

Testing the seesaw mechanism and leptogenesis with gravitational waves: Majorana versus Dirac cases

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We investigate the $B - L$ gauge extension of the Standard Model that the Dirac seesaw mechanism with thermal leptogenesis leaves imprints in the early Universe via the stochastic gravitational background emanating from a network of cosmic strings when $B - L$ symmetry is broken. With right-handed neutrino mass lighter than the typical scale of grand unification, the $B - L$ symmetry protecting the right-handed neutrinos leads to constraints on the Yukawa couplings for both Dirac and Majorana scenarios. Estimating the predicted gravitational wave background we find that future space-borne missions could probe the range concerning thermal Dirac leptogenesis. In a comparative analysis between such probes of gravitational waves sourced from cosmic strings in Dirac and Majorana leptogenesis in the $B - L$ extension, based on the energy scales of the leptogenesis, for instance, gravitational waves detectors will be able to probe the scale of Dirac leptogenesis up to 10^9 GeV, while for Majorana leptogenesis it would be upto 10^{12} GeV.

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I. INTRODUCTION

The evidence of going beyond the Standard Model (SM) of particle physics arises undisputedly from neutrino oscillation experiments involving solar neutrinos [1–6], atmospheric neutrinos [7,8], and reactor neutrinos [9–12]. We infer from these experiments that neutrinos have mass and the flavor states mix due to the propagation of multiple mass eigenstates. This however only quantifies the squared mass differences of the neutrinos and not the absolute mass scales involved. On the other hand, the β -decay experiment KATRIN [13] gives a stringent direct limit on the absolute value of neutrino mass scale, $m_\nu < 0.8$ eV.

On the cosmological frontiers, measurements of the cosmic microwave background radiation by Planck 2018 [14] and large-scale structure constrain the sum of all neutrino masses to $\sum_i m_{\nu_i} < 0.12$ eV [14,15]. On the other

hand, the observed baryon asymmetry of the Universe (BAU) [14,16], often expressed in terms of the baryon to photon ratio is, according to the Planck data [14],¹

$$\eta_B^{\text{CMB}} = \frac{n_B - n_{\bar{B}}}{n_\gamma} = (6.21 \pm 0.16) \times 10^{-10}. \quad (1)$$

This agrees with the values of abundances of light elements extracted from big bang nucleosynthesis data [17].

Just augmenting the Standard Model with two or more right-handed Majorana neutrinos (RHNs), one can explain the tiny neutrino masses, generated via the well-known type-I seesaw mechanism [18–21], while baryon asymmetry can be explained using the baryogenesis via the leptogenesis mechanism [22], provides a unified explanation. Nonetheless this mechanism demands the lightest RHN mass to exceed the Davidson-Ibarra bound [23], $M_1 \gtrsim 10^9$ GeV, assuming the RHN mass spectrum is hierarchical, along with no lepton flavor effects [24–27]. This deflates any hope of direct detection of the high-scale seesaw and leptogenesis since they lie well beyond the energy reaches of the current and future laboratory experiments.

¹It is also expressed as $Y_B = \frac{n_B}{s} = (0.82-0.92) \times 10^{-10}$. We use the central value $Y_B = \frac{n_B}{s} = 8.66 \times 10^{-11}$ in our estimates.

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Nonetheless, certain indirect signatures of new physics like lepton number violation processes through neutrinoless double beta decay [28] or CP violation in neutrino oscillation [29] are searched for in laboratories. There can be certain theoretical constraints on the low energy values of couplings, whose consistency conditions for UV completion such as in $SO(10)$ grand unified theories (GUT) [30–36] or with the demand of electroweak vacuum (meta)stability in the early Universe [37,38] sometimes provide some useful bounds on the large parameter space involved in seesaw and leptogenesis scenarios.

The discovery of Gravitational Waves (GWs) from black hole mergers by the LIGO, Virgo, KAGRA Collaborations [39,40] and evidence for measurements of stochastic GW background by pulsar timing array (PTA) observations [41–47] have led to several new physics cases to detect GWs of primordial and cosmological origins as well. In the context of baryogenesis via Leptogenesis scenarios investigated so far, cosmological pathways to probe such high-scale physics involve predictions of the CMB spectral indices [48] or gravitational waves from local cosmic strings [49–52], global cosmic strings [53], domain walls [54,55], nucleating and colliding vacuum bubbles [56–58], graviton bremsstrahlung [59–61], inflationary tensor perturbations of first order [62,63] and second order [64,65], primordial black holes [66–73], and utilizing features of the primordial non-Gaussianity and curvature bispectrum and trispectrum in CMB [74,75]. Particularly, it has been shown that one may probe directly as a possible footprint of high-scale leptogenesis with GW from cosmic strings, via investigating only the RHN mass sector, the origin of which may occur at a more fundamental level due to symmetry breaking [49,52,67,76,77]. For instance, the RHNs become massive because of the spontaneous breaking of a $U(1)_{B-L}$ gauge invariance, a symmetry that quite naturally gets embedded in several GUTs [78–81], and a breaking of which inevitably produces a cosmic string network that radiates gravitational wave signals. For any $U(1)_{B-L}$ invariant RHN mass term, the Yukawa coupling $y_N \bar{N}_R \Phi N_R^C$, where y_N is the Yukawa coupling, N_R represents the RHN particle and Φ denotes the $B-L$ scalar with vacuum expectation value v_Φ . When the $U(1)_{B-L}$ symmetry is broken, the RHNs obtain mass given by $M_N = y_N v_\Phi$, thereby relating to the properties (like string tension) of cosmic strings that are formed after the symmetry breaking [82–84] leading to production of GWs with a almost scale-invariant spectrum with an overall GW amplitude $\Omega_{\text{GW}} \propto v_\Phi$ [85,86].

Under this circumstance, while most of the leptogenesis scenarios consider the Majorana nature of light neutrinos, one equally appealing alternative is to consider the Dirac nature of light neutrinos. As originally proposed in Refs. [87,88], one can have successful leptogenesis even with light Dirac neutrino scenarios where total lepton number or $B-L$ is conserved just like it is conserved in

the SM of particle physics.² This mechanism is popularly known as the Dirac leptogenesis. What this mechanism exploits is the creation of an equal and opposite amount of lepton asymmetry in left-handed and right-handed neutrino sectors, which is subsequently followed by the conversion of left sector asymmetry into baryon asymmetry via the electroweak sphalerons. This is made possible since the right-handed sector is invisible to the sphalerons inside the SM plasma, only the left-handed asymmetry is converted into baryons. The asymmetry within the ν_R sector, and indeed the ν_R particles themselves, are well-nigh impossible to observe due to their extreme weak interactions. Needless to mention here, the effective Dirac Yukawa couplings are naturally very tiny, which guarantees that the lepton asymmetries in the left- and right-handed sectors do not equilibrate and are washed out. Myriad possible implementations of this idea have been considered, see some in Refs. [94–106]. Particularly, in a few related works [107–109], violation of $B-L$ symmetry was accommodated in a way to preserve the Dirac nature of light neutrinos, while generating lepton asymmetry simultaneously.

In this paper we first discuss Dirac leptogenesis, with thermalized RHNs and therefore suitable within the neutrinophilic [110] two-Higgs-doublet, which naturally explains the smallness of Dirac masses [111–114]. Moreover we follow the framework of lepton-number-violating (LNV) Dirac neutrinos [107,115] to create a lepton asymmetry from the CP -violating decay of a heavy particle that is present in the seesaw model.³ For this reason, the mechanism is actually more close to standard leptogenesis than neutrinogenesis, despite the fact that it contains Dirac neutrinos. Finally we do a detailed investigation of the parameter space of leptogenesis, in $B-L$ extensions of the SM, comparing Majorana and Dirac leptogenesis. In terms of scales of seesaw and leptogenesis, we provide quantitative estimates regarding various GW detectors like LISA and ET, which are able to measure GW from $B-L$ symmetry breaking cosmic strings, probing large regions of such parameter space consisting of different correlations between the seesaw scale and the scales of Dirac and Majorana leptogenesis. In the current scenario,

²The recently proposed mechanism of wash-in leptogenesis [89,90] advocates an approach of getting rid of CP violation in the RHN sector altogether and separates the energy scales of CP violation and $B-L$ violation altogether in the presence of nontrivial chemical background configurations. Like Dirac leptogenesis, here the RHNs are mere spectator processes that reprocess the chemical potentials of the SM particle species in the thermal bath. The action of the electroweak sphalerons on the chemical composition of the SM plasma then results in the usual violation of baryon-plus-lepton number $B+L$, while the action of the RHN interactions results in the violation of $B-L$, also see Refs. [91–93] for gravitational production during inflation.

³Prior to Ref. [115], Ref. [116] mentioned roles of LNV Dirac neutrinos in the early Universe.

we do not consider the lepton flavor effects [24–26], throughout the article, for simplicity.

The paper is organized as follows: In Sec. II, we give a brief introduction to the Dirac leptogenesis model and in the following Sec. III, we discuss the Dirac neutrinos and Dirac leptogenesis. The origin of the gravitational waves from the cosmic strings is discussed in Sec. IV. The numerical results while comparing the Dirac and Majorana leptogenesis with the gravitational wave signal are discussed in Sec. V. Finally we conclude in Sec. VI.

II. THE $B-L$ EXTENDED MODEL

We extend the SM symmetry with a well-motivated beyond the SM framework based on the gauged $U(1)_{B-L}$ symmetry, where B and L are the baryon and lepton numbers, respectively. Considering the gauged $U(1)_{B-L}$ symmetry brings unwanted triangular anomalies to the framework. These anomalies break gauge symmetry at the quantum level, caused by chiral fermions running in triangle loop diagrams involving gauge bosons.⁴ These anomalies can be canceled by introducing three right-handed neutrinos to the framework,⁵ which is also essential for realizing the neutrino masses [117–121]. The overall particle content of our model extends with three right-handed neutrinos ν_{R_α} , $\alpha = \{1, 2, 3\}$, and two singlet scalars φ and ξ , under the SM gauge group. The scalar φ is introduced in order to spontaneously break the $B-L$ symmetry by four units. The ξ acts as a mediator that is responsible for the creation of asymmetry in the right-handed neutrinos [107]. The $B-L$ charge assignment of the new particle content can be seen in Table I.

The relevant Lagrangian takes the following form:

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{kinetic}} + \mathcal{L}_{Z'} - V(H, \varphi, \xi) + \left(y_{\alpha\beta} \bar{L}_\alpha H \nu_{R,\beta} + \frac{1}{2} \kappa_{\alpha\beta} \xi \bar{\nu}_{R,\alpha} \nu_{R,\beta}^c + \text{H.c.} \right), \quad (2)$$

where

$$V(H, \varphi, \xi) \equiv \sum_{X=H, \varphi, \xi} (\mu_X^2 |X|^2 + \lambda_X |X|^4) + \sum_{\substack{X, Y=H, \varphi, \xi \\ X \neq Y}} \lambda_{XY} |X|^2 |Y|^2 - \mu(\varphi \xi^2 + \text{H.c.}), \quad (3)$$

The $\mathcal{L}_{\text{SM}}$, $\mathcal{L}_{\text{kinetic}}$, $\mathcal{L}_{Z'}$ are the all SM particle interactions, kinetic terms of new particle content, and Z' boson

TABLE I. Charge assignment of the additional fields under the extended gauge symmetry.

Fields	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$
ν_R	1	0	-1
φ	1	0	4
ξ	1	0	-2

interaction terms, respectively. H is the SM Higgs doublet. We consider $\mu_H^2, \mu_\varphi^2 < 0 < \mu_\xi^2$. Hence the potential has minimum with $\langle \xi \rangle = 0$, $\langle H \rangle \neq 0 \neq \langle \varphi \rangle$, which breaks the $SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ to $U(1)_{\text{EM}} \times Z_4^L$. Under Z_4^L the leptons transform as $\ell \rightarrow -i\ell$ and $\xi \rightarrow -\xi$. In principle the μ coupling in the last term in Eq. (3) can induce a nonzero vacuum expectation value (VEV) for ξ after φ acquires a VEV. If $\langle \xi \rangle \neq 0$, it spontaneously breaks the residual symmetry Z_4^L , which would generate Majorana mass terms for ν_R , which leads to a standard Majorana leptogenesis scenario, rather than the intended Dirac leptogenesis framework. Since ξ does not acquire a VEV and forbids a Majorana mass term it allows a small Dirac neutrino mass $m_{\nu_{\alpha\beta}} = y_{\alpha\beta} \langle H \rangle$ with the Yukawa coupling of the order $\lesssim 10^{-11}$. The Yukawa coupling $\kappa_{\alpha\beta}$ is a complex symmetric matrix, in general. The $\kappa_{\alpha\beta}$ plays a crucial role in generating the ν_R asymmetry. We define $\xi = (\zeta_1 + i\zeta_2)/\sqrt{2}$, where $\zeta_{1,2}$ are real scalars. The μ term present in the potential induces a mass splitting between the real scalars $\zeta_{1,2}$,

$$M_1^2 = M_c^2 - 2\mu \langle \varphi \rangle, \quad M_2^2 = M_c^2 + 2\mu \langle \varphi \rangle, \quad (4)$$

where $\langle \varphi \rangle \equiv \eta$ is the VEV of φ . The M_c is a common mass term for the real scalars $\zeta_{1,2}$. The M_c decides the mass scale of the M_{ζ_i} . The μ is the coupling strength of φ with ξ .

The mass scale of Z_{B-L} boson is very heavy, therefore it decouples from the thermal bath early in the Universe, and hence it does not thermalize the right-handed neutrinos. The gauge interactions mediated by the heavy Z_{B-L} boson or the scalar interactions involving φ are not sufficient enough to thermalize the RHN.

III. DIRAC NEUTRINOS AND LEPTON NUMBER VIOLATING DIRAC LEPTOGENESIS

In our scenario the neutrinos are Dirac particles, yet the $B-L$ number is violated. A heavy scalar particle, i.e., ξ , generates an asymmetry in the right-handed neutrinos in its CP -violating out-of-equilibrium decay. To generate such an asymmetry we need a second copy of ξ , i.e., $\xi_{1,2}$. The ξ_2 is defined as $\xi_2 = (\zeta_3 + i\zeta_4)/\sqrt{2}$ and also has a mass splitting similar to Eq. (4),

⁴The triangle anomalies for both $U(1)_{B-L}^3$ and the mixed $U(1)_{B-L} - (\text{gravity})^2$ diagrams are nonzero. These triangle anomalies for the SM fermion content turn out to be $\mathcal{A}^{\text{SM}}[U(1)_{B-L}^3] = -3$, $\mathcal{A}^{\text{SM}}[(\text{gravity})^2 \times U(1)_{B-L}] = -3$.

⁵The right-handed neutrinos contribute $\mathcal{A}^{\text{New}}[U(1)_{B-L}^3] = 3$, $\mathcal{A}^{\text{New}}[(\text{gravity})^2 \times U(1)_{B-L}] = 3$ leading to a vanishing total of triangle anomalies [117].

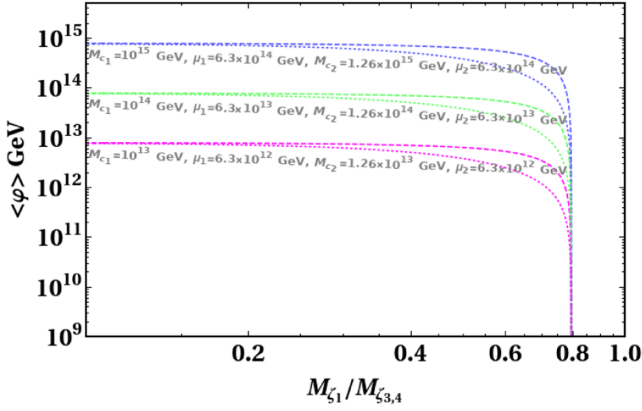


FIG. 1. Variation of mass scales with respect to $\langle\varphi\rangle$ in the plane of ratios M_{ζ_1}/M_{ζ_3} (dashed) and M_{ζ_1}/M_{ζ_4} (dotted).

$$M_3^2 = M_c^2 - 2\mu'\langle\varphi\rangle, \quad M_4^2 = M_c^2 + 2\mu'\langle\varphi\rangle. \quad (5)$$

In Fig. 1 we have shown the dependence of the mass ratios of M_{ζ_1}/M_{ζ_3} and M_{ζ_1}/M_{ζ_4} on $\langle\varphi\rangle$. As we mentioned earlier, the M_c and M'_c dictate the mass scale of the $M_{\zeta_{1,2}}$ and $M_{\zeta_{3,4}}$, respectively. The μ and μ' are the coupling strengths of φ with ξ . We fix the parameters M_c, M'_c, μ , and μ' as given in the description of the figure. By fixing these parameters, we choose to vary $\langle\varphi\rangle$ from 10^9 to 10^{15} GeV. For small values of $\langle\varphi\rangle$, the masses are similarly equal $M_{\zeta_1} \simeq M_{\zeta_2}$ and $M_{\zeta_3} \simeq M_{\zeta_4}$, since the mass splitting is small. Consequently the mass ratios M_{ζ_1}/M_{ζ_3} and M_{ζ_1}/M_{ζ_4} are almost constant until $\langle\varphi\rangle$ approaches the mass scale of M_c and M'_c . Once the $\langle\varphi\rangle$ approaches the mass scale of M_c and M'_c the mass ratios gradually decrease, since the mass splitting increases. Comparatively, the M_{ζ_1}/M_{ζ_4} falls faster than the M_{ζ_1}/M_{ζ_3} , see Fig. 1. The dashed and dotted lines correspond to the mass ratios M_{ζ_1}/M_{ζ_3} and M_{ζ_1}/M_{ζ_4} , respectively.

In presence of both $\xi_{1,2}$ the potential can have mixing terms, *i.e.*, $M_{12}^2 \xi_1^\dagger \xi_2 + \mu_{12} \varphi \xi_1^\dagger \xi_2 + \text{H.c.}$, which leads to mixing of the four fields ζ_i after breaking of $B-L$. The diagonalization of the 4×4 mixing mass matrix leads to the redefinition of the couplings $\kappa_{\alpha\beta}^i$ with the right-handed neutrinos, let us represent it as $\lambda_{\alpha\beta}^i$. For simplicity we neglect the mixing terms between the ξ_1 and ξ_2 and work with the four real scalar fields ζ_i with masses M_{ζ_i} and complex symmetric Yukawa couplings $\lambda_{\alpha\beta}^i = \lambda_{\beta\alpha}^i$. The corresponding Lagrangian can be rewritten as,

$$\mathcal{L} \supset \frac{1}{2} \lambda_{\alpha\beta}^j \zeta_j \bar{\nu}_{R,\alpha} \nu_{R,\beta}^c + \frac{1}{2} \bar{\lambda}_{\alpha\beta}^j \zeta_j \bar{\nu}_{R,\alpha}^c \nu_{R,\beta}. \quad (6)$$

In Fig. 2, we show the vertex and self-energy loop diagrams that contribute to the nonzero CP violation. The decay of the lightest ζ_i leads to a CP asymmetry [107],

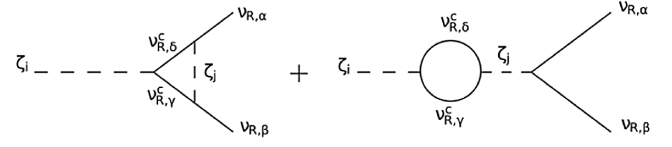


FIG. 2. Vertex and self-energy loop diagrams of ζ decay: $\zeta_i \rightarrow \nu_{R,\alpha} \nu_{R,\beta}$.

$$\epsilon_i \equiv 2 \frac{\Gamma(\zeta_i \rightarrow \nu_R \nu_R) - \Gamma(\zeta_i \rightarrow \nu_R^c \nu_R^c)}{\Gamma(\zeta_i \rightarrow \nu_R \nu_R) + \Gamma(\zeta_i \rightarrow \nu_R^c \nu_R^c)}. \quad (7)$$

The asymmetric contribution from the vertex (ϵ_i^v) and self-energy correction (ϵ_i^s),

$$\begin{aligned} \epsilon_i^v &= \frac{1}{4\pi} \frac{1}{\text{Tr}(\lambda^{i\dagger} \lambda^i)} \sum_{k \neq i} F(x_k) \text{Im}[\text{Tr}(\lambda^{i\dagger} \lambda^k \lambda^{i\dagger} \lambda^k)], \\ \epsilon_i^s &= -\frac{1}{24\pi} \frac{1}{\text{Tr}(\lambda^{i\dagger} \lambda^i)} \sum_{k \neq i} G(x_k) \text{Im}[\{\text{Tr}(\lambda^{i\dagger} \lambda^k)\}^2], \end{aligned} \quad (8)$$

with $x_k \equiv M_i^2/M_k^2 < 1$ and the functions

$$\begin{aligned} F(x) &\equiv \frac{x - \log(1+x)}{x} = \frac{x}{2} + \mathcal{O}(x^2), \\ G(x) &\equiv \frac{x}{1-x} = x + \mathcal{O}(x^2). \end{aligned} \quad (9)$$

The presence of only one scalar ξ_1 would lead to vanishing asymmetry $\epsilon^v = 0 = \epsilon^s$, since $\lambda^2 = i\lambda^1$. For a similar reason ξ_2 alone does not contribute to the CP asymmetry. Thus we require two copies of ξ . The total contribution to the total CP asymmetry comes from the lightest ζ_1 with the ζ_3 and ζ_4 mediation, *i.e.*, say, ϵ_3 and ϵ_4 , respectively. The nonzero contribution comes from ζ_3 and ζ_4 and they are opposite in sign, so that $\epsilon^v \propto F(x_3) - F(x_4)$ and $\epsilon^s \propto G(x_3) - G(x_4)$. For the low scales of $\langle\varphi\rangle$, the contributions of ϵ_3 and ϵ_4 are similarly equal in magnitude. They subtract with each other and the resultant total CP asymmetry is ϵ suppressed, since $\epsilon = \epsilon_3 - \epsilon_4$. The dependence of total ϵ on $\langle\varphi\rangle$ is given in Fig. 3. The blue (dashed), green (dotted), and black (solid) lines are ϵ_3 , ϵ_4 , and ϵ , respectively. Initially, the self-energy contribution ϵ_i^s dominates over vertex contribution ϵ_i^v . When they are equal a suppression occurs as shown in Fig. 3 for ϵ_3 and ϵ_4 . The left- and right-hand side of the suppression are negative and positive contributions, respectively. The subtraction of both magnitudes of ϵ_3 and ϵ_4 gives a negative contribution to the total CP asymmetry on the left side of the solid black line suppression. Note that when both the magnitudes of ϵ_3 and ϵ_4 become equal for the first instance (before black solid line suppression, *i.e.*, $\sim 1.5 \times 10^{14}$ GeV), they add up to give the maximum magnitude for ϵ , which is still a negative. When both ϵ_3 and ϵ_4 magnitudes become equal for the second instance ($\sim 5 \times 10^{14}$ GeV), they subtract

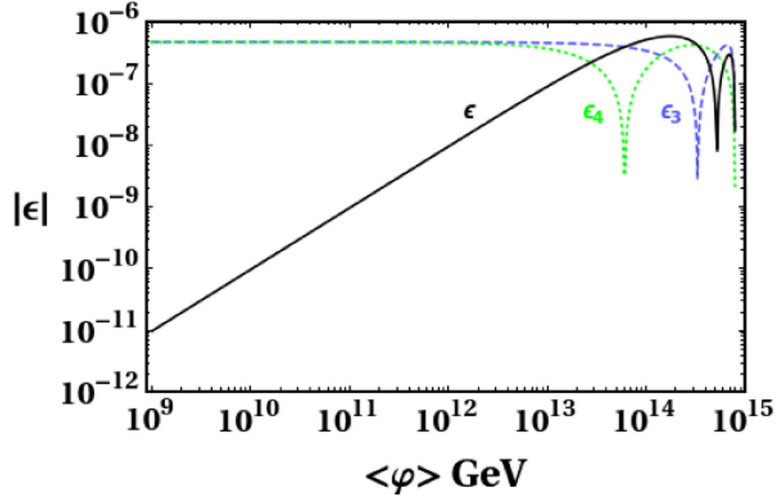


FIG. 3. Dependence of CP asymmetry with respect to $\langle\varphi\rangle$. The parameter choices are: $M_c = 1 \times 10^{15}$ GeV, $\mu = 6.3 \times 10^{14}$ GeV, $M'_c = 1.26 \times 10^{15}$ GeV, $\mu' = 7.94 \times 10^{14}$ GeV, and $\text{Tr}[\lambda^\dagger \lambda] = 10^{-4}$.

each other and give a suppression for the ϵ . At the solid black line suppression point, the CP asymmetry is zero hence there is no CP violation. The choice of parameter space is given in the description of Fig. 3.⁶

The lightest of ζ_i decays and creates an asymmetry in the right-handed neutrinos. The abundance of ζ_1 and the generated right-handed neutrino asymmetry can be estimated by solving the following Boltzmann equations,

$$\begin{aligned} \frac{dY_{\zeta_1}}{dz} &= -\frac{z}{H(M_{\zeta_1})} \Gamma_{\zeta_1}^{\nu_R \nu_R} [Y_{\zeta_1} - Y_{\zeta_1}^{eq}], \\ \frac{dY_{\nu_R}}{dz} &= \epsilon_1 \frac{z}{H(M_{\zeta_1})} \Gamma_{\zeta_1}^{\nu_R \nu_R} [Y_{\zeta_1} - Y_{\zeta_1}^{eq}], \end{aligned} \quad (10)$$

where $Y_{\nu_R} \equiv n_{\nu_R}/s$ is the total lepton asymmetry, *i.e.*, the right-handed neutrino number density n_{ν_R} relative to the entropy density $s = (2\pi^2/45)g_*T^3$. The $Y_{\zeta_1}^{eq} = (45/4\pi^4 g_*)z^2 \mathcal{K}_2(z)$, where $\mathcal{K}_2(z)$ is the modified Bessel function of the second kind. The $\Gamma_{\zeta_1}^{\nu_R \nu_R}$ is the total decay width of the ζ_1 decaying to right-handed neutrinos,

$$\Gamma_{\zeta_i}^{\nu_R \nu_R} = \frac{\text{Tr}(\lambda^{i\dagger} \lambda^i)}{4\pi} M_{\zeta_i}. \quad (11)$$

The $H(M_{\zeta_1}) = 1.66\sqrt{g_*} \frac{M_{\zeta_1}^2}{M_{\text{Pl}}}$ is the Hubble expansion rate at $T = M_{\zeta_1}$.

The out-of-equilibrium condition for the decay of the lightest ζ_i can be estimated as

⁶When $2\mu\langle\varphi\rangle$ approach the M_c value $M_{\zeta_1} \rightarrow 0$, which leads to a nonviable scenario.

$$\Gamma(\zeta_i \rightarrow \nu_R \nu_R, \nu_R^c \nu_R^c) \ll H(T \sim M_{\zeta_i}) \simeq 1.66\sqrt{g_*} \frac{M_{\zeta_i}^2}{M_{\text{Pl}}}. \quad (12)$$

This condition leads to

$$\frac{\text{Tr}(\lambda^{i\dagger} \lambda^i)}{10^{-6}} \ll \frac{M_{\zeta_i}}{10^{11} \text{ GeV}}. \quad (13)$$

The asymmetry generates in the right-handed neutrinos Y_{ν_R} after the breaking of $B-L$ and before the electroweak phase transition, so that sphalerons transfer the lepton symmetry to the baryon asymmetry. Initially, the ζ_1 creates the right-handed asymmetry,⁷ on which the sphalerons do not act. Here we require another term that can transfer the right-handed asymmetry efficiently to left-handed sector. This role is played by the term $w_{\alpha\beta} \bar{L}_\alpha \Phi \nu_{R,\beta}$, where Φ is the second Higgs doublet.⁸ This term thermalizes⁹ the ν_R and transfers the asymmetry from right-handed sector to left-handed sector, *i.e.*, $\Delta_{\nu_R} \rightarrow \Delta_L$. The second scalar doublet couples with neutrino sector due to an additional $U(1)$ symmetry under which only Φ and ν_R are charged. We solve the Boltzmann equations, given in Eq. (10), to

⁷We assume the initial abundance of the right-handed neutrinos to be 0.

⁸We can consider that the second Higgs doublet does not acquire a VEV, in which case the neutrinos get their mass via SM Higgs H . If we consider a small VEV to the second Higgs doublet, it can give rise to neutrino mass naturally by softly breaking the new $U(1)$ symmetry via the term $\mu_{12}^2 H^\dagger \Phi$ in the potential. This term induces a small VEV for Φ , *i.e.*, $\langle\Phi\rangle = \mu_{12}^2/M_\Phi^2 \langle H\rangle$, which gives rise to small Dirac neutrino masses $M_{\alpha\beta} = w_{\alpha\beta} |\langle\Phi\rangle|$.

⁹If $U(1)_{B-L}$ is a gauge symmetry, the Z' interactions (and Φ) can thermalize the ν_R . If $U(1)_{B-L}$ is a global symmetry, the Φ can thermalize the ν_R .

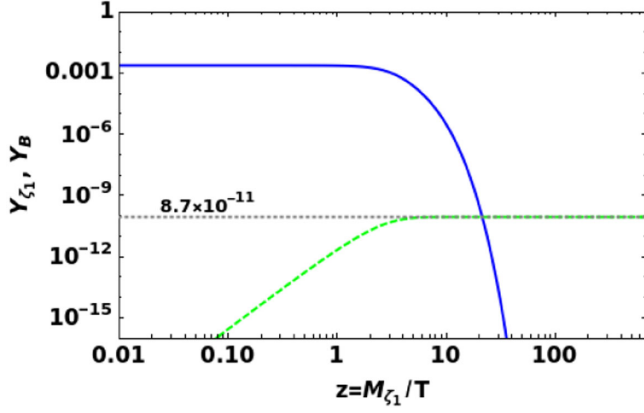


FIG. 4. Abundances of Y_{ζ_1} (blue) and Y_B (green). We consider $M_c = 8 \times 10^{13}$ GeV, $\mu = 10^{13}$ GeV, $M'_c = 10^{14}$ GeV, $\mu' = 10^{13}$ GeV, $\langle \phi \rangle = 10^{12}$ GeV, and $\text{Tr}[\lambda^\dagger \lambda] = 7.2 \times 10^{-4}$. The horizontal dashed line corresponds to the observed baryon asymmetry of the Universe.

estimate the abundance of Y_{ζ_1} and the baryon asymmetry Y_B , shown in Fig. 4. The total lepton asymmetry can be expressed as

$$Y_{\nu_R} \equiv \frac{n_{\nu_R}}{s} = \epsilon_i Y_{\zeta_1}^{eq} \sim \frac{\epsilon_i^v + \epsilon_i^s}{g_*}. \quad (14)$$

We assume equilibrium of the right-handed neutrino with the SM particles as well as sphalerons, and we use the chemical potentials to describe the plasma including the chemical potential for the right-handed neutrinos, resulting in the equilibrium condition $3B + L = 0$ or

$$Y_B = \frac{1}{4} Y_{B-L}, \quad Y_L = -\frac{3}{4} Y_{B-L}, \quad (15)$$

for three generations (and an arbitrary number of Higgs doublets). Once the right-handed neutrino asymmetry is produced a (1/4) part of it transfers to the baryon asymmetry of the Universe. The orange (solid) and blue (dashed) lines are the abundance of the ζ_1 and the baryon asymmetry Y_B , respectively.

IV. PRIMORDIAL GW SPECTRUM FROM LOCAL COSMIC STRINGS

Cosmic strings (CSs) may originate as fundamental objects such as in string theory [122,123] as well as topological defects in field theory during a symmetry-breaking phase transition. We will focus on cosmic strings that were formed in the early Universe due to the spontaneous breaking of a local $U(1)$ symmetry, this is sometimes called field theory strings. The simplest model where cosmic strings may arise is a $U(1)$ local gauge or global theory in the context of a complex scalar field denoted by ϕ . The Lagrangian density for such a theory can be written as,

$$\mathcal{L} = D_\mu \phi^* D^\mu \phi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - V(\phi) \quad (16)$$

where $D_\mu \phi = \partial_\mu \phi + ieA_\mu \phi$ is the covariant derivative and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the field tensor.

$V(\phi)$ is the typical Mexican hat potential written as,

$$V(\phi) = \frac{1}{2} \lambda' \left(\phi^* \phi - \frac{1}{2} \eta^2 \right)^2 \quad (17)$$

where λ' is the self-quartic coupling and $\langle \phi \rangle = \eta$ is the VEV of the scalar field.

Cosmic strings are essentially field configurations at the top of this Mexican hat potential. The CSs have a negligible thickness [with widths given by approximately $1/(\sqrt{\lambda}\eta)$] and hence their dynamics in an expanding Universe are well described by something known as the Nambu-Goto action, in its zero-width approximation. The mass per unit length μ of a local CS is given by [124],

$$\mu = 2\pi n \eta^2 \quad (18)$$

where η is the vacuum expectation value of the scalar field constituting the CS and n is the (integer-valued) winding number of the cosmic string formed. Only $n = 1$ is considered to be a stable configuration [125,126]. The local CS network starts forming at a temperature,

$$T_{\text{CS}} \simeq \eta \simeq 10^{11} \text{ GeV} \left(\frac{G\mu}{10^{-15}} \right)^{1/2}, \quad (19)$$

which is also the temperature of a $U(1)$ -breaking phase transition. G is Newton's constant and the dimensionless parameter $G\mu$ is known as the string tension.

Straight infinitely long strings are stable against decay due to their topological nature, therefore they contribute negligibly toward the GW spectra compared to cosmic string loops. Local CS loops, formed due to intercommutation events of CS, are free of any topological charge, so they contribute significantly toward GW radiation emission [127]. The GW radiation power for CS loops is given by,

$$P_{\text{GW}} = \Gamma G\mu^2 \quad (20)$$

where Γ is the total GW emission efficiency of loops and is determined typically from the Nambu-Goto simulation. This has been estimated to be $\Gamma \simeq 70$ [128].

The GW spectrum from local CS loops observed today is given follows [124]:

$$\Omega_{\text{GW}}(f) \equiv \frac{f}{\rho_c} \left| \frac{d\rho_{\text{GW}}}{df} \right| = \sum_k \Omega_{\text{GW}}^{(k)}(f) \quad (21)$$

where

$$\Omega_{\text{GW}}(f) = \sum_k \frac{1}{\rho_c} \int_{t_{\text{osc}}}^{t_0} d\tilde{t} \int_0^1 d\alpha \cdot \Theta \left[t_i - \frac{l_*}{\alpha} \right] \cdot \Theta[t_i - t_{\text{osc}}] \cdot \left[\frac{a(\tilde{t})}{a(t_0)} \right]^4 \cdot P_{\text{GW}}^{(k)} \times \left[\frac{a(t_i)}{a(\tilde{t})} \right]^3 \cdot \mathcal{P}_{\text{loop}}(\alpha) \cdot \frac{dt_i}{df} \cdot \frac{dn_{\text{loop}}}{dt_i}. \quad (22)$$

The first Heaviside function $\Theta[t_i - \frac{l_*}{\alpha}]$ ensures that loops smaller than critical length l_* will not contribute toward the GW spectrum as their main decay channel is through massive particle production. $l_* \equiv l_c$ is used for cusps and l_k is used for kinks.

$$l_c \equiv \beta_c \frac{\mu^{-1/2}}{(\Gamma G \mu)^2}; \quad l_k \equiv \beta_k \frac{\mu^{-1/2}}{\Gamma G \mu} \quad (23)$$

β_c and β_k are some $\mathcal{O}(1)$ numbers. The second Heaviside function $\Theta[t_i - t_{\text{osc}}]$ gets rid of all those loops that were formed before the network formation [Eq. (19)] or that form during the friction dominated era,

$$T \lesssim T_{\text{fric}} = \frac{4 \times 10^9 \text{ GeV}}{\beta} \left(\frac{g_*}{100} \right)^{1/2} \left(\frac{G\mu}{10^{-11}} \right) \quad (24)$$

and $t_{\text{osc}} = \text{Max}[t_{\text{fric}}, t_F]$, where t_F is the time of CS network formation, defined as $\sqrt{\rho_{\text{tot}}(t_F)} \equiv \mu$ where ρ_{tot} is the total energy budget of the Universe. In the presence of friction, the string motion is damped at high temperature until the time t_{fric} . dn_{loop}/dt_i expresses the rate of formation of loops with a distribution size $\mathcal{P}_{\text{loop}}(\alpha)$. These objects initially redshift as a^{-3} before radiating GW with power $P_{\text{GW}}^{(k)}$ after which they dilute as a^{-4} . Using such simplifications, one may express Eq. (22) as [124],

$$\Omega_{\text{GW}}^{(k)}(f) = \frac{1}{\rho_c} \cdot \frac{2k}{f} \cdot \frac{\mathcal{F}_\alpha}{\alpha(\alpha + \Gamma G \mu)} \int_{t_{\text{osc}}}^{t_0} d\tilde{t} \frac{C_{\text{eff}}(t_i)}{t_i^4} \times \left[\frac{a(\tilde{t})}{a(t_0)} \right]^5 \left[\frac{a(t_i)}{a(\tilde{t})} \right]^3 \Theta(t_i - t_{\text{osc}}) \Theta \left(t_i - \frac{l_*}{\alpha} \right) \quad (25)$$

$\alpha()$ is the loop length at the time of formation and $\mathcal{F}_\alpha$ gives the fraction of loops that forms with size α . t_i is the time of loop production, which is dependent on the emission time $\tilde{t}$ leading the observed frequency today to be,

$$t_i(f, \tilde{t}) = \frac{1}{\alpha + \Gamma G \mu} \left[\frac{2k}{f} \frac{a(\tilde{t})}{a(t_0)} + \Gamma G \mu \tilde{t} \right]. \quad (26)$$

t_0 is the present time. $C_{\text{eff}}(t_i)$ is the loop-production efficiency given by,

$$\tilde{C}_{\text{eff}} \equiv \sqrt{2} C_{\text{eff}}(t) = \frac{\tilde{c} \bar{v}(t)}{\xi^3(t)} \quad (27)$$

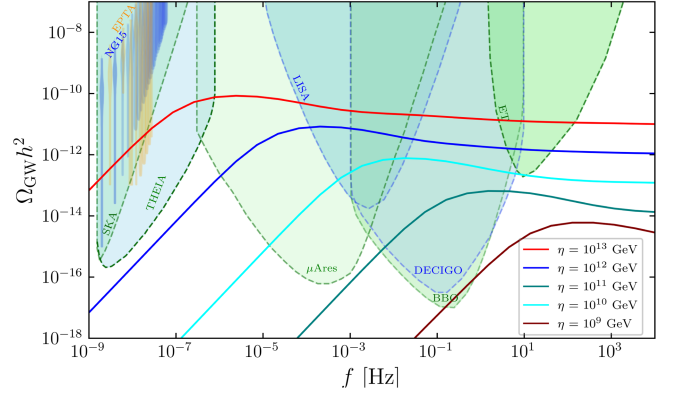


FIG. 5. Variation of the Stochastic gravitational wave background spectra from local cosmic strings with respect to the VEV $\langle \varphi \rangle = \eta$.

where $\tilde{c} = 0.23 \pm 0.04$ is a phenomenological parameter that quantifies the loop chopping efficiency factor. $\bar{v}$ is the root mean square speed of the string loops and ξ is defined as the quantity $\xi \equiv L/t$, where L is the correlation length of the cosmic string.

Figure 5 shows the full GW spectrum due to local CS computed using Eq. (25) assuming a Λ cold dark matter universe assuming the strings were formed after the end of primordial cosmic inflation. The GW spectrum can be broadly classified into a flat high-frequency regime and the radiation-to-matter transition at the low-frequency regime. Higher frequencies correspond to earlier and earlier in times, i.e., they correspond to a radiation-dominated era. The local CS loops that were produced and started GW emission during the radiation-dominated era have a nearly flat GW spectrum with amplitude typically given by,

$$\Omega_{\text{std}}^{\text{CS}} h^2 \simeq \Omega_r h^2 \mathcal{G}(\tilde{T}_M) \left(\frac{\eta}{M_{\text{pl}}} \right) \quad (28)$$

where $\Omega_r h^2 \simeq 4.2 \times 10^{-5}$ [129]. $\tilde{T}_M$ is the temperature of the Universe at the time of the maximal GW emission $\tilde{t}_M$,

$$\tilde{t}_M = \frac{\alpha}{2\Gamma G \mu} t_i. \quad (29)$$

Interestingly, the deviation from the nearly flat spectrum occurs due to a change in the number of relativistic degrees of freedom, which is approximated by,

$$\mathcal{G}(T) \equiv \left(\frac{g_*(T)}{g_*(T_0)} \right) \left(\frac{g_{*S}(T_0)}{g_{*S}(T)} \right)^{4/3} = 0.39 \left(\frac{106.75}{g_*(T)} \right)^{1/3}. \quad (30)$$

V. NUMERICAL RESULTS

All kinds of gravitational wave experimental efforts can be sorted into:

- (1) Ground based interferometers: Laser Interferometer Gravitational-wave Observatory (LIGO) [39,40,130–133], Advanced LIGO [134,135], Einstein Telescope (ET) [136,137], Cosmic Explorer (CE) [138].
- (2) Space based interferometers: μ -ARES [139], Laser Interferometer Space Antenna (LISA) [140,141], Big-Bang Observer [142,143], Deci-Hertz Interferometer Gravitational-wave Observatory (DECIGO) [144], Upgraded DECIGO [145–147].
- (3) PTAs: European Pulsar Timing Array; these are recasts of star surveys [148–150], Square Kilometre Array (SKA) [41–43], North American Nanohertz Observatory for Gravitational Waves [151–155].

A. Gravitational wave as probe of Dirac leptogenesis

The VEV of the scalar $\langle\phi\rangle$ plays a significant role in Dirac leptogenesis and the gravitational waves. Another important parameter that plays a key role in Dirac leptogenesis is $\text{Tr}[\lambda^\dagger\lambda]$, which quantifies the amount of CP violation. To study the allowed parameter space of the baryon asymmetry we plot Fig. 6, where the contours of baryon asymmetry are drawn in the plane of $\text{Tr}[\lambda^\dagger\lambda]$ and $\langle\phi\rangle$. We see the observed baryon asymmetry of the Universe is allowed for a wide range of $\text{Tr}[\lambda^\dagger\lambda]$ and $\langle\phi\rangle$. Here we fix the parameter as: $M_c = 10^{16}$ GeV, $\mu = 6.31 \times 10^{15}$ GeV, $M'_c = 1.26 \times 10^{16}$ GeV, and $\mu' = 7.94 \times 10^{15}$ GeV. As the CP asymmetry parameter ϵ is proportional to $\text{Tr}[\lambda^\dagger\lambda]$, for a particular contour of baryon asymmetry if we decrease $\text{Tr}[\lambda^\dagger\lambda]$ the resultant baryon asymmetry decreases. To compensate that the $\langle\phi\rangle$ has to increase so that the ϵ increases, as discussed in Fig. 3, which results in the desired contour of baryon asymmetry as shown in the figure. For the larger values of $\text{Tr}[\lambda^\dagger\lambda]$, the $\langle\phi\rangle$ takes lower values for the same reason. Note that for higher values of $\langle\phi\rangle$ the contour of baryon asymmetry moves toward the larger values of $\text{Tr}[\lambda^\dagger\lambda]$ due to a suppression occurring in the total CP asymmetry, as discussed in Fig. 3. Hence to compensate for that $\text{Tr}[\lambda^\dagger\lambda]$ increases. At the suppression point the CP asymmetry is zero, hence no baryon asymmetry. The white patch region shows the excessive amount of baryon asymmetry. For a decrease of the mass scales described in the caption of the figure¹⁰ the contours shift down by an order while leaving $\text{Tr}[\lambda^\dagger\lambda]$ intact.

In Fig. 6 we also show various GW detectors sensitivity limits come from DECIGO, μ ARES, LISA, ET, CE, SKA, etc.. The region above the particular sensitivity limit can be probed by the corresponding GW detectors with signal-to-noise ratio (SNR) estimations as shown in the Appendix. We observe that the Dirac leptogenesis can be probed by most GW detectors like DECIGO, μ ARES, LISA, ET, CE,

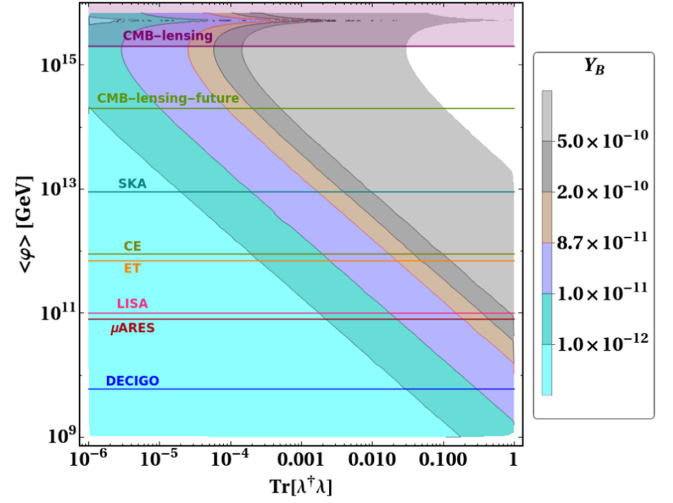


FIG. 6. Dirac leptogenesis: Parameter space probed by individual GW detectors in the plane of $\text{Tr}[\lambda^\dagger\lambda]$ and $\langle\phi\rangle$. We fix the parameters: $M_c = 10^{16}$ GeV, $\mu = 6.31 \times 10^{15}$ GeV, $M'_c = 1.26 \times 10^{16}$ GeV, and $\mu' = 7.94 \times 10^{15}$ GeV.

SKA and CMB-lensing-future. The region above CMB-lensing is ruled out.

B. Gravitational wave as probe of Majorana leptogenesis

In the type-I seesaw mechanism, the heavy right-handed neutrinos N_i , $\{i = 1, 2, 3\}$ are introduced to address the smallness of the light neutrino masses. The relevant Lagrangian is given as follows,

$$\mathcal{L} \supset y_{i\alpha} \bar{N}_i H^\dagger L_\alpha + \frac{1}{2} M_N \bar{N} N + \text{H.c.}, \quad (31)$$

where y is the Yukawa coupling and $M_N = y_m \langle\phi\rangle$ is a diagonal mass matrix of the right-handed neutrinos. In standard Majorana leptogenesis the heavy right-handed neutrino couples with L and H . The lightest of the heavy right-handed neutrinos decays $N_1 \rightarrow L + H$, while fulfilling the Sakharov conditions, and generates lepton asymmetry. A part of the lepton asymmetry transfers to the baryon asymmetry by sphaleron transitions. The Boltzmann equations governing the generation of the lepton asymmetry can be given as follows,

$$\frac{dY_{N_1}}{dz} = -D(z)(Y_{N_1} - Y_{N_1}^{eq}), \quad (32)$$

$$\frac{dY_{B-L}}{dz} = -\epsilon_1 D(z)(Y_{N_1} - Y_{N_1}^{eq}) - W_{ID}(z)Y_{B-L}, \quad (33)$$

where

$$z = M_1/T, \quad (34)$$

¹⁰The choice of high mass scales in Fig. 6 is to consider the scale of $\langle\phi\rangle$ to cover observations involving CMB-lensing, CMB-lensing-future, etc.

$$D(z) = K_1 z \frac{\mathcal{K}_1(z)}{\mathcal{K}_2(z)}, \quad (35)$$

$$W_{ID}(z) = \frac{1}{4} K_1 z^3 \mathcal{K}_1(z), \quad (36)$$

$$K_1 = \frac{\Gamma_{N_1}}{H(M_{N_1})}, \quad (37)$$

where $\mathcal{K}_{1,2}$ are modified Bessel functions of first and second kind. The equilibrium number density defined as

$$Y_{N_1}^{eq} = \frac{3}{4} \zeta(3) \frac{45}{2\pi^4 g_*} z^2 \mathcal{K}_2(z), \quad (38)$$

where g_* is the relativistic degrees of freedom. For a hierarchical right-handed neutrino masses the amount of CP violation is calculated by the CP -asymmetry parameter ϵ_1 as,

$$\epsilon_1 = \frac{3}{16\pi (y^\dagger y)_{ii}} \sum_{j \neq i} \text{Im} [(y^\dagger y)_{ji}^2] \left(\frac{M_{N_i}}{M_{N_j}} \right). \quad (39)$$

Once the lepton asymmetry has been generated, part of it can be transferred to the baryon asymmetry,

$$Y_B = -\frac{28}{79} Y_{B-L}. \quad (40)$$

We examine the Majorana leptogenesis in Fig. 7, the baryon asymmetry portrayed as contours in the plane of $\text{Tr}[\lambda^\dagger \lambda]$ and $\langle \varphi \rangle$ while imposing various GW detector constraints.¹¹ For large values of $\text{Tr}[\lambda^\dagger \lambda]$ and $\langle \varphi \rangle$, we observe the correct baryon asymmetry of the Universe. At higher scales of $\langle \varphi \rangle$ the contours of the baryon asymmetry behave as independent of $\text{Tr}[\lambda^\dagger \lambda]$. But as we move toward lower $\langle \varphi \rangle$ the washouts dominates and the baryon asymmetry decreases. To achieve the desired baryon asymmetry contour at low $\langle \varphi \rangle$ the coupling $\text{Tr}[\lambda^\dagger \lambda]$ has to increase, therefore contour lines move toward the right. Meanwhile for large values of $\text{Tr}[\lambda^\dagger \lambda]$ the washouts increase too and to compensate that the $\langle \varphi \rangle$ has to increase hence the contour line rises slightly at higher $\text{Tr}[\lambda^\dagger \lambda]$. The parameter choice is given in the description of Fig. 7. The contours of BAU shift down to low scales of $\langle \varphi \rangle$ and $\text{Tr}[\lambda^\dagger \lambda]$ upon decreasing the mass ratio of RHNs, M_{N_2}/M_{N_1} . In Fig. 7 we show various GW detectors sensitivity limits come from DECIGO, μARES , LISA, ET, CE, SKA, etc. The region above the particular sensitivity limit can be probed by the corresponding GW detectors with SNR estimations as shown in the Appendix.

¹¹We consider the scale of $B-L$ breaking, $\langle \varphi \rangle \gtrsim 10^9$ GeV, to make the M_N scale consistent with the Davidson-Ibarra bound [23].

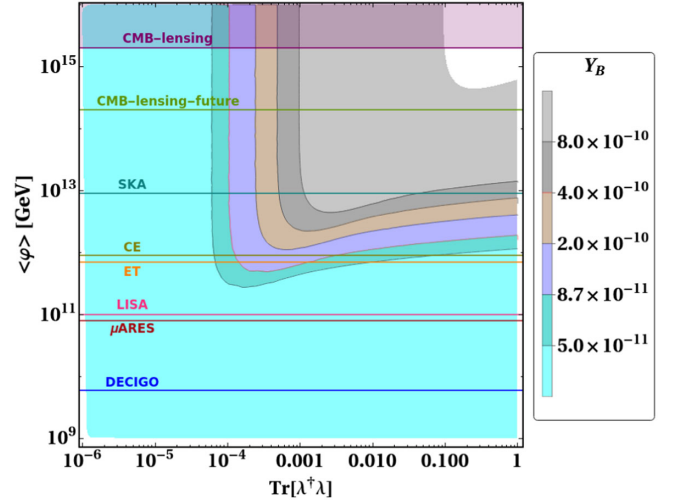


FIG. 7. Majorana leptogenesis: Parameter space probed by individual GW detectors in the plane of $\text{Tr}[\lambda^\dagger \lambda]$ and $\langle \varphi \rangle$. We fix the parameters: $y_m = 1$, $M_{N_2}/M_{N_1} = 10^2$.

C. Comparison: Dirac versus Majorana leptogenesis

In general, the $B-L$ number is violated in Majorana leptogenesis, whereas in Dirac leptogenesis the $B-L$ is conserved [87,106]. In the current scenario, the Dirac leptogenesis $B-L$ number violates similar to the Majorana scenario [107]. The origin of lepton asymmetry is different in both scenarios. In the Dirac leptogenesis case a heavy scalar generates asymmetry in right-handed neutrinos and then transfers to the lepton asymmetry, whereas in the Majorana leptogenesis case a heavy right-handed neutrino generates the lepton asymmetry. In the following, to establish a comparison between these scenarios we compare the mass scales of the leptogenesis.

Figure 8 depicts the mass scale of Dirac leptogenesis M_{ξ_1} with respect to $\langle \varphi \rangle$ with the contours of various choices of M_c . Note that mass scale of M_{ξ_1} depends on the choice of M_c . For a fixed value of M_c , the M_{ξ_1} is barely constant with the increment of $\langle \varphi \rangle$, but as $\langle \varphi \rangle$ approaches M_c the M_{ξ_1} gradually decreases, see Eq. (4). When the $2\mu\langle \varphi \rangle$ similar to the M_c^2 the M_{ξ_1} drops drastically and can approach zero with high fine-tuning.¹² For lower values of M_c the allowed mass scale for M_{ξ_1} decreases. The white patch region on the top left is redundant. The μ is taken one order lesser than M_c , i.e., $\mu = 0.63 \times M_c$ as an arbitrary choice.

We provide three benchmark points in Table II that are well agreement with the GWs originated from the cosmic strings and the Dirac leptogenesis. We consider three representative points for the $\langle \varphi \rangle = \{10^{13}(\text{yellow}), 10^{12}(\text{red}), 10^{11}(\text{black})\}$ GeV, from Fig. 5, the GW signal, which

¹²We use the fine-tuning of $2\mu\langle \varphi \rangle$ with M_c^2 of the order $\sim 10^{-5}$ to get such low M_{ξ_1} shown in Fig. 8.

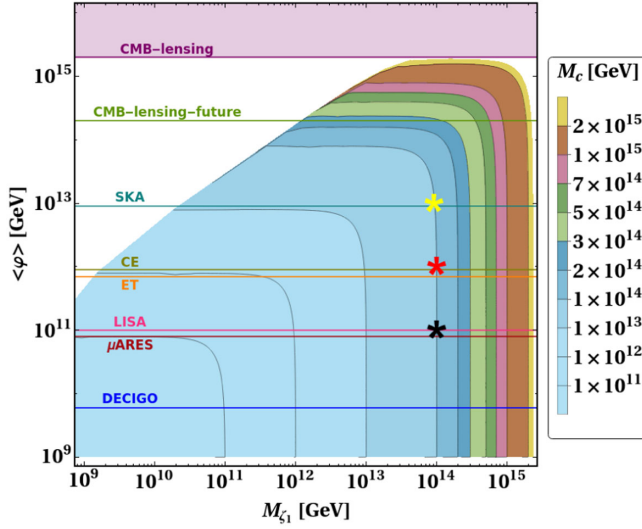


FIG. 8. Dirac leptogenesis: Parameter space probed by individual GW detectors in the plane of M_{ζ_1} and $\langle\varphi\rangle$. We fix the parameters: $\mu = 0.63 \times M_c$ GeV, $M'_c = 1.26 \times M_c$, and $\mu' = 0.63 \times M'_c$ GeV. The black, red, and magenta star points represent the benchmark points (BPs) provided in Table II. The plot represents the correlation between the scale of leptogenesis and the scale of seesaw and satisfies Eq. (4).

overlaps with the LISA and ET sensitivity. Based on $\langle\varphi\rangle$ we consider three representative points for the M_c that give different mass scales for the M_{ζ_1} . The other parameter choices are taken as specified in the description of the Fig. 8. We observe that three BPs well agree with the observed BAU. These BPs fall near the sensitivity of SKA, CE, ET, LISA and μ ARES, future detection regions. The BPs can be detectable by SKA, CE, ET, LISA, μ ARES, and DECIGO SNR threshold sensitivities and can give a hint toward the Dirac nature of light neutrinos and leptogenesis.

We also provide three BPs for Majorana leptogenesis for a comparison with the Dirac leptogenesis parameter space. Here also we consider three representative points for the $\langle\varphi\rangle = \{10^{13}(\text{yellow}), 10^{12}(\text{red}), 10^{11}(\text{black})\}$ GeV, from Fig. 5, the GW signal which overlaps with the LISA and ET. For comparison we consider the mass scale of the right-handed neutrino M_{N_1} is similar to the mass scale of Dirac leptogenesis, i.e., M_{ζ_1} . In Fig. 9, we show the mass scale of the Majorana leptogenesis with respect to φ with the contours of Yukawa coupling y_m . The contours are

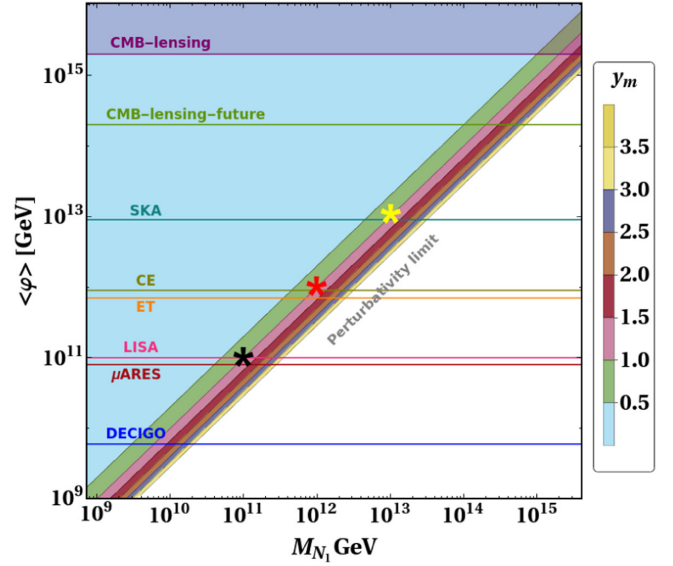


FIG. 9. Majorana leptogenesis: Parameter space probed by individual GW detectors in the plane of M_{N_1} and $\langle\varphi\rangle$. The mass ratio: $M_{N_2}/M_{N_1} = 10^2$. The black, red, and yellow star points represent the BPs provided in Table III. The plot represents the correlation between the scale of leptogenesis and the scale of seesaw.

straight lines, since the M_{N_1} and $\langle\varphi\rangle$ are proportional. We observe that two BPs with $\langle\varphi\rangle = 10^{13}, 10^{12}$ GeV well agree with the observed BAU, for the mass hierarchy $M_{N_2}/M_{N_1} = 10^2$, whereas the $\langle\varphi\rangle = 10^{11}$ GeV gives insufficient baryon asymmetry. The two BPs fall near the sensitivity of SKA, CE, ET, LISA, μ ARES, and DECIGO SNR threshold sensitivities and can give a hint toward the Dirac nature of light neutrinos and leptogenesis.

We observe that in the Dirac case for the low scale of $M_{\zeta_1}(M_c)$ the BAU allowed parameter space is restricted to low scales of $\langle\varphi\rangle$ and the Dirac leptogenesis allowed parameter space may leave hints in detector sensitivities like LISA, μ ARES, and DECIGO only, whereas for the high scale of $M_{\zeta_1}(M_c)$ the Dirac leptogenesis allowed parameter space may leave hints at various GW detectors like DECIGO, μ ARES, LISA, ET, CE, SKA, and CMB-lensing-future. In the Majorana case for the low scale of M_{N_1} the baryon asymmetry is inefficient to get the correct

TABLE II. Dirac leptogenesis: BPs are corresponding to the star (*) symbol points italic (BP 1), bold (BP 2), and bold italics (BP 3) represented in Fig. 8.

Parameters	M_{ζ_1} (GeV)	M_c (GeV)	μ (GeV)	$\langle\varphi\rangle$ (GeV)	$\text{Tr}(\lambda^\dagger \lambda)$	Y_B
BP 1 (*)	9.3×10^{13}	10^{14}	6.3×10^{13}	10^{13}	2.8×10^{-5}	8.7×10^{-11}
BP 2 (*)	9.9×10^{13}	10^{14}	6.3×10^{13}	10^{12}	1.6×10^{-4}	8.7×10^{-11}
BP 3 (*)	9.9×10^{13}	10^{14}	6.3×10^{13}	10^{11}	1.5×10^{-3}	8.7×10^{-11}

TABLE III. Majorana leptogenesis: BPs are corresponding to the star symbol points italic (BP 1), bold (BP 2), and bold italics (BP 3) represented in Fig. 9.

Parameters	M_{N_1} (GeV)	y_m	$\langle\varphi\rangle$ (GeV)	$\text{Tr}(\lambda^\dagger\lambda)$	Y_B
BP 1 (*)	10^{13}	1	10^{13}	1×10^{-4}	8.7×10^{-11}
BP 2 (*)	10^{12}	1	10^{12}	1.2×10^{-4}	8.7×10^{-11}
BP 3 (*)	10^{11}	1	10^{11}	5.2×10^{-5}	1.8×10^{-11}

abundance, based on the mass ratio of RHNs, M_{N_2}/M_{N_1} , whereas for high scale M_{N_1} the BAU may be achievable while leaving hints at DECIGO, μ ARES, LISA, ET, CE, and SKA. For both Dirac and Majorana cases the BPs corresponding to $\langle\varphi\rangle = \{10^{13}, 10^{12}, 10^{11}\}$ GeV show the high scale leptogenesis may be probed by the stochastic gravitational waves emanating from a network of cosmic strings.

VI. DISCUSSION AND CONCLUSION

Thermal Dirac leptogenesis through the Dirac seesaw mechanism gives an elegant and minimal explanation for two outstanding puzzles in the standard model: namely the generation of tiny neutrino masses and matter-antimatter asymmetry of the Universe involving no LNV on the whole.¹³ Unfortunately, the scales of new physics, namely the seesaw scale and the leptogenesis scale, are naturally well beyond what we can directly test in terrestrial laboratory experiments. Given the fundamental nature of these puzzles, indirect cosmological tests of Dirac leptogenesis are of great value and we propose cosmic strings as a powerful witness of the paradigm. We find that, if present, gravitational wave radiation from cosmic strings will carry the imprints of thermal Dirac leptogenesis. If there is $U(1)_{B-L}$ surviving below from GUT symmetry breaking scales to these scales to protect the right-handed neutrino mass, then we show that cosmic strings appear through the breaking of this $B-L$ symmetry predicting a spectrum of stochastic gravitational waves and, with future detectors being able to probe the entire mass range relevant to the paradigm of Dirac leptogenesis. If cosmic strings GW spectrum is discovered and falls into the energy scales relevant for the Dirac seesaw mechanism and leptogenesis, it would provide exciting hints of dynamics in the lepton sector at high energy scales. We showed a comparative analysis between such probes of GW in the context of Majorana and Dirac leptogenesis within the $B-L$ extension paradigm, and based on the energy scales of seesaw and leptogenesis. We have shown, for instance, the LISA detector will be able to probe the scale of seesaw $\langle\varphi\rangle = 10^{11}$ GeV or higher while ET will be able to probe $\langle\varphi\rangle = 10^{12}$ GeV or higher. Along with these, Figs. 6 and 7,

respectively, also tell us about the scales of Dirac and Majorana seesaws, for which correct BAU can be satisfied by varying the Yukawas ($\text{Tr}[\lambda^\dagger\lambda]$) that control the CP violation. It can be inferred that for the Majorana case, $\langle\varphi\rangle \geq 10^{12}$ GeV agree with observed BAU while below $\langle\varphi\rangle \leq 10^{12}$ GeV, the $Y_B \leq 8.66 \times 10^{-11}$ given $y_m = 1$, $\frac{M_{N_2}}{M_{N_1}} = 10^2$. This situation is however relaxed for the Dirac seesaw case, and the observed BAU can be satisfied for a wide range of values of $\langle\varphi\rangle$. Given the correlations between the scale of seesaw and scale of leptogenesis, which are different for Dirac leptogenesis and Majorana leptogenesis, the ET detector will be able to probe the scale of Dirac leptogenesis $M_{\zeta_1} \sim 10^9$ GeV or higher (see Fig. 8), while the same for Majorana leptogenesis $M_{N_1} \geq 10^{12}$ GeV within perturbative coupling limits (see Fig. 9). LISA on the other hand will be able to probe even up to lower scales of Dirac leptogenesis, $M_{\zeta_1} \sim 5 \times 10^8$ GeV, and that of Majorana leptogenesis $M_{N_1} \gtrsim 8 \times 10^{11}$ GeV.

Sources of primordial gravitational waves, like those from topological defects like cosmic strings, are a useful search for new physics especially high energy scales like that of leptogenesis, which is otherwise very challenging to test, if not impossible, in laboratory experiments. Gravitational wave astronomy aspires to achieve precision measurements with the current and planned network of GW detectors worldwide that are orders of magnitude better than the current detectors. This new era will make the dream of testing theories related to explanations for neutrino mass generation and matter-antimatter asymmetry of the Universe a reality in near future.

ACKNOWLEDGMENTS

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DATA AVAILABILITY

No data were created or analyzed in this study.

APPENDIX: SNR

Gravitational wave interferometers actually measure displacements of detector arms in terms of dimensionless strain-noise $h_{\text{GW}}(f)$, which is related to the GW amplitude. This can be straightforwardly converted into the corresponding energy density [156]

$$\Omega_{\text{exp}}(f)h^2 = \frac{2\pi^2 f^2}{3H_0^2} h_{\text{GW}}(f)^2 h^2, \quad (\text{A1})$$

¹³In $B-L$ extended Dirac leptogenesis one may have LNV.

where $H_0 = h \times 100$ (km/s)/Mpc is the Hubble expansion rate today. The SNR for a projected experimental sensitivity $\Omega_{\text{exp}}(f)h^2$ is then estimated in order to assess the detection probability of the primordial GW background originating from the global cosmic string background following the prescription [157,158]

$$\text{SNR} \equiv \sqrt{\tau \int_{f_{\min}}^{f_{\max}} df \left(\frac{\Omega_{\text{GW}}(f)h^2}{\Omega_{\text{exp}}(f)h^2} \right)^2}, \quad (\text{A2})$$

where $h = 0.7$ and $\tau = 4$ years is the observation time. Usually this is chosen to be $\text{SNR} \geq 10$ as the detection threshold for each individual GW detector.

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