Explaining the 3.5 keV X-ray Line in a $L_\mu - L_\tau$ Extension of the Inert Doublet Model

Anirban Biswas*,1, Sandhya Choubey,1, Laura Covi,3 and Sarif Khan1,4

1Harish-Chandra Research Institute, HBNI, Chhatnag Road, Jhunsi, Allahabad 211 019, India
2Department of Physics, School of Engineering Sciences, KTH Royal Institute of Technology, AlbaNova University Center, 106 91 Stockholm, Sweden
3Institute for Theoretical Physics, Georg-August University Göttingen, Friedrich-Hund-Platz 1, Göttingen, D-37077 Germany

Abstract

We explain the existence of neutrino masses and their flavour structure, dark matter relic abundance and the observed 3.5 keV X-ray line within the framework of a gauged $U(1)_{L_\mu - L_\tau}$ extension of the “scotogenic” model. In the $U(1)_{L_\mu - L_\tau}$ symmetric limit, two of the RH neutrinos are degenerate in mass, while the third is heavier. The $U(1)_{L_\mu - L_\tau}$ symmetry is broken spontaneously. Firstly, this breaks the $\mu - \tau$ symmetry in the light neutrino sector. Secondly, this results in mild splitting of the two degenerate RH neutrinos, with their mass difference given in terms of the $U(1)_{L_\mu - L_\tau}$ breaking parameter. Finally, we get a massive $Z_{\mu\tau}$ gauge boson. Due to the added $Z_2$ symmetry under which the RH neutrinos and the inert doublet are odd, the canonical Type-I seesaw is forbidden and the tiny neutrino masses are generated radiatively at one loop. The same $Z_2$ symmetry also ensures that the lightest RH neutrino is stable and the other two can only decay into the lightest one. This makes the two nearly-degenerate lighter neutrinos a two-component dark matter, which in our model are produced by the freeze-in mechanism via the decay of the $Z_{\mu\tau}$ gauge boson in the early universe. We show that the next-to-lightest RH neutrino has a very long lifetime and decays into the lightest one at the present epoch explaining the observed 3.5 keV line.

* Presently at Department of Physics, IIT Guwahati, Guwahati, Assam, India 781039
1Electronic address: anirban.biswas.sinp@gmail.com
2Electronic address: sandhya@hri.res.in
3Electronic address: laura.covi@theorie.physik.uni-goettingen.de
4Electronic address: sarifkhan@hri.res.in
I. INTRODUCTION

Experimental proof of non-zero neutrino masses and mixing as well as the dark matter in the universe remain the two most compelling evidences of the existence of physics beyond the standard model. Different neutrino oscillation experiments have confirmed the existence of flavour oscillations which can be explained only if neutrinos have tiny masses and mixing [1–3]. On the other hand, observations of the flatness of the galaxy rotation curve [4], gravitational lensing [5], cosmic microwave background anisotropy [6, 7] and more recently the observation of bullet cluster by NASA’s Chandra Satellite [8] demand that there must a non-baryonic component of matter in the universe, usually referred to as dark matter (DM). One of the most promising particle DM candidate is the Weakly Interacting Massive Particles (WIMP). However, the null result from the different earth and satellite-based direct and indirect DM searches have put severe constraints on the WIMP paradigm [9, 10]. One of the popular ways of shaking off the constraints from the direct and indirect DM searches is to postulate that the interaction strength of the DM with the standard model particles is extremely feeble. Such DM candidates go by the generic name Feebly Interacting Massive Particles (FIMP) [11–23]. Since their coupling with the standard model particles is feeble, they remain out of thermal equilibrium during the early universe when they are produced. Hence, these non-thermal particles are produced by the so-called freeze-in [11] mechanism instead of the freeze-out process which produces thermal relics.

More recently, the observation of an unknown 3.5 keV X-ray line in galaxies clusters [24, 25] and from the Galactic centre (GC) has been under much debate [26]. This excess has been confirmed by both the Chandra as well as NuSTAR satellites [26]. It has been argued that this signal can come from iron line background and S XVI charge exchange. Also such line has not been observed instead in stacked dwarf spheroidal galaxies [27], nor in galaxy spectra [28]. Nevertheless this signal has excited a lot of theoretical activity and can be explained by a plethora of theoretical models [29–58]. Generically it points towards a very weakly interacting Dark Matter like a light sterile neutrino, decaying into active neutrino and photon [59], but it can also arise from heavier DM particles in presence of mass degeneracy or from DM annihilation.

We address the issue of the observed neutrino masses and mixing, dark matter abundance of the universe and the 3.5 keV line within a BSM (Beyond SM) model, where we have naturally a two component Dark Matter and a nearly degenerate long-lived state. We extend the SM gauge group by an anomaly free local $U(1)_{L_\mu-L_\tau}$ symmetry [60, 63]. We break this gauge symmetry spontaneously by introducing in the model a SM singlet scalar charged under $U(1)_{L_\mu-L_\tau}$. The mass of the resultant neutral gauge boson is given in terms of the new gauge coupling and vacuum expectation value (VEV) of this scalar. Also included in the model are three RH neutrinos and a SM (inert) doublet scalar, both of which carry $-1$ charge with respect to an additional $Z_2$ symmetry, while all other particles carry charge $+1$. This forbids all Yukawa couplings of this
doublet with the SM fermions (thereby earning the name, inert doublet) and the only Yukawa
term where it appears is the one with the RH neutrinos. The $Z_2$ symmetry also forbids
the normal Yukawa coupling involving the lepton doublets, RH neutrinos and the SM Higgs doublet.
On the other hand, the allowed Yukawa coupling between the lepton doublets, RH neutrinos and
the inert doublet does not lead to a Dirac-like mass term since the inert doublet does not take
a VEV. As a result, light neutrino masses via Type-I seesaw is forbidden. However, the light
neutrinos get mass radiatively at one-loop, where the RH neutrinos and the inert doublet run
in the loop [61]. The RH neutrinos protected by the $Z_2$ symmetry become the dark matter of
the universe. The $Z_2$ symmetry allows the RH neutrinos to be coupled only to the Higgs sector
and the $Z_{\mu\tau}$. We invoke a non-thermal production mechanism for the generation of DM in the
early universe via the freeze-in mechanism [11] wherein the RH neutrinos are mainly produced
by out-of-equilibrium decays of $Z_{\mu\tau}$ gauge bosons.

The 3.5 keV $\gamma$ line can be explained by the decay of a heavy RH neutrino to another RH
neutrino if the two states are nearly degenerate and the mass splitting is 3.5 keV [47, 49].
Moreover the lifetime of the next-to-lightest neutrino has to be sufficiently long. Both conditions
are naturally realised in our scenario. Indeed we will see that in the $U(1)_{L_\mu-L_\tau}$ symmetric limit,
the $L_\mu - L_\tau$ symmetry enforces two completely degenerate states and one heavier state for the
RH mass spectrum in our model. The two lighter degenerate RH neutrino states play the role of
a two-component dark matter. The spontaneous breaking of $U(1)_{L_\mu-L_\tau}$ results in a small mass
splitting between the two degenerate RH neutrinos, determined by the symmetry breaking scale
and Yukawa couplings of the RH neutrinos. The lifetime of the heavier state is longer than the
age of the Universe due both to the phase-space suppression and to the small parameters needed
to explain the light neutrino masses.

The rest of the article is organised in the following way. In Section II we describe the model
in detail. In Section III we discuss the effect of $U(1)_{L_\mu-L_\tau}$ and its breaking on the RH neutrino
mass spectrum. In Section IV we present our formalism for the freeze-in production of the RH
neutrino DM. In section we show our results on the DM relic abundance and discuss all aspects
related to it. In Section V and Section VI we will present our DM results and explain the origin
of 3.5 keV line in our model from the RH neutrino decay respectively. We finally conclude in
Section VII.

II. MODEL

The complete gauge group in our model is, $SU(2)_L \times U(1)_Y \times U(1)_{L_\mu-L_\tau}$. In addition to the
SM particles, we augment our model with a SM scalar doublet, a SM scalar singlet and three RH
neutrinos. We also impose a $Z_2$ symmetry to make the additional doublet inert. The $Z_2$ charge
of the RH neutrinos are also taken to be $-1$ to keep them stable, such that they could be dark
matter candidates. The complete fermionic and scalar particle content of the model and their corresponding charges under the different symmetry groups are shown in Tables I and II.

| Gauge Group | Baryon Fields | Lepton Fields | Scalar Fields |
|-------------|---------------|---------------|---------------|
| SU(2)$_L$  | $Q_i^L = (u_i^L, d_i^L)^T$ | $L_i^L = (\nu_i^L, e_i^L)^T$ | $\phi_h, \phi_H, \eta$ |
| $U(1)_Y$   | 2 1 1        | 2 1 1         | 2 1 2         |
| $\mathbb{Z}_2$ | + + +       | + + +         | + + +         |

Table I: Particle contents and their corresponding charges under SM gauge group and discrete group $\mathbb{Z}_2$.

| Gauge Group | Baryonic Fields | Lepton Fields | Scalar Fields |
|-------------|----------------|---------------|---------------|
| $U(1)_{\mu - \tau}$ | $(Q_L^\mu, u_R^\mu, d_R^\mu)$ | $(L_L^\mu, \nu_R^\mu, N_R^\mu)$ | $\phi_h, \phi_H, \eta$ |
|             | 0 0 1         | -1           | 0 1 0         |

Table II: Particle contents and their corresponding charges under $U(1)_{\mu - \tau}$.

The complete Lagrangian $\mathcal{L}$ for the present model is as follows,

$$\mathcal{L} = \mathcal{L}_{SM} + \mathcal{L}_N + (D_\mu \phi_H)\bar{\phi}_H + (D_\mu \eta)\bar{\eta} + \sum_{j=\mu, \tau} Q_j^L L_j^\gamma \gamma_\mu L_j Z_{\mu\tau}^\rho - \frac{1}{4} F_{\mu\sigma \rho\sigma} F_{\mu\sigma} - V(\phi_h, \phi_H, \eta),$$

where $\phi_h$ and $\eta$ are two SU(2)$_L$ doublets while $\phi_H$ is a scalar singlet. Moreover, $Q_j^L = 1(-1)$ for $j = \mu(\tau)$ where $L_j = (\nu_j, j)^T$. Here, one of the scalar doublets namely $\eta$ which is odd under $\mathbb{Z}_2$ symmetry, does not have any Yukawa interaction involving only SM fermions and acts like an inert doublet. For the same symmetry reason it does not have any VEV. The field strength tensor for the extra neutral gauge field $Z_{\mu\tau}$ corresponding to gauge group $U(1)_{\mu - \tau}$ is denoted by $F_{\mu\tau}$. In principle we should include a mixing term between the SM neutral gauge boson $Z$ and the new neutral gauge boson $Z_{\mu\tau}$. The experimental bound restricts this mixing to be $< 10^{-3}$ br the LEP II [65, 66]. In this work we assume no mixing between the neutral gauge bosons of SM and $U(1)_{\mu - \tau}$. Indeed, if such mixing is generated at the loop level, we expect it to be
obtains two physical scalar states \( h \) and \( \eta \). The mass matrix are proportional to the parameter \( h \) and have the following form,

\[
\mathcal{L}_N = \sum_{i=e,\mu,\tau} \frac{i}{2} \bar{N}_i \gamma^\mu D_\mu N_i - \frac{1}{2} M_{ee} \bar{N}_e^c N_e - \frac{1}{2} M_{\mu\tau} (\bar{N}_{\mu}^c N_\tau + \bar{N}_{\tau}^c N_\mu) - \frac{1}{2} h_{e\mu}(\bar{N}_e^c N_\mu + \bar{N}_\mu^c N_e) \phi_H^\dagger - \frac{1}{2} h_{e\tau}(\bar{N}_e^c N_\tau + \bar{N}_\tau^c N_e) \phi_H - \sum_{\alpha=e,\mu,\tau} h_\alpha \bar{N}_\alpha \eta N_\alpha + h.c., \tag{2}
\]

where \( \eta = i\sigma_2 \eta^* \). The potential \( V(\phi_h, \phi_H, \eta) \) in Eq. (1) contains all possible interaction terms involving the two SM scalar doublets and one SM scalar singlet,

\[
V(\phi_h, \phi_H, \eta) = -\mu_H^2 \phi_H^\dagger \phi_H - \mu_h^2 \phi_h^\dagger \phi_h + \mu_\eta^2 \eta^\dagger \eta + \lambda_1 (\phi_h^\dagger \phi_h)^2 + \lambda_2 (\eta^\dagger \eta)^2 + \lambda_3 (\phi_H^\dagger \phi_H)^2 + \lambda_{12} (\phi_h^\dagger \phi_h) (\eta^\dagger \eta) + \lambda_{13} (\phi_h^\dagger \phi_h) (\phi_H^\dagger H H) + \lambda_{25} (\phi_H^\dagger \phi_H) (\eta^\dagger \eta) + \lambda_4 (\phi_h^\dagger \phi_h) (\eta^\dagger \eta) + \lambda_5 (\phi_h^\dagger \phi_h) (\eta^\dagger \eta) + h.c. \tag{3}
\]

After spontaneous breaking of \( U(1)_{L_\mu - L_\tau} \) and \( SU(2)_L \times U(1)_Y \), the scalars take the following form,

\[
\phi_h = \left( \begin{array}{c} 0 \\ v + H \\ \sqrt{2} \end{array} \right), \quad \phi_H = \left( \frac{\nu_{\mu\tau} + H_{\mu\tau}}{\sqrt{2}} \right), \quad \eta = \left( \begin{array}{c} \eta^+ \\ \eta_R^0 + i \eta_I^0 \\ \sqrt{2} \end{array} \right). \tag{4}
\]

There is mixing between the neutral components of \( \phi_h \) and \( \phi_H \), and the off diagonal elements of the mass matrix are proportional to the parameter \( \lambda_{13} \). After diagonalising the mass matrix one obtains two physical scalar states \( h_1 \) and \( h_2 \). Masses of \( h_1 \), \( h_2 \) and mixing angle \( \alpha \) are given by

\[
M_{h_1}^2 = \lambda_1 v^2 + \lambda_3 v_{\mu\tau}^2 - \sqrt{(\lambda_3 v_{\mu\tau}^2 - \lambda_1 v^2)^2 + (\lambda_{13} v v_{\mu\tau})^2}, \tag{5}
\]

\[
M_{h_2}^2 = \lambda_1 v^2 + \lambda_3 v_{\mu\tau}^2 + \sqrt{(\lambda_3 v_{\mu\tau}^2 - \lambda_1 v^2)^2 + (\lambda_{13} v v_{\mu\tau})^2}, \tag{6}
\]

\[
\tan 2\alpha = \frac{\lambda_{13} v v_{\mu\tau}}{\lambda_3 v_{\mu\tau}^2 - \lambda_1 v^2}. \tag{7}
\]

The lighter Higgs state \( h_1 \), for small mixing angle \( \alpha \) and \( v_{\mu\tau} \gg v \), behaves as the Standard Model Higgs observed at the LHC \[67, 68\] and therefore we will take its mass to be 125.5 GeV. From the

\[1\] In this work, to maintain the nonthermal nature of our DM candidates we consider \( g_{\mu\tau} \sim 10^{-11} \) (see Section [V]).
above Eq. (5)-(7), we can also write down the quartic couplings in terms of the physical masses of the Higgses $M_{h_1}$ and $M_{h_2}$ and the mixing angle $\alpha$. The expressions are as follows,

$$\lambda_3 = \frac{M_{h_1}^2 + M_{h_2}^2 + (M_{h_2}^2 - M_{h_1}^2) \cos 2\alpha}{4 v_{\mu\tau}^2},$$

$$\lambda_1 = \frac{M_{h_1}^2 + M_{h_2}^2 + (M_{h_1}^2 - M_{h_2}^2) \cos 2\alpha}{4 v^2},$$

$$\lambda_{13} = \frac{(M_{h_2}^2 - M_{h_1}^2) \cos \alpha \sin \alpha}{v v_{\mu\tau}},$$

(8)

In order to obtained a stable vacuum, the quartic couplings need to satisfy the following inequalities,

$$\lambda_1 \geq 0, \lambda_2 \geq 0, \lambda_3 \geq 0,$$

$$\lambda_{12} \geq -2\sqrt{\lambda_1 \lambda_2},$$

$$\lambda_{13} \geq -2\sqrt{\lambda_1 \lambda_3},$$

$$\lambda_{23} \geq -2\sqrt{\lambda_2 \lambda_3},$$

$$\lambda_{12} + \lambda_4 - |\lambda_5| \geq -2\sqrt{\lambda_1 \lambda_2},$$

$$\sqrt{\lambda_{13}} + 2\sqrt{\lambda_1 \lambda_3} \sqrt{\lambda_{23}} + 2\sqrt{\lambda_2 \lambda_3} \sqrt{\lambda_{12} + \lambda_4 - |\lambda_5|} \sqrt{\lambda_{13} + 2\sqrt{\lambda_1 \lambda_3} \sqrt{\lambda_{23} + 2\sqrt{\lambda_2 \lambda_3} \sqrt{\lambda_{12} + 2\sqrt{\lambda_1 \lambda_2}}}} \geq 0,$$

$$\sqrt{\lambda_{13}} + 2\sqrt{\lambda_1 \lambda_3} \sqrt{\lambda_{23}} + 2\sqrt{\lambda_2 \lambda_3} \sqrt{\lambda_{12} + 2\sqrt{\lambda_1 \lambda_2}} \geq 0 \ .$$

(9)

As we will see in the result section (Section V), in our analysis the value of the extra singlet scalar vev is around $10^{14}$ GeV, mass of BSM Higgs $M_{h_2} = 5$ TeV and the mixing angle between the neutral components of Higgses $\alpha = 0.01$. Hence, we get the following values for the quartic couplings by using the Eq. (8),

$$\lambda_1 = 0.15, \lambda_3 = 1.25 \times 10^{-21} \text{ and } \lambda_{13} = 1.01 \times 10^{-11}. \quad (10)$$

All the values of the quartic couplings as shown above are positive and in the present case the quartic couplings which are related to the inert doublet are free parameters (except $\lambda_5$, which we have considered $\sim 10^{-3}$ to obtain light neutrino masses in sub-eV range), hence all the inequalities as prescribed in Eq. (9) are inevitably satisfied.

On the other hand the masses of the inert doublet components after symmetry breaking can
be expressed in the following form,

\[
\begin{align*}
M_{\eta^\pm}^2 &= \mu_{\eta}^2 + \frac{1}{2}(\lambda_{12}v^2 + \lambda_{23}v_{\mu\tau}^2), \\
M_{\eta_R}^2 &= \mu_{\eta}^2 + \frac{1}{2}\lambda_{23}v_{\mu\tau}^2 + \frac{1}{2}(\lambda_{12} + \lambda_4 + \lambda_5)v^2, \\
M_{\eta_L}^2 &= \mu_{\eta}^2 + \frac{1}{2}\lambda_{23}v_{\mu\tau}^2 + \frac{1}{2}(\lambda_{12} + \lambda_4 - \lambda_5)v^2,
\end{align*}
\]

(11)

The mass term for the extra neutral gauge boson \( Z_{\mu\tau} \) is also generated when \( \phi_H \) acquires a nonzero VEV \( v_{\mu\tau} \) such that

\[
M_{Z_{\mu\tau}} = g_{\mu\tau}v_{\mu\tau},
\]

(12)

where \( g_{\mu\tau} \) is the gauge coupling corresponding to gauge group \( U(1)_{L_{\mu}-L_{\tau}} \). In this model all three RH neutrinos are odd under the \( \mathbb{Z}_2 \) symmetry. However, the mass of \( N_1 \) comes out to be higher than that of \( N_2 \) and \( N_3 \), so that \( N_1 \) can decay to the lighter RH neutrinos. Also, we will see in Section [III] that the masses of \( N_2 \) and \( N_3 \) are nearly degenerate because of the \( L_{\mu}-L_{\tau} \) symmetry, so that both can play the role of dark matter candidate. Furthermore, in Section [IV] we will show that the RH neutrinos can be produced by the freeze-in mechanism in the early Universe, which requires a tiny gauge coupling \( g_{\mu\tau} \sim \mathcal{O}(10^{-11}) \). Thus, in order to have a TeV scale gauge boson \( Z_{\mu\tau} \) we need large \( v_{\mu\tau} \). Therefore, by choosing appropriate values of the relevant model parameters we can make the masses of inert doublet components higher than the reheat temperature of the universe so that their effect on the production of \( N_2 \) and \( N_3 \) can be safely neglected.

III. HEAVY AND LIGHT NEUTRINO MASSES

In this section we will show how the \( U(1)_{L_{\mu}-L_{\tau}} \) symmetry determines the mass spectrum and mixing angles of all the six neutrinos, the three heavy ones as well as the three light ones. The relevant part of the Lagrangian was given in Eq. (2) where the first term gives the kinetic part while the rest give the mass terms and Yukawa terms involving the neutrinos. After \( U(1)_{L_{\mu}-L_{\tau}} \) and electroweak symmetry breaking the mass matrix for the RH neutrinos is given by

\[
\mathcal{M}_R = \begin{pmatrix}
M_{ee} & \frac{v_{\mu\tau}}{\sqrt{2}}h_{e\mu} & \frac{v_{\mu\tau}}{\sqrt{2}}h_{e\tau} \\
\frac{v_{\mu\tau}}{\sqrt{2}}h_{e\mu} & 0 & M_{\mu\tau}e^{i\xi} \\
\frac{v_{\mu\tau}}{\sqrt{2}}h_{e\tau} & M_{\mu\tau}e^{i\xi} & 0
\end{pmatrix},
\]

(13)
where the terms involving the VEV $v_{\mu\tau}$ appear after $U(1)_{L_{\mu}-L_{\tau}}$ breaking. In the limit that $U(1)_{L_{\mu}-L_{\tau}}$ is unbroken, the RH neutrino mass matrix is given by

$$
M_R = \begin{pmatrix}
M_{ee} & 0 & 0 \\
0 & 0 & M_{\mu\tau} e^{i\xi} \\
0 & M_{\mu\tau} e^{i\xi} & 0
\end{pmatrix}.
$$

(14)

Eigenvalues of Eq. (14) are

$$
M_{2/3}' = \pm M_{\mu\tau} e^{i\xi}
$$

$$
M_1' = M_{ee},
$$

(15)

giving very naturally two degenerate RH neutrinos with opposite parity. The $U(1)_{L_{\mu}-L_{\tau}}$ breaking terms in Eq. (13) brings corrections to the RH neutrino mass spectrum, breaking the degeneracy between $N_2$ and $N_3$. The mass splitting between them is given at first order for $M_{ee} \gg M_{\mu\tau}$ by

$$
\Delta M_{23} = \frac{(h_{e\mu} + h_{e\tau})^2 v_{\mu\tau}^2}{2 M_{ee}}.
$$

(16)

Hence, the mass splitting between $N_2$ and $N_3$ depends on the $U(1)_{L_{\mu}-L_{\tau}}$ breaking VEV $v_{\mu\tau}$ and the Yukawa couplings $h_{e\mu}$ and $h_{e\tau}$. In what follows, we will see that $v_{\mu\tau}$ will be determined by the choice of the $Z_{\mu\tau}$ gauge boson. However, the Yukawa couplings $h_{e\mu}$ and $h_{e\tau}$ can be suitably adjusted to yield a mass splitting of 3.5 keV, needed to explain the 3.5 keV X-ray line from $N_2 \to N_3 \gamma$ decay.

Despite having the RH neutrinos in this model, the masses for light neutrinos cannot be generated by the Type-I seesaw mechanism since the normal Yukawa term involving the RH neutrinos, lepton doublets and the standard model Higgs $\phi_h$ is forbidden by the $\mathbb{Z}_2$ symmetry. The other Yukawa term between the RH neutrinos, lepton doublets and inert doublet $\eta$ is allowed, but $\eta$ does not take any VEV. Hence, there is no mass term for the light neutrinos at the tree-level. However, masses for the light neutrinos gets generated radiatively at the one-loop level through the diagram shown in Fig. 1, giving the following mass matrix for the light neutrinos

$$
M_{ij}^\nu = \sum_k \frac{y_{ij} y_{jk} M_k}{16 \pi^2} \left[ \frac{M_{\eta R}^2}{M_{\eta R}^2 - M_k^2} \ln \frac{M_{\eta R}^2}{M_k^2} - \frac{M_{\eta L}^2}{M_{\eta L}^2 - M_k^2} \ln \frac{M_{\eta L}^2}{M_k^2} \right],
$$

(17)

where $M_k$ is the mass of $k^{th}$ RH neutrino while $M_{\eta R}^0, \eta_{\eta R}^0$ is the mass of $\eta^0_{R,I}$. The quantities $y_{ji} = h_j U_{ji}$, where $h_j$ are the Yukawa couplings in the last term of Eq. (2) and $U_{ji}$ are the
elements of the RH neutrino mixing matrix since the flavour basis \((N_{\alpha}, \alpha = 1, 2, 3)\) of the RH neutrinos and their mass basis \((N_{i}, i = 1, 2, 3)\) are related by a unitary transformation, 
\[N_{\alpha} = \sum U_{\alpha i} N_{i}.\]
If we put this relation into the last term of Eq. (2), one can write the Yukawa term involving SM leptons and RH neutrinos in the following way
\[
\mathcal{L}_{N} \supset h_{j} \bar{L}_{j} \tilde{\eta} U_{ji} L_{i} = y_{ji} \bar{\nu}_{j} \tilde{\eta} N_{i}.\] (18)

If we consider the mass square difference between \(\eta_{0}^{R}\) and \(\eta_{0}^{I}\) i.e. \(M_{\eta_{0}^{R}}^{2} - M_{\eta_{0}^{I}}^{2} = \lambda_{5} v^{2} < < M_{0}^{2}\) where
\[M_{0}^{2} = (M_{\eta_{0}^{R}}^{2} + M_{\eta_{0}^{I}}^{2})/2\]
then the above expression reduces to the following form,
\[
M_{ij}^{\nu} = \frac{\lambda_{5} v^{2}}{16\pi^{2}} \sum_{k} y_{ik} y_{jk} \frac{M_{k}}{M_{0}^{2} - M_{k}^{2}} \left[1 - \frac{M_{k}^{2}}{M_{0}^{2} - M_{k}^{2}} \ln \frac{M_{0}^{2}}{M_{k}^{2}}\right].\] (19)

In this work we have considered the masses of inert scalars greater than the reheat temperature of the Universe, i.e. \(M_{\eta_{0}^{R,I}} \sim 10^{6}\) GeV. The masses of RH neutrinos we consider to be around \(\sim 100\) GeV. If we take the parameter \(\lambda_{5} \sim 10^{-3}\) and \(v = 246\) GeV, then to obtain the neutrino masses of the order of \(M_{\nu} \sim 10^{-11}\) GeV, we need \(y_{ji}^{2} \sim 10^{-1}\) which can be easily obtained. The \(U(1)_{L_{\mu} - L_{\tau}}\) breaking ensures that the mixing angle \(\theta_{13}\) is non-zero and \(\theta_{23}\) is non-maximal.

\[\text{IV. PRODUCTION OF DARK MATTER}\]

We consider the non-thermal production of dark matter candidates. Hence, the initial number densities of these particles are assumed to be negligibly small and their interactions with the particles in the thermal bath are also extremely feeble. As mentioned before, the lighter RH neutrino states \(N_{2}\) and \(N_{3}\) are our dark matter candidates, stabilised by the \(Z_{2}\) symmetry. Because of their gauge and \(Z_{2}\) charges they could be produced only through the decay of \(Z_{\mu\tau}\) and \(h_{1}\) \footnote{Since the mass of the SM-like Higgs has to be kept at 125.5 GeV, the decay channel \(h_{1} \rightarrow N_{i}N_{j}\) will be kinematically allowed only for lighter \(N_{i}/N_{j}\) masses.} and \(h_{2}\) bosons. In what follows, we will see that the dominant production channel
for the RH neutrinos is via the decay of $Z_{\mu\tau}$. In order for the total abundances of $N_2$, $N_3$ to match the observed DM relic density at the present epoch, the gauge coupling has to be small $g_{\mu\tau} \lesssim 10^{-11}$. Since all the interactions of $Z_{\mu\tau}$ are proportional to the gauge coupling $g_{\mu\tau}$, the requirement of such a tiny gauge coupling makes the additional neutral gauge boson $Z_{\mu\tau}$ also decoupled from the thermal bath. Therefore, before computing the DM number density we first need to know the distribution function of mother particle $Z_{\mu\tau}$ by solving the relevant Boltzmann equation. The most general form of the Boltzmann equation describing the distribution function of any species can be expressed as,

$$\hat{L}[f] = C[f]$$

(20)

where $\hat{L}$ is the Liouville operator and $f$ is the distribution function which we want to compute while in the RHS the term $C$ contains interaction processes which are responsible for changing the number density of the species under considering. $C$ is known as the collision term. If one considers an isotropic and homogeneous Universe then using the FRW metric, the Liouville operator $^3$ takes the following form,

$$\hat{L} = \frac{\partial}{\partial t} - H p \frac{\partial}{\partial p},$$

(21)

where $p$ is magnitude of three momentum and $H$ is the Hubble parameter. Now, we change the variables $(p, t)$ to a new set of variables $(\xi_p, r)$ using a transformation as mentioned in Ref. $^{19}$

$$r = \frac{M_{sc}}{T}, \quad \xi_p = \left(\frac{g_s(T_0)}{g_s(T)}\right)^{1/3} \frac{p}{T}.$$  

(22)

$M_{sc}$ is some reference mass scale. Using the time-Temperature relationship $\frac{dT}{dt} = -HT \left(1 + \frac{Tg'_s(T)}{3g_s(T)}\right)^{-1}$, the Liouville operator defined in Eq. (21) can be reduced to the following form containing a derivative with respect to a single variable, i.e.

$$\hat{L} = rH \left(1 + \frac{Tg'_s}{3g_s}\right)^{-1} \frac{\partial}{\partial r}$$

(23)

where $g_s(T)$ and $g'_s(T)$ are the effective number of degrees of freedom (d.o.f) related to entropy of the Universe and its derivative with respect to the temperature $T$.

The Boltzmann equation to determine the distribution function $(f_{Z_{\mu\tau}})$ of $Z_{\mu\tau}$ is then given by,

$$\hat{L} f_{Z_{\mu\tau}} = \sum_{i=1,2} C_{h_i \rightarrow Z_{\mu\tau} Z_{\mu\tau}} + C_{Z_{\mu\tau} \rightarrow \text{all}},$$

(24)

$^3$ General form of the Liouville operator is, $\hat{L} = p^\alpha \frac{\partial}{\partial x^\alpha} - \Gamma^\alpha_{\beta\gamma} p^\beta p^\gamma \frac{\partial}{\partial p^\alpha}$ where $p^\alpha$ is the four momentum and $\Gamma^\alpha_{\beta\gamma}$ is the affine connection by which gravitational interaction enters in the equation.
where the first term in the RHS represents the production of $Z_{\mu\tau}$ from the decays of scalars $h_1$ and $h_2$ while the second term describing the depletion of $Z_{\mu\tau}$ due to its all possible decay modes. The expressions of collision terms $C^{h_i \rightarrow Z_{\mu\tau} Z_{\mu\tau}}$ and $C^{Z_{\mu\tau} \rightarrow all}$ are given in Appendix A. Note that generically also scattering processes, which change the $Z_{\mu\tau}$ number, are present, but those give a subleading contribution compared to the decay (see e.g. the Appendix of [51] for a discussion).

Once we numerically evaluate the non thermal momentum distribution of the gauge boson $Z_{\mu\tau}$, we can easily determine the number density of $Z_{\mu\tau}$ using following relation

$$n_{Z_{\mu\tau}}(r) = \frac{g T^3}{2 \pi^2} B(r)^3 \int d\xi \frac{\xi^2}{f_{Z_{\mu\tau}}(\xi, r)} ,$$

(25)

where

$$B(r) = \left( \frac{g_s(T_0)}{g_s(T)} \right)^{1/3} = \left( \frac{g_s(M_{sc}/r)}{g_s(M_{sc}/r_0)} \right)^{1/3} .$$

(26)

Here $T_0$ is the initial temperature and $M_{sc}$ is some reference mass scale. In this work we take $T_0 = 10$ TeV and $M_{sc} = M_{h_1} = 125.5$ GeV, the mass of SM Higgs boson. The entropy density of the Universe is given by [69],

$$s = \frac{2 \pi^2}{45} g_s(T) T^3 .$$

(27)

Therefore, after determining the number density of $Z_{\mu\tau}$ and the entropy of the Universe one can determine the comoving number density using the following relation,

$$Y_{Z_{\mu\tau}} = \frac{n_{Z_{\mu\tau}}}{s} .$$

(28)

Finally, to determine the comoving number densities of DM components $N_2$ and $N_3$, we need to solve the relevant Boltzmann equation for $N_2$ and $N_3$, which can be written in a generic form,

$$\frac{dY_{N_i}}{dr} = V_{ij} M_{pl} r \sqrt{g_s(r)} \left[ \sum_{k=1,2} \sum_{i=1,2,3} \langle \Gamma_{h_k \rightarrow N_j N_i} \rangle (Y_{h_k} - Y_{N_j} Y_{N_i}) \right]$$

$$+ \frac{V_{ij}}{1.66 M_{sc}^2 g_s(r)} \sum_{i=1,2,3} \langle \Gamma_{Z_{\mu\tau} \rightarrow N_j N_i} \rangle NTH \ (Y_{Z_{\mu\tau}} - Y_{N_j} Y_{N_i}) ,$$

(29)

where $M_{pl}$ is the Planck mass while $g_s(r) = \frac{g_s(r)}{\sqrt{g_{\rho}(r)}} \left( 1 - \frac{1}{3} \frac{d \ln g_s(r)}{d \ln r} \right)$ is a function of $g_{\rho}(r)$ and $g_s(r)$. The parameter $V_{ij} = 2$ for $i = j$ and equal to 1 otherwise. The first term in the above equation represents the production of $N_j$ from the decays of scalar fields $h_1$ and $h_2$. Since these scalar fields remain in thermal equilibrium throughout their cosmological evolution, one can consider
their distribution function as Maxwell-Boltzmann distribution. Therefore the thermal averaged decay width for a process \( h(k) \to N_j N_i \) is given by \[70\]

\[
⟨Γ_{h_k→N_j N_i}⟩ = Γ_{h_k→N_j N_i} \frac{K_1 \left( r \frac{M_h}{M_{sc}} \right)}{K_2 \left( r \frac{M_h}{M_{sc}} \right)},
\]

(30)

where \( K_i \) is the Modified Bessel function of \( i^{th} \) kind. As the neutral gauge boson \( Z_{\mu\tau} \) is not in thermal equilibrium (due to very small value of \( g_{\mu\tau} \)), one cannot assume a Maxwell-Boltzmann distribution function for \( Z_{\mu\tau} \). The distribution \( f_{Z_{\mu\tau}} \) of \( Z_{\mu\tau} \) can be found by solving Eq. (24) and we have shown it in Fig. 2. Although the shape of the distribution is similar in both cases but they differ by magnitude because in the current case \( Z_{\mu\tau} \) is always out of equilibrium and never attains equilibrium value. Once we get the distribution function \( f_{Z_{\mu\tau}} \) the non-thermal average of decay width for the process \( Z_{\mu\tau} \to N_j N_i \) can be computed as follows

\[
⟨Γ_{Z_{\mu\tau}→N_j N_i}⟩_{NT} = M_{Z_{\mu\tau}} Γ_{Z_{\mu\tau}→N_j N_i} \frac{\int f_{Z_{\mu\tau}}(p) d^3p}{\int f_{Z_{\mu\tau}}(p) d^3p}.
\]

(31)

All the relevant decay widths of \( h_2 \) and \( Z_{\mu\tau} \) needed in Eq. (29) are given in Appendix A in detail. After solving the above Boltzmann equations for \( j=2 \) and \( j=3 \), we can determine the comoving

Figure 2: Thermal and Non-thermal distribution function of \( Z_{\mu\tau} \) gauge boson.
number density of the DM candidates $N_2$ and $N_3$. Therefore, one can easily determine the total DM relic density for $N_2$ and $N_3$ candidates by using the following relation [71],

$$
\Omega_{DM} h^2 = 2.755 \times 10^8 \left( \frac{M_{N_2}}{\text{GeV}} \right) Y_{N_2}(T_{\text{Now}}) + 2.755 \times 10^8 \left( \frac{M_{N_3}}{\text{GeV}} \right) Y_{N_3}(T_{\text{Now}}) .
$$

(32)

V. RESULTS

Figure 3: Left panel: Variation of relic density with $r$ and contributions from $h_2$ and $Z_{\mu\tau}$ in the DM production. Right panel: Variation of comoving number density of $Z_{\mu\tau}$ and $N_2$, $N_3$ with $r$ for three different values of gauge boson mass. Other parameters have been kept fixed at $g_{\mu\tau} = 1.01 \times 10^{-11}$, mixing angle $\alpha = 0.01$, gauge boson mass $M_{Z_{\mu\tau}} = 1$ TeV, DM mass $M_{DM} = 100$ GeV, BSM Higgs mass $M_{h_2} = 5$ TeV and RH neutrinos masses $M_{N_1} = 150$ GeV and $M_{DM} = M_{N_2} \simeq M_{N_3} = 100$ GeV.

Using Eqs. (29), (30), (31) and (32) we numerically compute the DM abundance. In the left panel of Fig. 3 we show the time evolution of the DM relic density with $r(= M_{h_1}/T)$. The left panel of the this figure shows the comparative contribution for the two DM production channels, $Z_{\mu\tau} \to N_i N_j$ and $h_2 \to N_i N_j$. We have taken masses of the RH neutrinos $N_2$ and $N_3$ as 100 GeV and hence the decay of SM-like Higgs $h_1$ to a pair of RH neutrinos is kinematically forbidden. From the left panel we see that for the large value of BSM Higgs mass $(M_{h_2} \sim 5$ TeV), the DM production at low $r$ (which corresponds to high $T$) is dominated by $h_2$ decay. However, as the temperature of the universe falls and goes below the mass of the $Z_{\mu\tau}$ gauge bosons, they get produced, and for high value of $r$ (which corresponds to comparably lower temperature of the universe), the DM production via the $Z_{\mu\tau}$ decay channel dominates. The reason for this
dominance can be understood as follows. From Eqs. (A3) and (A7) given in Appendix A, we see that the decay width \( \Gamma_{Z_{\mu\tau} \to N_i N_j} \propto M_{Z_{\mu\tau}} \mu_{\tau} \) while \( \Gamma_{h_2 \to N_i N_j} \propto M_{h_2} h_{e\alpha} h_{e\beta} \), where \( h_{e\alpha} h_{e\beta} \) are products of two any of the Yukawa couplings \( h_{e\mu} \) and \( h_{e\tau} \) that appeared in Eq. (2). Since we have chosen \( M_{Z_{\mu\tau}} \sim M_{h_2} \) we can write

\[
\frac{\Gamma_{Z_{\mu\tau} \to N_i N_j}}{\Gamma_{h_2 \to N_i N_j}} \propto \frac{g_{\mu\tau}^2}{h_{e\alpha} h_{e\beta}},
\]

(33)

Since the Yukawa couplings \( h_{e\alpha} \) appear as the \( U(1)_{L_{\mu} - L_{\tau}} \) breaking terms in the RH neutrino mass matrix which instruments the splitting of 3.5 keV between \( N_2 \) and \( N_3 \) we have from Eq. (13)

\[
V_{e\alpha} = \frac{h_{e\alpha} \mu_{\tau}}{\sqrt{2}} \sim 0.1 \text{ GeV}.
\]

(34)

Inserting this in Eq. (33) and using the relation \( M_{Z_{\mu\tau}} = g_{\mu\tau} \mu_{\tau} \) we get

\[
\frac{\Gamma_{Z_{\mu\tau} \to N_i N_j}}{\Gamma_{h_2 \to N_i N_j}} \propto \frac{M_{Z_{\mu\tau}}^2}{V_{e\alpha}^2},
\]

(35)

explaining the dominance of the \( Z_{\mu\tau} \) decay channel.

In the right panel of Fig. 3 we show the variation of the comoving number densities of the \( Z_{\mu\tau} \) gauge boson \( \text{vis-a-vis} \) that of the sum of \( N_2 \) and \( N_3 \). We show this as function of \( r \) for three different values of the gauge boson mass \( M_{Z_{\mu\tau}} \).

The abundance \( Y_{Z_{\mu\tau}} \) (indicated by the dash line) has an initial rise, then flattens and finally decays. One can see from Eq. (24) that there are two collision terms in the Boltzmann Equation, one for \( Z_{\mu\tau} \) production and another one for its decay to all possible channels and they are active at different times. Note that the maximal abundance of \( Z_{\mu\tau} \) can be easily estimated also by the analytic formula for FIMP production, i.e. for \( M_{Z_{\mu\tau}} \ll M_{h_2} \)

\[
\Omega^{FI} h^2 = 1.09 \times 10^{27} \frac{g}{g_*^{3/2}} \frac{M_{Z_{\mu\tau}}}{M_{h_2}^2} \Gamma_{h_2 \to Z_{\mu\tau} Z_{\mu\tau}} \sim 2.18 \times 10^{24} \frac{g_{\mu\tau}^2 M_{h_2}}{32\pi M_{Z_{\mu\tau}}} = 8.54;\]

(36)

where \( g \) counts the number of internal degrees of freedom of the mother particle. According to eq. (32) this corresponds to \( Y_{Z_{\mu\tau}} = 0.3 \times 10^{-10} \) and is in perfect agreement with the plateau in Fig. 3. One interesting point to note is that as we increase the \( Z_{\mu\tau} \) mass \( M_{Z_{\mu\tau}} \), keeping \( g_{\mu\tau} \) fixed, the DM abundance decreases instead of increasing, as explained by the relation above. In the same figure also the production of dark matter as a result of the out-of-equilibrium decay of \( Z_{\mu\tau} \) can be seen beautifully. Less production of \( Z_{\mu\tau} \) results in lower DM abundance, since practically every \( Z_{\mu\tau} \) produces two Dark Matter particles.

The left panel of Fig. 4 shows the variation of relic density with the parameter \( r \) for different initial temperature \( T_{ini} \) (temperature where DM relic density is taken as zero). Important point
Figure 4: Left (Right) panel: Variation of relic density with $r$ for different initial temperature (for different gauge coupling values), while the other parameters have been kept fixed at $g_{\mu\tau} = 1.01 \times 10^{-11}$, ($T_{ini} = 10$ TeV), mixing angle $\alpha = 0.01$, gauge boson mass $M_{Z_{\mu\tau}} = 1$ TeV, BSM Higgs mass $M_{h_2} = 5$ TeV and RH neutrinos masses $M_{N_1} = 150$ GeV, $M_{N_2} \simeq M_{N_3} = 100$ GeV.

to note here that as long as the initial temperature is above the mass of the gauge boson $Z_{\mu\tau}$, final relic density remains the same. However, when we reduce the initial temperature below the $Z_{\mu\tau}$ mass (shown by the cyan color curve) then final abundance reduces significantly due to the Boltzmann suppression factor. In the right panel we show the variation of DM relic density with $r$ for different gauge coupling values ($g_{\mu\tau}$). One can see from the figure that if we increase the value of the gauge coupling, the DM production rate as well as the total DM abundance increases. The reason can be understood from Eq. (A3) which shows that the DM production rate, which is almost the same as the $Z_{\mu\tau}$ decay rate, is proportional to the second power of $g_{\mu\tau}$.

In the present model for $g_{\mu\tau} = 1.01 \times 10^{-11}$ we achieve the correct DM relic density value of the universe. In both the panels of Fig. 4, the horizontal magenta line corresponds to the present day correct DM relic density value of the universe. For the rest of the analysis, we have fixed the initial temperature of the universe at 10 TeV.

In the left panel of Fig. 5, we present the variation of the DM relic density for three different values of the DM mass $M_{DM} (=M_{N_2}, M_{N_3})$. As shown in Eq. (32) that DM relic density is proportional to the DM mass $M_{N_2}$ and $M_{N_3}$ and this dependence is evident in the left panel of Fig. 5. For the chosen value of the parameters (mentioned in the caption), we have obtained correct relic density value (indicated by the horizontal line) of the universe for DM mass value $M_{DM} = M_{N_2} \simeq M_{N_3} = 100$ GeV, this value will be different for different set of values of the other parameters. In the right panel of Fig. 5, we show the decay contributions of $Z_{\mu\tau}$ in different
channels. The relative contributions among the different channels is seen to differ significantly and the decay rate into $N_2 N_3$ dominates naturally producing equal populations of the two Dark Matter candidates. Indeed, to produce degenerate neutrinos i.e. $M_{N_2} \simeq M_{N_3}$, we have considered relatively small values of $\frac{h_{e\mu} v_{e\mu}}{\sqrt{2}}$ and $\frac{h_{e\tau} v_{e\tau}}{\sqrt{2}}$ ($\sim 0.1$), as discussed before. Therefore, the elements of the unitary matrix which relate the flavour and mass basis of the RH neutrinos take the following form, $U_{11} \sim 1$, $U_{12}, U_{13}, U_{21}, U_{31} \sim 0.01$, $U_{22} = U_{23} = \frac{1}{\sqrt{2}}$ and $U_{32} = -U_{33} = -\frac{1}{\sqrt{2}}$. Therefore, it is clear from the couplings (as listed in Eq. (A4)) that the dominant channel for DM production is $Z_{\mu\tau} \rightarrow N_2 N_3$, while the other channels will be suppressed which is clearly visible in the right panel of Fig. 5. Similar considerations will also be true for the $N_3$ DM production channels.

VI. 3.5 KEV $\gamma$ RAY LINE

Finally, we come to the explanation of the 3.5 keV $\gamma$-ray line from the RH neutrino radiative decay $N_2 \rightarrow N_3 \gamma$. Since the photon flux for a decaying Dark Matter candidate is given by

$$\Phi = \frac{1}{4\pi M_{N_2} \tau_{N_2}} \int_{l.o.s.} \rho_{N_2}(\vec{r}) d\vec{r}$$

(37)

where the last integral over the $N_2$ density is computed along the line of sight and $\tau_{N_2}$ is the lifetime of the heavier DM particle $N_2$. In order to explain the 3.5 keV line from a decay such
as $N_2 \rightarrow N_3 \gamma$, we need not only a mass splitting between the two fermion states of $\sim 3.5$ keV, but also a decay width of the unstable DM given as,

$$\Gamma(N_2 \rightarrow N_3 \gamma) = (0.72 - 6.6) \times 10^{-52} \text{GeV} \left( \frac{M_{N_2}}{3.5 \text{keV}} \right) = (0.2 - 1.9) \times 10^{-44} \text{GeV} \left( \frac{M_{N_2}}{100 \text{GeV}} \right).$$

(38)

Here we are assuming that the density of $N_2$ is approximately half of the DM density and rescaled the result of [24] accordingly.

The relevant decay diagrams for $N_2$ are shown in Fig. 6. We consider $N_2$ to be slightly heavier than $N_3$ ($\sim 3.5$ keV) so that it can produce the 3.5 keV $\gamma$-ray line. As discussed before, the 3.5 keV mass splitting between nearly-degenerate $N_2$ and $N_3$ can be easily achieved in our model via the $U(1)_{L_\mu - L_\tau}$ symmetry and its breaking parameters. So we take $V_{e\alpha} = \frac{h_{e\alpha} v_{\mu\tau}}{\sqrt{2}} \sim 0.1$ GeV ($\alpha = \mu, \tau$) and by suitably adjusting the $V_{e\alpha}$ parameters we can generate the 3.5 keV mass gap between $N_2$ and $N_3$. For the $U(1)_{L_\mu - L_\tau}$ conserving leading terms in Eq. (13) we take the values $M_{ee} = 11$ TeV and $M_{\mu\tau} = 100$ GeV which gives us $M_{N_2}$ and $M_{N_3} \sim 100$ GeV with opposite CP parities [59]. Ref. [59] has pointed out that if $N_2$ and $N_3$ have opposite CP, then the transition from $N_2$ to $N_3$ is governed only by the magnetic moment term ($\mu_{23}$), generated at one loop level as shown in Figure 6. Therefore, the effective Lagrangian for the decay process $N_2 \rightarrow N_3 \gamma$ is given as

$$\mathcal{L}_{\text{eff}} \approx i \frac{\mu_{23}}{2} \bar{N}_2 \sigma^{\mu\nu} N_3 F_{\mu\nu}.$$ 

(39)

In determining the expression for the above decay process we consider the ratio of lepton mass to RH neutrino mass to be very small ($\frac{M_l}{M_{N_2}} \ll 1$). Also, the ratio of the RH neutrino mass and the inert doublet mass is very small i.e. $\frac{M_{N_2}}{M_\eta} \ll 1$. The decay width of $N_2$ comes out as [72],

$$\Gamma(N_2 \rightarrow N_3 \gamma) = \frac{\mu_{23}^2}{4\pi} \delta^3 \left( 1 - P \frac{M_{N_3}}{M_{N_2}} \right)^2,$$

(40)
\[ \delta = \frac{M_{N_2}^2}{2} \left(1 - \frac{M_{N_2}^2}{M_{N_2}^2} \right), \]

\[ P \] gives the relative CP of the two neutrino states, which in the present model is \( P = -1 \). The magnetic moment coefficient \( \mu_{23} \) in our model is given by

\[ \mu_{23} = \sum_i \frac{e}{2} \frac{1}{(4\pi)^2} \frac{M_{N_2}}{M_{\eta}^2} (y_{i2}^2 y_{i3}^2), \]  

(41)

where \( y_{ij} = h_i U_{ij} \) being the derived Yukawa couplings given in Eq. (18). The values of the parameters appearing in the \( N_2 \) decay width are intimately related with those that determine the light neutrino masses. In Section III, we had set the parameter values to explain the tiny neutrino mass in the following order,

\[ M_{\eta} = 10^6 \text{ GeV}, M_{N_2} = 100 \text{ GeV}, (y_{ij})^2 = 10^{-1}. \]  

(42)

Using these in the Eq. (18) we get \( \mu_{23} \sim \mathcal{O}(10^{-14}) \text{ GeV}^{-1} \). Using Eq. (40), for DM mass around 100 GeV, \( \delta \simeq 3.5 \text{ keV} \) and \( \mu_{23} \sim 10^{-14} \text{ GeV}^{-1} \), we get the lifetime of \( N_2 \) of the order \( \mathcal{O}(10^{-44}) \) GeV, which is exactly what is needed to give the 3.5 keV line. Note that the lifetime of \( N_2 \) is then around \( 10^{19} \text{ sec} \) and hence greater than the age of the universe (\( 10^{17} \text{ sec} \)). Hence the present model can naturally explain the origin of the claimed 3.5 keV line.

VII. CONCLUSION

In the present work we extended the SM gauge group by a local \( U(1)_{L_\mu - L_\tau} \) gauge group and a \( Z_2 \) discrete symmetry. The particles spectrum was extended by three RH neutrinos, one inert doublet and one SM gauge singlet scalar. We showed that this model explains the observed 3.5 keV line consistently with the relic dark matter abundance in the framework of a model that generates light neutrino masses radiatively. The Type I seesaw in this model is forbidden by the \( Z_2 \) symmetry but tiny neutrino masses are generated via a one-loop diagram involving the RH neutrino and the inert doublet which does not take any VEV. We considered inert scalar masses \( \sim 10^6 \text{ GeV} \), which is higher than the reheat temperature, and RH neutrino masses \( \sim 100 \text{ GeV} \). Then for parameter choices \( \lambda_5 \sim 10^{-3} \) and Yukawa couplings \( y_{ji}^2 \sim 10^{-1} \) we can get light neutrino masses \( M_\nu \sim 0.01 \text{ eV} \). The RH neutrino mass matrix in our model is non-diagonal and carries the \( L_\mu - L_\tau \) flavour structure which ensures that two of the RH neutrino remain degenerate in the \( U(1)_{L_\mu - L_\tau} \) symmetric limit. The spontaneous breaking of the \( U(1)_{L_\mu - L_\tau} \) gauge symmetry generates terms in the RH neutrino mass matrix that splits the two degenerate RH neutrinos by 3.5 keV, while the third one remains heavier. The two nearly degenerate neutrinos form the two-component DM in our model. We showed that the RH neutrinos are predominately produced by the decay of the extra neutral gauge boson \( Z_{\mu\tau} \), which are taken in the 1 TeV mass range in our model. The production of RH neutrinos from decay of the additional scalar \( h_2 \) is subdominant,
while the annihilation channels have negligible effect. We showed that the peculiar structure of the unitary matrix \( U \) which relates the flavour and mass basis of the RH neutrinos ensures that the decay mode \( Z_{\mu\tau} \to N_2 N_3 \) is the dominant one among the other channels. Since the associated gauge coupling \( g_{\mu\tau} \) is taken to be very small here, the \( Z_{\mu\tau} \) stays out of equilibrium in the early universe and the RH neutrinos are produced by the freeze-in mechanism. We solved the coupled Boltzmann equation numerically and showed the dependence of the DM relic abundance on initial temperature \( T_{\text{ini}}, g_{\mu\tau}, M_{Z_{\mu\tau}}, M_{\text{DM}} \). Finally, we showed that the heavier of the two DM component \( N_2 \) can decay into the lighter \( N_3 \) \( (N_2 \to N_3 \gamma) \) through one loop diagram, thus producing the 3.5 keV X-ray line that was observed by Chandra satellite. The model parameter values which determine the lifetime of \( N_2 \) were obtained through constraints from the light neutrino mass sector and gave a decay rate of \( 10^{-44} \) GeV for \( N_2 \). So the lifetime of the heavier Dark Matter particle is consistent with both the age of the universe as well as the strength of the observed 3.5 keV line.

Regarding collider observables, this model unfortunately does not give many promising signatures. Indeed all the particles of the gauged \( \mu - \tau \) sector interact with the Standard Model only via the very small coupling \( g_{\mu\tau} \sim 10^{-11} \), so that their production at LHC or their effect on precision observables is very suppressed. If one would be able to produce those states, a long lifetime and possibly displaced vertices could be the characteristic signature \([49, 51, 73]\). On the other hand, more substantial can be the production cross-section for the heavier Higgs boson \( h_2 \), depending on its mass the mixing angle \( \alpha \). Unfortunately in this case, its dominant decay channels are those in Standard Model states through the mixing with the Higgs doublet and so the connection of this heavy state with the neutrino sector and the \( U(1)_{L_\mu-L_\tau} \) will be difficult to prove.

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Appendix

Appendix A: Analytical Expression of relevant Decay width and Collision terms

If we consider a generic process \( \chi(\tilde{p}) \to a(\tilde{p}_1) b(\tilde{p}_2) \) (where \( \tilde{p} = (E_p, \bar{p}) \)) then the collision term will take the following form \[69, 70\],

\[
C[f_\chi(p)] = \frac{1}{2 E_p} \int \frac{g_a d^3 p_1}{(2\pi)^3 2 E_{p_1}} \frac{g_b d^3 p_2}{(2\pi)^3 2 E_{p_2}} (2\pi)^4 \delta^4(\tilde{p} - \tilde{p}_1 - \tilde{p}_2) \times |M|^2 \times [f_a f_b (1 \pm f_\chi) - f_\chi (1 \pm f_a) (1 \pm f_b)]. \tag{A1}
\]

Now the full expressions of the collision terms in Eq. (24) are as follows \[19, 21\],

- \( C_{Z_{\mu\tau} \to all} \): Collision term for the extra gauge boson \( Z_{\mu\tau} \) decay can be written in the following way in terms of the parameters which we have introduced in Section IV.

\[
C_{Z_{\mu\tau} \to all} = -f_{Z_{\mu\tau}}(\xi_p) \times \Gamma_{Z_{\mu\tau} \to all} \times \frac{r_{Z_{\mu\tau}}}{\sqrt{\xi_p B(r)^2 + r_{Z_{\mu\tau}}^2}}. \tag{A2}
\]

where \( \Gamma_{Z_{\mu\tau} \to all} = \Gamma_{Z_{\mu\tau} \to f \bar{f}} + \Gamma_{Z_{\mu\tau} \to N_i N_j} \) and the expression for the each decay terms are as follows,

\[
\Gamma_{Z_{\mu\tau} \to f \bar{f}} = \frac{M_{Z_{\mu\tau}} g_{\mu\tau}^2}{12 \pi} \left( 1 + \frac{2M_f^2}{M_{Z_{\mu\tau}}^2} \right) \sqrt{1 - \frac{4M_f^2}{M_{Z_{\mu\tau}}^2}}
\]

\[
\Gamma_{Z_{\mu\tau} \to N_i N_j} = \frac{M_{Z_{\mu\tau}} g_{\mu\tau}^2}{12 \pi S_{ij}} \left( 1 - \frac{(M_{N_i} + M_{N_j})^2}{M_{Z_{\mu\tau}}^2} \right)^{3/2} \times \left( 1 - \frac{(M_{N_i} - M_{N_j})^2}{2 M_{Z_{\mu\tau}}^2} \right)^{1/2} \times \left( 1 - \frac{(M_{N_i} - M_{N_j})^2}{2 M_{Z_{\mu\tau}}^2} \right) \tag{A3}
\]

where \( f = \nu_\mu, \nu_\tau, \mu^\pm \) and \( \tau^\pm \) because of the \( (L_\mu - L_\tau) \) symmetry of the present model and the couplings take the following form depending on RH neutrinos,

\[
g_{Z_{\mu\tau} N_2 N_2} = -\frac{g_{\mu\tau}}{2} (U_{22}^2 - U_{32}^2)
\]

\[
g_{Z_{\mu\tau} N_3 N_3} = -\frac{g_{\mu\tau}}{2} (U_{23}^2 - U_{33}^2)
\]

\[
g_{Z_{\mu\tau} N_1 N_2} = -\frac{g_{\mu\tau}}{2} (U_{21} U_{22} - U_{31} U_{32})
\]

\[
g_{Z_{\mu\tau} N_1 N_3} = -\frac{g_{\mu\tau}}{2} (U_{21} U_{23} - U_{31} U_{33})
\]

\[
g_{Z_{\mu\tau} N_2 N_3} = -\frac{g_{\mu\tau}}{2} (U_{22} U_{23} - U_{32} U_{33}) \tag{A4}
\]

The statistical factor \( S_{ij} = 2 \) for \( i = j \) and \( 1 \) for \( i \neq j \).
• $C_{h_2 \rightarrow Z_{\mu \tau} Z_{\mu \tau}}$: The collision term for the extra gauge boson production $Z_{\mu \tau}$ from the decay of BSM Higgs $h_2$ takes the following form,

$$C_{h_2 \rightarrow Z_{\mu \tau} Z_{\mu \tau}} = \frac{r}{8\pi M_{sc}} \frac{B^{-1}(r)}{\xi_p \sqrt{\xi_p^2 B(r)^2 + \left( \frac{M_{Z_{\mu \tau}} r}{M_{sc}} \right)^2}} \frac{g_{h_2 Z_{\mu \tau} Z_{\mu \tau}}^2}{6} \left( 2 + \frac{(M_{h_2}^2 - 2M_{Z_{\mu \tau}}^2)^2}{4M_{Z_{\mu \tau}}^4} \right) \left( e^{-\sqrt{\left( \xi_{k_{\min}} \right)^2 B(r)^2 + \left( \frac{M_{h_2} r}{M_{sc}} \right)^2}} - e^{-\sqrt{\left( \xi_{k_{\max}} \right)^2 B(r)^2 + \left( \frac{M_{h_2} r}{M_{sc}} \right)^2}} \right).$$

(A5)

where

$$g_{h_2 Z_{\mu \tau} Z_{\mu \tau}} = \frac{2M_{Z_{\mu \tau}}^2 \cos \alpha}{v_{\nu}} ,$$

$$\xi_{k_{\min}}(\xi_p, r) = \frac{M_{sc}}{2B(r) r M_{Z_{\mu \tau}}} \left| \eta(\xi_p, r) - \frac{B(r) \times M_{h_2}}{M_{Z_{\mu \tau}} \times M_{sc}} \right| ,$$

$$\xi_{k_{\max}}(\xi_p, r) = \frac{M_{sc}}{2B(r) r M_{Z_{\mu \tau}}} \left( \eta(\xi_p, r) + \frac{B(r) \times M_{h_2}}{M_{Z_{\mu \tau}} \times M_{sc}} \right) ,$$

$$\eta(\xi_p, r) = \left( \frac{M_{h_2}}{M_{sc}} \right) \sqrt{\frac{M_{h_2}^2}{M_{Z_{\mu \tau}}^2} - 4 \sqrt{\xi_p^2 B(r)^2 + \left( \frac{M_{Z_{\mu \tau}} r}{M_{sc}} \right)^2}} .$$

(A6)

• $\Gamma_{h_k \rightarrow N_i N_j}$: Decay width for the SM like Higgs ($h_1$) and BSM Higgs ($h_2$) take the following form,

$$\Gamma_{h_k \rightarrow N_i N_j} = \frac{M_{h_k} g_{h_k N_i N_j}^2}{8 \pi S_{ij}} \left( 1 - \frac{(M_{N_i} + M_{N_j})^2}{M_{h_k}^2} \right)^{3/2} \times \left( 1 - \frac{(M_{N_i} - M_{N_j})^2}{M_{h_k}^2} \right)^{1/2} .$$

(A7)
where the couplings take the following form,

\[
\begin{align*}
    g_{h_2(1)N_1N_2} &= -\frac{\sqrt{2} \cos \alpha (\sin \alpha)}{4} (U_{11}U_{22}h_{\mu} + U_{12}U_{21}h_{\mu} + U_{11}U_{32}h_{\tau} + U_{12}U_{31}h_{\tau}) \\
    g_{h_2(1)N_1N_3} &= -\frac{\sqrt{2} \cos \alpha (\sin \alpha)}{4} (U_{11}U_{23}h_{\mu} + U_{13}U_{21}h_{\mu} + U_{11}U_{33}h_{\tau} + U_{13}U_{31}h_{\tau}) \\
    g_{h_2(1)N_2N_3} &= -\frac{\sqrt{2} \cos \alpha (\sin \alpha)}{4} (U_{12}U_{23}h_{\mu} + U_{13}U_{22}h_{\mu} + U_{12}U_{33}h_{\tau} + U_{13}U_{32}h_{\tau}) \\
    g_{h_2(1)N_2N_2} &= -\frac{\sqrt{2} \cos \alpha (\sin \alpha)}{2} (U_{12}U_{22}h_{\mu} + U_{12}U_{32}h_{\tau}) \\
    g_{h_2(1)N_3N_3} &= -\frac{\sqrt{2} \cos \alpha (\sin \alpha)}{2} (U_{13}U_{23}h_{\mu} + U_{13}U_{33}h_{\tau})
\end{align*}
\]
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