Experimentally Verifiable $U(1)_{B-L}$ Symmetric Model with Type-II Seesaw and Dark Matter

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Abstract

In an endeavor to explain the light neutrino masses and dark matter (DM) simultaneously, we study a gauged $U(1)_{B-L}$ extension of the standard model (SM). The neutrino masses are generated through a variant of type-II seesaw mechanism in which one of the scalar triplets has a mass in a scale that is accessible at the present generation colliders. Three right chiral fermions $\chi_{iR}(i = e, \mu, \tau)$ with $B-L$ charges -4, -4, +5 are invoked to cancel the $B-L$ gauge anomalies and the lightest one among these three fermions becomes a viable DM candidate as their stability is guaranteed by a remnant $Z_2$ symmetry to which $U(1)_{B-L}$ gauge symmetry gets spontaneously broken. Interestingly in this scenario, the neutrino mass and the co-annihilation of DM are interlinked through the breaking of $U(1)_{B-L}$ symmetry. Apart from giving rise to the observed neutrino mass and dark matter abundance, the model also predicts exciting signals at the colliders especially regarding the discovery of the triplet scalar in presence of the $B-L$ gauge boson. We see a $(34-54)\%$ enhancement in the production of the TeV scale doubly charged scalar in presence of the $Z_{BL}$ gauge boson in a mass range $2.5\text{ TeV}$ to $4.4\text{ TeV}$. We discuss all the relevant constraints on model parameters from observed DM abundance and null detection of DM at direct and indirect search experiments as well as the constraints on the $B-L$ gauge boson from recent colliders.

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I. INTRODUCTION

Out of all the lacunae afflicting the Standard Model (SM) of particle physics, the identity of DM and the origin of tiny but nonzero neutrino masses are the most irking ones. It is well established by now, thanks to numerous irrefutable observational evidences from astrophysics and cosmology like galaxy rotation curves, gravitational lensing, Cosmic Microwave Background (CMB) acoustic oscillations etc. [1–6], that a mysterious, non-luminous and non-baryonic form of matter exists called as dark matter (DM) which constitutes almost 85% of the total matter content and around 26.8% of the total energy density of the present Universe. In terms of density parameter $h = (\text{Hubble Parameter})/(100 \text{ km s}^{-1}\text{Mpc}^{-1})$, the present DM abundance is conventionally reported as [5, 6]

$$\Omega_{\text{DM}}h^2 = 0.120 \pm 0.001$$ (1)

But still we have no answer to the question what DM actually is, as none of the SM particle has the properties that of a DM particle is expected to have. Thus over the years, various beyond SM (BSM) scenarios have been considered to explain the puzzle of DM, with additional field content and augmented symmetry. The most popular among these ideas is something known as the weakly interacting massive particle (WIMP) paradigm. In this WIMP scenario, a DM candidate typically having a mass similar to electroweak (EW) scale and interaction rate analogous to EW interactions can give rise to the correct DM relic abundance, an astounding coincidence referred to as the WIMP Miracle [7, 8]. The sizeable interactions of WIMP DM with the SM particles has many phenomenological implications. Along with giving the correct relic abundance of DM through thermal freeze-out, it also leads to other phenomenological implications like optimistic direct and indirect detection prospects of DM which makes it more appealing. Several direct detection experiments like LUX, PandaX-II and XENON1T [9–13] and indirect detection experiments like space based telescopes Fermi-LAT and ground based telescopes MAGIC [14, 15] have been looking for signals of DM and have put constraints on DM-nucleon scattering cross-sections and DM annihilation cross-section to SM particles respectively.

Apart from the identity of DM, another appealing motivation for the BSM is the origin of neutrino masses. Despite compelling evidences for existence of light neutrino masses, from various oscillation experiments [16–20] and cosmological data [21–24], the origin of light neutrino masses is still unknown. The oscillation data is only sensitive to the difference in
mass-squareds\textsuperscript{[21, 25]}, but the absolute mass scale is constrained to \(\sum_i |m(\nu_i)| < 0.12 \text{ eV}\) \textsuperscript{[21]} from cosmological data. This also implies that we need new physics in BSM to incorporate the light neutrino masses as the Higgs field, which lies at the origin of all massive particles in the SM, can not have any Yukawa coupling with the neutrinos due to the absence of its right-handed counterpart.

Assuming that the neutrinos to be of Majorana type (which violates lepton number by two units), the origin of the tiny but non-zero neutrino mass is usually explained by the seesaw mechanisms (Type-I \textsuperscript{[26–29]}, Type-II \textsuperscript{[30–34]} and Type-III \textsuperscript{[35]}) which are the ultraviolet completed realizations of the dimension five Weinberg operator \(O_5 = y_{ij} \frac{\bar{T}_i L_j HH}{\Lambda}\), where \(L\) and \(H\) are the lepton and Higgs doublets of the SM and \(\Lambda\) is the scale of new physics \textsuperscript{[36, 37]}. In the type-I seesaw heavy singlet RHNs are introduced while in type-II and type-III case, a triplet scalar(\(\Delta\)) of hyper-charge 2 and triplet fermions \(\Sigma\) of hyper-charge 0 are introduced respectively such that new Yukawa terms can be incorporated in the theory. Tuning the Yukawa coupling and the cut-off scale (\(\Lambda\)) and adopting a necessary structure for the mass matrix, the correct masses and mixings of the neutrinos can be obtained.

In the conventional type-II seesaw, the relevant terms in the Lagrangian violating lepton number by two units are \(f_{ij} \Delta L_i L_j + \mu \Delta^\dagger HH\), where \(\Delta\) does not acquire an explicit vacuum expectation value(vev). However, after the electro-weak phase transition, a small induced vev of \(\Delta\) can be obtained as: \(\langle \Delta \rangle = -\frac{\mu \langle H^2 \rangle}{M_\Delta}\). Thus for \(\mu \sim M_\Delta \sim 10^{14}\text{ GeV}\), one can get \(M_\nu = f \langle \Delta \rangle \approx f \frac{\langle H^2 \rangle}{M_\Delta^2}\) of order \(\mathcal{O}(0.1)\text{ eV}\) for \(f \sim 1\). But such models lack verifiability as the mass scale of the scalar triplet is much larger than the energy attainable at current generation colliders.

In an alternative fashion, neutrino mass can be generated in a modified type-II seesaw if one introduces two scalar triplets: \(\Delta\) and \(\xi\) with \(M_\Delta \sim 10^{14}\text{ GeV}\) and \(M_\xi \sim \text{TeV} \ll M_\Delta\)\textsuperscript{[38]}. In this case imposition of additional B - L gauge symmetry \textsuperscript{[40]} allows for \(\mu \Delta^\dagger HH + f \xi LL + g \Phi_{\text{BL}}^2 \Delta^\dagger \xi\) terms in the Lagrangian (with proper choice of gauge charges for the scalars) where \(\Phi_{\text{BL}}\) is the scalar field responsible for B - L symmetry breaking at TeV scale. As is clear from the Lagrangian terms, once the \(\Phi_{\text{BL}}\) acquires a vev, it creates a small mixing between \(\Delta\) and \(\xi\) of the order \(\theta \sim \frac{(\Phi_{\text{BL}})^2}{M_\Delta^2} \sim 10^{-18}\). Thus the coupling of \(\xi\) with Higgs becomes extremely suppressed but \(\xi LL\) coupling can be large. In this scenario,

\textsuperscript{1} See also ref. \textsuperscript{[39]} for a modified double type-II seesaw with TeV scale scalar triplet
Δ being super heavy gets decoupled from the low energy effective theory but ξ can have mass from several hundred GeV to a few TeV and having large Di-lepton coupling can be probed at colliders through the same sign Di-lepton signature [41–47].

If this type-II seesaw framework is implemented in a gauged $U(1)_{B-L}$ symmetric model, then one needs to invoke additional fields to make the model anomaly free as gauging the $U(1)_{B-L}$ introduces non-trivial gauge and gravitational anomalies. With only the SM particle content all triangle anomalies in the gauged $B-L$ theory vanishes except for $\sum [U(1)_{B-L}]^3 = 3$ and $\sum [Grav.]^2 \times [U(1)_{B-L}] = 3$. Thus the only way to cancel these anomalies is by introducing new fermions in such a way that sum of their $B-L$ quantum numbers is $-3$. In this paper we adopt the choice of $U(1)_{B-L}$ charges $-4, -4, +5$ for three right chiral fermions $\chi_{iR} (i = e, \mu, \tau)$ introduced for anomaly cancellation, such that $\sum_{i=1}^3 (Y_{B-L})_i = -3$ [48–58]. For the details of the anomaly cancellation in a $B-L$ model, please see appendix A.

Interestingly, the lightest one among these three exotic fermions becomes a viable candidate of DM, thanks to the remnant $Z_2$ symmetry after $U(1)_{B-L}$ breaking, under which $\chi_{iR} (i = e, \mu, \tau)$ are odd while all other particles are even. 2

The origin of neutrino mass and DM is hitherto not known. Any connection between them is also not established yet. However, it will really be interesting if neutrino mass and the DM phenomenology have an inter-connection between them. In light of this, it is worth mentioning here that the framework we considered in this paper, the spontaneous breaking of $U(1)_{B-L}$ gauge symmetry via the vev of $\Phi_{BL}$ not only generates sub-eV masses of light neutrinos, but also give rise to co-annihilations among the dark fermions which in turn allows a larger parameter space in contrast to the gauged $B-L$ models with right-handed neutrino DM giving rise required relic only at the $Z_{B-L}$-resonance [56, 60–62].

The rest of the paper is organized as follows. In section II, we describe the proposed model, the neutrino mass generation through a variant of type-II seesaw, the scalar masses and mixing. We then discuss how the particles introduced for anomaly cancellation become viable DM candidate and study the relic density in section III. In section IV, we studied all the relevant constraints from direct, indirect search of DM on our parameter space as well as scrutinized it with respect to the constraint from colliders. We briefly summarize

2 Right handed neutrinos with $B-L$ charge -1 can also serve the purpose of $B-L$ anomaly cancellation and be viable DM candidate provided one introduces an additional ad hoc $Z_2$ symmetry to guarantee their stability. For instance see [59, 60]
the collider search strategies of the model in section V and finally conclude in section VI.

II. THE MODEL

| Fields                  | $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_{B-L}$ |
|-------------------------|-------------------------------------------------------------|
| Exotic Fermions         |                                                             |
| $\chi_{eR}, \chi_{\mu R}$ | 1 1 0 -4                                                  |
| $\chi_{\tau R}$        | 1 1 0 5                                                  |
| Heavy Scalars           |                                                             |
| $\Delta = \left( \begin{array}{cc} \delta^+ & \delta^{++} \\ \delta^0 & -\delta^+ \end{array} \right)$ | 1 3 2 0                                           |
| $\Phi_{BL}$             | 1 1 0 -1                                                 |
| $\Phi_{12}$             | 1 1 0 8                                                  |
| Light Scalars           |                                                             |
| $\xi = \left( \begin{array}{cc} \xi^+ & \xi^{++} \\ \xi^0 & -\xi^+ \end{array} \right)$ | 1 3 2 2                                           |
| $\Phi_3$                | 1 1 0 -10                                                |

TABLE I: Charge assignment of BSM fields under the gauge group $G \equiv G_{SM} \otimes U(1)_{B-L}$, where $G_{SM} \equiv SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$.

In this work, the model under consideration is a very well motivated BSM framework based on the gauged $U(1)_{B-L}$ symmetry [63–67] in which we implement a modified type-II seesaw to explain the sub-eV neutrino mass by introducing two triplet scalars $\Delta$ and $\xi$. $\Delta$ is super heavy with $M_\Delta \sim 10^{14}$ GeV and $M_\xi \sim \text{TeV} << M_\Delta$ and the $B-L$ charges of $\Delta$ and $\xi$ are 0 and 2 respectively. As already discussed in the previous section, the additional $U(1)_{B-L}$ gauge symmetry introduces $B-L$ anomalies in the theory. And to cancel these $B-L$ anomalies we introduce three right chiral fermions $\chi_{iR}$ ($i = e, \mu, \tau$), where the $B-L$ charges of $\chi_{eR}, \chi_{\mu R}$ and $\chi_{\tau R}$ are -4, -4, +5 respectively. Note that such unconventional $B-L$ charge assignment of the $\chi_{iR}$ ($i = e, \mu, \tau$) forbids their Yukawa couplings with the SM particles. Also three singlet scalars: $\Phi_{BL}$, $\Phi_{12}$ and $\Phi_3$ with $B-L$ charges -1, +8, -10 are introduced. As a result of which $\Phi_{12}$ and $\Phi_3$ couples to $\chi_{eR, \mu R}$ and $\chi_{\tau R}$ respectively.
through Yukawa terms and the vevs of $\Phi_{12}$ and $\Phi_3$ provides Majorana masses to these exotic fermions. The vev of $\Phi_{BL}$ provides a small mixing between $\Delta$ and $\xi$ which plays a crucial role in generating sub-eV masses of neutrinos. This vev $\langle \Phi_{BL} \rangle$ is also instrumental in controlling the co-annihilations among the dark sector particles and hence is crucial for DM phenomenology too. As a consequence this establishes an interesting correlation between the neutrino mass and DM. The particle content and their charge assignments are listed in Table I.

The Lagrangian involving the BSM fields consistent with the extended symmetry is given by:

$$\mathcal{L} = \mathcal{L}^{\text{SM+BSM Scalar}} + \mathcal{L}^{\text{DM}}$$

(2)

where

$$\mathcal{L}^{\text{SM+BSM Scalar}} \supset |D_\mu H|^2 + |D_\mu \Phi_3|^2 + |D_\mu \Phi_{BL}|^2 + |D_\mu \Phi_{12}|^2$$

$$+ Tr[(D_\mu \Delta)^\dagger (D^\mu \Delta)] + Tr[(D_\mu \xi)^\dagger (D^\mu \xi)]$$

$$- Y_{ij} \overline{\chi}_i \sigma_i^j \xi_j L_j$$

$$- V^L(H, \xi, \Phi_3) - V^H(\Delta, \Phi_{BL}, \Phi_{12}) - V^{LH}.$$  (3)

Here $i, j$ runs over all three lepton generations. In the above Lagrangian, $V^L$ is the scalar potential involving scalars in sub-TeV mass range $(H, \xi, \Phi_3)$, $V^H$ stands for scalar potential of the heavy fields $(\Delta, \Phi_{BL}, \Phi_{12})$ and the scalar potential which involves both sub-TeV and super heavy fields is defined by $V^{LH}$.

The covariant derivatives for these fields can be written as:

$$D_\mu H = \partial_\mu H + ig T^a W^a_\mu H + ig' B_\mu H$$

$$D_\mu F = \partial_\mu F + ig [T^a W^a_\mu, F] + ig' Y B_\mu F + ig_{BL} Y_{BL}(Z_{BL})_\mu F$$  where $F = \{\xi, \Delta\}$

$$D_\mu G = \partial_\mu G + ig_{BL} Y_{BL}(Z_{BL})_\mu G$$  where $G = \{\Phi_3, \Phi_{BL}, \Phi_{12}\}$. 

The Lagrangian of the dark sector can be written as:

$$\mathcal{L}^{\text{DM}} = \overline{\chi}_e i\gamma^\mu D_\mu \chi_e + \overline{\chi}_{\mu R} i\gamma^\mu D_\mu \chi_{\mu R} + \overline{\chi}_{\tau R} i\gamma^\mu D_\mu \chi_{\tau R}$$

$$+ Y_{11} \phi_{12} (\chi_{e R})^c \chi_{e R} + Y_{22} \phi_{12} (\chi_{\mu R})^c \chi_{\mu R} + Y_{12} \phi_{12} (\chi_{\tau R})^c \chi_{\tau R}$$

$$+ Y_{13} \phi_{BL}(\chi_{e R})^c \chi_{\tau R} + Y_{23} \phi_{BL}(\chi_{\mu R})^c \chi_{\tau R}$$

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\[ + Y_{33} \Phi_3 \left( \chi_{\tau R} \right)^c \chi_{\tau R} + \text{h.c.} \]  

(4)

where

\[ D_\mu \chi = \partial_\mu \chi + i g_{bl} Y_{bl}(Z_{bl})_\mu \chi. \]

The gauge coupling associated with \( U(1)_{B-L} \) is \( g_{bl} \) and \( Z_{bl} \) is the corresponding gauge boson. The scalar potentials which are mentioned in the Lagrangian in Eq. 3 can be written as:

\[ V^L(H, \xi, \Phi_3) = -\mu_H^2 H^\dagger H + \lambda_H(H^\dagger H)^2 + M_\xi^2 Tr[\xi^\dagger \xi] + \lambda_\xi (Tr[\xi^\dagger \xi])^2 + \lambda_\xi' Tr[(\xi^\dagger \xi)^2] + \lambda_\xi H Tr[\xi^\dagger \xi] (H^\dagger H) + \lambda_\xi' H\xi^\dagger H \]

\[ - \mu_{\Phi_3}^2 \Phi_3^\dagger \Phi_3 + \lambda_{\Phi_3} (\Phi_3^\dagger \Phi_3)^2 \]

(5)

\[ V^H(\Delta, \Phi_{bl}, \Phi_{12}) = M_\Delta^2 Tr[\Delta^\dagger \Delta] + \lambda_\Delta (Tr[\Delta^\dagger \Delta])^2 + \lambda_\Delta' Tr[(\Delta^\dagger \Delta)^2] \]

\[ - \mu_{\Phi_{bl}}^2 \Phi_{bl}^\dagger \Phi_{bl} + \lambda_{\Phi_{bl}} (\Phi_{bl}^\dagger \Phi_{bl})^2 + \lambda_{\Delta \Phi_{bl}} Tr[\Delta^\dagger \Delta] (\Phi_{bl}^\dagger \Phi_{bl}) \]

\[ - \mu_{\Phi_{12}}^2 \Phi_{12}^\dagger \Phi_{12} + \lambda_{\Phi_{12}} (\Phi_{12}^\dagger \Phi_{12})^2 + \lambda_{\Delta \Phi_{12}} Tr[\Delta^\dagger \Delta] (\Phi_{12}^\dagger \Phi_{12}) \]

(6)

\[ V^{LH} = \lambda_{\Delta H} Tr[\Delta^\dagger \Delta] (H^\dagger H) + \lambda_{\Delta H}' H^\dagger H (H^\dagger H) + [\mu_H (H^T i \sigma^2 \Delta^\dagger H) + \text{h.c.}] \]

\[ + \lambda_{H \Phi_{bl}} (H^\dagger H) (\Phi_{bl}^\dagger \Phi_{bl}) + \lambda_{H \Phi_{12}} (H^\dagger H) (\Phi_{12}^\dagger \Phi_{12}) + \lambda_{\Phi_{bl} \Phi_3} (\Phi_{bl}^\dagger \Phi_3) (\Phi_{3}^\dagger \Phi_3) \]

\[ + \lambda_{\Phi_{12} \Phi_3} (\Phi_{12}^\dagger \Phi_3) (\Phi_{3}^\dagger \Phi_3) + \lambda_{\xi \Phi_{bl}} Tr[\xi^\dagger \xi] (\Phi_{bl}^\dagger \Phi_{bl}) + \lambda_{\xi \Phi_{12}} Tr[\xi^\dagger \xi] (\Phi_{12}^\dagger \Phi_{12}) \]

\[ + \lambda_{\Delta \xi} Tr[\Delta^\dagger \Delta] Tr[\xi^\dagger \xi] + \lambda_{\Delta \xi} Tr[\Delta^\dagger \Delta] Tr[\xi^\dagger \Delta] \]

\[ + \lambda_{\Phi_{bl}}^2 Tr[\Delta^\dagger \xi] + \lambda_{\Phi_{bl}}^2 (\Phi_{bl}^\dagger \Phi_{12})^2 + \text{h.c.} \]  

(7)

Here it is worth mentioning that the mass squared terms of \( \Delta \) and \( \xi \) are chosen to be positive so they do not get any spontaneous vev. Only the neutral components of \( H, \Phi_{12}, \Phi_{bl} \) and \( \Phi_3 \) acquire non-zero vevs. However, after electroweak phase transition, \( \Delta \) and \( \xi \) acquire induced vevs.

For simplicity, we assume certain mass hierarchy among the scalars. The masses of \( H, \Phi_3 \) and \( \xi \) are of similar order in sub-TeV range, while the masses of \( \Phi_{bl} \) and \( \Phi_{12} \) are in Several
TeV scale. To make the analysis simpler we decouple the light scalar sector from the heavy scalar sector by considering all quartic couplings in the scalar potential $V^{LH}$ to be negligible. It is worth mentioning that this assumption does not affect our DM phenomenology.

We parameterize the neutral scalars in the low-energy scale as:

$$H^0 = \frac{v_H + h_H + i G_H}{\sqrt{2}}, \quad \xi^0 = \frac{v_\xi + h_\xi + i G_\xi}{\sqrt{2}}, \quad \Phi_3 = \frac{v_3 + h_3 + i G_{\Phi_3}}{\sqrt{2}}.$$  

A. NEUTRINO MASS

![Diagram of neutrino mass generation through the modified Type-II seesaw](image)

FIG. 1: Generation of neutrino mass through the modified Type-II seesaw.

The relevant Feynman diagram of this modified type-II seesaw which gives rise to light neutrino masses is as shown in Fig 1. In this modified version of type-II seesaw, the conventional heavy triplet scalar $\Delta$ can not generate Majorana masses for light neutrinos as the $B-L$ quantum number of $\Delta$ is zero. However this super heavy scalar $\Delta$ can mix with the TeV scalar triplet $\xi$, once $\Phi_{BL}$ acquires vev and breaks the $U(1)_{B-L}$ symmetry spontaneously. By the virtue of the trilinear term of $\Delta$ with SM Higgs doublet $H$, it gets an induced vev after electroweak symmetry breaking similar to the case of traditional type-II seesaw. The induced vev acquired by $\Delta$ after EW phase transition is given by

$$\langle \Delta \rangle = v_\Delta \simeq -\frac{\mu_{\Delta} v_H^2}{2\sqrt{2} M_\Delta^2}.$$  

The vev of $\Delta$ is required to satisfy:

$$\rho \equiv \frac{M_W^2}{M_Z^2 \cos^2 \theta} = \frac{1 + 2x^2}{1 + 4x^2} \approx 1$$  

where $x = \frac{v_\xi}{v}$.
where $x = v_\Delta/v_H$. The above constraint implies that $|v_\Delta| \sim \mathcal{O}(1)$ GeV.

Since $\xi$ mixes with $\Delta$ after $U(1)_{B-L}$ breaking, it also acquires an induced vev after EW symmetry breaking which is given by

$$
\langle \xi \rangle = v_\xi = -\frac{\lambda_P v_{BL}^2}{4M_\xi^2}v_\Delta.
$$

(10)

Assuming $\lambda_P v_{BL}^2 \sim M_\xi^2$, we obtain $v_\Delta \simeq v_\xi$, even if $\xi$ and $\Delta$ have several orders of magnitude difference in their masses. After integrating out the heavy degrees of freedom in the Feynman diagram given in Fig.1, we get the Majorana mass matrix of the light neutrinos to be

$$
(M_\nu)_{ij} = Y_{ij}^\xi v_\xi = -Y_{ij}^\xi \frac{\lambda_P v_{BL}^2}{4M_\xi^2}v_\Delta.
$$

(11)

As $v_\Delta \simeq v_\xi \sim \mathcal{O}(1)$ GeV, we obtain sub-eV neutrino masses. Here it is worth noticing that the mixing between the super heavy triplet scalar $\Delta$ and the TeV scale scalar triplet $\xi$ gives rise to the neutrino mass. Essentially this set up can be thought of in an effective manner. After the $U(1)_{B-L}$ breaking, $\xi$ develops an effective trilinear coupling with the SM Higgs i.e. $\mu_\xi \xi^\dagger H H$ where $\mu_\xi$ is given by $\mu_\xi = \mu_\Delta \langle \Phi_{BL} \rangle^2/M_\Delta^2$. And this effective coupling is similar to the conventional type-II seesaw which leads to the generation of neutrino mass in this scenario.

**B. SCALAR MASSES AND MIXING**

As already discussed in the section II, the only significant mixing relevant for low energy phenomenological aspects is the mixing between $H$, $\xi$ and $\Phi_3$ since all other mixings are insignificant and can be neglected. In this section we only consider the light scalar sector, $V^L(H, \xi, \Phi_3)$ which is relevant for low energy phenomenology.

The minimization conditions for the scalar potential are given by:

$$
\mu_H^2 = \frac{1}{2} \left( \lambda_H \Phi_3 v_3^2 + 2\lambda_H v_H^2 + v_\xi^2(\lambda_{\xi H} + \lambda'_{\xi H}) - 2\sqrt{2}\mu_\xi v_\xi \right)
$$

$$
M_\xi^2 = \frac{1}{2} \left( -\lambda_{\xi H} v_3^2 - v_H^2(\lambda_{\xi H} + \lambda'_{\xi H}) + \frac{\sqrt{2}\mu_\xi v_H^2}{v_\xi} - 2v_\xi^2(\lambda_\xi + \lambda'_\xi) \right)
$$

$$
\mu_{\Phi_3}^2 = \frac{1}{2} \left( 2\lambda_{\Phi_3} v_3^2 + \lambda_H \Phi_3 v_H^2 + \lambda_\xi \Phi_3 v_\xi^2 \right)
$$

(12)
CP even scalar sector:

\[
\mathcal{L}_{\text{mass}} = \frac{1}{2} (h_H \ h_\xi \ h_3) \begin{pmatrix}
2\lambda_H v_H^2 & v_H v_\xi (\lambda_{\xi H} + \lambda'_{\xi H}) & \lambda_{H\Phi_3} v_3 v_H \\
v_H v_\xi (\lambda_{\xi H} + \lambda'_{\xi H}) & \mu_{\xi}^2 v_\xi^2 + 2v_\xi^2 (\lambda + \lambda'_\xi) & \lambda_{\xi\Phi_3} v_3 v_\xi \\
\lambda_{H\Phi_3} v_3 v_H & \lambda_{\xi\Phi_3} v_3 v_\xi & 2\lambda_{\Phi_3} v_3^2
\end{pmatrix}
\begin{pmatrix}
(h_H) \\
(h_\xi) \\
(h_3)
\end{pmatrix}
\]

\[
= \frac{1}{2} (h_H \ h_\xi \ h_3) \mathcal{M}_{\text{CP Even}}^{CP}
\begin{pmatrix}
h_H \\
h_\xi \\
h_3
\end{pmatrix}
\]

\[
= \frac{1}{2} (H_1 \ H_2 \ H_3) \begin{pmatrix}
\mathcal{O}^T \mathcal{M}_{\text{CP-Even}} \mathcal{O}
\end{pmatrix}
\begin{pmatrix}
H_1 \\
H_2 \\
H_3
\end{pmatrix}
\]

\[
= \frac{1}{2} (H_1 \ H_2 \ H_3) \begin{pmatrix}
m_{H_1}^2 & 0 & 0 \\
0 & m_{H_2}^2 & 0 \\
0 & 0 & m_{H_3}^2
\end{pmatrix}
\begin{pmatrix}
H_1 \\
H_2 \\
H_3
\end{pmatrix}
\]

(13)

Here \( \mathcal{O} \) is the orthogonal matrix which diagonalises the CP-even scalar mass matrix. Thus the flavor eigen states and the mass eigen states of these scalars are related by:

\[
\begin{pmatrix}
h_H \\
h_\xi \\
h_3
\end{pmatrix}
= \begin{pmatrix}
c_{12} c_{13} & c_{13} s_{12} & s_{13} \\
-c_{12} s_{13} s_{23} - c_{23} s_{12} & c_{12} c_{23} - s_{12} s_{13} s_{23} & c_{13} s_{23} \\
s_{12} s_{23} - c_{12} c_{23} s_{13} & -c_{12} s_{23} - c_{23} s_{12} s_{13} & c_{13} c_{23}
\end{pmatrix}
\begin{pmatrix}
H_1 \\
H_2 \\
H_3
\end{pmatrix}
\]

(14)

where we abbreviated \( \cos \beta_{ij} = c_{ij} \) and \( \sin \beta_{ij} = s_{ij} \), with \( \{ij : 12, 13, 23\} \).

CP odd scalar:

Denoting the massive CP odd scalar state that emerges as \( A^0 \), its mass is given by:

\[
m_{A^0}^2 = \frac{\mu_\xi (v_H^2 + 4v_\xi^2)}{\sqrt{2}v_\xi}
\]

(15)

Singly charged scalar:

The mass of the massive singly charged scalar state, \( H^\pm \) is given by:

\[
m_{H^\pm}^2 = \frac{2\sqrt{2}\mu_\xi v_H^2 - \lambda'_{\xi H} v_\xi v_H^2 - 2\lambda'_{\xi H} v_\xi^3 + 4\sqrt{2}\mu_\xi v_\xi^2}{4v_\xi}
\]

(16)

Doubly charged scalar:
The mass of the doubly charged scalar, $H^{±±}$ is given by:

$$m_{H^{±±}}^2 = -\frac{\lambda_{H} v_{H}^2}{2} + \frac{\mu_{H} v_{H}^2}{\sqrt{2} v_{ξ}} - \frac{\lambda_{ξ} v_{ξ}^2}{2}$$  \hspace{1cm} (17)

$U(1)_{\text{B-L}}$ Gauge Boson mass: The $Z_{\text{BL}}$ boson acquires mass through the vevs of $\Phi_{\text{BL}}, \Phi_{12}, \Phi_{3}$ which are charged under $U(1)_{\text{B-L}}$ and is given by:

$$M_{Z_{\text{BL}}}^2 \simeq g_{\text{BL}}^2 (v_{\text{BL}}^2 + 64v_{t2}^2 + 100v_{3}^2).$$  \hspace{1cm} (18)

The quartic couplings of scalars are expressed in term of physical masses and are expressed as:

$$\lambda_{H} = \frac{c^2_{13} (c^2_{12} m_{H_1}^2 + m_{H_2}^2 s_{12}^2) + m_{H_3}^2 s_{13}^2}{2 v_{H}^2}$$

$$\lambda_{ξH} = \frac{c_{13} s_{13} s_{23} (-c^2_{12} m_{H_1}^2 - m_{H_2}^2 s_{12}^2 + m_{H_3}^2) + c_{12} c_{13} c_{23} s_{12} (m_{H_2}^2 - m_{H_3}^2) + 4 m_{H_±}^2}{v_{H} v_{ξ}} - \frac{2 m_{λ0}^2}{v_{H}^2 + 4 v_{ξ}^2}$$

$$\lambda_{ξ} = \frac{s_{23}^2 (c^2_{12} m_{H_1}^2 + m_{H_2}^2 s_{12}^2) + c^2_{13} m_{H_3}^2}{2 v_{ξ}^2} + \frac{2 c_{12} c_{23} s_{12} s_{23} (m_{H_1} - m_{H_2}) (m_{H_1} + m_{H_2})}{2 v_{ξ}^2} + \frac{4 m_{H_±}^2}{v_{H}^2 + 4 v_{ξ}^2} - \frac{2 m_{λ0}^2}{v_{H}^2 + 4 v_{ξ}^2}$$

$$\lambda'_{ξH} = \frac{4 m_{λ0}^2}{v_{H}^2 + 4 v_{ξ}^2}$$

$$\lambda_{φ_3} = \frac{m_{H_1}^2 (c_{12} c_{23} s_{13} - s_{12} s_{23})^2 + m_{H_2}^2 (c_{12} s_{23} + c_{23} s_{12} s_{13})^2 + c_{13} c_{23} m_{H_3}^2}{2 v_{3}^2}$$

$$\lambda_{Hφ_3} = \frac{c_{13} c_{23} s_{13} (-c_{12} m_{H_1}^2 - m_{H_2}^2 s_{12}^2 + m_{H_3}^2) + c_{12} c_{13} s_{12} s_{23} (m_{H_1} - m_{H_2}) (m_{H_1} + m_{H_2})}{v_{3} v_{H}}$$

$$\lambda_{ξφ_3} = \frac{m_{H_1}^2 (c_{12} s_{13} s_{23} + c_{23} s_{12}) (c_{12} c_{23} s_{13} - s_{12} s_{23}) - m_{H_2}^2 (c_{12} s_{23} + c_{23} s_{12} s_{13}) (c_{12} c_{23} - s_{12} s_{13} s_{23})}{v_{3} v_{ξ}} + \frac{c_{13} c_{23} m_{H_3}^2 s_{12}}{v_{3} v_{ξ}}$$  \hspace{1cm} (19)

**Constraints on scalar sector:**

Based on the measurement of the $ρ$ parameter $ρ = 1.0008^{+0.0017}_{-0.0016}$ ([21, 46]), as discussed in Sec.II.A, the triplet $ξ$ vev can have an upper bound of order 2.5-4.6 GeV ([68]). Also the mixing angle between the SM Higgs and the triplet scalar is constrained from Higgs decay measurement. As obtained by [69], this mixing angle $\sin β_{12}$ is bounded above, in
particular, $\sin \beta_1 \lesssim 0.05$ to be consistent with experimental observation of $H_1 \to WW^*$ [46, 69]. There are similar bounds on singlet scalar mixing with the SM Higgs boson. Such bounds come from both theoretical and experimental constraints [70, 71]. The upper bound on singlet scalar-SM Higgs mixing angle $\sin \beta_3$ comes from W boson mass correction [72] at NLO. For $250 \text{ GeV} < m_{H_3} < 850 \text{ GeV}$, $\sin \beta_3$ is constrained to be $\sin \beta_3 < 0.2 - 0.3$ where $m_{H_3}$ is the mass of the third physical Higgs.

For our further discussion we consider the following benchmark points where all the above mentioned constraints are satisfied:

\[ \{ m_{H_2} = 331.779, \ m_{H_3} = 366.784, \ m_{A^0} = 33, \ m_{H^\pm} = 369.841, \ m_{H^{\pm\pm}} = 404.343 \ v_\xi = 2.951 \ (\text{in GeV}); \ s_{12} = 0.03, \ s_{23} = s_{13} = 0.01 \} \quad (20) \]

III. DARK MATTER

The $U(1)_{B-L}$ gauge symmetry gets spontaneously broken down by the vev of $\Phi_{12}, \Phi_{B-L}$ and $\Phi_3$ to a remnant $Z_2$ symmetry under which the exotic fermions: $\chi_{iR} \ (i = e, \mu, \tau)$ are assumed to be odd, while all other particles transforms trivially. As a result the lightest among these fermions becomes a viable candidate of DM and can give rise to the observed relic density by thermal freeze-out mechanism.

A. THE EXOTIC FERMIONS AND THEIR INTERACTIONS

From Eq. 3, the mass matrix for $\chi_{iR} \ (i = e, \mu, \tau)$ in the effective theory can be written as:

\[ -\mathcal{L}_X^{\text{mass}} = \frac{1}{2} \begin{pmatrix} (\chi_{eR})^c & (\chi_{\mu R})^c & (\chi_{\tau R})^c \end{pmatrix} \mathcal{M} \begin{pmatrix} \chi_{eR} \\ \chi_{\mu R} \\ \chi_{\tau R} \end{pmatrix} \]

(21)

where

\[ \mathcal{M} = \frac{1}{\sqrt{2}} \begin{pmatrix} Y_{11}v_{12} & Y_{12}v_{12} & Y_{13}v_{bl} \\ Y_{12}v_{12} & Y_{22}v_{12} & Y_{23}v_{bl} \\ Y_{13}v_{bl} & Y_{23}v_{bl} & Y_{33}v_3 \end{pmatrix} = \begin{pmatrix} [M_{12}] & [M'] \\ [M']^T & M_3 \end{pmatrix} \]

(22)

Here $M_{12}, M', M_3$ are:
For simplicity we assume $Y_{11} = Y_{22}$ and $Y_{13} = Y_{23}$. As a result the above Majorana fermion mass matrix $\mathcal{M}$ can be exactly diagonalized by an orthogonal rotation $\mathcal{R} = \mathcal{R}_{13}(\theta)\mathcal{R}_{23}(\theta_{23} = 0)\mathcal{R}_{12}(\theta_{12} = \frac{\pi}{4})$ which is essentially characterized by only one parameter $\theta$ [73]. So we diagonalized the mass matrix $\mathcal{M}$ as $\mathcal{R} \mathcal{M} \mathcal{R}^T = \mathcal{M}_{\text{Diag.}}$, where the $\mathcal{R}$ is given by:

$$
\mathcal{R} = \begin{pmatrix}
\frac{1}{\sqrt{2}} \cos \theta & \frac{1}{\sqrt{2}} \cos \theta & \sin \theta \\
-\frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} & 0 \\
-\frac{1}{\sqrt{2}} \sin \theta & -\frac{1}{\sqrt{2}} \sin \theta & \cos \theta 
\end{pmatrix}
$$

(23)

Where the rotation parameter $\theta$ required for the diagonalization is given by:

$$
\tan 2\theta = \frac{2\sqrt{2} Y_{13} v_{\text{BL}}}{(Y_{11} + Y_{12}) v_{12} - Y_{33} v_{3}}
$$

(24)

Thus the physical states of the exotic fermions are $\chi_i = \frac{\chi_{1R} + (\chi_{1R})^c}{\sqrt{2}}$ and are related to the flavor eigen states by the following linear combinations:

$$
\chi_{1R} = \frac{1}{\sqrt{2}} \cos \theta \chi_{eR} + \frac{1}{\sqrt{2}} \cos \theta \chi_{\mu R} + \sin \theta \chi_{\tau R}
$$

$$
\chi_{2R} = -\frac{1}{\sqrt{2}} \chi_{eR} + \frac{1}{\sqrt{2}} \chi_{\mu R}
$$

$$
\chi_{3R} = -\frac{1}{\sqrt{2}} \sin \theta \chi_{eR} - \frac{1}{\sqrt{2}} \sin \theta \chi_{\mu R} + \cos \theta \chi_{\tau R}
$$

(25)

And the corresponding mass eigen values are given by:

$$
M_1 = \frac{1}{2\sqrt{2}} \left[ (Y_{11} + Y_{12}) v_{12} + Y_{33} v_{3} + \sqrt{((Y_{11} + Y_{12}) v_{12} - Y_{33} v_{3})^2 + 8(Y_{13} v_{\text{BL}})^2} \right]
$$

$$
M_2 = \frac{1}{\sqrt{2}} (Y_{11} - Y_{12}) v_{12}
$$

$$
M_3 = \frac{1}{2\sqrt{2}} \left[ (Y_{11} + Y_{12}) v_{12} + Y_{33} v_{3} - \sqrt{((Y_{11} + Y_{12}) v_{12} - Y_{33} v_{3})^2 + 8(Y_{13} v_{\text{BL}})^2} \right]
$$

(26)

Here it is worthy to mention that in the limit of $Y_{13} \to 0$ i.e. $\theta \to 0$, we get the mass eigen values of the DM particles as $M_{1,2} = \frac{1}{\sqrt{2}} (Y_{11} \pm Y_{12}) v_{12}$ and $M_3 = \frac{1}{\sqrt{2}} Y_{33} v_{3}$ and the corresponding mass eigen states are $\chi_{1R,2R} = \frac{1}{\sqrt{2}} (\chi_{\mu R} \pm \chi_{eR})$ and $\chi_{3R} = \chi_{\tau R}$. If we assume that the off diagonal Yukawa coupling in $M_{12}$ i.e. $Y_{12} << 1$, then $\chi_1$ and $\chi_2$ become almost degenerate (i.e $M_1 \simeq M_2$).
Assuming $\chi_3$ to be the lightest state makes it the viable DM candidate and $\chi_1$ and $\chi_2$ are the next to lightest stable particles (NLSP). Using the relation $R\cdot M\cdot R^T = M_{\text{Diag.}}$, one can express the following relevant parameters in terms of the physical masses $M_1, M_3$ and the mixing angle $\sin \theta$ as

$$v_3 = \frac{\sqrt{2}}{Y_{33}} \left( M_1 \sin^2 \theta + M_3 \cos^2 \theta \right)$$
$$v_{12} = \frac{\sqrt{2}}{Y_{33} + Y_{12}} \left( M_1 \cos^2 \theta + M_3 \sin^2 \theta \right)$$
$$Y_{13} = \frac{\Delta M \sin 2\theta}{2v_{\text{BL}}}.$$ \hspace{1cm} (27)

where $\Delta M$ is the mass splitting between the DM and NLSPs i.e. $\Delta M = M_1 - M_3$.

The gauge coupling $g_{\text{BL}}$ can be expressed as

$$g_{\text{BL}} \simeq \frac{M_{Z_{\text{BL}}}}{\sqrt{(v_{\text{BL}}^2 + 64v_{12}^2 + 100v_3^2)}}.$$ \hspace{1cm} (28)

The flavor eigen states can be expressed in terms of the physical eigen states as follows:

$$\chi_{e_R} = \frac{1}{\sqrt{2}} \cos \theta \chi_{1_R} - \frac{1}{\sqrt{2}} \chi_{2_R} - \frac{1}{\sqrt{2}} \sin \theta \chi_{3_R}$$
$$\chi_{\mu_R} = \frac{1}{\sqrt{2}} \cos \theta \chi_{1_R} + \frac{1}{\sqrt{2}} \chi_{2_R} - \frac{1}{\sqrt{2}} \sin \theta \chi_{3_R}$$
$$\chi_{\tau_R} = \sin \theta \chi_{1_R} + \cos \theta \chi_{3_R}.$$ \hspace{1cm} (29)

**DM Interactions**

The Yukawa and gauge interactions of DM relevant for the calculation of relic density can be written in the physical eigen states as follows:

$$\mathcal{L}_\text{Yuk.} = Y_{33} h_3 (\chi_{\tau_R})^c \chi_{\tau_R}$$
$$= Y_{33} \left[ (s_{12} s_{23} - c_{12} c_{23} s_{13})H_1 - (c_{12} s_{23} + c_{23} s_{12} s_{13})H_2 + (c_{13} c_{23})H_3 \right] (\chi_{\tau_R})^c \chi_{\tau_R}$$
$$= Y_{33} \left[ (s_{12} s_{23} - c_{12} c_{23} s_{13})H_1 - (c_{12} s_{23} + c_{23} s_{12} s_{13})H_2 + (c_{13} c_{23})H_3 \right]$$
$$\times \left[ \sin^2 \theta (\chi_{1_R})^c \chi_{1_R} + \cos^2 \theta (\chi_{3_R})^c \chi_{3_R} + \sin \theta \cos \theta \left( (\chi_{1_R})^c \chi_{3_R} + (\chi_{3_R})^c \chi_{1_R} \right) \right]$$ \hspace{1cm} (30)
and

\[
\mathcal{L}_{Z_{BL}} = g_{BL} \left[ (4 \cos^2 \theta + 5 \sin^2 \theta) \bar{\chi}_1 R \gamma^\mu \chi_1 R + 4 \bar{\chi}_2 R \gamma^\mu \chi_2 R \right. \\
\left. + (4 \sin^2 \theta + 5 \cos^2 \theta) \bar{\chi}_3 R \gamma^\mu \chi_3 R + \cos \theta \sin \theta \left( \bar{\chi}_1 R \gamma^\mu \chi_3 R + \bar{\chi}_3 R \gamma^\mu \chi_1 R \right) \right] (Z_{BL})_\mu \quad (31)
\]

Note that there is no co-annihilation of \( \chi_3 \) with \( \chi_2 \).

The dominant annihilation and co-annihilation channels for DM are shown in Fig 2-5 and Fig 6-9 respectively.

![Feynman diagrams for DM Annihilation: \( \chi_3 \chi_3 \rightarrow f\bar{f} \).](image1)

![Feynman diagrams for DM Annihilation: \( \chi_3 \chi_3 \rightarrow H_i H_j \) \((i,j,k = 1,2,3)\).](image2)

![Feynman diagrams for DM Annihilation: \( \chi_3 \chi_3 \rightarrow H_i Z_{BL} \) \((i = 1,2,3)\).](image3)
FIG. 5: Feynman diagrams for DM Annihilation: $\chi_3 \chi_3 \rightarrow Z_{BL} Z_{BL}; Z \ Z; W^+ W^-.$

FIG. 6: Feynman diagrams for DM Co-Annihilation: $\chi_3 \chi_1 \rightarrow f \bar{f}.$

FIG. 7: Feynman diagrams for DM Co-Annihilation: $\chi_3 \chi_1 \rightarrow H_i H_j (i,j,k = 1,2,3).$

FIG. 8: Feynman diagrams for DM Co-Annihilation: $\chi_3 \chi_1 \rightarrow H_i Z_{BL} (i = 1,2,3).$
FIG. 9: Feynman diagrams for DM Co-Annihilation: $\chi_3 \chi_1 \rightarrow Z_{BL} Z_{BL}; Z Z; W^+ W^-$. 

B. RELIC ABUNDANCE OF DM

The DM phenomenology is mainly governed by the following independent parameters:

$$\{ M_3 \equiv M_{DM}, \Delta M \equiv (M_1 - M_3) \simeq (M_2 - M_3), \sin \theta, Y_{33}, M_{Z_{BL}} \}.$$  \hspace{1cm} (32)

while the other independent parameters that are kept fixed are: $v_{BL} = 10$ TeV, $Y_{12} = 10^{-6}$, and the dependent parameters are $g_{BL}, v_3, v_{12}$ and $Y_{13}$ as mentioned in Eq 27,28.

The relic density of DM in this scenario can be estimated by solving the Boltzmann equation in the following form:

$$\frac{dn}{dt} + 3Hn = -\langle \sigma v \rangle_{eff} \left( n^2 - n_{eq}^2 \right),$$  \hspace{1cm} (33)

where $n$ denotes number density of DM, i.e. $n \sim n_{\chi_3}$ and $n_{eq} = g(M_{DM} T^3)/(2\pi^2) \exp(-M_{DM}/T)$ is equilibrium distribution. The DM freezes out giving us the thermal relic depending on $\langle \sigma v \rangle_{eff}$, which takes into account all number changing process listed in Fig 2-5 and Fig 6-9.

This can be written as:

$$\langle \sigma v \rangle_{eff} = \frac{g_2^3}{g_{eff}^2} \langle \sigma v \rangle_{\chi_3 \chi_3} + \frac{2g_3 g_1}{g_{eff}^2} \langle \sigma v \rangle_{\chi_3 \chi_1} \left( 1 + \frac{\Delta M}{M_{DM}} \right)^{\frac{3}{2}} \exp(-\frac{\Delta M}{M_{DM}})$$

$$+ \frac{g_2^2}{g_{eff}^2} \langle \sigma v \rangle_{\chi_1 \chi_1} \left( 1 + \frac{\Delta M}{M_{DM}} \right)^{\frac{3}{2}} \exp(-2x \frac{\Delta M}{M_{DM}}) + \frac{g_2^2}{g_{eff}^2} \langle \sigma v \rangle_{\chi_2 \chi_2} \left( 1 + \frac{\Delta M}{M_{DM}} \right)^{\frac{3}{2}} \exp(-2x \frac{\Delta M}{M_{DM}})$$  \hspace{1cm} (34)

Writing this equation in a precise form for convenience in discussion:

$$\langle \sigma v \rangle_{eff} = \frac{g_2^3}{g_{eff}^2} \langle \sigma v \rangle_{\chi_3 \chi_3} + \langle \sigma v \rangle_{\chi_3 \chi_1} f(\Delta M, M_{DM})$$

$$+ \langle \sigma v \rangle_{\chi_1 \chi_1} h_1(\Delta M, M_{DM}) + \langle \sigma v \rangle_{\chi_2 \chi_2} h_2(\Delta M, M_{DM})$$  \hspace{1cm} (35)

where $f, h_1$ and $h_2$ are the factors multiplied to the co-annihilation cross-sections which are functions of $\Delta M$. 

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Here $g_{\text{eff}}$ is the effective degrees of freedom which can be expressed as,

$$g_{\text{eff}} = g_3 + g_1 \left(1 + \frac{\Delta M}{M_{DM}}\right) \frac{3}{2} \exp(-x \frac{\Delta M}{M_{DM}}) + g_2 \left(1 + \frac{\Delta M}{M_{DM}}\right) \frac{3}{2} \exp(-x \frac{\Delta M}{M_{DM}})$$

(36)

where $g_3$, $g_1$, $g_2$ are the internal degrees of freedom of $\chi_3$, $\chi_1$ and $\chi_2$ respectively. The dimensionless parameter $x$ is defined as $x = \frac{M_{DM}}{M_f} = \frac{M_1}{T_f}$.

The relic density of the DM ($\chi_3$) then can be given by [74–76]:

$$\Omega_{\chi_3} h^2 = \frac{1.09 \times 10^9 \text{GeV}^{-1}}{g_*^{1/2} M_{Pl}} \frac{1}{J(x_f)}$$

(37)

where $g_* = 106.7$ and $J(x_f)$ is given by

$$J(x_f) = \int_{x_f}^{\infty} \frac{\langle \sigma v \rangle_{\text{eff}}}{x^2} dx$$

(38)

Here $x_f = \frac{M_{DM}}{T_f}$, and $T_f$ denotes the freeze-out temperature of the DM $\chi_3$. We may note here that for correct relic $x_f \simeq 25$.

It is worth mentioning here that we used the package MicrOmegas [77] for computing annihilation cross-section and relic density, after generating the model files using LanHEP [78].

1. PARAMETER SPACE SCAN

To understand the DM relic density and the specific role of the model parameters in giving rise to the observed relic density, we performed several analysis and scan for allowed parameter space. As discussed in section III A, the important relevant parameters controlling the relic abundance of DM are: the mass of DM ($M_{DM}$), mass splitting($\Delta M$) between the DM ($\chi_3$) and the next to lightest stable particle ($\chi_2$ and $\chi_1$ as $M_1 \simeq M_2$), and the mixing angle $\sin \theta$. Apart from these three, another crucial parameter which has a noteworthy effect on DM relic as well as other phenomenological aspects is the Yukawa coupling $Y_{33}$. We also keep the $B - L$ gauge boson mass ($M_{Z_{BL}}$) as a free parameter. The dependent parameters have already been mentioned in Eqs. 27 and 28. The other parameters that are kept fixed judiciously during the analysis are $v_{BL} = 10$ TeV and $Y_{12} = 10^{-6}$.

We show the variation of relic density of DM $\chi_3$ in Fig. 10 as a function of its mass $M_{DM}$ for different choices of $\Delta M$: 1-10 GeV, 20-50 GeV and 100-200 GeV shown by different colored points as mentioned in the inset of the figure. The dips in the relic density plots are essentially due to resonances corresponding to SM-Higgs, second Higgs and $Z_{BL}$ gauge bosons.
FIG. 10: Variation of relic density as a function of DM mass. Other parameters are kept fixed as mentioned inset of each figure.

respectively. In the top-left and top-right panel, $Y_{33}$ is varied in a range $0.1 \leq Y_{33} \leq 0.2$ where as in the bottom-left and bottom-right panel it is varied in an interval $0.9 \leq Y_{33} \leq 1$. Clearly as $Y_{33}$ increases, the effective annihilation cross-section increases which decreases the relic density. For a properly chosen value of $Y_{33}$ (with other parameters fixed), we can obtain the correct relic density for a wide range of DM mass ($M_{DM}$).

We can also analyze the effect of mixing angle $\sin \theta$ and mass splitting ($\Delta M$) from the results in Fig 10. Here it is worth mentioning that the parameter which decides the contribution of co-annihilations of DM to the relic density is $\sin \theta$ which can be understood by looking at Eq. 30 and 31. If $\sin \theta$ is small then contribution from annihilation of DM will dominate over all co-annihilation effects but for larger $\sin \theta$, co-annihilation contributions will be more as compared to the annihilations. The value of $\sin \theta$ predominantly decides the relative contribution of annihilation and co-annihilations of DM for the calculation of relic density. However the mass-splitting also plays a crucial role in the effect of annihilations and co-annihilations of DM. In the top and bottom right panel of Fig 10, $\sin \theta$ is randomly varied
FIG. 11: [Left]: Variation of $g_{\text{eff}}$ with $\Delta M$, [Right]: Variation of $f(\Delta M) = \frac{2g_1g_3}{g_{\text{eff}}^2} \left( 1 + \frac{\Delta M}{M_{DM}} \right)^3 \exp(-x\frac{\Delta M}{M_{DM}})$ and $h_{1,2}(\Delta M) = \frac{g_{1,2}^2}{g_{\text{eff}}^2} \left( 1 + \frac{\Delta M}{M_{DM}} \right)^3 \exp(-2x\frac{\Delta M}{M_{DM}})$ with $\Delta M$.

in a range $0.85 \leq \sin \theta \leq 0.98$ for two different ranges of $Y_{33}$. For such a large $\sin \theta$, DM annihilates very weakly, so the co-annihilations essentially decides the effective annihilation cross-section and hence the relic density. This means in Eq. 34, the first term is negligible as compared to the other terms. In such a case, as $\Delta M$ increases, these co-annihilations become less and less effective thus decreasing the effective annihilation cross-section hence increasing the relic density. This trend is clearly observed in the right panel plots of Fig 10. The effect of mass splitting in such a case can also be understood by looking at the right panel of Fig 11 where the multiplying functions (mentioned as $f(\Delta M)$ and $h(\Delta M)$) in the co-annihilation terms of the effective annihilation cross-section in Eq. 35, are plotted as a function of mass-splitting $\Delta M$. As $\Delta M$ increases, these factors decreases drastically consequently decreasing the overall effective annihilation cross-section and hence increasing the relic density of DM.

However if we consider the case of smaller $\sin \theta$ as considered for the left panel plots of Fig 10 (i.e. $0.1 \leq \sin \theta \leq 0.2$), here DM annihilation is the most effective and hence dominantly decides the relic density and except the first term in Eq. 34, other terms are negligible. In this case, with increase in mass splitting the effective thermal averaged cross-section increases and relic density decreases. This is due to the fact that, when $\Delta M$ increases,
the effective degrees of freedom $g_{\text{eff}}$ decreases which is shown in the left panel plot of Fig 11 for a benchmark value of $M_{DM}$. This in turn increases the $\langle \sigma v \rangle_{\text{eff}}$ and hence results in decrease in the DM relic abundance.

FIG. 12: Relic density allowed parameters space in $M_{DM} - Y_{33}$ plane for different intervals of $\sin \theta_{13}$.

To make the analysis more robust, in Fig. 12, the correct relic density allowed parameter space has been shown in the plane of $Y_{33}$ vs $M_{DM}$ for wide range of mixing angle \{\sin \theta = 0.1 - 0.3, 0.3 - 0.5, 0.5 - 0.7, 0.7 - 0.98\}, indicated by different colors. To carry out this scan of parameter space, $\Delta M$ is varied randomly within 1 to 1000 GeV. We can see that at the resonance regions, any value of $Y_{33}$ can give rise to correct relic abundance. However for other values of $M_{DM}$, a relatively larger $Y_{33}$ is required to obtain the correct relic.

In Fig 13, a scan similar to Fig 12 is carried out for correct relic density allowed parameter space in the $M_{DM}$ vs $Y_{33}$ plane, with $\Delta M$ varied randomly in different intervals for 3 different range of $\sin \theta$.

From this analysis, the riveting feature of this model that comes out is that, as compared to the earlier $U(1)_{B-L}$ models with right handed neutrino (RHN) dark matter where the correct relic density is usually achieved for DM mass near the resonance (i.e. $M_{DM} \sim M_{Z_{\text{BL}}}/2$)[56, 60–62], the crucial difference here is; in addition to the resonances, we obtain a large parameter space satisfying correct relic density for larger Yukawa coupling $Y_{33}$ due to co-annihilations among the dark sector particles.
IV. DETECTION PROSPECTS OF DM

A. DIRECT DETECTION

There are various attempts to detect DM. One such major experimental procedure is the Direct detection of the DM at terrestrial laboratories through elastic scattering of the DM off nuclei. Several experiments put strict bounds on the dark matter nucleon cross section like LUX [9], PandaX-II [10, 11] and XENON-1T [12, 13]. In this model, the DM-nucleon scattering is possible via Higgs-mediated interaction represented by the Feynman diagram shown in Fig 14. Here, it is worth mentioning that the DM being a Majorana fermion has only the off diagonal (axial vector) couplings with the $Z_{BL}$ boson and therefore do not contribute to spin independent direct search.

The cross section per nucleon for the spin-independent (SI) DM-nucleon interaction is given by:

$$\sigma_{SI} = \frac{1}{\pi A^2} \mu_r^2 |\mathcal{M}|^2,$$  \hspace{1cm} (39)
FIG. 14: Higgs-mediated DM-nucleon scattering

where $A$ is the mass number of the target nucleus, $\mu_r$ is the reduced mass of the DM-nucleon system and $\mathcal{M}$ is the amplitude for the DM-nucleon interaction, which can be written as:

$$\mathcal{M} = \left[Z f_p + (A - Z) f_n \right], $$

where $f_p$ and $f_n$ denote effective interaction strengths of DM with proton and neutron of the target used with $A$ being mass number and $Z$ being atomic number. The effective interaction strength can then further be decomposed in terms of interaction with partons as:

$$f^{p,n}_{p,n} = \sum_{i=1}^{3} \alpha^i_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{p,n}^{T_G} \sum_{Q=c,t,b} \alpha_Q \frac{m_{p,n}}{m_Q}, $$

with

$$\begin{align*}
\alpha^1_q &= -Y_{33} \cos^2 \theta \frac{m_q}{v_H} \left[ \frac{(s_{12}s_{23} - c_{12}c_{23}s_{13})^2}{m_{H_1}^2} \right] \\
\alpha^2_q &= -Y_{33} \cos^2 \theta \frac{m_q}{v_H} \left[ \frac{(c_{12}s_{23} + c_{23}s_{12}s_{13})^2}{m_{H_2}^2} \right] \\
\alpha^3_q &= -Y_{33} \cos^2 \theta \frac{m_q}{v_H} \left[ \frac{(c_{13}c_{23})^2}{m_{H_3}^2} \right]
\end{align*}$$

coming from DM interaction with SM via Higgs portal coupling. In Eq.41, the different coupling strengths between DM and light quarks are given in ref. [1, 79] as $f_{T_u}^p = 0.020 \pm 0.004, f_{T_d}^p = 0.026 \pm 0.005, f_{T_s}^p = 0.014 \pm 0.062, f_{T_u}^n = 0.020 \pm 0.004, f_{T_d}^n = 0.036 \pm 0.005, f_{T_s}^n = 0.118 \pm 0.062$. The coupling of DM with the gluons in target nuclei is parameterized by:

$$f_{T_G}^{p,n} = 1 - \sum_{q=u,d,s} f_{T_q}^{p,n}.$$
In the context of DM direct search, the model parameters that enter the DM-nucleon direct search cross-section, are the Higgs-DM Yukawa coupling ($Y_{33}$) and the mixing angle ($\sin \theta$), which can be constrained by requiring that $\sigma_{SI}$ is less than the current DM-nucleon cross-sections dictated by non-observation of DM in current direct search data. In Fig. 15, we show the DM-nucleon cross section mediated by scalars in comparison to the latest XENON1T bound. In left panel of Fig. 15, we confronted the points satisfying relic density with the spin independent DM-nucleon elastic scattering cross section obtained for the model as a function of DM mass. The XENON1T bound is shown by dashed black line. Thus the region below this line satisfy both relic density as well as direct detection constraint. We can see that, though for DM mass at the resonance regions, $\sin \theta$ values $0.1 - 0.98$ can satisfy the direct detection constraint but for DM masses other than at the resonances, only larger $\sin \theta$ values ($0.7 \leq \sin \theta \leq 0.98$) are favored which is indicated by the orange points. As we have already discussed that in the larger $\sin \theta$ regime, the relic density is governed predominantly through the co-annihilations of DM, so this result interestingly implies that the co-annihilation effect essentially enhances the parameter space that satisfies the direct search constraints other than the resonance regions.

---

**FIG. 15:** [Left]: Spin-independent direct detection cross section of DM with nucleon as a function of DM mass (in GeV) confronted with XENON-1T data over and above relic density constraint from PLANCK. [Right]: Points satisfying both relic density constraint and Direct Detection constraint are projected in the $Y_{33} - M_{DM}$ plane. Different colors depict the different range of values of $\sin \theta$ as mentioned in the figure inset.

We again show the points satisfying relic density as well as direct detection constraint from XENON-1T in the plane of $Y_{33}$ vs $M_{DM}$ in the right panel of Fig. 15. Except for the
resonances, where all values of $Y_{33}$ \textit{(i.e.} 0.1 – 0.98), satisfies both relic and DD constraints, for other values of DM mass, only with larger $Y_{33}$; 0.6 $\leq Y_{33} \leq$ 1 these constraints can be satisfied. This plot again reinforces the fact that, though for DM mass near the resonances, all sin$\theta$ values 0.1 – 0.98 are allowed by the direct detection constraint but for DM mass other than at the resonances, only larger sin$\theta$ values (0.7 $\leq$ sin$\theta$ $\leq$ 0.98) are favored.

**B. INDIRECT DETECTION**

Apart from direct detection experiments, DM can also be probed at different indirect detection experiments which essentially search for SM particles produced through DM annihilations. Among these final states, photon and neutrinos, being neutral and stable can reach the indirect detection experiments without getting affected much by intermediate medium between the source and the detector. For DM in the WIMP paradigm, these photons lie in the gamma ray regime and hence can be measured at space based telescopes like the Fermi Large Area Telescope (FermiLAT) or ground based telescopes like MAGIC. Measuring the gamma ray flux and using the standard astrophysical inputs, one can constrain the DM annihilation into different final states like $\mu^+\mu^-$, $\tau^+\tau^-$, $W^+W^-$, $b\bar{b}$. Since DM can not couple to photons directly, gamma rays can be produced from such charged final states. Using the bounds on DM annihilation to these final states from the indirect detection bounds arising from the global analysis of the Fermi-LAT and MAGIC observations of dSphs \cite{14, 15}, we check for the constraints on our DM parameters. Since there are multiple annihilation channels to different final states, the Fermi-LAT constraints on individual final states are weak for most of the cases. In Fig. 16(Left panel), we show the points satisfying both relic constraint and direct search constraint confronted with the constraints from indirect detection from MAGIC+FermiLAT for annihilation of DM to $W^+W^-$ which is the most stringent as compared to DM annihilation to other channels. The combined bound from MAGIC and FermiLAT is shown by the black dotted line. The points below this line are allowed by relic, direct and indirect search constraints. In the right panel of Fig. 16, we represent the parameter space allowed by relic, direct as well as indirect search constraints in the plane of $Y_{33}$ vs $M_{DM}$. Here also we see that the parameter space that satisfies relic and direct detection governed through the co-annihilation effects, are also consistent with the constraints from indirect search of DM.
FIG. 16: [Left] $\langle \sigma v \rangle_{\chi^3 \chi^3 \rightarrow W^+ W^-}$ is shown as a function of $M_{DM}$. Only the points satisfying both Relic and DD constraints are shown. [Right] Relic+DD+ID allowed parameters are plotted in $M_{DM} - Y_{33}$ plane. Other parameters are kept fixed as mentioned in the inset of each figure.

C. COLLIDER CONSTRAINTS ON $g_{BL} - M_{Z_{BL}}$

Apart from constraints from relic density and direct, indirect search of DM, there exists stringent experimental constraints on the $B-L$ gauge sector from colliders like ATLAS, CMS and LEP-II. There exists a lower bound on the ratio of new gauge boson mass to the new gauge coupling $M_{Z'}/g' \geq 7$ TeV from LEP-II data [80, 81]. However the bounds from the current LHC experiments have already surpassed the LEP II limits. In particular, search for high mass Di-lepton resonances have put strict bounds on such additional gauge sector. The latest bounds from the ATLAS experiment [82, 83] and the CMS experiment [84] at the LHC rule out such gauge boson masses below 4.3 TeV for $g_{BL}$ of the same order as that of SM coupling, from analysis of 13 TeV data. However such bounds get weaker, if the corresponding gauge couplings are smaller [82] than the electroweak gauge couplings.

In Fig. 17, we show a parameter scan in the plane of $g_{BL}$ vs $M_{DM}$ to scrutinize our parameter space with respect to the constraints from ATLAS and LEP-II. The bounds on $g_{BL}$ for a fixed $M_{Z_{BL}} (= 1000$ GeV) from both LEP-II and ATLAS are shown by dotted black lines. It is clear that the constraint from LEP-II is much weaker than the constraints from ATLAS. Only those points which lie below this black dotted line is allowed from all the relevant constraints (i.e. Relic + Direct Detection + Indirect Detection + ATLAS). The different colored points depict different $Y_{33}$ values. From this analysis, as it is clear from Fig. 17, it can be inferred that all the relevant constraints as discussed above are getting
FIG. 17: $g_{BL}$ vs $M_{DM}$ plot. ATLAS and LEP-II bounds are shown for $M_{Z_{B-L}} = 1$ TeV.

satisfied at the resonance regions. However, some of the points corresponding to larger $Y_{33}$ values beyond 0.4 are also allowed from all these constraints for $M_{DM}$ beyond 1 TeV.

$1 \text{ GeV} \leq M_{Z_{BL}} \leq 4000 \text{ GeV}$:

So far whatever analysis we have done is with a fixed mass of the $B-L$ gauge boson $M_{Z_{BL}} = 1000$ GeV. But if we vary the mass of $Z_{BL}$ boson in a range $1 \text{ GeV} \leq M_{Z_{BL}} \leq 4000$ GeV, then the parameter space is shown in Fig. 18. This figure is similar to the Fig. 15, the only difference being the variation of $M_{Z_{BL}}$ in the later one. All the points below the black dotted line, are allowed by the relic as well as the direct detection constraint. On the right panel of Fig. 18, the parameter space consistent with all the constraints except from colliders (ATLAS) is shown in the plane of $Y_{33}$ vs $M_{DM}$. Different colors are used to specify the value of $\sin \theta$ in different intervals during the scan.

We now turn to find the allowed parameter space in light of ATLAS bound on $g_{BL} - M_{Z_{BL}}$. The constraint on $g_{BL}$ for corresponding values of $M_{Z_{BL}}$ coming from the non-observation of a new gauge boson ($Z_{BL}$) at LHC from ATLAS [82] analysis is shown by the black thick dotted line in right panel of Fig. 19. This indicates that points below the line with smaller $g_{BL}$ is allowed, while those above the line are ruled out. The plot shows points which satisfy relic density constraint, direct as well as indirect search constraints. Different colors indicate
FIG. 18: [Left]: Spin-independent direct detection cross section of DM with nucleon as a function of DM mass (in GeV) confronted with XENON-1T data with $M_{Z_{BL}}$ varied in a range of 1 GeV to 4000 GeV. All points are satisfied the relic density constraint from PLANCK. [Right]: Points satisfying all these constraints (Relic+ Direct Detection+ Indirect Detection) projected in the $Y_{33}$ vs $M_{DM}$ plane.

ranges of $Y_{33}$ as mentioned in figure inset.

FIG. 19: Relic, direct detection and indirect detection satisfied points are shown in $M_{DM} - M_{Z_{BL}}$ plane in left figure and $g_{BL} - M_{Z_{BL}}$ plane in right figure with different range of $Y_{33}$. The thick black dotted line in right panel figure shows the ATLAS bound on $g_{BL}$ vs $M_{Z_{BL}}$ plane from non-observation of $Z_{BL}$ at colliders.

We then showcase the final parameter space in the plane of $M_{DM}$ vs $M_{Z_{BL}}$ and $Y_{33}$ vs $M_{DM}$ after imposing the bounds from correct relic density of DM, direct and indirect detection of DM and search for $B - L$ gauge boson at ATLAS which are shown in Fig. 20. If we compare the plot of Fig. 18 and Fig. 19 with the plot of Fig. 20, we can infer how much parameter space
FIG. 20: Parameter space satisfying relic density, direct and indirect detection bound as well as $g_{BL} - M_{Z_{BL}}$ constraint from ATLAS is shown in the plane of $M_{DM}$ vs $M_{Z_{BL}}$ (left) and $Y_{33}$ vs $M_{DM}$ (Right).

gets ruled out by the imposition of the constraints from ATLAS. We also show the parameter space satisfying correct relic, direct and indirect constraint along with the constraint from ATLAS on $g_{BL} - M_{Z_{BL}}$ in the plane of $M_{DM}$ vs $M_{Z_{BL}}$ (left) and $Y_{33}$ vs $M_{DM}$ (right) in Fig 20. The points along the diagonal of the plot essentially corresponds to the points at the $Z_{BL}$ resonance which occurs for $M_{DM} \sim M_{Z_{BL}}/2$. But as we can see clearly that apart from these points, all relevant constraints are also getting satisfied for a DM Mass range $M_{DM} > M_{Z_{BL}}$.

This enhanced parameter space is the result of additional annihilation and co-annihilations possible in this considered model. The different colors of the points corresponds to different ranges of the Yukawa coupling $Y_{33}$. Clearly as the DM mass increases, a larger $Y_{33}$ can lead to correct relic as well as all other constraints discussed in this context.

V. COLLIDER SIGNATURE OF DOUBLY CHARGED SCALAR IN PRESENCE OF $Z_{BL}$

The light doubly charged scalar in this model offers novel multi-lepton signatures with missing energy and jets. It is worthy to mention here that the dark sector which contains the gauge singlet Majorana fermions do not have any promising collider signatures as the mono-X type signal processes, arising out of initial state radiation are extremely suppressed. The doubly charged scalar, $H^{\pm\pm}$ which is also charged under $U(1)_{B-L}$ can be produced at Large Hadron Collider (LHC) via Higgs ($H_{1,2,3}$) and gauge bosons ($\gamma$, $Z$, $Z_{BL}$) mediations.
Further decay of $H^{\pm\pm}$ to $W^\pm W^\pm$ pair (assumed $m_{H^{\pm\pm}} \geq 2m_W$) with almost 100% branching ratio for $v_\xi \sim 2.951$ GeV yields $W^+W^+W^-W^-$ final state. As a result the four $W$ final state offers: $4\ell + E_T$ and $m_\ell + n_j + E_T$ signatures at collider. For details of branching fraction and partial decay widths of $H^{++}$ with $v_\xi$, please see appendix B. Although this type of signatures have been studied in the context of Type -II see-saw model, the triplet scalar $\xi$ considered in this model also have $U(1)_{B-L}$ charge on top SM gauge charges and that makes this model different from the usual type-II seesaw scenario. In this section we will briefly highlight the effect of additional gauge boson $Z_{BL}$ on the pair production cross-section of doubly charged scalar. The corresponding Feynman diagram of this type process is shown in Fig.21.

The pair production cross-section of doubly charged scalar, $H^{++}H^{--}$ as function of mass, $m_{H^{\pm\pm}}$ for fixed value of $M_{Z_{BL}} = 4.5$ TeV with $\sqrt{s} = 13$ TeV is shown in Fig.22. The production cross-sections are computed in MicrOmegas using the NNPDF23 parton distributions. The black solid line corresponds to the case where $U(1)_{B-L}$ gauge boson, $Z_{BL}$ is absent and the scenario resembles the usual type-II see-saw scenario. And in that case the $H^{++}H^{--}$ pair can be produced via SM Higgs and SM gauge boson $(\gamma, Z)$ mediated Drell-Yan processes. However in a gauged $B-L$ scenario, the presence of the additional gauge boson $Z_{BL}$ can affect this pair production cross-section of $H^{++}H^{--}$. The effects of $U(1)_{B-L}$ gauge boson on top of SM gauge bosons are shown by dotted lines in the Fig.22 for two different values of gauge couplings: $g_{BL} = 0.44$ (purple line) and 0.55 (blue line). It is important to
FIG. 22: Production cross-section for $p p \rightarrow H^{++} H^{--}$ as a function of doubly charged scalar mass $m_{H^{\pm\pm}}$ considering $\text{Br}(H^{\pm\pm} \rightarrow W^{\pm}W^{\pm}) \sim 100\%$ at $\sqrt{s} = 14$ TeV. The black solid line corresponds to the usual type-II seesaw scenario where $Z_{BL}$ gauge mediated diagrams are absent. The effects of $Z_{BL}$ on the production cross-section are shown dashed in dashed lines: purple line ($M_{Z_{BL}} = 4.5$ TeV, $g_{BL} = 0.44$) and blue line ($M_{Z_{BL}} = 4.5$ TeV, $g_{BL} = 0.55$). Other parameters are fixed as mentioned inset of the figure. The shaded region is excluded from ATLAS data on doubly charged scalar mass, $m_{H^{\pm\pm}}$ for $\sqrt{s} = 13$ TeV and luminosity $139 \text{ fb}^{-1}$.

Note here that the above values of the $g_{BL}$ can be obtained using the Eqn.28 keeping the other parameters fixed as mentioned in the inset of the figure. For illustration purpose we considered two moderate values of $g_{BL}$: 0.44 (purple dashed line) and 0.55 (blue dashed line) which are in agreement with the current ATLAS bound $g_{BL} \leq 0.57$ for $M_{Z_{BL}} = 4.5$ TeV. It is noticeable from the graph that the presence of $Z_{BL}$ enhances the production cross-section towards the heavy mass region of doubly charge scalar with moderate value of $g_{BL}$ compared to the case without $U(1)_{B-L}$ augmentation. It is because of the on-shell decay of $Z_{BL}$ to $H^{++}H^{--}$ pair as $M_{Z_{BL}} > 2m_{H^{\pm\pm}}$ and there is a constructive interference between $Z_{BL}$ and the SM gauge bosons. The orange dashed line shows the observed and expected upper limit on $H^{++}H^{--}$ pair production cross-section as a function of doubly charged scalar mass $m_{H^{\pm\pm}}$ at 95% CL which is obtained from the combination of multi-leptons with jets plus missing...
energy search at ATLAS with $\sqrt{s} = 13$ TeV and integrated luminosity, $\mathcal{L} = 139$ fb$^{-1}$ [85]. This upper limit of production cross-section excludes the region of doubly charged triplet mass, $m_{H^{\pm\pm}} \leq 350$ GeV as shown by the shaded region in Fig. 22.

FIG. 23: The model parameters are projected against the production cross-section of doubly charged scalar ($H^{\pm\pm}$) as a function of $B-L$ gauge boson mass $M_{Z_{BL}}$. The blue points are satisfying the Relic+Direct Detection+ Indirect Detection constraints. Red points are consistent with the ATLAS constraint on $g_{BL} - M_{Z_{BL}}$.

In figure 23, we projected the points satisfying all the relevant constraints against the doubly charged scalar ($H^{\pm\pm}$) production cross-section as a function of $B-L$ gauge boson mass $M_{Z_{BL}}$ with $\sqrt{s} = 13$ TeV for a benchmark value of $m_{H^{\pm\pm}} = 404.343$ GeV. The orange dashed line shows the upper limit on the production cross-section from ATLAS [85]. The blue points show the parameters that satisfy all the relevant constraints like correct relic density, direct and indirect search of DM. We can see that in presence of the $B-L$ gauge boson, the production cross-section $\sigma_{pp \rightarrow H^{++}H^{--}}$ can get a distinctive enhancement as compared to the case where production of $H^{\pm\pm}$ happens through SM gauge bosons ($\gamma^*, Z^*$) mediation only which is shown as the dashed black line for easy comparison. But since the gauge coupling, $g_{BL}$ is strongly constrained for a corresponding value of $M_{Z_{BL}}$ from non-observation of new gauge boson at LHC, when we impose this constraint on $g_{BL} - M_{Z_{BL}}$, we see that there is not much enhancement in the production cross-section as compared to the SM. It is clear from the red colored points in fig 23 that even in the presence of additional gauge boson $Z_{BL}$, we
can see only a 5% enhancement (i.e. $\sigma_{pp\rightarrow H^{++}H^{--}} = 4.2$ fb in presence of $Z_{BL}$ as compared to 4 fb predicted by SM.) in the production of the doubly charged scalar of mass around 400 GeV. Thus it is conclusive to say that the contribution from $Z_{BL}$ mediated production is negligible as compared to the SM gauge boson mediation for $m_{H^{\pm\pm}} \sim 350 - 800$ GeV or so.

![FIG. 24: The production cross-section of doubly charged scalar ($H^{\pm\pm}$) as a function of $M_{Z_{BL}}$. The orange points correspond to gauge coupling: $0.001 \leq g_{BL} \leq 0.50$. The blue points are allowed from LEP exclusion bound. Red points are consistent with ATLAS constraint on $g_{BL} - M_{Z_{BL}}$.](image)

However if we consider the doubly charged scalar mass in the TeV scale, requiring a higher $M_{Z_{BL}}$ ($> 2$ TeV) for the resonance enhancement to happen, we can see a perceptible signal at the collider. Thus to demonstrate this fact, we considered the doubly charged scalar mass $m_{H^{\pm\pm}} = 1$ TeV. In figure 24, we show the production cross-section of doubly charged scalar ($H^{\pm\pm}$) as a function of $M_{Z_{BL}}$ considering the gauge coupling within the interval $0.001 \leq g_{BL} \leq 0.50$ with $\sqrt{s} = 14$ TeV shown by the orange points. Though the constraints from the current LHC experiments have already surpassed the LEP II limits on $g_{BL} - M_{Z_{BL}}$, for comparison we show the blue points in the plot which depicts the maximum increase in $\sigma_{pp\rightarrow H^{++}H^{--}}$ when the constraint from LEP on $g_{BL} - M_{Z_{BL}}$ is incorporated into the calculation. However, even after imposing the most stringent constraint from ATLAS on $g_{BL} - M_{Z_{BL}}$, we observe that there is noteworthy enhancement in the production of $H^{\pm\pm}$ as compared to the value predicted by SM. The production cross-section $\sigma_{pp\rightarrow H^{++}H^{--}}$ increases

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by almost 34% (0.063 fb in presence of $Z_{BL}$ as compared to 0.047 fb predicted by SM) for $M_{Z_{BL}}$ around 2.5 TeV and it further gets enhanced by almost 54% (0.074 fb in presence of $Z_{BL}$ as compared to 0.047 fb predicted by SM) if $M_{Z_{BL}}$ is around 4.4 TeV. This feature is evident from the red points in fig 24. This is the crucial evidence of the scenario considered here that can be probed by the near and future colliders and hence the feasibility of this model can be verified.

This also establishes an interesting connection between the dark sector and the generation of neutrino mass via the modified type-II seesaw in a gauged $B-L$ setup that we discussed.

VI. CONCLUSION

In this paper, we have studied a very well motivated gauge extension of the standard model by augmenting the SM gauge group with a $U(1)_{B-L}$ symmetry which happens to be an accidental symmetry of SM to simultaneously address non-zero masses and of light neutrinos as well as a phenomenologically viable dark matter component of the universe. The neutrino mass is explained through an alternative type-II see-saw mechanism by incorporating two scalar triplets in the model, one being super heavy and the other in the TeV scale and thus having the potential to be detected at present and future colliders through its interesting signatures. We minimally extend the fermion particle content of the model by adding three exotic right chiral fermions $\chi_{iR}$ with $B-L$ charges $-4, -4$ and $+5$ in order to cancel the gauge and gravitational anomalies that arise when one gauges the $B-L$ symmetry. The stability of these fermions is owed to the remnant $Z_2$ symmetry after the $U(1)_{B-L}$ breaking which distinguishes the added fermions from the SM as $\chi'_{iR}$s are odd under $Z_2$ while all other particles are even. Thus the dark matter emerges as the lightest Majorana fermion from the mixture of these exotic fermions.

A very interesting and important aspect of the model is the correlation between neutrino mass and dark sector. As generation of neutrino mass is explained through a modified type-II seesaw at TeV scale by introducing two $SU(2)_L$ triplet scalars $\Delta$ and $\xi$. $\Delta$ is super heavy and doesn’t have a coupling with the lepton and hence can not generate the neutrino mass even after acquiring an induced vev after the EW symmetry breaking. Thus the neutrino mass is essentially generated through the $\xi_{LL}$ coupling as given in Eq. 11. In the limit $v_{BL} \to 0$, which essentially means vanishing mixing between $\Delta$ and $\xi$, the neutrino mass also vanishes.
Also Eqs. 3, 22 and 24 implies that the interaction between $\chi_{\tau R}$ and $\chi_{e R}, \chi_{\mu R}$ is established through the scalar $\Phi_{BL}$. In the limit of $\langle \Phi_{BL} \rangle \to 0$, which essentially implies $\sin \theta \to 0$, the DM candidate $\chi_3$ decouples from the heavier dark particles $\chi_1$ and $\chi_2$. In this limit there will be no co-annihilations among the dark sector particles. Thus only if $\langle \Phi_{BL} \rangle \neq 0$, we get non-zero masses of light neutrinos as well as it switches on the co-annihilations of DM and hence enlarges the parameter space satisfying all relevant constraints.

We studied the model parameter space by taking into account all annihilation and co-annihilation channels for DM mass ranging from 1 GeV to 2 TeV. We confronted our results with recent data from PLANCK and XENON-1T to obtain the parameter space satisfying relic density as well as direct detection constraints. The DM being Majorana in nature, it escapes from the gauge boson mediated direct detection constraint. We also checked for the constraints on our model parameters from the indirect search of DM using the recent data from Fermi-LAT and MAGIC which we found to be relatively weaker than other constraints. We also imposed the constraint on $g_{BL} - M_{Z_{BL}}$ from current LHC data to obtain the final parameter space allowed from all constraints including correct relic, direct and indirect detection of DM as well as the constraints from colliders on the $B - L$ gauge boson and the corresponding coupling.

We also studied the detection prospects of the doubly charged scalar triplet which can have novel signatures at the colliders with multi-leptons along with missing energy and jets. We showed how in the presence of the $B - L$ gauge boson $Z_{BL}$, the pair production cross-section of $H^{++}H^{--}$ can get enhanced and also depicted how the dark parameters satisfying all the relevant constraints can affect the production of this doubly charged scalar.

Acknowledgments

SM would like to acknowledge Alexander Belyaev and Alexander Pukhov for useful discussions. PG would like to acknowledge the support from DAE, India for the Regional Centre for Accelerator based Particle Physics (RECAPP), Harish Chandra Research Institute.
Appendix A: Anomaly Cancellation

In any chiral gauge theory the anomaly coefficient is given by \[86\]:

\[
\mathcal{A} = Tr[T_aT_bT_{c}]_R - Tr[T_aT_bT_{c}]_L ,
\]

where \( T \) denotes the generators of the gauge group and \( R, L \) represent the interactions of right and left chiral fermions with the gauge bosons.

Gauging of \( U(1)_{B-L} \) symmetry within the SM lead to the following triangle anomalies:

\[
\mathcal{A}_1[\{U(1)_{B-L}\}]^{3} = 3
\]

\[
\mathcal{A}_2[\{Gravity\}^{2} \times U(1)_{B-L}] = 3 .
\]

The natural choice to make the gauged B – L model anomaly free is by introducing three right handed neutrinos, each of having B – L charge \(-1\) such that they result in \( \mathcal{A}_1[\{U(1)_{B-L}\}] = -3 \) and \( \mathcal{A}_2[\{Gravity\}^{2} \times U(1)_{B-L}] = -3 \) which lead to cancellation of above mentioned gauge anomalies.

However we can have alternative ways of constructing anomaly free versions of \( U(1)_{B-L} \) extension of the SM. In particular, three right chiral fermions with exotic B – L charges \(-4,-4,+5\) can also give rise to vanishing B – L anomalies.

\[
\mathcal{A}_1[\{U(1)_{B-L}\}] = \mathcal{A}_1^{SM}[\{U(1)_{B-L}\}] + \mathcal{A}_1^{New}[\{U(1)_{B-L}\}] = 3 + [(-4)^3 + (-4)^3 + (5)^3] = 0
\]

\[
\mathcal{A}_2[\{Gravity\}^{2} \times U(1)_{B-L}] = \mathcal{A}_2^{SM}[\{Gravity\}^{2} \times U(1)_{B-L}] + \mathcal{A}_2^{New}[\{Gravity\}^{2} \times U(1)_{B-L}]
\]

\[
= 3 + [(-4) + (-4) + (5)] = 0
\]

Appendix B: Decay of Doubly Charged Scalar

The partial decay widths of the doubly charged scalar\((H^{++})\) are given as:

\[
\Gamma(H^{++} \rightarrow l_l^{+}l_l^{+}) = \frac{m_{H^{++}}}{4\pi v^2}(1 \delta_{\alpha\beta})^2 |(M_{\nu})_{\alpha\beta}|^2 (B1)
\]

and

\[
\Gamma(H^{++} \rightarrow W^+W^+) = g^4 v^2 m^3_{H^{++}} \left[ 1 - 4 \left( \frac{m_W}{m_{H^{++}}} \right)^2 \right] \left[ 1 - 4 \left( \frac{m_W}{m_{H^{++}}} \right)^2 + \left( \frac{m_W}{m_{H^{++}}} \right)^4 \right] (B2)
\]
FIG. 25: The contour for $R$ in the plane of $v_\xi$ vs $M_\xi$.

This can be well analyzed by plotting contours of the ratio

$$R = \frac{\Gamma(H^{++} \to W^+W^+)}{\Gamma(H^{++} \to l_\alpha^+l_\beta^+)}$$

in the plane of $m_{H^{++}}$ vs $v_\xi$ as shown in the Fig. 25. It can be easily inferred from Fig 25 that if $v_\xi$ is a few hundred MeV or more, then $H^{++}$ dominantly decays to $W^+W^+$.

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