Scalar dark matter, leptogenesis and \(0\nu\beta\beta\) in minimal scotogenic model

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Abstract

We study the minimal scotogenic model constituting an additional inert Higgs doublet and three sets of right-handed neutrinos. This model connects dark matter, baryon asymmetry of the Universe, neutrinoless double beta decay \(0\nu\beta\beta\) in light of the latest experimental data. In view of the recent constraints from Planck data, baryogenesis is obtained for TeV scale heavy neutral singlet fermion \((N_1)\) with the active neutrino masses satisfying experimental bound from KamLAND-Zen. We primarily focus on the intermediate-mass region of dark matter within \(M_W < M_{DM} \leq 550\) GeV, where observed relic density is suppressed due to co-annihilation processes. We consider thermal as well as the non-thermal approach of dark matter production and explore the possibility of the lightest stable candidate being a dark matter candidate. Within the IHDM desert, we explore a new allowed region of dark matter masses for the non-thermal generation of dark matter with a mass splitting of 10 GeV among the inert scalars, keeping intact constraints from Planck limit as well as direct detection experiment XENON1T.

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I. INTRODUCTION

The Standard Model (SM) of particle physics is an affluent and self-consistent one in the current scenario. However, it is not accountable for explaining various problems persisting in the Universe. Among all the anomalies, baryon asymmetry of the Universe [1, 2], absolute neutrino mass [3], dark matter [4, 5] are the ones drawing much of the attention in the ongoing studies nowadays.

Successively, there has been significant growth in the past few years in providing pieces of evidence to these mysterious and yet interesting form of non-baryonic matter, commonly termed as dark matter (DM) in the present Universe. The significant lines of evidence of DM include observations in galaxy cluster by Fritz Zwicky [6] in 1933, gravitational lensing (which could allow galaxy cluster to act as gravitational lenses as postulated by Zwicky in 1937) [7], galaxy rotation curves in 1970 [8], cosmic microwave background [9] and the most recent cosmology data given by Planck satellite [10] are some of the most remarkable ones. From the recent Planck satellite data, it is certain that approximately 27% of the present Universe is comprised of DM, which is about five times more than the baryonic matter. The conditions required to be fulfilled by particle candidates for DM is found in this paper [11], from which it is confirmed that the possibility of SM particle to be a DM candidate is ruled out. This has resulted in the extension of the SM, of which the weakly interacting massive particle (WIMP) paradigm is the most discussed framework.

A notable co-occurrence frequently termed as the WIMP miracle [12] is feasible in the WIMP paradigm, where a dark matter candidate typically with an electroweak scale mass and electroweak alike interactions can produce correct dark matter relic abundance. WIMPs can be thermally produced in the early Universe as the interactions governing them are of electroweak scale. Thus, relic abundance of a thermal DM candidate can be generated while the interactions freeze out, ensuing the expansion as well as the cooling of the Universe. Also, the WIMP paradigm foretells the observable DM nucleon scattering cross-section through the same interactions that were operational at the time of freeze-out. However, many dark matter direct detection experiments like LUX [13], PandaX-II [14], and XENON1T [15] have reported their null results. Therefore, the exclusion curve in the mass-cross section plane is lowered. Similar null results have been obtained from the Large Hadron Collider (LHC), which further gives an upper bound on the DM interaction with the SM particles. A strict constraint on the WIMP parameter space can be summarized from the different null results.
Besides DM, the baryon asymmetry of the Universe is another puzzle, which is the observed imbalance in the baryonic matter and anti-baryonic matter in the observable Universe. A particle to create baryon asymmetry, it must satisfy the Sakharov conditions [16], which demands baryon number (B) violation, C and CP violation, and departure from thermal equilibrium. As these conditions cannot be fulfilled within the SM in an adequate amount, we need formalism beyond the SM. Of these criteria, the out-of-equilibrium decay of a heavy particle leading to the generation of the baryon asymmetry of the Universe (BAU) has so far been a widely known mechanism for baryogenesis [17, 18]. We can incorporate such a mechanism via leptogenesis [19], where a net leptonic asymmetry is generated first, which further gets converted into baryogenesis through $(B + L)$ violating electroweak sphaleron phase transitions [20]. A rich literature is available for various leptogenesis processes [21–26]. In the case of an elementary scenario, mostly referred to as vanilla leptogenesis, where the lower mass bound, by the allowance of flavor effect, comes down to be about $M_{1}^{min} = 10^8$ GeV [27, 28]. Owing to the fact that the CP asymmetry in RHN decays is a consequence of the active and sterile neutrino masses along with the necessity of tiny SM neutrino masses, the high mass scale of RHN is needed [29, 30]. Nevertheless, such a high mass scale of RHN is disagreeable for several reasons, of which, the mere possibility of detecting the dynamics of leptogenesis in future collider experiments [31] is of much significance. Another reason is that the high-scale leptogenesis may be precluded due to the future detection of lepton number violation at low energies [32]. Thus, these observations act as a catalyst to opt for other alternatives to the archetype of standard thermal leptogenesis in the type-I seesaw mechanism that copes up to produce the BAU at much lower RHN mass order. Implementing the idea of low mass RHN as mentioned in [27], the study of thermal leptogenesis in Ernest Ma’s scotogenic model [33, 34], which is considered to be the simplest model of radiative neutrino masses is carried out in this work. The scotogenic model is of much significance as we can relate the light SM neutrino mass with the physics of dark matter [33].

This work primarily focuses on the inert Higgs doublet model (IHDM) desert, i.e., $M_{W} < M_{DM} \leq 550$ GeV, wherein the generation of the relic abundance is prohibited as mentioned in various literatures [35, 36]. The core reason behind this discrepancy is that in the IHDM desert, the annihilation cross-section of the dark matter is large compared to the amount necessary to produce the correct relic abundance via the freeze-out mechanism. Thus, we get an underabundant DM
in this regime due to the large annihilation rates. Though the lower bound of the IHDM desert is rigid, the upper bound can be a little flexible depending on the choice of parameters such as the DM-Higgs coupling and the mass splitting between the inert scalars. Thus, we try to see the viability of the IHDM desert, concentrating on the upper bound satisfying the relic abundance with the latest restriction from the direct detection experiment XENON1T [15]. The production of a correct relic in this regime can be possible by fine-tuning of the DM-Higgs coupling and suitable mass splitting of the other inert scalars.

Motivated by these factors, in this model, the SM is extended by a Higgs doublet field ($\eta$) and three singlet neutral fermions ($N_k$), which are odd under $Z_2$ symmetry, in contradiction to the SM particles which are $Z_2$ even. The possibility of a DM candidate comes from the $Z_2$ odd lightest particle. Whereas, leptogenesis is a result of the $Z_2$ odd fermions, i.e., the heavy RHN, which occurs via the out-of-equilibrium decay into the SM leptons and the inert Higgs doublet [37]. The entire work is carried out keeping the dark matter mass in the intermediate dark matter mass range, also known as IHDM desert, which lies between $M_W < M_{DM} \leq 550$ GeV. Leptogenesis is obtained for this very range of dark matter mass with the heavy RHN on a low mass scale. Also, an important criterion that is kept intact while generating the TeV scale leptogenesis is the sum of neutrino masses and its effective mass being consistent with the constraints from Planck data and neutrinoless double beta decay experiment, KamLAND-Zen. We also check the relic abundance of the dark matter candidate (lightest of $\eta$) for different choices of mass splitting between the scalars of the inert scalar doublet. We further investigate the parameter space, i.e. the values of DM-Higgs coupling and dark matter mass for which it satisfies the bounds from relic abundance and direct detection experiment. Furthermore, we also study the mixture of thermal and non-thermal production of DM abundance for various masses within the IHDM desert. In one of the cases, we have considered mass splitting of the scalars in the inert doublet to be 10 GeV and studied the criteria that satisfy the observed relic for higher DM masses within the IHDM desert via purely thermal production as well as non-thermal production. The rest of the paper is divided into five sections, where section(II) includes a brief introduction of the scotogenic model involving the generation of neutrino mass. Section(III) and section(IV) constitutes discussions on baryogenesis in scotogenic model and neutrinoless double beta decay, respectively. Thermal and non-thermal production of dark matter is discussed in section(V). A detailed numerical analysis, along with
results, is shown in section (VI) followed by the conclusion given in section (VII).

II. SCOTOGENIC MODEL

Scotogenic model is an extension of the Inert Higgs Doublet Model (IHDM) [37] and the IHDM is nothing but a minimal extension of the SM by a Higgs field which is a doublet under $SU(2)_L$ gauge symmetry with hypercharge $Y = 1$ and a built-in discrete $Z_2$ symmetry [35, 36, 38–48]. The necessity of this modification took place as the IHDM could only accommodate dark matter, whereas it failed in explaining the origin of neutrino masses at a renormalizable level [46]. In this model, we add three neutral singlet fermions $N_i$ with $i = 1, 2, 3$ in order to generate the neutrino masses and assign them with a discrete $Z_2$ symmetry. In view of $N_i$, the neutrinos can get masses in two ways. One of the ways is similar to the type-I seesaw mechanism [22–26], where the neutrino masses arise as a result of $N_i$ being $Z_2$ even. Also, it is limited to show no dark matter phenomenology of the IHDM and keeps the neutrino masses decoupled from the DM characteristics. Therefore, we opt for the other way in which $N_i$ is odd under $Z_2$ symmetry, whereas the SM fields remain $Z_2$ even. Symbolic transformation of the particles under $Z_2$ symmetry is given by,

$$N_i ightarrow -N_i, \eta \rightarrow -\eta, \Phi \rightarrow \Phi, \Psi \rightarrow \Psi,$$

(2.1)

where $\eta$ is the inert Higgs doublet, $\Phi$ is the SM Higgs doublet and $\Psi$ denotes the SM fermions. The new leptonic and scalar particle content can thereafter be represented as follows under the group of symmetries $SU(2) \times U(1)_Y \times Z_2$:

$$\begin{pmatrix} \nu_{\alpha} \\ l_{\alpha} \end{pmatrix}_L \sim (2, -\frac{1}{2}, +), \kappa_{\alpha} \sim (1, 1, +), \begin{pmatrix} \Phi^+ \\ \Phi^0 \end{pmatrix} \sim (2, \frac{1}{2}, +), N_i \sim (1, 1, -),$$

$$\begin{pmatrix} \eta^+ \\ \eta^0 \end{pmatrix} \sim (2, \frac{1}{2}, -).$$

(2.2)

The scalar doublets are written as follows:

$$\eta = \begin{pmatrix} \eta^\pm \\ \eta^0_R + i\eta^0_I \end{pmatrix}, \Phi = \begin{pmatrix} \Phi^+ \\ h + i\xi \end{pmatrix}.$$  

(2.3)
We have no Dirac mass term with $\nu$ and $N$; however, the similar Yukawa-like coupling involving $\eta$ is allowed. Nevertheless, the scalar cannot get a VEV. The neutrino mass can be generated through a one-loop mechanism, which is based on the exchange of $\eta$ particle and a heavy neutrino. In Fig. [1] we see two Higgs fields $\phi^0$ are involved. They will not propagate but will acquire VEV after the EWSB.

The lagrangian involving the newly added field is:

$$\mathcal{L} \supset \frac{1}{2} (M_N)_{ij} N_i N_j + Y_{ij} \bar{L} \eta N_j + h.c$$

where, the 1st term is the Majorana mass term for the neutrino singlet and the 2nd term is the Yukawa interactions of the lepton. The new potential on addition of the new inert scalar doublet is:

$$V_{Scalar} = m_2^{2} (\Phi^{+} \Phi) + m_2^{2} \eta^{+} \eta + \frac{1}{2} \lambda_1 (\Phi^{+} \Phi)^2 + \frac{1}{2} \lambda_2 (\eta^{+} \eta)^2 + \lambda_3 (\Phi^{+} \Phi)(\eta^{+} \eta) + \lambda_4 (\Phi^{+} \eta)(\eta^{+} \Phi) + \frac{1}{2} \lambda_5 [(\Phi^{+} \eta)^2 + h.c.]$$

All the parameters in Eq. (2.5) are real by hermicity of the Lagrangian, except for $\lambda_5$. Since, the bilinear term $(\Phi^{+} \eta)$ is forbidden by the exact $Z_2$ symmetry, therefore one can always choose $\lambda_5$ real by rotating the relative phase between $\Phi$ and $\eta$. Furthermore, after the spontaneous symmetry breaking like in SM, we are left with one physical Higgs boson $h$ which resembles the SM Higgs boson, as well as four dark scalars: one CP even($\eta_R^0$), one CP odd($\eta_I^0$) and a pair of charged ones.
$(\eta^\pm)$. The masses of these physical scalars are:

\begin{align}
    m_h^2 &= -m_1^2 = 2\lambda_1 v^2, \\
    m_{\eta^\pm}^2 &= m_2^2 + \lambda_3 v^2, \\
    m_{\eta^0_R}^2 &= m_2^2 + (\lambda_3 + \lambda_4 + \lambda_5) v^2, \\
    m_{\eta^0_I}^2 &= m_2^2 + (\lambda_3 + \lambda_4 - \lambda_5) v^2. \\
\end{align}

(2.6)

It is clear from the above equations that all the scalar couplings are written in terms of physical scalar masses and $m_2$, thereby providing six independent parameters of the model to be: \( \{m_2, m_h, m_{\eta^0_R}, m_{\eta^0_I}, m_{\eta^\pm}, \lambda_2\} \). Here, $m_h$ is the mass of SM-Higgs, $m_{\eta^0_R}, m_{\eta^0_I}$ and $m_{\eta^\pm}$ are the masses of CP-even, CP-odd and charged scalars of the inert doublet respectively. In this work, as we have considered the CP-even scalar to be the lightest particle and a probable DM candidate, so we consider $\lambda_5 < 0$ without any loss of generality. Also, the limit $\lambda_5 \rightarrow 0$ leads to the mass degeneracy of the neutral components of the inert doublet. Following the ‘t Hooft scenario [49], the smallness of $\lambda_5$ to obtain the lepton asymmetry, which would have been lost if considered to be zero, is acceptably natural. We have a simplified diagram that can be split further into two diagrams and from which the mass can be easily calculated by considering mechanism after EWSB. Calculation on the basis of one diagram is sufficient and considered as other would be same except

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{diagram.png}
\caption{One-loop diagram with exchange of $\eta^0_R$ and $\eta^0_I$. $\nu_i$ and $\nu_j$ representing two different generations of active neutrinos. $N_k$ is the right handed neutrino.}
\end{figure}

for $\eta^0_R$ replaced by $\eta^0_I$. The neutrino mass matrix arising from the radiative mass model is given by [41, 50]:

\begin{align}
    M_{ij}^\nu &= \frac{1}{16\pi^2} \sum_k h_{ik} h_{jk} M_k \left[ \frac{m_{\eta^0_R}^2}{m_{\eta^0_R}^2 - M_k^2} \ln \frac{m_{\eta^0_R}^2}{M_k^2} - \frac{m_{\eta^0_I}^2}{m_{\eta^0_I}^2 - M_k^2} \ln \frac{m_{\eta^0_I}^2}{M_k^2} \right] \\
    &\equiv \frac{1}{16\pi^2} \sum_k h_{ik} h_{jk} M_k L_k(m_{\eta^0_R}^2) - L_k(m_{\eta^0_I}^2), \\
\end{align}

(2.7)
where $M_k$ represents the mass eigenvalue of the mass eigenstate $N_k$ of the neutral singlet fermion $N_k$ in the internal line with indices $j=1,2,3$ running over the three neutrino generation with three copies of $N_k$. The function $L_k(m^2)$ used in Eq. (2.7) is given by:

$$L_k(m^2) = \frac{m^2}{m^2 - M_k^2} \ln \frac{m^2}{M_k^2}$$  \hspace{1cm} (2.8)

In our study, we calculate the Yukawa couplings by the incorporation of the constraints on the sum of neutrino masses [51] and the neutrino oscillation data [52]. For simplicity of the Yukawa coupling calculation, we write the mass formula given by Eq. (2.7), in the form similar to type-I seesaw formula [53]:

$$M_\nu = Y \Lambda^{-1} Y^T,$$  \hspace{1cm} (2.9)

where $\Lambda$ is a diagonal matrix represented by:

$$\Lambda_k = \frac{2\pi^2}{\lambda_5} \frac{2M_k}{\bar v^2},$$  \hspace{1cm} (2.10)

with,

$$s_k = \left(\frac{M_k^2}{8(m^2_{\nu_R} - m^2_{\nu_I})}[L_k(m^2_{\nu_R}) - L_k(m^2_{\nu_I})]\right)^{-1}$$  \hspace{1cm} (2.11)

The light neutrino mass matrix (2.7) can be diagonalised by an unitary matrix known as the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix.

The diagonal light neutrino mass matrix can be written as:

$$M_\nu^{\text{diag}} = U^\dagger M_\nu U^*$$  \hspace{1cm} (2.12)

Also, we use a special yet one of the most popular types of parametrization known as the Casas-Ibarra parametrization [74] in order to link the Yukawa coupling with the light neutrino parameters.

$$Y = U \sqrt{M_\nu^{\text{diag}}} R^\dagger \sqrt{\Lambda},$$  \hspace{1cm} (2.13)

where $R$ is an complex orthogonal matrix satisfying the condition $R^\dagger R = 1$. We also parameterized the $R$ matrix as per our convenience and the orthogonal complex matrix $R$ takes the form,

$$R = \frac{\begin{pmatrix} 0 & \cos Z & \sin Z \\ 0 & -\sin Z & \cos Z \\ 1 & 0 & 0 \end{pmatrix}}{2 \sqrt{\cos^2 Z + \sin^2 Z}}.$$  \hspace{1cm} (2.14)
where, $Z = (z_R + iz_I)$ with $z_R$ and $z_I$ bearing the values 1.42 and 1.6232 respectively for normal hierarchy (NH). In the case of inverted hierarchy, we arbitrarily choose lower values of $z_R = 0.22$ and $z_I = 0.58$, which contributes to a slight difference in the baryogenesis plot as a function of RHN $N_1$. This choice of the orthogonal matrix $R$ is made to calculate the Yukawa couplings related by the Casas- Ibarra parametrization given in Eq.(2.13), in order to obtain a non-zero complex term for $(Y^\dagger Y)_{11}$ which is inversely proportional to the CP asymmetry $\epsilon_1$. Since $\epsilon_1$ is directly dependent on $(Y^\dagger Y)_{12}$ and $(Y^\dagger Y)_{13}$ as well, the requirement of these quantities to be non-zero is a must. Therefore, such a choice of $R$ as in Eq.(2.14) is adequate in fulfilling the foresaid criteria.

III. BARYOGENESIS IN SCOTOGENIC MODEL

A fascinating way to dynamically produce the Baryon asymmetry of the Universe (BAU) in near the beginning of the Universe is baryogenesis via leptogenesis [55–58]. There arises an intrinsic limitation of the standard thermal leptogenesis, which is due to the requirement of a very high right-handed neutrino (RHN) mass scale. In the most generic scenario, occasionally known as the vanilla leptogenesis, there exists an absolute lower bound on the mass of the lightest RHN to be $M_1 \simeq 10^9$ GeV [59, 60]. Whereas, in the case of the scotogenic model, with three $Z_2$ odd SM singlet fermions, one can bring down the limit on the lightest RHN mass scale to be as low as 10 TeV [27]. In our work, we have taken the lightest RHN mass scale of the range $10^4 - 10^5$ GeV, and that of the heavier RHNs $N_2$ and $N_3$ of the range $10^7 - 10^8$ GeV and $10^9 - 10^{10}$ GeV respectively for generating the required baryogenesis. If kinematically allowed via the Yukawa interactions, the SM singlet neutral fermions decay into the SM leptons, and the inert Higgs doublet $\eta$. The final lepton asymmetry is generated only because of the asymmetry created by $N_1$ decays, which are the most pertinent compared to that produced by decays of $N_2$ and $N_3$. This leptogenesis is further converted into the baryon asymmetry of the Universe (BAU) by the electro-weak sphaleron phase transition [61]. The simultaneous Boltzmann equations for $N_1$ decay and formation of $N_{B-L}$ are to be solved to obtain the results for baryogenesis. The B-L calculation is mainly governed on the comparison between the Hubble parameter and the decay rates for $N_1 \rightarrow l\eta, \bar{\eta}^*$ processes which will have a certain impact on the asymmetry, as well as on the CP-asymmetry parameter $\epsilon_1$. We now further look into the various expressions and quantities that are required for the
calculation of thermal leptogenesis in the scotogenic model. As essential in thermal leptogenesis, we need to distinguish between a weak washout and a strong washout regime. The differentiation is characterized based on the values of the decay parameter,

\[ K_1 = \frac{\Gamma_1}{H(z_1 = 1)}, \]  

(3.1)

where, \( \Gamma_1 \) is the total \( N_1 \) decay width, \( H \) being the Hubble parameter and \( z_1 = \frac{M_1}{T} \) with temperature \( T \) of the photon bath. Leptogenesis occurs above the electroweak scale during the era of radiation domination. The Hubble parameter can therefore be expressed in terms of \( T \) as follows:

\[ H = \sqrt{\frac{8\pi^3 g_* T^2}{90 M_{Pl}}}, \]  

(3.2)

where \( g_* \) is the effective number of relativistic degrees of freedom and \( M_{Pl} \approx 1.22 \times 10^{19} \text{ GeV} \) is the Planck mass. With the varied choice of parameters, i.e., \( M_{N_1}, M_{DM} \) and most crucially value of the lightest active neutrino mass in the range \( m_l = 10^{-3} - 1 \text{ eV} \) compels the 3RHN scenario to fall in the strong washout regime similar to 2RHN case or type-I leptogenesis\[27\]. Thus we get a significantly large value of \( K_1 \), which is \( K_1 \simeq 10^3 \) and above. This further enables us to assume \( N_1 \) dominated leptogenesis and neglect washout via scattering effects. Thereby, for a large value of \( K_1 \), we can use the approximation for the efficiency factor in the strong washout regime as,

\[ \kappa_1(K_1) \approx \frac{1}{1.2K_1[\ln K_1]^{0.8}}. \]  

(3.3)

The \( N_1 \) decay rate incorporating the Yukawa coupling is given by,

\[ \Gamma_1 = \frac{M_1}{8\pi} (Y^\dagger Y)_{11} \left[ 1 - \left( \frac{m_{DM}}{M_1} \right)^2 \right]^2 = \frac{M_1}{8\pi} (Y^\dagger Y)_{11}(1 - \eta_1)^2. \]  

(3.4)

The CP asymmetry parameter \( \epsilon_1 \) for the decays \( N_1 \rightarrow l\eta, \bar{l}\eta^* \) is given by,

\[ \epsilon_1 = \frac{1}{8\pi (Y^\dagger Y)_{11}} \sum_{j \neq 1} \text{Im}[(Y^\dagger Y)^2]_{1j} \left[ f(r_{j1}, \eta_1) - \frac{\sqrt{r_{j1}}}{r_{j1} - 1}(1 - \eta_1)^2 \right], \]  

(3.5)

where, the term \( f(r_{j1}, \eta_1) \) is expressed as,

\[ f(r_{j1}, \eta_1) = \sqrt{r_{j1}} \left[ 1 + \frac{(1 - 2\eta_1 + r_{j1})}{(1 - \eta_1)^2} \ln \left( \frac{r_{j1} - \eta_1^2}{1 - 2\eta_1 + r_{j1}} \right) \right], \]  

(3.6)
with \( r_{j1} = \frac{M_j}{M_1} \), \( \eta_1 \equiv \left( \frac{m_{DM}}{M_1} \right)^2 \). The frequently appearing \( Y^\dagger Y \) in the above equations can be expressed using the CI-parametrization\[54]\,

\[
(Y^\dagger Y)_{ij} = \sqrt{\Lambda_i \Lambda_j} (RM_{\nu}^{\text{diag}} R^\dagger)_{ij}.
\] (3.7)

An exciting piece of information regarding the \( Y^\dagger Y \) is that it is independent of the PMNS matrix. This ensures that the CP-violating phases applicable for leptogenesis is independent of the CP-violating phases in the PMNS matrix. Again, starting with the initial thermal abundance of \( N_1 \), wherein its rate of interaction is above the Hubble rate, we solve the Boltzmann equations. It is only feasible if the Yukawa couplings corresponding to \( N_1 \) are not very small. In our work, we calculate the Yukawa coupling, which falls in the range applicable to generate the observed baryon asymmetry.

The Boltzmann equations for the number densities of \( N_1 \) and \( N_{B-L} \), given by\[62]\,

\[
\frac{dn_{N_1}}{dz} = -D_1 (n_{N_1} - n_{N_1}^{\text{eq}}),
\] (3.8)

\[
\frac{dn_{B-L}}{dz} = -\epsilon_1 D_1 (n_{N_1} - n_{N_1}^{\text{eq}}) - W_1 n_{B-L},
\] (3.9)

respectively. The equilibrium number density of \( N_1 \) is given by \( n_{N_1}^{\text{eq}} = \frac{z^2}{2} K_2(z) \), where \( K_i(z) \) is the modified Bessel function of \( i^{th} \) type and

\[
D_1 \equiv \frac{\Gamma_1}{Hz} = K_N z \frac{K_1(z)}{K_2(z)}
\] (3.10)

is the measure of the total decay rate with respect to the Hubble rate, and \( W_1 \) is the total washout rate given by \( W_1 = \frac{\Gamma_1}{Hz} \). The total washout term \( W_1 \) is the sum of the washout due to inverse decays \( l\eta, \bar{l}\eta^* \to N_1 \) and the washout due to the \( \Delta L = 2 \) scatterings \( l\eta \leftrightarrow \bar{l}\eta^*, ll \leftrightarrow \eta^*\eta^* \), i.e. \( W_1 = W_{1D} + W_{\Delta L=2} \)\[27\], where \( W_{1D} = \frac{1}{4} K_{N_1} z^3 K_1(z) \) and,

\[
W_{\Delta L=2} \simeq \frac{18\sqrt{10}}{\pi^2 g_t \sqrt{g} z^2 v^4} \left( \frac{2\pi^2}{\lambda_5} \right)^2 M_1 \bar{m}_\chi^2.
\] (3.11)

In Eq.\(3.11\), \( g_t \) stands for the internal degrees of freedom for the SM leptons, and \( \bar{m}_\chi \) is the effective neutrino mass parameter, defined by:

\[
\bar{m}_\chi^2 \simeq 4 \bar{s}_1 m_1^2 + \bar{s}_2 m_2^2 + \bar{s}_3 m_3^2,
\] (3.12)
with \( m_i' \)'s being the light neutrino mass eigenvalues and \( \varsigma_k \) is as defined in Eq.\((2.11)\). We assess the final B-L asymmetry \( n^{f}_{B-L} \) just before sphaleron freeze-out by numerically solving the Eqs.\((3.8)\) and \((3.9)\), which is further converted into the baryon-to-photon ratio as,

\[
n_B = \frac{3}{4} g_*^0 a_{sph} n^{f}_{B-L} \simeq 9.2 \times 10^{-3} n^{f}_{B-L}, \tag{3.13}
\]

where \( a_{sph} = \frac{8}{23} \) is the sphaleron conversion factor with the consideration of two Higgs doublet. \( g_* = 110.75 \) is the effective relativistic degrees of freedom at the time of final lepton asymmetry production, and \( g_*^0 = \frac{43}{11} \) is the effective degrees of freedom at the recombination epoch. In this work, we have studied the effects on leptogenesis by the variation of parameters such as quartic coupling in the range \( 10^{-2} - 10^{-5} \), the probable DM candidate mass in the intermediate-mass regime, i.e., \( M_W < M_{DM} \leq 550 \) GeV. From this choice of parameters, along with the mass of the light neutrino in the range \( 10^{-3} - 1 \) eV, we calculate the Yukawa couplings for which we achieve \( n^{obs}_B \) inferred from the Planck limit 2018, i.e., \( (6.04 \pm 0.08) \times 10^{-10} \) at 68% C.L. \cite{63}. Therefore, we get baryogenesis keeping intact the light neutrino mass satisfying the neutrino oscillation data.

**IV. NEUTRINOLESS DOUBLE BETA DECAY**

With the light neutrino parameters considered in our work, we can make connections with observable in the on-going experiments. A well known and significant experimental technique of detecting neutrino mass is the neutrinoless double beta decay\((0\nu\beta\beta)\), with experiments such as KamLAND-Zen, GERDA, KATRIN. In such experiments, what is measured is the effective neutrino mass \( |m_{\beta\beta}| \) which can be determined by the formula,

\[
|m_{\beta\beta}| = \sum_{k=1}^{3} m_k U^2_{ek} \tag{4.1}
\]

where, \( U^2_{ek} \) are the elements of the PMNS matrix with \( k \) holding up the generation index. This eq.\((4.1)\) can be further expressed as,

\[
|m_{\beta\beta}| = |c_{12}^2 c_{13}^2 m_1 + s_{12}^2 c_{13}^2 m_2 e^{2i\alpha} + s_{13}^2 m_3 e^{2i\beta}| \tag{4.2}
\]

where, \( c_{ij} = \cos \theta_{ij} \) and \( s_{ij} = \sin \theta_{ij} \). It is important to check the satisfying bound of the effective mass with the lightest neutrino mass so that we can relate the current light neutrino parameters giving correct hints to ongoing experiments and their future sensitivity.
V. DARK MATTER IN SCOTOGENIC MODEL

The dark matter, which was in chemical and thermal equilibrium in the early Universe, loses its equilibrium state when the pair annihilation rate becomes less than the expansion rate of the Universe, eventually leading the particles to decouple from the cosmic plasma. The relic densities of such thermally produced dark matter candidates can be calculated by solving the Boltzmann equation [12, 64]:

\[
\dot{n}_{DM} + 3Hn_{DM} = - <\sigma v>(n_{DM}^2 - (n_{DM}^{eq})^2),
\]

where, \(n_{DM}\) is the number density of the dark matter candidate and \(n_{DM}^{eq}\) is the number density of the dark matter candidate in thermal equilibrium. The numerical solution of the Boltzmann equation in terms of partial wave expansion, \(<\sigma v> = a + bv^2\) is of the form,

\[
\Omega h^2 \approx 1.04 \times 10^9 x_f\frac{m_{DM}}{M_{Pl} \sqrt{g^*}}\frac{1}{g_*(a + 3b/x_f)},
\]

where, \(x_f = \frac{m_{DM}}{T_f}\), \(T_f\) is the freeze-out temperature, \(m_{DM}\) is the mass of dark matter, \(g_*\) is the number of relativistic degrees of freedom at the time of freeze-out, and \(M_{Pl} \approx 2.4 \times 10^{18} \text{ GeV}\) is the Planck mass. Furthermore, we can also express this above expression in a simpler analytical form for the approximation of DM relic abundance as [65],

\[
\Omega h^2 \approx 3 \times 10^{-27} \text{ cm}^3 \text{ s}^{-1} \frac{<\sigma v>}{\sqrt{s}}
\]

(5.3)

The corresponding thermal averaged annihilation cross section is therefore given by [66]:

\[
<\sigma v> = \frac{1}{8m_{DM}^4 K_2^2(m_{DM}/T)} \int_{4m_{DM}^2}^{\infty} \sigma(s - 4m_{DM}^2)\sqrt{s}K_1(\sqrt{s}/T)ds,
\]

(5.4)

where, \(K_1\) and \(K_2\) are the modified Bessel functions, \(m_{DM}\) is the mass of dark matter candidate and \(T\) is the temperature. In our model, we have considered one of the neutral component of the scalar doublet \(\eta\), i.e \(\eta^0\) to be the dark matter candidate which resembles that with the inert doublet model discussed in the papers [33, 35, 36, 39–48]. From the literature [67], we can express the effective cross-section as,

\[
\sigma_{eff} = \sum_{i,j}^N <\sigma_{ij}v> \frac{g_ig_j}{g_{eff}^2}(1 + \Delta_i)^{3/2}(1 + \Delta_j)^{3/2}e^{(-x_f(\Delta_i + \Delta_j))},
\]

(5.5)

with, \(\Delta_i = \frac{m_i - m_{DM}}{m_{DM}}\) and \(g_{eff} = \sum_{i=1}^N g_i(1 + \Delta_i)^{3/2}e^{-x_f\Delta_i}.\)
In the above equation, \( m_i \) denotes the mass of the heavier inert Higgs doublet. Therefore, the expression for the thermally averaged cross section is given by

\[
<\sigma_{ij}v> = \frac{x_f}{8m_i^2m_j^2m_{DM}}K_2\left(\frac{m_{ij}}{m_{DM}}\right)K_2\left(\frac{m_{ij}}{m_{DM}}\right) \times \int_{(m_i+m_j)^2}^{\infty} \sigma_{ij}(s-2(m_i^2+m_j^2))\sqrt{s}K_1\left(\frac{\sqrt{s}x_f}{m_{DM}}\right)ds. \tag{5.6}
\]

The only parameters mainly affecting the relic is the DM-Higgs coupling (\( \lambda_L \)) and the mass differences between the inert scalars. By appropriate choice of \( \lambda_L \) and mass splitting, it is possible to generate the correct relic abundance for DM mass around 500 GeV. However, it is impossible to get the observed relic density below 500 GeV of dark matter mass, if the dark matter is produced thermally. Hence, we approach the non-thermal production of dark matter production mechanisms and study its consequences within the IHDM desert.

A non-thermal contribution in the production of relic abundance can be useful in generating the correct relic for masses of dark matter within the IHDM desert. The addition of the non-thermal part can enhance the under-abundant relic, which was observed in the IHDM desert to satisfy the Planck limit. This can be actually achieved by the late decay of the heavy particle, in our case \( N_1 \) decays to DM and SM leptons, i.e. \( N_1 \rightarrow \eta, \bar{\eta}^{*} \), resulting in the production of a correct relic of the DM candidate(\( \eta_R^0 \)). We proceed with the method as discussed in [68], and solve the coupled Boltzmann equations shown below to calculate the number densities of DM candidate and \( N_1 \):

\[
\frac{dn_{DM}}{dt} + 3Hn_{DM} = -<\sigma v>(n_{DM}^2 - (n_{DM}^eq)^2) + N\Gamma_{N_1}n_{N_1}. \tag{5.7}
\]

\[
\frac{dn_{N_1}}{dt} + 3Hn_{N_1} = -\Gamma_{N_1}n_{N_1},
\]

where \( N \) is the average number of DM particles produced on the decay of \( N_1 \), and \( \Gamma_{N_1} \) is the decay width of \( N_1 \). We then move towards the analytical solution of the Boltzmann equation for \( n_{N_1} \) by taking into consideration some of the crucial assumptions, that the co-moving entropy density(\( g_{*s} \)) and co-moving energy density(\( g_* \)) is almost constant. We now transform the above equation interns of \( Y_{DM} \) and \( Y_{N_1} \) by using the relation \( Y_{DM} = \frac{n_{DM}}{s} \) and \( Y_{N_1} = \frac{n_{N_1}}{s} \) where \( s = \frac{2\pi^2g_*T^4}{45} \) is the entropy density. The final equation we obtain on changing the variable \( t \) to \( x = \frac{M_{DM}T}{s} \) and also inserting the above variables:

\[
\frac{dY_{DM}}{dx} = -\frac{<\sigma v>}{Hx}Y_{DM}^2 - (Y_{DM}^eq)^2 + Nr_xY_{N_1}(x_0)exp(-\frac{r}{2}(x^2 - x_0^2)). \tag{5.8}
\]

In eq.(5.8), \( r = \frac{\Gamma_{N_1}}{Hx^2} = \left(\frac{\Gamma_{N_1}M_{Pl}}{\pi M_{DM}^2}\right)\sqrt{\frac{90}{g_*}} \) is a constant depending upon the decay width of the heavy decaying particle and \( Y_{N_1}(x_0) \) is the initial abundance of \( N_1 \). After finding the numerical solution
of eq.\((5.8)\), we obtain the present day abundance of DM and further we implement this solution in calculating the relic abundance of DM in the present Universe using the equation:

\[
\Omega h^2 = \frac{M_{DM} Y_0 s_0}{\rho_c},
\] (5.9)

where, \(\rho_c \sim 1.05 \times 10^{-5} h^2 \text{ GeVcm}^{-3}\) is the critical density of the Universe, \(s_0 \sim 2891.2 \text{ cm}^{-3}\) is the current entropy density and \(h = 0.72\) is the Hubble parameter.

As we know, that the decay of \(N_1\) may release entropy resulting in some discrepancy in the ratio of the abundance of the light particles, which were confirmed to match with the standard ΛCDM cosmology. Hence, the decay of \(N_1\) must be restricted to occur during or after the epoch of the big-bang nucleosynthesis (BBN) \([46]\). Thus, we get a constraint on the minimum value of decay width of \(N_1\), i.e., \(\Gamma_{N_1} \geq \Gamma_{N_1,\text{min}} \equiv 6.58 \times 10^{-25} \text{ GeV}\), arising from the consideration that the decay lifetime of \(N_1\) should be less than 1 second. Again, an upper bound on the decay width, i.e. \(\Gamma_{N_1} \leq \Gamma_{N_1,\text{max}} \equiv \frac{M_{DM}}{x_0} \times 10^{-18} \text{ GeV}\) is a manifestation of the fact that the decay of \(N_1\) should take part mostly after the DM candidate freezes out thermally so as to give adequate contribution towards the relic abundance. Thus, we investigate the limitations that we encountered during the thermal production of the relic and see for what benchmark values of the free parameters \(\Gamma_{N_1}\) and \(Y_{N_1}(x_0)\) we can have correct relic abundance within the IHDM desert even for high mass splitting.

As we have considered the lightest stable scalar particle to be a probable dark matter candidate, thus, the spin independent scattering cross section of the SM Higgs is expressed by\([69]\):

\[
\sigma_{SI} = \frac{\lambda_L^2 f^2 m^2_{\mu} m^2_n}{4\pi m^4_h M^2_{DM}},
\] (5.10)

where, \(\lambda_L = (\lambda_3 + \lambda_4 + \lambda_5)/2\) is the quartic coupling taking part in the DM-Higgs interaction, \(m^2_{\mu} = m_n M_{DM}/m_n + M_{DM}\) is the DM-nucleon reduced mass and \(f\) is the Higgs-nucleon coupling which is estimated to be \(f = 0.32\) \([70]\). There also can be a Higgs portal coupling independent DM-nucleon scattering cross-section at a one-loop level \([71]\). However, by appropriate choice of the mass splitting between the scalar components, we can generate spin-independent scattering cross-section much lower than that obtained from direct detection experiment XENON1T.
FIG. 3: Plots in the first-row show baryon asymmetry as a function of dark matter mass ($M_{DM}$), the second-row show baryon asymmetry as a function of right-handed neutrino mass ($M_{N_1}$), in third-row baryon asymmetry as a function of lightest neutrino mass with the red vertical line signifying Planck limit for the sum of the light neutrino masses, is shown. The fourth row depicts baryon asymmetry as a function of the absolute value of quartic coupling ($|\lambda_5|$) for NH and IH, respectively. The black horizontal line gives the current Planck limit for BAU.
FIG. 4: Effective mass as a function of lightest neutrino mass for NH/IH.

FIG. 5: Baryon Asymmetry as a function of effective mass of neutrino for NH/IH. The horizontal (black) line is the Planck limit for BAU and the vertical (red) line depicts the KamLAND-Zen limit for $0\nu\beta\beta$.

VI. NUMERICAL ANALYSIS AND RESULTS

In this study, we choose the dark matter mass in the intermediate-mass range, $M_W < M_{DM} \leq 550$ GeV, and study the consequences of neutrino mass, neutrinoless double beta decay and baryon asymmetry of the Universe. The plot in the first row of fig. 3 depicts that the observed baryogenesis is satisfied for almost the entire IHDM desert for NH, whereas, in case of IH, baryogenesis is obtained for dark mass above 100 GeV. Furthermore, for $N_1$ leptogenesis in the scotogenic model, we can consider the mass of the lightest RHN as low as 10 TeV. Hence, in our work, we have chosen the RHN masses $M_{N_1}, M_{N_2}$ and $M_{N_3}$ in the range $10^4 - 10^5$ GeV, $10^7 - 10^8$ GeV and $10^9 - 5 \times 10^9$
GeV respectively.

The first row of fig. 3 shows the variation between the baryon asymmetry of the Universe and the dark matter mass (\(M_{DM}\)). In contrast, the second row shows the variation of BAU results with the mass of the lightest right-handed neutrino \(M_{N_1}\) for both NH and IH and we obtain the parameter space of \(M_{DM}\) and \(M_{N_1}\) that satisfies currently observed value of BAU in both the mass ordering. From the results of \(M_{N_1}\) vs. \(\eta_B\), one can get the lower bound on the lightest RH neutrino, which is almost \(3 \times 10^4\) GeV for NH, whereas in the case for IH, \(M_{N_1} \sim 2 \times 10^4\) GeV is successful in generating the observed BAU. As we realize WIMP type of dark matter, we seek the Yukawa couplings to be much larger than that required in FIMP type dark matter. We can generate large

![FIG. 6: Variation of relic abundance of DM in the intermediate dark matter mass range.](image)

![FIG. 7: The allowed region of parameter space in \(\lambda_L-M_{DM}\) plane from the requirement of satisfying the relic abundance and depiction of the strict constraints from dark matter direct detection experiment, XENON1T.](image)
values of Yukawa couplings by taking the lightest active neutrino mass \( m_l \) in the range of \( 10^{-3} - 1 \) eV. Such a mass range of \( m_l \) along with the fact that \( M_{DM} \neq M_{N_1} (\eta_1 \neq 1) \) is crucial in the strong washout regime. Thus, even for the 3RHN scenario, we are in a strong washout regime similar to the 2RHN case. For such a span of lightest active neutrino mass, a plot of baryogenesis vs. \( m_l \) is shown in the third row of fig. 3 where the left panel shows the variation for NH and the right panel for IH. Though, in NH we see few points with \( m_l \geq 10^{-2} \) eV satisfying the baryogenesis bound however for IH, \( m_l \sim 10^{-3} \) eV and above satisfies baryogenesis. The entire work is carried out for a fixed range of the quartic coupling \( |\lambda_5| \), i.e., \( 10^{-5} - 10^{-2} \). Therefore, we analyze the parameter space of the quartic coupling satisfying the observed baryon asymmetry, which can be estimated to be \( \mathcal{O}(10^{-3} - 10^{-2}) \) as shown in the last row of fig. 3. As we have also studied \( 0\nu\beta\beta \) in this work and the variation of \( m_{\beta\beta} \) vs. \( m_l \) for NH and IH are shown in fig. 4. Moreover, a correlative analysis of the points satisfying both effective mass and baryogenesis is also shown in fig. 5.

As, it is a \( N_1 \) dominated leptogenesis, the probable candidate of DM will be the lightest particle among the inert Higgs doublet. In our study, \( \eta_{R}^0 \) is considered to be a source of DM, with the assumption of it being the lightest of all scalars. Therefore, it’s relic abundance is calculated by implementing first this minimal scotogenic model in Feynrules\(^72\) and then using the computational package MicrOmega 5.0.4\(^73\). The relic abundance as a function of the DM mass \( M_{DM} \) is manifested in fig. 6 where, the DM-Higgs coupling is taken to be as low as \( \lambda_L = 0.0001 \) and the mass differences \( \Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 1 \) GeV(left panel). Also, in fig. 6 we have shown a similar plot of relic vs. \( M_{DM} \) for higher values of \( \Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 10 \) GeV(right panel). From fig. 6, we can anticipate that for low mass splitting between the scalars i.e. \( \Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 1 \) GeV, the relic is suppressed in the low mass regime due to the increase in co-annihilation between the different components of inert scalar doublet. Whereas, the high mass regime has a suppressed relic because the annihilation contribution of the electroweak bosons that increases with the mass square differences among the inert scalars.

Furthermore, instead of fixing the DM-Higgs coupling, we show the allowed region of parameter space in the \( \lambda_L - M_{DM} \) plane from the obligation of satisfying the correct relic abundance depicted in fig. 7. With the relic abundance bound on the \( \lambda_L - M_{DM} \) plane, there also exist strict constraint from the dark matter direct detection experiment XENON1T. The scattered points in fig. 7 corresponds to the values of \( M_{DM} \) and \( \lambda_L \), which are allowed from the direct detection bound of XENON1T and
FIG. 8: Variation of relic abundance for three different values of dark matter mass with values of $\Gamma_{N_1}$ and $Y_N(x_0)$ given in plot. The corresponding parameters which contribute in determining relic are kept fixed with values: $\Delta M_{\eta^\pm} = \Delta M_{\eta_0^I} = 1$ GeV(left panel) and $\Delta M_{\eta^\pm} = \Delta M_{\eta_0^I} = 10$ GeV(right panel) , $\lambda_L = 0.0001$, $\lambda_2 = 0.2$ and $M_h = 125.5$ GeV.

the small dark portion refer to the points allowed by the current value of relic density. Thus, we can see that there exists a coincidence of both points signifying the parameter space, which obeys constraints from both the cosmological aspects mentioned above. We see a significant difference in the parameter space of $\lambda_L$ w.r.t. the mass difference of the scalars. Hence, we can confirm the choice of mass difference is of utmost importance in determining the relic abundance of dark matter [71].

Dark matter relic density primarily depends on the dark matter mass, Higgs portal coupling, and mass differences with the LSP and nLSP\textsuperscript{1}. In the low mass region for $M_{DM} < 10$ GeV, most dominating DM annihilation processes are to the SM fermions only, and due to small coupling strength and mass, we get an overabundance of the relic density. Moreover, the dominant part of the points ruled out by the Higgs/Z invisible decay width and direct detection constraints for the low mass. Within IHDM, irrespective of the choice of parameter spaces, the region in between $M_W < M_{DM} \leq 530$ GeV does not give observed relic abundance value due to the very high annihilation rate of $DM \rightarrow W^\pm W^\pm, ZZ$ [74][76]. However, by considering different production mechanisms as discussed by [46][77], we can work out the on the IHDM desert region to get

\textsuperscript{1} Lightest stable particle and next to lightest stable particle.
FIG. 9: Variation of relic abundance with DM mass fixed at $M_{DM} = 430$ GeV and $\Gamma_{N_1} = 1.2 \times 10^{-19}$ for three different values of $Y_N(x_0)$. The corresponding parameters which contribute in determining relic are kept fixed with values: $\Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 1$ GeV (left panel) and $\Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 10$ GeV (right panel), $\lambda_L = 0.0001$, $\lambda_2 = 0.2$ and $M_h = 125.5$ GeV.

FIG. 10: Relic abundance for dark matter mass 530 GeV with $Y_N(x_0) = 10^{-11}$ and values of $\Gamma_{N_1}$ given in plot. The corresponding parameters which contribute in determining relic are kept fixed with values: $\Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 10$ GeV, $\lambda_L = 0.0001$, $\lambda_2 = 0.2$ and $M_h = 125.5$ GeV.

observed relic abundance. Here we consider the decay of a particle $N_1$, which produces dark matter non-thermally, and by adjusting suitable decay width and initial abundance of dark matter candidate, we can generate observed relic density within the IHDM desert. From fig.6, we can see
the deviation in relic abundance, taking into consideration the crucial parameter, i.e., the mass splitting among the scalars of the inert doublet. For $\Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 1$ GeV, we get the correct relic abundance corresponding to $M_{DM} = 530$ GeV, whereas for $\Delta M_{\eta^\pm} = \Delta M_{\eta^0} = 10$ GeV, we fail to generating the relic. Therefore, we proceed with the non-thermal production of dark matter to see if the desired relic is obtained for the above value of dark matter and masses even lower than it. We consider the low mass splitting case with $M_{DM} = 530$ GeV and by appropriate choice of the decay width($\Gamma_{N_1}$), we see that for $Y_N(x_0) = 10^{-11}$ GeV, it produces the correct relic abundance. Again for high mass splitting, we show the deviation in thermal and non-thermal production of the relic. We verify the result obtained in the right panel of fig.6, again by fig.10, that dark matter is underabundant thermally. It shows that for $M_{DM}=530$ GeV and the choice of other parameters, $\Gamma_N = 2.5 \times 10^{-19}$ and $Y_N(x_0) = 10^{-11}$, we obtain the relic, whereas for $\Gamma_N = 0$ GeV, there is an underabundant production of relic. We also do a relative study for three benchmark values of dark matter in the low mass splitting as well as the high mass splitting scenario depicted in fig.8. We now fine-tune the decay width in order to obtain the correct relic abundance for $M_{DM} = 430$ GeV, which was underabundant for the values shown in fig.8. Thus, fig.9 showcases the two different mass splitting scenarios for $M_{DM} = 430$, and investigate the values of $\Gamma_{N_1}$ and $Y_N(x_0)$ which satisfies the correct relic abundance. Therefore, we can see a distinct variation of $\Gamma_{N_1}$ and $Y_N(x_0)$ w.r.t. dark matter mass resulting in the production of correct relic abundance.

VII. CONCLUSION

In this paper, we study an extension of the SM popularly known as the scotogenic model, which is extended by a Higgs doublet ($\eta$) and three singlet neutral fermions ($N_k$). An additional $Z_2$ charge is assigned in the model, and all the SM particles are ever under it while additional fields are odd. The possibility of a DM candidate comes from the $Z_2$ odd lightest particle. We carry out this work with the dark matter mass strictly focusing in the intermediate dark matter mass range, also known as the inert Higgs doublet model (IHDM) desert, which lies between $M_W < M_{DM} \leq 550$ GeV. Along with DM, baryogenesis via the mechanism of thermal leptogenesis and neutrinoless double beta decay is also addressed in this work. Leptogenesis is a result of the decay of $Z_2$ odd fermions, i.e, the heavy RHN, which occurs via the out-of-equilibrium decay into the SM leptons and the
inert Higgs doublet. The out-of-equilibrium decay of \( N_1 \rightarrow l\eta, \bar{l}\eta^* \), where \( \eta \) is the inert Higgs doublet constituting the dark matter candidate \( \eta^0_R \) that generates the observed baryon asymmetry of the Universe. For two different choice of mass splitting between the DM (LSP) and the next heavier scalar (nLSP), we study the relic abundance of the dark matter candidate (lightest of \( \eta \)). We also study the mixture of thermal and non-thermal production of DM abundance for various masses within the IHDM desert.

As our model is compatible with baryogenesis studied in the IHDM desert, we are successfully able to show co-relation plot of dark matter mass \( (M_{DM}) \), RHN mass \( (M_{N_1}) \), lightest neutrino mass \( (m_1) \) and quartic coupling parameter \( (\lambda_5) \) with the latest observed value of BAU. We obtain a particular range of quartic coupling, between \( 10^{-3} - 10^{-2} \), which is accountable for reproducing the observed baryon asymmetry of the Universe by the decay of \( N_1 \) with a mass in the range \( 10^4 - 10^5 \) GeV. We also consider the lightest neutrino mass in the range \( 10^{-3} - 1 \) eV and check its consistency with the experimental bounds obtained from KamLAND-Zen by the neutrinoless double beta decay method. The correlation between the BAU result and \( 0\nu\beta\beta \) has a very constrained space in our work for both the mass ordering. From the synchronous study of \( 0\nu\beta\beta \) and baryogenesis, it is evident that both the observable are loosely co-related in our model. Moreover, the light neutrino masses ranging from approximately 0.05 eV to 0.1 eV are more likely to satisfy the KamLAND-Zen limit for \( m_{\beta\beta} \), and at the same time, they obey Planck limit for generating the observed BAU.

The significant conclusion we observe from our analysis is that the mass splitting, \( \Delta M_{\eta^\pm} = \Delta M_{\eta^0_R} \) plays a vital role in the production of relic abundance. As, in our work, we could generate relic for \( \Delta M_{\eta^\pm} = \Delta M_{\eta^0_R} = 1 \) GeV but failed in the case of \( \Delta M_{\eta^\pm} = \Delta M_{\eta^0_R} = 10 \) GeV for the same value of \( \lambda_L = 0.0001 \), which therefore satisfies the LEP constraints \[78\] as it rules out values of mass splitting greater than 8 GeV. This draws attention to how effective the mass splitting could be in the IHDM. It also motivates us to study the non-thermal production of dark matter. For non-thermal production of dark matter, we observe current relic abundance for the appropriate choice of decay with and coupling parameters with \( \Delta M = 10 \) GeV.
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