Thermal keV neutrino dark matter in minimal gauged B-L model with cosmic inflation

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We investigate the possibility of keV scale thermal dark matter in a minimal gauged B − L extension of the standard model with three right-handed neutrinos in the context of cosmic inflation.

The complex singlet scalar field responsible for the spontaneous breaking of B − L gauge symmetry is non-minimally coupled to gravity and serves the role of inflaton. The lightest right-handed neutrino \( N_1 \) can be a dark matter candidate if its couplings to leptons are sufficiently suppressed or forbidden. While keV scale \( N_1 \) gives rise to the possibility of warm dark matter, its thermal production leads to over-abundance. The subsequent entropy dilution due to heavier right-handed neutrino decay can bring the DM abundance within the observed limit. We constrain the model parameters from the requirement of producing sufficient entropy dilution, inflationary observables along with other phenomenological constraints. While these requirements cannot satisfy light neutrino data, suitable extension of the minimal model can accommodate it.

I. Introduction

The existence of dark matter (DM) occupying a significant amount in the energy budget of the universe is firmly established after the observations of different astrophysical and cosmological experiments, as summarised in review articles [1–3]. The observed anisotropies in the cosmic microwave background (CMB) measurements provide a precise measurement of the DM abundance quoted as \( \Omega_{DM} h^2 = 0.120 \pm 0.001 \) at 68% CL [4, 5], where \( \Omega \) refers to density parameter and \( h \) is the reduced Hubble constant in the unit of 100 km/sec/Mpc. Since none of the standard model (SM) particles has the required properties of a DM particle\(^1\), construction of beyond Standard model (BSM) frameworks has become a necessity. Among different BSM frameworks for such particle DM, weak interacting massive particle (WIMP) is the most popular one where a DM particle having mass and interactions similar to those around the electroweak scale gives rise to the observed relic after thermal freeze-out, a remarkable coincidence often referred to as the WIMP Miracle [6]. Such interactions enable the WIMP DM to be produced in thermal equilibrium in the early universe and eventually its number density gets frozen out when the rate of expansion of the Universe takes over the interaction rates. Such DM candidates typically remain non-relativistic at the epochs of freeze-out as well as matter-radiation equality and belong to the category of Cold Dark Matter (CDM). However, the absence of any observational hints for WIMP at direct search experiments namely LUX [7], PandaX [8] and Xenon1T [9] or collider experiments such as the large hadron collider (LHC) [10, 11] has also encouraged the particle physics community to pursue other viable alternatives.

As mentioned above, the HDM paradigm is disfavoured due to its relativistic nature giving rise to a large free streaming length (FSL) [12, 13] which can erase small-scale structures. One intermediate possibility is the so-called Warm Dark Matter (WDM) scenario where DM can remain semi-relativistic during the epoch of matter radiation equality. The FSL of WDM also falls in the intermediate regime between those of CDM and HDM. Typically, WDM candidates have masses in the keV regime which once again stays intermediate between sub-eV scale masses of HDM and GeV-TeV scale masses of CDM. A popular choice of WDM candidate is a right-handed sterile neutrino, singlet under the SM gauge symmetry. A comprehensive review of keV sterile neutrino DM can be found in [14]. To assure the stability of DM, the sterile neutrino should have tiny or zero mixing with the SM neutrinos leading to a lifetime larger than the age of the Universe. The lower bound on WDM mass arises from observations of the Dwarf spheroidal galaxies and Lyman-\( \alpha \) forest which is around 2 keV [12, 15–17]. Although WDM may not be traced at typical direct search experiments or the LHC, it can have interesting signatures at indirect search experiments. For example, a sterile neutrino WDM candidate having mass 7.1 keV can decay on cosmological scales to a photon and a SM neutrino, indicating an origin to the unidentified 3.55 keV X-ray line reported by analyses of refs. [18] and [19] with the data collected by the XMM-Newton X-ray telescope. On the other hand, WDM can have interesting astrophysical signatures as it has the potential to solve some of the small-scale structure problems of the CDM paradigm [20]. For example, due to small FSL, the CDM paradigm gives rise to formation of structures at a scale as low as that of the solar system, giving rise to tensions with astrophysical observations. The WDM paradigm,

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\(^1\) It is worth mentioning that neutrinos in the SM satisfy some of the criteria for being a good DM candidate, but due to their sub-eV masses, they remain relativistic at the epoch of freeze-out as well as matter radiation equality, giving rise to hot dark matter (HDM) which is ruled out by both astrophysics and cosmology observations.
due to its intermediate FSL can bring the predictions closer to observations.

The CMB experiments also reveal that our universe is homogeneous and isotropic on large scales up to an impressive accuracy. Such observations lead to the so-called horizon and flatness problems which remain unexplained in the description of standard cosmology. Theory of cosmic inflation which accounts for the presence of a rapid accelerated expansion phase in the early universe, was proposed in order to alleviate these problems [21–23]. The Higgs field from the SM sector [24], has the ability to serve the role of inflaton, but this particular description suffers from the breakdown of perturbative Unitarity [25]. The simple alternative is to extend the SM of particle physics by a gauge singlet scalar which can be identified with the inflaton.

Motivated by the pursuit of finding a common framework for cosmic inflation and WDM, in this work, we adopt one of the most economical extensions of the SM namely, the gauged $B − L$ scenario which includes three right handed neutrinos (RHN) required to cancel the gauge anomalies. The model also contains a complex singlet scalar to spontaneously break the $B − L$ gauge symmetry while simultaneously generating RHN masses. The model is known to naturally accommodate active neutrino masses (which SM by itself can not explain) via type-I seesaw. The study of cosmic inflation with the scalar singlet as the inflaton in such a minimal framework has been worked out earlier in [26, 27]. Further extension of the minimal $B − L$ model with a discrete symmetry $\mathbb{Z}_2$ can provide a stable dark matter candidate which is the lightest RH neutrino. An attempt to simultaneous realisation of inflation and cold dark matter (both thermal and non-thermal) has been performed in [28]. Here, we envisage the scope of realizing warm dark matter in the same inflationary framework which also offers correct magnitude of spectral index and tensor to scalar ratio view of most recent Planck 2018+BICEP/Keck data [29] published in 2021.

The thermally produced keV scale warm dark matter in minimal gauged $U(1)_{B−L}$ model turns overabundant. As a remedy, the relic can be brought down to the observed limit by late time entropy production due to decay of another long lived RH neutrino. Earlier, such dilution mechanism of keV scale dark matter abundance from long lived heavier RHN decay has been studied in a broader class of models having $U(1)_{B−L}$ gauge symmetry without accommodating cosmic inflation by the authors of [30–35]. Here we revisit the analysis in the minimal gauged $B−L$ model by including the cosmic inflation. We consider the lightest RHN $N_1$ to be the WDM candidate. We assume $N_2$ to be in thermal equilibrium at the early Universe and its late decay to deplete the overproduced thermal keV scale DM abundance by entropy dilution. The requirement of successful inflation provides strong bounds on additional parameters namely, the $B−L$ gauge coupling and the singlet scalar-RHN coupling and combination of them. Such constraints have important impact on thermal history of $N_2$ and subsequently on WDM phenomenology. These constraints, in general, can affect the realisation of thermalisation condition for $N_2$ (also $N_1$) as well as the decoupling temperature of $N_2$. The decoupling temperature of $N_2$ or the freeze out abundance of $N_2$ is very sensitive to the amount of late entropy injection which in turn, dilutes the overproduced thermal WDM relic. In addition, the mass scales of the additional gauge boson and the heavier RHNs can not be arbitrary in view of constraints arising from inflationary dynamics. With these impressions, the present study is important to combine thermally produced keV scale dark matter phenomenology in presence of long-lived $N_2$ with cosmic inflation in minimal $B−L$ gauged model which is not performed earlier to the best of our knowledge. The constraints from these requirements are so severe that the minimal model fails to satisfy the light neutrino data requiring further extensions.

The paper is organized as follows. In section II, we briefly discuss the minimal gauged $B−L$ model followed by summary of scalar singlet inflation via non-minimal coupling to gravity in section III. In section IV, we discuss the procedure to calculate the relic density of keV RHN DM followed by discussion of results in section V. Finally, we conclude in section VI.

II. Gauged $B−L$ Model

As mentioned earlier, gauged $B−L$ extension of the SM [36–41] is one of the most popular BSM frameworks. In the minimal version of this model, the SM particle content is extended with three right handed neutrinos ($N_R$) and one complex singlet scalar ($\Phi$) all of which are singlet under the SM gauge symmetry. The requirement of triangle anomaly cancellation fixes the $B-L$ charge for each of the RHNs as -1. The complex singlet scalar is having $B−L$ charge 2 and not only leads to spontaneous breaking of gauge symmetry but also generate RHN masses dynamically.

The gauge invariant Lagrangian of the model can be written as

$$\mathcal{L} = \mathcal{L}_{SM} - \frac{1}{4} B'_{\alpha\beta} B'^{\alpha\beta} + \mathcal{L}_{\text{scalar}} + \mathcal{L}_{\text{fermion}} ,$$

(1)

where $\mathcal{L}_{SM}$ represents the SM Lagrangian involving quarks, gluons, charged leptons, left handed neutrinos and electroweak gauge bosons. The second term in the $\mathcal{L}$ indicates the kinetic term of $B−L$ gauge boson ($Z_{BL}$), expressed in terms of field strength tensor $B'^{\alpha\beta} = \partial^\mu Z^\beta_{BL} - \partial^\beta Z^\alpha_{BL}$. The gauge invariant scalar Lagrangian of the model (involving SM Higgs $H_S$ and singlet scalar $\Phi$) is given by,

$$\mathcal{L}_{\text{scalar}} = (D_\mu H_S)^\dagger (D_\mu H_S) + (D_\mu \Phi)^\dagger (D^\mu \Phi) - V(H_S, \Phi) ,$$

(2)
where,
\[ V(H_S, \Phi) = -\mu_1^2 |H_S|^2 - \mu_2^2 |\Phi|^2 + \lambda_1 |H_S|^4 \\
+ \lambda_2 |\Phi|^4 + \lambda_3 |H_S|^2 |\Phi|^2. \] (3)

The covariant derivatives of scalar fields are written as,
\[ D_\mu H_S = \left( \partial_\mu + i \frac{g_1}{2} \sigma_\mu W_\mu \right) H_S, \]
\[ D_\mu \Phi = \left( \partial_\mu + i 2g_{BL} Z_{BL\mu} \right) \Phi, \] (5)

with \( g_1 \) and \( g_2 \) being the gauge couplings of \( SU(2)_L \) and \( U(1)_Y \) respectively and \( W_\mu^a \) \((a = 1, 2, 3)\), \( B_\mu \) are the corresponding gauge fields. On the other hand \( Z_{BL}, g_{BL} \) are the gauge boson and gauge coupling respectively for \( U(1)_{B-L} \) gauge group. There can also be a kinetic mixing term between \( U(1)_Y \) of SM and \( U(1)_{B-L} \) of the form \( \epsilon B_{\alpha\beta} B_{\alpha\beta} \) with \( \epsilon \) being the mixing parameter. While this mixing can be assumed to be vanishing at tree level, it can arise at one-loop level as \( \epsilon \approx g_{BL} g_2/(16\pi^2) \) [40]. As we will see in upcoming sections, the \( B-L \) gauge coupling \( g_{BL} \) has tight upper bound from inflationary dynamics, and hence the one-loop kinetic mixing can be neglected in comparison to other relevant couplings and processes. Therefore, for simplicity, we ignore such kinetic mixing for the rest of our analysis.

The gauge invariant Lagrangian involving RHNs can be written as
\[ \mathcal{L}_{\text{fermion}} = i \sum_{k=1}^3 N_{R_k} \overline{Q}^{iR} \gamma^\nu (Q_{kR}) N_{R_k} - \sum_{j=1}^3 Y_D^{jj} \overline{D}^{j} \dot{H}_S N_{R_j}^c \]
\[ - \sum_{i,j=1}^3 Y_{N_{ij}} \Phi \overline{N}_{R_i} N_{R_j} + \text{h.c.} \] (6)

The covariant derivative for \( N_{R_k} \) is defined as
\[ \overline{D}\Phi(Q_{kR}) N_{R_k} = \gamma^\mu \left( \partial_\mu + i g_{BL} Q^{kR}_{BL\mu} \right) N_{R_k}, \] (7)

with \( Q^{kR} = -1 \) is the \( B-L \) charge of right handed neutrino \( N_{R_k} \). Hereafter, we denote the RHNs by \( N_i, i = 1, 2, 3 \) only without explicitly specifying their chirality. The lightest RHN namely, \( N_1 \) is identified as the DM candidate. If we consider an additional \( Z_2 \) symmetry to stabilise the DM, \( N_1 \) will have no mixing with heavier RHNs \( N_{2,3} \) as well as with SM leptons. If we do not have such additional discrete symmetry, we can tune the corresponding couplings of \( N_1 \) with other fermions appropriately in order to make it long-lived cosmologically.

After spontaneous breaking of both \( B-L \) symmetry and electroweak symmetry, the SM Higgs doublet and singlet scalar fields are expressed as,
\[ H_S = \left( \begin{array}{c} H_S^+ \\ h + v/\sqrt{2} \end{array} \right), \]
\[ \Phi = \left( \begin{array}{c} \phi + v_{BL} + iA' \\ v_{BL} \end{array} \right) \] (8)

where \( v \) and \( v_{BL} \) are vacuum expectation values (VEVs) of \( H_S \) and \( \Phi \) respectively. The right handed neutrinos and \( Z_{BL} \) acquire masses after the \( U(1)_{B-L} \) breaking as,
\[ M_{Z_{BL}} = 2g_{BL} v_{BL}, \] (9)
\[ M_{N_i} = \sqrt{g_{BL}} Y_{N_i} v_{BL}. \] (10)

We have considered diagonal Yukawa matrix \( Y_N \) in the \((N_1, N_2, N_3)\) basis. Using Eq.(9) and Eq.(10), it is possible to relate \( M_{Z_{BL}} \) and \( M_{N_i} \) by,
\[ M_{N_i} = \frac{1}{\sqrt{g_{BL}}} Y_{N_i} M_{Z_{BL}}. \] (11)

After the spontaneous breaking of \( SU(2)_L \times U(1)_Y \times U(1)_{B-L} \) gauge symmetry, the mixing between scalar fields \( h \) and \( \phi \) appears and can be related to the physical mass eigenstates \( H_1 \) and \( H_2 \) by a rotation matrix as,
\[ \begin{pmatrix} H_1 \\ H_2 \end{pmatrix} = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} h \\ \phi \end{pmatrix}, \] (12)

where the scalar mixing angle \( \theta \) is found to be
\[ \tan 2\theta = \frac{-3v v_{BL}}{(\lambda_3 v v_{BL} - 2\lambda_3 v v_{BL})}. \] (13)

The mass eigenvalues of the physical scalars are given by,
\[ M_{H_1}^2 = 2\lambda_1 v^2 \cos^2 \theta + 2\lambda_2 v_{BL}^2 \sin^2 \theta - 2\lambda_3 v v_{BL} \sin \theta \cos \theta, \] (14)
\[ M_{H_2}^2 = 2\lambda_1 v^2 \sin^2 \theta + 2\lambda_2 v_{BL}^2 \cos^2 \theta + 2\lambda_3 v v_{BL} \sin \theta \cos \theta. \] (15)

Here \( M_{H_1} \) is identified as the SM Higgs mass whereas \( M_{H_2} \) is the singlet scalar mass.

One of the strong motivations of the minimal \( U(1)_{B-L} \) model is the presence of heavy RHNs which can yield correct light neutrino mass via type I seesaw mechanism. The analytical expression for the light neutrino mass matrix is
\[ m_\nu = m_D^T M_N^{-1} m_D, \] (16)

where \( m_D = Y_D v/\sqrt{2} \). We consider the right handed neutrino mass \( M_N \) to be diagonal. Since in our case \( N_1 \) does not interact with SM leptons (or interacts very feebly), the lightest active neutrino would be massless (or vanishingly small).

There are several theoretical and experimental constraints that restrict the model parameters of the minimal \( U(1)_{B-L} \) model. To begin with, the criteria to ensure the scalar potential bounded from below yields following conditions involving the quartic couplings,
\[ \lambda_{1,2,3} \geq 0, \lambda_3 + \sqrt{\lambda_1 \lambda_2} \geq 0 \] (17)

On the other hand, to avoid perturbative breakdown of the model, all dimensionless couplings must obey the following limits at any energy scale,
\[ \lambda_{1,2,3} < 4\pi, |Y_D, Y_N| < \sqrt{4\pi}, |g_1, g_2, g_{BL}| < \sqrt{4\pi}. \] (18)
The non-observation of the extra neutral gauge boson in the LEP experiment [43, 44] imposes the following constraint on the ratio of $M_{Z_{BL}}$ and $g_{BL}$:

$$\frac{M_{Z_{BL}}}{g_{BL}} \geq 7 \text{ TeV.} \quad (19)$$

The recent bounds from the ATLAS experiment [45, 46] and the CMS experiment [47] at the LHC rule out additional gauge boson masses below 4-5 TeV from analysis of 13 TeV centre of mass energy data. However, such limits are derived by considering the corresponding gauge coupling $g_{BL}$ to be similar to the ones in electroweak theory and hence the bounds become less stringent for smaller gauge couplings.

Additionally, the parameters associated with the singlet scalar of the model are also constrained [48, 49] due to the non-zero scalar-SM Higgs mixing. The bounds on scalar singlet-SM Higgs mixing angle arise from several factors namely $W$ boson mass correction [50] at NLO, requirement of perturbativity and unitarity of the theory, the LHC and LEP direct search [51, 52] and Higgs signal strength measurement [52]. If the singlet scalar turns lighter than SM Higgs mass, SM Higgs can decay into a pair of singlet scalars. Latest measurements by the ATLAS collaboration restrict such SM Higgs decay branching ratio into invisible particles to be below 13% [53] at 95% CL.

In our case, we work with very small singlet scalar-SM Higgs mixing and considered all the scalar quartic couplings to be positive. These help in evading the bounds on scalar singlet mixing angle and the boundedness of the scalar potential from below. We choose the magnitude of the relevant couplings below their respective perturbativity limits. The LEP (LHC) bound on $M_{Z_{BL}}$ and $g_{BL}$ is also taken care of.

### III. Cosmic Inflation

Here we briefly talk about the dynamics of inflation in view of most recent data from combination of Planck and BICEP/Keck [29]. For a detailed discussion of inflation in the minimal B-L model, we refer [26, 27]. We identify the real part of $\phi$ of singlet scalar field $\Phi$ as the inflaton. Along with the renormalisable potential in Eq.(3), we also assume the presence of non-minimal coupling of $\Phi$ to gravity. The potential that governs the inflation is given by

$$V_{\text{inf}}(\phi) = \frac{\lambda_2}{4} \phi^4 + \xi \frac{\phi^2}{2} R,$$  \quad (20)

where $R$ represents the Ricci scalar and $\xi$ is a dimensionless coupling of singlet scalar to gravity. We have neglected the contribution of $v_{BL}$ in Eq.(20) by considering it to be much lower than the Planck mass scale ($M_P$). With this form of potential, the action for $\phi$ in Jordan frame is expressed as,

$$S_J = \int d^4x \sqrt{-g} \left[ -\frac{M_P^2}{2} \Omega(\phi)^2 R + \frac{1}{2} (D_\mu \phi)^\dagger (D^\mu \phi) - \frac{\lambda_2}{4} \phi^4 \right],$$  \quad (21)

where $\Omega(\phi)^2 = 1 + \frac{\xi \phi^2}{M_P^4}$, $g$ is the spacetime metric in the $(-,+,+,+)$ convention, $D_\mu \phi$ stands for the covariant derivative of $\phi$ containing couplings with the gauge bosons which reduces to the normal derivative $D_\mu \rightarrow \partial_\mu$ (since during inflation, the SM and BSM fields except the inflaton are non-dynamical).

Following the standard prescription, we use the following conformal transformation to write the action $S_J$ in the Einstein frame [54, 55] as

$$\tilde{g}_{\mu\nu} = \Omega^2 g_{\mu\nu}, \quad \sqrt{-\tilde{g}} = \Omega^4 \sqrt{-g},$$  \quad (22)

so that it resembles a regular field theory action of minimal gravity. In the above equation, $\tilde{g}$ represents the metric in the Einstein frame. Furthermore, to make the kinetic term of the inflaton appear canonical, we transform the $\phi$ by

$$\frac{d\chi}{d\phi} = \sqrt{\frac{\Omega^2 + \frac{\xi \phi^2}{M_P^4}}{\Omega^4}} = Z(\phi),$$  \quad (23)

where $\chi$ is the canonical field. Using these inputs, the inflationary potential in the Einstein frame can be written as,

$$V_E(\phi(\chi)) = \frac{V_J(\phi(\chi))}{\Omega(\phi(\chi))^4} = \frac{1}{4} \frac{\lambda_2 \phi^4}{(1 + \frac{\xi \phi^2}{M_P^4})^2},$$  \quad (24)

where $V_J(\phi)$ is identical to $V_{\text{inf}}(\phi)$ in Eq. (20). We then make another redefinition: $\tilde{\Phi} = \frac{\phi}{\sqrt{1 + \frac{\xi \phi^2}{M_P^4}}}$ and reach at a much simpler from of $V_E$ given by

$$V_E(\Phi) = \frac{1}{4} \lambda_2 \Phi^4.$$  \quad (25)

Note that for an accurate analysis, one should work with renormalisation group (RG) improved potential and in that case, $\lambda_2$ in Eq. (25) will be function of $\Phi$ such that,

$$V_E(\Phi) = \frac{1}{4} \lambda_2(\Phi) \Phi^4.$$  \quad (26)

The full set of one loop renormalisation group evolution (RGE) equations of the relevant parameters associated with the inflationary dynamics can be found in [27]. We consider a diagonal RH neutrino mass matrix with the hierarchy $M_{N_1} \ll M_{N_2} < M_{N_3}$. This implies $Y_{N_1} \ll Y_{N_2} < Y_{N_3}$, where $M_{N_i} = \frac{Y_{N_i} v}{v}$. A simplified form for the RGE equation of $\lambda_2$ can be written by assuming $Y_{N_{2,3}} g_{BL} \gg \lambda_2, \lambda_3$ leading to the following beta function,

$$\beta_{\lambda_2} = 96 g_{BL}^4 - Y_{N_2}^4 - Y_{N_3}^4.$$  \quad (27)
Below we provide the three important conditions to realise a successful RGE improved inflation in minimal gauged $B - L$ model.

- In general, for $\xi \lesssim 1$, the self-quartic coupling of inflaton must be very small to be in agreement with the experimental bounds on inflationary observables [56]. Since $\lambda_2$ is very small, any deviation of $\beta_{\lambda_2}$ from zero (due to larger $g_{BL}$ and $Y_N$), if sufficient, may cause sharp changes in $\lambda_2$ value from its initial magnitude during the RGE running. This may trigger unwanted instability to the inflationary potential [27, 28]. Therefore keeping the $\beta_{\lambda_2}$ in the vicinity of zero during inflation is a desired condition for successful inflation. To ensure $\beta_{\lambda_2} \sim 0$, the equality $\Delta = 96g_{BL}^2 - 82Y_N^4 \sim 0$ needs to be maintained, where we have assumed $Y_N = 3Y_{N_2}$ and $g_{BL}^2 Y_{N_2}^2 \gg \lambda_2$.

- In addition to the stability ($\beta_{\lambda_2} \sim 0$) confirmation, the inflationary potential should also be monotonically increasing function of the inflaton field value which implies $\frac{d^2 \lambda_2}{d \Phi^2} > 0$ during inflation. It is reported in [27, 28] that with the increase of $g_{BL}$, a local minimum appears (due to the violation of $\frac{d^2 \lambda_2}{d \Phi^2} > 0$) within the inflationary trajectory, which can stop the inflaton from rolling. This poses an upper bound on the size of $B - L$ gauge coupling.

- Finally, another constraint on $g_{BL}$ comes from the criteria that the inflationary predictions (e.g. spectral index and tensor to scalar ratio) stay within the $1 - \sigma$ allowed range as provided by Planck+BICEP experiment. This particular bound on $g_{BL}$ is stronger than the one arising due to appearance of a local minimum along the inflationary trajectory as we shall show in a while.

Next, we use the standard definitions of slow roll parameters ($\epsilon$ and $\eta$) [57] while calculating the magnitude of spectral index ($n_s = 1 - 6\epsilon + 2\eta$) and tensor to scalar ratio ($r \sim 16\epsilon$). We consider the number of e-folds ($N_e$) as 60. We also impose the condition $\Delta = 0$ at the scale of horizon exit of inflaton while estimating the inflationary observables.

We perform a numerical scan over $g_{BL}$ and $\xi$ to estimate the inflationary observables $n_s$ and $r$ considering $\Delta = 0$. We have observed that the factor $\lambda_2$ does not alter the slow roll parameters, rather it is fixed by the measured value of scalar perturbation spectrum ($P_S = 2.4 \times 10^{-9}$) at horizontal exit of inflaton. It also turns out that the value of $r$ does not change much with the variation of $g_{BL}$ for a constant value of $\xi$ since $\beta_{\lambda_2}, \Delta \sim 0$ at inflationary energy scale. Contrary to this, value of $n_s$ is quite sensitive to $g_{BL}$ as it involves second order derivative of the inflationary potential. In the left panel of Fig. 1, we show the $n_s - r$ predictions for different $g_{BL}$ lines (with varying $\xi$) and check its viability against the improved version of Planck+BICEP/Keck data [29] published in 2021. In the right panel of Fig. 1, we constrain the $\xi - g_{BL}$ plane by using the criteria of yielding correct values for $n_s$ and $r$ allowed by the combined Planck+BICEP/Keck data [29]. The maximum permitted value (green dashed line) of $g_{BL}$ as function of $\xi$ is also depicted in the same figure which corresponds to non-appearance of any local minimum along the inflationary trajectory. The parameter $\lambda_2$ at inflationary energy scale can be fixed with the observed value of scalar perturbation spectrum as earlier mentioned. As an example, we find $\lambda_2^{\text{inf}} \simeq O(10^{-10})$ considering $\xi = 1$. We shall use this particular reference point in our DM analysis.

After inflation ends, the oscillation regime of inflaton starts and the universe enters into the reheating phase. Since the minute details of reheating phase is not much relevant for the present work, we discuss it briefly here. We followed the instantaneous perturbative reheating mechanism and compute the reheating temper-
ature. In principle, such description is not fully complete since the particle production from inflaton can also happen during very early stages of oscillation which is known as preheating mechanism. In general, this approach always makes the true reheating temperature larger than what we find by using the perturbative approximation. The estimate of reheating temperature is important to ensure that it is bigger than the mass scale $M_{Z_{BL}}$ since we talk about thermal dark matter and thermalisation of few other heavier particles (with masses $\lesssim M_{Z_{BL}}$) as well. Therefore, for any parameter space, if we can satisfy the thermalisation criteria using the perturbative approach, it automatically means that a more involved analysis of reheating dynamics would not change our conclusion. We adopt this conservative approach in our work. The inflaton can decay to SM fields owing to its small mixing with the SM Higgs. In addition, inflaton decaying to BSM fields ($Z_{BL}, N_1$) is also possible if kinematically allowed, with these BSM fields further decaying into SM fields. We have found that for the concerned ranges of the relevant parameters in DM phenomenology (to be discussed in upcoming sections), the reheating temperature always turns larger than $M_{Z_{BL}}$ with $\xi = 1$ and $\lambda_3 \sim O(10^{-7})$.

### IV. WDM relic

In this section, we discuss the details of WDM relic calculation. As mentioned before, keV scale dark matter $N_1$ gets thermally overproduced requiring late entropy injection from heavier RHNs (we consider late decay of $N_2$). Since entropy is not conserved, this requires us to solve coupled Boltzmann equations for dark matter candidate $N_1$, the heavier right-handed neutrino $N_2$ and temperature of the universe ($T$) which can be written as function of scale factor $a$ as follows [58, 59],

$$\frac{dE_{N_1}}{da} = \frac{\langle \sigma v \rangle_1}{H a^4} \left( (E_{N_1}^{eq})^2 - E_{N_1}^2 \right),$$

$$\frac{dE_{N_2}}{da} = \frac{\langle \sigma v \rangle_2}{H a^4} \left( (E_{N_2}^{eq})^2 - E_{N_2}^2 \right) - \frac{\Gamma_{N_2}}{H a} E_{N_2},$$

$$\frac{dT}{da} = \left( 1 + T \frac{dg_{*s}}{dT} \right)^{-1} \left[ - \frac{T}{a} + \frac{\Gamma_{N_2} M_{N_2}^2}{3H} \right],$$

where the entropy density is expressed by $s = \frac{2\pi^2}{45} g_{*s} T^3$, with $g_{*s}$ being the relativistic entropy degrees of freedom. The Boltzmann equations are solved using the approximation that $N_2$ freezes out while being non-relativistic in a way similar to WIMP type DM belonging to CDM category. This can be ensured if $N_2$ is heavy enough $M_{N_2} \sim O(M_{Z_{BL}})$. The quantities $E_{N_{1,2}}$ in the above equations are defined as the co-moving number densities respectively, i.e., $E_{N_{1,2}} = n_{1,2} g_{*s}$ where $n_{1,2}$ are the usual number densities. The Hubble parameter is denoted by $H$ and $\Gamma_{N_2}$ corresponds the decay width of $N_2$. The interaction cross-sections for a light dark matter ($N_1$) is dominated by the $Z_{BL}$ mediated annihilations into SM fermions [60]. On the other hand, the heavier RHN $N_2$ can annihilate to SM particles as well as $\{Z_{BL}Z_{BL}, Z_{BL}H_2, Z_{BL}H_2, H_2H_2, N_1N_1\}$ final states. We incorporate all the relevant processes in thermally averaged annihilation cross sections for both $N_1$ and $N_2$ as denoted by $\langle \sigma v \rangle_1$ and $\langle \sigma v \rangle_2$ respectively. Here we do not include the heavier RHN $N_3$ considering it decays promptly at early stage and has negligible impact on DM phenomenology. Considering long-lived $N_3$ will not change our conclusion significantly and hence we stick to the simple scenario mentioned above.

The energy density and the corresponding Hubble parameter in a radiation dominated universe are defined as

$$\rho_R = \frac{\pi^2}{30} g_{*s} T^4, \quad H = \sqrt{\frac{\rho_R}{3 M_P^2}},$$

with $g_{*s}$ is the effective relativistic degrees of freedom and $M_P$ denotes the reduced Planck scale. At earlier epoch, $N_2$ was part of radiation bath by virtue of its gauge and Higgs portal interactions with the SM particles. When $N_2$ becomes non-relativistic and goes out of equilibrium it can be safely treated as matter. In that case the Hubble parameter is redefined as,

$$H = \sqrt{\frac{M_N E_{N_2}}{a^3} + \frac{\rho_R}{3 M_P^2}}. \quad (31)$$

Now, the dark matter $N_1$ freezes out when it is relativistic. The decoupling temperature of $N_1$ is determined by the strength of the interactions, mainly governed by $g_{BL}$ and $M_{Z_{BL}}$. In the left panel of Fig. 2, we show the evolution of the DM interaction rate ($\Gamma$) and the expansion rate all the relevant processes in thermally averaged annihilation of $N_1$, i.e.,

$$\left( \frac{\sigma v_1}{\rho_{DM}} \right) = 76.4 \times \left( \frac{\rho_{DM}}{M_{N_1} \text{ keV}} \right), \quad (32)$$

where $g$ is the number of internal degrees of freedom of the DM candidate, $g_*(x_f)$ represents the number of relativistic entropy degrees of freedom at the instant of $N_1$ decoupling $x_f = \frac{M_{DM}}{T_f}$, where $T_f$ is the freeze-out temperature. Considering the DM to decouple around the EW scale (with $g_*(x_f) \approx 106.75$), we find

$$\Omega_{DM} h^2 = 76.4 \left( \frac{3 \times 2}{4 \times 106.75} \right) \left( \frac{M_{N_1}}{\text{keV}} \right) \approx 0.3. \quad (33)$$
Thus, DM remains overproduced roughly by one order of magnitude or by a factor of around 50. This overabundance can be brought down by the late decay of $N_2$ after the DM freeze-out, which injects entropy ($s$) into the thermal bath [58]. In order to realise this possibility of sufficient entropy dilution, it is necessary for the long-lived $N_2$ to dominate the energy density of the universe at late time over the radiation component. Importantly, the validity of Eq. (32) breaks down if the universe has such non-standard history where $N_2$ dominates the energy density of the universe for some epoch (similar to early matter domination scenarios [61]). In that case, we need to solve the system of Boltzmann equations numerically (for analytical approach, see [62]) and obtain the relic abundance of dark matter using

$$\Omega_{\text{DM}} h^2 = 2.745 \times 10^9 m_{\text{DM}} Y_{a \rightarrow \infty},$$  \hspace{1cm} (34)$$

where $Y_{a \rightarrow \infty} = \frac{E_{N_1}}{m_{a \rightarrow \infty}}$.

Earlier we found that the additional gauge coupling $g_{BL}$ is restricted by the inflationary requirements. The strongest bound on $g_{BL}$ ($\lesssim 0.07$) comes from providing correct values of scalar spectral index $n_s$ and tensor-to-scalar ratio $r$ for $\xi = 1$. Note that these bounds can be relaxed for very large values of $\xi \gg 1$ which we do not consider here, keeping all dimensionless parameters within order one. Moreover, in order to ensure the stability of the inflationary potential, we require the condition $96g_{BL}^4 - Y_{N_2}^4 - Y_{N_3}^4 \sim 0$ to hold true at inflationary energy scale. We consider $Y_{N_3} = 3Y_{N_2}$ in our analysis. It is also possible that $Y_{N_1}, Y_{N_2} \ll Y_{N_3}$ and $\Delta \approx 96g_{BL}^4 - Y_{N_1}^4$. In this regime, $N_2$ is significantly lighter than $N_3$ and can freeze out while being relativistic in the early universe. As shown in [24] for left-right symmetric model (LRSM), such light $N_2$ can lead to larger entropy dilution. However, in minimal $B - L$ model, we have much less free parameters and decay modes of RHNs compared to that of LRSM and hence we do not consider this possibility here and stick to heavy $N_2$ which freezes out like WIMP. Also, since our effort in this work is to utilise the maximal constraints arising from inflation sector, we use the compressed limit $M_{N_2} = 3M_{N_3}$ at inflationary energy scale which also fixes the value of $Y_{N_2}$ from the stability condition as defined earlier. We then consider the RGE running of all the relevant couplings from inflationary energy scale up to the $B - L$ breaking scale and obtain the values for $M_{N_2}, M_{Z_{BL}}$ and other mass parameters. It turns out that the RGE running effects on $g_{BL}$ and $Y_{N_2}$ are very minimal in view of bringing any noticeable change to the DM phenomenology.

\[\text{V. Results and Discussions}\]

We solve the system of Boltzmann equations (Eq. (28), (29), (30)) in the post-reheating era as function of the scale factor ranging from $a_i$ to $a_f$, where we define $T_i/a_i = 1$. The relevant independent parameters which are crucial for DM phenomenology are given by

$$\left\{ g_{BL}, M_{Z_{BL}}, M_{N_1}, \Gamma_{N_2} \right\} \hspace{1cm} (35)$$

The mass scale $M_{N_2}$ and $M_{Z_{BL}}$ are connected through the stability condition $\Delta \sim 0$ at inflationary energy scale. The scalar sector couplings are considered to be small and have little impact on the DM phenomenology except the $\lambda_2$ which fixes $M_{H_2}$. The decay of $N_2$ can occur dominantly at tree level to SM Higgs and leptons in the final states. The decay width of the $N_2$ is proportional to $|Y_{N_1}|^2M_{N_3}$, where we have always considered $M_{N_2} > M_{H_2} \approx 125$ GeV.

We first stick to a benchmark scenario (as given in table I) to explain the dynamics of the DM phenomenology. Recall that the Yukawa coupling $Y_{N_2}$ (or mass $M_{N_2}$) can not be chosen arbitrarily since it is fixed by $\Delta = 0$ condition at inflationary energy scale. The solution of the
FIG. 3. Left panel: A comparison between the energy densities of radiation and $N_2$ are shown from early time to late epoch. Right panel: Enhancement of radiation energy density is shown due to $N_2$ decay at late time.

FIG. 4. Left panel: Evolution of $E_{N_2}^{eq}$ (solid red) and $E_{N_2}$ (solid black) as function of scale factor. Right panel: The evolution of comoving DM density (dashed green) as function of scale factor where the effect of late time entropy dilution can be clearly observed. The black solid line corresponds to the same in absence of any entropy injection while the solid blue line corresponds to the required comoving number density at present epoch from observations.

| Parameters | Values |
|------------|--------|
| $g_{BL}$   | 0.01   |
| $M_{\nu_{BL}}$ | 100 TeV |
| $M_{N_1}$  | 5 keV  |
| $\Gamma_{N_2}$ | $10^{-21}$ GeV |

TABLE I. Benchmark point used for Figs. 3-5.

coupled Boltzmann equations for the benchmark point in table I are shown in Figs. 3-5.

- The left panel of Fig. 3 shows the variation of radiation and $N_2$ energy densities from early time to late epoch. Initially the radiation dominates over $N_2$ abundance as usual, followed by an intermediate epoch where $N_2$ starts dominating. The radiation domination is recovered again after $N_2$ starts decaying into SM particles. The enhancement of radiation energy density ($E_R = \rho_R a^4$) from out-of-equilibrium decay of $N_2$ at late time can also be observed from the right panel of Fig. 3.

- In the left panel of Fig. 4, evolution of $E_{N_2}$ as function of scale factor is shown. The solid black line indicates the scale factor dependence of $E_{N_2}^{eq}$ while the actual abundance $E_{N_2}$ is shown by the solid red line. Initially $N_2$ was part of radiation bath and freezes out when the interaction rates goes below the Hubble parameter. At late epochs (for large $a$) $N_2$ dominates the energy budget of the universe and finally decays into radiation. In the right panel of Fig. 4, the evolution of comoving DM density (dashed green) is shown as function of the scale factor.
factor which goes through late time dilution due to sizeable entropy production from \( N_2 \) decay. For the chosen benchmark point, the final relic abundance exactly matches with the measured one by Planck, shown by the solid blue line.

- In the left panel of Fig. 5, we draw a comparison between the interaction rate of \( N_1 \) and the Hubble parameter by considering an intermediate \( N_2 \) dominated era. This clearly shows that DM reaches thermal equilibrium at early epochs, prior to \( N_2 \) domination. Also, the decoupling temperature of \( N_2 \) is larger than the \( N_1 \). In addition to this, \( N_1 \) decouples before the \( N_2 \) domination in the Hubble parameter sets in. The right panel of Fig. 5 shows the temperature evolution with the scale factor which goes through late time dilution due to sizeable entropy production from \( N_2 \) decay. In between for a brief period, different pattern of the temperature evolution is observed which is due to \( N_2 \) domination and the entropy injection into the radiation bath.

We then perform a numerical scan by fixing \( \Gamma_{N_2} \) at two different benchmark values to constrain the \( g_{BL} - M_{Z_{BL}} \) plane from the condition of satisfying the five requirements namely, (i) \( N_2 \) thermalisation at early epochs, (ii) \( N_1 \) thermalisation at early epochs, (iii) bound on \( g_{BL} \) from inflationary dynamics, (iv) LHC bounds and (v)
correct final relic of DM. The resulting parameter space after applying all relevant bounds is shown in Fig. 6. The left panel is for $\Gamma_{N_2} = 10^{-16}$ GeV while in the right panel we have $\Gamma_{N_2} = 10^{-21}$ GeV. The mass of $N_2$ is varied in the range of $10^3 - 10^8$ GeV as shown in the figure. The light blue and light green regions are ruled out due to non-thermalisation of $N_1$ and $N_2$ respectively. The region corresponding to larger values of $g_{BL}$ are disfavoured by inflationary criteria mentioned earlier. The pink shaded region corresponds to DM overabundance or equivalently, insufficient entropy dilution from $N_2$ decay. The disfavored region from the LHC in the same plane is also highlighted by the brown shaded region. The relatively weaker bound from the LHC results from the search for high mass dilepton resonances. Note that the expression for $\Gamma_{N_2}$ contains $M_{N_2}$ which varies for each point (satisfying the condition $\Delta = 0$) in the $g_{BL} - M_{Z_{BL}}$ plane, however a constant $\Gamma_{N_2}$ can always be attained by tuning the neutrino Yukawa coupling $Y_{\nu_2}$. We identify the allowed region (white) where the first four constraints mentioned above are satisfied and also DM relic corresponds $\Omega h^2 \lesssim 0.12$. Along the boundary line between pink and white coloured regions corresponding to correct DM relic $\Omega_{N_2} h^2 \sim 0.12$, larger $g_{BL}$ requires heavier $Z_{BL}$. This attributes to the fact that, for a fixed $M_{Z_{BL}}$, a larger value of $g_{BL}$ makes the $N_2$ to decouple later from the radiation bath. This leads to suppressed freeze out abundance of $N_2$ and eventually reduced amount of entropy production from its decay. To obviate that, one needs to raise $M_{Z_{BL}}$ accordingly in order to decrease the interaction cross section for $N_2$ and successively obtaining the correct relic abundance. In the same context, a larger $\Gamma_{N_2}$ (early decay of $N_2$) for constant $g_{BL}$ and $M_{Z_{BL}}$ means lower impact on dilution of DM abundance. Hence we see less amount of allowed region (white) in the left panel of Fig. 6 compared to the right panel with lower $\Gamma_{N_2}$.

To show the dependence on $N_2$ decay width further, in Fig. 7, we obtain the relic satisfied region ($\Omega h^2 \lesssim 0.12$) in $M_{DM} - \Gamma_{N_2}$ plane considering two different sets of $(g_{BL}, M_{Z_{BL}})$. The allowed space shrinks significantly for larger $g_{BL}$ since the freeze-out of $N_2$ is delayed compared to the smaller $g_{BL}$ value leading to less entropy dilution. Also, for larger DM mass, the required order of $\Gamma_{N_2}$ is smaller to bring the DM abundance within the desired limit. This is due to the fact that larger DM mass yields enhanced relic and hence needs large amount of entropy injection to the SM bath at late epochs for sufficient dilution of overproduced thermal DM relic. It is pertinent to comment here that the decay width of $N_2$ cannot be arbitrary small. Decay of $N_2$ around or after the epoch of big bang nucleosynthesis (BBN) may raise the neutrino temperature [63] which is strictly restricted by the number of relativistic degrees of freedom during BBN as measured by Planck [5]. In view of this, it is safe to keep lifetime ($\tau_{N_2}$) of $N_2$ typically above one second. We have highlighted the corresponding bound on $\Gamma_{N_2}$ (red) in Fig. 7.

**Neutrino mass:** We have seen that, it requires very small decay width of the $N_2$ for adequate entropy production that dilutes the DM relic. This, in turn, will require tiny Dirac Yukawa coupling of $N_2$ with SM leptons. Note that out of the three RHNs in minimal $B - L$ model, we have already made $N_1$ effectively decoupled from neutrino mass generation in order to be a stable or long-lived DM candidate. While two RHNs can also generate light neutrino masses and mixing successfully, the requirement of feeble $N_2$ coupling with SM leptons may lead to conflict with light neutrino data. In Fig. 8, we have shown the relic satisfied region ($\Omega h^2 \lesssim 0.12$) for two $g_{BL}$ values with $M_{N_1} = 2$ keV. The higher $g_{BL}$ re-
quires larger $M_{ZBL}$ to obey the relic bound in consistent

with the results discussed earlier. We have also shown the effective light neutrino mass line ($\nu_L \sim 0.01$ eV) generated by type-I seesaw considering only $N_2$ as RHN. Here the dependence of $\nu_L$ on $g_{BL}$ goes away due to the inflationary condition $\Delta = 0$ mentioned before. As clearly seen from this figure, the neutrino mass line never enters into the relic allowed region. Thus simultaneous realisation of neutrino mass with cosmic inflation and thermal keV dark matter phenomenology is not possible in the minimal $B - L$ gauged model. Lowering the $g_{BL}$ further, may reduce the gap between the relic and the light neutrino mass line, however such case may be disfavored by thermalisation of both $N_1$ and $N_2$ in the early universe. One can of course have a scenario where keV scale DM is purely non-thermal and produced from the freeze-in mechanism [28] at the cost of very fine-tuned gauge coupling $g_{BL}$ where not only the observational aspects get limited, but the inflationary dynamics also get decoupled from the gauge sector, unlike in our present setup.

As mentioned earlier, keV scale warm dark matter can alleviate some of the small scale structure issues of cold dark matter paradigm. Due to larger free-streaming length, such WDM scenarios can give rise to different structure formation rates which can be probed by several galaxy survey experiments. For example, the authors of [64] constrained thermal WDM mass around keV scale by comparing the predictions of DM sub-structures in Milky Way with the estimates of total satellite galaxy population. Similar bounds on thermal WDM mass were also derived in [65] by using constraints from stellar streams. In addition to these astrophysical aspects, such keV scale DM can also have interesting indirect detection aspects. In the absence of any additional $Z_2$ symmetry, DM can have tiny Yukawa with SM leptons by virtue of which keV scale $N_1$ can decay into a light neutrino and a photon [66] radiatively, providing a plausible explanation for the 3.55 keV X-ray line reported by analyses [18] and [19] of the data collected by the XMM-Newton X-ray telescope.

VI. Conclusion

We have revisited the widely studied minimal gauged $B - L$ model in order to check the possibility of realizing cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass. A compatible picture of cosmic inflation, keV scale warm dark matter and light neutrino mass.

We have also observed that to achieve correct relic abundance of dark matter the model prefers heavier $Z_{BL}$ ($\gtrsim O(1)$ TeV ) with larger lifetime of $N_2$. A long lived $N_2$ automatically implies the associated neutrino Dirac Yukawa coupling to be small, insufficient to produce light neutrino mass of correct order. Suitable extensions of this minimal model can accommodate light neutrino data. One can suitably extend the minimal model by a scalar triplet which can satisfy light neutrino data within a type-II seesaw mechanism. We leave exploration of such non-minimal models to future works.

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