

Thermal leptogenesis, dark matter, and gravitational waves from an extended canonical seesaw scenario

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In a canonical type-I seesaw scenario, the Standard Model is extended with three singlet right-handed neutrinos (RHNs) N_i , $i = 1, 2, 3$ with masses M_i , $i = 1, 2, 3$ to simultaneously explain sub-eV masses of light neutrinos and baryon asymmetry of the Universe at high scales. In this paper, we show that a relatively low-scale thermal leptogenesis accompanied by gravitational wave signatures is possible when the type-I seesaw is extended with a singlet fermion (S) and a singlet scalar (ρ), where S and ρ are odd under a discrete Z_2 symmetry. We also add a vectorlike fermion doublet Ψ and impose a Z'_2 symmetry under which both N_1 and Ψ are odd while all other particles are even. This gives rise to a singlet-doublet Majorana fermion dark matter in our setup. At a high scale, the Z_2 symmetry is broken spontaneously by the vacuum expectation value of ρ and leads to (i) mixing between RHNs (N_2, N_3) and S and (ii) formation of Domain walls (DWs). In the former case, the final lepton asymmetry is generated by the out-of-equilibrium decay of S , which dominantly mixes with N_2 . We show that the scale of thermal leptogenesis can be lowered to $M_S \sim 2 \times 10^6$ GeV, which is 3 orders of magnitude lower than the thermal leptogenesis in canonical type-I seesaw. In the latter case, the disappearance of the DWs gives observable gravitational wave signatures, which can be probed at LISA, DECIGO, μ ARES etc.

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I. INTRODUCTION

Recently NANOGrav [1,2], EPTA [3], PPTA [4] have reported positive evidence for stochastic gravitational wave (GW) background in the nanohertz (nHz) frequency range. This stochastic GW background could originate from various sources. One such possibility is the domain walls (DWs), which are two-dimensional topological defects that emerge when a discrete symmetry is spontaneously broken in the early Universe [5]. In cosmology, the formation of DWs possess a significant challenge because their energy density can quickly surpass the total energy density of the Universe, which contradicts current observational data [6]. However, it's possible that DWs are inherently unstable and collapse before they can dominate the total energy density of the Universe. This instability could be due to an explicit breaking of the discrete

symmetry in the underlying theory [7–9]. If this is the case, a considerable amount of GWs could be generated during the collisions and annihilations of these domain walls [10–16]. These GWs might persist as a stochastic background in the universe today. Detecting these GWs would offer insights into early cosmic events and provide a novel method for exploring physics at extremely high energies.

In this paper, we try to connect the nonzero neutrino mass, dark matter (DM), observed baryon asymmetry of the Universe, and the GW sourced by the DWs in a common framework as shown pictorially in Fig. 1. With this motivation, we extend the canonical type-I seesaw model [17–21] with a singlet fermion S and one singlet scalar ρ . We impose a discrete Z_2 symmetry under which S and ρ are odd while all other particles are even. This forbids the direct coupling of S with the L, H . We also introduce a vectorlike fermion doublet $\Psi \equiv (\psi^0 \psi^-)^T$ and impose a discrete symmetry Z'_2 under which Ψ and the lightest RHN (N_1) are odd while all other particles are even. This gives rise to a singlet-doublet Majorana fermion DM [22–45]. At a high scale, when the singlet scalar ρ obtains a vacuum expectation value (vev), the Z_2 symmetry gets broken and leads to (i) mixing between the singlet fermion S and RHNs (N_2, N_3) and (ii) formation of Domain walls (DWs) [6,7,46–49]. In the former case, the S decays to the Standard Model (SM)

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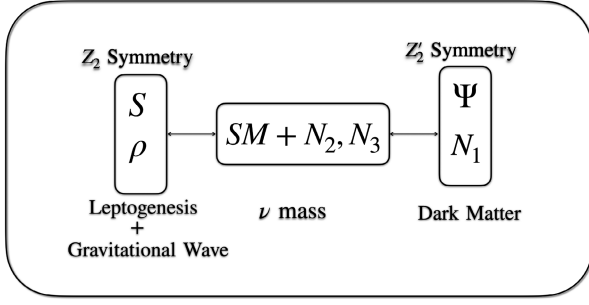


FIG. 1. The schematic picture of our setup.

lepton, L and H via the mixing with $N_{2,3}$. Assuming a hierarchy among the S and $N_{2,3}$, the final lepton asymmetry is established by the CP violating out-of-equilibrium decay of S . Due to the mixing, the decay width of $S \rightarrow LH$ will be suppressed compared to the usual type-I case $N \rightarrow LH$. As a result, the out-of-equilibrium condition can be satisfied for lower values of the mass of S . The CP asymmetry parameter does not have any suppression. It depends on the value of M_S unlike M_N in the usual case. However, the value of M_S cannot be arbitrarily small since it will make the CP asymmetry parameter ϵ_S very small. We find that the leptogenesis scale can be brought down to $\mathcal{O}(10^6)$ GeV, which is 3 orders of magnitude smaller than the thermal leptogenesis in canonical type-I seesaw. For low scale leptogenesis in various beyond Standard Model frameworks, see, e.g., incorporating flavor effects [50], by adding extra scalar fields [51–54], resonant leptogenesis [55], right-handed sector leptogenesis [56,57], wash-in leptogenesis [58], and by decoupling neutrino mass and leptogenesis [59–62].

Another intriguing consequence of the discrete Z_2 symmetry breaking is the formation of DWs. We analyze the dynamics of these DWs. To ensure their instability, we introduce an explicit Z_2 -breaking term in the scalar potential, creating a pressure difference across the walls. This causes the DWs to collapse and annihilate, leading to the production of stochastic GWs. Crucially, in our framework, it is the vev of ρ that establishes the connection between leptogenesis and GW signals. We identify the parameter space where both leptogenesis and GW production occur simultaneously.

The paper is organized as follows. We discuss the model in Sec. II. In Sec. III, we discuss a low energy thermal leptogenesis scenario arising from a Z_2 symmetry breaking. The domain wall dynamics and signature of gravitational waves are discussed in Sec. IV. The DM phenomenology is discussed in Sec. V. We finally conclude in Sec. VI.

II. THE MODEL

We extend the SM with three RHNs, N_i , $i = 1, 2, 3$ with masses M_i , $i = 1, 2, 3$, one singlet fermion S with mass M_S , and one singlet scalar ρ .

TABLE I. Particles and their charge assignment under the imposed symmetry.

Field	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	Z_2	Z'_2
N_1	1	1	0	+	–
$N_{2,3}$	1	1	0	+	+
S	1	1	0	–	+
ρ	1	1	0	–	+
Ψ	1	2	–1	+	–

We impose a discrete symmetry Z_2 under which S and ρ are odd, while all other particles are even as listed in Table I. As a result, the direct coupling of S with L , H is forbidden. We also introduce a vectorlike fermion doublet $\Psi \equiv (\psi^0 \psi^-)^T$ and impose an additional discrete symmetry Z'_2 under which Ψ and N_1 are odd, while all other particles are even.¹ The neutral component of Ψ combines with N_1 to give rise to a singlet-doublet fermion dark matter. The relevant Lagrangian involving N , S , and ρ is given as

$$\mathcal{L} = \bar{N} i \gamma_\mu \partial^\mu N + \bar{S} i \gamma_\mu \partial^\mu S - y_{N1} \bar{L} \tilde{H} N - y_{NS} \bar{N} \rho S - \frac{1}{2} M_S \bar{S} S - \frac{1}{2} M_N \bar{N} N + \text{H.c.}, \quad (1)$$

where N represents N_2, N_3 .

The scalar potential involving the SM Higgs doublet H and singlet scalar ρ can be written as

$$V(H, \rho) = -\mu_H^2 H^\dagger H - \mu_\rho^2 \rho^2 + \lambda_H (H^\dagger H)^2 + \lambda_\rho \rho^4 + \lambda_{H\rho} (H^\dagger H) \rho^2. \quad (2)$$

Due to negative mass-squared term of ρ , the latter acquires a vev: $\langle \rho \rangle \neq 0$, which leads to the spontaneous breaking of Z_2 symmetry. The minimum of the potential is at

$$\langle \rho \rangle = v_\rho = \pm \sqrt{\frac{\mu_\rho^2}{\lambda_\rho}}. \text{ Here, } \mu_\rho = \sqrt{\lambda_\rho v_\rho^2}. \text{ At finite temperature, interaction of the scalar field } \rho \text{ with others fields may induce a thermal correction to the potential, which leads to thermal mass of } \rho. \text{ In our model, } \rho \text{ interacts with other fermions via the Yukawa interaction term } y_{NS} N S \rho. \text{ This leads to an effective potential of the form [63,64]}$$

$$V_{\text{eff}}(\rho) = -\mu_\rho^2 \rho^2 + \lambda_\rho \rho^4 + c T^2 \rho^2 = (-\mu_\rho^2 + c T^2) \rho^2 + \lambda_\rho \rho^4, \quad (3)$$

where $c = \frac{y_{NS}^2}{4} + \frac{\lambda_{H\rho}}{12} + \frac{\lambda_\rho}{2}$. At high temperature, the $c T^2$ term dominates over the $-\mu_\rho^2$ term, and the sign of the effective mass term changes when

¹This restricts the $N_1 L H$ coupling, thus resulting in only two nonzero SM light neutrino masses.

$$m_{\text{eff}}^2(T) = -\mu_\rho^2 + cT^2 > 0. \quad (4)$$

This leads to a critical temperature for symmetry restoration, which is given by

$$T_C = \sqrt{\frac{\mu_\rho^2}{c}}. \quad (5)$$

Thus, for $T > T_C$, the quadratic coefficient of the field becomes positive, which sets $v_\rho = 0$ and restores the Z_2 symmetry. For $T < T_C$, $v_\rho \neq 0$ and the symmetry is broken. The vev of the ρ is then given at different temperatures as

$$v_\rho(T) = \begin{cases} 0 & T \geq T_C \\ \sqrt{\frac{\mu_\rho^2 - cT^2}{\lambda_\rho}} & T < T_C \\ \sqrt{\frac{\mu_\rho^2}{\lambda_\rho}} \equiv v_\rho & T = 0. \end{cases} \quad (6)$$

Therefore, the mass of the ρ at different temperatures can be given as

$$M_\rho(T) = \begin{cases} 0 & T \geq T_C \\ \sqrt{2(\mu_\rho^2 - cT^2)} & T < T_C \\ \sqrt{2}\mu_\rho & T = 0. \end{cases} \quad (7)$$

The electroweak symmetry breaking (EWSB) occurs when the SM Higgs obtains a vev v_h . The vacuum fluctuations of the scalar fields at zero temperature are given as

$$H = \frac{1}{\sqrt{2}}(0, h + v_h)^T, \quad \rho = \frac{\rho' + v_\rho}{\sqrt{2}}, \quad (8)$$

where the vevs are $v_h = \sqrt{\frac{\mu_H^2}{\lambda_H}}$, and $v_\rho = \sqrt{\frac{\mu_\rho^2}{\lambda_\rho}}$.

A. Neutrino mass

Above electroweak phase transition (EWPT), the singlet scalar ρ acquires a vev: v_ρ and breaks the Z_2 symmetry spontaneously. This allows the singlet fermion S to mix with the RHNs N_2 and N_3 . In the effective theory, the 6×6 fermion mass matrix can be written in the basis $[L_i, N_{2,3}, S]$ as

$$\mathcal{M} = \begin{pmatrix} 0 & m_D & 0 \\ m_D & M_N & d \\ 0 & d & M_S \end{pmatrix}, \quad (9)$$

where $d = y_{NS}v_\rho/\sqrt{2}$, $m_D = y_{NI}v_h/\sqrt{2}$.

Diagonalizing this mass matrix, we obtain the masses of heavy eigenstates as

$$M'_N \simeq M_N + \frac{d^2}{M_N - M_S}, \quad (10)$$

$$M'_S \simeq M_S - \frac{d^2}{M_N - M_S}, \quad (11)$$

where the mixing angle is given as

$$\theta \simeq \frac{d}{M_N - M_S}. \quad (12)$$

The light neutrino mass matrix is then given as

$$(m_\nu)_{ij} \simeq -\sum_k (m_D)_{ik} \left(M_k + \frac{d^2}{M_k - M_S} \right)^{-1} (m_D)_{kj}. \quad (13)$$

The canonical type-I seesaw can be restored in the limit $d \rightarrow 0$.

III. THERMAL LEPTOGENESIS FROM Z_2 SYMMETRY BREAKING

In the type-I seesaw, the CP violating out-of-equilibrium decay of heavy neutrinos to the SM sector (leptons and Higgs) can generate a net lepton asymmetry, which can be converted to the baryon asymmetry via the electroweak sphalerons. In our scenario, the lepton asymmetry can be generated from the direct decay of N_2, N_3 as well as from the decay of S via its mixing with N_2 , where we assume the mixing with N_3 is negligible. We are interested in the region where the final asymmetry is produced mainly due to mixing suppressed decay of S . However, while calculating the lepton asymmetry, we simultaneously take into account the effects of N_2 and N_3 along with S . We have also taken into account the scale dependence of the coupling by solving the renormalization group (RG) equations as given in Appendix B. We used SARAH 4.13.0 [65] to calculate the RG equations. The values of the couplings listed in Table II are taken at the EW scale. The decay width of the $N_{2,3}$ is given by

$$\Gamma_i = \frac{(y_{NI}^\dagger y_{NI})_{ii}}{8\pi} M_i, \quad (14)$$

where $M_i, i = 2, 3$ are the masses of N_2 and N_3 , respectively. The decay width of the singlet fermion S is then given by

$$\Gamma_S = \theta_s(T)^2 \frac{(y_{NI}^\dagger y_{NI})_{22}}{8\pi} M_S, \quad (15)$$

where $\theta_s(T)$ represents the mixing between $S - N_2$ at finite temperature and is given by

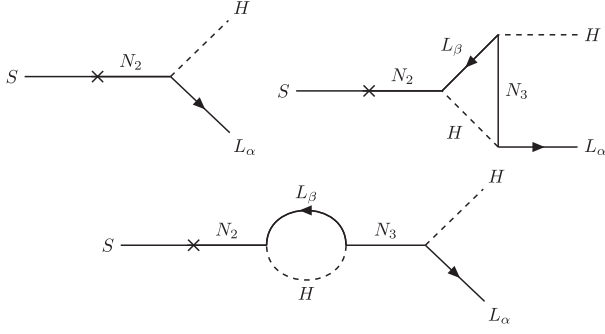


FIG. 2. Tree and one loop diagrams of S decay giving rise to CP asymmetry.

$$\theta_s(T) = \frac{y_{NS} v_\rho(T)}{\sqrt{2}(M_2 - M_S)}. \quad (16)$$

The Yukawa coupling, y_{NI} in Eq. (15) can be calculated by using the Casas-Ibarra parametrization as [66]

$$y_{NI} = i \frac{\sqrt{2}}{v_h} \left(U_{PMNS}^* \cdot \sqrt{\hat{m}_\nu} \cdot R^T \cdot \sqrt{\hat{M}_N} \right), \quad (17)$$

where U_{PMNS} is the lepton mixing matrix, $\hat{m}_\nu$ is 3×3 diagonal light neutrino mass matrix with eigenvalues $0, m_2$, and m_3 ; $\hat{M}_N$ is 3×3 diagonal RHN mass matrix with eigenvalues $0, M_2$, and M_3 ; R is an arbitrary rotation matrix. The CP asymmetry parameters from the decay of $N_{2,3} \rightarrow LH$ arises via the interference of the tree and one loop diagrams, and this is given by [59,60]

$$\epsilon_i = -\frac{1}{8\pi} \frac{1}{(y_{NI}^\dagger y_{NI})_{ii}} \sum_{j \neq i} \text{Im}\{(y_{NI}^\dagger y_{NI})_{ij}^2\} \times \left[f_v \left(\frac{M_j^2}{M_i^2} \right) + f_s \left(\frac{M_j^2}{M_i^2} \right) \right], \quad (18)$$

where the $f_v(x), f_s(x)$ denote the one loop vertex and self-energy corrections given as

$$f_v(x) = \sqrt{x} \left[1 - (1+x) \ln \left(\frac{x}{1+x} \right) \right]; f_s(x) = \frac{\sqrt{x}}{1-x}. \quad (19)$$

In a hierarchical limit, the CP asymmetry parameters can be estimated from Eq. (18) as

$$\epsilon_i = -\frac{3}{8\pi v_h^2} \sum_{j \neq i} \frac{M_i}{M_j} \frac{\text{Im}[(m_D^\dagger m_D)_{ij}]^2}{(m_D^\dagger m_D)_{ii}}. \quad (20)$$

The CP asymmetry (ϵ_S) generated by the decay of S , which comes from the interference of the tree and one loop diagrams, is shown in Fig. 2 and can be expressed as

$$\epsilon_S = -\frac{3}{8\pi v_h^2} \frac{M_S}{M_3} \frac{\text{Im}[(m_D^\dagger m_D)_{23}]^2}{(m_D^\dagger m_D)_{22}}. \quad (21)$$

Thus the CP asymmetry parameters of S and N_2 can be related as

$$\epsilon_S = (M_S/M_2)\epsilon_2. \quad (22)$$

From Eq. (22), we see that M_S cannot be arbitrarily small in order to produce correct lepton asymmetry via the decay of S through $S - N_2$ mixing. Above electroweak phase transition, the net lepton asymmetry produced by N_2, N_3 , and S can be converted to the observed baryon asymmetry via nonperturbative sphaleron processes. The resulting baryon asymmetry can be given as

$$\begin{aligned} \eta_B(z \rightarrow \infty) &= \frac{C_{L \rightarrow B}}{f} \sum_i Y_{\Delta L}(z \rightarrow \infty) \\ &= \frac{C_{L \rightarrow B}}{f} \sum_i \epsilon_i \kappa_i(z \rightarrow \infty) Y_i^{\text{eq}}(z \rightarrow 0) \\ &= -0.01394 \sum_i \epsilon_i \kappa_i(z \rightarrow \infty), \end{aligned} \quad (23)$$

where $i \in \{N_2, N_3, S\}$, $z = M_S/T$, κ_i is the efficiency factor, Y_i^{eq} is the equilibrium abundance of i th particle defined as $Y_i^{\text{eq}} = n_i^{\text{eq}}/n_\gamma$, n_i^{eq} is the equilibrium number density of i th species, n_γ is the photon number density, $f = 29.54^2$ is the dilution factor calculated assuming standard photon production from the onset of leptogenesis till recombination, $C_{L \rightarrow B}$ is the lepton to baryon asymmetry conversion factor and in our case the value of it is $C_{L \rightarrow B} = -0.54902$.³ We note that $Y_i^{\text{eq}}(z \rightarrow 0) = 3/4$. The required value of the baryon asymmetry is $\eta_B = (6.1 \pm 0.3) \times 10^{-10}$ [69], which translates to the required lepton asymmetry as $\sum_i Y_{\Delta L} \equiv \sum_i \epsilon_i \kappa_i Y_i^{\text{eq}}(z \rightarrow 0) \sim \mathcal{O}(10^{-8})$. As discussed above, the lepton asymmetry production due to S decay is via its mixing with N_2 . Therefore, if the out-of-equilibrium decay of S happens at late epoch, the washout processes will be no more active. In that case, the efficiency factor, κ_S will be ~ 1 . From Eq. (23), we see that S will contribute to the lepton asymmetry significantly when the contributions from N_2, N_3 are negligible. This can be realized only if $\epsilon_2 \kappa_2, \epsilon_3 \kappa_3 \ll 10^{-8}$ [70,71]. In that case, Eq. (23) will be reduced to

$$\eta_B = -0.01394 \epsilon_S \kappa_S. \quad (24)$$

² $g_s^* = 106.75 + \frac{7}{8} \times 2 + \frac{7}{8} \times (4 + 4) = 115.5$ is the relativistic d.o.f at the onset of leptogenesis, and $g_0^* = 3.91$ is the relativistic d.o.f today. The dilution factor f is then, $f = g_s^*/g_0^* = 29.54$.

³ $C_{L \rightarrow B} = -\frac{4m+8n}{9m+14n}$. For the steps of the calculation, refer to [67,68].

TABLE II. Benchmark points for leptogenesis.

BPs	M_S (GeV)	M_2/M_S	M_3/M_2	z_a	$(y_{Nl}^\dagger y_{Nl})_{22}$	y_{NS}	μ_ρ (GeV)	T_C (GeV)	λ_ρ
$BP1$	2×10^9	50.93	322.98	$0 - i0.542$	1.65×10^{-4}	0.35	10^9	2.62×10^9	1.5×10^{-5}
$BP2$	2×10^9	13.19	300.06	$6.42 \times 10^{-5} - i0.539$	4.21×10^{-5}	0.13	10^9	2.85×10^9	1.5×10^{-5}
$BP3$	4.1×10^8	18.84	13.27	$0 - i0.568$	1.43×10^{-5}	0.003	10^6	3.47×10^6	10^{-3}
$BP4$	2.8×10^6	2775.5	22.3	$0.402 - i1.955$	3.309×10^{-2}	0.015	10^6	3.47×10^6	10^{-3}

Thus, from Eqs. (18) and (24), we get a lower limit on the M_S when the $\kappa_S \sim 1$ (i.e., S is in weak washout regime, which is true due to the late decay of S when all the washouts processes are suppressed) to be

$$M_S \gtrsim 2.15 \times 10^6 \text{ GeV} \left(\frac{\eta_B}{6 \times 10^{-10}} \right) \left(\frac{2 \times 10^{-14} \text{ GeV}^{-1}}{\epsilon_2/M_2} \right), \quad (25)$$

where the conservative upper bound on ϵ_2/M_2 is obtained by choosing $\epsilon_2 \kappa_2 \leq 10^{-9}$. The details are given in Fig. 6.

To see the impact of N_2 and N_3 on the final lepton asymmetry, we solve the complete set of BEs for the evolution of lepton asymmetry as well as the abundances of S , N_2 , and N_3 which are given as

$$\frac{dY_{N_2}}{dz} = -\frac{\Gamma_2^D}{\mathcal{H}z} (Y_{N_2} - Y_{N_2}^{\text{eq}}) - \frac{\Gamma_1'}{\mathcal{H}z} (Y_{N_2} - Y_{N_2}^{\text{eq}}), \quad (26)$$

$$\frac{dY_{N_3}}{dz} = -\frac{\Gamma_3^D}{\mathcal{H}z} (Y_{N_3} - Y_{N_3}^{\text{eq}}) - \frac{\Gamma_1''}{\mathcal{H}z} (Y_{N_3} - Y_{N_3}^{\text{eq}}), \quad (27)$$

$$\begin{aligned} \frac{dY_S}{dz} = & -\frac{\Gamma_S^D}{\mathcal{H}z} (Y_S - Y_S^{\text{eq}}) - \frac{(\Gamma_1 + \Gamma_0)}{\mathcal{H}z} (Y_S - Y_S^{\text{eq}}) \\ & - \frac{\Gamma_0'}{\mathcal{H}z} \frac{(Y_S^2 - (Y_S^{\text{eq}})^2)}{Y_S^{\text{eq}}}, \end{aligned} \quad (28)$$

$$\begin{aligned} \frac{dY_{\Delta L}^{\text{all}}}{dz} = & \epsilon_S \frac{\Gamma_S^D}{\mathcal{H}z} (Y_S - Y_S^{\text{eq}}) + \epsilon_2 \frac{\Gamma_2^D}{\mathcal{H}z} (Y_{N_2} - Y_{N_2}^{\text{eq}}) \\ & + \epsilon_3 \frac{\Gamma_3^D}{\mathcal{H}z} (Y_{N_3} - Y_{N_3}^{\text{eq}}) \\ & - \left(\frac{1}{2} \frac{\Gamma_S^{ID} + \Gamma_{N_2}^{ID} + \Gamma_{N_3}^{ID}}{\mathcal{H}z} + \frac{\Gamma_2^W}{\mathcal{H}z} + \frac{\Gamma_S^{W_1} + \Gamma_{N_2}^{W_1} + \Gamma_{N_3}^{W_1}}{\mathcal{H}z} \right) Y_{\Delta L}, \end{aligned} \quad (29)$$

where $z = M_S/T$, Y_x is the abundance of x species defined as $Y_x = n_x/n_\gamma$, n_x is the number density of x , and $\mathcal{H}$ is the Hubble parameter. The evolution of N_2 , N_3 , and S are calculated by solving Eqs. (26), (27), (28). Equation (29) gives a net lepton asymmetry generated via the decay of N_2 , N_3 , and S . In the above equations, $\Gamma_2^D, \Gamma_3^D, \Gamma_S^D$ represent the thermal averaged decay width of N_2 , N_3 , and S , respectively. The inverse decay terms due to N_2, N_3, S are

$\Gamma_{N_2}^{ID}, \Gamma_{N_3}^{ID}, \Gamma_S^{ID}$. The $\Delta L = 1$ scattering interactions due to N_2, N_3 , and S are encoded in Γ_1', Γ_1'' , and Γ_1 , whereas the washout due to $\Delta L = 1$ scatterings are given by $\Gamma_{N_2}^{W_1}, \Gamma_{N_3}^{W_1}, \Gamma_S^{W_1}$ for N_2, N_3, S in Eq. (29). Γ_2^W represents the $\Delta L = 2$ washout processes. Here, Γ_0 denotes the $\Delta L = 0$ scattering processes like $S\rho \rightarrow \rho S$ in the t and s channel, and Γ_0' represents the $\Delta L = 0$ scattering processes where two S are in the initial state like $S\bar{S} \rightarrow \rho\rho, HH^\dagger, L\bar{L}$. These lepton number conserving processes bring the S into thermal equilibrium depending on the strength of the coupling y_{NS} . We also calculate lepton asymmetry considering only S contribution by solving Eq. (28) and

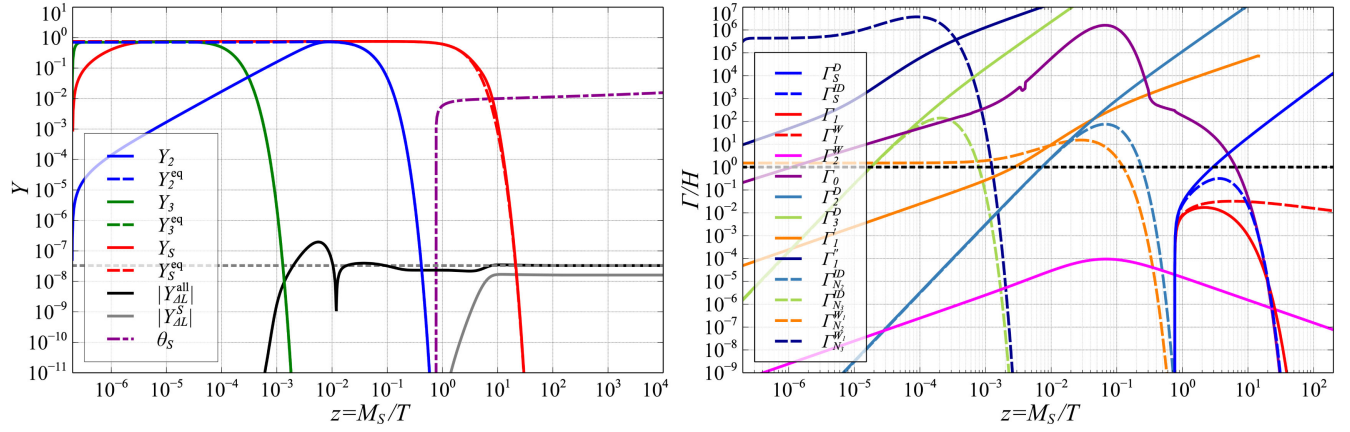
$$\frac{dY_{\Delta L}^S}{dz} = \epsilon_S \frac{\Gamma_S^D}{\mathcal{H}z} (Y_S - Y_S^{\text{eq}}) - \left(\frac{1}{2} \frac{\Gamma_{ID}}{\mathcal{H}z} + \frac{\Gamma_1^W + \Gamma_2^W}{\mathcal{H}z} \right) Y_{\Delta L}. \quad (30)$$

We solve the above equations simultaneously for four sets of benchmark points⁴ as mentioned in Table II to get the lepton asymmetry. Here, to evaluate the Yukawa couplings, we used the Casas-Ibarra (CI) parametrization as mentioned in Eq. (17).⁵

BP1: We first calculate the lepton asymmetry for $BP1$ from Table II, considering only the effect of S , i.e., by solving Eqs. (28) and (30), as shown in the left panel of Fig. 3. For the chosen parameters, we find the critical temperature to be $T_C = 2.62 \times 10^9$ GeV; i.e., above 2.62×10^9 GeV, the vev of ρ is zero. This implies that the mixing is absent for $z \lesssim 0.76$. In this regime, S reaches equilibrium via $\Delta L = 0$ scatterings that do not involve mixing. For $z > 0.76$, S starts to decay to L, H through mixing with N_2 . It should be noted that the mixing angle, $\theta_S(T)$ evolves with time, as evident from Eq. (16). The evolution of $\theta_S(T)$ is shown with a magenta dashed-dotted line. For $BP1$, the CP asymmetry parameters are $\epsilon_S = 1.97 \times 10^{-7}$, $\epsilon_2 = 10^{-5}$, and $\epsilon_3 = 1.07 \times 10^{-7}$. The resulting lepton asymmetry from the decay of S is calculated to be $|Y_{\Delta L}^S| = 1.61 \times 10^{-8}$, shown by the gray solid line. The gray dashed line represents the observed lepton asymmetry.

⁴We fix $\lambda_{H\rho} = 0.5$ ensuring the stability of the potential.

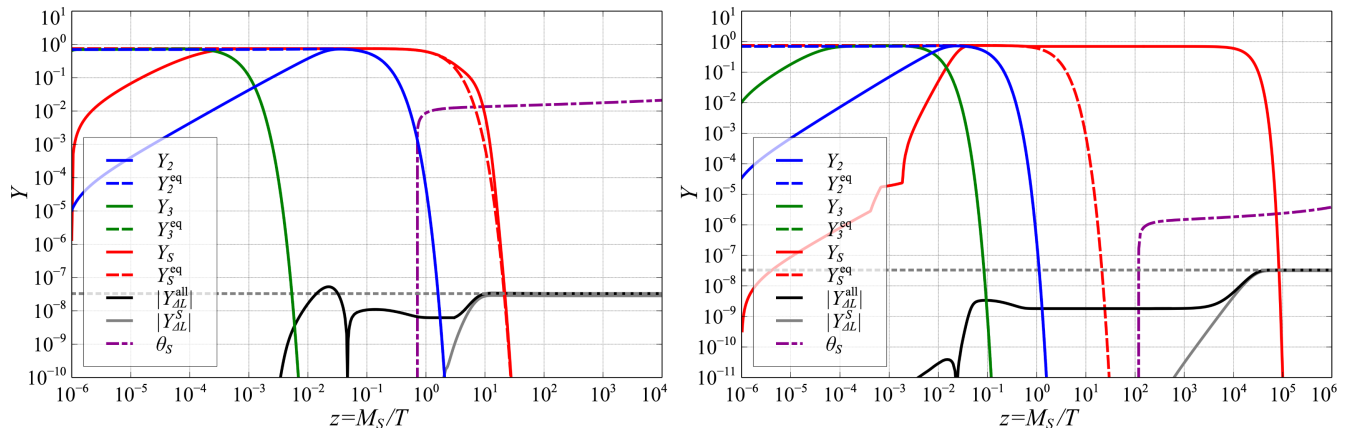
⁵For evaluating the Yukawa matrix using the CI parametrization, we use the best fit values of the neutrino oscillation parameters [72]. In Eq. (17), we use the R matrix as $R_{12}R_{23}R_{13}$ with rotation angle z_a .



We now include the effects of N_2 and N_3 by solving Eqs. (26), (27), (28), and (29) for the same BPI . The final asymmetry is shown by the black solid line. We observe that N_2, N_3 produce an asymmetry of $|Y_{\Delta L}| \sim 2.33 \times 10^{-8}$, which acts as an initial condition for the lepton asymmetry production via the decay of S . The final asymmetry after S fully decays is $|Y_{\Delta L}^{\text{all}}| \sim 3.30 \times 10^{-8}$. The percentage increase in the final asymmetry compared to the S -only contribution is 104.94%. In this case, $\epsilon_2 \kappa_2 \sim 3 \times 10^{-8}$ and $\epsilon_3 \kappa_3 \ll \epsilon_2 \kappa_2$. This can be interpreted as an N_2 -dominated leptogenesis scenario, as the asymmetry produced by N_3 is entirely washed out by N_2 , and S alone fails to generate the correct lepton asymmetry—clearly visible from the *left* panel of Fig. 3. In the *right* panel of Fig. 3, the interaction rates of all relevant processes are shown with respect to the Hubble parameter. The different colored lines represent

various interaction rates involved in leptogenesis, as indicated in the figure's inset.

BP2: We then move to *BP2*, where we set the mass ratios as $M_2/M_S = 13.19$ and $M_3/M_2 = 3957.09$ and choose the Yukawa coupling value as $y_{NS} = 0.13$. The *CP* asymmetry parameters for this *BP2* are $\epsilon_S = 1.95 \times 10^{-7}$, $\epsilon_2 = 2.57 \times 10^{-6}$, and $\epsilon_3 = 2.86 \times 10^{-8}$. The mixing becomes nonzero after $z \sim 0.7$. The lepton asymmetry is generated from the decay of *S*, and the final asymmetry saturates at $|Y_{\Delta L}^S| \sim 2.84 \times 10^{-8}$. The evolution of the asymmetry is shown in the *left* panel of Fig. 4 with a gray solid line. After including the contributions from N_2 , N_3 , the asymmetry once N_2 has completely decayed is 6.25×10^{-9} , which acts as a initial condition for the lepton asymmetry produce via *S* decay. The total final asymmetry after the decay of *S* is $|Y_{\Delta L}^{\text{all}}| \sim 3.32 \times 10^{-8}$. This corresponds to a change of only



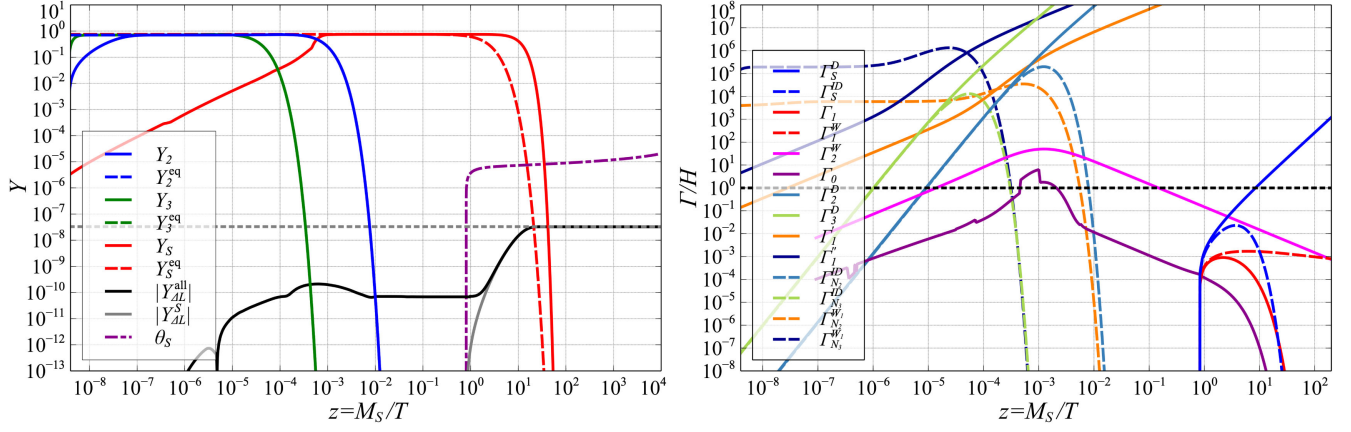


FIG. 5. Left: Cosmological evolution of the lepton asymmetry and abundances for *BP4* from Table II. Right: The comparison of interaction rates of all the processes involved are shown with respect to the Hubble parameter for the same *BP4*.

16.97% compared to the *S*-only contribution. Note that in this case, $\epsilon_2 \kappa_2 \sim 8.33 \times 10^{-9}$ and $\epsilon_3 \kappa_3 \ll \epsilon_2 \kappa_2$.

BP3: In the *right* panel of Fig. 4, we present the evolution of the asymmetries corresponding to *BP3* from Table II. The associated *CP* asymmetry parameters are $\epsilon_S = 4.49 \times 10^{-8}$, $\epsilon_2 = 8.45 \times 10^{-7}$, and $\epsilon_3 = 7.86 \times 10^{-8}$. For this parameter choice, the critical temperature is found to be 3.47×10^6 GeV. Initially, *S* reaches equilibrium through lepton number-conserving scatterings. Beyond $z \sim 119$, mixing with N_2 sets in, enabling *S* to decay into *LH* and generate a lepton asymmetry, as depicted by the gray solid line. This yields a final asymmetry from *S* decay alone to be $|Y_{\Delta L}^S| \sim 3.14 \times 10^{-8}$. To include the full dynamics, we next incorporate the effects of N_2 and N_3 along with *S*, with results shown by the black solid line. The asymmetry left behind after the complete decay of N_2 is 1.8×10^{-9} , with $\epsilon_2 \kappa_2 \sim 2.4 \times 10^{-9}$. Subsequently, as *S* decays (starting near $z \sim 119$), additional asymmetry is generated via *S* – N_2 mixing. The total final lepton asymmetry from the decays of N_2 , N_3 and *S* is estimated to be $|Y_{\Delta L}^{\text{all}}| \sim 3.32 \times 10^{-8}$. Compared to the *S*-only contribution, this gives rise an enhancement of about 5.7%.

BP4: In Fig. 5, we compute the lepton asymmetry for the *BP4*, as listed in Table II. The *CP* asymmetry parameters for this case are $\epsilon_S = 4.46 \times 10^{-8}$, $\epsilon_2 = 1.24 \times 10^{-4}$, and $\epsilon_3 = 2.39 \times 10^{-4}$. The critical temperature is found to be 3.47×10^6 GeV. Consequently, *S* begins decaying to *LH* around $z \sim 0.81$, generating a lepton asymmetry of $|Y_{\Delta L}^S| \sim 3.25 \times 10^{-8}$. When the contributions from N_2 and N_3 are included, the final asymmetry increases by only 0.18%, indicating that *S* dominates the generation of lepton asymmetry in this scenario. A comparison of the relevant interaction rates is shown in the *right* panel of Fig. 5.

We now show the parameter space where *S* is the sole degree of freedom contributing to the lepton asymmetry, in the $\epsilon_S \kappa_S - M_S$ plane, as illustrated in Fig. 6. The free parameters are scanned over the ranges:

$M_2 \in [2.4 \times 10^9, 10^{14}]$ GeV, $M_3/M_2 \in [10, 10^4]$, and $z_a \in [-\pi, \pi] + i[-\pi, \pi]$, subject to the condition $M_{2,3} > M_S$. We identify regions in the parameter space where the contributions from N_2 and N_3 are negligible. This condition is satisfied when $\epsilon_2 \kappa_2, \epsilon_3 \kappa_3 \ll 10^{-8}$. These points are then plotted in the $\epsilon_S \kappa_S - M_S$ plane using the Eq. (18). Each line in Fig. 6 (distinguished by color shading) corresponds to a specific set of parameters $\{z_a, \epsilon_2, \epsilon_3, M_2, M_3\}$ for which the lepton asymmetry from N_2 and N_3 remains negligible. The horizontal dashed lines indicate the observed baryon asymmetry, which translates into a required value of the *CP* asymmetry times efficiency factor as $\epsilon_S \kappa_S = (4.38 \pm 0.22) \times 10^{-8}$. By appropriately choosing the other parameters $y_{NS}, M_S, \lambda_\rho, \mu_\rho, \lambda_{H\rho}$, this asymmetry can be reproduced solely through the contribution from *S*. To better understand the regime of *S* dominance, we take one

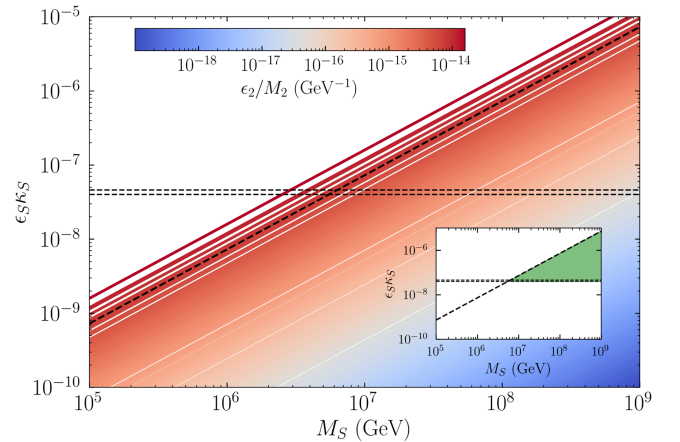


FIG. 6. Values of $\epsilon_S \kappa_S$ as a function of M_S . Each of the lines corresponds to a set of parameters when $\epsilon_2 \kappa_2 \leq 10^{-9}$. The color code represents the values of ϵ_2/M_2 . The dashed horizontal lines represent the observed baryon asymmetry, which translates to $\epsilon_S \kappa_S = (4.38 \pm 0.22) \times 10^{-8}$. The lowest allowed value of M_S corresponds to $\epsilon_2/M_2 \sim 2 \times 10^{-14}$ GeV $^{-1}$.

set of parameters as $\{z_a = -3.05 \times 10^{-6} - i1.91, \epsilon_2 = 2.17 \times 10^{-4}, \epsilon_3 = 1.39 \times 10^{-4}, M_2 = 2.979 \times 10^{10} \text{ GeV}, M_3 = 2.51 \times 10^{12} \text{ GeV}\}$ for which $\epsilon_2/M_2 = 7.27 \times 10^{-15}$, which is shown with the black dashed line. Now for this set of parameters, the green-shaded triangular region, in the plane of $\epsilon_S \kappa_S - M_S$, will be the S dominated regime, which is shown in the inset figure. In other words, in the green shaded region, the contribution from N_2, N_3 are negligible, and the final asymmetry will be produced solely due to the decay of S .

From the plot, we observe that the minimum value of M_S that can account for the observed baryon asymmetry is $M_S \simeq 2.8 \times 10^6 \text{ GeV}$. For $M_S < 2.8 \times 10^6 \text{ GeV}$, values of ϵ_S alone are insufficient to generate the correct lepton asymmetry. Notably, this lower bound on M_S is 3 orders of magnitude below the Davidson–Ibarra (DI) bound for the scale of thermal leptogenesis, consistent with Eq. (25).

IV. DOMAIN WALLS AND SIGNATURES OF GRAVITATIONAL WAVES FROM Z_2 SYMMETRY BREAKING

The spontaneous breaking of Z_2 symmetry giving rise to lepton asymmetry also leads to formation of domain walls (DWs) in the early Universe. The energy density of the DWs falls with the cosmological scale factor as $\sim R^{-1}$, which is much slower than the matter ($\sim R^{-3}$) and radiation ($\sim R^{-4}$). Thus, the DWs may over close the Universe if they are stable. This problem can be solved by making the DW unstable, and it will disappear eventually in the early Universe.

We demonstrate by considering the potential for the scalar fields as

$$V(\rho) = -\frac{\mu_\rho^2}{2}\rho^2 + \frac{\lambda_\rho}{4}\rho^4 + \frac{cT^2}{2}\rho^2 = \frac{\lambda_\rho}{4}(\rho^2 - v_\rho(T)^2)^2. \quad (31)$$

The potential has two degenerate minima at $\langle \rho \rangle = \pm v_\rho(T)$. The field can occupy any one of the two minima after the symmetry breaking resulting in two different domains, separated by a wall. We consider a static planar DW perpendicular to the x axis in the Minkowski space, $\rho = \rho(x)$.

The equation of motion for the DW is [5,8,9,73,74]

$$\frac{d^2\rho}{dx^2} - \frac{dV}{d\rho} = 0, \quad (32)$$

with the boundary condition

$$\lim_{x \rightarrow \pm\infty} \rho(x) = \pm v_\rho. \quad (33)$$

After solving the equation of motion, we obtain

$$\rho(x) = v_\rho \tanh(\alpha x), \quad (34)$$

where $\alpha \simeq \sqrt{\frac{\lambda_\rho}{2}} v_\rho$.

The DW is extended along $x = 0$ plane, and the two vacua are realized at $x \rightarrow \pm\infty$. The width of the DW is estimated as $\delta \sim \left(\frac{\sqrt{\lambda_\rho} v_\rho}{\sqrt{2}}\right)^{-1}$. The surface energy density, also referred as tension of the DWs, is calculated to be

$$\sigma(T) = \frac{4}{3} \sqrt{\frac{\lambda_\rho}{2}} v_\rho(T)^3 \simeq \frac{2}{3} M_\rho(T) v_\rho(T)^2, \quad (35)$$

where $M_\rho(T) = \sqrt{2\lambda_\rho} v_\rho(T)$.

As discussed earlier without a soft Z_2 breaking term, the DW will be stable and will over close the energy density of the Universe. In order to over come this problem, we introduce an energy bias in the potential as $\mu_b^3 \rho / \sqrt{2}$, which breaks the Z_2 symmetry explicitly. Here, μ_b is a mass dimension one coupling. Equation (31) then becomes

$$\mathcal{V} = V(\rho) + \mu_b^3 \rho / \sqrt{2}. \quad (36)$$

As a result, the degeneracy of the minima is lifted by

$$V_{\text{bias}} \equiv |\mathcal{V}(-v_\rho) - \mathcal{V}(v_\rho)| = \sqrt{2} \mu_b^3 v_\rho(T). \quad (37)$$

This creates a pressure difference across the wall [8,9,73]. We assume the annihilation happens in the radiation dominated era. The energy bias has to be large enough, so that the DW can disappear before the BBN epoch, i.e., $t_{\text{ann}} < t_{\text{BBN}}$, where

$$t_{\text{ann}} = C_{\text{ann}} \frac{\mathcal{A} \sigma(T_{\text{ann}})}{V_{\text{bias}}}, \quad (38)$$

where C_{ann} is a coefficient of $\mathcal{O}(1)$, $\mathcal{A} \simeq 0.8 \pm 0.1$ [13] is area parameter, and t_{BBN} is the BBN timescale. This gives a lower bound on the V_{bias} as

$$V_{\text{bias}} > 6.58 \times 10^{-14} \text{ GeV}^4 C_{\text{ann}} \mathcal{A} \left(\frac{10^{-2} \text{ sec}}{t_{\text{BBN}}} \right) \left(\frac{\sigma}{1 \text{ TeV}^3} \right). \quad (39)$$

Equation (39) can be written in terms of the Z_2 breaking parameters μ_b as

$$\mu_b > 1.45839 \times 10^{-4} \text{ GeV} C_{\text{ann}}^{\frac{1}{3}} \mathcal{A}^{\frac{1}{3}} \left(\frac{10^{-2} \text{ sec}}{t_{\text{BBN}}} \right)^{\frac{1}{3}} \times \left(\frac{M_\rho}{1 \text{ TeV}} \right)^{\frac{1}{3}} \left(\frac{v_\rho}{10^5 \text{ TeV}} \right)^{\frac{1}{3}}. \quad (40)$$

The DWs have to disappear before they could start to dominate the energy density of the Universe, i.e. $t_{\text{ann}} < t_{\text{dom}}$, where

$$t_{\text{dom}} = \frac{3}{32\pi} \frac{M_{\text{pl}}^2}{\mathcal{A}\sigma}. \quad (41)$$

This puts a lower bound on the annihilation temperature as

$$T_{\text{ann}} > 1.34772 \text{ GeV} \mathcal{A}^{1/2} \left(\frac{g_*(T_{\text{ann}})}{10} \right)^{-\frac{1}{4}} \left(\frac{v_\rho}{10^5 \text{ TeV}} \right) \times \left(\frac{M_\rho}{1 \text{ TeV}} \right)^{\frac{1}{2}}. \quad (42)$$

In terms of μ_b , it can be expressed as

$$\mu_b > 4.1404 \times 10^{-3} \text{ GeV} \mathcal{C}_{\text{ann}}^{\frac{1}{3}} \mathcal{A}^{\frac{2}{3}} \left(\frac{M_\rho}{1 \text{ TeV}} \right)^{\frac{2}{3}} \left(\frac{v_\rho}{10^5 \text{ TeV}} \right). \quad (43)$$

The DWs can then annihilate and emit their energy in the form of stochastic gravitational waves (GWs), which can be detectable at present time.

The peak amplitude of the GW spectrum at the present time, t_0 , is given by [73]

$$\Omega_{\text{GW}} h^2(t_0)|_{\text{peak}} = 7.18824 \times 10^{-18} \mathcal{A}^2 \tilde{\epsilon}_{\text{GW}} \left(\frac{\sigma}{1 \text{ TeV}^3} \right)^2 \times \left(\frac{g_{*s}(T_{\text{ann}})}{10} \right)^{-\frac{4}{3}} \left(\frac{T_{\text{ann}}}{10^{-2} \text{ GeV}} \right)^{-4}, \quad (44)$$

where $\tilde{\epsilon}_{\text{GW}} \simeq 0.7 \pm 0.4$ [13] is the efficiency parameter, T_{ann} is the temperature at which the DWs annihilate, and

$g_{*s}(T_{\text{ann}})$ is the relativistic entropy degrees of freedom at the epoch of DWs annihilation.

From Eq. (44), we see that the peak of GW spectrum is directly proportional to σ^2 and inversely proportional to T_{ann}^4 . In our setup, even though DW are produced before EW phase transition, they can sustain until late epoch to give rise larger peak amplitude of the GW spectrum. If the DW annihilation is happening at an earlier time, i.e., at a large temperature, $T_1 > T_2$, then the amplitude of the GW will be larger for T_2 as compared to T_1 .

Assuming the DW disappear at temperature T_{ann} , the peak frequency of the GW spectrum at present time is estimated as

$$f_{\text{peak}}(t_0) = 1.78648 \times 10^{-10} \text{ Hz} \left(\frac{g_{*s}(T_{\text{ann}})}{10} \right)^{-\frac{1}{3}} \times \left(\frac{g_*(T_{\text{ann}})}{10} \right)^{\frac{1}{2}} \left(\frac{T_{\text{ann}}}{10^{-2} \text{ GeV}} \right). \quad (45)$$

The amplitude of the GW for any frequency at the present time varies as

$$\Omega_{\text{GW}} h^2(t_0, f) = \Omega_{\text{GW}} h^2(t_0)|_{\text{peak}} \begin{cases} \frac{f_{\text{peak}}}{f} & f > f_{\text{peak}} \\ \left(\frac{f}{f_{\text{peak}}} \right)^3 & f < f_{\text{peak}} \end{cases} \quad (46)$$

It is worth mentioning that v_ρ is the only parameter that is sensitive to both the leptogenesis and GW spectrum. In Fig. 7, we have illustrated the GW spectrum for two benchmarks as mentioned in Table III. We have shown different sensitivities from experiments BBO [75], CE,

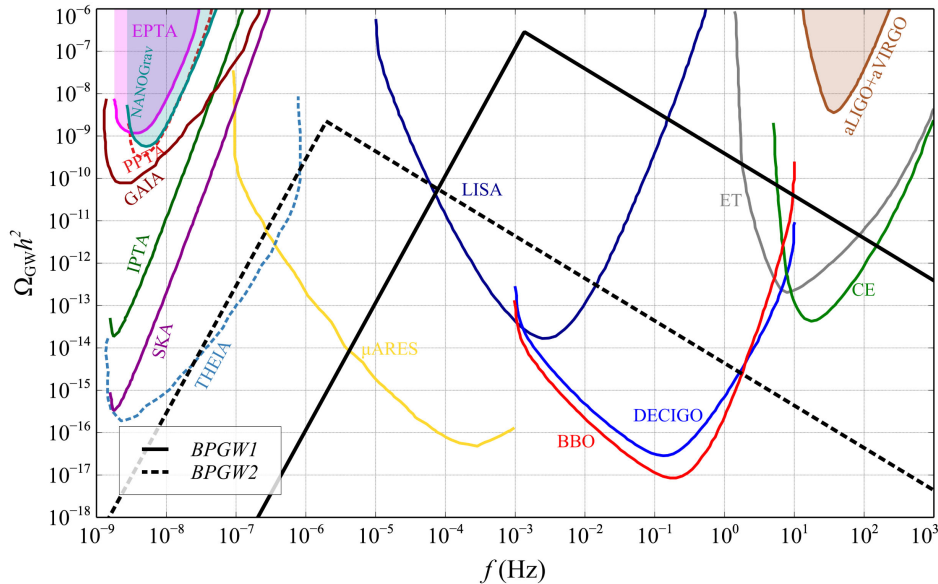


FIG. 7. Gravitational wave spectrum from the annihilating domain walls for four benchmark points as mentioned in Table III. Sensitivities from different gravitational wave search experiments have been shown with different colors.

TABLE III. Benchmark points for gravitational wave.

BPs	M_S (GeV)	M_2/M_S	T_C (GeV)	T_{ann} (GeV)	$\lambda_\rho(T_{\text{ann}})$	$M_\rho(T_{\text{ann}})$ (GeV)	$v_\rho(T_{\text{ann}})$ (GeV)	$\sigma(\text{TeV}^3)$
<i>BPGW1</i>	2×10^9	13.19	2.85×10^9	5.1×10^4	2.53×10^{-2}	1.42×10^9	6.31×10^9	3.76×10^{19}
<i>BPGW2</i>	2.8×10^6	277.5	3.47×10^6	75	1.36×10^{-4}	1.41×10^6	8.56×10^7	6.91×10^{12}

DECIGO [76], NANOGrav [1,2], EPTA [3], CPTA [77], PPTA [4], ET [78], GAIA [79], IPTA [80], LISA, SKA [81], THEIA [79], aLIGO [82], aVIRGO, and μ ARES [83]. For *BPGW1*, which corresponds to *BP2* in Table II, the lower bound on the domain wall (DW) annihilation temperature is given by Eq. (42) as 5.0171×10^4 GeV. We choose 5.1×10^4 GeV as the annihilation temperature for this case, leading to a surface energy density of 3.76×10^{19} TeV^3 and a GW peak frequency of ~ 0.0013 Hz. This benchmark lies inside the sensitivity range of space-based interferometers such as LISA, DECIGO, as well as μ ARES, ET, and CE. Moving to *BPGW2* (which is same as *BP4* in Table II), the lower bound on the DW annihilation temperature is found to be 72 GeV. We take 75 GeV as the annihilation temperature, resulting in a surface energy density of 6.91×10^{12} TeV^3 . Although this energy density is much lower than that of *BPGW1*, the reduced annihilation temperature compensates a little for the decrease in GW amplitude, as described by Eq. (44). The peak GW frequency is 1.9×10^{-6} Hz, placing this benchmark inside the sensitivity range of experiments such as THEIA, μ ARES, LISA, BBO, and DECIGO. At this juncture, we note that the recent study [84] has pointed out that BBN can be sensitive to domain wall annihilation occurring near the BBN epoch (MeV scale). However, in our scenario, the domain wall annihilation takes place at

$T_{\text{ann}} = 75$ GeV, which is well before the onset of BBN, making it safe from such constraints.

Figure 8 illustrates the sensitivity reaches of various experiments in the v_ρ versus T_{ann} plane corresponding to *BPGW1* in Table III. Here, we treat the vev of ρ as an independent parameter. The parameter space is constrained by two key requirements: (i) domain walls (DWs) must annihilate before big bang nucleosynthesis, and (ii) they must do so before dominating the energy density of the Universe.

V. DARK MATTER PHENOMENOLOGY

We turn to comment on DM in our setup. Due to the unbroken Z'_2 symmetry, combination of N_1 and $\Psi = (\psi^0 \psi^-)^T \equiv (\psi_L^0 + \psi_R^0 \psi^-)^T$ can give rise to a singlet-doublet Majorana DM [41]. The relevant DM Lagrangian reads as

$$\mathcal{L}_{\text{DM}} = \bar{\Psi} i \gamma^\mu \mathcal{D}_\mu \Psi - M \bar{\Psi} \Psi + \bar{N}_1 i \gamma^\mu \partial_\mu N_1 - \frac{1}{2} M_1 \bar{N}_1^c N_1 - \frac{y_1}{\sqrt{2}} \bar{\Psi} \tilde{H} (N_1 + N_1^c) + \text{H.c.} \quad (47)$$

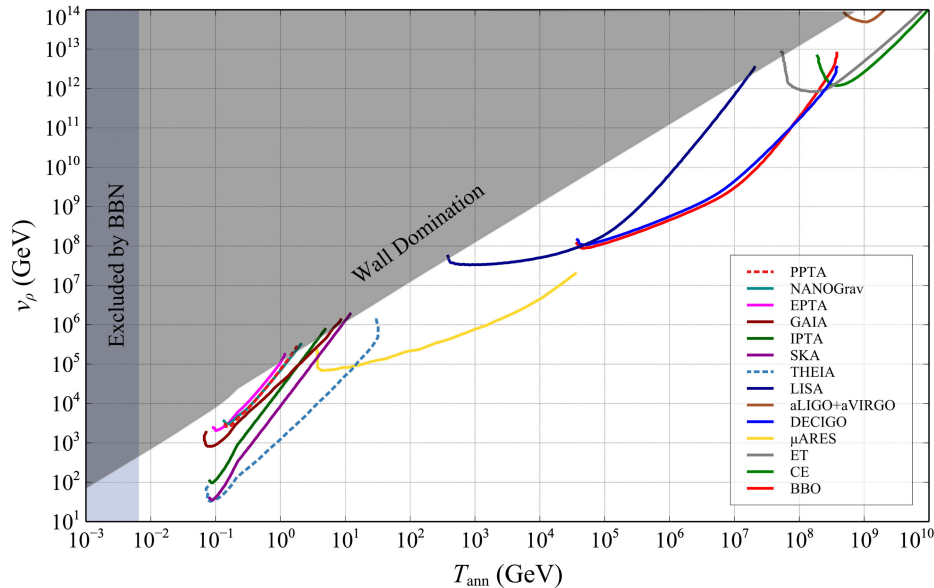


FIG. 8. The sensitivity of the gravitational wave search experiments in the plane of $v_\rho - T_{\text{ann}}$. These contours are obtained for peak frequency and peak amplitude of the GW as per Eqs. (45) and (44), respectively.

The neutral fermion mass matrix can be written in the basis $((\psi_R^0)^c, \psi_L^0, (N_1)^c)^T$ as

$$\begin{pmatrix} 0 & M & \frac{m_D}{\sqrt{2}} \\ M & 0 & \frac{m_D}{\sqrt{2}} \\ \frac{m_D}{\sqrt{2}} & \frac{m_D}{\sqrt{2}} & M_1 \end{pmatrix}, \quad (48)$$

where $m_D = y_1 v_h / \sqrt{2}$. The mass matrix can be diagonalized with a unitary matrix of the form $U(\theta) = U_{13}(\theta_{13} = \theta) \cdot U_{23}(\theta_{23} = 0) \cdot U_{12}(\theta_{12} = \frac{\pi}{4})$. The three neutral states mix and gives three Majorana states as $\chi_i = \frac{\chi_{iL} + \chi_{iL}^c}{\sqrt{2}}$, where

$$\begin{aligned} \chi_{1L} &= \frac{\cos \theta}{\sqrt{2}} (\psi_L^0 + (\psi_R^0)^c) + \sin \theta N_1^c, \\ \chi_{2L} &= \frac{i}{\sqrt{2}} (\psi_L^0 - (\psi_R^0)^c), \\ \chi_{3L} &= -\frac{\sin \theta}{\sqrt{2}} (\psi_L^0 + (\psi_R^0)^c) + \cos \theta N_1^c. \end{aligned} \quad (49)$$

The corresponding mass eigenvalues are

$$\begin{aligned} M_{\chi_1} &= M \cos^2 \theta + M_1 \sin^2 \theta + m_D \sin 2\theta, \\ M_{\chi_2} &= M, \\ M_{\chi_3} &= M \sin^2 \theta + M_1 \cos^2 \theta - m_D \sin 2\theta, \end{aligned} \quad (50)$$

where the mixing angle is given as

$$\tan 2\theta = \frac{2m_D}{M - M_1}. \quad (51)$$

Here, we identify the χ_3 be DM candidate. The Yukawa coupling can be expressed as

$$y_1 = \frac{\Delta M \sin 2\theta}{\sqrt{2} v_h}, \quad (52)$$

where ΔM is the mass splitting between DM and the next heavy neutral fermion state. The free parameters in the DM phenomenology are $\{M_{\chi_3} \equiv M_{\text{DM}}, \Delta M = M_{\chi_1} - M_{\chi_3} \approx M_{\chi_2} - M_{\chi_3}, \sin \theta\}$.

The DM relic is decided by the freeze-out of various annihilation and coannihilation processes in the early Universe. We compute the relic density and the DM spin-independent cross-section using the micrOMEGAs package [85]. In Fig. 9 (left), we show the mass splitting between the DM and the light neutral fermion as a function of DM mass for correct relic density. As the DM mass increases, the annihilation cross-section decreases, and as a result, the relic increases. The coannihilation plays important role in bringing down the relic to correct ball park. The coannihilation is large when the mass splitting between the

DM and the next Z'_2 odd sector particle is smaller. This feature is clearly visible in Fig. 9 (left). When the mass splitting ΔM is large, coannihilation becomes negligible. The Higgs mediated annihilation processes decide the relic density mainly. Thus, the mixing angle $\sin \theta$ becomes important here as the Yukawa coupling is $\propto \Delta M \sin 2\theta$. The dependence of mixing angle for large mass splitting is prominent in Fig. 9 (left). For fixed $\sin \theta$, as the mass splitting increases, the coupling increases, which leads to an increase in annihilation cross-section, and thus, the relic decreases. To get the correct relic, the DM mass has to be large, which makes the cross-section smaller giving the correct relic. Larger $\sin \theta$ requires smaller ΔM to give the correct relic. We take these correct relic data points and impose the direct detection constraint from LZ experiment [86] and show in Fig. 9 (right). The direct detection is possible via the Higgs portal. The spin-independent DM-nucleon cross-section is $\sigma_{\text{DM-N}}^{\text{SI}} \propto y_1^2 \sin^2 2\theta \propto \Delta M^2 \sin^4 2\theta$. A larger mixing angle as well as larger mass splitting will result in larger direct detection cross-section. However, larger mass splitting is allowed for DM mass at Higgs resonance. We also impose the LEP bound on the mass of the charged component of the doublet, which is $M_{\Psi^+} > 102.7 \text{ GeV}$ [87].

VI. CONCLUSION

In this paper, we explored the potential for generating successful thermal leptogenesis at a scale lower than the Davidson Ibarra bound in an extended type-I seesaw framework along with nonzero neutrino mass, dark matter, and gravitational wave. In our setup, the lightest RHN N_1 mixes with the neutral component of a vectorlike fermion doublet Ψ , which are odd under an unbroken Z'_2 symmetry. As a result, we get singlet-doublet Majorana DM in a wide range of parameter space. We also added a singlet scalar ρ and a singlet fermion S in the canonical type-I seesaw, which are odd under an imposed discrete Z_2 symmetry. At high scale, typically above the EWPT, ρ acquires a vev and breaks the Z_2 symmetry spontaneously. As a result, we got S mixed up with N_2 such that late decay of S could give rise to a relatively low scale thermal leptogenesis. In particular, we saw that successful thermal leptogenesis requires $M_S \gtrsim 2.8 \times 10^6 \text{ GeV}$. We note that this is 3 orders magnitude smaller than the usual Davidson Ibarra bound ($M_N > 2.4 \times 10^9 \text{ GeV}$). The spontaneous breaking of the Z_2 symmetry also gave rise to DWs in the early Universe. We discussed the evolution of DWs, which annihilate by emitting stochastic GWs. In the appropriate parameter space, these GWs can produce detectable signatures at space-based interferometers such as LISA, UDECIGO, BBO as well as ET, CE, THEIA, SKA, μARES . In our scenario, the GW frequencies lie in the nHz to kHz range, making them relevant for high-frequency GW detectors as well as low-frequency detectors.

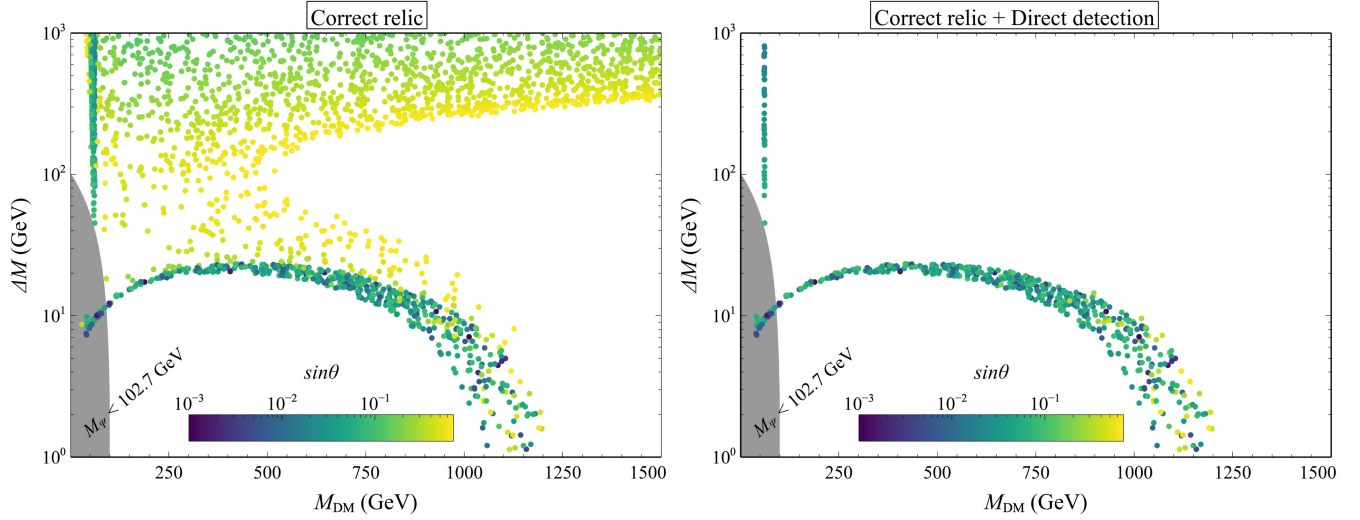


FIG. 9. Left: DM parameter space satisfying correct relic in the plane of ΔM vs M_{DM} . Right: Parameter space consistent with both relic density and direct search constraints (LZ) in the ΔM versus M_{DM} plane. The gray shaded region in the bottom left corner is ruled out by the LEP exclusion bound on charged fermion mass $M_{\Psi^+} \equiv M > 102.7$ GeV.

The scenario we discussed here can be implemented in a gauged $U(1)_{B-L}$ symmetric model where the connection between the DM, leptogenesis and GW is more coherent. In this case, the presence of three RHNs, $N_{1,2,3}$ is required to cancel the gauge anomaly. The minimal $U(1)_{B-L}$ model can be extended with a vectorlike singlet fermion S , a vectorlike doublet fermion Ψ , and one singlet scalar ρ . Additionally a Z_2 symmetry is imposed under which S and ρ are odd, while all other particles are even. The lightest RHN, N_1 along with Ψ , can be made a viable DM candidate by assigning another discrete symmetry Z'_2 under which they are odd and all the other particles are even. All the fermions ($N_{1,2,3}$, S) will get their masses from the $U(1)_{B-L}$ breaking scalar Φ . The leptogenesis will be generated by the decay of S , while the neutrino mass can be realized via the type-I seesaw mechanism anchored by $N_{2,3}$. In this scenario, there will be two sources of GWs: (i) the breaking of $U(1)_{B-L}$ symmetry and (ii) the annihilation of DWs that arise due to breaking of the Z_2 symmetry by the vev of ρ .

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: LOWER BOUND ON THE LEPTOGENESIS SCALE IN CANONICAL TYPE-I SEESAW

In the hierarchical scenario of the canonical type-I leptogenesis, the lightest RHN decays to L , H and $\bar{L}, H^\dagger$. The interference between these tree level and one loop processes can give rise to a CP asymmetry given as

$$\epsilon_{N_1} = -\frac{3}{8\pi v_h^2} M_1 \sum_{j=2,3} \frac{1}{M_j} \frac{\text{Im}[(m_D^\dagger m_D)_{1j}]^2}{(m_D^\dagger m_D)_{11}}, \quad (\text{A1})$$

where $v_h = 246$ GeV, is the SM Higgs vev.

The decay width of N_1 is calculated to be

$$\Gamma_{N_1} = \frac{(m_D^\dagger m_D)_{11}}{8\pi v_h^2} M_1. \quad (\text{A2})$$

A net lepton asymmetry can be generated once this decay rate falls below the Hubble expansion rate of the Universe,

$$\mathcal{H}(T) = 1.66\sqrt{g_*} \frac{T^2}{M_{\text{pl}}}, \quad (\text{A3})$$

where g_* is the effective number of relativistic degrees of freedom, and $M_{\text{pl}} = 1.22 \times 10^{19}$ GeV is the Planck mass.

The CP asymmetry is bounded from above as [88]

$$\epsilon_{N_1} \leq \frac{3}{8\pi v_h^2} M_1 \sqrt{\Delta m_{\text{atm}}^2}. \quad (\text{A4})$$

In the zero initial abundance case, the lower bound on the lightest RHN mass is found to be [89]

$$M_1 \gtrsim 2.4 \times 10^9 \text{ GeV.} \quad (\text{A5})$$

This is mainly because the same coupling is responsible to give neutrino mass as well as leptogenesis. Various attempts have been made to lower this leptogenesis scale,⁶ e.g., incorporating flavor effects [50], by adding extra scalar fields [51,52,54], resonant leptogenesis [55], and by decoupling neutrino mass and leptogenesis [59,60].

APPENDIX B: RENORMALIZATION GROUP EQUATIONS

The RG equations are given as

$$\beta_{g_1}^{(1)} = \frac{41}{10} g_1^3, \quad (\text{B1})$$

$$\beta_{g_2}^{(1)} = -\frac{19}{6} g_2^3, \quad (\text{B2})$$

$$\beta_{g_3}^{(1)} = -7g_3^3, \quad (\text{B3})$$

$$\beta_{\lambda_\rho}^{(1)} = 2(4\lambda_\rho y_{NS}^2 - 4y_{NS}^4 + 9\lambda_\rho^2 + \lambda_{H\rho}^2), \quad (\text{B4})$$

$$\begin{aligned} \beta_{\lambda_{H\rho}}^{(1)} = & 4\lambda_{H\rho} y_{NS}^2 - 8y_{NS}^2 y_{NI}^2 + \frac{1}{10} \lambda_{H\rho} (120\lambda_H + 20y_{NI}^2 \\ & + 40\lambda_{H\rho} - 45g_2^2 + 60\lambda_\rho - 9g_1^2 + 60y_t^2), \end{aligned} \quad (\text{B5})$$

$$\begin{aligned} \beta_{\lambda_H}^{(1)} = & \frac{27}{200} g_1^4 + \frac{9}{20} g_1^2 g_2^2 + \frac{9}{8} g_2^4 - \frac{9}{5} g_1^2 \lambda_H - 9g_2^2 \lambda_H + 24\lambda_H^2 \\ & + \frac{1}{2} \lambda_{H\rho}^2 + 4\lambda_H y_{NI}^2 - 2y_{NI}^4 + 12\lambda_H y_t^2 - 6y_t^4, \end{aligned} \quad (\text{B6})$$

$$\beta_{y_{NS}}^{(1)} = 5y_{NS}^3 + y_{NS} y_{NI}^2, \quad (\text{B7})$$

$$\beta_{y_{NI}}^{(1)} = \frac{1}{2} (y_{NS}^2 y_{NI} + 3y_{NI}^3) + y_{NI} \left(-\frac{9}{20} g_1^2 - \frac{9}{4} g_2^2 + y_{NI}^2 + 3y_t^2 \right),$$

$$\beta_{y_t}^{(1)} = \frac{3}{2} y_t^3 + y_t \left(-\frac{17}{20} g_1^2 - \frac{9}{4} g_2^2 - 8g_3^2 + y_{NI}^2 + 3y_t^2 \right) \quad (\text{B8})$$

⁶The corresponding super symmetric theory requires a low reheating temperature [90–95] due to the overproduction of gravitino, which is in conflict with the lower bound on the lightest RHN mass $M_1 \gtrsim 2.4 \times 10^9 \text{ GeV}$.

Here, the beta functions are defined as $\beta_x^{(1)} = 16\pi^2 \frac{dx}{d\log \mu}$.

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