Common origin of baryon asymmetry, Dark Matter and neutrino mass

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ABSTRACT: In this work, we explain three beyond standard model (BSM) phenomena, namely neutrino masses, the baryon asymmetry of the Universe and Dark Matter, within a single model and in each explanation the right handed (RH) neutrinos play the prime role. Indeed by just introducing two RH neutrinos we can generate the neutrino masses by the Type-I seesaw mechanism. The baryon asymmetry of the Universe can arise from thermal leptogenesis from the decay of lightest RH neutrino before the decoupling of the electroweak sphaleron transitions, which redistribute the $B - L$ number into a baryon number. At the same time, the decay of the RH neutrino can produce the Dark Matter (DM) as an asymmetric Dark Matter component. The source of CP violation in the two sectors is exactly the same, related to the complex couplings of the neutrinos. By determining the comoving number density for different values of the CP violation in the DM sector, we obtain a particular value of the DM mass after satisfying the relic density bound. We also give prediction for the DM direct detection (DD) in the near future by different ongoing DD experiments.

KEYWORDS: Cosmology of Theories beyond the SM, Neutrino Physics

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1 Introduction

The Standard model (SM) is a concrete and successful theory which describes beautifully all the past and present particle physics measurements at colliders and at low energy experiments. After the recent discovery of the Higgs boson, all the particles of the SM have been detected and are fully consistent so far with the SM predictions. In spite of this terrific success, we know though that we still do not have a complete picture of particle physics. Indeed on one side we have clear evidence for neutrino masses, by observing the neutrino oscillation among the different flavours in atmospheric, solar and reactor neutrinos [1–11]. On the other hand, we know from astrophysical and cosmological data that in the Universe we need an additional matter component, apart for baryonic matter, to make up around 26% of the total energy budget. There are many evidences which support the existence of DM, mainly flatness of the galaxy rotation curves, the observation of bullet cluster and the CMB anisotropy [12–18]. In order to satisfy the observations DM has to be electrically neutral and stable or the decay life time has to be much longer than the age of the Universe (e.g. see ten point test of becoming the DM candidate in [19]). No particle of the SM has the right characteristics to provide the main Dark Matter component. Indeed neutrinos, which are neutral and stable, are unfortunately too light and cannot be the sole Dark Matter particle, see e.g. [20].

Another big puzzle is that there exist an excess of baryonic matter over anti-baryonic matter in the Universe and the baryon asymmetry has been measured very precisely by the satellite-based experiments WMAP and Planck [16–18]. It turns out that the baryonic matter makes up approximately 5% of the present energy density of the Universe and is approximately a factor 5 less abundant than Dark Matter. In order to generate the baryon asymmetry of the Universe, we need to satisfy the Sakharov conditions [21] including sufficient C and CP violation and deviation from thermal equilibrium. These conditions are difficult to realize within the SM and require generically new physics.

In this work, we mainly address the above mentioned three puzzles and try to solve them in a unified manner through the presence of RH handed-neutrinos, which mediate with the Dark Matter sector. Indeed the introduction of at least two RH neutrinos, singlets under the SM gauge group, allows to generate Majorana masses for the light neutrinos by the Type-I seesaw mechanism [22, 23]. Moreover, it is well known that RH neutrinos can also generate the baryon asymmetry of the Universe via leptogenesis [24, 25]. The lepton number asymmetry generated in the lightest RH neutrino decay can in fact be partially converted into a baryon asymmetry by sphaleron processes, which remain in thermal equilibrium until the electroweak phase transition [26].

Regarding the Dark Matter, the most popular ways to generate the appropriate density, like the freeze-out or freeze-in mechanisms, are usually independent from the neutrino sector and from baryogenesis. But in our work, we will instead follow the paradigm of asymmetric Dark Matter (ADM) also in order to explain the comparable densities of baryons and DM. Indeed if the RH neutrinos not only have a Yukawa coupling with the light neutrinos and the SM Higgs, but also couple with the Dark Matter sector, they can decay also in the Dark Sector generating an asymmetry of a similar order. Many models of asymmetric Dark Matter have been proposed in the literature [27–52], but in this work we will give a new realization of the scenario and exploit and explore more in depth the connection to neutrino physics.

In order to solve the above mentioned problems in a common way, we extend the SM both in the particle content as well as in the gauge group structure. We add to the gauge group a local $SU(2)_D$ interaction¹ in the Dark sector and discrete $\mathbb{Z}_3$, $\mathbb{Z}_2$ groups, ensuring the Dark Matter stability as well as forbidding the Majorana mass terms among the dark sector particles and limiting the number of model parameters. The particle list is also enlarged to include two dark sector fermionic left handed doublets and their RH counterparts, singlet under the dark $SU(2)_D$, two scalar doublets of $SU(2)_D$ and two singlet RH neutrinos. We must introduce two $SU(2)_D$ fermionic doublets in order to cancel the Witten anomaly [54]. As we will see in the result section, the RH neutrinos take part in generating the lepton asymmetry, Dark Matter asymmetry and neutrino mass. One of the scalar doublets takes a vacuum expectation value, generating a mass for the exotic fermions and also mixing with the SM Higgs. This opens up the possibility to have DM scatterings mediated by the Higgs fields, which may be detected in different ongoing and proposed direct detection experiments [55–61]. The other exotic scalar doublet does not

¹WIMP type DM obeying $SU(2)$ gauge symmetry has been studied earlier in [53].

obtain a non-zero vacuum expectation value and plays a role similar to the inert doublet, participating in our case to the DM production.

The paper is organised in the following way. In section 2, we describe all the details of our model. The generation of the neutrino mass is discussed in section 3. In section 4, we give the Boltzmann equations for both the dark sector and the leptonic sector, while section 5, contains the full numerical result of the Boltzmann equations. In section 6, the DM direct detection is addressed. Finally, in section 7, we conclude our work with an outlook on possible signature at colliders.

2 Model

In this article, we consider a hidden sector which has a local $SU(2)_D$ gauge invariance. In this hidden sector we introduce four fermions ψ_i ($i = 1$ to 4) whose left handed components transform like a doublet under $SU(2)_D$ while the right handed counterparts are $SU(2)_D$ singlets. Therefore, we have a pair of $SU(2)_D$ doublets $\Psi_{\alpha L}$ ($\alpha = 1, 2$). In table 1, we show the complete list of particles in the present model. Here we want to point that, since we have an even number of fermionic $SU(2)_D$ doublets, our model is free from the Witten anomaly [54]. In addition, we have two scalar doublets in the hidden sector as well. One of the scalar doublets η_D does not get any vacuum expectation value (VEV) while the remaining one (ϕ_D) has a nonzero VEV and thus mixes with the SM Higgs doublet ϕ_h . Moreover, in order to have a stable DM candidate, we also impose a discrete $\mathbb{Z}_3$ symmetry, and keep all the hidden sector fermions as well as the inert doublet charged under $\mathbb{Z}_3$. These symmetries allow Majorana mass terms among the extra $SU(2)_D$ singlet fermions, which after the breaking of the dark symmetry could switch on the conversion of the DM to anti-DM. Therefore, to be on the safe side we introduce an additional $\mathbb{Z}_2$ symmetry to forbid the Majorana mass terms among the extra fermions and reduce the possible couplings of the second fermionic state. Under $\mathbb{Z}_2$, ψ_{2L} , ψ_{3R} and ψ_{4R} are odd and the rest of the particles including the Standard Model particles are even. Furthermore, we have two right handed (RH) fermions N_i ($i = 1, 2$), singlets under both SM gauge group as well as $SU(2)_D$. These singlet fermions play the role of the RH neutrino and are the only *connector* between the visible and the hidden sector, as long as the electroweak and dark $SU(2)_D$ are unbroken and the mixing in the scalar sector vanishes. In this sense our model is a special case in the class of neutrino(+Higgs) portal models [62–67].

2.1 Particle spectrum

2.2 Lagrangian

The $SU(3)_c \times SU(2)_L \times SU(2)_D \times U(1)_Y \times \mathbb{Z}_3$ invariant Lagrangian for our present model is given by,

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + i \bar{\Psi}_k \gamma^\mu D_\mu^k \Psi_k + (D_\mu^D \phi_D)^\dagger (D^{D\mu} \phi_D) + (D_\mu^D \eta_D)^\dagger (D^{D\mu} \eta_D) + \left(y_{ij} \bar{L}_i \tilde{\phi}_h N_{jR} + h.c. \right) \\ & + \left(\lambda_1 \bar{\Psi}_{1L} \tilde{\phi}_D \psi_{1R} + \lambda_2 \bar{\Psi}_{1L} \phi_D \psi_{2R} + \lambda_3 \bar{\Psi}_{2L} \tilde{\phi}_D \psi_{3R} + \lambda_4 \bar{\Psi}_{2L} \phi_D \psi_{4R} + h.c. \right) \\ & - \alpha_j \bar{\Psi}_{1L} \eta_D N_{jR} + i \bar{N}_{jR} \not{D} N_{jR} - M_j \bar{N}_{jR}^c N_{jR} - \mathcal{V}(\phi_h, \phi_D, \eta_D), \end{aligned} \quad (2.1)$$

Gauge Group	Fermion Fields							Scalar Fields		
	$\Psi_{1L}=(\psi_1, \psi_2)_L^T$	ψ_{1R}	ψ_{2R}	$\Psi_{2L}=(\psi_3, \psi_4)_L^T$	ψ_{3R}	ψ_{4R}	N_i	ϕ_h	ϕ_D	η_D
$SU(3)_c$	1	1	1	1	1	1	1	1	1	1
$SU(2)_L$	1	1	1	1	1	1	1	2	1	1
$SU(2)_D$	2	1	1	2	1	1	1	1	2	2
$\mathbb{Z}_3 \times \mathbb{Z}_2$	$(\omega, 1)$	$(\omega, 1)$	$(\omega, 1)$	$(\omega^2, -1)$	$(\omega^2, -1)$	$(\omega^2, -1)$	$(1, 1)$	$(1, 1)$	$(1, 1)$	$(\omega, 1)$

Table 1. List of hidden sector particles and *connector* particles and their corresponding charges under various symmetry groups. All the particles listed above have zero hypercharge except SM Higgs doublet ϕ_h which has hypercharge $Y = 1/2$.

where $i = 1$ to 3 while j runs from 1 to 2 and D_μ^D is the covariant derivative for $SU(2)_D$. We assign the $\mathbb{Z}_3$ charges to the different fields in such a way so that the lightest component of the dark doublet Ψ_{1L} becomes stable and a viable DM candidate. Moreover, in the present work we are interested in ψ_1 ,² production from the decays of RH-neutrinos. Note that we can consider this state to have all real couplings by absorbing all the phases in $\lambda_{1,2}$, in the RH states $\psi_{1,2R}$, while the phases of $\alpha_{1,2}$ can be absorbed into the doublets Ψ_{1L} and η_D . Without loss of generality we can also consider the heavy Majorana masses $M_{1,2}$ to be real and positive, but redefining accordingly the heavy RH neutrino fields $N_{1,2R}$.

In the model there is also a second state, possibly ψ_3 , which is stable due to the Z_2 and Z_3 symmetries and could contribute to the Dark Matter density. To avoid both a substantial ψ_3 freeze-out density and a symmetric DM component from ψ_1 , the presence of the $SU(2)_D$ gauge symmetry is crucial since it allows for an efficient annihilation of the fermions, as long as they are not too heavy. We will discuss the precise value of the allowed mass range later on. In this way, the main component of the DM in the Universe will be generated by the DM asymmetry in the RH neutrino's decays.

Given the symmetries discussed above, the gauge invariant scalar potential has the following form

$$\begin{aligned}
 \mathcal{V}(\phi_h, \phi_D, \eta_D) = & -\mu_h^2(\phi_h^\dagger \phi_h) + \lambda_h(\phi_h^\dagger \phi_h)^2 - \mu_D^2(\phi_D^\dagger \phi_D) + \lambda_D(\phi_D^\dagger \phi_D)^2 \\
 & + \mu_\eta^2(\eta_D^\dagger \eta_D) + \lambda_\eta(\eta_D^\dagger \eta_D)^2 + \lambda_{hD}(\phi_h^\dagger \phi_h)(\phi_D^\dagger \phi_D) + \lambda_{h\eta}(\phi_h^\dagger \phi_h)(\eta_D^\dagger \eta_D) \\
 & + \lambda_{D1}(\phi_D^\dagger \phi_D)(\eta_D^\dagger \eta_D) + \lambda_{D2}(\phi_D^\dagger \eta_D)(\eta_D^\dagger \phi_D) + \lambda_{D3}(\phi_D \eta_D^3 + h.c.). \quad (2.2)
 \end{aligned}$$

Here, both the doublets ϕ_h and ϕ_D acquire VEVs and generate masses to SM particles and the hidden sector fermions after spontaneous breaking of $SU(2)_L \times U(1)_Y$ and $SU(2)_D$ symmetries, respectively. In the unitary gauge, scalar doublets ϕ_h and ϕ_D take the following form after symmetry breaking,

$$\phi_h = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix}, \quad \phi_D = \begin{pmatrix} 0 \\ \frac{v_D+H}{\sqrt{2}} \end{pmatrix}. \quad (2.3)$$

²Like the SM lepton doublets, in dark sector too, we have assumed that between the components of a dark doublet the component with isospin $+1/2$ is the lightest one.

In the scalar potential there is a mixing term between the CP even neutral components of the two doublets ϕ_h and ϕ_D , hence the gauge basis and mass eigenbasis will be different. In h, H basis the mass square mixing matrix will be as follows

$$M_{scalar}^2 = \begin{pmatrix} 2\lambda_h v^2 & \lambda_{hD} v v_D \\ \lambda_{hD} v v_D & 2\lambda_D v_D^2 \end{pmatrix}. \quad (2.4)$$

After diagonalising the above mass matrix we get the physical masses and the corresponding physical states which are linear combinations of gauge basis in the following manner

$$\begin{aligned} h_1 &= h \cos \zeta - H \sin \zeta, \\ h_2 &= h \sin \zeta + H \cos \zeta, \end{aligned} \quad (2.5)$$

where ζ is the mixing angle between h_1, h_2 and the mixing angle can be expressed in terms of the Lagrangian parameters in following way

$$\tan 2\zeta = \frac{\lambda_{hD} v v_D}{\lambda_D v_D^2 - \lambda_h v^2}. \quad (2.6)$$

As mentioned above, after diagonalising the scalar mass matrix in eq. (2.4), we get the physical masses for the two neutral scalars as

$$\begin{aligned} M_{h_1}^2 &= \lambda_h v^2 + \lambda_D v_D^2 - \sqrt{(\lambda_D v_D^2 - \lambda_h v^2)^2 + (\lambda_{hD} v v_D)^2}, \\ M_{h_2}^2 &= \lambda_h v^2 + \lambda_D v_D^2 + \sqrt{(\lambda_D v_D^2 - \lambda_h v^2)^2 + (\lambda_{hD} v v_D)^2}. \end{aligned} \quad (2.7)$$

We identify the lighter Higgs scalar as the SM-like Higgs observed at the LHC. Therefore, we take $M_{h_1} = 126 \text{ GeV}$ and consider small mixing angle $\sin \zeta \leq 0.1$ in order to ensure agreement with the Higgs signal strengths at the LHC [68–70].

Now, we can express all the quartic couplings in terms of the physical Higgs masses as,

$$\begin{aligned} \lambda_D &= \frac{M_{h_2}^2 + M_{h_1}^2 + (M_{h_2}^2 - M_{h_1}^2) \cos 2\zeta}{4v_D^2}, \\ \lambda_h &= \frac{M_{h_2}^2 + M_{h_1}^2 - (M_{h_2}^2 - M_{h_1}^2) \cos 2\zeta}{4v_D^2}, \\ \lambda_{hD} &= \frac{(M_{h_2}^2 - M_{h_1}^2) \sin 2\zeta}{2v v_D}, \\ \mu_h^2 &= \lambda_h v^2 + \lambda_{hD} \frac{v_D^2}{2}, \\ \mu_D^2 &= \lambda_D v_D^2 + \lambda_{hD} \frac{v^2}{2}. \end{aligned} \quad (2.8)$$

In the above expressions, all the quartic couplings have to be within the perturbative regime which is $\lambda_i < 4\pi$.

After the $SU(2)_D$ symmetry breaking, the DM candidate $\psi_1 (= \psi_{1L} \oplus \psi_{1R})$ will get mass which is

$$M_{\psi_1} = M_{\text{DM}} = \frac{\lambda_1 v_D}{\sqrt{2}}. \quad (2.9)$$

The other scalar doublet η_D , which has a nonzero $\mathbb{Z}_3$ charge, will also get mass after breaking of both SM and hidden sector gauge symmetries. In the present model, among the $\mathbb{Z}_3$ charged particles i.e. hidden sector fermions and scalar η_D , we consider the fermion ψ_1 as the lightest one. This is always possible by tuning the couplings related to η_D and λ_α s ($\alpha = 2$ to 4) so that heavier $\mathbb{Z}_3$ charged particles decay to the lightest one and thereby ψ_1 becomes stable. Thus, ψ_1 will be a viable DM candidate in our model. Moreover, the invariance of the $\mathbb{Z}_3$ and $\mathbb{Z}_2$ symmetry actually make the hidden sector fermion mass matrix diagonal, which means unlike the quark or neutrino mixing in the SM, there is no mixing between the fermions in different $SU(2)_D$ doublets.

3 Neutrino mass

In this work, as mentioned in the Model section (section 2), instead of three right handed neutrinos we consider a minimal setting and we add only two, which is sufficient to explain the current neutrino oscillation data. Therefore, in this framework, light neutrino masses are generated by the well known Type-I seesaw mechanism, where the light neutrino mass matrix is related to the Dirac and Majorana mass matrices in the following way

$$m_\nu = -M_D M_R^{-1} M_D^T, \quad (3.1)$$

where M_D is the Dirac mass matrix (3×2) while the Majorana mass matrix for the heavy right handed neutrinos N_{jR} ($j = 1, 2$) is denoted by a 2×2 matrix M_R . In this work, for simplicity and without loss of generality, we assume M_R to be a diagonal matrix i.e. $M_R = \text{diag}(M_{N_1}, M_{N_2})$. One can also choose M_{N_1} , M_{N_2} real and positive by redefining the phases of the spinors N_1 and N_2 in the mass eigenstate basis. Similarly, by redefining the phases of the left handed neutrinos in the flavour basis, we can remove the phases of one entire column of M_D matrix. So we consider all the elements of the first column of the Dirac mass matrix (M_D) as real. Therefore, in matrix form it looks like as follows,

$$M_D = \frac{y_{ij} v}{\sqrt{2}} = \frac{v}{\sqrt{2}} \begin{pmatrix} y_{ee} & y_{e\mu}^R - iy_{e\mu}^I \\ y_{\mu e} & y_{\mu\mu}^R - iy_{\mu\mu}^I \\ y_{\tau e} & y_{\tau\mu}^R - iy_{\tau\mu}^I \end{pmatrix}. \quad (3.2)$$

We have computed the physical masses of light neutrinos (eigenvalues) and mixing angles (eigenvectors) by diagonalising the complex symmetric Majorana mass matrix m_ν . All the elements of m_ν are explicitly given in the appendix A. In studying the neutrino phenomenology, we take into account the observed values of neutrino oscillation parameters [71]. In this work, we mostly focus on the normal hierarchy (NH) of the light neutrino masses, but a similar study can be done for the inverted hierarchical scenario as well. Bounds which we have considered to constrain the elements of m_ν matrix are two mass square differences,

the three oscillation angles and the cosmological upper bound on the sum of three light neutrino masses. These bounds are as follows

- from the neutrino oscillation experiments there are tight constraints on two mass square differences. For NH 3σ bounds are as follows [71],

$$6.93 \leq \frac{\Delta m_{21}^2}{10^{-5}} (\text{eV}^2) \leq 7.97 \quad \text{and} \quad 2.37 \leq \frac{\Delta m_{31}^2}{10^{-3}} (\text{eV}^2) \leq 2.63. \quad (3.3)$$

- Three mixing angles namely, solar mixing angle θ_{12} , atmospheric mixing angle θ_{23} and reactor mixing angle θ_{13} are also very well measured now. The allowed values of mixing angles for NH in 3σ are listed below [71]

$$\begin{aligned} 30^\circ &\leq \theta_{12} \leq 36.51^\circ, \\ 7.82^\circ &\leq \theta_{13} \leq 9.02^\circ, \\ 37.99^\circ &\leq \theta_{23} \leq 51.71^\circ, \end{aligned} \quad (3.4)$$

- From the Planck data [17] there is a bound on the sum of light neutrinos masses which is $\sum_{i=1,2,3} m_{\nu_i} \leq 0.23 \text{ eV}$.

Moreover, we have not applied any bound on the CP violating phase δ , which is yet to be measured accurately by the different ongoing and upcoming oscillation experiments like T2K [72], T2HK [73], DUNE [74–77] and INO [78]. There is a hint of maximal CP violation ($\delta \sim -\frac{\pi}{2}$) from the T2K experiment [79], which excludes the CP conserving values of $\delta = 0$ or π , at 90% C.L. However, there is some tension between the observed value of δ from T2K and NOvA [80].

In satisfying the above mentioned neutrino oscillation parameters, we vary the model parameters, specifically, the elements of the Dirac mass matrix (eq. (3.2)) and the RH-neutrino masses in the following range,

$$\begin{aligned} 10^{-4} &\leq y_{ee}, y_{\mu e}, y_{\tau e}, y_{e\mu}^R \leq 10^{-2}, \\ 10^{-3} &\leq y_{\mu\mu}^R, y_{\mu\mu}^I, y_{\tau\mu}^R, y_{\tau\mu}^I \leq 1, \\ 10^{-6} &\leq y_{e\mu}^I \leq 10^{-4}, \\ 10^8 (10^{13}) &\leq M_{N_1} (M_{N_2}) \text{ GeV} \leq 10^{11} (10^{15}). \end{aligned} \quad (3.5)$$

In eq. (3.2), all the elements in the first column are real and positive (can be made by phase rotation) while the elements in the second column can have both positive as well as negative values. However, to make the CP asymmetry parameter positive we consider only positive values of all the elements in the second column of the Dirac mass matrix, i.e. in the above $y_{\alpha\mu}^{R,I} = |y_{\alpha\mu}^{R,I}|$, ($\alpha = e, \mu, \tau$). Here, we take hierarchical RH-neutrino masses, which is evident from the ranges of M_{N_1} and M_{N_2} that we have chosen in eq. (3.5). The hierarchical scenario for the RH-neutrino masses implies that we have to consider the lepton as well as DM asymmetry generated from the decay of lightest RH-neutrino (N_1) only. Indeed the asymmetry generated from the decay of N_2 will be washed out by the decay as well as

inverse decay of N_1 which is still in thermal equilibrium during $T \sim M_{N_2}$. In the result section (section 5), we will see that there exist correlations among the parameters of the neutrino sector and hence not the entire adopted ranges in eq. (3.5) are allowed by the neutrino oscillation data and the requirement of successful leptogenesis.

4 Boltzmann equation for studying lepton asymmetry and Dark Matter asymmetry

As we have already mentioned earlier, in this work our goal is to generate an asymmetry in the dark sector following the idea of leptogenesis in the visible sector. In other words, the asymmetries in both sectors may have a common source i.e. they can be generated from the CP-violating out-of-equilibrium decay of the lightest RH-neutrino N_1 . Therefore, we need to solve (at least) three Boltzmann equations simultaneously: one gives the number density for the lightest RH-neutrino N_1 and the other two will govern the asymmetries of the visible and dark sectors. As mentioned in [29], depending on the decay width and the mass of the RH neutrinos, the source of dark and visible sector asymmetry production can be divided in to two regimes. In the case $\Gamma_{N_1} \ll M_{N_1}$, we are in the narrow width approximation and the RH neutrino couples very weakly to the thermal bath, so that it is strongly out of equilibrium. On the other hand, for $\Gamma_{N_1} \simeq M_{N_1}$, we are in the large washout or transfer regime where the wash-out processes mediated by N_1 are not negligible. In the current work, we mostly focus on the RH-neutrino mass around 10^9 GeV and from the light neutrino mass constraint which is $\sim 10^{-11}$ GeV we have very small Yukawa couplings in the Dirac mass, hence we are in the narrow width approximation regime with $\Gamma_{N_1} \ll M_{N_1}$. Moreover, for this assumption we neglect the transfer diagrams [29].

The relevant Boltzmann equations which we need to solve for the narrow width approximation regime are as follows,

$$\frac{dY_{N_1}}{dz} = -\frac{z}{sH(M_1)} \left[\left(\frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \times (\gamma_{D_1} + 2\gamma_{\phi,s}^1 + 4\gamma_{\phi,t}^1) \right], \quad (4.1)$$

$$\frac{dY_{\Delta l}}{dz} = -\frac{\Gamma_{N_1}}{H(M_1)} \left[\epsilon_l \frac{zK_1(z)}{K_2(z)} (Y_{N_1}^{eq} - Y_{N_1}) + \text{Br}_l \frac{z^3 K_1(z)}{4} Y_{\Delta l} \right], \quad (4.2)$$

$$\frac{dY_D}{dz} = -\frac{\Gamma_{N_1}}{H(M_1)} \left[\epsilon_D \frac{zK_1(z)}{K_2(z)} (Y_{N_1}^{eq} - Y_{N_1}) + \text{Br}_D \frac{z^3 K_1(z)}{4} Y_D \right], \quad (4.3)$$

where the first equation represents the evolution of the comoving yield Y_{N_1} of N_1 . The yield of a species is defined as the actual number density of that species divided by the entropy density of the Universe. If there is no interaction then the yield of a species remains unaltered, as the expansion of the Universe dilutes the number density and the entropy density in the same way.

The R.H.S. of the Boltzmann equation for N_1 describes the possible ways to change the number density of N_1 . The quantity γ_{D_1} is related to the total decay width of N_1 i.e. the decay of N_1 into both visible and dark sectors. On the other hand, $\gamma_{\phi,s}^1$ and $\gamma_{\phi,t}^1$ are related to s -channel and t -channel scattering of N_1 mediated by ϕ_h , which can also lead to the destruction or production of N_1 . The expressions of γ_{D_1} , $\gamma_{\phi_h,s}^1$ and $\gamma_{\phi_h,t}^1$ are given in

the appendix B. The second and third equations are the evolution equations for the lepton asymmetry and Dark Matter asymmetry, respectively. The first term in the R.H.S. of eq. (4.2) (eq. (4.3)) is the source term of lepton (Dark Matter) asymmetry from N_1 decay, while the second term represents the washout effects on the created lepton (Dark Matter) asymmetry due to the inverse decays of N_1 . In eqs. (4.2), (4.3), Br_l and Br_D are the branching ratios of RH neutrino N_1 decay to leptonic sector and dark sector, respectively. In the above equations the CP asymmetry parameters $\epsilon_{l,D}$ are zero at tree level. However, by considering both tree level and one loop level diagrams (vertex correction and wave function correction, see figure 2 of [82]), non-zero values for the CP-asymmetry parameters $\epsilon_{l,D}$ arise due to the interference between tree level and one loop level diagrams. The CP asymmetry parameter for the visible sector is defined as [25, 29]

$$\begin{aligned}\epsilon_l &= \frac{\Gamma(N_1 \rightarrow L\phi_h) - \Gamma(N_1 \rightarrow \bar{L}\phi_h^\dagger)}{\Gamma_{N_1}}, \\ &= \frac{M_{N_1}}{16\pi M_{N_2}} \frac{\text{Im} \left[3((y^\dagger y)_{12}^*)^2 + 2\alpha_1^* \alpha_2 (y^\dagger y)_{12}^* \right]}{[(y^\dagger y)_{11} + \alpha_1 \alpha_1^*]},\end{aligned}\quad (4.4)$$

where we have normalized to the total RH neutrino decay rate and summed over the lepton flavours, i.e. $\epsilon_l = \sum_\alpha \epsilon_l^\alpha$. In eq. (4.4) we include contributions from the vertex and wave-function diagrams with virtual SM states as in classic leptogenesis, see e.g. [25], and also the contribution from the wave-function diagram with virtual dark sector states.

Similarly, the CP asymmetry in dark sector is defined as [29]

$$\begin{aligned}\epsilon_D &= \frac{\Gamma(N_1 \rightarrow \psi_{1L} \eta_D) - \Gamma(N_1 \rightarrow \overline{\psi_{1L}} \eta_D^\dagger)}{\Gamma_{N_1}}, \\ &= \frac{M_{N_1}}{16\pi M_{N_2}} \frac{\text{Im} \left[2\alpha_1^* \alpha_2 (y^\dagger y)_{12}^* + 3(\alpha_1^* \alpha_2)^2 \right]}{[(y^\dagger y)_{11} + \alpha_1 \alpha_1^*]}.\end{aligned}\quad (4.5)$$

In the case of Dark Matter, we have in an analogous way included contributions from the vertex and wave-function diagrams from the dark sector and only the wave-function diagram from the leptons. The total decay width of the RH-neutrino N_1 is given by

$$\Gamma_{N_1} = \frac{M_{N_1}}{8\pi} \left[(y^\dagger y)_{11} + |\alpha_1|^2 \right]. \quad (4.6)$$

From eqs. (4.4), (4.5) we see that both ϵ_l and ϵ_D are determined by the Yukawa couplings y_{ij} and RH-neutrino masses. One very important thing to stress again is that in the dark sector we can absorb the phases of the couplings α_j ($j = 1, 2$) by redefining the phases of ψ_{1L} and the complex scalar doublet η_D . Indeed, as in the case of a single generation of fermions, no CP phase is physical in the DM sector. Therefore the only source of CP violation for both the leptonic and DM sectors are the imaginary parts of the lepton-neutrino Yukawa couplings. We have then

$$\frac{\epsilon_l}{\epsilon_D} = \frac{\text{Im} \left[3((y^\dagger y)_{12}^*)^2 + 2\alpha_1 \alpha_2 (y^\dagger y)_{12}^* \right]}{2\alpha_1 \alpha_2 \text{Im} [(y^\dagger y)_{12}^*]} = 1 + \frac{\text{Im} \left[3((y^\dagger y)_{12}^*)^2 \right]}{2\alpha_1 \alpha_2 \text{Im} [(y^\dagger y)_{12}^*]} = 1 + 3 \frac{\sum_\beta y_{\beta e} y_{\beta \mu}^R}{\alpha_1 \alpha_2}. \quad (4.7)$$

We see that in the particular case when $(y^\dagger y)_{12}^*$ is purely imaginary, i.e. $y^R = 0$, or generically when $\alpha_1 \alpha_2 \gg \sum_\beta y_{\beta e} y_{\beta \mu}^R$, the two CP violation parameters are equal and we can expect a similar asymmetry in the two sectors, as long as the wash-out processes are negligible. This is indeed not difficult to achieve as the Yukawas connected to the first generation of the SM $y_{\beta e}$ have to be small to give the small mixing angle θ_{13} . From the matrix in eq. (3.2), we have that

$$(y^\dagger y)_{12}^* = y_{ee}(y_{e\mu}^R + iy_{e\mu}^I) + y_{\mu e}(y_{\mu\mu}^R + iy_{\mu\mu}^I) + y_{\tau e}(y_{\tau\mu}^R + iy_{\tau\mu}^I) \quad (4.8)$$

so that this quantity is purely imaginary when the second column of the Dirac mass matrix is purely imaginary and only six real Yukawa parameters remain. Note that in this limit, the CP violation in the RH neutrino decay can still be large, while the light neutrino mass matrix is real and the Dirac phase is therefore vanishing. We have checked that with only imaginary components in the second column of the Dirac mass matrix (see eq. (3.2)) and also in the same range of the parameters value as given in eq. (3.5), one can easily obtain the three neutrino mixing angles and the two mass square differences in their observed 3σ ranges. For this particular choice of parameters, one can easily estimate the value of the lepton CP asymmetry parameter (ϵ_l). In this case ϵ_l takes the following form,

$$\epsilon_l = \frac{M_{N_1}}{8\pi M_{N_2}} \frac{\alpha_2}{\alpha_1} \text{Im} \left[(y^\dagger y)_{12}^* \right] \quad (4.9)$$

for large α_2/α_1 and $(y^\dagger y)_{11} < \alpha_1^2$. If we take $\text{Im}(y^\dagger y)_{12}^* \sim 2 \times 10^{-3}$, $M_{N_1}/M_{N_2} \sim 10^{-4}$ and $\alpha_2/\alpha_1 \sim 10^2$, then we obtain $\epsilon_l \sim 10^{-6}$ sufficient to generate the lepton asymmetry of the Universe, as we will show later. One important conclusion we can draw for this scenario is that although there is no low scale CP violation,³ there still exist a sufficiently large high scale CP violation by which we can generate lepton asymmetry of the Universe and henceforth the observed baryon asymmetry of the Universe.

Another limiting case is when the real and imaginary parts $y_{i\mu}^{R,I}$ in the second column of the Yukawa matrix are equal and large. In that case it is $((y^\dagger y)_{12}^*)^2$ which becomes purely imaginary and can even dominate the CP violation in the leptonic sector. In that case the dark CP violation parameter ϵ_D can be substantially smaller, but we can compensate the smaller number density of the Dark Matter by increasing its mass and still satisfy the CMB constraint. Moreover, note that we can take large value of α_2 ($\sim \mathcal{O}(1)$) without violating any constraint, but a large value of α_1 will violate the narrow width approximation ($\Gamma_{N_1} \ll M_{N_1}$) via eq. (4.6). Therefore, remaining within the narrow width approximation, but at the same time aiming to increase the CP asymmetry parameter ϵ_l for the production of sufficient lepton asymmetry, we vary the two parameters α_1 and α_2 in the following ranges:

$$\begin{aligned} 10^{-3} &\leq \alpha_1 \leq 10^{-2}, \\ 0.3 &\leq \alpha_2 \leq 1.0. \end{aligned} \quad (4.10)$$

³If the light neutrino mass matrix is real, the Dirac phase δ is vanishing, but the Majorana phases can still be non-trivial if the eigenvalues of the mass matrix have different sign. This indeed happens if one column of the Yukawa matrix is imaginary, see the discussion in appendix A.

By solving the above three Boltzmann equations, we obtain the lepton and DM asymmetries as a function of the temperature of the Universe. For leptogenesis, we need to calculate the lepton asymmetry before the sphaleron decoupling temperature $T_{sph} \sim 150 \text{ GeV}$ [81, 82] because the produced lepton asymmetry has to be converted into the observed baryon asymmetry via sphaleron transitions in equilibrium. The conversion factor from $Y_{\Delta l}$ to Y_B is given by [83]

$$Y_B = \frac{8N_f + 4N_{\phi_h}}{22N_f + 13N_{\phi_h}} Y_{\Delta l}, \quad (4.11)$$

where N_f is number of generations of quarks and leptons while N_{ϕ_h} being number of scalar doublets in the visible sector.

Regarding the Dark Matter, we need to ensure that the symmetric DM component efficiently annihilates away leaving only the asymmetric component. As the Dark Matter is charged under a Dark SU(2) gauge group, this is not difficult to achieve as long as the Dark Matter mass is not too large. Indeed the dominant annihilation channel for an SU(2) doublet is into the corresponding gauge bosons, that subsequently annihilate/decay through the scalar sector into SM states.⁴ The generic expectation for the annihilation rate of a doublet is similar to that of the Higgsino in supersymmetric models and it is enhanced by coannihilations [84] and even the Sommerfeld effect. In the case of Dark Matter mass of similar order to the gauge boson mass and $\lambda_2 \gtrsim \lambda_1$, the coannihilation effect is dominant and we obtain [85]

$$\sigma_{\psi\bar{\psi} \rightarrow W_D W_D} \sim \frac{g_D^4}{128\pi M_{DM}^2} = 2.5 \times 10^{-9} \text{ GeV}^{-2} g_D^4 \left(\frac{M_{DM}}{1 \text{ TeV}} \right)^{-2} \quad (4.12)$$

which is slightly larger than the thermal cross section $\langle \sigma v \rangle \sim 10^{-9} \text{ GeV}^{-2}$. So for a Dark Gauge coupling of order one, the symmetric Dark Matter component becomes important only at masses above 1 TeV. Even heavier masses can be allowed if the gauge bosons remain light and the Sommerfeld effect increases the cross section further [86].

For masses below the TeV range, where the annihilation is strong enough, we can determine the Dark Matter relic density by computing Dark Matter asymmetry at the present epoch $T = T_0 \sim \mathcal{O}(10^{-13}) \text{ GeV}$ ($z \rightarrow \infty$) and using the following relation [84],

$$\Omega h^2 = 2.755 \times 10^8 \left(\frac{M_{DM}}{\text{GeV}} \right) Y_D(z \rightarrow \infty). \quad (4.13)$$

Earlier WMAP and now the Planck satellite have measured the DM relic density very precisely and its present value at the 68% C.L. is [17]

$$0.1172 \leq \Omega h^2 \leq 0.1226. \quad (4.14)$$

Equating eq. (4.13) and eq. (4.14) one can find the mass range for the Dark Matter particle M_{DM} which reproduces the observed Dark Matter abundance.

⁴In the same way, the symmetric component of ψ_3 annihilates away.

Figure 1. Variation of baryon asymmetry of the Universe and the dark sector asymmetry with z . In generating the figure we have assumed the following values $M_{N_1} = 10^{10}$ GeV, $M_{N_2} = 10^{13}$ GeV, $y^\dagger y = 10^{-6}$, $\Gamma_{N_1} \sim 398$ GeV, $Br_\chi = 0.2$, $Br_l = 0.8$.

Let us finally comment on the presence of the inert doublet η_D , which is also produced in the RH neutrino decay and, as it is heavier, could regenerate a ψ_1 population after freeze-out and after the $SU(2)_D, SU(2)_{EW}$ gauge symmetries are broken, by decaying into a DM fermion and a lepton. The number density of η_D can be efficiently reduced by the cubic interaction with ϕ_D , which allows for semiannihilation [87] through the process $\eta_D + \eta_D \rightarrow \eta_D^* + h_{1,2}$. In this case semiannihilation is determined by the potential parameter λ_{D3} , which does not affect the stability of the vacuum and can be chosen large to allow for a sufficiently large semiannihilation cross section. We assume here that due to such process only a negligible number density of the inert doublet is left after freeze-out.

5 Results

In figure 1, we show the production of the baryon asymmetry and the Dark Matter asymmetry from the decay of RH neutrino N_1 . When the RH neutrino starts decaying, both the lepton sector asymmetry and the dark sector asymmetry grow fast to their final values. For $\epsilon_l = 4 \times 10^{-7}$ and taking the other parameters values as mentioned in the caption, we obtain the correct value of the matter antimatter asymmetry of the Universe which lies within the value measured by the Planck satellite. For two values of $\epsilon_D = 3.5 \times 10^{-7}$ and 3.5×10^{-9} , we can provide for a dark sector asymmetry which fulfils the total Dark Matter

Figure 2. Left panel (LP) shows the variation in the $Y_D - \epsilon_D$ plane and in the right panel (RP) variation in $(Y_B - \epsilon_l)$ plane has been shown. All the points satisfy neutrino oscillation data as mentioned in section 3. All the parameters have been varied in the range as shown in section 3.

abundance for a Dark Matter mass of 0.76 GeV and 76 GeV, respectively. One interesting thing to note here is that for $\epsilon_D = 3.5 \times 10^{-7}$ (which is less than ϵ_l) we are producing more dark asymmetry than the lepton asymmetry. This is because we have considered $Br_D < Br_l$ *i.e.* washout effects for the dark sector are weaker than for the visible sector. So even for CP violation of the same order, we can obtain an enhancement of the Dark Matter density compared to the baryon density giving naturally $\Omega_{DM} > \Omega_b$ for DM masses in the range of the charm/bottom quark mass. In the subsequent figures we will see the dependence of the baryon and dark sector asymmetry on the model parameters and how they are related with neutrino oscillations.

In the LP of figure 2, we show the variation of the Dark Matter yield with ϵ_D . All the points satisfy the neutrino oscillation data. The figure shows a sharp correlation between ϵ_D and Y_D , as expected in the narrow width regime. For the particular value of the yield we determine the Dark Matter mass such that the DM energy density coincides with the observed value using the expression as given in eq. (4.13). We see that the present model can accommodate the right Dark Matter abundance in the mass range $M_{DM} = 1 - 10^6$ GeV, which is, at least in the lower mass range, within the sensitivity of ongoing direct detection experiments. The large mass region above the TeV is disfavoured by the fact that the CP violation parameter has to be very suppressed and by the possible presence of a substantial symmetric DM component.

In the RP, we show the allowed region in the $Y_B - \epsilon_l$ plane. As expected, here also a direct correlation exist between these two quantities. The narrow magenta band is the present day accepted value of the baryon anti-baryon asymmetry of the Universe. The parameter values for which the baryon asymmetry lies above the magenta line are ruled out, while below the line leptogenesis cannot provide the full matter-antimatter asymmetry

Figure 3. LP (RP): scatter plot in the $y_{ee} - M_{N_1}$ ($y_{\mu\mu}^R - M_{N_2}$) plane after satisfying neutrino oscillation data as mentioned in section 3. All the parameters have been varied in the range as shown in section 3.

of the Universe. We see that we need a CP violation parameter $\epsilon_l \sim 10^{-6}$ in order to obtain the observed baryon asymmetry. Moreover, comparing the LP with the RP, it is clear that we can indeed achieve the right abundance of Dark Matter and baryons for the case $\epsilon_l \sim \epsilon_D$ and then the Dark Matter mass is in the range 1–10 GeV as expected.

In the LP and RP of figure 3, we show the correlation among the RH neutrino masses and the elements of the Dirac mass matrix, which is required in order to satisfy the oscillation data. We see that the ratio $\frac{y_{ee}^2}{M_{N_1}}$ is fixed by neutrino oscillations, as it provides one of the mass scales in the light neutrino mass matrix. Similar correlations of M_{N_1} are present also for the other Yukawa parameters in the first column, $y_{\mu e}, y_{\tau e}$. Here we are generating the neutrino mass by the Type-I seesaw mechanism, hence, the elements in the first column of the Dirac mass matrix (see eq. (3.2) and appendix A) are directly related to M_{N_1} . Similarly, the elements of the second column of the Dirac mass matrix are always suppressed by $1/M_{N_2}$ when they appear in the light neutrino mass matrix. Therefore, for the RP we also see a similar kind of correlation with elements of the second column of Dirac mass matrix and M_{N_2} , but in this case the correlation is less sharp. Since the RH neutrinos N_1 and N_2 are hierarchical, a difference in the magnitude of the elements of the first and second column of the Yukawa matrix are arranged such that both the light neutrino mass squared differences we obtain are consistent with the neutrino oscillation data. In these plots, all points are allowed by neutrino oscillation data and some of them (indicated by blue colour) produce the observed matter-antimatter asymmetry of the Universe and for rest we need extra sources of baryogenesis. Indeed we have adjusted the imaginary part of the elements of the second column of the Yukawa matrix such that both the lepton asymmetry as well as the light neutrino mass matrix are obtained, consistent with their respective measured values.

Figure 4. LP (RP): scatter plot in the $y_{ee} - y_{\mu e}$ ($y_{ee} - y_{\tau e}$) plane after satisfying neutrino oscillation data as mentioned in section 3. All the parameters have been varied in the range as shown in section 3.

Figure 5. LP (RP): scatter plot in the $y_{\mu\mu}^R - y_{\mu\mu}^I$ ($\epsilon_l - \epsilon_D$) plane after satisfying neutrino oscillation data as mentioned in section 3. All the parameters have been varied in the range as shown in section 3.

In figure 4, we show the allowed region after satisfying the neutrino oscillation data in both LP and RP. A clear correlation exist among the parameters in order to satisfy neutrino oscillation data. Only the blue points are close to the current value of the Universe's baryon asymmetry, while the red and green points give a too low value.

The LP of figure 5 shows explicitly the correlation between the real and imaginary part of the same element of the Dirac mass matrix, i.e. $y_{\mu\mu}^R$ and $y_{\mu\mu}^I$. Either of them can give

Figure 6. Scatter plot in the $m_\nu^{ee} - Y_B$ plane after satisfying neutrino oscillation data (N.O.D) as mentioned in section 3. All the parameters have been varied in the range as shown in section 3.

the right contribution to the light neutrino mass matrix to fit the oscillation data, but only a substantial imaginary part allows for non-vanishing CP violation and the production of a sufficiently large lepton asymmetry. This plot shows that the correct baryon asymmetry can be obtained both if the real and imaginary part are equal and large or when the real part is negligible and the imaginary part provides a substantial contribution to both the CP violation and the neutrino mass. This two cases correspond to the limiting cases discussed earlier.

In RP of the same figure we show the allowed region in the $\epsilon_l - \epsilon_D$ plane. We see that our predicted baryon asymmetry comes close to the measured value only for reasonably high values of ϵ_l . Since there exists a sharp correlation between Y_B and ϵ_l , lower ϵ_l results in production of lower lepton asymmetry. As we discussed earlier, we expect ϵ_D to be equal or less than ϵ_l and indeed this is also reproduced in this figure. Note that while the baryon asymmetry is correctly given only in a quite narrow region of the parameter space, a much wider range of ϵ_D is allowed as we can adjust the DM mass to match the Dark Matter abundance. Generically the mass of the Dark Matter state is given by the VEV of the Dark $SU(2)_D$ scalar doublet and can vary compared to the scale of SM fermion masses. Nevertheless a natural range for the mechanism to work are DM masses in the GeV to tens of GeV range, well below the WIMP mass scale around 1 TeV.

In figure 6, we give the allowed model points in the $m_\nu^{ee} - Y_B$ plane. Both the blue and magenta points satisfy fully the neutrino oscillation data and the only difference between them is the choice of the Dirac mass matrix elements in eq. (3.2). The blue points correspond to the general M_D matrix with complex second column, whereas the magenta points are obtained for a purely imaginary second column of M_D (i.e. $y_{i\mu}^R = 0$, $i = e, \mu, \tau$). Taking the Majorana phase angle convention for the light neutrinos as $\text{diag}(1, e^{i\alpha_{21}/2}, e^{i\alpha_{31}/2})$ and

one zero mass eigenstate $m_1 = 0$, we can write down the (1,1) element of the neutrino mass matrix (m_ν), which generates the $0\nu\beta\beta$ decay, as

$$m_\nu^{ee} = [s_{12}^4 c_{13}^4 m_2^2 + 2m_2 m_3 s_{12}^2 c_{13}^2 s_{13}^2 \cos(\alpha_{31} - \alpha_{21} - 2\delta_{\text{CP}}) + m_3^2 s_{13}^4]^{1/2} \quad (5.1)$$

where m_2, m_3 are the light neutrino masses, $c_{ij} = \cos\theta_{ij}$, $s_{ij} = \sin\theta_{ij}$ are the mixing angles, while δ_{CP} is Dirac CP phase and α_{31}, α_{21} are the Majorana phases.⁵ Varying the mixing angles and the phases to cover $\cos(\alpha_{31} - \alpha_{21} - 2\delta_{\text{CP}}) = \pm 1$, we get a range of m_ν^{ee} for the normal hierarchy which lies in between 1 meV and 4.3 meV. This range is in complete agreement with the figure 6 which we obtain after satisfying the neutrino oscillation data (as given in section 3) at the time of diagonalising neutrino mass matrix m_ν . Magenta points with a real light neutrino mass matrix correspond to $\delta_{\text{CP}} = 0$ and $(\alpha_{31} - \alpha_{21}) = \pi$, so they point towards the lowest possible value of the neutrinoless double beta decay rate, while the small spread in the $Y_B - m_\nu^{ee}$ plane is mainly due to the variation of θ_{12} and m_2 within their experimental ranges. Hence, with a purely imaginary second column of M_D , we can satisfy both neutrino oscillation data and produce the baryon and DM asymmetry of the Universe with the Dirac CP phase $\delta_{\text{CP}} = 0$ and $\epsilon_l = \epsilon_D$. This is the case of the minimal number of parameters in the neutrino sector.

In the case of a general M_D matrix (shown by the blue points), we can obtain higher values of m_ν^{ee} and hence it will be easier for $0\nu\beta\beta$ experiments to test this scenario. The higher values of $m_\nu^{ee} \sim 4.2$ meV imply as well a Dirac CP phase $\delta_{\text{CP}} = -\frac{\pi}{2}$ (in accordance with the recent data from the neutrino oscillation experiments) when $(\alpha_{31} - \alpha_{21}) = \pi$. Therefore, with a generic M_D matrix, we can satisfy the recent data of the neutrino oscillation parameters and produce the baryon and DM asymmetry of the Universe also with large Dirac CP phase, $\delta_{\text{CP}} \sim -\frac{\pi}{2}$.

6 Direct detection of Dark Matter

In our model, although DM only has a dark $\text{SU}(2)_D$ charge, it can still *talk* to the visible sector through the exchange of SM-like Higgs boson h_1 and dark sector Higgs h_2 respectively and the corresponding Feynman diagram is shown in figure 7.

The expression for the spin-independent scattering cross section between DM and nucleon mediated by scalars h_1, h_2 is given by

$$\sigma_{\text{SI}} = \frac{\mu_{\text{red}}^2}{\pi} \left[\frac{M_N f_N}{v} \left(\frac{g_{\psi_1 \psi_1 h_2} \sin \zeta}{M_{h_2}^2} + \frac{g_{\psi_1 \psi_1 h_1} \cos \zeta}{M_{h_1}^2} \right) \right]^2, \quad (6.1)$$

where $\mu_{\text{red}} = \frac{M_N M_{\text{DM}}}{M_N + M_{\text{DM}}}$ is the reduced mass and $f_N \sim 0.3$ [88]. The DM couplings with the scalars have the following form

$$\begin{aligned} g_{\psi_1 \psi_1 h_1} &= -\frac{\lambda_1}{\sqrt{2}} \sin \zeta = -\frac{m_{\text{DM}}}{v_D} \sin \zeta \\ g_{\psi_1 \psi_1 h_2} &= \frac{\lambda_1}{\sqrt{2}} \cos \zeta = \frac{m_{\text{DM}}}{v_D} \cos \zeta, \end{aligned} \quad (6.2)$$

⁵Note that for one massless neutrino eigenstate, only one Majorana phase, given as the difference between α_{31} and α_{21} , is physical.

Figure 7. Feynman diagram for the spin-independent scattering cross section of Dark Matter with nucleon mediated by both SM-like Higgs h_1 and hidden sector Higgs h_2 .

so we see that we have a negative interference and full cancellation for equal masses of the dark and SM Higgs fields, where the mixing in the scalar sector vanishes. Indeed we obtain

$$\sigma_{\text{SI}} = \frac{\mu_{\text{red}}^2}{\pi} \left[\frac{M_N f_N}{v} \frac{|\lambda_1| \sin 2\zeta}{2\sqrt{2} M_{h_1}^2} \left(1 - \frac{M_{h_1}^2}{M_{h_2}^2} \right) \right]^2. \quad (6.3)$$

The masses M_{h_1} , M_{h_2} and mixing angle ζ are given in eqs. (2.7), (2.6) while λ_1 is the Yukawa coupling between Ψ_{1L} , $\tilde{\phi}_D$ and ψ_{1R} , related as well to the Dark Matter mass by eq. (2.9).

In figure 8, we plot the variation of the spin-independent DM-nucleon scattering cross section with the DM mass. In generating the plot, we have kept λ_1 fixed at unity while DM mass has been varied in the range 10 GeV to 10^4 GeV by appropriately adjusting v_D . The mixing angle ζ has been scanned over its present allowed range i.e $\zeta < 10^{-1}$ rad. From the figure 7, we see that a part of the parameter space, corresponding to large ζ and heavy M_{h_2} , is already ruled out by the null results at various direct detection experiments like PandaX-II [61] and Xenon1T [60], however enough parameter space is still open and can be tested in the future experiments like Darwin [59].

7 Conclusion

In this work we have tried to solve three major puzzles of cosmology by the presence of two RH neutrinos and a Dark Sector charged under an $SU(2)_D$. The hidden sector of the model is chosen to resemble the SM electroweak sector, but with just two non mixing families, so that the mass of the DM particles could be similar to the SM fermions and the presence of the $SU(2)_D$ interaction is crucial for annihilating away all the symmetric Dark Matter components.

We generate the neutrino mass through the Type-I seesaw mechanism, while both the lepton (later processed into baryons) and a DM asymmetries are produced from the decay of the lightest RH neutrino. Hence, all the three BSM phenomena have a common origin. Moreover, the CP violation in both sectors is related to complex entries in the neutrino Yukawa couplings. Two RH neutrinos with hierarchical masses are sufficient to accommodate all the present oscillation data and satisfy successfully the present day

Figure 8. Variation of σ_{SI} with DM mass where the scalar mixing angle ζ has been varied in the range as shown in legend. Other parameters have been kept fixed $M_{h_1} = 125.5$ GeV, $M_{h_2} = 1.5$ TeV and $\lambda_1 = 1.0$.

bounds on the sum of the light neutrino masses. We have shown that the limited number of parameters in the model results in strong correlations among some of the entries in the Dirac and Majorana mass matrices.

Due to the strong hierarchy, we can generate the baryon asymmetry of the Universe and an asymmetric Dark Matter component at the same time from the decay of the lighter RH neutrino state. Also in this case, two RH neutrino states are enough to give both, as long as they have similar Yukawa coupling with the light neutrinos and the additional dark fermions, so that the two decays can have naturally a similar decay rate and branching fraction. As CP violation is sourced in both sectors by the neutrino Yukawa y_{ij} , comparable values of ϵ for leptons and the Dark Matter are expected in most parameter space. Indeed, the two CP asymmetries in the decays are exactly equal if the neutrino Dirac mass has a purely imaginary column, which gives a real Majorana mass matrix for the light neutrinos and a vanishing Dirac phase. We expect nevertheless to have one non-trivial Majorana phase, that could lead to a partial cancellation in the matrix element for neutrinoless $\beta\beta$ decay. Similarly the CP asymmetry are similar as well in the case $\alpha_2 \gg \alpha_1$ for the Dark sector Yukawa with the RH neutrinos. In other cases we have $\epsilon_D < \epsilon_l$, but this can be partially compensated by the presence of less effective wash-out processes or in any case by a larger DM mass. For specific choices of the parameters, i.e. the imaginary parts of the Yukawas and Dark Matter mass, we can simultaneously satisfy neutrino oscillation

parameters bounds and produce the correct value of the baryon asymmetry and of the Dark Matter energy density.

In this scenario the visible and dark sector do not communicate only through the neutrino portal: indeed after the electroweak and dark $SU(2)_D$ symmetries are broken, a mixing appears also in the scalar sector, so that the physical Higgs field contains also a small component of Dark Higgs and can couple to Dark Matter. For light Dark Matter we therefore expect an invisible contribution to the Higgs width proportional to the mixing angle in the scalar sector. If the Dark Matter is heavier than half of the Higgs mass, an observable signal could still appear in the Direct Detection experiments and from the production of Dark Matter at colliders through an off-shell Higgs. Moreover, as in many models with an extended Higgs sector, we can also expect to detect exotic scalars at colliders. Regarding the heavier Higgs state h_2 , Direct Detection and collider experiments are in our model highly complementary, as for heavy M_{h_2} , whereas the production of h_2 at colliders is suppressed, the scattering cross section with nucleons becomes larger for fixed mixing angle ζ . Note as well that in this type of model, as in all models with extended Higgs sector, the Higgs self-couplings are modified in a characteristic way and that may be observable even when all the Dark Sector particles are beyond the present collider reach, see e.g. [89, 90].

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A Expression for the Majorana mass matrix of light neutrinos

Here we have given the expression of all the elements of the light neutrino mass matrix m_ν (using eq. (3.1)) in terms of the Yukawa couplings and the RH neutrino masses.

$$\begin{aligned}
 (\tilde{m}_\nu)_{11} &= \frac{y_{ee}^2}{M_{N_1}} + \frac{(y_{e\mu}^R)^2 - (y_{e\mu}^I)^2}{M_{N_2}} - i \frac{2 y_{e\mu}^R y_{e\mu}^I}{M_{N_2}}, \\
 (\tilde{m}_\nu)_{12} &= \frac{y_{ee} y_{\mu e}}{M_{N_1}} + \frac{y_{e\mu}^R y_{\mu\mu}^R - y_{e\mu}^I y_{\mu\mu}^I}{M_{N_2}} - i \frac{y_{e\mu}^R y_{\mu\mu}^I + y_{e\mu}^I y_{\mu\mu}^R}{M_{N_2}}, \\
 (\tilde{m}_\nu)_{13} &= \frac{y_{ee} y_{\tau e}}{M_{N_1}} + \frac{y_{e\mu}^R y_{\tau\mu}^R - y_{e\mu}^I y_{\tau\mu}^I}{M_{N_2}} - i \frac{y_{e\mu}^R y_{\tau\mu}^I + y_{e\mu}^I y_{\tau\mu}^R}{M_{N_2}}, \\
 (\tilde{m}_\nu)_{21} &= (\tilde{m}_\nu)_{12}, \\
 (\tilde{m}_\nu)_{22} &= \frac{y_{\mu e}^2}{M_{N_1}} + \frac{(y_{\mu\mu}^R)^2 - (y_{\mu\mu}^I)^2}{M_{N_2}} - i \frac{2 y_{\mu\mu}^R y_{\mu\mu}^I}{M_{N_2}}, \\
 (\tilde{m}_\nu)_{23} &= \frac{y_{\mu e} y_{\tau e}}{M_{N_1}} + \frac{y_{\mu\mu}^R y_{\tau\mu}^R - y_{\mu\mu}^I y_{\tau\mu}^I}{M_{N_2}} - i \frac{y_{\mu\mu}^R y_{\tau\mu}^I + y_{\mu\mu}^I y_{\tau\mu}^R}{M_{N_2}}, \\
 (\tilde{m}_\nu)_{31} &= (\tilde{m}_\nu)_{13}, \\
 (\tilde{m}_\nu)_{32} &= (\tilde{m}_\nu)_{23}, \\
 (\tilde{m}_\nu)_{33} &= \frac{y_{\tau e}^2}{M_{N_1}} + \frac{(y_{\tau\mu}^R)^2 - (y_{\tau\mu}^I)^2}{M_{N_2}} - i \frac{2 y_{\tau\mu}^R y_{\tau\mu}^I}{M_{N_2}}, \tag{A.1}
 \end{aligned}$$

$$m_\nu = -\frac{v^2}{2} \begin{pmatrix} (\tilde{m}_\nu)_{11} & (\tilde{m}_\nu)_{12} & (\tilde{m}_\nu)_{13} \\ (\tilde{m}_\nu)_{21} & (\tilde{m}_\nu)_{22} & (\tilde{m}_\nu)_{23} \\ (\tilde{m}_\nu)_{31} & (\tilde{m}_\nu)_{32} & (\tilde{m}_\nu)_{33} \end{pmatrix}. \tag{A.2}$$

We see from these expressions that if the second column of the Yukawa matrix is imaginary, i.e. for $y_{i\mu}^R = 0$, the light neutrino mass is the sum of two degenerated real matrices, each with a single non-zero eigenvalue and opposite sign. If the massive eigenvectors of the two matrices are orthogonal to each other, we have then simply two opposite-sign mass eigenstates and one zero mass eigenstate as:

$$m_3 = -\frac{v^2}{2} \frac{\sum_i y_{ie}^2}{M_1}, \tag{A.3}$$

$$m_2 = \frac{v^2}{2} \frac{\sum_i (y_{i\mu}^I)^2}{M_2}, \tag{A.4}$$

$$m_1 = 0. \tag{A.5}$$

So we have to choose the hierarchy in M_i and the Yukawa couplings appropriately in order to match the measured mass differences. If the two massive eigenvectors are not orthogonal, a more complex mixing pattern appears and the two mass eigenstates obtain contributions from both heavy RH neutrinos, nevertheless for hierarchical masses and not so strongly hierarchical Yukawas, still the heaviest mass m_3 is mostly determined by the lighter mass M_1 . Generically for real m_ν , the mixing matrix is real, so that the Dirac phase is exactly vanishing, but the Majorana phases are not, as they have to be chosen

to give positive light neutrino masses, i.e. we obtain for the mass eigenstates above the Majorana phases $\alpha_{31} = \pi, \alpha_{21} = 0$. We expect also in the generic case with a purely imaginary Yukawa column to have two eigenstates with different Majorana phases, leading to a partial cancellation in the matrix element for neutrinoless double beta decay:

$$m_{\beta\beta} = \left| \sum_i m_i U_{ei}^2 \right| = |m_3 \sin^2 \theta_{13} - m_2 \cos^2 \theta_{13} \sin^2 \theta_{12}| \sim 10^{-2} \text{eV} . \quad (\text{A.6})$$

B Expression of γ_{D_1} , $\gamma_{\phi_h, s}^1$ and $\gamma_{\phi_h, t}^1$:

Expression of γ_{D_1} takes the following form [81],

$$\gamma_{D_1} = n_{N_1}^{eq} \frac{K_1(Z)}{K_2(z)} \Gamma_{N_1} , \quad (\text{B.1})$$

where $n_{N_1}^{eq}$ is the equilibrium number density of the RH neutrino N_1 and $K_n(z)$ is the n th order modified Bessel function of second kind while Γ_{N_1} is the total decay width of N_1 . The expression of Γ_{N_1} is given in eq. (4.6).

Further, the general expression of $\gamma(a + b \leftrightarrow i + j + \dots)$ for a two body scattering process $a + b \leftrightarrow i + j + \dots$ is given by [81],

$$\gamma(a + b \leftrightarrow i + j + \dots) = \frac{T}{64\pi^4} \int_{(M_a + M_b)^2}^{\infty} ds \hat{\sigma}(s) \sqrt{s} K_1 \left(\frac{\sqrt{s}}{T} \right) , \quad (\text{B.2})$$

where, s is one of the Mandelstam variables which physically represents square of the centre of mass energy for a scattering process in centre of momentum frame. Moreover, $\hat{\sigma}(s)$ is the reduced cross section for the scattering process $a + b \leftrightarrow i + j + \dots$, which is related to the actual cross section by the following relation

$$\hat{\sigma}(s) = \frac{8}{s} \left[(p_a \cdot p_b)^2 - M_a^2 M_b^2 \right] \sigma(s) . \quad (\text{B.3})$$

Here, p_i and M_i are the three momentum and mass of the species i respectively. The expressions of reduced cross sections for the processes $N_1 + l \rightarrow \bar{t} + q$ (s -channel process mediated by ϕ_h) and $N_1 + t \rightarrow \bar{l} + q$ (t -channel process mediated by ϕ_h) are given as

$$\begin{aligned} \hat{\sigma}_{\phi_h, s} &= \frac{3\pi\alpha^2 M_t^2}{M_W^4 \sin^4 \theta_w} (M_D^\dagger M_D)_{11} \left[\frac{s - M_{N_1}^2}{s} \right]^2 , \\ \hat{\sigma}_{\phi_h, t} &= \frac{3\pi\alpha^2 M_t^2}{M_W^4 \sin^4 \theta_w} (M_D^\dagger M_D)_{11} \left[\frac{s - M_{N_1}^2}{s} + \frac{M_{N_1}^2}{s} \ln \left(\frac{s - M_{N_1}^2 + M_{h_1}^2}{M_{h_1}^2} \right) \right] , \end{aligned} \quad (\text{B.4})$$

where $\alpha = \frac{g^2 \sin^2 \theta_w}{4\pi}$, g being the $\text{SU}(2)_L$ gauge coupling and θ_w is the weak mixing angle (Weinberg angle).

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