Left Right Symmetric Models with a Mixture of keV-TeV Dark Matter

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Abstract

We discuss the possibility of realising a multi-component dark matter scenario with widely separated dark matter masses: one having keV scale mass and the other with GeV-TeV scale mass, within the framework of left right symmetric models. To reduce the level of fine tuning, we consider a version of left right model with universal seesaw where keV-GeV scale right handed neutrinos having tiny mixing with light neutrinos can be naturally realised. Due to gauge interactions, both the dark matter candidates are produced thermally in the early Universe but overproducing the keV mass candidate. We consider one of the right handed neutrinos to be decaying at late epochs, just before the big bang nucleosynthesis, in order to dilute the thermally overproduced keV dark matter. We constrain the parameter space from the requirement of producing sub-dominant keV-TeV dark matter, satisfying indirect detection constraints from gamma ray searches and producing the tantalising 3.55 keV monochromatic X-ray line, reported by several groups to be present in galaxy and galaxy cluster data, from the decay of a 7.1 keV dark matter on cosmological scales. We find that these requirements constrain the right sector gauge boson masses to be much heavier than the present collider sensitivity, but can have interesting signatures at experiments like the IceCube.

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I. INTRODUCTION

In the last few decades, there have been significant amount of hints and evidences suggesting the presence of non-luminous and non-baryonic form of matter (popularly known as dark matter (DM)) in the present Universe. Starting from the galaxy cluster observations by Fritz Zwicky [1] back in 1933, observations of galaxy rotation curves in 1970’s [2], the more recent observation of the bullet cluster [3] to the latest cosmology data provided by the Planck satellite [4], the astrophysics, cosmology as well as the particle physics community have a come a long way. The latest data from the Planck mission suggest that around 26% of the present Universe’s energy density is in the form of dark matter. In terms of density parameter and \( h = \frac{\text{Hubble Parameter}}{100 \text{ km s}^{-1} \text{Mpc}^{-1}} \), the present dark matter abundance is conventionally reported as [4]

\[
\Omega_{\text{DM}}h^2 = 0.1198 \pm 0.0015. \tag{1}
\]

In spite of these convincing evidences from astrophysics and cosmology confirming the presence of dark matter in the Universe, we do not yet know the particle nature of DM. In fact, none of these observations tell us anything about the particle nature of DM. However, there exists some requirements for a particle to be a DM candidate, a comprehensive list of which can be found in [5]. These criteria in fact, rules out all the particles in the standard model (SM) of particle physics from being DM candidates which has lead to a plethora of DM models within several beyond standard model (BSM) frameworks. Although the SM neutrinos satisfy some of these criteria, yet they remain relativistic at the epoch of freeze-out as well as matter radiation equality, giving rise to Hot Dark Matter (HDM) which is ruled out by both astrophysics and cosmology observations. Among different BSM frameworks for viable DM candidates, the most popular or the most widely studied scenario perhaps, is the so called weakly interacting massive particle (WIMP) paradigm. In this framework, a dark matter candidate typically with electroweak scale mass and interaction rate similar to electroweak interactions can give rise to the correct dark matter relic abundance, a remarkable coincidence often referred to as the \textit{WIMP Miracle}. Such interactions kept the WIMP DM in thermal equilibrium in the early Universe and eventually its number density gets frozen out when the rate of expansion of the Universe takes over the interaction rates. Such DM candidates typically remain non-relativistic at the epoch of freeze-out as well as matter radiation equality and belong to the category of Cold Dark Matter (CDM).
The sizeable interactions of WIMP DM with other SM particles can not only generate its relic abundance after thermal freeze-out naturally, but also enhances the testability of it as such a DM particle can scatter off nuclei kept in a typical detector. However, till date no such DM-nucleon scattering has been observed in any of the experiments. The most recent dark matter direct detection experiments like LUX, PandaX-II and Xenon1T have also reported their null results [6–9]. Similar null results have been also reported by other direct search experiments like the large hadron collider (LHC) giving upper limits on DM interactions with the SM particles. A recent summary of collider searches for DM can be found in [10]. Although such null results could indicate a very constrained region of WIMP parameter space, they have also motivated the particle physics community to look for beyond the thermal WIMP paradigm. One interesting scenario is the kind of DM which remains mildly relativistic at the epoch of matter radiation equality, keeping it at intermediate stage between HDM and CDM and referred to as Warm Dark Matter (WDM). They typically have mass in the keV range, in contrast to HDM with sub-eV mass and CDM with GeV-TeV scale mass. For a recent review on WDM, one can refer to [11]. Such a scenario is particularly interesting as it can address several challenges like the missing satellite problem, too big to fail problem related to small scale structure formation, that arise in a CDM framework. For a recent review on these small scale challenges, please refer to [12]. The classification of Hot, Warm and Cold DM is primarily done on the basis of their free streaming lengths which is roughly the distance for which the DM particles can freely propagate. For detailed calculation of free streaming lengths, please refer to [13, 14]. Typically, the free streaming length $\lambda_{FS} = 0.1$ Mpc, about the size of a dwarf galaxy, acts as a boundary line between HDM ($\lambda_{FS} > 0.1$ Mpc) and WDM ($\lambda_{FS} < 0.1$ Mpc). For CDM, on the other hand, the free streaming lengths are considerably smaller than this value. Therefore, CDM structures keep forming till scales as small as the solar system which gives rise to disagreement with observations at small scales [12]. HDM, on the other hand, erases all small scale structure due to its large free streaming length, disfavouring the bottom up approach of structure formation. WDM can therefore act as a balance between the already ruled out HDM possibility and the CDM paradigm having issues with small scale structures.

Apart from these motivations, there are motivations from indirect detection experiments as well. There have been many efforts to look for indirect dark matter signatures at different experiments with the hope that even though dark matter may not scatter off nuclei signifi-
cantly as indicated by the null results at direct detection experiments, but they may decay or
annihilate into the standard model particles on cosmological scales and leave some indirect
signatures. Interestingly, there have been some recent observations at some of these indirect
detection experiments which could have possible dark matter origins. One promising indirect
signature of dark matter was reported by two independent analysis [15] and [16] of the data
collected by the XMM-Newton X-ray telescope. Their analysis hinted towards the existence
of a monochromatic X-ray line with energy 3.55 keV in the spectrum of 73 galaxy clusters.
The analysis [15] also claimed the presence of the same line in the Chandra observations of
the Perseus cluster. Later on, the same line was also found in the Milky Way by analysing
the XMM-Newton data [17]. Although the analysis of the preliminary data collected by
the Hitomi satellite (before its unfortunate crash) do not confirm such a monochromatic
line [18], one still needs to wait for a more sensitive observation with future experiments to
have a final word on it. Interestingly, the authors of [19] considered a specific dark matter
model to show consistency among Hitomi, XMM-Newton and Chandra observations. More
recently, the authors of [20] have reported a $3\sigma$ detection of a 3.55 keV emission line in the
spectrum of the Cosmic X-ray background using Chandra observations towards the COS-
MOS Legacy and CDFS survey fields. Such a signal, if confirmed in future experiments,
can be naturally explained by a keV scale sterile neutrino WDM that has mixing with the
SM neutrinos of the order $\approx 10^{-11} - 10^{-10}$ [15, 16]. Different possible keV DM scenarios
that can give rise to such an X-ray line were also discussed, for example [21]. One can also
generate such a signal in typical WIMP DM models if there are two quasi-degenerate DM
candidates having mass splitting of 3.55 keV, allowing the heavier one to decay into the
lighter one and a photon. One such work can be found in [22]. The alternative possibility of
keV dark matter annihilation into monochromatic photons was also discussed very recently
by the authors of [23]. Although such keV scale WDM can not be detected in typical direct
detection experiments like LUX, PandaX-II and Xenon1T, there have been some interesting
proposals for direct detection of such light DM candidates. For example, one may refer to
this recent article on direct detection prospects of sub-MeV DM [24].

Instead of completely giving up on the CDM framework due to the negative results at
dark matter direct detection as well as collider experiments, here we consider an exotic
scenario where the dark sector consists of both cold and warm components. Such a mixed
dark matter model can be very interesting from astrophysical structure point of view, as it
may provide a way to solve the small scale structure problem [25]. There have been several proposals for multi-component WIMP dark matter in the last few years, some of which can be found in [26–32]. Such multi-component DM scenarios even if both the DM candidates are of the same type, can have very interesting signatures at direct as well as indirect detection experiments, as have been discussed extensively in several works [22, 33–46]. However, there have not been many works regarding mixed DM scenarios with different thermal histories.

In [47], the authors considered a mixed DM scenario where one candidate is of WIMP type whereas the other has a non-thermal origin due to its feeble interactions. Such a scenario has more optimistic detection prospects as it can be probed at both types of experiments: sensitive to sub-MeV as well as electroweak scale DM. From model building perspective, it may however be challenging to come up with realistic models that can naturally account for such multi-component DM scenario. Since, both the DM components should be long lived or stable on cosmological scales, one may require exotic or non-minimal symmetries to guarantee that. For example a discrete unbroken $Z_2 \times Z_2$ symmetry can stabilise two DM components. Apart from the stability issue, another important aspect such models should have is a consistent production mechanism of DM. The CDM component, if belongs to the WIMP type DM, can be thermally produced in the early Universe. On the other hand, the production mechanism of WDM depends on the particular realisation and can have either thermal or non-thermal origin. For a summary of these production mechanisms, one can refer to this recent review article [11]. Here we consider a particle physics framework which naturally takes care of both the stability and production issue of the DM components. This is based on the left right symmetric model (LRSM) framework where the SM gauge symmetry is extended to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ which has been studied very extensively in the last few decades. Apart from the usual motivations for LRSM, here we have more motivations from DM point of view. Such an enlarged gauge symmetry can not only guarantee the stability of DM but also can ensure their productions in the early Universe by virtue of their gauge interactions. We consider CDM belonging to both left and right sectors of LRSM but keep the WDM part to the right sector only. This choice is particularly made in order to avoid severe electroweak precision constraints on introducing new keV scale particles having electroweak gauge interactions. We check how the requirement of generating a specific percentage of dark matter in terms of WDM affects the CDM parameter space and vice versa. This gives rise to completely new region of parameter space compared to single
component DM. We also discuss how such a scenario can have interesting indirect detection prospects in both gamma-ray and X-ray experiments.

This article is organised as follows. In section II, we briefly discuss the minimal LRSM and then briefly discuss the possibility of minimal dark matter in LRSM in section III. In section IV we discuss LRSM with universal seesaw for all fermions. In section V, we discuss the possibilities of mixed dark matter in LRSM with universal seesaw. We then discuss in section VI, the relic abundance calculation of both cold and warm dark matter in these models. We discuss our results in section VII and finally conclude in section VIII.

II. MINIMAL LEFT-RIGHT SYMMETRIC MODEL (MLRSM)

Left-Right Symmetric Model [48–53] is one of the well studied and well motivated BSM frameworks where the gauge symmetry of the electroweak theory is extended to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The right handed fermions are doublets under $SU(2)_R$ similar to the way left handed fermions transform as doublets under $SU(2)_L$, in order to treat them on equal footing. The requirement of an anomaly free $U(1)_{B-L}$ makes the presence of three right handed neutrinos a necessity rather than a choice. This is in contrast with the type I seesaw models where three right handed singlet neutrinos are added by hand in order to generate light neutrino masses through seesaw mechanism. In MLRSM, to allow Dirac Yukawa couplings between $SU(2)_{L,R}$ doublet fermions, the Higgs field has to transform as a bidoublet under $SU(2)_{L,R}$ gauge symmetry. In order to break the gauge symmetry of the model to that of the SM spontaneously, scalar triplet fields with non-zero $U(1)_{B-L}$ charges are introduced, which also give Majorana masses to the left and right handed neutrinos.

The fermion content of the MLRSM is given by

$$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix} \sim (3, 2, 1, \frac{1}{3}), \quad Q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \sim (3^*, 1, 2, \frac{1}{3}),$$

$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \sim (1, 2, 1, -1), \quad \ell_R = \begin{pmatrix} \nu_R \\ e_R \end{pmatrix} \sim (1, 1, 2, -1).$$

Similarly, the scalar content of the MLRSM is

$$\Phi = \begin{pmatrix} \phi^0_{11} & \phi^+_{11} \\ \phi^-_{12} & \phi^0_{12} \end{pmatrix} \sim (1, 2, 2, 0)$$
\[
\Delta_L = \begin{pmatrix}
\frac{\delta_L^+}{\sqrt{2}} & \frac{\delta_L^{++}}{\sqrt{2}} \\
\delta_L^0 & -\frac{\delta_L^-}{\sqrt{2}}
\end{pmatrix} \sim (1, 3, 1, 2), \quad \Delta_R = \begin{pmatrix}
\frac{\delta_R^+}{\sqrt{2}} & \frac{\delta_R^{++}}{\sqrt{2}} \\
\delta_R^0 & -\frac{\delta_R^-}{\sqrt{2}}
\end{pmatrix} \sim (1, 1, 3, 2)
\]

where the numbers in brackets denote the transformations of the fields under the gauge group \(SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}\) of the model. During the spontaneous symmetry breaking of MLRSM gauge group down to the SM gauge group, the neutral component of the Higgs triplet \(\Delta_R\) acquires a non-zero vacuum expectation value (vev) after which the neutral components of Higgs bidoublet \(\Phi\) acquire non-zero vev’s to break the SM gauge symmetry into the \(U(1)\) of electromagnetism. This symmetry breaking chain can be denoted as:

\[
SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow{\langle \Delta_R \rangle} SU(2)_L \times U(1)_Y \xrightarrow{\langle \Phi \rangle} U(1)_{em}
\]

Denoting the vev of the two neutral components of the bidoublet as \(k_1, k_2\) and that of triplets \(\Delta_{L,R}\) as \(v_{L,R}\) and considering \(g_L = g_R, k_2 \sim v_L \approx 0\) and \(v_R \gg k_1\), the approximate expressions for gauge boson masses after symmetry breaking can be written as

\[
M_{W_L}^2 = \frac{g^2}{4} k_1^2, \quad M_{W_R}^2 = \frac{g^2}{2} v_R^2
\]

\[
M_{Z_L}^2 = \frac{g^2 k_1^2}{4 \cos^2 \theta_w} \left(1 - \frac{\cos^2 2\theta_w}{2 \cos^4 \theta_w} \frac{k_1^2}{v_R^2}\right), \quad M_{Z_R}^2 = \frac{g^2 v_R^2 \cos^2 \theta_w}{\cos 2\theta_w}
\]

where \(\theta_w\) is the Weinberg angle. If we consider tiny but non-zero \(k_2\), it gives rise to a left-right mixing between \(W_L - W_R\) given by

\[
\tan \theta_{LR} = -\frac{2k_1 k_2}{v_R^2}
\]

Even if we switch off the vev \(k_2\), then also there can be non-zero \(W_L - W_R\) mixing, generated at one loop level. This can be calculated as [54, 55]

\[
\sin 2\theta_{LR} = \frac{2W_{LR}}{\sqrt{\left(M_{W_R}^2 - M_{W_L}^2\right)^2 + 4W_{LR}^2}}
\]

\[
W_{LR} = \frac{4\pi\alpha}{\sin^2 \theta_W} \sum_{u,d} m_u m_d V_{u,d} V_{u,d}^* f(x_{u,d}); \quad x_{i,j} = \frac{m_i^2}{m_j^2}
\]

\[
f(x_{i,j}) = \frac{1}{16\pi^2} \left[\frac{x_{i,j} \ln(x_{i,j}) + 1 - x_{i,j}}{1 - x_{i,j}} + \ln \left(\frac{\mu^2}{m_j^2}\right)\right]
\]

The relevant Yukawa couplings for fermion masses can be written as

\[
\mathcal{L}_{\nu}^{II} = y_{ij} \bar{\ell}_i L \Phi_{i,j \nu} + y'_{ij} \bar{\ell}_i L \Phi'_{i,j \nu} + Y_{ij} \bar{q}_i L \Phi_{i,j \nu} + Y'_{ij} \bar{q}_i L \Phi'_{i,j \nu} + \text{h.c.}
\]

\[
+ f_{ij} \left(\bar{q}_{i \nu} C i\sigma_2 \Delta_R \ell_{j \nu} + (R \leftrightarrow L)\right) + \text{h.c.}
\]
where $\tilde{\Phi} = \tau_2 \Phi^* \tau_2$. In the above Yukawa Lagrangian, the indices $i, j = 1, 2, 3$ correspond to the three generations of fermions. The Majorana Yukawa couplings $f$ is same for both left and right handed neutrinos because of the inbuilt left-right symmetry ($f_L = f_R$). These couplings $f$ give rise to the Majorana mass terms of both left handed and right handed neutrinos after the triplet Higgs fields $\Delta_{L,R}$ acquire non-zero vev. Although it is the $\Delta_R$ field which gets a vev at high scale breaking the left-right symmetry, the subsequent electroweak symmetry breaking induces a non-zero vev to the left handed counterpart. The induced vev for the left-handed triplet $v_L$ can be shown for generic LRSM to be

$$v_L = \gamma \frac{M^2_{WL}}{v_R}$$

with $M_{WL} \sim 80.4$ GeV being the weak boson mass such that

$$|v_L| << M_{WL} << |v_R|$$

In general $\gamma$ is a function of various couplings in the scalar potential of generic LRSM. Using the results from Deshpande et al., [53], $\gamma$ is given by

$$\gamma = \frac{\beta_2 k_1^2 + \beta_1 k_1 k_2 + \beta_3 k_2^2}{(2 \rho_1 - \rho_3)(k_1^2 + k_2^2)}$$

(5)

where $\beta, \rho$ are dimensionless parameters of the scalar potential. Without any fine tuning $\gamma$ is expected to be of the order unity ($\gamma \sim 1$). However, for TeV scale type I+II seesaw, $\gamma$ has to be fine-tuned as we discuss later. The type II seesaw formula for light neutrino masses can now be expressed as

$$M_\nu = \gamma M^2_{WL}/v_R M_{RR} - m_{LR} M^{-1}_{RR} m^T_{LR}$$

(6)

If we consider the first term on the right hand side of the above expression, one can make an estimate of neutrino mass for TeV scale LRSM. Considering $v_R \sim 6$ TeV, the type II seesaw term will be of the order of light neutrino mass $M_\nu \sim 0.1$ eV if

$$\gamma \approx \frac{5.6 \times 10^{-7}}{M_R}$$

where $M_R$ is the right handed neutrino mass. Thus, for TeV scale right handed neutrino masses, the dimensionless parameter $\gamma$ is fine tuned to the level of $10^{-10} - 10^{-9}$ in order to get correct order of neutrino masses. This will involve unnatural fine tuning of the scalar potential parameters appearing in the expression for $\gamma$ given in (5). Similar but slightly less
fine tuning is involved in the type I seesaw term for TeV scale $M_{RR}$. The Dirac Yukawa couplings should be fine tuned to around $10^{-6} - 10^{-5}$ in order to get light neutrino mass of order 0.1 eV.

III. DARK MATTER IN LRSM

The minimal LRSM discussed above does not have a stable cold dark matter candidate. One can however, minimally extend the model by including additional scalar or fermionic multiplets in the spirit of minimal dark matter scenario [56–58]. Such minimal dark matter scenario in LRSM has been studied recently by the authors of [59, 60]. In these models, the dark matter candidate is stabilised either by a $Z_2 = (-1)^{B-L}$ subgroup of the $U(1)_{B-L}$ gauge symmetry or due to an accidental symmetry at the renormalisable level due to the absence of any renormalisable operator leading to dark matter decay [54]. Some more studies on left-right dark matter also appeared in the recent works [45, 46, 55, 61, 62]. The possibility of right handed neutrino dark matter in a different version of LRSM where the right handed lepton doublets do not contain the usual charged leptons, was also studied in the recent works [63–65].

The possibility of warm dark matter within minimal LRSM was also studied in the works [66, 67]. In these works, the lightest right handed neutrino with keV scale mass was considered to be the WDM candidate. Such a WDM candidate can decay into a light neutrino and a photon at one loop level and can be cosmologically long-lived if the mixing with the light active neutrinos are appropriately tuned. Such a keV scale right handed neutrino typically gets overproduced in the early Universe, by virtue of its gauge interactions with the standard model particles. In the above mentioned works, the abundance of WDM was brought to the observed DM limits by late time entropy dilution mechanism due to the late decay of heavier right handed neutrinos [68]. In such scenarios, we need to fine tune several Yukawa couplings in order to keep the mixing of WDM with light neutrinos small as well as to allow the late decay of heavier right handed neutrinos. Also, if the WDM in such models are responsible for the origin of the $3.55$ keV monochromatic X-ray line, then one requires even smaller mixing angle, requiring unnatural fine-tuning.

We first discuss the phenomenology of combining CDM and WDM within simple extensions of minimal LRSM. After that we also consider a different version of LRSM where
the dark matter sector is much richer and at the same time requires much less fine-tuning compared to the scenarios based on extensions of the minimal LRSM.

A. Mixed Dark Matter in Minimal LRSM

One can have a mixture of CDM and WDM by extending the minimal LRSM with a pair of scalar doublets. One can also include higher multiplets, but we stick to doublet as it is the minimal non-trivial multiplet of $SU(2)$. Scalar doublet DM in LRSM was studied in [60] considering only their gauge interactions so that, both of them can be stable. However, if scalar interactions are taken into account, only the lighter of them will be stable and can be a CDM candidate. If the lightest right handed neutrino is in the keV range, it can be long-lived and serve as a WDM candidate.

To have a minimal mixed dark matter scenario in this model, we can either add a pair of scalar doublets $\eta_L(1,2,1,-1), \eta_R(1,1,2,-1)$ or a pair of fermion triplets $\Sigma_L \equiv (1,3,1,0), \Sigma_R \equiv (1,1,3,0)$ to the minimal LRSM discussed above. Here, the numbers in brackets correspond to the quantum numbers under the gauge symmetry of the model $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The neutral components of these multiplets can be good CDM candidates. While the discussion of left scalar dark matter is similar to the inert scalar doublet model studied extensively in the literature [56, 69–78], the right handed scalar dark matter $\eta^0_R$ was studied recently by the authors of [54, 55, 60]. The left fermion triplet dark matter was studied a few years back [79] whereas the right fermion dark matter was studied more recently within the context of LRSM in [45, 46, 59, 60, 80, 81]. As discussed in these works, the phenomenology of fermion triplet CDM is much richer due to the fact that both the left and right handed fermion can be stable giving rise to a multi-component CDM scenario. On the other hand, among the scalar doublets, only the lighter of them will be stable as the heavier one can decay into the lighter one and standard model particles through its interaction with the light Higgs.

IV. NON-MINIMAL LRSM WITH UNIVERSAL SEESAW

In the minimal model, one requires fine-tuning in the couplings of the right handed neutrinos, so that the lightest of them can be a keV scale WDM whereas the heavier ones
can be long-lived so that their late decay can release entropy and bring the abundance of WDM into the correct regime. In this section, we consider a different version of LRSM where such fine tunings can be controlled naturally. This is based on the version of LRSM without the usual Higgs bidoublet, popularly known as the LRSM with universal seesaw [82–86]. Recently, this model was also studied in the context of light Dirac (pseudo-Dirac) neutrinos [54, 55] and $SU(3)_c \times SU(3)_R \times SU(3)_L \times U(1)_X$ gauge symmetry [87]. This model has been studied in the context of cosmology [88] and neutrinoless double beta decay in several works [89–91]. The model also received attention [92–95] in the context of the 750 GeV di-photon excess at LHC [96] (which however, did not get confirmed in subsequent measurements by the same experiments). The model has the following additional fermion content compared to the minimal LRSM

\[
U_L(3, 1, 1, \frac{4}{3}), \quad U_R(3^*, 1, 1, \frac{4}{3}), \quad D_L(3, 1, 1, -\frac{2}{3}), \quad D_R(3^*, 1, 1, -\frac{2}{3}),
\]

\[
E_{L,R}(1, 1, 1, -2), \quad N_{L,R}(1, 1, 1, 0)
\]

per generation of quarks and leptons. Instead of Higgs bidoublet, this model has a pair of Higgs doublets

\[
H_L(1, 2, 1, -1), \quad H_R(1, 1, 2, -1)
\]

Due to the absence of the usual bidoublet, the left and right handed fermion doublets of the MSLRM can not directly couple to each other. However, they can couple to the scalar fields $H_{L,R}$ via the additional vector like fermions.

\[
\mathcal{L} \supset Y_U(\overline{Q}_L H_L U_R + \overline{Q}_R H_R U_L) + Y_D(\overline{Q}_L H_L^\dagger D_R + \overline{Q}_R H_R^\dagger D_L) + M_{U,L} \overline{U}_L U_R + M_{D,L} \overline{D}_L D_R \\
+ Y_E(\overline{t}_L H_L^\dagger E_R + \overline{t}_R H_R^\dagger E_L) + Y_\nu(\overline{\ell}_L H_L N_R + \overline{\ell}_R H_R N_L) + M_{E,L} \overline{E}_L E_R + M_{N,L} \overline{N}_L N_R \\
+ \frac{1}{2} M_N^M (N_L N_L + N_R N_R) + \text{h.c.}
\]

(7)

The scalar potential of the model can be written as

\[
V = \mu_H^2 \left( H_L^\dagger H_L + H_R^\dagger H_R \right) + \lambda_H \left( (H_L^\dagger H_L)^2 + (H_R^\dagger H_R)^2 \right) + \lambda' (H_L^\dagger H_L)(H_R^\dagger H_R)
\]

(8)

The scalar fields can acquire non-zero vev as

\[
\langle H_R \rangle = \left( \frac{\nu_R}{\sqrt{2}} \right), \quad \langle H_L \rangle = \left( \frac{\nu_L}{\sqrt{2}} \right)
\]

(9)
The vev of the neutral component of $H_R$ spontaneously breaks the symmetry of the LRSM
to that of the SM while the vev of the neutral component of $H_L$ gives rise to the usual
electroweak symmetry breaking. In other words, the desired symmetry breaking chain is

$$SU(2)_L \times SU(2)_R \times U(1)_{B-L} \quad \langle H^0_R \rangle \quad SU(2)_L \times U(1)_Y \quad \langle H^0_L \rangle \quad U(1)_Q$$

After integrating out the heavy fermions, the charged fermions of the standard model develop
Yukawa couplings to the scalar doublet $H_L$ as follows

$$y_u = Y_u \frac{v_R}{M_U} Y^T_U, \quad y_d = Y_D \frac{v_R}{M_D} Y^T_D, \quad y_e = Y_E \frac{v_R}{M_E} Y^T_E$$

The apparent seesaw then can explain the observed mass hierarchies among the three generations of fermions. The neutrino mass arises in a more complicated seesaw due to additional Majorana mass terms for $N_L, N_R$ shown above. As shown in [88], the neutrino mass matrix in the basis $(\nu_L, \nu_R)$ can have three independent terms

$$m_L = -Y_\nu \frac{v_L}{M_N} Y^T_\nu v_L$$

$$m_R = -Y_\nu \frac{v_R}{M_N} Y^T_\nu v_R$$

$$m_D = Y_\nu \frac{1}{M^D_N} (M_N^D) \frac{1}{M^M_N} Y^\dagger_\nu v_L v_R$$

after integrating out the heavy neutral fermions $N_{L,R}$. This mass matrix can further be
diagonalised to get the type I seesaw formula for neutrino mass which however, is suppressed
by $(\frac{M^D_N}{M^M_N})^2$ compared to $m_L$. Thus, the left and right handed neutrinos have Majorana mass
terms which are related by

$$m_R = m_L \frac{v^2_R}{v^2_L} \quad (10)$$

If we have a TeV scale left-right symmetry breaking say $v_R \approx 6 \text{ TeV}$, then the above relation
says $m_R \approx 595 m_L$. Thus for $m_L \leq 0.1 \text{ eV}$, the right handed neutrino masses are $m_R \leq 59.5 \text{ eV}$. Since the right and left handed neutrino masses are different by a factor $\frac{v_R}{v_L}$ where $v_L = 246 \text{ GeV}$ is the electroweak vev, one can easily adjust the scale $v_R$ in order to have
the desired right handed neutrino spectrum. For example, if we want the heaviest sterile
neutrino mass to be around a few keV, suitable for warm dark matter candidates, we need
to push the left-right breaking scale $v_R$ to beyond 60 TeV. Although such a high scale
symmetry breaking is beyond the reach of present collider experiment, they can naturally
give rise to keV scale right handed neutrinos that can be a warm dark matter candidate, having interesting indirect detection signatures as we discuss in details below. One can similarly derive the mass spectrum of gauge bosons as well, which can be written as

\[ M_{W_L} = \frac{g}{2} v_L, \quad M_{W_R} = \frac{g}{2} v_R, \quad M_{Z_L} = \frac{g}{2} v_L \sqrt{1 + \frac{g_1^2}{g^2 + g_1^2}}, \quad M_{Z_R} = \frac{v_R}{2} \sqrt{(g^2 + g_1^2)}. \]

V. MIXED DARK MATTER IN NON-MINIMAL LRSM

A. Scalar/Fermion CDM and Fermion WDM

Since the light and heavy neutrino masses are related to each other in the LRSM with universal seesaw, through the symmetry breaking scales, the model gets quite restricted once we choose the spectrum of right handed neutrinos. For example, if the lightest right handed neutrino is the WDM with keV scale mass, and the next to lightest right handed neutrino is responsible for late time entropy generation, it puts a tight lower bound on the right handed sector gauge boson masses. Using the relations between heavy and light neutrinos derived in previous section, it is straightforward to see that if we want to keep the lightest right handed neutrino mass at keV scale, the left-right symmetry breaking scale has to be at least \( v_R \sim 60 \) TeV. Assuming the decay of \( N_2 \) to produce the required entropy, if we consider \( M_{N_2} \geq 1 \text{ GeV} \) in order to produce the desired entropy \([67]\), it gives a bound on the left-right symmetry breaking scale as

\[ M_{N_2} = m_2 \frac{v_R^2}{v_L^2} = \sqrt{m_1^2 + \Delta m_{21}^2} v_R \frac{v_R^2}{v_L^2} \geq 1 \text{ GeV} \]

\[ \Rightarrow v_R \geq \frac{1 \text{ GeV}^{1/2} v_L}{(\Delta m_{21}^2)^{1/4}} \approx 82 \text{ PeV} \quad (11) \]

where we have used \( v_L = 246 \text{ GeV} \), \( \Delta m_{21}^2 \approx 8 \times 10^{-5} \text{ eV}^2 \). Thus, setting the requirement for the lightest right handed neutrino to be around a keV, one can constrain the light as well as heavy neutrino spectrum along with heavy gauge boson masses from the requirement of generating the correct abundance of WDM. Although the scale of left-right symmetry breaking scale is pushed to a high scale, one can naturally explain the long lifetime of heavy neutrino \( N_2 \) by virtue of its small mixing with light neutrinos that can be achieved naturally in this model due to suppressed Dirac mass term discussed above. One can however, bring the scale \( v_R \) lower, by giving a new contribution to light neutrino mass so that the constrained
relation between right and left handed neutrino masses can be made weaker. As we will see later, even for a mass of $W_R$ around a few tens of TeV, we can achieve the correct entropy dilution due to $N_2$, in case $N_1$ is the WDM candidate.

On the other hand, if we add a pair of scalar doublets $\eta_{L,R}$ to this model so that the lightest of the neutral components can be a CDM candidate, their stability issue arises as they possess the same quantum numbers as $H_{L,R}$ taking part in the symmetry breaking. The stability issue can be taken care of by considering a discrete $Z_2$ symmetry under which $\eta_{L,R} \rightarrow -\eta_{L,R}$ so that the lighter of them will be stable [55]. This stability issue does not arise if we consider a pair of fermion triplet dark matter $\Sigma_{L,R}$ since there is no renormalisable coupling between them and the other particles of the model. In such a case, additional $Z_2$ symmetry is not necessary for CDM stability.

![FIG. 1: Radiative generation of heavy-light neutrino mixing.](image)

**B. Scalar/Fermion CDM and Scalar WDM**

One can also have a more exotic version of the LRSM discussed in the previous section where the warm dark matter component can be a fundamental scalar. Although there involves an issue of fine-tuning in generating keV scale or smaller mass of a scalar, there are some advantages of this scenario. Firstly, the lower bound on dark matter mass is not applicable like it is there in case of fermion DM from the galactic phase space criteria. If a fermion DM is to constitute the entire DM in a galaxy, then below a certain mass, the phase space density of DM particles that would be required by the observed amount of DM
in dwarf galaxies, would violate the Pauli exclusion principle. This lower bound on fermion DM mass (around 0.4 keV) was calculated long back by Tremaine-Gunn [97]. For scalar DM, this bound is relaxed as the Pauli exclusion principle does not apply there.

Since the universal seesaw for all fermions results in a rather strict mass spectrum of light and heavy neutrinos, we suitably extend the model by a pair of triplet scalars $\Delta_L \equiv (1, 3, 1, 2), \Delta_R \equiv (1, 1, 3, 2)$ similar to the minimal LRSM discussed before. This will decouple the origin of light and heavy neutrino mass and enable us to get rid of the tight constraint on the left-right symmetry breaking scale (from equation (10)) which exists in the LRSM with universal seesaw for all fermions. However, the coupling between light and heavy neutrinos can remain suppressed due to the smallness of $\left(\frac{M_D^2}{M_N^2}\right)^2$ factor discussed before. Instead of choosing a small $\left(\frac{M_D^2}{M_N^2}\right)^2$ factor, we can also generate heavy-light neutrino mixing at one loop order (as seen from figure 1), if we consider the heavy singlet fermions to be odd under the same $Z_2$ symmetry protecting the CDM candidate (lighter of $\eta_{L,R}^0$ and $\eta_{L,R}^0$) stable. Thus, under the $Z_2$ symmetry we have $\eta_{L,R} \rightarrow -\eta_{L,R}, N_{L,R} \rightarrow -N_{L,R}$. In such a case, the lightest right handed neutrino can be long lived both due to kinematical suppression as well as one loop suppressed coupling with the light neutrinos.

The scalar doublets $\eta_{L,R}$ can not give rise to two stable or long-lived DM candidates due to the presence of quartic terms like $\lambda_{LR}\eta_L^\dagger \eta_R H_L^1 H_R$ in the Lagrangian which can not be prevented by the additional discrete $Z_2$ symmetry or in fact any other symmetries unless it explicitly breaks the in-built left-right symmetry by assigning different charge or transformations to left and right scalar doublets. We therefore introduce a pair of scalar triplets $\Omega_L \equiv (1, 3, 1, 0), \Omega_R \equiv (1, 1, 3, 0)$. In such a setup, the lighter of $\eta_{L,R}^0$ can be CDM while $\Omega_{L,R}^0$ can be WDM.

VI. DARK MATTER RELIC CALCULATION

Several astrophysical and cosmological evidences suggest the presence of dark matter (DM) in our Universe. The latest data collected by the Planck experiment suggests around 26% of the present Universe’s energy density being made up of dark matter [4] as mentioned earlier in (1). According to the list of criteria, a dark matter candidate must fulfil [5], none of the SM particles can qualify for it. In this section, we outline the standard procedures to calculate the abundance of both keV and TeV-ish DM candidates.
A. Cold Thermal Relic

The relic abundance of a dark matter particle $\text{DM}$, which was in thermal equilibrium at some earlier epoch can be calculated by solving the Boltzmann equation

$$\frac{dn_{\text{DM}}}{dt} + 3Hn_{\text{DM}} = -\langle \sigma v \rangle (n_{\text{DM}}^2 - (n_{\text{eq}}^{\text{DM}})^2)$$  \hspace{1cm} (12)$$

where $n_{\text{DM}}$ is the number density of the dark matter particle $\text{DM}$ and $n_{\text{eq}}^{\text{DM}}$ is the number density when DM was in thermal equilibrium. $H$ is the Hubble expansion rate of the Universe and $\langle \sigma v \rangle$ is the thermally averaged annihilation cross section of the dark matter particle DM. In terms of partial wave expansion $\langle \sigma v \rangle = a + bv^2$. Numerical solution of the Boltzmann equation above gives $[98, 99]$

$$\Omega_{\text{DM}}h^2 \approx 1.04 \times 10^9 x_F$$  \hspace{1cm} (13)$$

where $x_F = M_{\text{DM}}/T_F$, $T_F$ is the freeze-out temperature, $M_{\text{DM}}$ is the mass of dark matter, $g_*$ is the number of relativistic degrees of freedom at the time of freeze-out and and $M_{\text{Pl}} \approx 2.4 \times 10^{18}$ GeV is the Planck mass. Dark matter particles with electroweak scale mass and couplings freeze out at temperatures approximately in the range $x_F \approx 20 - 30$. More generally, $x_F$ can be calculated from the relation

$$x_F = \ln \frac{0.038g_{\text{Pl}}M_{\text{DM}} < \sigma v >}{g_*^{1/2} x_F^{1/2}}$$  \hspace{1cm} (14)$$

which can be derived from the equality condition of DM interaction rate $\Gamma = n_{\text{DM}} \langle \sigma v \rangle$ with the rate of expansion of the Universe $H \approx g_*^{1/2} \frac{T^2}{M_{\text{Pl}}}$. There also exists a simpler analytical formula for the approximate DM relic abundance $[100]$

$$\Omega_{\text{DM}}h^2 \approx 3 \times 10^{-27} \text{cm}^3\text{s}^{-1} \langle \sigma v \rangle$$  \hspace{1cm} (15)$$

The thermal averaged annihilation cross section $\langle \sigma v \rangle$ is given by $[101]$

$$\langle \sigma v \rangle = \frac{1}{8m^4TK_i^2(m/T)} \int_{4m^2}^{\infty} \sigma(s - 4m^2)\sqrt{s}K_1(\sqrt{s/T})ds$$  \hspace{1cm} (16)$$

where $K_i$’s are modified Bessel functions of order $i$, $m$ is the mass of Dark Matter particle and $T$ is the temperature.

If there exists some additional particles having mass difference close to that of DM, then they can be thermally accessible during the epoch of DM freeze out. The can give rise
to additional channels through which DM can coannihilate with such additional particles and produce SM particles in the final states. This type of coannihilation effects on dark matter relic abundance were studied by several authors in [102–104]. Here we summarise the analysis of [102] for the calculation of the effective annihilation cross section in such a case. The effective cross section can given as

$$
\sigma_{\text{eff}} = \sum_{i,j} \langle \sigma_{ij} v \rangle r_i r_j

= \sum_{i,j} \langle \sigma_{ij} v \rangle \frac{g_i g_j}{g_{\text{eff}}} \left(1 + \Delta_i\right)^{3/2} \left(1 + \Delta_j\right)^{3/2} e^{-x_F (\Delta_i + \Delta_j)}
$$

(17)

where $x_F = \frac{m_{\text{DM}}}{T_F}$ and $\Delta_i = \frac{m_i - M_{\text{DM}}}{M_{\text{DM}}}$ and

$$
g_{\text{eff}} = \sum_{i=1}^{N} g_i (1 + \Delta_i)^{3/2} e^{-x_F \Delta_i}
$$

(18)

The masses of the heavier components of the inert Higgs doublet are denoted by $m_i$. The thermally averaged cross section can be written as

$$
\langle \sigma_{ij} v \rangle = \frac{x_F}{8m_i^2 m_j^2 M_{\text{DM}} K_2((m_i/M_{\text{DM}}) x_F) K_2((m_j/M_{\text{DM}}) x_F)} \times

\int_{(m_i + m_j)^2}^{\infty} ds \sigma_{ij}(s - 2(m_i^2 + m_j^2)) \sqrt{s} K_1(\sqrt{s} x_F / M_{\text{DM}})
$$

(19)

B. Warm Thermal Relic

The lightest right handed neutrino can be long lived if it has a mass below the mass of an electron, since it can decay only at loop level into lighter particles like standard model neutrinos and photon. Since the right handed neutrino $N_1$ has gauge interactions in LRSM, they can be in thermal equilibrium in the early Universe freezing out subsequently around

$$
T_{fN_1} \approx g_{*}^{1/6} \left(\frac{M_{W_R}}{M_{W_L}}\right)^{4/3} T_{f\nu}
$$

(20)

with $g_{*f}$ being the relativistic degrees of freedom at $T = T_{fN_1}$. It is defined as

$$
g_{*} = \sum_{i \in \text{boson}} \left(\frac{T_i}{T}\right)^4 g_i + \frac{7}{8} \sum_{i \in \text{fermion}} \left(\frac{T_i}{T}\right)^4 g_i.
$$
If all relativistic particles are in equilibrium with each other, it can simply be written as

\[ g_* = \sum_{i \in \text{boson}} g_i + \frac{7}{8} \sum_{i \in \text{fermion}} g_i. \]

In the above equation (20), \( T_{f\nu} \sim 1 - 2 \text{ MeV} \) is the freeze-out temperature of light neutrinos. Thus, for TeV scale \( W_R \), the right handed neutrino can remain in equilibrium until late epochs corresponding to a temperature of a few hundred MeV’s. At such high temperatures, a keV right handed neutrino can behave like a relativistic species whose number and entropy densities can be given as

\[ n = g_* \frac{\zeta(3)}{\pi^2} g_{\nu} T^3, \quad s = \frac{2\pi^4}{45} g_{\text{eff}} T^3 \]

where \( g_* = 1, \frac{3}{4} \) for boson, fermion respectively, and \( g_{\text{eff}} \) is given by

\[ g_{\text{eff}} = \sum_{i \in \text{boson}} \left( \frac{T_i}{T} \right)^3 g_i + \frac{7}{8} \sum_{i \in \text{fermion}} \left( \frac{T_i}{T} \right)^3 g_i. \]

Before QCD phase transition temperature (\( \sim \) a few hundred MeV), since all relativistic species are in equilibrium with each other \( (T_i = T, \forall i) \) we can write the effective relativistic degrees of freedom for entropy density as

\[ g_{\text{eff}} = \sum_{i \in \text{boson}} g_i + \frac{7}{8} \sum_{i \in \text{fermion}} g_i. \]

The \( N_1 \) number density to entropy density after freeze-out is given by

\[ \frac{n_{N_1}}{s} |_f = \frac{1}{g_{*f}} \frac{135\zeta(3)}{4\pi^4} \]

The present abundance of \( N_1 \) in comparison to the total DM abundance is

\[ \frac{\Omega_{N_1}}{\Omega_{DM}} = \frac{n_{N_1}}{s} |_f \frac{M_{N_1} s_0}{\Omega_{DM} \rho_c} \approx \frac{1}{g_{*f}} 7.59 \times 10^3 \left( \frac{M_{N_1}}{7 \text{ keV}} \right) \]

with \( s_0, \rho_c = \frac{3H_0^2}{8\pi G} \) being the entropy density and critical density of the present Universe. Thus, even if the freeze-out occurs above the electroweak symmetry breaking so that \( g_{*f} \approx 107 \), the abundance of \( N_1 \) will be much more than the observed DM, overclosing the Universe. For decoupling temperature of a few hundred GeVs for which \( g_{*f} \approx 60 \), we can normalise the abundance of \( N_1 \) as

\[ \frac{\Omega_{N_1}}{\Omega_{DM}} \approx 1.265 \times 10^2 \left( \frac{60}{g_{*f}} \right) \left( \frac{M_{N_1}}{7 \text{ keV}} \right) \]
Similarly, the keV scalar DM can also remain in thermal equilibrium by virtue of gauge interactions. If the neutral component of $\Omega_R$ is the keV scalar, then it has interactions with $W_R$ bosons. Since, a neutral scalar can not have three point interactions with $W_R$,$Z_R$ bosons, the only interactions it can have is the four point ones of the type $W^+_R W^-_R \Omega^0_R \Omega^0_R$. The interaction that can keep $\Omega^0_R$ in equilibrium until late epochs is $\gamma\gamma \rightarrow \Omega^0_R \Omega^0_R$ through a $W_R$ boson loop. The cross section can be estimated as

$$\sigma(\gamma\gamma \rightarrow \Omega^0_R \Omega^0_R) = \frac{E_{\Omega_R}^2 F_W^2}{64\pi} \left( \frac{e^2 g^2}{32\pi^2 M_{W_R}^2} \right)^2$$

(25)

Here $E_{\Omega_R} = \rho_{\Omega_R}/n_{\Omega_R} = 2.7T$ and $F_W = 7$ is a loop function. To find the decoupling temperature, we equate the interaction rate $\Gamma$ with the Hubble expansion rate $H$ as follows.

$$\Gamma = n_\gamma \sigma v = H(T_{f\Omega}) = 1.66 \sqrt{g_{*f}} \frac{T_{f\Omega}^2}{M_{Pl}}$$

(26)

$$\implies \frac{2(3)(3)}{\pi^2} T_{f\Omega}^3 \frac{E_{\Omega_R}^2 F_W^2}{64\pi} \left( \frac{e^2 g^2}{32\pi^2 M_{W_R}^2} \right)^2 = 1.66 \sqrt{g_{*f}} \frac{T_{f\Omega}^2}{M_{Pl}}$$

(27)

$$\implies T_{f\Omega} = 3.58 \times 10^{-4} g_{*f}^{1/6} \left( \frac{M_{W_R}}{\text{GeV}} \right)^{4/3}$$

Thus, even if we take the lowest possible value of $M_{W_R} \sim 3$ TeV corresponding to $g_{*f}$ value of at least 107 (same as that of SM particles at high temperatures), the keV scalar DM freezes out at around $T_{f\Omega} \approx 33$ GeV. Since the scalar WDM decouples while being relativistic, it is straightforward to calculate the present abundance.

The abundance of $\Omega^0_R$ can be written in terms of the ratio of number density to entropy density as

$$Y_{\Omega_R} = \frac{n_{\Omega_R}}{s}$$

(28)

Using the expressions for number and entropy densities for relativistic species, we can write it as

$$Y_{\Omega_R} = \frac{45(3)}{2\pi^4} \frac{g_{*R}}{g_{eff}} \frac{\rho_{\Omega^0_R}}{\rho_c}$$

(29)

Since $Y_{\Omega_R}$ since decoupling is conserved as the Universe evolves, the present abundance can be written as

$$\Omega_{\Omega_R} = Y_{\Omega_R} m_{\Omega_R} \rho_c$$

(30)

where $s_0 \approx 2.89 \times 10^3$ cm$^{-3}$ is the entropy density and $\rho_c \approx 1.05 \times 10^{-5} h^2$ GeV cm$^{-3}$ is the critical density of the Universe at present. Using $g_{*R} = 1, h = 0.68$ we can find

$$\Omega_{\Omega_R} = 1.645 \times 10^8 \frac{1}{g_{eff}} \left( \frac{m_{\Omega_R}}{\text{GeV}} \right).$$

(31)
Using appropriate normalisations, we can rewrite it as

$$\Omega_{\Omega_R}^0 = 11.54 \left( \frac{100}{g_{\text{eff}}} \right) \left( \frac{m_{\Omega_R}}{7 \text{ keV}} \right).$$

(32)

Here $g_{\text{eff}} \approx 100$ is the appropriate relativistic degrees of freedom at freeze-out temperature $T_{f\Omega} \approx 33$ GeV corresponding to $M_{W_R} \sim 3$ TeV. Also, the mass of the scalar is normalised to 7 keV, which has interesting implications for the origin of 3.55 keV X-ray line as we discuss below. Thus, for this generic normalisations, the thermal abundance of keV scalar DM comes out to be around 43 times the required DM abundance.

This requires entropy dilution after freeze-out to bring down the abundance of $\Omega_{N_1} \leq \Omega_{\text{DM}}$. Late decay of heavier right handed neutrinos like $N_2$ can release such entropy. Such a decay should however occur before the big bang nucleosynthesis (BBN) temperature $T_{\text{BBN}} \sim \mathcal{O}(\text{MeV})$ in order to be consistent with successful BBN predictions. Such late decay of long lived particles can release extra entropy and dilute the abundance of keV dark matter to bring it into the observed limit [68]. The dilution factor due to the decay of such a heavy long lived particle $N_2$ is given by [68]

$$d = \frac{s_{\text{before}}}{s_{\text{after}}} \approx 0.58 \left[ g_*(T_r) \right]^{-1/4} \frac{\sqrt{\Gamma_{N_2} M_{\text{Pl}}}}{M_{N_2} Y_{N_2}},$$

(33)

where $\Gamma_{N_2}$ is the decay width of the heavy particle with mass $M_{N_2}$ and

$$Y_{N_2} = \frac{n}{s} = \frac{135}{4\pi^4} \frac{\zeta(3)}{g_*(T_{fN_2})}$$

is the initial abundance of the particle $N_2$ before it started to decay. Also, $g_*(T_r)$ is the relativistic degrees of freedom at a temperature $T_r$ just after the decay of $N_2$. This temperature to which the Universe cools down to following the release of entropy due to the decay of $N_2$ can be approximated as

$$T_r \approx 0.78 \left[ g_*(T_r) \right]^{-1/4} \sqrt{\Gamma_{N_2} M_{\text{Pl}}}.$$  

(34)

Also, $g_*(T_{fN_2})$ is the relativistic degrees of freedom at the epoch of $N_2$ freeze-out. For maximum dilution or minimum value of $d$, it is desirable to have $g_*(T_r)$ minimum ($\approx 10.75$), equal to the value of $g_*$ just before BBN.
VII. RESULTS AND DISCUSSION

A. Relic Abundance

We first calculate the relic abundance of different CDM candidates discussed in the work. Since we are considering a mixture of CDM and WDM, we find out the parameter space that gives rise to under-abundant CDM and compare it with the parameter space that gives 100% CDM. We show the parameter space for left scalar doublet DM for two different values of mass splitting within the components of the scalar doublet in figure 2. The results are similar to the inert scalar doublet model discussed extensively in the literature [56, 69–78], but shown here for different fractions of total DM abundance. The DM relic abundance is primarily governed by three parameters: the Higgs-DM coupling $\lambda_L$, DM mass $M_{\eta_L}$ and mass splitting $\Delta M$ between different components of left scalar doublet $\eta_L$. For simplicity, we consider same mass splitting between lightest neutral scalar and the charged as well as neutral pseudoscalar components of the doublet. There are two mass regions which gives rise to correct relic abundance or a sizeable fraction of it, as seen from the right panel plot of figure 2. The thermal DM abundance remains very much suppressed in the intermediate mass regime due to very large annihilations to electroweak gauge bosons. In the left panel plot of figure 2, the low mass regime disappears as the DM coannihilations are very large due to smaller mass splitting $\Delta M$. We also show the freeze-out temperature of left scalar doublet DM in figure 3. This is important for the comparison with the freeze-out temperature of light keV DM, which is expected to be below the freeze-out of heavier DM, so that their relic can be estimated independently, in a simple manner. It is observed that, the freeze-out temperature is same irrespective of the total DM fraction, the scalar doublet contributes to.

We show the parameter space giving rise to different fraction of total DM abundance for right scalar doublet DM in figure 4. The results are similar to the ones shown in [54, 55, 60], but extended here for different fractions of total DM abundance. Similar to these works, here also we do not take the Higgs portal interactions of right scalar doublet DM so that the DM relic abundance is more sensitive to the $SU(2)_R$ gauge sector. In this approximation, the DM relic abundance is mainly governed by the DM coannihilations through $W_R, Z_R$ bosons, specially when the DM mass is less than the gauge boson masses. We show the parameter space in terms of gauge boson mass $M_{W_R}$ and DM mass for two different mass
FIG. 2: Parameter space giving rise to left scalar doublet dark matter including both scalar and gauge interactions. Different coloured lines correspond to different fraction of total DM abundance.

FIG. 3: Freeze-out temperatures of left scalar doublet DM for different mass splittings. Different coloured lines correspond to different fraction of total DM abundance.

splittings $\Delta M$ in figure 4. The corresponding freeze-out temperatures are shown in figure 5.

We also show the parameter space for right fermion triplet DM in figure 6. Since the fermion triplet can annihilate only through gauge interactions, the number of free parameters affecting the relic abundance is less, in fact only two namely, the DM mass and the right
FIG. 4: Parameter space giving rise to right scalar doublet dark matter including gauge interactions. Different coloured lines correspond to different fraction of total DM abundance.

FIG. 5: Freeze-out temperatures of right scalar doublet DM for different mass splittings. Different coloured lines correspond to different fraction of total DM abundance.

handed gauge boson masses. Unlike scalar DM, here the mass splitting between charged and neutral components of the fermion triplet $\Sigma_R$ is not a free parameter but generated at one loop level through gauge boson corrections. This mass splittings between charged and
FIG. 6: The left panel shows the parameter space giving rise to right fermion triplet dark matter abundance. The right panel shows the corresponding freeze-out temperature. Different coloured lines correspond to different fraction of total DM abundance.

Neutral components of right-handed triplet fermion is given by

\[ M_{\Sigma^\pm_R} - M_{\Sigma^0_R} \approx \frac{\alpha_2}{4\pi} g_R^2 g_L^2 \left[ f(r_{W_R}) - c_M^2 f(r_{Z_R}) - s_W^2 s_M^2 f(r_{Z_L}) - c_W^2 s_M^2 f(r_{\gamma}) \right]. \]  

(35)

Here the one loop self-energy corrections through mediations of gauge bosons are presented within the square bracket. For example, the mass splitting with the approximation \( M_\Sigma \gg M_{W_R} \) goes as \( \alpha_2 (M_{W_R} - c_M^2 M_{Z_R}) / 2 \). The sine and cosine of different angles \( c_M, c_W, s_M, s_W \) etc. correspond to the angles involved in the rotation of neutral gauge bosons given by

\[
\begin{pmatrix}
W^3_{L\mu} \\
W^3_{R\mu} \\
B_{\mu}
\end{pmatrix}
= \begin{pmatrix}
c_W c_\phi & c_W s_\phi & s_W \\
- s_W s_M c_\phi - c_M s_\phi & - s_W s_M s_\phi + c_M c_\phi & c_W s_M \\
- s_W c_M c_\phi + s_M s_\phi & - s_W c_M s_\phi - s_M c_\phi & c_W c_M
\end{pmatrix}
\begin{pmatrix}
Z_{L\mu} \\
Z_{R\mu} \\
A_{\mu}
\end{pmatrix}.
\]  

(36)

Also, \( r_B = \frac{M_B}{M_\Sigma} \) and the loop function \( f(r) \) being given as

\[ f(r) \equiv 2 \int_0^1 dt (1 + t) \ln \left[ t^2 + (1 - t) r^2 \right]. \]  

(37)

Since the mass splitting decides the amount of coannihilations, choosing the mass of right handed gauge bosons and DM mass is enough to predict the thermal abundance of DM shown in the left panel plot of figure 6. The right panel plot shows the corresponding freeze-out temperatures. One can similarly calculate the abundance of right fermion triplet DM.
as well for which the mass splitting due to electroweak gauge corrections is

\[ M_{\Sigma^\pm} - M_{\Sigma^0} \simeq \alpha_2 M_W \sin^2(\theta_W/2) + \mathcal{O}(M_W^2/M_\Sigma^2). \]  

(38)

Since there is only one free parameter that decides fermion triplet abundance which is its mass, we do not have any parameter space to show for it. As shown first by the authors of [79], such a triplet satisfies correct relic abundance only for DM mass around 3 TeV. It is important to emphasise that, for such heavy fermion triplet DM, Sommerfeld effects [105–107] are important as the corresponding gauge bosons masses are very small compared to DM mass. Including such effects, pushes the allowed DM mass slightly beyond 3 TeV [107]. Such effects are not important in the right fermion triplet DM, as long as the triplet mass is comparable to right handed gauge boson masses. Detailed calculation of such effects is beyond the scope of the present work and can be found in above references and also in [60] within the context of LRSM.

\[ \Omega_{N_1} / \Omega_{DM} \approx 1.0 \left( \frac{107}{g_{*f}} \right) \left( \frac{M_{N_1}}{1 \text{ keV}} \right) \left( \frac{1 \text{ sec}}{7_{N_2}} \right)^{1/2} \left( \frac{1 \text{ GeV}}{M_{N_2}} \right) \left( \frac{g_* T_{fN_2}}{60} \right) \]  

(39)

FIG. 7: Lifetime of right handed neutrinos including only the gauge boson mediated decay channels. The left panel shows the lifetime of the lightest right handed neutrino \( N_1 \) while the right panel shows the lifetime of the next to lightest right handed neutrino \( N_2 \) keeping \( N_1 \) mass fixed at 7.1 keV.

After showing the parameter space giving rise to the desired relic abundance of CDM, we move on to discussing the requirements for keV scale WDM. For fermion WDM that is \( N_1 \), the final relic abundance after entropy dilution due to \( N_2 \) decay discussed above is given by
Similarly, the abundance of scalar WDM after entropy dilution due to the decay of $N_2$ is

$$\frac{\Omega_{N_2}}{\Omega_{\text{DM}}} \approx 1.0 \left( \frac{107}{g_{*f}} \right) \left( \frac{M_{N_2}}{1 \text{ keV}} \right) \left( \frac{1 \text{ sec}}{\tau_{N_1}} \right)^{1/2} \left( \frac{1 \text{ GeV}}{M_{N_1}} \right) \left( \frac{g_{*}(T_{fN_1})}{84} \right)$$

(40)

It can be seen from figure 7 that the required lifetime of decaying particles releasing entropy can be achieved for suitable values of right handed gauge boson masses, taking into account of gauge mediated decay channels only. If we consider non-zero left-right neutrino mixing $\theta_{\nu}$, the decay lifetime can be even shorter, insufficient for the correct entropy dilution. Figure 8, shows the lifetime as a function of $W_L - W_R$ mixing $\theta_{LR}$ (whose one radiative origin is given in (3)) for a fixed value of $\theta_{\nu}$. Such a small value of $\theta_{\nu}$ can be naturally generated either at tree level or one loop level without much fine tuning, as discussed before.
B. Indirect Detection

Due to the existence of two DM components with widely separated mass scales, the models discussed in this work can have very interesting indirect detection signatures different from single component DM. Since the CDM has mass in the GeV-TeV scale and WDM has keV scale mass in this setup, they can annihilate or decay into SM particles with different energies. Among such SM final state particles, photons and neutrinos, being electromagnetically neutral, have the potential to reach the indirect detection experiments without getting affected in the intermediate regions. If DM is of CDM type with typical masses in the GeV-TeV scale, such photons lie in the gamma ray regime whereas for keV scale WDM they correspond to X-ray part of the electromagnetic spectrum.

\[ \sigma v \left[ \text{cm}^3/\text{s} \right] \]

\[ M \eta_0 \]

\[ (\text{GeV}) \]

\[ \lambda_L = 0.1 \]

\[ \lambda_L = 10^{-3} \]

\[ \Delta M = 1 \text{GeV} \]

\[ \Delta M = 10 \text{ GeV} \]

\[ \Delta M = 50 \text{ GeV} \]

\[ 10^{-30} \]

\[ 10^{-29} \]

\[ 10^{-28} \]

\[ 10^{-27} \]

\[ 10^{-26} \]

\[ 10^{-25} \]

\[ 10^{-24} \]

\[ 10^{-23} \]

\[ 10^{-22} \]

\[ 10^{-21} \]

\[ 10^{-20} \]

\[ 10^{-19} \]

\[ 10^{-18} \]

\[ 10^{-17} \]

\[ 10^{-16} \]

\[ 10^{-15} \]

\[ 10^{-14} \]

\[ 10^{-13} \]

\[ 10^{-12} \]

\[ 10^{-11} \]

\[ 10^{-10} \]

\[ 10^{-9} \]

\[ 10^{-8} \]

\[ 10^{-7} \]

\[ 10^{-6} \]

\[ 10^{-5} \]

\[ 10^{-4} \]

\[ 10^{-3} \]

\[ 10^{-2} \]

\[ 10^{-1} \]

\[ 10^{0} \]

\[ 10^{1} \]

\[ 10^{2} \]

\[ 10^{3} \]

\[ 10^{4} \]

\[ 10^{5} \]

\[ 10^{6} \]

\[ 10^{7} \]

\[ 10^{8} \]

\[ 10^{9} \]

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\[ 10^{11} \]

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\[ 10^{30} \]

\[ \Delta \Omega \] is the solid angle corresponding to the observed region of the sky, \( \langle \sigma v \rangle \) is the thermally averaged DM annihilation cross section, \( dN/dE \) is the average gamma ray spectrum.

FIG. 9: Left scalar doublet DM annihilation into \( W^+W^- \) final states for mass splitting 1 GeV and 10 GeV respectively. Thick solid lines show the limits obtained by combining Fermi-LAT observations of 15 dSphs with MAGIC observations of Segue 1. The thin-dotted line, green and yellow bands show, respectively, the median and the symmetrical, two-sided 68% and 95% containment bands for the distribution of limits under the null hypothesis, details of which can be found in [108].

The observed differential gamma ray flux produced due to DM annihilations is given by

\[ \frac{d\Phi}{dE}(\Delta \Omega) = \frac{1}{4\pi} \langle \sigma v \rangle J(\Delta \Omega) \frac{dN}{2M_{DM}^2 dE} \]  \hspace{1cm} (41)
FIG. 10: Left scalar doublet DM annihilation into $W_L^+ W_L^-$ final states for mass splitting 1 GeV and 10 GeV respectively, after incorporating the constraints from relic abundance criteria. Thick solid lines show the limits obtained by combining Fermi-LAT observations of 15 dSphs with MAGIC observations of Segue 1. The thin-dotted line, green and yellow bands show, respectively, the median and the symmetrical, two-sided 68% and 95% containment bands for the distribution of limits under the null hypothesis, details of which can be found in [108].

Per annihilation process and the astrophysical $J$ factor is given by

$$J(\Delta \Omega) = \int_{\Delta \Omega} d\Omega' \int_{LOS} dl \rho^2(l, \Omega').$$  \hspace{1cm} (42)$$

In the above expression, $\rho$ is the DM density and LOS corresponds to line of sight. Therefore, measuring the gamma ray flux and using the standard astrophysical inputs, one can constrain the CDM annihilation into different charged final states like $\mu^+ \mu^-, \tau^+ \tau^-, W_L^+ W_L^-, b\bar{b}$. Since DM can not couple to photons directly, gamma rays can be produced from such charged final states. Using the bounds on DM annihilation to these final states arising from the global analysis of the Fermi-LAT and MAGIC observations of dSphs [108], we show the status of our model for different benchmark values of parameters. It turns out that some region of parameter space for left scalar doublet CDM studied in this work, the indirect detection constraints are rather severe, if CDM constitutes 100% of the DM of the Universe. Here we see, how a multi-component DM scenario can help us to relax such bounds.

We first show the CDM (left scalar doublet) annihilations into $W_L^+ W_L^-$ final states for
FIG. 11: Left scalar doublet DM annihilation into $b\bar{b},\tau^+\tau^-$ final states, the vertical lines corresponding to the values of DM mass where total relic abundance criteria is satisfied. Thick solid lines show the limits obtained by combining Fermi-LAT observations of 15 dSphs with MAGIC observations of Segue 1. The thin-dotted line, green and yellow bands show, respectively, the median and the symmetrical, two-sided 68% and 95% containment bands for the distribution of limits under the null hypothesis, details of which can be found in [108].

two different values of DM-Higgs coupling $\lambda_L$ in figure 9. It can seen that, as we increase the mass splitting between the components of the scalar doublet, the annihilation cross section to $W_L^+W_L^-$ final states increases. In fact, the mass splitting of 50 GeV is ruled out by the Fermi-LAT+MAGIC bounds even beyond DM mass of 1 TeV. While these plots show the overall behaviour of the annihilation cross section as a function of mass, it does not incorporate the constraints from relic abundance criteria. In figure 10, we show the same annihilation channel by taking relic abundance criteria into account. As can be seen, the constraints become weaker for sub-dominant left scalar doublet DM. As the mass splitting is increased, the relic abundance gets more and more suppressed in the high mass regime of left scalar doublet DM. It can be seen from the right panel of figure 10 that the line corresponding to 100% contribution to DM disappears if the mass splitting is 10 GeV. For 50 GeV mass splitting, the relic abundance remains much suppressed, far below even 25% of the total abundance, and hence not shown here. Similarly, the annihilation cross sections to $b\bar{b},\tau^+\tau^-$ respectively, are shown in figure 11, for mass splitting of 10 GeV.
coupling $\lambda_L$ is increased, the cross section to these final states also increase as expected. The vertical lines in the plots shown in figure 11 corresponds to the values of DM mass where total relic abundance criteria is satisfied. It can be seen that, the indirect detection constraints on DM annihilations to $b\bar{b}$ final states almost rule out the DM-Higgs coupling $\lambda_L = 0.1$. However, this will be allowed, if we have a multicomponent DM scenario. For example, if the left scalar doublet contributes only 50% of the total DM abundance, then the indirect detection line (or bands) will go up by a factor of 4 as annihilation rate is proportional to DM density squared. The constraints for $\mu^+\mu^-$ final states will be weaker due to smaller Yukawa couplings for muons with the SM Higgs and hence not shown here.

For keV scale WDM, one can have another interesting signature at indirect detection experiments in terms of a monochromatic x-ray line. Such a monochromatic x-ray line could have already been seen by the XMM-Newton X-ray telescope as mentioned earlier. Though one requires future data to confirm this claim as discovery, it nevertheless motivates one to look for rich particle physics explanations. In the present models, such a decay can occur due to the long lived nature of both fermion and scalar WDM. The fermion WDM, which is the lightest right handed neutrino, can decay into a photon and a light neutrino at one loop level through the diagrams shown in figure 12. Since the light neutrinos can be considered to be almost massless, such a decay can lead to the final states carrying energy $M_{WDM}/2$ each. The two processes arise due to heavy-light neutrino mixing and $W_L - W_R$ mixing. It is interesting to note that, both these mixings can arise at one loop level in LRSM with universal seesaw mentioned above, making such decay widths naturally small, required for long livedness of WDM. The one loop $W_L - W_R$ mixing is given in (3). In figure 12, the heavy and light neutrino mass eigenstates are denoted by $N_i, \nu_i$ respectively whereas $W_i, (i = 1, 2)$ correspond to the physical mass eigenstates of $W_L, W_R$ gauge bosons. The decay width is given as

$$\Gamma_{N_1 \rightarrow \gamma\nu} = \frac{(m_{N_1}^2 - m_{\nu_i}^2)^3}{16\pi m_{N_1}^3} (|\sigma_L|^2 + |\sigma_R|^2)$$

(43)
where

\[
\sigma_L = \frac{ieg^2}{16\pi^2} \left[ \frac{\sin(2\theta_\nu)}{2} \left( \frac{\cos^2(\theta_{LR}) g(m_{W_i}, t_{W_i}) - \sin^2(\theta_{LR}) g(m_{W_2}, t_{W_2})}{m_{W_1}^2} \right) - \cos^2(\theta_\nu) \frac{\sin(2\theta_\nu)}{2} \left( \frac{f(m_{W_i}, t_{W_i}) - f(m_{W_2}, t_{W_2})}{m_{W_1}^2} \right) + \sin^2(\theta_\nu) \frac{\sin(2\theta_\nu)}{2} \left( \frac{h(m_{W_i}, t_{W_i}) - h(m_{W_2}, t_{W_2})}{m_{W_1}^2} \right) \right]
\]

\[
\sigma_R = -\frac{ieg^2}{16\pi^2} \left[ \frac{\sin(2\theta_\nu)}{2} \left( \frac{\cos^2(\theta_{LR}) g(m_{W_i}, t_{W_i}) - \sin^2(\theta_{LR}) g(m_{W_2}, t_{W_2})}{m_{W_1}^2} \right) + \sin^2(\theta_\nu) \frac{\sin(2\theta_\nu)}{2} \left( \frac{f(m_{W_i}, t_{W_i}) - f(m_{W_2}, t_{W_2})}{m_{W_1}^2} \right) - \cos^2(\theta_\nu) \frac{\sin(2\theta_\nu)}{2} \left( \frac{h(m_{W_i}, t_{W_i}) - h(m_{W_2}, t_{W_2})}{m_{W_1}^2} \right) \right]
\]

\[
g(m_{B,t_B}) = \frac{(m_{N_1} - m_\nu)}{4m_B^2(t_B - 1)^4} \left( 2 \log(t_B) \left( m_{N_1} m_\nu \left( t_B^2 + 1 \right) + m_B^2(t_B((5 - 3t_B)t_B - 16) + 8) \right) - 3(t_B - 1)^2 \left( m_{N_1} m_\nu(t_B + 1) + m_B^2((t_B - 5)t_B + 2) \right) \right)
\]

\[
f(m_{B,t_B}) = -\frac{m_{N_1} m_\nu \sqrt{t_B}((t_B - 2)t_B - 6 \log(t_B) + 13)}{6m_B(t_B - 1)^3}
\]

\[
h(m_{B,t_B}) = \frac{\sqrt{t_B}}{4m_B(t_B - 1)^4} \left( 2 \log(t_B) \left( t_B(t_B + 1) \left( m_{N_1}^2 + m_\nu^2 \right) + 2m_B^2(t_B - 1)(t_B(t_B - 1)(2t_B + 5) + 8) \right) - (t_B - 1)^2 \left( m_{N_1}^2(t_B + 1) \right) + m_\nu^2(t_B + 1) + 8m_B^2(t_B(t_B + 2) - 2) \right)
\]

and \( t_B = (m_\nu/m_B)^2 \).

FIG. 12: Heavy neutrino decay into a light one and a photon at one loop.

To give rise to the 3.55 keV line [15–17] from the decay of a 7.1 keV DM particle, one requires a lifetime of the order of \( 10^{28} \) s. For lightest right handed neutrino dark matter,
such a lifetime can be generated for the mixing angles $\theta_{LR}, \theta_\nu$ shown in the left panel plot of figure 13. While making the left panel plot of figure 13, the $W_R$ (or rather $W_2$) boson mass was fixed at 5 TeV, but the shape of the line does not change much even if we increase the mass to a high value, implying that the decay is happening mostly via left-right mixing. Although the decay lifetime is mostly governed by these mixing, one can related the mixing $\theta_{LR}$ with the $W_R$ mass using the expression relating them, given in (3). This is shown in the right panel plot of figure 13. This shows that, one requires $W_R$ mass beyond 1 PeV, for the correct mixing $\theta_{LR}$ required for long lifetime of fermion WDM. Such a preference for higher $W_R$ mass is similar to the one required for correct entropy dilution, discussed before.

![Graph 1](image1.png)

**FIG. 13:** Left scalar doublet DM annihilation into $W^+W^-$ final states for mass splitting 1 GeV and 10 GeV respectively, after incorporating the constraints from relic abundance criteria.

![Graph 2](image2.png)

**FIG. 14:** Scalar WDM decaying into two photons at one loop.
FIG. 15: Parameter space giving rise to a long lived scalar WDM decaying into two photons at one loop.

For scalar WDM on the other hand, such a decay can occur at one loop level through charged component of CDM doublet as seen from figure 14. The Decay of the neutral component of the triplet $\Omega_R$ to two photons are given as

$$\Gamma_{\Omega_R^0 \rightarrow \gamma \gamma} = \frac{\mu^2 e^4}{16\pi m_{\Omega_R^0}} |I|^2$$  \hspace{1cm} (49)
where

\[ |\mathcal{I}|^2 = \frac{1}{256\pi^4} \left( 4|A|^2 + \Re(A^*B) \right) \] (50)

\[ \mathcal{I}^{\mu\nu} = \frac{i}{16\pi^2} \left( A g^{\mu\nu} + \frac{k_1^{\mu} k_1^{\nu}}{m_\Omega^2} B \right) \] (51)

\[ A = \frac{i}{16\pi^2} \left( 1 + t \ln \left( \frac{2t - 1 + \sqrt{1 - 4t}}{2t} \right)^2 \right) \] (52)

\[ B = -2A \] (53)

with \( t = m_{\eta R}^2 / m_{\Omega R}^2 \). We fix the mass of \( \Omega_R^0 \) at 7.1 keV and vary the other two free parameters, namely \( \mu, m_{\eta R} \) from the requirement of the lifetime of \( \Omega_R^0 \) to be around \( 10^{28} \) s. The resulting parameter space is shown in figure 15. Such a small trilinear coupling between \( \Omega_R \) and \( \eta_R \) by invoking the presence of additional symmetries.

**VIII. CONCLUSION**

We have studied a class of left right symmetric models where the dark matter sector can consist of a keV scale warm component and a GeV-TeV scale cold component. Since both the DM components have gauge interactions, they can be thermally produced in the early Universe. While the cold component’s relic can be produced through the usual WIMP freeze-out mechanism, the warm component typically gets overproduced. This requires late time entropy dilution to bring the overproduced warm dark matter relic density to the observed or under-abundant regime. This requires a mother particle with a relatively long lifetime of 1 s, which we consider to be one of the right handed neutrinos.

While the minimal LRSM can accommodate cold DM component in a straightforward manner, there exists significant amount of fine tuning in order to have a keV DM component. We then mention another version of LRSM where the fermion masses can be generated through a common universal seesaw. Here, one can naturally achieve a keV scale fermion DM without much of a fine tuning. Another advantage of this model is the generation of tiny mixing between right and left handed neutrinos. This can happen either through a seesaw type term or at one loop if additional \( Z_2 \) symmetries are incorporated. If the lightest right handed neutrino is the keV DM candidate, such tiny mixing with the left handed neutrinos help it to acquire a long lifetime, required for a DM candidate. This also helps in entropy
dilution required to dilute the overproduced keV DM, by generating a long lifetime ($\sim 1$ s) of the decaying particle responsible for generating entropy.

We study three different cold dark matter candidates namely, a left scalar doublet, a right scalar doublet and a right fermion triplet and find the parameter space that can give rise to total as well as sub-dominant DM density. Since left fermion triplet relic abundance depends only on its mass (a fixed value around 3 TeV \cite{79}), we do not pursue it in this work. For other CDM candidates, we constrain the right sector gauge boson mass as well as DM-Higgs portal couplings from the relic abundance criteria. We then calculate the thermal relic abundance of keV fermion DM (lightest right handed neutrino), and the required entropy dilution due to the decay of the next to lightest right handed neutrino. This requires the right handed gauge boson mass to be a few tens of TeV. We also do the analysis for a keV scalar DM for the sake of completeness, though generating a keV scalar mass requires severe fine tuning.

We finally show the most interesting aspect of such a scenario that is, the indirect detection prospects of such mixed DM scenario. While the CDM component can have interesting signatures at gamma ray telescopes like the Fermi-LAT, the keV DM can give rise to monochromatic X-ray line if it decays on cosmological scales at radiative level. We find the parameter space for cold dark matter that can saturate the latest gamma ray bounds from experiments like the Fermi-LAT. We also find the relevant parameter space such that either a fermion or a scalar WDM with 7.1 keV mass can give rise to a monochromatic 3.55 keV X-ray line, as claimed to be present in the XMM-Newton telescope data.

It is interesting to observe that both the criteria of having the correct abundance of keV DM as well as producing the correct 3.55 keV X-ray line restricts the right sector gauge boson mass similarly. The similar bounds from the cold DM sector is somewhat lose as CDM has other couplings (like Higgs portal), independent of the right sector gauge bosons. Such high scale right sector gauge bosons are natural in LRSM with universal seesaw, once the lightest right handed neutrino mass is fixed to keV scale. Such a high scale right handed gauge sector, though out of reach from present collider experiments, can have interesting signatures at experiments like the IceCube \cite{46}.  

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