Scalar dark matter, Neutrino mass, Leptogenesis and rare B decays in a U(1)$_{B-L}$ model

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Abstract

We investigate the phenomenology of singlet scalar dark matter in a simple U(1)$_{B-L}$ gauge extension of standard model, made anomaly free with four exotic fermions. The enriched scalar sector and the new gauge boson $Z'$, associated with U(1) gauge extension, connect the dark sector to the visible sector. We compute relic density, consistent with Planck limit and $Z'$ mediated dark matter-nucleon cross section, compatible with PandaX bound. The mass of $Z'$ and the corresponding gauge coupling are constrained from LEP-II and LHC dilepton searches. We also briefly scrutinize the neutrino mass generation through radiative mechanism. Additionally, we restrict the new gauge parameters by using the existing data on branching ratios of rare $B(\tau)$ decay modes. With a slight modification of the same model, we discuss resonant leptogenesis phenomena with TeV scale exotic fermions to produce the observed baryon asymmetry of the Universe.

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I. INTRODUCTION

Standard Model (SM) of particle physics has produced a remarkable success in explaining physics of the fundamental particles below electroweak scale. However, it does not accommodate the explanation for existence of dark matter (DM), observed matter-asymmetry and few anomalies associated with B-sector. The experimental detection of dark matter signal is one of the most awaited event to happen, ever since it was proposed by Fritz Zwicky in early 1930’s [1, 2]. The theoretical proposal of Weakly Interacting Massive Particle (WIMP) has received decent attention in the recent past, where it can produce the correct relic density by freeze-out mechanism. Numerous beyond SM scenarios were realized with WIMP kind of dark matter, and explored immensely in literature [3, 4]. The interaction of WIMP with SM particles opens the scope of its detection prospects, through the production of DM particles in colliders or the direct scattering with the nucleus.

On the other hand, the LHCb as well as Belle and BaBar experiments have reported discrepancy in the angular observables of rare decay modes, induced by the quark level transitions, $b \to s l^+ l^−$ and $b \to c l \bar{\nu}$ over the last few years. These measurements include disagreements at the level of $\sim 3\sigma$ in the decay distribution [5, 6] and $P'_5$ observable of $B \to K^* \mu^+ \mu^−$ [6, 7]. The decay rate of $B_s \to \phi \mu^+ \mu^−$ also show $3\sigma$ discrepancy in the high recoil limit [8, 9]. Additionally, the lepton universality violating ratios, $R_K \equiv \Gamma(B^+ \to K^+ \mu^+ \mu^−)/\Gamma(B^+ \to K^+ e^+ e^−)$ along with $R_{K^*} \equiv \Gamma(B^0 \to K^{*0} \mu^+ \mu^−)/\Gamma(B^0 \to K^{*0} e^+ e^−)$ deviates at $\sim 2.5\sigma$ level [10-13] and the $R_{D^{(*)}} \equiv \Gamma(B \to D^{(*)} \tau \bar{\nu})/\Gamma(B \to D^{(*)} l \bar{\nu})$ ($R_{J/\psi} \equiv \Gamma(B \to J/\psi \tau \bar{\nu})/\Gamma(B \to J/\psi l \bar{\nu})$), where $l = e, \mu$ ratios disagrees with the SM at the level of $\sim 3.08\sigma$ (1.7\sigma) [16-19].

Moreover, Baryon Asymmetry of the Universe (BAU) being a mysterious problem, needs to be investigated in detail in the growing astro-particle experiments. With the necessity of Sakharov’s conditions for baryogenesis, leptogenesis is the most preferable way to fit with the current cosmological observation of the baryon asymmetry, $\Omega_B h^2 = 0.0223 \pm 0.0002$ [20], which corresponds to $Y_B \equiv \eta_B/s \approx 0.86 \times 10^{−10}$. Generation of lepton asymmetry comes from the CP violating out of equilibrium decay of heavy particle, which later converts to the baryon asymmetry through sphaleron transitions. In general, lepton asymmetry produced by the decay of right handed neutrinos has been widely studied in the literature [21-23]. But with one flavor approximation, the lower limit on right-handed neutrino mass ($\gtrsim 10^9$ GeV) corresponds to the Ibarra bound, which is quite impossible to have any experimental signature in coming decades. Of the many attempts made in literature, resonant leptogenesis is the simplest and well known way to generate a successful asymmetry, by bringing down the mass scale, also compatible with the current neutrino oscillation data [24].

To resolve the above issues in a common theory, the SM needs to be extended with additional symmetries or particles. Among many beyond SM frameworks, $U(1)$ extensions stand in the front row, when it comes to simplicity. They are fruitful in phenomenological perspective, with minimal particle and parameter content. These kind of models also provide new scalar and gauge bosonic type mediator particles, that communicate visible sector to the additional particle spectrum. This article includes a minimal $U(1)_{B-L}$ gauge extension of the SM in two different scenario to address these experimental conflicts in a model dependent
framework. To avoid triangle gauge anomalies, these extensions require neutral fermions with appropriate $B - L$ charges. A solution of adding three right-handed fermions, each with $B - L$ charge as $-1$ [25], has been explored in [26]. The inclusion of three exotic fermions, charged $-4, -4$ and $+5$ was proposed [25] and explored in [27–33]. In the present context, we go for the choice of adding four exotic fermions [34] with fractional $B - L$ charges $-1/3, -2/3, -2/3$ and $4/3$. We explore scalar singlet DM, whose $B - L$ charge is such that it ensures the stability. We discuss two different scenario of the current framework. In the first scenario we add a scalar triplet and study the DM phenomenology, radiative neutrino mass. However, this seems to be insufficient to explain leptogenesis. So, in the second context we replace the scalar triplet with a scalar singlet to explain the resonant leptogenesis after the breaking of $B - L$ symmetry. We also scrutinize the rare $B$ decay modes at one-loop level (via $Z'$ boson) in the present framework, which is independent of both the scenarios.

The plan of the paper is as follows. In section II, we describe the model along with the relevant interaction Lagrangian. We discuss the symmetry breaking pattern, particle mass spectrum and radiative neutrino mass in section III. Section IV gives a detail study of dark matter phenomenology in relic density and direct detection perspective and also impose constraints from collider studies. Then in section V, we additionally constrain the new gauge parameters from $B$ and $\tau$ sectors. Resonant leptogenesis with quasi degenerate right-handed fermions and the solutions to Boltzmann equations are discussed in section VI. Summarization of the model is provided in section VII.

II. THE MODEL FRAMEWORK

We study scalar dark matter in an uncomplicated $U(1)_{B-L}$ gauge extension of SM. Apart from the existing SM particle content, four exotic fermions ($N_i$’s, where $i = 1, 2, 3, 4$), assigned with $B - L$ charges $-1/3, -2/3, -2/3$ and $4/3$ are added to avoid the unwanted triangle gauge anomalies. This choice of $B - L$ charges was first proposed by [34] and later explored in various works [35, 36]. We add two scalar singlets $\phi_1, \phi_2$, in the process of breaking $B - L$ gauge symmetry spontaneously and also generate mass terms to all the new particles. An inert scalar singlet $S_{DM}$, qualifies as a dark matter in the present model, whose stability is ensured by the $B - L$ gauge symmetry itself [26, 37]. Finally, we include a scalar triplet $\Delta$, to generate mass splitting in the CP-odd and CP-even components of DM singlet, useful in generating neutrino mass at one-loop. The complete field content along with their corresponding charges under $SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ are provided in Table I.
The Yukawa interaction for the present model is given by

\[
L_{\text{Yuk}} = \sum_{\alpha} Y_{\alpha} \overline{Q}_{L}^{\alpha} \chi_{u}^{\alpha} \left( \partial_{\mu} + \frac{i g}{2} \cdot \tilde{W}_{\mu} \right) \tilde{H} N_{L} \phi_{\text{DM}} + \text{H.c}.
\]  

The relevant terms in the fermion interaction Lagrangian is given by

\[
\begin{align*}
L_{\text{Kin.}}^f &= \overline{Q}_{L}^{\alpha} \gamma^\mu \left( \partial_{\mu} + i g \frac{\tau^a}{2} \cdot \tilde{W}_{\mu} + \frac{1}{6} g' B_{\mu} + \frac{1}{3} g B_{\mu} Z'_{\mu} \right) Q_{L} \\
+ \overline{u}_{R}^{\alpha} \gamma^\mu \left( \partial_{\mu} + \frac{2}{3} g' B_{\mu} + \frac{1}{3} g B_{\mu} Z'_{\mu} \right) u_{R} \\
+ \overline{d}_{R}^{\alpha} \gamma^\mu \left( \partial_{\mu} - \frac{1}{3} g' B_{\mu} + \frac{1}{3} g B_{\mu} Z'_{\mu} \right) d_{R} \\
+ \overline{\ell}_{L}^{\alpha} \gamma^\mu \left( \partial_{\mu} + i g \frac{\tau^a}{2} \cdot \tilde{W}_{\mu} - \frac{1}{2} g' B_{\mu} - \frac{1}{3} g B_{\mu} Z'_{\mu} \right) \ell_{L} \\
+ \overline{e}_{R}^{\alpha} \gamma^\mu \left( \partial_{\mu} - i g' B_{\mu} - i g B_{\mu} Z'_{\mu} \right) e_{R} \\
+ \overline{N}_{1R}^{\alpha} \gamma^\mu \left( \partial_{\mu} - \frac{1}{3} i g B_{\mu} Z'_{\mu} \right) N_{1R} + \overline{N}_{2R}^{\alpha} \gamma^\mu \left( \partial_{\mu} - \frac{2}{3} i g B_{\mu} Z'_{\mu} \right) N_{2R} \\
+ \overline{N}_{3R}^{\alpha} \gamma^\mu \left( \partial_{\mu} - \frac{4}{3} i g B_{\mu} Z'_{\mu} \right) N_{3R} + \overline{N}_{4R}^{\alpha} \gamma^\mu \left( \partial_{\mu} - \frac{4}{3} i g B_{\mu} Z'_{\mu} \right) N_{4R}.
\end{align*}
\]  

The Yukawa interaction for the present model is given by

\[
\begin{align*}
\mathcal{L}_{\text{Yuk}} &= Y_{u} \overline{Q}_{L} \tilde{H} u_{R} + Y_{d} \overline{Q}_{L} H d_{R} + Y_{e} \overline{\ell}_{L} H e_{R} + H.c + \left( \sum_{\beta=2,3,\alpha} \frac{Y_{\alpha\beta}}{\Lambda} \overline{\ell}_{\alpha\beta} \tilde{H} N_{\beta R} \phi_{\text{DM}} + H.c \right) \\
+ \sum_{\alpha} \left( \frac{Y_{\alpha 4}}{\Lambda} \overline{\ell}_{\alpha 4} \tilde{H} N_{4R} \phi_{\text{DM}} + H.c \right) + \sum_{\beta=2,3} \left( h_{\beta 1} \phi_{1} N_{\beta R} N_{1R} + h_{\beta 2} \phi_{2} N_{\beta R} N_{4R} \right),
\end{align*}
\]  

with \( \tilde{H} = i \sigma_{2} H^{\ast} \). The interaction Lagrangian for the scalar sector is as follows

\[
\mathcal{L}_{\text{scalar}} = \left( D_{\mu} H \right)^{\dagger} \left( D^{\mu} H \right) + \left( D_{\mu} \phi_{\text{DM}} \right)^{\dagger} \left( D^{\mu} \phi_{\text{DM}} \right) + \left( D_{\mu} \phi_{1} \right)^{\dagger} \left( D^{\mu} \phi_{1} \right) \\
+ \left( D_{\mu} \phi_{2} \right)^{\dagger} \left( D^{\mu} \phi_{2} \right) + \text{Tr} \left[ \left( D_{\mu} \Delta \right)^{\dagger} \left( D^{\mu} \Delta \right) \right] + V \left( H, \phi_{\text{DM}}, \phi_{1}, \phi_{2}, \Delta \right),
\]  

TABLE I: Particle spectrum and their charges of the proposed U(1)$_{B-L}$ model.
where the covariant derivatives are

\[ D_\mu H = \partial_\mu H + i g \tilde{W}_\mu L \cdot \frac{\tau}{2} H + i g' B_\mu H, \]

\[ D_\mu \phi_{DM} = \partial_\mu \phi_{DM} - ig_{BL} \frac{1}{3} Z'_\mu \phi_{DM}, \]

\[ D_\mu \phi_1 = \partial_\mu \phi_1 + ig_{BL} Z'_\mu \phi_1, \]

\[ D_\mu \phi_2 = \partial_\mu \phi_2 + 2ig_{BL} Z'_\mu \phi_2, \]

\[ D_\mu \Delta = \partial_\mu \Delta + i g \tilde{W}_{\mu a} \cdot \frac{\tau}{2} \Delta + ig' B_{\mu a} \Delta - \frac{2}{3} ig_{BL} Z'_\mu \Delta, \] (4)

where, \( \Delta \) can be written in the isospin basis as

\[ \Delta = \begin{pmatrix} \Delta_+ \Delta_0 \Delta_- \end{pmatrix}. \]

And the scalar potential takes the form

\[ V(H, \phi_1, \phi_2, \phi_{DM}) = \mu_1^2 H^\dagger H + \mu_2^2 \phi_1^\dagger \phi_1 + \mu_3 \phi_1^\dagger \phi_1 + \mu_4 \phi_2^\dagger \phi_2 + \mu_5 \phi_1^\dagger \phi_2 + \mu_6 \phi_2^\dagger \phi_1 + \mu_7 \phi_1^\dagger \phi_1 + \mu_8 \phi_2^\dagger \phi_2 + \mu_9 \phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2, \]

\[ + \lambda_1 \phi_1^\dagger \phi_1 + \lambda_2 \phi_2^\dagger \phi_2 + \lambda_3 \phi_1^\dagger \phi_2 + \lambda_4 \phi_2^\dagger \phi_1 + \lambda_5 \phi_1^\dagger \phi_1 + \lambda_6 \phi_2^\dagger \phi_2 + \lambda_7 \phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2, \]

\[ + \lambda_{semi} ((\phi_{DM})^3 \phi_1 + (\phi_{DM})^3 \phi_2), \] (5)

\[ V_\Delta = \mu_1^2 \text{Tr}(\Delta^\dagger \Delta) + \lambda_1 \text{Tr}(\Delta^\dagger \Delta)^2 + \lambda_2 \text{Tr}(\Delta^\dagger \Delta)(H^\dagger H) + \lambda_3 \text{Tr}(\Delta^\dagger \Delta)(\phi_1^\dagger \phi_1) + \lambda_4 \text{Tr}(\Delta^\dagger \Delta)(\phi_2^\dagger \phi_2) + \lambda_5 \text{Tr}(\Delta^\dagger \Delta)(\phi_{DM}^\dagger \phi_{DM}) + \frac{\lambda_{6p}}{\lambda} \phi_{DM}^2 H^\dagger H. \] (6)

Full potential of this model is given by

\[ V = V(H, \phi_1, \phi_2, \phi_{DM}) + V_\Delta. \] (7)

Here, \( \phi_{DM} = \frac{\phi_{SM} + \phi_{DM}}{\sqrt{2}} \) is the DM singlet in the present model. The stability of the potential is assured by the copositive criteria, given as [26]

\[ \lambda_1, \lambda_2, \lambda_{DM}, \lambda_T \geq 0, \quad \lambda_{H1} + \sqrt{\lambda_H \lambda_1} \geq 0, \]

\[ \lambda_{H2} + \sqrt{\lambda_H \lambda_2} \geq 0, \quad \lambda_{HD} + \sqrt{\lambda_H \lambda_D} \geq 0, \]

\[ \lambda_{HT} + \sqrt{\lambda_H \lambda_T} \geq 0, \quad \lambda_{12} + \sqrt{\lambda_1 \lambda_2} \geq 0, \]

\[ \lambda_{D1} + \sqrt{\lambda_D \lambda_1} \geq 0, \quad \lambda_{D2} + \sqrt{\lambda_D \lambda_2} \geq 0, \]

\[ \lambda_{T1} + \sqrt{\lambda_T \lambda_1} \geq 0, \quad \lambda_{T2} + \sqrt{\lambda_T \lambda_2} \geq 0, \quad \lambda_{DT} + \sqrt{\lambda_D \lambda_T} \geq 0. \] (8)

III. SPONTANEOUS SYMMETRY BREAKING AND MIXING

Spontaneous symmetry breaking of SU(2)_L × U(1)_Y × U(1)_{B−L} to SU(2)_L × U(1)_Y is realized by assigning non-zero VEV to the scalar singlets \( \phi_1 \) and \( \phi_2 \). Later, the SM gauge group gets spontaneously to low energy theory by the SM Higgs doublet \( H \). The scalar
sector can be written in terms of CP-even and CP-odd components as
\[ H^0 = \frac{1}{\sqrt{2}}(v + h) + \frac{i}{\sqrt{2}} A^0, \]
\[ \phi_1 = \frac{1}{\sqrt{2}}(v_1 + h_1) + \frac{i}{\sqrt{2}} A_1, \]
\[ \phi_2 = \frac{1}{\sqrt{2}}(v_2 + h_2) + \frac{i}{\sqrt{2}} A_2, \]
where, \( \langle H \rangle = (0, v/\sqrt{2})^T, \langle \phi_1 \rangle = v_1/\sqrt{2}, \langle \phi_2 \rangle = v_2/\sqrt{2} \) and the VEV of scalar triplet is given as \( \langle \Delta \rangle = \frac{v_T}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ 1 & 0 \end{pmatrix} \).

### A. Mixing in scalar sector

The minimisation conditions of the scalar potential correspond to
\[ \mu^2_H = -\frac{1}{2}\left[2\lambda_H v^2 + \lambda_{H1} v_1^2 + \lambda_{H2} v_2^2 + \lambda_{HT} v_T^2 \right], \]
\[ \mu^2_1 = -\frac{1}{2}\left[2\lambda_1 v_1^2 + \lambda_{12} v_2^2 + \lambda_{H2} v_2^2 + \lambda_{T2} v_T^2 v_2 + \sqrt{2}\mu_{12} v_2 \right], \]
\[ \mu^2_2 = -\frac{1}{2}\frac{v_2}{2v^2}\left[2\lambda_2 v_2^3 + \lambda_{12} v_1^2 v_2 + \lambda_{H1} v_1^2 + \lambda_{T1} v_T^2 + \sqrt{2}\mu_{12} v_2 \right], \]
\[ \mu^2_T = -\frac{1}{2}\left[2\lambda_{T1} v_1^2 + \lambda_{T2} v_T^2 + 2\lambda_{HT} v^2 + 2\lambda_{T} v_T^2 \right]. \] (9)

After the electroweak symmetry breaking, the CP-even scalar mass matrix takes the form
\[ M_E^2 = \begin{pmatrix} 2\lambda_H v^2 & \lambda_{H1} v_1 v \\ \lambda_{H1} v_1 v & 2\lambda_1 v_1^2 \end{pmatrix} \begin{pmatrix} \lambda_{H2} v_2 v \\ \lambda_{H2} v_2 v + v_1 (\lambda_{12} v_2 + \sqrt{2}\mu_{12}) \end{pmatrix}. \] (10)

We assume that the third neutral Higgs (\( \phi_2 \)) and the neutral component of scalar triplet to be very heavy and decoupled. We consider mixing between only \( H \) and \( \phi_1 \) (in the small mixing limit) as
\[ M_E^2 = \begin{pmatrix} 2\lambda_H v^2 & \lambda_{H1} v_1 v \\ \lambda_{H1} v_1 v & 2\lambda_1 v_1^2 \end{pmatrix}. \] (11)

This 2 \( \times \) 2 matrix can be diagonalized by the following usual rotation matrix as \( U_E^\dagger M_E U_E = \text{diag} [M^2_{H1}, M^2_{H2}] \), where
\[ U_E = \begin{pmatrix} \cos \beta & \sin \beta \\ -\sin \beta & \cos \beta \end{pmatrix}. \] (12)

The involved scalar couplings can be written as
\[ \lambda_H = \frac{1}{2v^2} \left( \sin^2 \beta M^2_{H1} + \cos^2 \beta M^2_{H2} \right), \]
\[ \lambda_1 = \frac{1}{2v_1^2} \left( \cos^2 \beta M^2_{H1} + \sin^2 \beta M^2_{H2} \right), \]
\[ \lambda_{H1} = \frac{1}{v v_1} \left( \cos \beta \sin \beta (M^2_{H2} - M^2_{H1}) \right). \]
The mass eigenstate $H_1$ is considered to be observed Higgs at LHC ($M_{H_1} = 125\text{ GeV}$) and $H_2$ (with mass $M_{H_2}$) is the heavy physical scalar in the present model. This mixing is taken to be minimal ($\beta < 0.1$), such that it does not violate the LHC bounds on the observed Higgs. The mass matrix of the CP-odd components takes the form

$$M_O = \begin{pmatrix} -2\sqrt{2}\mu_{12}v_2 & \sqrt{2}\mu_{12}v_1 \\ \sqrt{2}\mu_{12}v_1 & -\mu_{12}v_2 \end{pmatrix}. \quad (13)$$

This matrix gives one zero mass eigenvalue and the corresponding eigenstate gets converted into the longitudinal mode of the new $U(1)$ gauge boson $Z'$. The other orthogonal combination $A_{NG}$, obtains the mass

$$M_{2_{NG}} = -\mu_{12}(v_1^2 + 4v_2^2)/\sqrt{2}v_2.$$ 

The mass of $Z'$ is given by

$$M_{Z'} = g_{BL}\sqrt{v_1^2 + 4v_2^2}.$$ 

### B. Comments on Neutrino Masses

![FIG. 1: Generation of radiative neutrino mass in this model.](image)

The one loop neutrino mass arising from the diagram in Fig. 1 can be estimated as

$$(m_\nu)_{ij} = \frac{Y_{ik}Y_{jk}M_kv^2}{16\pi^2\Lambda^2} \left( \frac{M_S^2}{M_S^2 - M_k^2} \ln \frac{M_S^2}{M_k^2} - \frac{M_A^2}{M_A^2 - M_k^2} \ln \frac{M_A^2}{M_k^2} \right). \quad (14)$$

Here, $M_S$ ($M_A$) is the mass of CP even (odd) part of the singlet DM and $M_k$ is the mass of singlet fermions $N_k$ in the loop. For $M_S^2 + M_A^2 \approx M_k^2$, the above expression can be simply written as

$$(m_\nu)_{ij} \approx \frac{(M_S^2 - M_A^2)v^2Y_{ik}Y_{jk}}{32\pi^2\Lambda^2 M_k}. \quad (15)$$

The mass splitting between the CP odd and CP even component is given by

$$M_A^2 - M_S^2 = \frac{\lambda_{sp}}{\Lambda} v^2 v_T. \quad (16)$$

The flavor structure of $Y$ and right-handed neutrino mass $M_R$ in the flavor basis can be written from the Yukawa Lagrangian in Eq.(2) as

$$Y = \begin{pmatrix} 0 & Y_{12} & Y_{13} & Y_{14} \\ 0 & Y_{22} & Y_{23} & Y_{24} \\ 0 & Y_{32} & Y_{33} & Y_{34} \end{pmatrix}, \quad M_R = \begin{pmatrix} 0 & h_{21}v_1 & h_{31}v_1 & 0 \\ h_{21}v_1 & 0 & 0 & h_{24}v_2 \\ h_{31}v_1 & 0 & 0 & h_{34}v_2 \\ 0 & h_{24}v_2 & h_{34}v_2 & 0 \end{pmatrix}. \quad (17)$$
Assuming \( h_{21} = h_{34} = af, \ h_{31} = h_{24} = bf \) and \( v_1 = v_2 = v' \), we will have double fold degenerate right handed neutrinos with masses \( (a - b)fv', \ -(a - b)fv', \ (a + b)fv' \) and \(-(a + b)fv'.\) The eigenvector matrix that diagonalizes \( M_R \), i.e \( M_k = U_N M_R U_N^T \), is given by
\[
U_N = \begin{pmatrix}
\frac{1}{2} & -\frac{1}{2} & -\frac{1}{2} & -\frac{1}{2} \\
-\frac{1}{2} & \frac{1}{2} & \frac{1}{2} & -\frac{1}{2} \\
-\frac{1}{2} & \frac{1}{2} & -\frac{1}{2} & \frac{1}{2} \\
\frac{1}{2} & \frac{1}{2} & \frac{1}{2} & \frac{1}{2}
\end{pmatrix}.
\]
Therefore the mass eigenstates of the singlet fermions are given as follows
\[
N_{D_1} = \frac{1}{2} (N_1 - N_2 - N_3 + N_4),
\]
\[
N_{D_2} = -\frac{1}{2} (N_1 - N_2 + N_3 + N_4),
\]
\[
N_{D_3} = -\frac{1}{2} (N_1 + N_2 - N_3 + N_4),
\]
\[
N_{D_4} = \frac{1}{2} (N_1 + N_2 + N_3 + N_4).
\]

The sample benchmark points (BP) are given in the Table II below to achieve a neutrino mass as per the observed cosmological bound.

| Parameters | \( M_k \, [\text{GeV}] \) | \( Y_{ij} \) | \( \lambda_{sp} \) | \( m_{\nu} \, [\text{eV}] \) |
|------------|----------------|-------------|----------------|----------------|
| BP1        | \( 10^3 \)    | 0.1         | \( \approx 0.2 \) | 0.02           |
| BP2        | \( 10^3 \)    | 0.05        | \( \approx 0.33 \) | 0.01           |

TABLE II: Benchmark points for radiatively generated neutrino masses.

IV. PHENOMENOLOGY OF SINGLET SCALAR DARK MATTER

A. Relic density

The relic abundance of DM, can be obtained by solving the Boltzmann equation
\[
\frac{dn_{DM}}{dt} + 3Hn_{DM} = -\langle \sigma v \rangle (n_{DM}^2 - (n_{DM}^{\text{eq}})^2),
\]
where, \( n_{DM}^{\text{eq}} \) is the equilibrium number density of DM, \( H \) denotes the Hubble expansion rate of the Universe. \( \langle \sigma v \rangle \) is the thermally averaged annihilation cross section of DM and can be written in terms of partial wave expansion as \( \langle \sigma v \rangle = a + bv^2 \). Numerical solution of the above Boltzmann equation gives \cite{39,40}
\[
\Omega_{DM} h^2 \approx \frac{1.04 \times 10^9 x_F}{M_{Pl} \sqrt{g_*(a + 3b/x_F)}},
\]
where \( x_F = M_{DM}/T_F \), \( T_F \) is the freeze-out temperature, \( M_{DM} \) is the mass of dark matter, \( g_* \) is the total relativistic degrees of freedom at the time of freeze-out and and \( M_{Pl} \approx 1.22 \times 10^{19} \) GeV is the Planck mass. DM with electroweak scale mass and couplings freeze out at
temperatures approximately in the range $x_F \approx 20 - 30$. Where, $x_F$ can be calculated from the relation

$$x_F = \ln \frac{0.038g_{\text{Pl}}M_{\text{DM}}\langle \sigma v \rangle}{g_1^{1/2}x_F^{1/2}},$$

which can be derived from the condition $\Gamma_H = 1$. Here, $\Gamma(= n_{\text{DM}}\langle \sigma v \rangle)$ is the DM interaction rate and the Hubble expansion $H \approx g_1^{1/2}T^2 M_{\text{Pl}}$. The analytically approximated expression for DM relic abundance [41] is given by

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

(23)

The thermal averaged annihilation cross section $\langle \sigma v \rangle$ is given by [42]

$$\langle \sigma v \rangle = \frac{1}{8M_{\text{DM}}^4TK_3^2(M_{\text{DM}}/T)} \int_{4M_{\text{DM}}^2}^{\infty} \sigma(s - 4M_{\text{DM}}^2)\sqrt{s}K_1(\sqrt{s}/T)ds,$$

where, $K_i$’s are modified Bessel functions of order $i$ and $T$ stands for the temperature.

If there exists some additional particles having mass difference close to that of DM, then they can be thermally accessible during the epoch of DM freeze out. This can give rise to additional channels through which DM can co-annihilate with such additional particles and produce SM particles in the final states. This type of co-annihilation effects on dark matter relic abundance were studied by several authors in [43–45].

This model accommodates a singlet scalar DM, which has both gauge and Higgs portal annihilation channels which are provided in Fig. [2]. The mass difference between the CP odd and CP even component of the singlet scalar is generated by the Higgs triplet. Thus the co-annihilation channels also contribute to the DM relic density. We analyze the contribution of different annihilation and co-annihilation channels within different DM mass regime. We consider $M_{Z'} = 1.4 \text{ TeV}$ and $M_{H_2} = 1 \text{ TeV}$ and show the behavior of relic density as a function of DM mass. In the left panel of Fig. [3] we fix $\lambda_{\text{semi}} = 0.3$ and vary the gauge coupling $g_{BL} = 0.1, 0.3, 0.6$, we found that for DM mass below 75 GeV, the annihilation to fermions maximally contribute and after this mass regime the dominant contribution comes from the annihilation to SM gauge bosons and SM Higgs. After $M_{\text{DM}} = 980 \text{ GeV}$, the Higgs
FIG. 3: Left panel represents the variation relic density with DM mass for different values of gauge coupling $g_{BL}$ by fixing $\lambda_{semi}$ parameter and the horizontal dashed lines represent Planck 3σ limit. Right panel shows the relic density as a function of DM mass showing the impact of semi-annihilation.

FIG. 4: t-channel scattering of DM with nucleus.

mediated annihilation channels contribute maximally and then the annihilation of DM to $Z'$ pairs adds to the relic density for $M_{DM} > 1.4$ TeV. Because of s-channel annihilation, the resonances (dips) are observed near $M_{DM} = \frac{1}{2} M_{H_1,H_2,Z'}$. Due to non-zero $\lambda_{semi}$, the semi-annihilation channels of DM also contribute to the relic density. In the right panel of Fig. 3, we interpret the behavior of relic density for small gauge coupling $g_{BL}$ by varying $\lambda_{semi}$. In this case, the $Z'$ mediated processes are suppressed and the resonance near $M_{DM} = \frac{M_{Z'}}{2}$ diminishes. The scalar mediated annihilation and semi-annihilation processes significantly contribute to the relic density.

B. Direct searches

Now we look for the constraints on the model parameters due to direct detection limits. The effective Lagrangian for $Z'$-mediated t-channel process shown in Fig. 4 is given as

$$\mathcal{L}^{V}_{\text{eff}} \supset - \frac{n_{DM} g_{BL}^2}{3 M_{Z'}^2} (S \partial^\mu A - A \partial^\mu S) \bar{u} \gamma_\mu u - \frac{n_{DM} g_{BL}^2}{3 M_{Z'}^2} (S \partial^\mu A - A \partial^\mu S) \bar{d} \gamma_\mu d.$$ (25)

The corresponding spin independent (SI) WIMP-nucleon cross section turns out to be

$$\sigma_{Z'} = \frac{\mu^2}{\pi} \frac{n_{DM} g_{BL}^4}{M_{Z'}^4}. \quad (26)$$
FIG. 5: Colored lines represent the dilepton signal cross section as a function of $M_{Z'}$ for different values of $g_{BL}$ with the black dashed line points to ATLAS bound \[46\].

Here, $\mu$ denotes the reduced mass of DM-nucleon system. The t-channel scalar exchange i.e via $H_1, H_2$, can also give a SI contribution, but this is not relevant for the purpose of our study.

C. Collider constraints

ATLAS and CMS experiments are searching for new heavy resonances in both dilepton and dijet signals. It is found in the recent past, that these two experiments provide lower limit on $Z'$ boson with dilepton signature, resulting a stronger bound than dijets due to relatively fewer background events. The investigation for $Z'$, through dilepton signals from ATLAS experiment \[46\] concluded with stringent limit on the ratio of $Z'$ mass ($M_{Z'}$) and the gauge coupling ($g_{BL}$). We use CalcHEP \[47, 48\] to calculate the production cross section of $Z'$ to dilepton ($e^+e^-, \mu^+\mu^-$) in final states. The variation of $Z'$ production cross section times the branching of dilepton as a function of $M_{Z'}$ is shown in Fig. 5. From this plot we can interpret that, for $g_{BL} = 0.01$, $M_{Z'} < 0.5$ TeV regime is excluded by the ATLAS bound. Similarly for $g_{BL} = 0.03$, the mass regime for $M_{Z'} < 1.4$ TeV is not allowed. We found with a little larger values of $g_{BL} = 0.1, 0.3$, the allowed mass regime for $M_{Z'}$ should be greater than 2.7 TeV and 3.7 TeV, respectively. Furthermore, there is also a lower limit on the ratio $\frac{M_{Z'}}{g_{BL}}$ from LEP-II \[49\], i.e., 6.9 TeV.

For the parameter scan, we vary the DM mass $M_{DM}$ from 50 GeV to 2 TeV, $M_{Z'}$ in the range 0.5 to 4 TeV and the gauge coupling $g_{BL}$ between 0 to 1. The left panel Fig. 6 shows the $M_{Z'} - g_{BL}$ parameter space, consistent with 3$\sigma$ range of Planck limit on relic density. LEP-II and ATLAS exclusion bounds are denoted with magenta and orange dashed lines respectively. Here, the green data points violate the stringent upper limit on WIMP-nucleon SI cross section set by PandaX-II \[50\] (visible from right panel). Therefore, the viable region of gauge parameters that survives all the experimental limits is the blue data points below ATLAS exclusion limit in the left panel.
FIG. 6: Left panel depicts the $M_{Z'} - g_{BL}$ parameter space consistent with $3\sigma$ region of Planck relic density limit. Dashed lines represent the ATLAS \cite{46} and LEP-II \cite{49} limits. Right panels shows the SI WIMP-nucleon cross section for the parameter space shown in the left panel. Dashed lines represent the recent bounds dictated by PandaX-II \cite{50}, XENON1T \cite{51} and LUX \cite{52}.

V. CONSTRAINTS ON NEW GAUGE PARAMETERS FROM QUARK AND LEPTON SECTORS

Since the $Z'q_iq_j$ interaction term is not allowed in the proposed model, the leptonic/semileptonic $B, D, K$ modes involving the quark level transitions $q_i \rightarrow q_jll(\nu_l\bar{\nu}_l)$ ($q_i = b, c, s$, $q_j = u, d, s$ and $l$ is any charged lepton) can only occur at one loop level via $Z'$ boson as shown in Fig. 7. We mainly focus on the existing data on the branching ratios of $B$ meson channels to constrain the $M_{Z'} - g_{BL}$ plane. The Lepton Flavor Violating (LFV) decay modes of $B$ meson and $\tau(\mu)$ lepton are not allowed due to the absence of $Z'\ell_i\ell_j$ coupling, thus we use the branching ratio of only possible $\tau \rightarrow \mu\nu_\tau\bar{\nu}_\mu$ process for this purpose. In SM, the explicit form of the effective Hamiltonian which is responsible for leptonic/semileptonic $b \rightarrow q(= s, d)ll$ transitions is given by \cite{53,54}

$$H_{\text{eff}} = -\frac{4G_F}{\sqrt{2}}V_{tb}V_{tq}^* \left( \sum_{i=1,\cdots,10,S,P} C_i O_i + \sum_{i=7,\cdots,10,S,P} C'_i O'_i \right), \quad (27)$$

FIG. 7: One loop penguin diagram of $q_i \rightarrow q_jll(\nu_l\bar{\nu}_l)$ mediated by $Z'$ gauge boson.
where $G_F$ is the Fermi constant, $V_{qq'}$ is the product of CKM matrix elements. Here $O_i$'s are the effective operators, defined as

\begin{align*}
O_7^{(i)} & = \frac{e}{16\pi^2} \left( \bar{q}\sigma_{\mu\nu}(m_q P_L + m_b P_R)L_b \right) F^{\mu\nu}, \\
O_9^{(i)} & = \frac{\alpha_{\text{em}}}{4\pi} (\bar{q}\gamma^\mu P_L L_b)(\bar{l}\gamma_{\mu\nu}), \\
O_{10}^{(i)} & = \frac{\alpha_{\text{em}}}{4\pi} (\bar{q}\gamma^\mu P_L b)(\bar{l}\gamma_{\mu\nu}),
\end{align*}

(28)

with $\alpha_{\text{em}}$ is the fine structure constant, $P_{L,R} = (1 \mp \gamma_5)/2$ are the projection operators and $C_i^{(i)}$'s are the corresponding Wilson coefficients. The values of Wilson coefficients in the SM are taken from $[54, 57–59]$. The primed operators are absent in the SM, however the respective coefficients may be non-zero in the presence of $Z'$ boson arising due to the $U(1)_{B–L}$ gauge extension. Using the interaction terms of SM fermions with $Z'$ from Eq.(1), the effective Hamiltonian for $b \to qll$ processes is given by

\begin{equation}
\mathcal{H}^Z' = \frac{G_F}{12\sqrt{2}\pi^2} V_{tb} V_{tq}^* F \left( \frac{m_t^2}{M_Z^2} \right) \frac{g_{B_L}^2}{M_W^2} (\bar{q}\gamma^\mu P_L L_b)(\bar{l}\gamma_{\mu\nu}),
\end{equation}

(29)

where $F \left( \frac{m_t^2}{M_W^2} \right)$ is the loop function that is order one ($F \left( \frac{m_t^2}{M_W^2} \right) \approx 1$) by using $m_t$ and $M_W$ from PDG $[60]$. Now comparing Eq.(29) with (27), we obtain an additional Wilson coefficient contribution to the SM as

\begin{equation}
C_{9}^{NP} = -\frac{1}{12\pi\alpha_{\text{em}}} \frac{g_{B_L}^2}{M_Z^2}.
\end{equation}

(30)

The effective Hamiltonian for rare decay processes mediated by $b \to q(= d, s)\nu\bar{\nu}$ transitions are given by $[61]$

\begin{equation}
\mathcal{H}_\nu^{\text{eff}} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{tq}^* (C_L^{\nu} O_L^\nu + C_R^{\nu} O_R^\nu) + \text{h.c.},
\end{equation}

(31)

where the $O_{L(R)}^{\nu}$ effective operators are defined as

\begin{equation}
O_L^{\nu} = \frac{\alpha_{\text{em}}}{4\pi} (\bar{q}\gamma_{\mu} P_L b)(\bar{l}\gamma^\mu (1 - \gamma_5) \nu), \quad O_R^{\nu} = \frac{\alpha_{\text{em}}}{4\pi} (\bar{q}\gamma_{\mu} P_R b)(\bar{l}\gamma^\mu (1 - \gamma_5) \nu).
\end{equation}

(32)

Here the Wilson coefficient $C_L^{\nu} = -X \left( \frac{m_t^2}{M_W^2} \right) / \sin^2 \theta_W$ is calculated by using the loop function $X \left( \frac{m_t^2}{M_W^2} \right) [58] [59]$ and $C_R^{\nu}$ is negligible in the SM. The effective Hamiltonian in the presence of $Z'$ is

\begin{equation}
\mathcal{H}_\nu^{Z'} = \frac{G_F}{24\sqrt{2}\pi^2} V_{tb} V_{tq}^* F \left( \frac{m_t^2}{M_W^2} \right) \frac{g_{B_L}^2}{M_{Z'}^2} (\bar{q}\gamma_{\mu} P_L b)(\bar{l}\gamma^\mu (1 - \gamma_5) \nu),
\end{equation}

(33)

which in comparison with Eq.(31) provides new contribution to $C_L$ Wilson coefficient as

\begin{equation}
C_L^{\nu NP} = -\frac{1}{24\pi\alpha_{\text{em}}} \frac{g_{B_L}^2}{M_{Z'}^2}.
\end{equation}

(34)

After collecting an idea on new Wilson coefficient contribution, we now proceed to constrain the new gauge parameters from the flavor observables, to be presented in the subsequent subsections.
A. $B \to (\pi, K)ll$

The branching ratio of $B \to Kl\ell$ process with respect to $q^2$ is given by [12]

$$
\frac{dBr}{dq^2} = \tau_B \frac{G_F^2 \alpha_{em}^2 |V_{ub}V_{ts}^*|^2}{2^{\delta+3} M_B^4} \sqrt{\lambda(M_B^2, M_K^2, q^2)\beta_l f_+^2(q^2)(a_l(q^2) + \frac{c_l(q^2)}{3})},
$$

(35)

where,

$$
a_l(q^2) = q^2 |F_P|^2 + \frac{\lambda(M_B^2, M_K^2, q^2)}{4}(|F_A|^2 + |F_V|^2) + 2m_l(M_B^2 - M_K^2 + q^2)\text{Re}(F_P F_A^*) + 4m_l^2 M_B^2 |F_A|^2,
$$

$$
c_l(q^2) = -\frac{\lambda(M_B^2, M_K^2, q^2)}{4}\beta_l^2 (|F_A|^2 + |F_V|^2),
$$

(36)

with

$$
F_V = \frac{2m_b}{M_B} C_{\text{eff}}^{\text{eff}} + C_9^{\text{eff}} + C_9^{\text{NP}}, \quad F_A = C_{10},
$$

$$
F_P = m_l C_{10} \left[ \frac{M_B^2 - M_K^2}{q^2} \left( \frac{f_0(q^2)}{f_+(q^2)} - 1 \right) - 1 \right],
$$

(37)

and

$$
\lambda(a, b, c) = a^2 + b^2 + c^2 - 2(ab + bc + ca), \quad \beta_l = \sqrt{1 - 4m_l^2/q^2}.
$$

(38)

The $B \to \pi$ processes follow the same expression with proper replacement of particle mass, lifetime, CKM matrix elements and Wilson coefficients. To compute the branching ratios of $B^{+}(0) \to (\pi^{+}(0), K^{+}(0))l\ell$ in the SM, all the required input parameters are taken from [60]. The form factors for $B \to (\pi, K)$ in the light cone sum rule approach are considered from [62, 63].

B. $B \to (\pi, K)\nu_\ell\bar{\nu}_\ell$

The differential branching ratio of $B \to P\nu_\ell\bar{\nu}_\ell$ process, where, $P = \pi, K$ are pseudoscalar mesons, is given by [61]

$$
\frac{dBr}{ds_B} = \tau_B \frac{G_F^2 \alpha_{em}^2 |V_{ub}V_{ts}^*|^2}{2^{\delta+3} \pi^5} V_{tb} V_{ts}^* |M_B|^2 M_B^5 \lambda^{3/2}(s_B, \tilde{M}_P^2, 1) |f_+^P(s_B)|^2 |C_L^\nu + C_{LNP}^\nu|^2,
$$

(39)

where $\tilde{M}_P = M_P/M_B$ and $s_B = s/M_B^2$.

C. $B_{d(s)} \to (K^*, \phi, \rho)\nu_\ell\bar{\nu}_\ell$

The double differential decay rate of $B_{d(s)} \to V\nu_\ell\bar{\nu}_\ell$ processes, where, $V = K^*, \phi, \rho$ are the vector mesons, is given by [61]

$$
\frac{d^2\Gamma}{ds_B d\cos \theta} = \frac{3}{4} \frac{d\Gamma_T}{ds_B} \sin^2 \theta + \frac{3}{2} \frac{d\Gamma_L}{ds_B} \cos^2 \theta.
$$

(40)
Here \( \Gamma_L (T) \) are the longitudinal (transverse) part of decay rate

\[
\frac{d\Gamma_L}{ds_B} = 3M_B^2 |A_0|^2, \quad \frac{d\Gamma_T}{ds_B} = 3M_B^2 (|A_\perp|^2 + |A_\parallel|^2),
\]

where, the explicit expression for transversality amplitudes are given as

\[
A_\perp(s_B) = 2N^\nu \sqrt{2} \lambda^{1/2}(1, \tilde{M}_V^2, s_B)(C_L^\nu + C_L^{\nu\text{NP}}) \frac{V(s_B)}{(1 + \tilde{M}_V)},
\]

\[
A_\parallel(s_B) = -2N^\nu \sqrt{2}(1 + \tilde{M}_V)(C_L^\nu + C_L^{\nu\text{NP}})A_1(s_B),
\]

\[
A_0(s_B) = -N^\nu(C_L^\nu + C_L^{\nu\text{NP}}) \frac{M_V}{\sqrt{s_B}} \left[ (1 - \tilde{M}_V^2 - s_B)(1 + \tilde{M}_V)A_1(s_B) - \lambda(1, \tilde{M}_V^2, s_B) \frac{A_2(s_B)}{1 + \tilde{M}_V^2} \right],
\]

with

\[
N^\nu = V_{tb}V^*_{tq} \left[ \frac{G_F^2 M_B^3}{3 \cdot 2^{10} \pi^5} s_B \lambda^{1/2}(1, \tilde{M}_V^2, s_B) \right]^{1/2}, \quad \tilde{M}_V = M_V/M_B.
\]

For branching ratio computation in the SM, the \( B(s) \rightarrow V \) form factors are taken from [64] and remaining required input values from PDG [60].

**D. \( \tau \rightarrow \mu \nu_\tau \bar{\nu}_\mu \)**

The \( \tau \rightarrow \mu \nu_\tau \bar{\nu}_\mu \) process occurs via one loop box diagram in the presence of \( Z' \) boson as shown in Fig. 8.

![Box diagram of \( \tau \rightarrow \mu \nu_\tau \bar{\nu}_\mu \) process mediated by the \( Z' \) gauge boson.](image)

Including the \( Z' \) contribution, the total branching ratio of this process is given by [65]

\[
\text{Br}(\tau \rightarrow \mu \nu_\tau \bar{\nu}_\mu) = \text{Br}(\tau \rightarrow \mu \nu_\tau \bar{\nu}_\mu)^{\text{SM}} \left( 1 - \frac{3g_{BL}^2 \log(M_W^2/M_{Z'}^2)}{4\pi^2} \frac{1 - M_W^2/M_{Z'}^2}{1 - M_{Z'}^2/M_W^2} \right).
\]

The SM values and corresponding measurements of all the above defined processes involved in our analysis are presented in Table III. The exchange of \( Z' \) boson provides only \( C_{9\text{NP}}^{\nu\text{NP}} \) additional contributions to \( b \rightarrow (s, d)ll \), thus the leptonic \( B_{d,s} \rightarrow ll \) decays could not provide any strict bound on the new parameters. Since the considered model has no \( Z'l_i l_j \) couplings, the neutral and charged lepton flavor violating decay processes like \( B \rightarrow K^{(s)} l_i^+ l_j^\mp, l_i \rightarrow l_j \gamma, \)
TABLE III: The SM values and the respective experimental limits on the branching ratios of rare $B$ and $\tau$ decay modes.

| Decay modes                  | SM Values              | Experimental Limit |
|------------------------------|------------------------|--------------------|
| $B_d^0 \to \pi^0 \nu \bar{\nu}$ | $(3.252 \pm 0.26) \times 10^{-6}$ | $< 9 \times 10^{-6}$ |
| $B_d^0 \to K^0 \nu \bar{\nu}$ | $(4.55 \pm 0.34) \times 10^{-6}$ | $< 2.6 \times 10^{-5}$ |
| $B_d^0 \to K^0 \nu \bar{\nu}$ | $(9.54 \pm 0.66) \times 10^{-6}$ | $< 1.8 \times 10^{-5}$ |
| $B_d^0 \to K^0 \nu \bar{\nu}$ | $(1.11 \pm 0.09) \times 10^{-5}$ | $< 1.27 \times 10^{-4}$ |
| $B_u^0 \to \rho^0 \nu \bar{\nu}$ | $(7.624 \pm 0.587) \times 10^{-6}$ | $< 4.0 \times 10^{-5}$ |
| $B_u^+ \to \pi^+ \nu \bar{\nu}$ | $(1.25 \pm 0.099) \times 10^{-7}$ | $< 1.4 \times 10^{-5}$ |
| $B_u^+ \to K^+ \nu \bar{\nu}$ | $(1.752 \pm 0.128) \times 10^{-7}$ | $< 1.6 \times 10^{-5}$ |
| $B_u^+ \to K^+ \nu \bar{\nu}$ | $(1.03 \pm 0.08) \times 10^{-5}$ | $< 4.0 \times 10^{-5}$ |
| $B_u^+ \to \rho^+ \nu \bar{\nu}$ | $(8.16 \pm 0.653) \times 10^{-6}$ | $< 3.0 \times 10^{-5}$ |
| $B^0 \to \pi^0 e^+ e^-$ | $(7.66 \pm 0.62) \times 10^{-10}$ | $< 8.4 \times 10^{-8}$ |
| $B_d^0 \to \pi^0 \mu^+ \mu^-$ | $(7.67 \pm 0.575) \times 10^{-10}$ | $6.9 \times 10^{-8}$ |
| $B_d^0 \to K^+ e^+ e^-$ | $(1.53 \pm 0.1224) \times 10^{-7}$ | $(1.6^{+1.0}_{-0.8}) \times 10^{-7}$ |
| $B_d^0 \to K^0 \mu^+ \mu^-$ | $(1.51 \pm 0.1163) \times 10^{-7}$ | $(3.39 \pm 0.34) \times 10^{-7}$ |
| $B_u^+ \to \pi^+ e^+ e^-$ | $(8.26 \pm 0.645) \times 10^{-10}$ | $< 8.0 \times 10^{-8}$ |
| $B_u^+ \to \pi^+ \mu^+ \mu^-$ | $(8.27 \pm 0.579) \times 10^{-10}$ | $(1.76 \pm 0.23) \times 10^{-8}$ |
| $B_u^+ \to K^+ e^+ e^-$ | $(1.643 \pm 0.127) \times 10^{-7}$ | $(5.5 \pm 0.7) \times 10^{-7}$ |
| $B_u^+ \to K^0 \mu^+ \mu^-$ | $(1.626 \pm 0.122) \times 10^{-7}$ | $(4.41 \pm 0.23) \times 10^{-7}$ |
| $B_u^+ \to K^+ \tau^+ \tau^-$ | $(1.54 \pm 0.13) \times 10^{-7}$ | $< 2.25 \times 10^{-3}$ |
| $\tau \to \mu \nu_\tau \bar{\nu}_\mu$ | $(17.29 \pm 0.032)\%$ | $(17.39 \pm 0.04)\%$ |

$l_i \to l_j l_k \bar{l}_k$ do not play any role. Now using the existing limits on the branching ratios of allowed decay modes (Table III) and applying the relation $M_{Z'}/g_{BL} > 6.9$ TeV, the constraints on $g_{BL}$ and $M_{Z'}$ parameters are shown in orange color in Fig. 9 (Flavor). In this figure, the parameter space allowed by both DM and flavor studies (DM+Flavor) are graphically presented in cyan color. From Fig. 9, the bound on $\frac{M_{Z'}}{g_{BL}}$ is found to be greater than 7.14 (9.1) TeV from Flavor (DM+Flavor) case.

VI. REALIZATION OF LEPTOGENESIS IN THE PRESENT FRAMEWORK

So far in the current framework, we have discussed DM phenomenology, radiative neutrino mass and rare B decays. But this model does not accommodate the explanation for leptogenesis due to the exotic $B - L$ charges of the four heavy fermions, which effectively interact with the SM particles given in Eq. (2). Generation of asymmetry is only allowed after electroweak symmetry breaking, hence leptogenesis before sphaleron transition is not possible. But one can still explain the leptogenesis phenomena with a simple modification of the model by replacing the scalar triplet ($\Delta$) with a scalar singlet ($\phi_3$), which is assigned with a $B - L$ charge of $-\frac{1}{3}$. This allows an effective interaction of the exotic fermions with
The new potential for the scalar sector is given by

\[ V' = V(H, \phi_1, \phi_2, \phi_{DM}) + \mu_2 \phi_1^\dagger \phi_3 + \lambda_3 (\phi_3^\dagger \phi_3)^2 + \lambda_{H3} (H^\dagger H) (\phi_3^\dagger \phi_3) \]

\[ + \lambda_{13} (\phi_1^\dagger \phi_1) (\phi_3^\dagger \phi_3) + \lambda_{23} (\phi_2^\dagger \phi_2) (\phi_3^\dagger \phi_3) + \lambda_m ((\phi_3)^3 \phi_1 + (\phi_3^\dagger)^3 \phi_1^\dagger) \]

\[ + \lambda_{D3} (\phi_{DM}^\dagger \phi_{DM}) (\phi_3^\dagger \phi_3) + \mu_{D3} (\phi_{DM}^\dagger \phi_3 + \phi_{DM} \phi_3^\dagger) + \lambda_{DD} ((\phi_{DM}^\dagger \phi_3)^2 + (\phi_3^\dagger \phi_{DM})^2), \]

where, \( \phi_3 = \frac{v_3 + h_3^* + i A_3}{\sqrt{2}} \). The mixing between the CP even scalar sector is explained in detail in Eq. [10]. We neglect the mixing of CP even component of \( \phi_3 \) with other neutral Higgs due to mass decoupling and therefore the mass matrix remains same as in Eq. [11]. The mixing matrix of the CP odd sector is given by

\[ M_O^2 = \begin{pmatrix}
-2 \sqrt{2} \mu_{12} v_2 - \frac{\lambda_m v_3^2}{2 v_1} & \sqrt{2} \mu_{12} v_1 & -\frac{3}{2} \lambda_m v_3^2 \\
\sqrt{2} \mu_{12} v_1 & -\frac{\mu_{12} v_1^2}{2 v_2} & 0 \\
-\frac{3}{2} \lambda_m v_3^2 & 0 & -\frac{9}{2} \lambda_m v_1 v_3 
\end{pmatrix}. \]

To simplify the diagonalization, we assume \( v_2 = v_3 = v_1 = 1 \) TeV and \( \lambda_m = \frac{\mu_{12} v_1^2}{v_2^2} \). We then obtained two massive CP odd states with masses \( M_{A_1}^2 \approx 3 \mu_{12} v_2 \) and \( M_{A_2}^2 \approx 6 \mu_{12} v_2 \). The third eigenstate, \( A_3 \) remains massless and gets absorbed by the new gauge boson \( Z' \) to acquire mass, which is given by \( M_{Z'} = g_{BL} \sqrt{v_1^2 + 4 v_2^2 + \frac{1}{9} v_3^2} \).
A. Leptogenesis

Inclusion of a singlet scalar with retention of previous particle content (except scalar triplet) is sufficient to accommodate a detailed explanation for leptogenesis before sphaleron transition. Participation of this scalar singlet in B – L breaking at TeV scale leads to a two body decay of heavy exotic fermion to SM Higgs and lepton, with the interaction provided in Eq. (45). The asymmetry is generated in the final state leptons from the decay of these heavy fermions due to non-zero difference in particle and anti-particle decay width. Provided with the relation, \( Y_B = -\frac{8N_F + 4N_H}{22N_F + 13N_H} Y_L = -\frac{28}{79} Y_L \), a fraction of lepton asymmetry can be converted to baryon asymmetry through sphaleron transition. Here, \( N_F \) and \( N_H \) denote the number of fermion generations and Higgs doublets respectively. The tree and loop level decay of lightest heavy fermion is shown in Fig. 10. The general expression for the CP asymmetry is given by

\[
\epsilon_i = \frac{\sum_j \Gamma(N_i \to \ell_j H_1) - \Gamma(N_i \to \ell_j H_1^*)}{\sum_j \Gamma(N_i \to \ell_j H_1) + \Gamma(N_i \to \ell_j H_1^*)}.
\]

(48)

Flavor eigenstates of the heavy fermions are rotated with the mixing matrix \( U_N \) (section IIIB), which gives a double fold degenerate masses (\( M_1 = M_2, M_3 = M_4 \)), with two of them are lighter than the rest. Hence the asymmetry can be generated from the decay of any of the lightest mass eigenstate \( N_{D_1} \) or \( N_{D_2} \) to SM lepton and Higgs. We can neglect the contribution from the decay of heavier fermions, as they get washed out by the inverse decay of the lighter one. Therefore the final asymmetry will be summed up \( \epsilon_f = \epsilon_1 + \epsilon_2 = 2\epsilon_1 \), where, \( \epsilon_1 (\epsilon_2) \) is the CP asymmetry of \( N_{D_1} (N_{D_2}) \).

Interference of tree level decay with one loop self energy and vertex correction gives non-zero CP asymmetry, can be written as

\[
\epsilon_i = \frac{M_i}{M_j M_j} \frac{V}{2 + S} \text{Im} \left( \frac{\langle \tilde{Y} \gamma \tilde{Y}^\dagger \rangle_{ij}}{\langle \tilde{Y} \gamma \tilde{Y}^\dagger \rangle_{ii} (\tilde{Y} \gamma \tilde{Y}^\dagger)_{jj}} \right). 
\]

(49)
Here, the modified Yukawa coupling matrix $\tilde{Y}'$ is given by

$$\tilde{Y}' = Y' U_N, \quad Y' = \begin{pmatrix} 0 & Y'_{12} & Y'_{13} & Y'_{14} \\ 0 & Y'_{22} & Y'_{23} & Y'_{24} \\ 0 & Y'_{32} & Y'_{33} & Y'_{34} \end{pmatrix}, \quad \text{(50)}$$

$V$ and $S$ denote the vertex and self-energy contributions respectively, given by

$$V = \frac{2 M_j^2}{M_i^2} \left[ \left( 1 + \frac{M_j^2}{M_i^2} \right) \log \left( 1 + \frac{M_j^2}{M_i^2} \right) - 1 \right], \quad \text{(51)}$$

$$S = \frac{M_j^2 \Delta M_{ij}^2}{(\Delta M_{ij}^2)^2 + M_i^2 \Gamma_j^2}, \quad \frac{\Gamma_j}{M_j} = \frac{(Y'Y'^\dagger)_{jj} v^2}{8\pi\Lambda^2}, \quad \Delta M_{ij}^2 = M_j^2 - M_i^2. \quad \text{(52)}$$

In the above expression $\Gamma_j$ is the tree level decay width of the corresponding heavy fermion. Leptogenesis from the decay of heavy Majorana neutrinos with a hierarchical mass structure has been widely discussed in the literature [22, 68–70]. Those studies mainly focus on different cases like single flavor approximation and flavor consideration. In a high mass regime with one flavor approximation, the lower bound on heavy neutrino mass ($\gg \mathcal{O}(10^9)$ GeV) is obtained from the neutrino oscillation data [71], in order to explain the observed baryon asymmetry. But this mass scale of heavy fermions is very difficult to be tested in collider experiments. Above all these, a lot more work has been done on resonance enhancement of CP asymmetry in case of quasi degenerate Majorana neutrinos with a mass scale as low as TeV [24, 72–75], which seems to be more viable to test leptogenesis in future experiments.

In the current framework, CP asymmetry from the decay of lightest heavy fermion in one flavor regime is suppressed by $\frac{1}{\Lambda^2}$. Hence we consider the resonance enhancement by fixing the heavy fermion masses in TeV scale. We assume $M_1 = M_2 \approx M_3 (M_4)$ and introduce the new notation $\frac{\Delta M_{ij}^2}{M_j^2} = 1 - \frac{M_i^2}{M_j^2} \approx 0$ ($i = 1, 2, j = 3, 4$). With the approximation, $\Delta M_{ij}^2 \approx M_i \Gamma_j$, Eq.(52) gives maximum contribution from self energy with $S \approx \frac{M_i}{\Gamma_j} \gg 1$ and safely neglect the vertex contribution in this case. Thus CP asymmetry can be reduced...
to the form
\[ \epsilon_i \approx \frac{\text{Im} \left[ \left( \tilde{Y} \tilde{Y}^\dagger \right)_{ij} \right]}{\left( \tilde{Y} \tilde{Y}^\dagger \right)_{ii}} \frac{\left( \tilde{Y} \tilde{Y}^\dagger \right)_{jj}}{\left( \tilde{Y} \tilde{Y}^\dagger \right)_{jj}} \] (53)

Hence from the above expression, considering the Yukawa couplings in similar order, the CP parameter can be enhanced to achieve a value of order 1.

### B. Boltzmann Equation

The final baryon asymmetry depends on the efficiency of leptogenesis, which could be derived from the dynamics of relevant Boltzmann equations. When the gauge interaction rate is more than the Hubble expansion, particles attain thermal equilibrium and are subjected to the chemical equilibrium constraints. Hence the Boltzmann equations are so important to study the particle number density after the chemical or kinetic decoupling from the thermal bath in a specific temperature regime. In this model, we consider four right-handed exotic fermions, where, each two of them are degenerate. By considering the quasi degeneracy between the differently massed heavy fermions, resonant enhancement of the CP asymmetry is studied. As it has been discussed in the literature earlier that the lepton number violation demands the decay of the heavy fermion to be out of equilibrium to satisfy the Sakharov’s condition.

\[ H(T) = \frac{4\pi^3 g_* T^2}{45 M_{pl}} \quad Y_{\text{eq}}^N = \frac{135 \zeta(3) g_N}{16 \pi^4 g_*} z^2 K_2(z), \quad Y_{\text{eq}}^\ell = \frac{345 \zeta(3) g_\ell}{4 \pi^4 g_*}, \] (54)

where, \( g_* = 106.75 \), which is the total relativistic degree of freedom of the SM particles in the equilibrium. \( g_\ell = 2 \), \( g_N = 2 \) are the degrees of freedom of lepton and right-handed fermions respectively, and \( z = \frac{M}{T} \), with \( M \) being the mass of the decaying particle. The co-moving entropy density is given by \( s = (\frac{2\pi^2}{45}) g_* T^3 \), \( \zeta(3) \approx 1.202 \), and \( K_i(z) \) are the modified Bessel functions of type \( i \). The Boltzmann equations for the evolution of the number density of right-handed fermion and lepton are given by \[ 67, 76 \]

\[ \frac{dY_N}{dz} = -\frac{z}{sH(M_N)} \left[ \left( \frac{Y_N}{Y_{\text{eq}}^N} - 1 \right) (\gamma_D + 2\gamma_s + 4\gamma_t) + \left( \frac{Y_N^2}{Y_{\text{eq}}^2} - 1 \right) (\gamma_{Z'} + \gamma_{\phi_2}) \right], \] (55)

\[ \frac{dY_{B-L}}{dz} = -\frac{z}{sH(M_N)} \left[ \left( \frac{1}{2} \frac{Y_{B-L}}{Y_{\text{eq}}^\ell} + \epsilon_1 \left( \frac{Y_N}{Y_{\text{eq}}^N} - 1 \right) \right) \gamma_D + \frac{Y_{B-L}}{Y_{\text{eq}}^\ell} \gamma_w \right]. \] (56)

Here, \( \gamma_s \) and \( \gamma_t \) are the s and t-channel scattering cross sections of the lightest right-handed fermion and \( \gamma_w \) corresponds to the sum of all possible washout processes that reduce the lepton asymmetry. \( \gamma_{Z'} \) is the \( Z' \) mediated scattering processes and \( \gamma_{\phi_2} \) is the t-channel scattering of heavy fermion to pair of \( \phi_2 \), which maximally contribute to the heavy fermion number density. The relevant Feynman diagrams for these scattering processes are displayed in Fig. 11. The detailed expressions for \( \gamma_{Z'} \) and \( \gamma_{\phi_2} \) are provided in Appendix A. The left panel of Fig. 12 represents the freeze out of the lightest heavy fermion and the generation of lepton asymmetry of the order \( 10^{-9} \) for a Yukawa coupling of order \( \approx 10^{-5} \). In this case, we considered \( \gamma_{\phi_2} = 0 \) and only finite contribution from \( \gamma_{Z'} \). The value of lepton asymmetry
remains compatible with the observed baryon asymmetry $Y_B = kY_L \approx O(10^{-10})$, where, $k$ stands for the efficiency factor. Along with we fixed the CP violation parameter $\epsilon_i = 0.1$ and showed the impact of different values of Dirac Yukawa coupling on the baryon number density in the right panel of Fig. 12. Here, we considered the contribution from both $\gamma_{Z'}$ and $\gamma_{\phi_2}$, which play a significant role in reducing the heavy fermion and lepton number density. We analyze the Boltzmann equations by fixing $M_{Z'} = 4$ TeV, $M_i = O(1)$ TeV, $g_{BL} = 0.3$ and the corresponding Yukawa coupling of the lightest heavy fermion to be order of 0.1. We found that by switching off the dominant scattering mediated by $Z'$ boson, still $\gamma_{\phi_2}$ significantly reduces the number density of decaying fermion. In the whole analysis, we have considered only these two scattering process, while the rest are neglected due to higher dimension suppression.

C. Comments on dark matter and neutrino mass

We now compute the DM relic density with the contribution from new channels mediated by $\phi_3$, projected in Fig. 13. By fixing the new Higgs mass to be 2.5 TeV, the resonance is found at $M_{DM} = 1250$ GeV, which is displayed in Fig. 14. Coming to the neutrino mass, we can have a tree level Dirac mass for the active neutrinos, which can be constructed from the 5-dimension Yukawa coupling in Eq. (45). Mixing and
diagonalization of the Majorana fermion masses are already discussed in section III. Hence we can construct a tree level small neutrino Majorana mass matrix within type I seesaw framework. i.e

$$m_\nu = M_D M_R^{-1} M_D^T, \quad (57)$$

where,

$$M_D = \frac{v v_3}{\Lambda} \begin{pmatrix} 0 & Y_{12} & Y_{13} & Y_{14} \\ 0 & Y_{22} & Y_{23} & Y_{24} \\ 0 & Y_{32} & Y_{33} & Y_{34} \end{pmatrix}. \quad (58)$$

We show a sample benchmark by fixing $M_{iR} \approx 2$ TeV, $v_3 = 1$ TeV, $\Lambda = 10$ TeV and $Y_\nu = 10^{-5}$, gives a neutrino mass of order 0.03 eV.

VII. CONCLUSION

In this article, we have addressed dark matter phenomenology, lepton asymmetry, radiative neutrino masses and rare $B$ decay modes in a simple $U(1)_{B-L}$ gauge extension of standard model. Four exotic fermions with fractional $B-L$ charges are included to make the model free of triangle gauge anomalies. We have computed relic density and direct detection cross section of the singlet scalar, whose stability is assured by the $U(1)$ gauge symmetry. The channels contributing to relic density are mediated by scalars and $Z'$ boson. We have constrained the new parameters of the proposed model, by imposing Planck Satellite data on relic density and PandaX limit on spin independent DM-nucleon scattering cross-section. Along with, we obtained the constraints on the $Z'$ mass and the gauge coupling from LEP-II and ATLAS dilepton study. We have imposed additional constraint on the new gauge parameters from the available data in the quark and lepton sectors. Since there is no new contribution to $C_{10}$ coefficient, one could not constrain the new parameters from the leptonic $B$ decay modes. The $Z'$ boson has no lepton flavor violating couplings, thus the channels like $B \to K^{(*)} l_i l_j$, $l_i \to l_j \gamma$ and $l_i \to 3l_j$ do not play any role. Hence, we have only considered the branching ratios of rare semileptonic lepton flavor conserving $B$ and $\tau$ decays to compute the allowed parameter space. We have also explained the radiative mechanism of neutrino
mass at one loop level, where the mass splitting between the real and imaginary parts of the scalar DM is obtained from the effective interaction with the scalar triplet.

This four fermion model with scalar triplet extension is not sufficient to explain the leptogenesis phenomena before the electroweak symmetry breaking. Hence the asymmetry in the lepton sector cannot be converted to baryon sector through electroweak sphaleron process. But this interesting model can still accommodate the explanation for lepton asymmetry by a simple modification of extending with a scalar singlet instead of a scalar triplet. Thus one can explain the leptogenesis, without appreciable difference in the DM and flavor sectors. We discussed the resonant leptogenesis phenomena with doubly quasi degenerate heavy fermions. We have obtained a lepton asymmetry of the order $10^{-9}$ by solving the Boltzmann equations governing the particle dynamics, which is compatible with the observed baryon asymmetry ($\approx \mathcal{O}(10^{-10})$). To conclude, we have performed a combined study of the dark matter, neutrino mass and rare $B$ decays in a new variant of $U(1)_{B-L}$ gauge extended framework, along with the lepton asymmetry with a slight modification in the particle spectrum, which satisfy all the current respective experimental bounds.

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Appendix A: Decay and scattering cross section of heavy Majorana fermion

The general expression for the decay and thermally averaged cross sections of any fermion are given by [67]

\[
\gamma_D = \gamma^\psi_\psi(\psi \rightarrow i + j + \ldots) = n_\psi K_1(z) \Gamma_\psi, \tag{A1}
\]

\[
\gamma_{\psi,a}(\psi + a \rightarrow i + j + \ldots) = \frac{T}{64\pi^4} \int_{(m_\psi + m_a)^2}^{\infty} ds \hat{\sigma}(s) \sqrt{s} K_1 \left( \frac{\sqrt{s}}{T} \right). \tag{A2}
\]

Here, $K_1$ and $K_2$ are the modified Bessel functions of type 1 and 2 respectively. $n_\psi$ and $\Gamma_\psi$ are the equilibrium number density and the tree level decay rate of the decaying particle. The t-channel scattering cross section for the process $N_i + N_i \rightarrow \phi_2 + \phi_2$ is given by

\[
\hat{\sigma}_{N,\phi_2}(s) = \frac{y_N^2}{8\pi} \frac{x - 4}{x} \left( -2 + \frac{x}{2} + \frac{x^2 - 8x + 16}{x\sqrt{x(x-4)}} \log \frac{x - \sqrt{x(x-4)}}{x + \sqrt{x(x-4)}} \right), \tag{A3}
\]

where $x = \frac{s}{M_i^2}$ and $y_N$ is the heavy fermion Dirac like coupling. The $Z'$ mediated heavy fermion scattering cross section for the process $N_i + N_i \rightarrow f + \tilde{f}$ is given by

\[
\hat{\sigma}_{Z'}(s) = \frac{104\pi}{3} g_{BL}^2 \frac{\sqrt{x}}{(x-y)^2 + yc (x-4)^2}, \tag{A4}
\]

In the above expression, $y = \frac{M_{Z'}}{M_i^2}$ and $c = \left( \frac{\Gamma_{Z'}}{M_i} \right)^2$. The tree level decay width of $Z'$ is given by

\[
\Gamma_{Z'} = \frac{g_{BL} M_{Z'}}{6} \left( 3(1 - 4/y)^2 \theta(y - 4) + 13 \right). \tag{A5}
\]
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