

Probing flavored regimes of leptogenesis with gravitational waves from cosmic strings

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Cosmic strings radiate detectable gravitational waves in models featuring high-scale symmetry breaking, e.g., high-scale leptogenesis. In this Letter, for the first time, we show that different flavored regimes of high-scale leptogenesis can be tested with the spectral features in cosmic string-radiated gravitational waves. This is possible if the scalar field that makes right-handed neutrinos massive is feebly coupled to the Standard Model Higgs. Each flavored regime, sensitive to low-energy neutrino experiments, leaves a marked imprint on the gravitational waves spectrum. A three-flavor and a two-flavor regime could be probed by a characteristic fall-off of the gravitational wave spectrum at the LISA-DECIGO-ET frequency bands with preceding scale-invariant amplitudes bounded from above and below. We present gravitational waves windows for flavored regimes of leptogenesis testable in the upcoming experiments. We also provide the first construction of a leptogenesis framework where a testable distinction of flavor regimes is possible without constraining the flavor structure of the theory.

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Introduction. Within the framework of the seesaw mechanism of light neutrino masses, leptogenesis [1] is a simple process for obtaining the observed baryon asymmetry of the Universe [2]. In this process, CP asymmetric and out-of-equilibrium decays of heavy right-handed neutrinos (RHNs) to lepton and the Standard Model (SM) Higgs doublets, create a lepton asymmetry, which is then processed to baryon asymmetry by sphalerons [3]. In the thermal leptogenesis scenario, the lepton asymmetry is produced at a temperature $T_{\text{lepto}} \sim M_N$, where M_N is the mass of the decaying RHN—also referred to as the scale of thermal leptogenesis [4,5]. Despite being a simple mechanism, leptogenesis operating on scales $M_N > \mathcal{O}(\text{TeV})$ is not directly testable, e.g., with collider experiments.

Accurate incorporation of flavor effects [6–8] nonetheless opens up the opportunity to test high-scale leptogenesis with low-energy neutrino observables, specifically with the leptonic CP violation [9]—a measurement of which is one of the primary goals of the ongoing neutrino experiments such as T2K [10] and NO ν A [11]. Depending on the interaction strengths of the charged lepton flavors, there exist three distinct regimes of high-scale leptogenesis: $M_N \gtrsim 10^{12}$ GeV (one-flavor/vanilla, 1FL), 10^9 GeV $\lesssim M_N \lesssim 10^{12}$ GeV (two-flavor, 2FL), and $M_N \lesssim 10^9$ GeV (three-flavor, 3FL). While the 1FL regime is not sensitive to neutrino mixing parameters, the 2FL and 3FL regimes can offer low-energy neutrino phenomenology, which however may not be unequivocal, unless the flavor structure of the theory is constrained by invoking symmetries (see, e.g., Refs. [12,13]).

Lately, there have been efforts to obtain a direct possible signature of high-scale leptogenesis with primordial stochastic gravitational waves (GWs) background by focusing on the RHN sector. This is promising because, unlike electromagnetic radiation, GWs traverse the early Universe unimpeded with unerring information about their origin. Therefore, no matter how high the leptogenesis scale is, it can be probed with GWs. The idea arises from contemplating the origin of RHN mass at a more fundamental

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level [14–17]. Indeed, in Ref. [14] it has been argued that the RHNs become massive because of the spontaneous breaking of a $U(1)_{B-L}$ gauge invariance, which many grand unified theories (GUT) might embed [18–21]. Besides generating RHN mass, a gauged $U(1)_{B-L}$ breaking inevitably produces cosmic strings (here we consider only stable cosmic strings [22]) that radiate gravitational waves, which may be a signature of leptogenesis. Let us consider the $U(1)_{B-L}$ invariant RHN mass term $f_N \overline{N}_R \Phi N_R^C$, where f_N is the Yukawa coupling, N_R is the RHN field, and Φ is the $B-L$ scalar with vacuum expectation value v_Φ . After the $U(1)_{B-L}$ breaking, the RHNs obtain mass given by $M_N = f_N v_\Phi$, cosmic strings are formed [23–25] and radiate GWs with a scale-invariant spectrum at higher frequencies with an overall amplitude $\Omega_{\text{GW}} \propto v_\Phi$ [26,27]. While the scenario is interesting, note however, that the GW spectrum being dependent only on v_Φ is not so sensitive to M_N because of the additional free parameter f_N . As such, a GW spectrum for $v_\Phi = 10^{14}$ GeV may correspond to 3FL for $f_N = 10^{-6}$ as well as 2FL for $f_N = 10^{-4}$. Therefore, a more realistic GW signature of leptogenesis requires dependence on both parameters v_Φ and f_N , and hence also on M_N .

In this Letter, we show that the M_N -dependent GW spectrum and a concrete link between GWs and high-scale leptogenesis can be achieved in the $U(1)_{B-L}$ model with a peculiar scalar field dynamics [28]. After the phase transition, Φ rolls down to its true vacuum and oscillates coherently to provide an early matter-domination when it is long-lived. The lifetime of the scalar field is determined by v_Φ and f_N , and therefore, M_N controls the beginning and the end of the matter-domination. Differently from Ref. [14], the amplitude and the spectral features of the GWs radiated from cosmic strings are now imprinted by the M_N -dependent matter domination, thus carrying information on M_N , and therefore on different regimes of flavored leptogenesis. In particular, the scale invariance in the GW spectrum at higher frequencies breaks at a frequency f_{dec} , which is determined by the end of the matter domination at the temperature $T_{\text{dec}}(M_N)$ (see Fig. 1). We point out that such a scenario allows us to probe the RHN mass scale markedly, in agreement with all the relevant constraints on the parameter space (see Fig. 2). In particular, the prominent ones are: i) the bound from the recent observations of nHz GWs by pulsar timing arrays (PTA) [29–33], which disfavor GWs from stable cosmic strings [34]; ii) the condition that the cosmic strings radiate predominantly GWs rather than particles [35,36]; iii) the requirement that the Φ particles decay before the big-bang-nucleosynthesis (BBN) [37,38]; and iv) the scalar field Φ is long-lived. We highlight that if the break of the GW spectrum occurs in the LISA (DECIGO) frequency band, it is a signature (most likely signature) of a 3FL regime (see Fig. 3). On the other hand, a break of scale invariance in ET would correspond to both regimes, but with less signal strength for a 3FL.

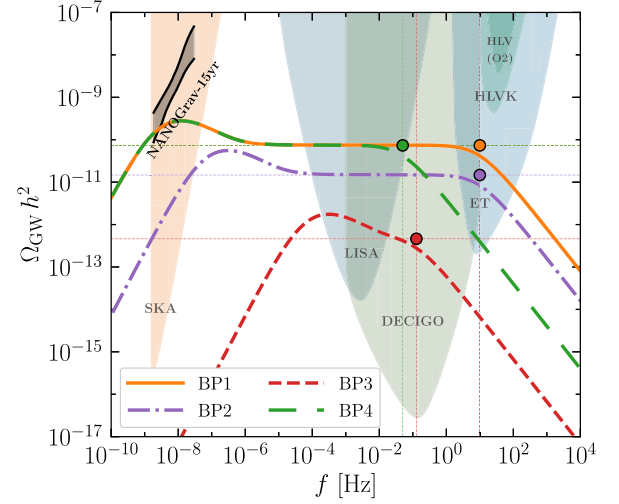


FIG. 1. GW spectrum from the cosmic strings predicted by our model for four benchmark cases with $g' = 10^{-4}$: BP1 (2FL), $M_N = 1.6 \times 10^9$ GeV, and $f_N = 1.4 \times 10^{-5}$; BP2 (3FL), 4.2×10^8 GeV, and $f_N = 1.8 \times 10^{-5}$; BP3 (3FL), 5.2×10^6 GeV, and $f_N = 6.8 \times 10^{-6}$; BP4 (3FL), 2.7×10^8 GeV, and $f_N = 2.2 \times 10^{-6}$. The points show the calculations of the GW plateau amplitude $\Omega_{\text{GW}}^{\text{pl}}$ and the break frequency f_{dec} according to Eqs. (8) and (9), respectively. BP1 and BP2 correspond to equal $\Omega_{\text{GW}}^{\text{pl}}$ but different f_{dec} , while BP1 and BP4 to equal f_{dec} but different $\Omega_{\text{GW}}^{\text{pl}}$. BP3 represents the weakest allowed signal for $g' = 10^{-4}$ (see Fig. 2).

Remarkably, the amplitudes and the spectral break frequencies of the GW spectra predicted by our model are very well delimited and they are all within reach of planned gravitational wave experiments such as LISA, DECIGO, and

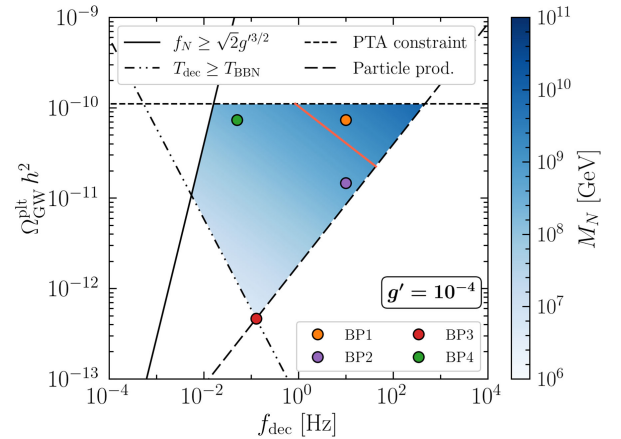


FIG. 2. Allowed parameter space (GWFR window) in the $f_{\text{dec}} - \Omega_{\text{GW}}^{\text{pl}} h^2$ plane for $g' = 10^{-4}$ with the leptogenesis scale M_N color coded. The different black lines display the constraints according to the discussed model requirements and the PTA measurements. The different points correspond to the benchmark cases displayed in Fig. 1. The orange solid line marks the transition between three-flavor and two-flavor leptogenesis regimes.

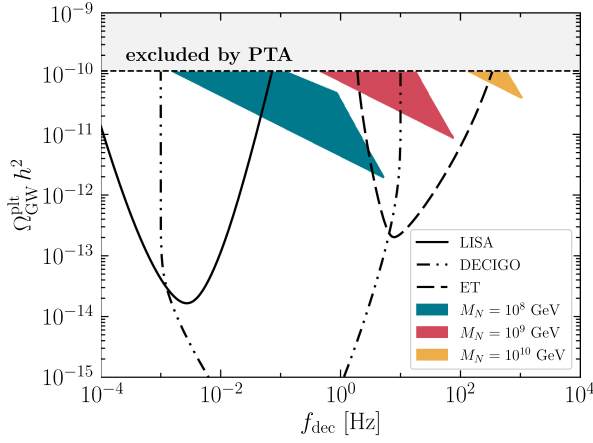


FIG. 3. GWFRL windows for different values of the leptogenesis scale M_N , scanning over the model parameters g' and f_N . The black lines show the sensitivity of LISA, DECIGO, and ET detectors (black lines). The grey shaded region is excluded by current PTA data.

ET [39–41]. Our results not only ensure testable GW signatures of different flavor regimes but also imply that the latter can be probed without constraining the flavor space of the leptogenesis Lagrangian.

The M_N -dependent matter domination and constraints. The seesaw Lagrangian for neutrino masses and leptogenesis is given by

$$-\Delta\mathcal{L} \subset f_{D\alpha} \overline{N}_{Ri} \tilde{H}^\dagger L_\alpha + \frac{1}{2} f_{Ni} \overline{N}_{Ri} \Phi N_{Ri}^C + \text{H.c.}, \quad (1)$$

where f_D is the neutrino Dirac Yukawa coupling and $H(L)$ is the SM Higgs (lepton doublet). The fields N_R and L_α have $B-L$ charge -1 , whereas the same charges for Φ and H are 2 and 0 , respectively. The scalar field dynamics is dictated by the finite temperature potential restoring the symmetry at higher temperatures given by [42,43]

$$V(\Phi, T) = \frac{\lambda}{4} \Phi^4 + D(T^2 - T_0^2) \Phi^2 - E T \Phi^3, \quad (2)$$

where D , E and T_0 are functions of $B-L$ gauge coupling g' , the self-interaction coupling λ , and the vacuum expectation value $v_\Phi = m/\sqrt{\lambda}$, which is determined by the zero temperature potential $V(\Phi, 0) = -\frac{m^2}{2} \Phi^2 + \frac{\lambda}{4} \Phi^4$. The last term in Eq. (2) generates a potential barrier causing a secondary minimum at $\Phi \neq 0$, which becomes degenerate with the $\Phi = 0$ one at $T = T_c$. The potential barrier vanishes at $T_0 (\lesssim T_c)$, and the minimum at $\Phi = 0$ becomes a maximum. The transition from $\Phi = 0$ to $\Phi = v_\Phi$ can be treated as a second-order transition with an extremely quick disappearance of the potential barrier if the order parameter $\Phi_c/T_c \ll 1$. Such a smooth transition can be obtained approximately within the ballpark of $\lambda \simeq g'^3$ and for $g' \lesssim 10^{-2}$ [15,44]. Once

Φ rolls down, it oscillates around v_Φ behaving as matter [28,45,46]. Three important observations are in order. First, the $U(1)_{B-L}$ breaking scale can be as high as the GUT scale Λ_{GUT} , implying $v_\Phi \lesssim 5 \times 10^{15}$ GeV. Second, the Hubble friction at T_c should also be negligible for the rolling of the field down to v_Φ ; this can be achieved by considering m_Φ larger than the Hubble parameter $\mathcal{H}(T_c)$. Third, RHNs become massive after the phase transition at $T = T_c$. Therefore, the scale of leptogenesis is bounded from above as $T_{\text{lepto}} \sim M_N \lesssim T_c \simeq 2\sqrt{g'} v_\Phi$.

We aim to make Φ long-lived to induce a matter-domination epoch with a lifetime determined by v_Φ and f_N . In high-scale leptogenesis scenario, the coupling f_N should be generically large to have M_N comparable with v_Φ . In such a case, if $m_\Phi > 2M_N$, tree-level decay of Φ to RHN pairs ($\Phi \rightarrow NN$) would be too quick to make Φ long lived. This restricts ourselves to the case $M_N > m_\Phi = \sqrt{2\lambda} v_\Phi$. Therefore, along with $M_N \lesssim T_c$, the Φ scalar is long-lived within the following range:

$$m_\Phi < M_N \lesssim T_c \Rightarrow \sqrt{2g'} g' < f_N \lesssim 2\sqrt{g'}, \quad (3)$$

where we have assumed $\lambda = g'^3$. In this case, the mass of Φ is suppressed by a factor $\sqrt{g'}$ compared to the $M_{Z'} = \sqrt{2} g' v_\Phi$, thus kinematically forbidding the tree-level decays of Φ to Z' pairs. Note interestingly that, even if defined to vanish at a very high scale, at lower energies, the seesaw Lagrangian in Eq. (1) can provide a non-vanishing effective coupling $\lambda_{\Phi H, 1\text{-loop}} \sim f_D^2 f_N^2 / (2\pi^2)$ at one loop to trigger $\Phi \rightarrow HH$ decay [28,47,48]. Therefore, despite f_N being large, we can have a lower decay rate proportional to $\lambda_{\Phi H, 1\text{-loop}}^2$. This decay involves both the RHN mass and the active neutrino mass m_ν via the Dirac Yukawa coupling $f_D = \sqrt{m_\nu f_N v_\Phi / v_H^2}$. In our discussion, we consider this radiative $\Phi \rightarrow HH$ decay with the assumption that $\lambda_{\Phi H, 1\text{-loop}} \gg \lambda_{\Phi H, \text{tree}}$ (this the crucial theoretical identification in the $U(1)_{B-L}$ model, which makes the GW spectrum sensitive to M_N). The decay rate is given by [28,47,48]

$$\Gamma_\Phi^{HH} \simeq \Gamma_0 \frac{f_N^6}{\lambda} \left(\frac{v_\Phi}{10^{13} \text{ GeV}} \right)^2 \left(\frac{m_\Phi}{10^8 \text{ GeV}} \right) \ln^2 \left(\frac{\Lambda}{\tilde{\Lambda}} \right), \quad (4)$$

where the running of the one-loop effective coupling is determined by the UV cut-off scale, e.g., $\Lambda \simeq \Lambda_{\text{GUT}}$ and the relevant physical scale $\tilde{\Lambda}$. We have $\Gamma_0 = 1.3 \times 10^{-2} \text{ GeV}$ assuming $m_\nu \simeq 0.01 \text{ eV}$ and the SM Higgs vacuum expectation value $v_H = 174 \text{ GeV}$.

The evolution of the energy density of Φ and the duration of the matter-dominated epoch, which ends when Φ decays [28], can be tracked by solving the Friedmann equations (see Supplemental Material I [49] for further details). We can compute the decay temperature of the

scalar field by recursively solving the implicit equation $\Gamma_{\Phi}^{HH}(T_{\text{dec}}) \simeq \mathcal{H}(T_{\text{dec}})$ taking $\tilde{\Lambda} \simeq T_{\text{dec}}$ in Eq. (4). The scalar must decay before BBN, i.e., $T_{\text{dec}} \gtrsim T_{\text{BBN}} \sim 10$ MeV. Moreover, because the scalar field is long-lived, it produces entropy, which dilutes any pre-existing relic. Therefore, in the computation of baryon asymmetry, the entropy dilution must be considered, because in this framework, the produced entropy can be as large as 10^6 (see Supplemental Material I [49]). Finally, we mention that we can safely neglect the additional one-loop decay $\Phi \rightarrow f\bar{f}V$ (the two-body decay $\Phi \rightarrow f\bar{f}$ is suppressed due to chirality flip [50]), where f and V are SM fermions and vector bosons. Indeed, this decay dominates over the $\Phi \rightarrow HH$ only for $g' \gtrsim 10^{-2}$ and smaller values of v_{Φ} , which are outside our parameter space.

Imprints of flavored leptogenesis on GWs from cosmic strings. Gravitational waves are radiated from cosmic string loops chopped off from the long strings resulting from the spontaneous breaking of the gauged $U(1)_{B-L}$ [26,27]. Long strings are described by a correlation length $L = \sqrt{\mu/\rho_{\infty}}$, where ρ_{∞} is the long string energy density and μ is the string tension defined as [51]

$$\mu = \pi v_{\Phi}^2 h(\lambda, g') \quad \text{with} \quad h \simeq \frac{1}{\ln(2g^2/\lambda)}, \quad (5)$$

for $\lambda \ll 2g^2$. In the present analysis, we take $\lambda = g^3$ therefore $h \simeq 1/\ln(2/g')$. The time evolution of a radiating loop of initial size $l_i = \alpha t_i$ is given by $l(t) = l_i - \Gamma G\mu(t - t_i)$, where $\Gamma \simeq 50$ [26,27], $\alpha \simeq 0.1$ [52,53], $G = M_{\text{Pl}}^{-2}$, and t_i is the initial time of loop creation. The total energy loss from a loop is decomposed into a set of normal-mode oscillations with frequencies $f_k = 2k/l_k = a(t_0)/a(t)f$, where $k = 1, 2, 3, \dots, k_{\text{max}}$, f is the present-day frequency at t_0 , and a is the scale factor. The total GW energy density is computed by summing all the k modes giving [52,53]

$$\Omega_{\text{GW}} = \sum_{k=1}^{k_{\text{max}}} \frac{2k\mathcal{F}_{\alpha}G\mu^2\Gamma_k}{f\rho_c} \int_{t_i}^{t_0} \left[\frac{a(t)}{a(t_0)} \right]^5 n_{\omega}(t, l_k) dt, \quad (6)$$

where ρ_c is the critical energy density of the Universe, $\mathcal{F}_{\alpha} \simeq 0.1$ is an efficiency factor [52], and $n_{\omega}(t, l_k)$ is the loop number density. The latter can be computed from the velocity-dependent-one-scale model as [54–57]

$$n_{\omega}(t, l_k) = \frac{A_{\beta}}{\alpha} \frac{(\alpha + \Gamma G\mu)^{3(1-\beta)}}{[l_k(t) + \Gamma G\mu t]^{4-3\beta} t^{\beta}}, \quad (7)$$

where $\beta = 2/3(1 + \omega)$ with ω being the equation of state parameter of the Universe, and $A_{\beta} = 5.4$ ($A_{\beta} = 0.39$) for the radiation-dominated (matter-dominated) universe [57]. The quantity $\Gamma_k = \Gamma k^{-\delta}/\zeta(\delta)$ quantifies the emitted power in the k -th mode, with $\delta = 4/3$ ($\delta = 5/3$) for loops containing

cusps (kinks) [58]. The integral in Eq. (6) is subjected to two Heaviside functions $\Theta(t_i - t_{\text{fric}})\Theta(t_i - l_c/\alpha)$, each setting a cutoff at a very high frequency f_* , beyond which the GW spectrum falls as f^{-1} . A similar cutoff owing to inflation-diluted strings [59] is not so relevant to our context because we consider thermal leptogenesis. The quantity t_{fric} represents the time until which the motion of the string network is damped by friction [60], and l_c is a critical length above which GW emission is dominant over particle production as shown by high-resolution numerical simulations [35,36] (see also Ref. [61]). The critical length is $l_c \simeq \delta_w(\Gamma G\mu)^{-\gamma}$, where $\delta_w = (\sqrt{\lambda}v_{\Phi})^{-1}$ is the width of the string, and $\gamma = 2$ ($\gamma = 1$) for loops containing cusps (kinks). We shall neglect the cutoffs due to the kinks and friction as they are weaker than those from cusps. Here we deal with small values of λ leading to “thick” strings, which are not conventional while studying GWs signatures from cosmic strings in Beyond-the-Standard-Model (BSM) models [14]. In this case, the particle production cutoff in the GW spectrum is expected to be much stronger than the one usually considered for “thin” strings ($\lambda \simeq 1$). Thus, our scenario is also the first concrete BSM model that discusses the results of Refs. [35,36] for small λ . Note however that we generalized the simulation results obtained for large couplings (λ, g') by replacing $\delta_w = (\sqrt{\lambda}v_{\Phi})^{-1}$ in l_c . But in principle, Eq. (11) of Ref. [35] would exhibit coupling dependence, which is not trivial to deduce without numerical simulation. Therefore, the particle production cutoff discussed above should be considered conservative. Note however that the strings predicted in this model effectively evolve as thin strings because the condition $\delta_w \ll \mathcal{H}^{-1}$ is always satisfied (see Supplemental Material II [49]). In the following, we present the results for the dominant $k = 1$ mode, which captures the qualitative features of our analysis.

In Fig. 1, we show the GW spectrum from Eq. (6) (for $k = 1$) predicted by our model for four benchmark scenarios mentioned in the caption. Without any intermediate matter-dominated epoch, the GW spectrum has two main features: *i*) a low-frequency peak owing to the GW radiation from the loops that originate in the radiation epoch and decay in the standard matter epoch, and *ii*) an almost scale-invariant plateau at high frequencies as [52,53,62]

$$\Omega_{\text{GW}}^{\text{plt}} = \frac{128\pi\mathcal{F}_{\alpha}G\mu A_R}{9\zeta(\delta)} \frac{A_R}{\epsilon_R} \Omega_R[(1 + \epsilon_R)^{3/2} - 1], \quad (8)$$

which arises from loop dynamics in the radiation-dominated epoch only. In Eq. (8), $\epsilon_R = \alpha/\Gamma G\mu \gg 1$, $A_R \equiv A_{\beta} \simeq 5.4$, and $\Omega_R \sim 9 \times 10^{-5}$. Note that $\Omega_{\text{GW}}^{\text{plt}} \propto \sqrt{\mu} \propto v_{\Phi}$, implying the larger the symmetry breaking scale, the stronger the amplitude of GW. In our scenario, which features instead an intermediate matter-dominated epoch, the plateau breaks at a high frequency f_{dec} , beyond which the spectrum falls as $\Omega_{\text{GW}}(f > f_{\text{dec}}) \propto f^{-1}$ (see Fig. 1). The

analytical expression for this spectral break frequency f_{dec} can be estimated as [63]

$$f_{\text{dec}} \simeq 0.45 \text{ Hz} \left(\frac{10^{-12}}{G\mu} \right)^{1/2} \left(\frac{T_{\text{dec}}}{\text{GeV}} \right). \quad (9)$$

To robustly claim the spectral fall as a signature of M_N -dependent matter domination, we must require $t_{\text{dec}} > l_c/\alpha$. This is equivalent to imposing a constraint on the parameter space such that $T_{\text{cusp}}/T_{\text{dec}} > 1$, where T_{cusp} can be derived as

$$T_{\text{cusp}} \simeq 25 \text{ GeV} \left(\frac{G\mu}{10^{-12}} \right)^{5/4} \left(\frac{g'}{10^{-4}} \right)^{3/4} h(g')^{-1/4}. \quad (10)$$

For loops containing kinks, one can obtain a similar constraint, which however is less stringent than Eq. (10). The parameter space of our model is also constrained by the recent PTA observation of stochastic GWs, which disfavors stable cosmic strings due to a different spectral slope [34]. Thus, PTA data require $\Omega_{\text{GW}} \lesssim \Omega_{\text{GW}}^{\text{PTA}}$ at $f = 30 \text{ nHz}$ —giving $G\mu \lesssim 7.0 \times 10^{-11}$.

We note that there could be an additional contribution to the GW spectrum at very high frequencies from the cosmic string loops in the first radiation epoch before Φ starts to dominate. However, the amplitude at higher frequencies is suppressed by the entropy production and the particle production cutoff.

Results and discussion. In Fig. 2, we show the allowed parameter space for $g' = 10^{-4}$, in agreement with all the model conditions and constraints discussed above. For a given gauge coupling, the spectral break frequency f_{dec} and the GW plateau amplitude $\Omega_{\text{GW}}^{\text{plt}}$ are uniquely determined by the model parameters M_N and f_N . Usually, they are limited from above by the PTA data. However, in our model, they are also bounded from below by the BBN constraint on the matter-dominated epoch and by the requirement that the spectral break is not due to the particle production from cosmic strings. Scanning over different values of the gauge coupling (see Supplemental Material III [49] for the gauge coupling dependence of the parameter space), we find that our scenario implies $10^{-3} \lesssim f_{\text{dec}}/\text{Hz} \lesssim 10^3$ and $10^{-13} \lesssim \Omega_{\text{GW}}^{\text{plt}} h^2 \lesssim 10^{-10}$, which we define as the gravitational waves windows for flavored regimes of leptogenesis (GWFRL). These ranges, within which a robust GW signature of high-scale leptogenesis is achieved, will be fully explored by planned GW experiments such as LISA, DECIGO, and ET. These GWFRL windows also trivially satisfy two important constraints. First, $G\mu < 6.3 \times 10^{-3} g' h(g')$ due to the Hubble friction being always less than m_Φ at $T \simeq T_c$. Second, the ratio of the Φ vacuum energy to the radiation at T_c : $\rho_\Phi(T_c)/\rho_R(T_c) = 4.7 \times 10^{-4} g'$ is always smaller than one, thus avoiding a second period of inflation.

In Fig. 3, we report the allowed regions for f_{dec} and $\Omega_{\text{GW}}^{\text{plt}}$, once we fix the value of the leptogenesis scale M_N (regions with different colors) and take the gauge coupling to vary from 10^{-7} to 10^{-1} . The black lines (from left to right) represent the sensitivity of LISA, DECIGO, and ET detectors, respectively. Irrespective of the values of the gauge coupling, going from light to heavy M_N , the maximum allowed value for the spectral break frequency f_{dec} increases. Remarkably, this trend allows us to point out that a spectral break in the LISA (DECIGO) frequency band would be (most likely) a striking signature of the 3FL leptogenesis regime ($M_N \leq 10^9 \text{ GeV}$). On the other hand, at the ET frequency band, both the 3FL and 2FL regimes leave their signatures distinct in signal strength, with a smaller GW amplitude for the 3FL case. As shown with the BP1 (2FL) and BP2 (3FL) cases, despite corresponding to the same f_{dec} , they have different plateau amplitudes (see Fig. 1). On the other hand, the 2FL regime (corresponding to $g' \gtrsim 10^{-3}$) is highly disfavored by the PTA data, which already excluded the 1FL case.

The GWFRL windows obtained in the present analysis are robust at the leading order, nevertheless, they can be further improved in future works. We have considered $\lambda = g'^3$, but a more rigorous parameter space scan could be performed to find a smooth transition of the field while being consistent with other gauge coupling dependent constraints (see Supplemental Material [49] for a brief discussion). Moreover, one can account for the GW emission from the higher k modes. However, we find that within an instantaneous scaling network (instantaneous change in A_β during the radiation-to-matter transition), summing a large number of modes makes the PTA constraint stronger by a factor of $\sim \zeta(4/3)$ and lowers the f_{dec} by a factor of ~ 2.5 . The latter has been derived by taking a 10% deviation from the near scale-invariant spectrum. In addition, the spectrum falls as $f^{-1/3}$ instead of f^{-1} at very high frequencies [15,17] with no impact on our results based on the detectability of the signal at lower frequencies.

In conclusion, we have demonstrated that different flavor regimes of high-scale leptogenesis could produce GW signatures with spectral features depending on the mass scale of the heavy neutrinos. This is an interesting aspect of the simple realization of the seesaw model based on $U(1)_{B-L}$ symmetry such that different flavored regimes could be probed without invoking any flavor symmetry to reduce the number of free parameters in the seesaw Lagrangian. On the other hand, we have established synergy between primordial gravitational waves and low-energy neutrino physics in seesaw models with flavor symmetries. It is important to emphasize that the predicted GW signal should not be regarded as a definitive indicator, as similar spectral features could arise in other BSM models involving matter domination. Furthermore, in this work, we do not delve into the high-frequency behavior of the GW

spectra pertaining to different flavor regimes, and neither do we perform a statistical analysis to address uncertainties in reconstructing weak GW signals (e.g., BP3 in Fig. 1). These aspects will be thoroughly explored in future works.

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