contribute to any $\ell_\alpha \to \ell_\beta \gamma$ processes.

JCAP06(2014)027

3 Dark matter relics

We have some DM candidates with odd under $\mathbb{Z}_2$ parity. They are the right-handed neutrinos N_i^c, singlet fermion S_i and neutral component of η. It is natural to choose the lightest right-handed neutrino as a DM since the mass is generated at one-loop level and lighter than the other candidates. Hereafter we call the DM as N_1 with the mass m_{N_1}. In addition to the right-handed neutrino DM, we have an extra DM candidate χ. This is because after the breaking of the global $B-L$ symmetry, we still have remnant discrete $\mathbb{Z}_6$ symmetry under which χ is odd. This guarantees the stability of χ. The lighter one of χ_R and χ_I can be the second DM, and we assume χ_R is DM. Thus we have two component DM of N_1 and χ_R. The assumption of the mass hierarchy $m_{N_1} < m_{\chi_R}$ is reasonable from the mass generation mechanism. The DM χ_R can annihilate into the other DM (right-handed neutrino), but cannot decay into the SM particles with the renormalizable interactions. They cannot be taken care independently when one computes each relic density since one DM annihilates into the other DM. The set of Boltzmann equations is written as

$$\frac{dn_N}{dt} + 3Hn_N = -\langle\sigma_N v\rangle \left(n_N^2 - n_N^{\mathrm{eq}2}\right) + \langle\sigma_{\mathrm{ex}} v\rangle \left[n_\chi^2 - \left(\frac{n_\chi^{\mathrm{eq}}}{n_N^{\mathrm{eq}}}\right)^2 n_N^2\right], \tag{3.1}$$

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle\sigma_\chi v\rangle \left(n_\chi^2 - n_\chi^{\mathrm{eq}2}\right) - \langle\sigma_{\mathrm{ex}} v\rangle \left[n_\chi^2 - \left(\frac{n_\chi^{\mathrm{eq}}}{n_N^{\mathrm{eq}}}\right)^2 n_N^2\right], \tag{3.2}$$

where the time of universe is expressed by t, n_N and n_χ are the number density of N_1 and χ_R respectively. The thermally averaged annihilation cross section into all channels is written as $\langle\sigma_N v\rangle$ for N_1. For χ_R, the total cross section into the SM particles is written by $\langle\sigma_\chi v\rangle$, and $\langle\sigma_{\mathrm{ex}} v\rangle$ implies the DM exchange process $\chi_R\chi_R \to NN$. If $\langle\sigma_{\mathrm{ex}} v\rangle$ is negligible compared with $\langle\sigma_\chi v\rangle$, the simultaneous Boltzmann equation becomes independent of each other, and the total relic density should be a sum of N_1 and χ_R: $\Omega_{N_1}h^2 + \Omega_{\chi_R}h^2$. If $\langle\sigma_{\mathrm{ex}} v\rangle$ is a main channel of χ_R annihilation, the most of χ_R annihilates into the other DM N_1, but the relic density of N_1 almost does not depend on the DM exchange process because N_1 is still in thermal equilibrium when χ_R is frozen-out. Therefore the effect of the DM exchange process is small, and the DM system can be treated as two independent DM as a good approximation. We have checked it numerically and the fact that effect of semi-annihilation of DM is typically a few percent supports our result [66]. The contours of satisfying $\Omega_{N_1}h^2 + \Omega_{\chi_R}h^2 = 0.12$ are shown in figure 1. The x-axis and y-axis are the total cross section of N_1 and χ respectively.

There are two channels for N_1 annihilation, which are $N_1N_1 \to \ell\bar{\ell}, \nu\nu$. The annihilation cross section for $\ell\bar{\ell}$ and $\nu\nu$ are given as follows in the leading power of the DM relative velocity v_{rel}:

$$\sigma v_{\mathrm{rel}} \left(N_1N_1 \to \ell\bar{\ell}\right) = \frac{\left[\left(y_\nu y_\nu^\dagger\right)_{11}\right]^2}{48\pi m_{N_1}^2} \frac{m_{N_1}^4 \left(m_{N_1}^4 + m_\eta^4\right)}{(m_{N_1}^2 + m_\eta^2)^4} v_{\mathrm{rel}}^2, \tag{3.3}$$

$$\sigma v_{\mathrm{rel}} \left(N_1N_1 \to \nu\nu\right) = \frac{\left[\left(y_\nu y_\nu^\dagger\right)_{11}\right]^2}{24 m_{N_1}^2} \frac{m_{N_1}^4 (m_{N_1}^4 + m_0^4)}{(m_{N_1}^2 + m_0^2)^4} v_{\mathrm{rel}}^2, \tag{3.4}$$

where m_0^2 is average mass between $m_{\eta_R}^2$ and $m_{\eta_I}^2$ and they are assumed to be degenerate. Such a small mass difference is required to obtain a proper neutrino mass scale as have discussed in ref. [50]. Note that for the annihilation into neutrinos the factor 2 larger than

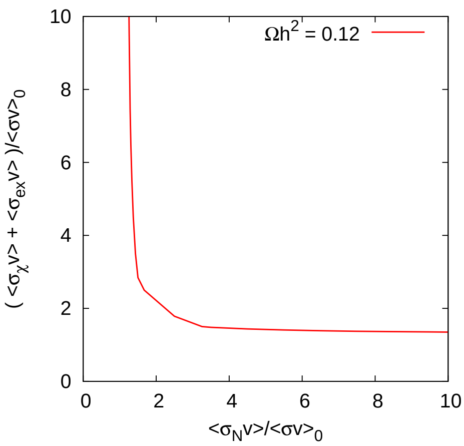

Figure 1. $\langle \sigma v \rangle_0$ is the typical scale of annihilation cross section 2.0×10^{-26} [cm^3/s].

for the charged leptons because of Majorana property of neutrinos. The mass matrix of the right-handed neutrinos m_N is generated at one loop level and the expression is found as

$$(m_N)_{ij} = \sum_k \frac{(y_N)_{ik}(y_N)_{jk} m_{Sk}}{(4\pi)^2} \left[\frac{m_{\chi R}^2}{m_{\chi R}^2 - m_{Sk}^2} \ln\left(\frac{m_{\chi R}^2}{m_{Sk}^2}\right) - \frac{m_{\chi I}^2}{m_{\chi I}^2 - m_{Sk}^2} \ln\left(\frac{m_{\chi I}^2}{m_{Sk}^2}\right) \right]. \quad (3.5)$$

The DM and the mediator η masses should be $10 \lesssim m_N \lesssim 60$ GeV and $100 \lesssim m_\eta, m_0 \lesssim 300$ GeV to reproduce the correct relic density of the observed value $\Omega h^2 \approx 0.12$ [57]. Otherwise the right-handed neutrino DM is overproduced since the cross section becomes too small as one can see from eq. (3.3) and (3.4). We also should note that there is the constraint from slepton search in SUSY models via the decay into lepton and missing energy at LHC [67–69]. In our case, the charged scalar η^+ has a similar properties with sleptons. The constraint on the mass of slepton is roughly $m_{\tilde{\ell}} \gtrsim 270$ GeV. Although this is not needed to be fully considered in our case, one minds such the matter.

For χ_R annihilation, there are five channels: $\chi_R\chi_R \to hh$, ZZ, W^+W^-, $f\bar{f}$, GG. Each cross section is written by

$$\begin{aligned}
\sigma v(\chi_R\chi_R \to ZZ) &= \frac{g_2^2 m_Z^2}{4\pi s}\sqrt{1 - \frac{4m_Z^2}{s}} \left[3 - \frac{s}{m_Z^2} + \frac{1}{4}\left(\frac{s}{m_Z^2}\right)^2\right] \\
&\quad \times \left| \frac{\mu_{\chi\chi h}\cos\alpha}{s - m_h^2 + i m_h \Gamma_h} + \frac{\mu_{\chi\chi H}\sin\alpha}{s - m_H^2 + i m_H \Gamma_H} \right|^2, \qquad (3.6) \\
\sigma v(\chi_R\chi_R \to WW) &= \frac{g_2^2 m_W^2}{2\pi s}\sqrt{1 - \frac{4m_W^2}{s}} \left[3 - \frac{s}{m_W^2} + \frac{1}{4}\left(\frac{s}{m_W^2}\right)^2\right] \\
&\quad \times \left| \frac{\mu_{\chi\chi h}\cos\alpha}{s - m_h^2 + i m_h \Gamma_h} + \frac{\mu_{\chi\chi H}\sin\alpha}{s - m_H^2 + i m_H \Gamma_H} \right|^2, \qquad (3.7) \\
\sigma v(\chi_R\chi_R \to f\bar{f}) &= \frac{y_f^2}{2\pi}\left(1 - \frac{4m_f^2}{s}\right)^{3/2} \left| \frac{\mu_{\chi\chi h}\cos\alpha}{s - m_h^2 + i m_h \Gamma_h} + \frac{\mu_{\chi\chi H}\sin\alpha}{s - m_H^2 + i m_H \Gamma_H} \right|^2, \qquad (3.8)
\end{aligned}$$

$$\sigma v(\chi_R\chi_R \to hh) = \frac{1}{64\pi^2 s}\int \left| \frac{12\mu_{\chi\chi h}\mu_{hhh}}{s-m_h^2+im_h\Gamma_h} + \frac{4\mu_{\chi\chi H}\mu_{hhH}}{s-m_H^2+im_H\Gamma_H} \right.$$
$$\left. +\lambda_{10}\cos^2\alpha + \lambda_{12}\sin^2\alpha + \frac{4\mu_{\chi\chi h}^2}{t-m_{\chi_R}^2} + \frac{4\mu_{\chi\chi h}^2}{u-m_{\chi_R}^2} \right|^2 d\Omega, \tag{3.9}$$

$$\sigma v(\chi_R\chi_R \to GG) = \frac{1}{16\pi^2 s}\int \left| \frac{\mu_{\chi\chi h}\sin\alpha}{s-m_h^2+im_h\Gamma_h}\frac{s}{v'} - \frac{\mu_{\chi\chi H}\cos\alpha}{s-m_H^2+im_H\Gamma_H}\frac{s}{v'} \right.$$
$$\left. +\frac{\sqrt{2}m_5}{v'} - \frac{2m_5^2}{t-m_{\chi_I}^2} - \frac{2m_5^2}{u-m_{\chi_I}^2} \right|^2 d\Omega, \tag{3.10}$$

where s, t, u are the Mandelstam variables, and the cubic couplings $\mu_{\chi\chi h}$, $\mu_{\chi\chi H}$, μ_{hhh} and μ_{hhH} are given by

$$\mu_{\chi\chi h} = -\frac{m_5}{\sqrt{2}}\sin\alpha + \frac{\lambda_{10}}{2}v\cos\alpha - \frac{\lambda_{12}}{2}v'\sin\alpha, \tag{3.11}$$
$$\mu_{\chi\chi H} = \frac{m_5}{\sqrt{2}}\cos\alpha + \frac{\lambda_{10}}{2}v\sin\alpha + \frac{\lambda_{12}}{2}v'\cos\alpha, \tag{3.12}$$
$$\mu_{hhh} = \lambda_1 v\cos^3\alpha - \lambda_6 v'\sin^3\alpha + \frac{\lambda_7}{2}v\sin^2\alpha\cos\alpha - \frac{\lambda_7}{2}v'\sin\alpha\cos^2\alpha, \tag{3.13}$$
$$\mu_{hhH} = 3\lambda_1 v\sin\alpha\cos^2\alpha + 3\lambda_6 v'\sin^2\alpha\cos\alpha$$
$$+\frac{\lambda_7}{2}v\sin^3\alpha - \lambda_7 v\sin\alpha\cos^2\alpha - \lambda_7 v'\sin^2\alpha\cos\alpha + \frac{\lambda_7}{2}v'\cos^3\alpha. \tag{3.14}$$

As have discussed in ref. [50], we need a large mass difference between χ_R and χ_I in order to obtain a proper scale of active neutrino masses. Hence the parameter relation is roughly estimated as

$$\frac{m_5 v'}{m_{\chi_R}^2} \gtrsim \mathcal{O}(1), \tag{3.15}$$

The origin of the mass difference is the cubic coupling m_5. Thus we can see that the cubic couplings $\mu_{\chi\chi h}$ and $\mu_{\chi\chi H}$ tend to be large compared with the other couplings. In this case, the annihilation channels into gauge bosons (ZZ and WW) become dominant over the other channels because of the longitudinal mode of the gauge bosons unless $\sin\alpha$ is extremely small. The cross section is roughly $\sigma v \sim 10^{-24}$ cm^3/s when $\sin\alpha \sim 1$. As we will discuss later, such a large mixing angle is excluded by direct detection of DM and the invisible decay mode of the SM-like Higgs. In this case, the most of the DM χ_R disappears at the early universe and only the right-handed neutrino DM remains. On the contrary, the annihilation channels into the Higgs and Goldstone boson become dominant when the mixing is small such as $\sin\alpha \lesssim 0.01$. The DM exchange channel $\chi_R\chi_R \to N_1N_1$ also may be a leading channel. The cross section of the process $\chi_R\chi_R \to N_1N_1$ is found as

$$\sigma_{\rm ex}v_{\rm rel}\left(\chi_R\chi_R \to N_1N_1\right) \approx \sum_i \left[\frac{(y_N)_{1i}^4}{8\pi m_{\chi_R}^2}\frac{\mu_i}{(1+\mu_i)^2} - \frac{(y_N)_{1i}^4}{24\pi m_{\chi_R}^2}\frac{\mu_i(1+3\mu_i)}{(1+\mu_i)^4}v_{\rm rel}^2\right], \tag{3.16}$$

where $\mu_i = m_{\chi_R}^2/m_{Si}^2$. Notice here that the above cross section is the massless limit of the final state particles.

Next we discuss detectability of the two DM candidates. For the case of the scalar DM, it would be possible to detect it by direct search if the cubic or quartic couplings in the scalar

JCAP06(2014)027

potential are $\mathcal{O}(1)$.[5] The Higgs exchange is a primary channel because the scalar DM does not have direct interactions with the SM particles except the Higgs potential. Thus χ_R is so-called Higgs portal DM [70–77]. Since the term with m_5 is dominant in eq. (3.11) and eq. (3.12), the spin independent elastic scattering cross section with proton is written by

$$\sigma_{p\text{-}\chi_R} \approx \frac{c}{8\pi}\frac{m_p^4 m_5^2 \sin^2 2\alpha}{(m_{\chi_R}+m_p)^2 v^2}\left(\frac{1}{m_h^2}-\frac{1}{m_H^2}\right)^2, \tag{3.17}$$

where $m_p = 938\,\text{MeV}$ is the proton mass and $c \approx 0.079$ is a coefficient that is determined by the lattice simulation [78, 79][6]. The stringent constraint can be obtained by LUX experiment that tells us $\sigma_{p\text{-}\chi_R} \lesssim 7.6 \times 10^{-46}$ cm^2 at $m_{\chi_R} \approx 33\,\text{GeV}$ [36] where we implicitly assumed χ_R is dominant component of DM in this estimation. When $m_p \ll m_{\chi_R}$ and $m_h \ll m_H$, the conservative limit is given by

$$\frac{m_5 \sin 2\alpha}{m_{\chi_R}} \lesssim 0.11. \tag{3.18}$$

It is not difficult to satisfy this relation because the mixing angle $\sin\alpha$ should be sufficiently small in order to be dominant in two DM system. Otherwise the DM χ_R becomes sub-dominant.

When $m_h \gg m_H$, further smaller mixing angle $\sin\alpha$ is required since the elastic cross section is enhanced by the light Higgs. However, it is interesting to consider such a light extra Higgs because it is correlated with the additional contribution of the Goldstone boson G to the effective number of neutrino species N_{eff} [58]. The discrepancy of the effective number of neutrino species ΔN_{eff} has been reported by several experiments such as Planck [57], WMAP9 polarization [81], and ground-based data [82, 83], which tell us $\Delta N_{\text{eff}} = 0.36 \pm 0.34$ at the 68 % confidence level. The Goldstone boson G may contribute to the effective neutrino number ΔN_{eff} if the period of freezing out of the particle is suitable. The appropriate era of freeze-out of the Goldstone boson is before muon annihilation while the other SM particles are decoupled, thus it corresponds to $T \approx m_\mu$ where T is the temperature of the universe. The scattering of the Goldstone boson with the SM particles occurs through the Higgs exchange. The interaction rate should be same order with the Hubble parameter H when $T \approx m_\mu$. From the rough evaluation of the reaction rate of G and the Hubble parameter, we obtain the condition

$$\frac{\sin^2 2\alpha (m_h^2 - m_H^2)^2}{4(vv')^2}\frac{m_\mu^7 m_{\text{pl}}}{m_h^4 m_H^4} \approx 1, \tag{3.19}$$

where $m_{\text{pl}} \approx 1.2 \times 10^{19}$ GeV is the Planck mass and m_μ is muon mass. Typically the extra Higgs boson should be light to satisfy this relation. As have discussed in ref. [58], the invisible decay mode $h \to GG$ also constrains the mixing angle $\sin\alpha$. However, we found that the constraint from the direct detection is stronger.

Combining with eq. (3.17) and (3.15), the following constraint on elastic cross section is obtained to get a certain value of ΔN_{eff}

$$\sigma_{p-\chi_R} \approx \frac{c}{2\pi}\frac{m_p^4 m_5^2 v'^2}{m_{\chi_R}^2 m_\mu^7 m_{\text{pl}}} \gtrsim \frac{c}{2\pi}\frac{m_p^4 m_{\chi_R}^2}{m_\mu^7 m_{\text{pl}}}. \tag{3.20}$$

[5] The right-handed neutrino DM also may be detected through one loop photon exchange interaction if the Yukawa matrix y_ν is complex and the mass is degenerate with the second right-handed neutrino [6].

[6] When $m_h \approx m_H$, there is cancellation between the two terms in eq. (3.17). And we can easily evade the direct detection bound when the mixing angle $\alpha \lesssim 0.4$ coming from the LHC Higgs searches [80].

JCAP06(2014)027

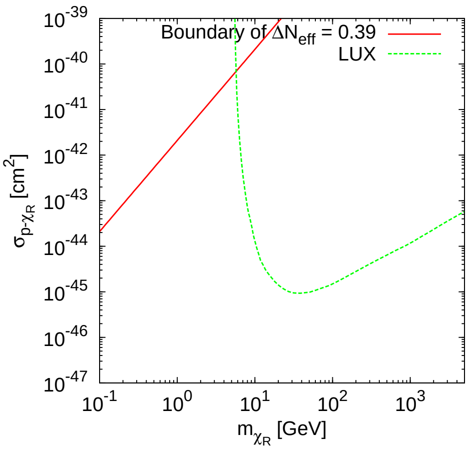

Figure 2. The effective number of neutrinos and the constraint from LUX experiment.

This requirement is shown with the limit of LUX experiment [36] in figure 2. The upper left region of the red line implies the region that $\Delta N_{\rm eff} \approx 0.39$ can be derived as ref. [58]. Such a large elastic cross section is obtained when the extra Higgs boson H is quite lighter than the SM-like Higgs h. The lower right region corresponds too fast deviation of the Goldstone boson from thermal bath and the contribution to $\Delta N_{\rm eff}$ is negligible. As the figure, when we consider the case of $m_H \ll m_h$, we can get upper bound on the DM mass $m_{\chi_R} \lesssim 5.5\,{\rm GeV}$ from the LUX experiment [36]. Therefore this result would contradict with the above discussion of thermal relics of DM since we have assumed $m_{N_1} < m_{\chi_R}$. However if χ_R is a sub-dominant component DM, the constraint of LUX experiment is moderated. Or, we could also consider a light DM scenario such as $m_{\chi_R} < m_N$. Although we need a little fine-tuning for quartic couplings of the Higgs potential is needed to obtain such a light Higgs mass, it is not difficult to obtain sizable $\Delta N_{\rm eff} \approx 0.39$ [58, 84].

4 Differences from the original model

This work includes some similar parts with our previous paper [50]. Thus it is better that the main differences with our previous work are made clear. In this paper, we suppose that the $\mathrm{U}(1)_{B-L}$ symmetry is global, unlike our previous paper [50] that has been discussed in the local gauged symmetry. We do not need to introduce additional $\overline{S}$ as the previous paper to cancel gauge anomaly, thus the new model is more economical. More detail DM phenomenology with two-component DM was discussed in this paper. In particular due to the global symmetry, we have a Goldstone boson G that provides the feasibility of the observed discrepancy of the effective number of neutrino species $\Delta N_{\rm eff} \approx 0.39$ in a similar way of Weinberg model [58].

In addition to the above things, there is another expectation for the global case. In ref. [50], the DM candidates have been N_1 and $\overline{S}$ which are fermion both. Then since their annihilation cross sections have been p-wave suppressed, there has been no detectability for indirect detection. On the other hand, in this paper the scalar DM χ_R is included and it has s-wave in general. The scalar DM χ_R has the mass of $\mathcal{O}(10)\,{\rm GeV}$ from the view of the neutrino effective number $\Delta N_{\rm eff}$. Therefore the recently discussing gamma-ray excess below $10\,{\rm GeV}$ would be explained well by the scalar DM if the dominant annihilation channel is $\chi_R\chi_R \to \tau\overline{\tau}$ [85, 86].

JCAP06(2014)027

5 Conclusions

We have constructed a two-loop radiative seesaw model with global $B - L$ symmetry at the TeV scale, which provides neutrino masses with more natural parameters. Various phenomenological constraints such as S-T parameters, neutrino masses and mixing and lepton flavor violation, stability of Higgs potential have been taken into account. The Casas-Ibarra parametrization for the neutrino Yukawa matrix have been used to describe the lepton flavor violating process $\mu \to e\gamma$.

We have studied the multi-component DM properties with fermion N_1 and scalar boson χ_R. The mass of N_1 is generated at one-loop level, thus $m_{N_1} < m_{\chi_R}$ is natural. The set of the Boltzmann equation for N_1 and χ_R is solved simultaneously. For the relic density of N_1, a large Yukawa coupling $\mathcal{O}(1)$ is required to reduce the abundance appropriately. On the other hand, the relic density of χ_R depends on the Higgs mixing angle α. In case of large mixing angle α, χ_R component DM can be sub-dominant since the cross section becomes quite large. In case of small α, χ_R component can be dominant.

It would be also possible to detect the scalar DM by direct search through Higgs exchange elastic scattering if the cubic or quartic couplings in the scalar potential are sufficiently large, since an elastic scattering occurs with quarks via Higgs exchange. The Higgs mixing angle α is extremely constrained by the latest direct search experiment LUX, in particular when the extra Higgs boson is much lighter than the SM-like Higgs.

At the end, we have discussed the discrepancy of the effective number of neutrino species, ΔN_{eff} between theory and experiments. We found that light extra Higgs and small mixing angle is needed to obtain $\Delta N_{\text{eff}} \approx 0.39$. Moreover, the scalar DM mass is quite limited as $m_{\chi_R} \lesssim 5.5\,\text{GeV}$ when we consider the current direct detection search of LUX.

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JCAP06(2014)027