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.

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I. 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 from 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 [27–49], 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 $SU(2)_{D}$ interaction \(^{1}\) in the Dark sector and a discrete $Z_{3}$ group, ensuring the Dark Matter stability. 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 also introduce two $SU(2)_{D}$ fermionic doublets just to overcome the Witten anomaly [51]. 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 [52–58]. The other exotic scalar doublet does not 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 II, we describe all the details of our model. The generation of the neutrino mass is discussed in section III. In section IV, we give the Boltzmann equations for both the dark sector and the leptonic sector, while section V contains the full numerical result of the Boltzmann equations. In section VI the DM direct detection is addressed. Finally, in section VII we conclude our work with an outlook on possible signature at colliders.

II. 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 $\psi_{i}$ ($i = 1$ 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 I we show the complete list of particles in the present model. Here we want to point that, since we have an even number

\(^{1}\) WIMP type DM obeying $SU(2)$ gauge symmetry has been studied earlier in [50].
of fermionic SU(2)$_D$ doublets, our model is free from the Witten anomaly [51]. In addition, we have two scalar doublets in the hidden sector as well. One of the scalar doublets $\eta_D$ does not get any vacuum expectation value (VEV) while the remaining one ($\phi_D$) has a nonzero VEV and thus mixes with the SM Higgs doublet $\phi_h$. 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 $\mathbb{Z}_3$. Furthermore, we have two right handed (RH) fermions $N_i$ ($i = 1, 2$), singlets under both SM gauge group as well as SU(3)$_D$. 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 [59–64].

A. Particle spectrum

| Gauge Group | Fermion Fields | Scalar Fields |
|-------------|----------------|--------------|
| SU(3)$_c$   | $\Psi_1L = (\psi_1, \psi_2)^T_L$ $\psi_1R$ $\psi_2R$ $\Psi_2L = (\psi_3, \psi_4)^T_L$ $\psi_3R$ $\psi_4R$ $N_i$ $\phi_h$ $\phi_D$ $\eta_D$ |
| SU(2)$_L$   | 1 1 1 1 1 1 1 |
| SU(2)$_D$   | 2 1 1 2 1 1 1 |
| $\mathbb{Z}_3$ | $\omega$ $\omega$ $\omega$ $\omega^2$ $\omega^2$ $\omega^2$ 1 1 1 |

Table I: 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 $\phi_h$ which has hypercharge $Y = 1/2$.

B. 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,

$$\mathcal{L} = \mathcal{L}_{SM} + i \bar{\Psi}_k \gamma^\mu D^k \Psi_k + (D^\mu \phi_D) \dagger (D^\mu \phi_D) + (D^\mu \eta_D) \dagger (D^\mu \eta_D) + \left( y_{ij} \bar{L_i} \tilde{\phi}_h N_{jR} + h.c. \right) + \left( \lambda_1 \bar{\Psi}_L \phi_D \psi_1 R + \lambda_2 \bar{\Psi}_L \phi_D \psi_2 R + \lambda_3 \bar{\Psi}_L \phi_D \psi_3 R + \lambda_4 \bar{\Psi}_L \phi_D \psi_4 R + h.c. \right) - \alpha_j \bar{\Psi}_L \eta_D N_{jR} - \beta_j \bar{\Psi}_L \eta_D N_{jR} + i \bar{N}_{jR} \phi N_{jR} - M_j N_{jR}^c N_{jR} - V(\phi_h, \phi_D, \eta_D), \tag{1}$$

where $i = 1$ to 3 while $j$ runs from 1 to 2 and $D^\mu$ is the covariant derivative for $SU(2)_D$. One important comment we would like to make here is that both the lightest components of dark
doublets \( \Psi_{1L} \) and \( \Psi_{2L} \) are stable by virtue of different \( \mathbb{Z}_3 \) charge assignments. Moreover, both \( \psi_1 \) and \( \psi_3 \) are produced in the same way (from the decays of RH-neutrinos) and have very similar phenomenology. Due to the fact that the only scalar field with non-trivial \( \mathbb{Z}_3 \) charge is also a doublet, we cannot have any Yukawa interaction between \( \psi_1 \) and \( \psi_3 \), so the only way to transfer the abundance of one type of DM into another is via the RH neutrinos. So assuming that the couplings of \( \psi_3 \) to \( N_{jR} \), i.e. \( \beta_j \), are sufficiently small, we can neglect the asymmetry produced in the second DM state and the transfer processes.

Under these circumstances, it is sufficient to study the phenomenology of one of the dark fermions, the one providing the larger contribution to the DM energy density. Indeed the case when both fermions have a similar couplings and energy densities are realised only in a very narrow region of the parameter space and it is not qualitatively very different in the physics predictions. Henceforth, for the sake of simplicity we will just consider the lightest component of \( \Psi_{1L} \) as our DM candidate, but a similar discussion also holds for the other DM candidate contained in the doublet \( \Psi_{2L} \). Note that in both cases the \( SU(2)_D \) gauge symmetry in the early Universe is crucial in order to allow for the annihilation of the symmetric Dark Matter component into the gauge and scalar sector, i.e. in particular the \( \phi_D \), which we will assume lighter than the Dark Matter.

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

\[
V(\phi_h, \phi_D, \eta_D) = -\mu_h^2 (\phi_h^\dagger \phi_h)^2 + \lambda_h (\phi_h^\dagger \phi_h)^2 - \mu_D^2 (\phi_D^\dagger \phi_D)^2 + \lambda_D (\phi_D^\dagger \phi_D)^2
\]

\[
+ \mu_{\eta}^2 (\eta_D^\dagger \eta_D)^2 + \lambda_{\eta} (\eta_D^\dagger \eta_D)^2 + \lambda_{h\eta} (\phi_h^\dagger \phi_h)(\eta_D^\dagger \eta_D)
\]

\[
+ \lambda_{D1} (\phi_D^\dagger \phi_D)(\eta_D^\dagger \eta_D) + \lambda_{D2} (\phi_D^\dagger \phi_D)(\eta_D^\dagger \eta_D) + \lambda_{D3} (\phi_D^\dagger \phi_D + h.c.).
\]

Here, both the doublets \( \phi_h \) and \( \phi_D \) 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 \( \phi_h \) and \( \phi_D \) take the following form after symmetry breaking,

\[
\phi_h = \begin{pmatrix} 0 \\ \frac{v_h}{\sqrt{2}} \end{pmatrix}, \quad \phi_D = \begin{pmatrix} 0 \\ \frac{v_D + H}{\sqrt{2}} \end{pmatrix}.
\]

In the scalar potential there is a mixing term between the CP even neutral components of the two doublets \( \phi_h \) and \( \phi_D \), hence the gauge basis and mass eigenbasis will be different. In \( h, H \)

\footnote{Like 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.}
basis the mass square mixing matrix will be as follows

\[ M_{\text{scalar}}^2 = \begin{pmatrix} 2\lambda_h v^2 & \lambda_{hD} v v_D \\ \lambda_{hD} v v_D & 2\lambda_D v_D^2 \end{pmatrix}, \]  

(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

\[ h_1 = h \cos \zeta - H \sin \zeta, \]
\[ h_2 = h \sin \zeta + H \cos \zeta, \]  

(5)

where \( \zeta \) 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

\[ \tan 2\zeta = \frac{\lambda_{hD} vv_D}{\lambda_D v_D^2 - \lambda_h v^2}. \]  

(6)

As mentioned above, after diagonalising the scalar mass matrix in Eq. (4), we get the physical masses for the two neutral scalars as

\[ M_{h_1}^2 = \lambda_h v^2 + \lambda_D v_D^2 - \sqrt{(\lambda_D v_D^2 - \lambda_h v^2)^2 + (\lambda_{hD} v v_D)^2}, \]
\[ M_{h_2}^2 = \lambda_h v^2 + \lambda_D v_D^2 + \sqrt{(\lambda_D v_D^2 - \lambda_h v^2)^2 + (\lambda_{hD} v v_D)^2}. \]  

(7)

We identify the lighter Higgs scalar as the SM-like Higgs observed at the LHC. Therefore, we take \( M_{h_1} = 126 \) GeV and consider small mixing angle \( \sin \zeta \leq 0.1 \) in order to ensure agreement with the Higgs signal strengths at the LHC [65–67].

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

\[ \lambda_D = \frac{M_{h_2}^2 + M_{h_1}^2 + (M_{h_2}^2 - M_{h_1}^2) \cos 2\zeta}{4v_D^2}, \]
\[ \lambda_h = \frac{M_{h_2}^2 + M_{h_1}^2 - (M_{h_2}^2 - M_{h_1}^2) \cos 2\zeta}{4v_D^2}, \]
\[ \lambda_{hD} = \frac{(M_{h_2}^2 - M_{h_1}^2) \sin 2\zeta}{2vv_D}, \]
\[ \mu_h^2 = \lambda_h v^2 + \lambda_{hD} v v_D^2, \]
\[ \mu_D^2 = \lambda_D v_D^2 + \lambda_{hD} v v_D^2. \]  

(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

$$M_{\psi_1} = M_{DM} = \frac{\lambda_1 v_D}{\sqrt{2}}.$$  \hspace{1cm} (9)

The other scalar doublet $\eta_D$, 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 $\eta_D$, we consider the fermion $\psi_1$ as the lightest one. This is always possible by tuning the couplings related to $\eta_D$ and $\lambda_\alpha s$ ($\alpha = 2$ to $4$) so that heavier $\mathbb{Z}_3$ charged particles decay to the lightest one and thereby $\psi_1$ becomes stable. Thus, $\psi_1$ will be a viable DM candidate in our model. Moreover, the invariance of the $\mathbb{Z}_3$ symmetry actually makes 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.

III. NEUTRINO MASS

In this work, as mentioned in the Model section (Section II), 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

$$m_\nu = -M_D M_R^{-1} M_D^T,$$  \hspace{1cm} (10)

where $M_D$ is the Dirac mass matrix ($3 \times 2$) while the Majorana mass matrix for the heavy right handed neutrinos $N_{jR}$ ($j = 1, 2$) is denoted by a $2 \times 2$ 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,

$$M_D = \frac{y_{ij} v}{\sqrt{2}} = \frac{v}{\sqrt{2}} \begin{pmatrix} y_{ee} & y_{e\mu}^R - iy_{e\mu}^I \\ y_{\mu e} & y_{\mu\mu}^R - iy_{\mu\mu}^I \\ y_{\tau e} & y_{\tau\mu}^R - iy_{\tau\mu}^I \end{pmatrix}.$$  \hspace{1cm} (11)

We have computed the physical masses of light neutrinos (eigenvalues) and mixing angles (eigen-vectors) by diagonalising the complex symmetric Majorana mass matrix $m_\nu$. All the elements
of $m_{\nu}$ are explicitly given in the Appendix. In studying the neutrino phenomenology, we take into account the observed values of neutrino oscillation parameters \[68\]. 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 $m_{\nu}$ 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 $3\sigma$ bounds are as follows \[68\],
  \[
  6.93 \leq \frac{\Delta m_{21}^2}{10^{-5}} \text{ (eV}^2) \leq 7.97 \quad \text{and} \quad 2.37 \leq \frac{\Delta m_{31}^2}{10^{-3}} \text{ (eV}^2) \leq 2.63 \, .
  \]

- Three mixing angles namely, solar mixing angle $\theta_{12}$, atmospheric mixing angle $\theta_{23}$ and reactor mixing angle $\theta_{13}$ are also very well measured now. The allowed values of mixing angles for NH in $3\sigma$ are listed below \[68\],
  \[
  30^0 \leq \theta_{12} \leq 36.51^0 , \quad 7.82^0 \leq \theta_{13} \leq 9.02^0 , \quad 37.99^0 \leq \theta_{23} \leq 51.71^0 ,
  \]

- 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 $\delta$, which is yet to be measured accurately by the different ongoing and upcoming oscillation experiments like T2K \[69\], T2HK \[70\], DUNE \[71-74\] and INO \[75\]. There is a hint of maximal CP violation ($\delta \sim -\frac{\pi}{2}$) from the T2K experiment \[76\], which excludes the CP conserving values of $\delta = 0$ or $\pi$, at $90\%$ C.L. However, there is some tension between the observed value of $\delta$ from T2K and NOvA \[77\].

In satisfying the above mentioned neutrino oscillation parameters, we vary the model parameters, specifically, the elements of the Dirac mass matrix (Eq. (11)) and the RH-neutrino masses in the following range,

\[
10^{-4} \leq y_{ee}, y_{\mu e}, y_{\tau e}, y_{e\mu}^R \leq 10^{-2} ,
\]

\[
10^{-3} \leq y_{\mu\mu}, y_{\mu\mu}^R, y_{\tau\mu}^R, y_{\tau\mu}^I \leq 1 ,
\]

\[
10^{-6} \leq y_{e\mu}^I \leq 10^{-4} ,
\]

\[
10^8 (10^{13}) \leq M_{N_1} (M_{N_2}) \text{ GeV} \leq 10^{11} (10^{15}) .
\]

In Eq. (11), 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 $y^{R,I}_{\alpha\mu} = |y^{R,I}_{\alpha\mu}|$, ($\alpha = e, \mu, \tau$). 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. (14). 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 V), we will see that there exist correlations among the parameters of the neutrino sector and hence not the entire adopted ranges in Eq (14) are allowed by the neutrino oscillation data and the requirement of successful leptogenesis.

**IV. 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,

\[
\frac{dY_{N_1}}{dz} = -\frac{z}{sH(M_1)} \left[ \left( \frac{Y_{N_1}}{Y_{eq}} - 1 \right) \times \left( \gamma_{D_1} + 2\gamma_{l,s}^1 + 4\gamma_{l,t}^1 \right) \right], \tag{15}
\]

\[
\frac{dY_{\Delta l}}{dz} = -\frac{\Gamma_{N_1}}{H(M_1)} \left[ \epsilon_{l_1} \frac{zK_1(z)}{K_2(z)} (Y_{eq} - Y_{N_1}) + Br_l z^3 K_1(z) Y_{\Delta l} \right], \tag{16}
\]

\[
\frac{dY_D}{dz} = -\frac{\Gamma_{N_1}}{H(M_1)} \left[ \epsilon_{D} \frac{zK_1(z)}{K_2(z)} (Y_{eq} - Y_{N_1}) + Br_D z^3 K_1(z) Y_D \right], \tag{17}
\]

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 \( \gamma_{D_1} \) 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, \( \gamma_{l,s}^1 \) and \( \gamma_{l,t}^1 \) are related to s-channel and t-channel scattering of \( N_1 \) mediated by \( \phi_{ph} \), which can also lead to the destruction or production of \( N_1 \). The expressions of \( \gamma_{D_1} \), \( \gamma_{l,s}^1 \) and \( \gamma_{l,t}^1 \) 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. (16) (Eq. (17)) 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. (16, 17), \( Br_l \) 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 Fig. 2 of [79]), 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]

\[
\epsilon_l = \frac{\Gamma(N_1 \to L\phi_h) - \Gamma(N_1 \to \bar{L}\phi_h^*)}{\Gamma_{N_1}},
\]

\[
= \frac{M_{N_1}}{16\pi M_{N_2}} \text{Im} \left[ \frac{3 ((y^\dagger y)_{12}^* )^2 + 2\alpha_1 \alpha_2 (y^\dagger y)_{12}^* }{((y^\dagger y)_{11} + \alpha_1 \alpha_1^*)} \right], \tag{18}
\]

where we have normalized to the total RH neutrino decay rate and summed over the lepton flavours, i.e. \( \epsilon_l = \sum_\alpha \epsilon_{l,\alpha}^l \). In Eq. (18) 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 

\[ \epsilon_D = \frac{\Gamma(N_1 \rightarrow \psi_1 \eta_D) - \Gamma(N_1 \rightarrow \psi_1 \eta_D^\dagger)}{\Gamma_{N_1}} \]

\[ = \frac{M_{N_1}}{16\pi M_{N_2}} \frac{\text{Im}[2\alpha_1^*\alpha_2(y^\dagger y)_{12}^* + 3(\alpha_1^*\alpha_2)^2]}{[(y^\dagger y)_{11} + \alpha_1\alpha_1^*]^2}. \]  

(19)

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 SM particles. The total decay width of the RH-neutrino \( N_1 \) is given by

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

(20)

From Eqs. (18, 19) we see that both \( \epsilon_l \) and \( \epsilon_D \) are determined by the Yukawa couplings \( y_{ij} \) and RH-neutrino masses. One important thing to note here is that in the dark sector we can absorb the phases of the couplings \( \alpha_j \) \((j = 1, 2)\) by redefining the phases of \( \psi_1 \) and the complex scalar doublet \( \eta_D \), as long as we consider the theory above the scale of \( SU(2)_D \) breaking. In that case the CP violation is sourced only by the imaginary part of the lepton-neutrino Yukawa couplings in both sectors. We have then

\[ \epsilon_l = \frac{\epsilon_D}{\epsilon_D} = \frac{\text{Im}[3((y^\dagger y)_{12}^*)^2 + 2\alpha_1\alpha_2(y^\dagger y)_{12}^*]}{2\alpha_1\alpha_2 \text{Im}[(y^\dagger y)_{12}^*]} = 1 + \frac{\text{Im}[3((y^\dagger y)_{12}^*)^2]}{2\alpha_1\alpha_2 \text{Im}[(y^\dagger y)_{12}^*]} \].

(21)

We see that in the particular case when \( (y^\dagger y)_{12}^* \) is purely imaginary or when \( \alpha_1\alpha_2 \gg |(y^\dagger y)_{12}| \), 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. From the matrix in Eq. (11), we have that

\[ (y^\dagger y)_{12} = y_{ee}(y_{e\mu}^R + iy_{e\mu}^I) + y_{\mu e}(y_{\mu\mu}^R + iy_{\mu\mu}^I) + y_{\tau e}(y_{\tau\mu}^R + iy_{\tau\mu}^I) \]

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. (11)) and also in the same range of the parameters value as given in Eq. (14), one can easily obtain the three neutrino mixing angles and the two mass square differences in their observed 3\( \sigma \) ranges. For this particular choice of parameters, one can easily estimate the value of the lepton CP asymmetry parameter \( (\epsilon_l) \). In this case \( \epsilon_l \) takes the following form,

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

(23)
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 \( y_{ij}^{R,I} \) 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 \( \epsilon_D \) 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.

Hence, from now onwards, we will treat \( \alpha_1 \) and \( \alpha_2 \) as real parameters and assume that the imaginary part needed for CP violation in the Dark sector arises from the RH neutrino sector. Moreover, note that we can take large value of \( \alpha_2 \) \( (\sim \mathcal{O}(1)) \) without violating any constraint, but a large value of \( \alpha_1 \) will violate the narrow width approximation \( (\Gamma_{N_1} \ll M_{N_1}) \) via Eq. (20). Therefore, remaining within the narrow width approximation, but at the same time aiming to increase the CP asymmetry parameter \( \epsilon_l \) for the production of sufficient lepton asymmetry, we vary the two parameters \( \alpha_1 \) and \( \alpha_2 \) in the following ranges:

\[
10^{-3} \leq \alpha_1 \leq 10^{-2},
0.3 \leq \alpha_2 \leq 1.0.
\] (24)

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 \( T_{sph} \sim 150 \text{ GeV} \) \([78, 79]\) 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 \([80]\)

\[
Y_B = \frac{8N_f + 4N_{\phi h}}{22N_f + 13N_{\phi h}} Y_{\Delta l},
\] (25)

where \( N_f \) is number of generations of quarks and leptons while \( N_{\phi h} \) 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

\(^3\) If the light neutrino mass matrix is real, the Dirac phase \( \delta \) 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.
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. 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 [81] 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 [82]

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

(26)

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 [83].

Assuming 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 [81],

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

(27)

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]

$$0.1172 \leq \Omega h^2 \leq 0.1226.$$  

(28)

Equating Eq. (27) and Eq. (28) one can find the mass range for the Dark Matter particle $M_{DM}$ which reproduces the observed Dark Matter abundance.

Let us finally comment on the presence of the inert doublet $\eta_D$, which is also produced in the RH neutrino decay and, as it is heavier, could regenerate a $\psi_1$ population after freeze-out and after the $SU(2)_D, SU(2)_{EW}$ gauge symmetries are broken, by decaying into a DM fermion, a lepton and a Higgs. The number density of $\eta_D$ can be efficiently reduced by the cubic interaction with $\phi_D$, which allows for semiannihilation [84] through the process $\eta_D + \eta_D \rightarrow \eta_D^* + h_{1,2}$. In this case semiannihilation is determined by the potential parameter $\lambda_{D3}$, 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.

V. RESULTS

In Fig. 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 $3.5 \times 10^{-9}$, we can provide for a dark sector asymmetry which fulfils the total Dark Matter 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 $\epsilon_l$) 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. 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 Fig.2 we show the variation of the Dark Matter yield with $\epsilon_D$. All the points satisfy the neutrino oscillation data. The figure shows a sharp correlation between $\epsilon_D$ 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. (27). 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 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 Fig. 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 $\frac{y_{e}}{M_{N_{1}}}$ is fixed by neutrino oscillations, as it provides one of the mass scales in the light neutrino mass matrix. Similar correlations of $M_{N_i}$ 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. (11))
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.

In Fig. 3 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 Fig. 3 shows explicitly the correlation between the real and imaginary part of the same element of the Dirac mass matrix, i.e. $y^R_{\mu\mu}$ and $y^I_{\mu\mu}$. Either of them can give 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

Figure 3: LP (RP): Scatter plot in the $y_{ee} - M_{N_1}$ ($y^R_{\mu\mu} - M_{N_2}$) plane after satisfying neutrino oscillation data as mentioned in section III. All the parameters have been varied in the range as shown in section III.
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 $\epsilon_l$. Since there exists a sharp correlation between $Y_B$ and $\epsilon_l$, lower $\epsilon_l$ results in production of lower lepton asymmetry. As we discussed earlier, we expect $\epsilon_D$ to be equal or less than $\epsilon_l$ 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 $\epsilon_D$ 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 be vary compared to the SM Higgs VEV.

We would like to finally to comment on the case when the second dark fermion contributes substantially to the Dark Matter energy density. Of course we would expect this to happen for similar parameters as the mass of the two states is determined by the same VEV $v_D$. Indeed for $\epsilon'_D \ll \epsilon_D$ and/or $Br'_D \ll Br_D$ (where the primed quantities are for the second fermion), such a contribution would result in negligible contribution from the second fermion. For equal contributions, the total Dark Matter density is realised in the model for half the value of the DM mass, while the overall scenario remains qualitatively the same.

VI. DIRECT DETECTION OF DARK MATTER

In our model, although DM only has a dark $SU(2)_D$ 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 Fig. 6.

![Feynman diagram](image)

Figure 6: 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 $h_2$.

The expression for the spin-independent scattering cross section between DM and nucleon
Figure 7: Variation of $\sigma_{SI}$ with DM mass where the scalar mixing angle $\zeta$ 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 $\lambda_1 = 1.0$.

mediated by scalars $h_1, h_2$ is given by

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

(29)

where $\mu_{red} = \frac{M_N M_{DM}}{M_N + M_{DM}}$ is the reduced mass and $f_N \sim 0.3$ [85]. The DM couplings with the scalars have the following form

$$g_{\psi_1 \psi_1 h_1} = -\frac{\lambda_1}{\sqrt{2}} \sin \zeta = -\frac{m_{DM}}{v_D} \sin \zeta$$
$$g_{\psi_1 \psi_1 h_2} = \frac{\lambda_1}{\sqrt{2}} \cos \zeta = \frac{m_{DM}}{v_D} \cos \zeta,$$

(30)

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

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

(31)

The masses $M_{h_1}, M_{h_2}$ and mixing angle $\zeta$ are given in Eqs. (7), (6) while $\lambda_1$ is the Yukawa coupling between $\Psi_{1L}, \phi_D$ and $\psi_{1R}$, related as well to the Dark Matter mass by Eq. (9).
In Fig. 7 we plot the variation of the spin-independent DM-nucleon scattering cross section with the DM mass. In generating the plot, we have kept $\lambda_1$ 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 $\zeta$ has been scanned over its present allowed range i.e. $\zeta \leq 10^{-1}$ rad. From the Fig. 6 we see that a part of the parameter space, corresponding to large $\zeta$ and heavy $M_{h_2}$, is already ruled out by the null results at various direct detection experiments like PandaX-II [58] and Xenon1T [57], however enough parameter space is still open and can be tested in the future experiments like Darwin [59].

**VII. CONCLUSION**

In this work we have studied a relatively minimal model to solve three major puzzles of cosmology by the presence of two RH neutrinos and a Dark Sector charged under an $SU(2)_D$. We generated 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 bounds on the sum of the light neutrino masses. We showed that the limited number of parameters in the model result 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 $\epsilon$ for leptons and the Dark Matter are quite natural. 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. 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 $\zeta$. 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. [86, 87].

VIII. ACKNOWLEDGEMENTS

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Appendix 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 $m_{\nu}$ (using Eq. 10) in terms of the Yukawa couplings and the RH neutrino masses.

\[
(\tilde{m}_{\nu})_{11} = \frac{y_{ee}}{M_{N_1}} + \frac{(y_{eR})^2 - (y_{eI})^2}{M_{N_2}} - i \frac{2 y_{eR} y_{eI}}{M_{N_2}},
\]

\[
(\tilde{m}_{\nu})_{12} = \frac{y_{ee} y_{\mu e}}{M_{N_1}} + \frac{y_{eR} y_{\mu R} - y_{eI} y_{\mu I}}{M_{N_2}} - i \frac{y_{eR} y_{\mu I} + y_{eI} y_{\mu R}}{M_{N_2}},
\]

\[
(\tilde{m}_{\nu})_{13} = \frac{y_{ee} y_{\tau e}}{M_{N_1}} + \frac{y_{eR} y_{\tau R} - y_{eI} y_{\tau I}}{M_{N_2}} - i \frac{y_{eR} y_{\tau I} + y_{eI} y_{\tau R}}{M_{N_2}},
\]

\[
(\tilde{m}_{\nu})_{21} = (\tilde{m}_{\nu})_{12},
\]

\[
(\tilde{m}_{\nu})_{22} = \frac{y_{\mu e}}{M_{N_1}} + \frac{(y_{\mu R})^2 - (y_{\mu I})^2}{M_{N_2}} - i \frac{2 y_{\mu R} y_{\mu I}}{M_{N_2}},
\]

\[
(\tilde{m}_{\nu})_{23} = \frac{y_{\mu e} y_{\tau e}}{M_{N_1}} + \frac{y_{\mu R} y_{\tau R} - y_{\mu I} y_{\tau I}}{M_{N_2}} - i \frac{y_{\mu R} y_{\tau I} + y_{\mu I} y_{\tau R}}{M_{N_2}},
\]

\[
(\tilde{m}_{\nu})_{31} = (\tilde{m}_{\nu})_{13},
\]

\[
(\tilde{m}_{\nu})_{32} = (\tilde{m}_{\nu})_{23},
\]

\[
(\tilde{m}_{\nu})_{33} = \frac{y_{\tau e}}{M_{N_1}} + \frac{(y_{\tau R})^2 - (y_{\tau I})^2}{M_{N_2}} - i \frac{2 y_{\tau R} y_{\tau I}}{M_{N_2}},
\]

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

We see from these expressions that if the second column of the Yukawa matrix is imaginary, i.e. for $y_{eI}^R = 0$, 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:

\[
m_3 = -\frac{v^2}{2} \sum_i \frac{y_{eI}^2}{M_i},
\]

\[
m_2 = \frac{v^2}{2} \sum_i (y_{\mu I}^2)^2,
\]

\[
m_1 = 0.
\]

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 $m_\nu$, 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 $\xi_3 = \pi/2, \xi_2 = 0$, while $\xi_1$ is undetermined
and can be chosen zero. 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:

$$m_{\beta\beta} = \sum_i m_i U_{ei}^2 = |m_3 \sin^2 \theta_{13} - m_2 \cos^2 \theta_{13} \sin^2 \theta_{12}| \sim 10^{-2} \text{eV}.$$  \text{(A6)}

**Appendix B: Expression of $\gamma_{D_1}$, $\gamma_{\phi_{h,s}}^1$ and $\gamma_{\phi_{h,t}}^1$**:

Expression of $\gamma_{D_1}$ takes the following form \([78]\),

$$\gamma_{D_1} = n_{N_1}^{eq} \frac{K_1(Z)}{K_2(z)} \Gamma_{N_1},$$  \text{(B1)}

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 $\Gamma_{N_1}$ is the total decay width of $N_1$. The expression
of $\Gamma_{N_1}$ is given in Eq. \((20)\).

Further, the general expression of $\gamma(a+b \leftrightarrow i + j + ...)$ for a two body scattering process
$a+b \leftrightarrow i + j + ...$ is given by \([78]\),

$$\gamma(a+b \leftrightarrow i + j + ...) = \frac{T}{64\pi^4} \int_{(M_a+M_b)^2}^{\infty} ds \hat{\sigma}(s) \sqrt{s} K_1 \left(\frac{\sqrt{s}}{T}\right),$$  \text{(B2)}

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 $a+b \leftrightarrow i + j + ...$, which is related to the actual cross
section by the following relation

$$\hat{\sigma}(s) = \frac{8}{s} \left[ (p_a.p_b)^2 - M_a^2 M_b^2 \right] \sigma(s).$$  \text{(B3)}

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 $\phi_h$)
and $N_1 + t \rightarrow \bar{t} + q$ ($t$-channel process mediated by $\phi_h$) are given as

$$\hat{\sigma}_{\phi_{h,s}} = \frac{3\pi\alpha^2 M_t^2}{M_W^4 \sin^4 \theta_w} (M_D^4 M_D)_{11} \left[ \frac{s - M_{N_1}^2}{s} \right]^2,$$

$$\hat{\sigma}_{\phi_{h,t}} = \frac{3\pi\alpha^2 M_t^2}{M_W^4 \sin^4 \theta_w} (M_D^4 M_D)_{11} \left[ \frac{s - M_{N_1}^2}{s} + \frac{M_{N_1}^2}{s} \ln \left( \frac{s - M_{N_1}^2 + M_{h_1}^2}{M_{h_1}^2} \right) \right].$$  \text{(B4)}
where $\alpha = \frac{g^2 \sin^2 \theta_w}{4\pi}$, $g$ being the SU(2)$_L$ gauge coupling and $\theta_w$ is the weak mixing angle (Weinberg angle).

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