

Exploring leptogenesis in the era of first order electroweak phase transition

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We present a novel approach for implementing baryogenesis via leptogenesis at low scale within the neutrino seesaw framework where a sufficient lepton asymmetry can be generated via out-of-equilibrium CP -violating decays of right-handed neutrinos (RHNs) even when their mass falls below the Standard Model (SM) Higgs mass. It becomes possible by keeping the sphaleron in equilibrium below its conventional decoupling temperature $T_{\text{sp}}^{\text{SM}} \sim 131.7$ GeV in SM so as to facilitate the conversion of lepton asymmetry to baryon asymmetry at such a low scale, thanks to the flexibility of the bubble nucleation temperature in case the electroweak phase transition (EWPT) is of first order. The scenario emerges as an exciting (and perhaps unique) possibility for low scale leptogenesis, particularly if the Universe attains a reheating temperature lower than 131.7 GeV. We show that a stochastic gravitational wave, characteristic of such first order EWPT, may be detected in near future detectors, while the presence of RHNs of mass as low as 35 GeV opens up an intriguing detection possibility at current and future accelerator experiments.

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Understanding the baryon asymmetry of the Universe (BAU) [1] remains a core problem in particle physics and cosmology. Among various proposals, generating a lepton asymmetry from the out-of-equilibrium decay of heavy Standard Model (SM) singlet right-handed neutrinos (RHNs) N_i into SM lepton (l_L) and Higgs (H) doublets in the early Universe, called leptogenesis [2–5], and transferring it to the baryon sector via sphaleron process [6–10] manifest itself as the most natural explanation for BAU. This is particularly because of its close connection to the neutrino mass generation mechanism via the Type I seesaw [11–16]. This minimal SM extension not only takes care of expounding the tiny neutrino mass and BAU, but also provides candidate for another unsolved conundrum [17], the dark matter [18,19].

In thermal leptogenesis, the mass of the decaying RHN is typically constrained by the Davidson-Ibarra bound [20] $M_N \gtrsim 10^9$ GeV for hierarchical RHNs. Alongside, with quasidegenerate RHNs, resonant enhancement [21–24] of lepton asymmetry production helps to achieve leptogenesis with lighter RHN masses around electroweak (EW) scale $\mathcal{O}(160$ GeV) [25,26] and above, making such scenarios testable at current and future colliders. The apparent possibility to enhance such detection prospect to next level

by lowering the RHN mass scale further down is limited by the fact that the sphalerons decouple below $T_{\text{sp}}^{\text{SM}} = 131.7$ GeV in the SM [25], preventing lepton-to-baryon asymmetry conversion.

Leptogenesis can also proceed via Higgs decay [27] (considering thermal effects) with (quasidegenerate) RHNs mass below the EW scale, in a typical temperature window between $T_{\text{sp}}^{\text{SM}}$ to EW symmetry breaking temperature $T_{\text{EW}} \sim \mathcal{O}(160$ GeV). Additionally, leptogenesis via oscillations [28] remains viable with RHNs having mass as low as in the GeV regime, occurring at temperatures well above $T_{\text{sp}}^{\text{SM}}$. Recently, we also showed [29] that temperature-dependent RHN masses can yield sufficient lepton asymmetry at high temperatures, even if their zero-temperature masses lie in the GeV regime.

In all scenarios involving sub-EW mass RHNs, leptogenesis requires the temperature of the Universe to exceed $T_{\text{sp}}^{\text{SM}}$. Again, any asymmetry (lepton number [L] or baryon number [B]) must arise in a postinflationary epoch (otherwise, a complete erasure is inevitable), the onset of which is usually (with instantaneous reheating) marked by the reheating temperature T_{RH} . The lower bound on T_{RH} being a few MeV only (from big bang nucleosynthesis [BBN]), an interesting and relevant question emerges: is it feasible to realize leptogenesis when the reheating temperature lies below the $T_{\text{sp}}^{\text{SM}}$, i.e., 131.7 GeV? In such scenarios, with the Universe never reaching $T_{\text{sp}}^{\text{SM}}$, lepton-to-baryon asymmetry conversion is inhibited. Hence, this intriguing possibility remains unexplored in the literature to the best of our knowledge.

In this work, we demonstrate that leptogenesis indeed remains a possibility at such a low scale (i.e., even below

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$T_{\text{sp}}^{\text{SM}}$) provided the EW phase transition (EWPT) is of strongly first order,¹ having a characteristic bubble nucleation temperature T_n , at which one bubble on an average is nucleated per horizon. Although bubbles of broken phase (inside which Higgs vev , $\langle H \rangle = v(T) \neq 0$) begin to form at a certain critical temperature T_c above T_n , they cannot grow beyond a critical size and initiate the conversion of the Universe from the symmetric phase ($\langle H \rangle = 0$) to the broken one until the temperature drops to T_n . Hence, the symmetric phase is prevalent in the entire Universe until T_n . Note that T_n can be kept below $T_{\text{sp}}^{\text{SM}}$ in scenarios with the first order EWPT (FOEWPT) in general. This may have interesting consequence in terms of leptogenesis as in such a circumstance, an RHN having mass M_N within a range $T_n < M_N < T_{\text{sp}}^{\text{SM}}$ can decay out of equilibrium ($N \rightarrow l_L + H$) and produces a lepton asymmetry in the Universe (in symmetric phase) that can still be converted into BAU via sphalerons. This is due to the fact that sphalerons (above T_n and below 131.7 GeV) remaining in the symmetric phase are in equilibrium as the expansion of the broken phase bubbles (beyond a critical size) pervading the Universe are yet to be started. Immediately below T_n , tunneling to the true vacuum ($v(T) \neq 0$) from the false one starts to proceed efficiently, leading to the nucleation of bubbles. During the nucleation, the baryon asymmetry (produced outside the bubble) is engulfed inside with the gradual expansion of the bubble. As the sphaleron rate is exponentially suppressed: $\Gamma_{\text{sph}}^{(\text{in})} \sim T^4 \exp[-8\pi v(T)/g_w T]$ [30] within the true vacuum bubbles, sphalerons decouple immediately and hence, the enclosed baryon asymmetry remains preserved.

Note that the pivotal role in materializing the above idea is played by the nucleation temperature T_n , associated to FOEWPT, being smaller than the sphaleron decoupling temperature of the SM. This can in general be possible within any framework of FOEWPT, making the proposal model independent. Hence, before going to evaluate the parameter space specific to such a possibility, we specify here the background of estimating the T_n first. The FOEWPT occurs through the nucleation of true-vacuum ($\langle H \rangle = v(T) \neq 0$) bubbles. Although broken-phase bubbles can begin to form at the critical temperature T_c , at which false ($\langle H \rangle = 0$) and true minima become degenerate, they cannot grow beyond a critical size and initiate the conversion of the Universe from the symmetric phase ($\langle H \rangle = 0$) to the broken one as their nucleation remain suppressed due to the small false vacuum decay rate. Hence, the symmetric phase is prevalent in the entire Universe till the temperature at which the probability of forming at least one bubble per horizon volume reaches order 1, and the transition from the false to the true vacuum effectively proceeds, initiating bubble nucleation. The

characteristic temperature at which this occurs is known as the nucleation temperature (T_n), which can be determined using the relation [31,32]

$$N_b(T_n) = \int_{T_n}^{T_c} \frac{dT}{T} \frac{\Gamma(T)}{\mathcal{H}(T)^4} = 1, \quad (1)$$

where N_b is the number of bubbles per horizon, and $\Gamma(T) \simeq T^4 \left(\frac{S_3}{2\pi T}\right)^{3/2} \exp(-S_3/T)$ is the false vacuum decay rate [33–35] with S_3 being the three-dimensional Euclidean action for the $O(3)$ -symmetric bounce solution [36], computed numerically using the package FindBounce [45], and $\mathcal{H}(T) = 0.33\sqrt{g_*} \frac{T^2}{M_p}$ ($M_p = 2.4 \times 10^{18}$ GeV is the reduced Planck mass) is the Hubble parameter. Though the exact nucleation temperature depends on the specific model parameters of the extensions of the SM (responsible for FOEWPT) such as scalar [46,47] or fermionic extensions [48], or via the inclusion of a nonrenormalizable dimension-six operator to the SM Higgs potential [49–54], it is in general plausible for T_n to remain significantly lower, even below $T_{\text{sp}}^{\text{SM}}$. In the following, as a case study, we adopt a minimal setup by incorporating a dimension-six operator to the SM Higgs potential to analyze the FOEWPT quantitatively and show to what extent T_n can be lowered in such a picture. Based on such findings, we plan to estimate low scale leptogenesis connected to a T_n below the $T_{\text{sp}}^{\text{SM}}$ to set the lower bound of RHN mass for successful leptogenesis.

To accommodate the proposal described above, where EWPT needs to be strongly first order, we must go beyond the SM, since within the SM, the EWPT remains a smooth crossover with the Higgs mass 125 GeV. From the minimality point of view, we therefore extend the SM Higgs doublet potential by introducing a nonrenormalizable dimension-six operator $(H^\dagger H)^3/\Lambda^2$ as proposed in various studies [49–54], where Λ is the cutoff scale and $H^T = [G_1 + iG_2, \phi + h + iG_3]/\sqrt{2}$ with G_i as the Goldstone bosons and h the physical Higgs field. Here, ϕ corresponds to the background (classical) field, the tree-level potential of which is given by

$$V_{(\phi)}^{\text{tree}} = -\frac{\mu_h^2}{2}\phi^2 + \frac{\lambda_h}{4}\phi^4 + \frac{1}{8}\frac{\phi^6}{\Lambda^2}. \quad (2)$$

In order to analyze the phase transition precisely, it is essential to incorporate the one-loop finite temperature correction V_T into it, which results in an effective potential $V_{\text{eff}} = V_{(\phi)}^{\text{tree}} + V_T(\phi, T)$ for ϕ as [36]

$$V_{\text{eff}} = \left(-\frac{\mu_h^2}{2} + \frac{1}{2}c_h T^2\right)\phi^2 + \left(\frac{\lambda_h}{4} + \frac{\lambda_1}{4}T^2\right)\phi^4 + \frac{1}{8}\frac{\phi^6}{\Lambda^2}, \quad (3)$$

¹Note that for the above mentioned low scale leptogenesis scenarios, the EWPT is considered to be a smooth crossover.

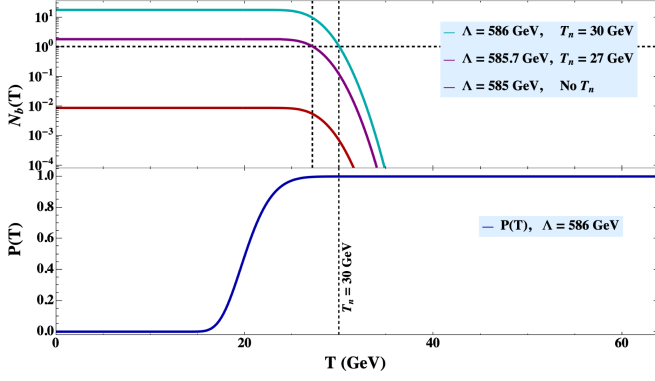


FIG. 1. Number of bubbles per horizon as a function of T .

where

$$c_h = \frac{1}{16} \left(\frac{4m_h^2}{v_0^2} + 3g_w^2 + g_Y^2 + 4y_t^2 - 12 \frac{v_0^2}{\Lambda^2} \right), \quad \lambda_1 = \frac{1}{\Lambda^2},$$

with $m_h = 125$ GeV and $v_0 = 246$ GeV as the Higgs mass and Higgs vev at zero-temperature, respectively. At this stage, we restrict ourselves with terms up to order T^2 and hence, daisy diagrams [55–58] are not incorporated. The one-loop Coleman-Weinberg correction [59] being negligible compared to the thermal correction [60], also is not included. Note that we adopt such approximated form of the effective Higgs potential, Eq. (3), to provide a simplified illustration of how the upper and lower bound of Λ can be obtained and how the lowest T_n for successful first order electroweak phase transition can be determined. However, in the subsequent part of analysis to demonstrate the phenomena of bubble nucleation and the transition of space volume from the false vacuum to the true vacuum phase, we employ the full effective thermal potential, including daisy resummation [36] to accurately estimate the final parameter space of our study.

As seen from the effective potential in Eq. (3), two degenerate minima form at the critical temperature T_c (corresponding to a particular Λ), separated by a potential barrier in between. The parameter Λ controls the barrier height between the local and global minima [36]. Increasing Λ diminishes the barrier height, weakening the first order nature of the transition. Employing the full effective thermal potential in the CosmoTransition package [61], we find that for $\Lambda > 810$ GeV, the transition is no longer first order. Thereby, $\Lambda = 810$ GeV serves as the upper limit. Conversely, a lower bound on Λ can be exercised by analyzing the dynamics of bubble nucleation.

In the upper panel of Fig. 1, we illustrate the $N_b(T)$ dependence on temperature, for three closely spaced values of Λ where the corresponding T_n values [as obtained from Eq. (1) using the high-temperature approximated potential in Eq. (3)] are indicated by the vertical black dashed lines. It turns out that for $\Lambda = 585$ GeV, N_b never reaches $\mathcal{O}(1)$, indicating no bubble nucleation and hence an incomplete

phase transition. In contrast, for $\Lambda \gtrsim 585.7$ GeV, the condition $N_b \geq 1$ is satisfied, revealing that the false vacuum decay rate surpasses the expansion rate of the Universe, i.e., $\Gamma(T) \gtrsim \mathcal{H}(T)^4$ at a specific T_n depending on Λ .

Upon incorporating the full effective thermal potential, including both the Coleman-Weinberg and daisy corrections, the above analysis establishes the lower limit $\Lambda = 571.6$ GeV. Combining this with the earlier upper bound from barrier height perspective, the viable range of Λ and correspondingly $T_n(\Lambda)$ (and T_c , too) for a strongly FOEWPT is

$$571.6 \text{ GeV} \lesssim \Lambda \lesssim 810 \text{ GeV}; \quad 33.8 \text{ GeV} \lesssim T_n \lesssim 119 \text{ GeV}, \quad (4)$$

where the criteria for successful bubble nucleation and the completion of the phase transition are satisfied.

Observing that $T_n(T_c)$ can be as low as 34 (79) GeV for a strongly FOEWPT, we now explore the implications of such low T_n on the sphaleron decoupling. Since the sphaleron rate remains unsuppressed in an environment of vanishing the SM Higgs vev (i.e., the Universe stays in a false vacuum), keeping the sphalerons in equilibrium, it is pertinent to understand the fate of the false vacuum around T_c and/or T_n . Unlike the smooth crossover in the SM, where the symmetric phase persists up to 160 GeV, here the Universe can remain trapped in the false vacuum until T_n when broken-phase bubbles begin to envisage the false vacuum space to convert it into the true vacuum inside which sphalerons decouple. To illustrate this, we compute the probability $P(T)$ of a spatial point remaining in the false vacuum at temperature T , expressed as [31,32]

$$P(T) = e^{-I(T)},$$

$$I(T) = \frac{4\pi}{3} \int_T^{T_c} dT' \frac{\Gamma(T')}{T'^4 \mathcal{H}(T')} \left(\int_T^{T'} d\tilde{T} \frac{v_w}{\mathcal{H}(\tilde{T})} \right)^3, \quad (5)$$

where $v_w (< c_s, \text{ the sound speed})$ denotes the bubble wall velocity.² The temperature dependence of $P(T)$ is shown in the lower panel of Fig. 1 for $\Lambda = 586$ GeV using the potential in Eq. (3). The result $P(T_n) \simeq 1$ clearly demonstrates that the Universe continues to exist in the false vacuum down to the nucleation temperature T_n . In particular, for our case, the Universe remains in the false vacuum down to the lowest temperature $T_n^{\min} \simeq 34$ GeV [Eq. (4)], keeping the sphalerons active and in equilibrium even at $\mathcal{O}(34 \text{ GeV})$.

²Determining the bubble wall velocity requires solving the coupled scalar field and plasma transport equations with hydrodynamic matching [62,63], which is beyond the scope of this work. We therefore treat v_w as a phenomenological parameter and assume a subsonic value [64]. Since the lepton asymmetry is generated before bubble nucleation in our framework, the baryon asymmetry is essentially independent of v_w , which mainly impacts the gravitational wave prediction.

The emergence of such a low sphaleron decoupling temperature in the context of FOEWPT, compared to $T_{\text{sp}}^{\text{SM}}$, while the Universe stays in the false minima is the key observation of our study, which has far-reaching implications for low scale leptogenesis. Firstly, the Universe being in false vacuum till T_n ($< T_{\text{sp}}^{\text{SM}}$), the SM fields are massless (apart from their thermal masses), and hence the RHNs of mass $M_N > T_n$ can decay out of equilibrium to lepton and Higgs doublets (via neutrino Yukawa interaction). This requires the satisfaction of the condition, $M_N > M_L(T) + M_H(T)$ where $M_L(T)$ and $M_H(T)$ represent the thermal masses of lepton and Higgs doublets, respectively, with $M_L(T) + M_H(T) \simeq 0.77T$. Hence, such decay would take place close to an RHN mass-equivalent temperature $T \sim M_N$, which automatically satisfies $M_N > 0.77T$. The exact temperature (or epoch) of occurrence of such RHN decay needs to be estimated by solving the Boltzmann equations, as we have established later in this work. Secondly, the sphalerons are able to convert the lepton asymmetry to a baryon one until a temperature at or above T_n . For example, T_n^{min} being 34 GeV in our scenario, the standard (resonant) leptogenesis can easily take place via the out-of-equilibrium decay of (quasidegenerate) RHNs having mass above T_n^{min} but below the SM gauge boson's masses $m_{W,Z}$ ($T_n^{\text{min}} < M_N < m_{W,Z}$) in the temperature window $T_n^{\text{min}} - M_N$, which can explain BAU. Note that producing the correct baryon asymmetry through RHNs out-of-equilibrium decay at such low scale would otherwise remain highly challenging for $M_N < T_{\text{sp}}^{\text{SM}} \sim 131.7$ GeV, even with resonant leptogenesis as conversion of lepton to baryon asymmetry is effectively switched off below $T_{\text{sp}}^{\text{SM}}$ in the regime of smooth crossover EWPT. Interestingly, our framework is still capable of generating baryon asymmetry in case the Universe attains a low reheating temperature³ T_{RH} below the $T_{\text{sp}}^{\text{SM}}$, quite plausible as the lower bound on reheating temperature is only a few MeV from BBN [67–70], contrary to other low scale leptogenesis scenarios such as Higgs-decay leptogenesis [27], leptogenesis via oscillations [28,29] where T_{RH} requires to be higher than $T_{\text{sp}}^{\text{SM}}$. It is worth noting that, in our framework, the nucleation temperature T_n sets a lower bound on the RHN mass M_N to reproduce the observed baryon asymmetry, while the cutoff scale Λ , which also affects T_n , imposes an upper bound on M_N from the requirement that the effective framework (responsible for rendering the electroweak phase transition strongly first order via dimension-six Higgs operator) remains valid.

To proceed in estimating the lepton asymmetry with such a low sphaleron decoupling temperature, we begin

³In case of noninstantaneous reheating [65,66], the maximum temperature T_{Max} being more than T_{RH} , we expect the present scenario would work with T_{RH} close to the BBN bound.

with the usual Type I seesaw Lagrangian⁴ (in the charged lepton and RHN mass diagonal bases), as follows:

$$-\mathcal{L}_1 = \bar{\ell}_{L\alpha}(Y_\nu)_{ai}\tilde{H}N_i + \frac{1}{2}\bar{N}_i^c M_i N_i + \text{H.c.}, \quad (6)$$

where $i = 1, 2$ (for minimal scenario) and $\alpha = e, \mu, \tau$ in general. With two quasidegenerate RHNs, one can evaluate the CP asymmetry,

$$\varepsilon_\ell^i = \sum_{j \neq i} \frac{\text{Im}(Y_\nu^\dagger Y_\nu)_{ij}^2}{(Y_\nu^\dagger Y_\nu)_{ii}(Y_\nu^\dagger Y_\nu)_{jj}} \frac{[M_i^2 - M_j^2]M_i\Gamma_{N_j}}{[M_i^2 - M_j^2]^2 + M_i^2\Gamma_{N_j}^2}, \quad (7)$$

which can be $\sim \mathcal{O}(1)$ if the resonance condition $\Delta M = M_2 - M_1 \sim \Gamma_{N_1}/2$ is satisfied. Here Γ_{N_1} is the decay rate of N_1 to lepton and Higgs doublets. Using Casas-Ibarra parametrization [71], the Y_ν matrix can be constructed using: $Y_\nu = -i\frac{\sqrt{2}}{v_0}UD\sqrt{m}\mathbf{R}D\sqrt{M}$ where U is the Pontecorvo-Maki-Nakagawa-Sakata matrix [72] which connects the flavor basis to the mass basis of light neutrinos. Here $D\sqrt{m} = \text{diag}(\sqrt{m_1}, \sqrt{m_2}, \sqrt{m_3})$ and $D\sqrt{M} = \text{diag}(\sqrt{M_1}, \sqrt{M_2})$ denote the diagonal matrices containing the square root of light neutrino masses and RHN masses, respectively, and $\mathbf{R}(\theta_R)$ represents a complex orthogonal matrix.

The Boltzmann equations for the abundance of RHNs ($Y_N = n_N/s$) and the yield of B-L asymmetry (Y_{B-L}) can be written as [2,73–75]

$$s\mathcal{H}z \frac{dY_{N_i}}{dz} = -\left(\frac{Y_{N_i}}{Y_{N_i}^{\text{eq}}} - 1\right)(\gamma_D^i + 2\gamma_S^i + 4\gamma_t^i), \quad (8)$$

$$s\mathcal{H}z \frac{dY_{B-L}}{dz} = \sum_{i=1}^2 \left[-\varepsilon_\ell^i \left(\frac{Y_{N_i}}{Y_{N_i}^{\text{eq}}} - 1 \right) \gamma_D^i - \frac{Y_{B-L}}{Y_{B-L}^{\text{eq}}} \right. \\ \left. \times \left(2\gamma_S^i + \gamma_S^i \frac{Y_{N_i}}{Y_{N_i}^{\text{eq}}} + 2\gamma_N \right) \right], \quad (9)$$

where, $z = M_1/T$, s is the entropy density, and $\gamma_D^i = n_{N_i}^{\text{eq}} \frac{K_1(M_i/T)}{K_2(M_i/T)} \Gamma_{N_i}$. Here, γ_S^i and γ_t^i denote the reaction rate densities for $\Delta L = 1$ scatterings (s and t channels, respectively) while γ_N represents the reaction rate density of $\Delta L = 2$ scattering processes [2,36,73–75]. Together with decays and inverse decays, these processes account for the relevant washout effects. Considering $T_n(\Lambda) \lesssim M_i < \Lambda$ (as discussed above) while using best fit values of the neutrino mixing angles and mass-square differences with $m_1 = 0$ for

⁴The nonrenormalizable *explicit* lepton-number breaking operator $c\ell_L\ell_L HH/\Lambda_L$ can in principle also be present. However, considering only the *soft* breaking of the lepton number as present in the Majorana mass of the RHNs in the Type I seesaw, the effect of this term on neutrino mass can be ignored by considering the coefficient c to be small enough to be justified with a UV complete picture that is beyond the present discussion.

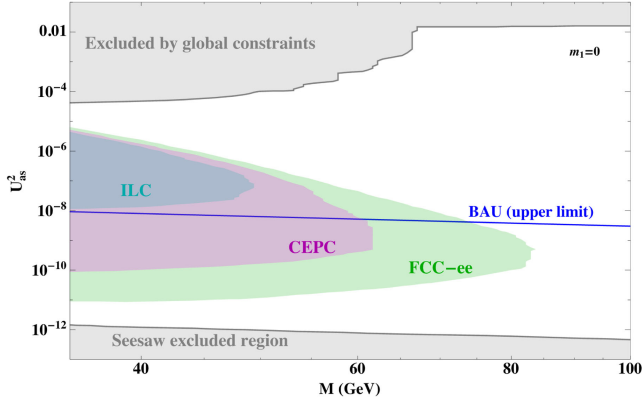


FIG. 2. Viable parameter space for leptogenesis and sensitivity regions of future lepton colliders [76] with the upper gray region excluded by global constraints [77] from direct and indirect search experiments.

normal hierarchy for getting Y_ν , we solve the above equations numerically with thermalized RHNs as the initial condition. It is found that the correct BAU corresponding to $M_1 \in [35, 100]$ GeV can be obtained by varying $\Delta M \in [3.16 \times 10^{-11}, 1.07 \times 10^{-7}]$ and $\text{Im}(\theta_R) \in [-4, -0.2]$ with $\text{Re}(\theta_R) = 0.8$. Correspondingly, the range of neutrino Yukawa coupling (specifically, the largest entry of $|Y_\nu|$) is found to be $[9.41 \times 10^{-8}, 8.75 \times 10^{-6}]$, while the RHN mass varies within the above specified range. The $\text{Im}(\theta_R)$ dependence can be conventionally parametrized by the active-sterile mixing angle, $|\Theta_{ai}|^2 = |(Y_\nu)_{ai}|^2 v_0^2 / M_i^2$. The result is depicted in Fig. 2 in the $M-U_{as}^2$ plane, with $M = (M_1 + M_2)/2$ and $U_{as}^2 = \sum_{i,\alpha} |\Theta_{ai}|^2$, where the blue solid line corresponds to the upper limit of U_{as}^2 for correct BAU (as well as neutrino oscillation data) while the region below it stands for more than required BAU, which can, however, be brought down to correct asymmetry easily by appropriate ΔM and θ_R .

Note that inclusion of $(H^\dagger H)^3 / \Lambda^2$ in the SM Higgs potential modifies the triple Higgs coupling as

$$\lambda_3 = \frac{1}{6} \frac{d^3 V_{\text{eff}}(\phi, T=0)}{d\phi^3} \Big|_{\phi=v_0} = \lambda_3^{\text{SM}} + \frac{v_0^3}{\Lambda^2}, \quad (10)$$

where $\lambda_3^{\text{SM}} = m_h^2 / (2v_0)$. Current bounds from the ATLAS and CMS experiments at 95% CL, expressed in terms of $k_3 = \lambda_3 / \lambda_3^{\text{SM}}$, are given by $k_3 \in [-1.2, 7.2]$ ATLAS [78], $k_3 \in [-1.2, 7.5]$ CMS [79]. The lowest value of Λ used in our analysis is 571.6 GeV, which corresponds to $k_3 = 2.44$, which lies well within the current experimental bounds of k_3 as reported by the ATLAS and CMS collaborations. Analogously, the upper limit of Λ (810 GeV) obtained in our framework results to $k_3 = 1.72$, hence falls within the allowed range. However, the High Luminosity (HL)-LHC is expected to constrain λ_3 within 40% of the SM value at 68% CL [80–82], providing an indirect collider probe of Λ and leptogenesis. We present $\Delta\lambda_3 = (\lambda_3 - \lambda_3^{\text{SM}}) / \lambda_3^{\text{SM}}$ as a

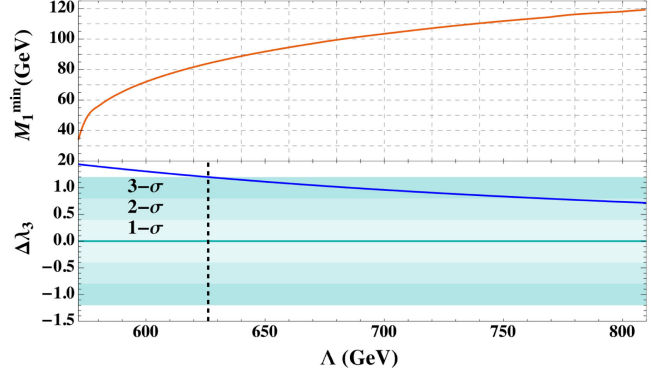


FIG. 3. Variation of $\Delta\lambda_3$ (bottom panel) and M_1^{min} (upper panel) with Λ .

function of the cutoff scale Λ , in the range of our interest from FOEWPT and leptogenesis, as shown in Fig. 3 (bottom panel), with the HL-LHC experimental reach at 1σ , 2σ , and 3σ indicated by the colored shaded regions. Since the value of Λ is intricately connected to the T_n playing a significant role in realizing low scale leptogenesis with a minimum mass of RHN M_1^{min} as shown in the upper panel of Fig. 3, this serves as an interesting future probe of the low scale leptogenesis, particularly for $\Lambda \gtrsim 625$ GeV (concluded from the vertical dashed line) or for RHN mass $\gtrsim \mathcal{O}(84)$ GeV, within 3σ detection prospects of λ_3 . Although the cutoff scale Λ in our framework is relatively low, of order $\mathcal{O}(600\text{--}800)$ GeV, this does not conflict with current collider constraints, as the new states (for the UV completion of the dimension-six operator involved) can naturally be heavier than Λ while remaining fully consistent with existing searches.

Furthermore, stochastic gravitational waves (GWs) produced during the strong FOEWPT could also be detectable in future GW detectors [83] primarily sourced by the sound waves in the plasma ($\Omega_{\text{sw}} h^2$), and magnetohydrodynamic turbulence ($\Omega_{\text{turb}} h^2$) in our case. The GW signals with different Λ values are included in Fig. 4 along with sensitivity regions of various proposed GW detectors like LISA [84], BBO [85–88], DECIGO [85,89], μARES [90], CE [91,92], ET [93–96] and THEIA [97]. The involvement of T_n in estimating such signals are elaborated in Ref. [36].

To sum up, we find that the sphaleron decoupling temperature can effectively be lowered compared to its SM value $T_{\text{sp}}^{\text{SM}} = 131.7$ GeV in the context of first order EWPT. This originates due to the existence of unsuppressed sphaleron transition rate until a relatively low bubble nucleation temperature $T_n < T_{\text{sp}}^{\text{SM}}$, characteristic of the new physics scale (involved in the dimension-six operator in the Higgs potential) responsible for realizing the EWPT of first order. This conclusion, however, continues to hold for other frameworks of FOEWPT, too. The strongly first order electroweak phase transition therefore plays a pivotal role in our framework by lowering the effective sphaleron decoupling temperature, thereby allowing

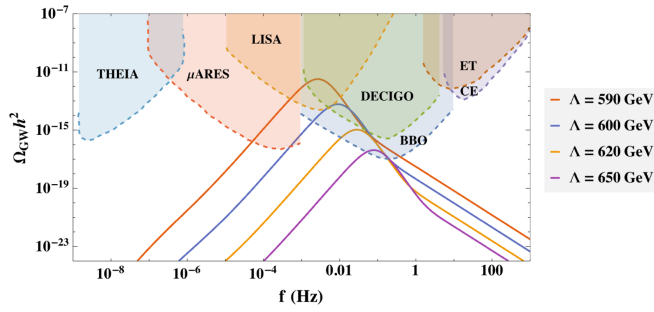


FIG. 4. Gravitational wave signals with sensitivity ranges for various proposed GW detectors.

successful low scale leptogenesis with right-handed neutrino masses below the electroweak scale, even lighter than the SM Higgs. In particular, such a finding paves the way of registering a low scale (resonant) leptogenesis scenario where a set of two quasidegenerate RHNs decay out of equilibrium and produce enough lepton asymmetry below $T < 131.7$ GeV, which can still be converted to baryon asymmetry. This not only enables the seesaw mechanism to be testable at colliders by lowering down the seesaw scale (the RHN mass M_N), but also makes it compatible with the low reheating temperature of the Universe, T_{RH} below 131.7 GeV. The latter realization features a unique finding

since other existing low scale leptogenesis scenarios, such as via oscillation or Higgs decay, require the reheating temperature of the Universe to be at or above 131.7 GeV. This proposal also carries profound importance in exploring the sub-EW mass RHNs at future lepton colliders like FCC-ee, CEPC, and ILC. The one-to-one correspondence between the new physics scale Λ and T_n (and/or the lower limit of RHN mass scale, $M_N \gtrsim 35$ GeV) exhibits a tantalizing possibility to prove leptogenesis and the seesaw mechanism by measuring the triple Higgs coupling at the HL-LHC, which may also illuminate the era of electroweak symmetry breaking in the early Universe with the complimentary information obtained from the detection of associated GWs.

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Data availability. The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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