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Second leptogenesis: Unraveling the baryon-lepton asymmetry discrepancy

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Abstract: We propose a novel scenario to explain the matter-antimatter asymmetry by twofold leptogenesis, wherein heavy Majorana neutrinos exhibit temperature-dependent masses and engage in CP-violating decays. This scenario envisages two distinct phases of leptogenesis: one occurring above the electroweak scale and the other below it. The sphaleron process converts the first lepton asymmetry to baryon asymmetry, but not the second one due to its decoupling. This mechanism potentially explains the significant discrepancy between baryon and lepton asymmetries, as suggested by recent observations of Helium-4. Furthermore, our model implies that the present masses of Majorana neutrinos are lighter than the electroweak scale, offering a tangible avenue for experimental verification in various terrestrial settings.

Keywords: Baryo-and Leptogenesis, Sterile or Heavy Neutrinos

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Contents

1 Introduction

The universe is baryon asymmetric as suggested by the measured baryon-to-photon ratio [1]

$$\eta_B \equiv \frac{n_B - n_{\bar{B}}}{n_\gamma} \simeq (6.14 \pm 0.25) \times 10^{-10}, \tag{1.1}$$

where n_B, $n_{\bar{B}}$, and n_γ are the number density of the baryon, antibaryon, and photon, respectively. Considering cosmic inflation, η_B must be generated after the inflation rather than being an initial condition, [2], a process known as the baryogenesis [3]. It is well-established that the standard model (SM) falls short in explaining this [4–9].

Leptogenesis is a plausible new physics scenario to explain η_B, where the heavy Majorana neutrino N possessing the Yukawa couplings to the lepton doublet ℓ and the Higgs doublet Φ, are added to the SM [10]. At the decoupling of N in the early universe, CP-violating decays $N \to \ell\Phi$ and $N \to \bar{\ell}\Phi^\dagger$ produce the lepton number (L). It is converted to the baryon number (B) via the sphaleron process [11, 12]. N also facilitates the explanation of tiny neutrino masses through the type-I seesaw mechanism [13–17].

Recently, the EMPRESS experiment reported a new result of the ^{4}He abundance observation [18]. It suggests the large degeneracy parameter of the electron neutrino; $\xi_e = 0.05^{+0.03}_{-0.02}$. This implies substantial lepton-to-photon ratio given by [19]

$$\eta_L \equiv \frac{n_L - n_{\bar{L}}}{n_\gamma} \simeq \sum_l \frac{g_l \pi^2}{12\zeta(3)} \left(\frac{T_l}{T_\gamma}\right)^3 \xi_l, \tag{1.2}$$

where l represents all leptons. g_l, ξ_l, and T_l are the degree of freedom, the degeneracy parameter, and the temperature of l, respectively. The EMPRESS result suggests

$$\eta_L \simeq \frac{\pi^2 \sum_{i=e,\mu,\tau} \xi_{\nu_i}}{6\zeta(3)} \left(\frac{T_\nu}{T_\gamma}\right)^3 \simeq (7.5^{+4.5}_{-3.0}) \times 10^{-2}, \tag{1.3}$$

where we have used $(T_\nu/T_\gamma)^3 = 4/11$ and the flavor universality ($\xi_{\nu_e} = \xi_{\nu_\mu} = \xi_{\nu_\tau}$) due to the neutrino oscillations. Given the universe's electrical neutrality, we disregard the charged

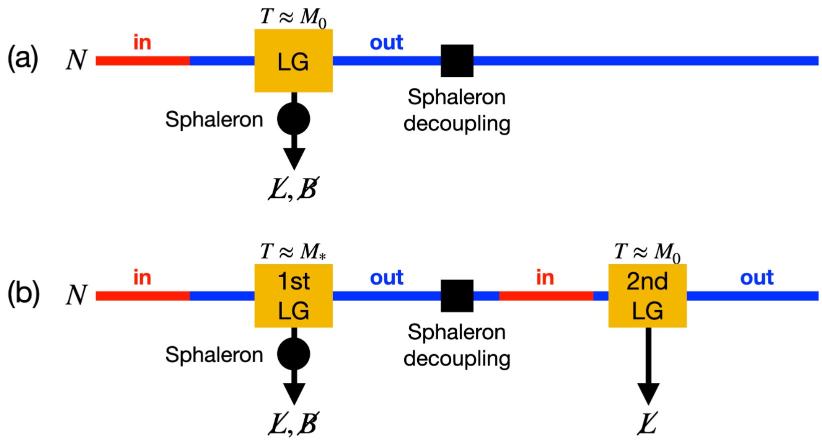

Figure 1. (a) The standard leptogenesis with the constant mass M_0. In the red (blue) region, N is in (out of) thermal equilibrium. Leptogenesis occurs in the blue region. (b) Leptogenesis may occur twice for the temperature-dependent mass, resulting in a larger lepton asymmetry than the baryon asymmetry.

lepton asymmetry of the same magnitude as η_B, as it does not match the observed lepton asymmetry. See refs. [20–28] for further discussions on the lepton asymmetry. While the deviation is at a 2.5σ level and may not appear excessively large, the suggested value has a significant phenomenological impact.

The lepton asymmetry in eq. (1.3) is much larger than the baryon asymmetry in eq. (1.1); $\eta_L/\eta_B \simeq 10^8$. If the sphaleron process occurs frequently, the baryon and lepton asymmetries are made to have the same size, which is proportional to the initial $B - L$ [29, 30]. Thus, the traditional baryogenesis scenario cannot explain this large discrepancy. The discrepancy seems to suggest that two asymmetries were individually generated by different new physics at different epochs in the early universe.

In this paper, we propose a novel, yet simple leptogenesis scenario to explain this discrepancy. It posits two occurrences of leptogenesis in the early universe, driven by the temperature-dependent mass of N. The first leptogenesis takes place prior to the sphaleron decoupling at the electroweak symmetry breaking ($T_{\rm sph} \simeq 100$ GeV) [31], and the generated lepton asymmetry is converted to the baryon asymmetry. On the other hand, the second leptogenesis happens below $T_{\rm sph}$, allowing the resultant lepton asymmetry to persist into the present universe.

2 Second leptogenesis

First, we describe the scenario of the twofold leptogenesis and how it can explain the suggested large discrepancy between the baryon and lepton asymmetries. This concept is visualized in figure 1, alongside a comparison with the standard leptogenesis scenario.

We consider the following case. In the early universe, the heavy neutrino N acquires the temperature-dependent Majorana mass by the new physics effect in addition to the bare Majorana mass M_0. The sums of them are denoted by $M(T)$. We assume that M_0 is smaller than the electroweak scale; $M_0 < 100$ GeV.

$M(T)$ behaves as a constant M_* ($\gg M_0$) above the temperature T_*. At $T < T_*$, it decreases as a function of temperature. When the temperature has dropped enough below a certain temperature T_N ($\ll T_*$), the new physics effect becomes negligibly small, and $M(T) \simeq M_0$ until the current universe. A specific new physics to realize this scenario will be discussed later in this paper.

At high temperatures $T > T_*$, N behaves as Majorana fermion with the mass M_* and is thermalized via the Yukawa interaction. At $T \simeq M_*$, the production rate is exponentially suppressed, and N begins to be decoupled. The first leptogenesis occurs at this stage. L is generated, and it is converted to B, which remains as the baryon asymmetry until the current universe.

At $T < T_*$, $M(T)$ begins to decrease as cooling of the universe; $M(T) \propto T^n$. If this decrease is faster than the temperature decrease ($n > 1$), $M(T)$ can be lower than T at some point, and N can be thermalized again.

As the universe cools further, $M(T)$ behaves as a constant again but much smaller than the mass at the high temperature; $M(T) \simeq M_0 \ll M_*$. At $T \simeq M_0$, the heavy neutrino is decoupled again. The second leptogenesis occurs at this stage, generating the extra lepton number ΔL. Since M_0 is lower than the electroweak scale, the sphaleron process has already decoupled, and ΔL remains as the additional lepton asymmetry until the current universe.

Since the baryon asymmetry is set by L, the size of ΔL has to be much larger than L to explain the large baryon-lepton asymmetry discrepancy. How can a large enhancement of ΔL be made at the second leptogenesis? It can be realized in a natural way. The size of the Yukawa coupling y is proportional to $\sqrt{M_0}$ to reproduce the neutrino mass matrix by the type-I seesaw mechanism [32]. Thus, the ratio of the production rate $\Gamma_{\rm prod} \propto y^2 M(T)$ and the Hubble parameter $H \propto T^2$ is given by $\Gamma_{\rm prod}/H \propto M(T)M_0/T^2$. At the first leptogenesis ($T \simeq M_*$), the production is much suppressed because $\Gamma_{\rm prod}/H \propto M_0/M_* \ll 1$. Such a case is referred to as the weak washout [33], and the generated lepton asymmetry is also suppressed. On the other hand, at the second leptogenesis ($T \simeq M_0$), the production rate is not suppressed because $\Gamma_{\rm prod}/H \propto M(T)M_0/T^2|_{T=M_0} \simeq 1$. This is the strong washout [33], and much larger lepton asymmetry can be generated via the second leptogenesis compared to the first one. It can naturally explain the large difference between the baryon and lepton asymmetries.

3 Realization in a wave dark matter model

Here, we discuss a specific realization of the temperature-dependent mass of N. The neutrino mass variation over cosmic time has been studied a lot in the context of the mass-varying neutrinos in the quintessence dark energy field, where the neutrinos may get their masses from the quintessence field [34, 35]. Lately, there have been many new studies of the mass-varying neutrinos [36–60], including those taking the scalar wave dark matter (DM) [61] in place of the quintessence scalar field. In this paper, we focus on the case that $M(T)$ is given by the coupling to the scalar wave DM.

The scalar wave DM ϕ obeys the following equation of motion with the assumption of spatial homogeneity;

$$\ddot{\phi} + 3H\dot{\phi} + m_\phi^2 \phi = 0, \tag{3.1}$$

where m_ϕ is the mass of ϕ and is constrained to $3 \times 10^{-21}\,\mathrm{eV} < m_\phi < 30\,\mathrm{eV}$ [61]. The lower bound arises from Lyman-α forest data, and the upper from the de Broglie wavelength exceeding inter-particle separation.

At high temperatures, ϕ is fixed at the nonzero initial value due to the large Hubble friction. Since the H decreases over time, it becomes comparable to m_ϕ at the temperature T_*; $H(T_*) = m_\phi$. We assume the universe is radiation-dominated in the following. It leads to $T_* \simeq \left(m_\phi M_{\mathrm{Pl}}\sqrt{90/(8\pi^3 g_*)}\right)^{1/2}$, where M_{Pl} is the Planck mass and $g_* = 106.75$ is the effective degree of freedom of the energy density [38, 40, 51, 59].

At $T < T_*$, ϕ coherently oscillates by the mass term;

$$\phi(t) = \frac{\sqrt{2\rho(t)}}{m_\phi}\cos(m_\phi t), \tag{3.2}$$

where $\rho(t) = \frac{1}{2}\dot{\phi}^2 + \frac{1}{2}m_\phi^2\phi^2$ is the energy density of ϕ. Since ρ behaves as the matter-like, $\rho \propto a^{-3}$ where a is the scale factor, the oscillation amplitude becomes smaller proportional to $a^{-3/2} \propto T^{3/2}$ as the temperature decreases [62–64].

At the current temperature $T_0 \simeq 2.73$ K [65], the oscillation energy density ρ_0 contributes to ρ_{DM}, the relic energy density of the DM. In this paper, we assume that the oscillating ϕ contributes to the entire DM. It requires the current oscillation amplitude ϕ_0 to be $\phi_0 = \sqrt{2\rho_{\mathrm{DM}}}/m_\phi$.

The interaction between particles and the wave DM provides the time-dependent mass of the particles. Some references have investigated leptogenesis with the varying neutrino mass by using quintessence dark energy [66, 67], the neutrino itself as dark energy [68], or other new physics [69]. Nevertheless, none of these works discussed the possibility of the second leptogenesis.

We assume three Majorana neutrinos N_i $(i = 1, 2, 3)$ which couple to the scalar wave DM [39, 51, 59]. The relevant part of the Lagrangian is given by

$$\mathcal{L} = -\frac{1}{2}(M_{0i} + g_i\phi)\overline{N}_i^c N_i + \mathrm{h.c.} \tag{3.3}$$

where M_{0i} and g_i are the bare Majorana masses and coupling constants of N_i, respectively. We do not consider off-diagonal couplings for simplicity. The cosmic scaling of the Majorana neutrino mass in the wave DM setup was also studied in refs. [39, 51, 59].

The second term in eq. (3.3) generates the time-dependent mass $M_i(t) = M_{0i} + g_i\phi(t)$. If the oscillation period is much shorter than the time scale of the relevant physics (in our case, N_i decay time), the oscillating term $\phi(t)$ can be approximated by its time average. As a result, we obtain the temperature-dependent mass of N_i as follows;

$$M_i(T) = \begin{cases} M_{0i} + \frac{g_i\phi_0}{\sqrt{2}}\left(\frac{T_*}{T_0}\right)^{3/2} & T > T_*, \\ M_{0i} + \frac{g_i\phi_0}{\sqrt{2}}\left(\frac{T}{T_0}\right)^{3/2} & T_* > T. \end{cases} \tag{3.4}$$

In order to have the twofold leptogenesis scenario described earlier, we consider the case that the first term of $M_i(T)$ is dominant at low temperatures ($T \ll T_*$); on the other hand, the second one is at high temperatures ($T \simeq T_*$). Then, we can find the temperature T_{N_i} ($\ll T_*$), below which the effect of the wave DM (the second term) is negligible compared with the bare mass (the first term). It is evaluated by $M_{0i} = g_i\phi_0(T_{N_i}/T_0)^{3/2}/\sqrt{2}$.

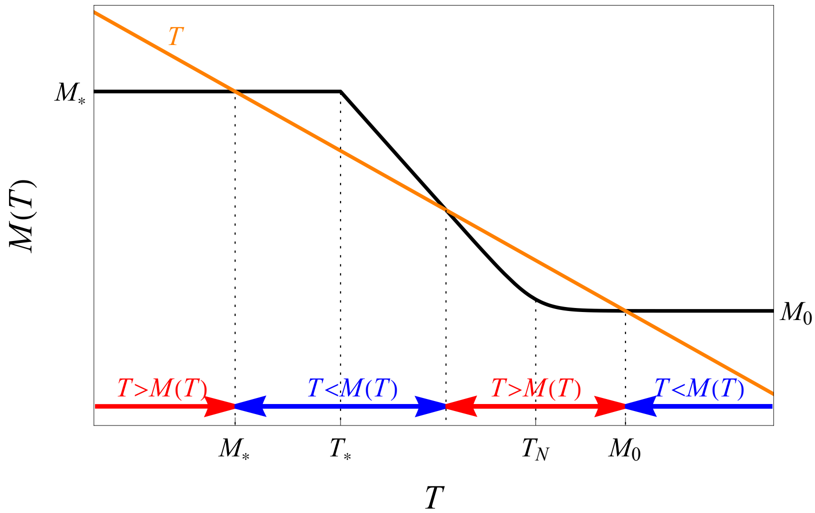

Figure 2. The temperature-dependent mass $M(T)$ in eq. (3.5) using the time-average approximation (black) in comparison with temperature T (orange). Leptogenesis occurs when the heavy neutrinos decouple (two blue periods). Indices i are suppressed here.

m_ϕ	Wave dark matter mass
M_{*i}	Majorana mass at T_*
M_{0i}	Bare Majorana mass
T_*	Temperature when oscillation starts
T_{N_i}	Temperature when $g_i\phi$ term becomes negligible
T_0	Temperature of the current universe

Table 1. The notations of masses and temperatures. The free parameters are only m_ϕ, M_{*i} and M_{0i}, and others are determined by these parameters.

Consequently, the behavior of $M_i(T)$ is described by

$$M_i(T) \simeq \begin{cases} M_{*i} & T > T_*, \\ M_{0i} + g_i \frac{\phi_0}{\sqrt{2}} \left(\frac{T}{T_0}\right)^{3/2} & T_* > T > T_{N_i}, \\ M_{0i} & T_{N_i} > T, \end{cases} \tag{3.5}$$

where $M_{*i} \equiv M_{0i} + g_i\phi_0(T_*/T_0)^{3/2}/\sqrt{2}$. In figure 2, we illustrate how $M_i(T)$ varies with temperature in the second leptogenesis scenario. The masses and temperatures are summarized in table 1.

4 Parameter regions for the second leptogenesis

Here, we consider constraints on the model parameters. The Lagrangian of the model includes three kinds of new parameters m_ϕ, M_{0i}, and g_i. We can also choose a more convenient set of the parameters: T_*, M_{0i}, and M_{*i}. Using these parameters, T_{N_i} is evaluated by

$$T_{N_i} = T_* \left(\frac{M_{0i}}{M_{*i} - M_{0i}}\right)^{2/3} \simeq T_* \left(\frac{M_{0i}}{M_{*i}}\right)^{2/3}. \tag{4.1}$$

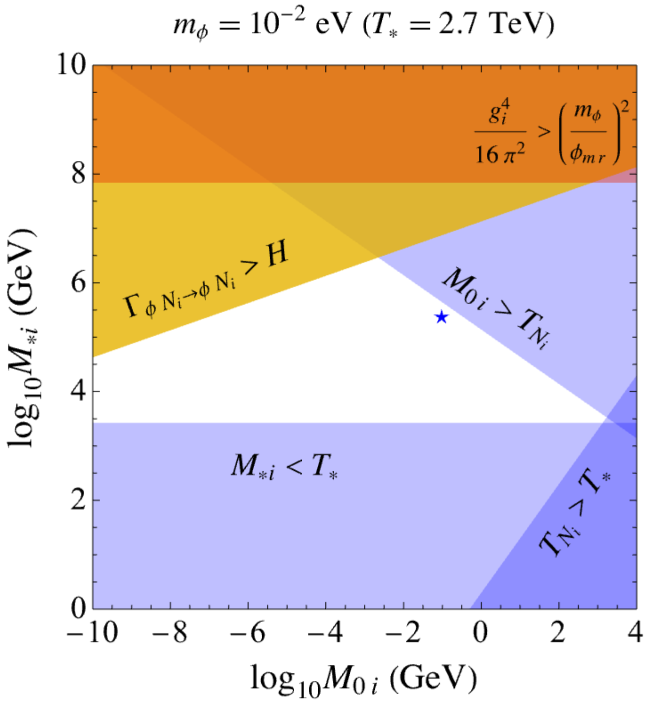

Figure 3. The allowed region for the second leptogenesis (blank region) for $m_\phi = 10^{-2}$ eV. In the blue regions, the necessary conditions for the second leptogenesis are not satisfied. The red and yellow regions are excluded by theoretical requirements to avoid a dominant quartic coupling and thermalization of ϕ, respectively. The blue star is the benchmark point for figure 4.

The necessary conditions for the second leptogenesis are as follows. For the lightest N_i, (i) $M_{*i} > T_*$: the first decoupling happens earlier than T_*. (ii) $T_{N_i} > M_{0i}$: the second decoupling occurs later than T_{N_i}. (iii) $T_* > T_{N_i}$: the time interval for the temperature-dependent mass should exist, which is equivalent to $M_{*i} > 2M_{0i}$ by eq. (4.1).

Next, we consider theoretical constraints. The quartic coupling of ϕ is radiatively induced by g_i, and it has to be smaller than the mass term at least at the matter-radiation equality $T_{mr} \simeq 1$ keV, otherwise $\rho(t)$ behaves as a^{-4} not a^{-3} [51, 57, 70]. This requires $m_\phi^2/\phi_{mr}^2 > g_i^4/(16\pi^2)$, where ϕ_{mr} is the oscillation amplitude at T_{mr}. In addition, in order to avoid thermalization of ϕ, the scattering rate has to be smaller than $H(T)$ [51]. We consider two scatterings $\phi\nu \to \phi\nu$ and $\phi N_i \to \phi N_i$. The former gives a weaker constraint because ν couples to ϕ only via tiny mixing. The scattering rate of the latter is roughly given by $\Gamma_{\phi N_i \to \phi N_i} \sim g_i^4 T$ when N_i is relativistic. Thus, we obtain $g_i^4 > \sqrt{8\pi^3 g_*/90}\; T/M_{\rm Pl}$.

The coupling g_i is also subject to constraints from various experimental studies, such as the Majoron emitting decay [71, 72], neutrino free-streaming on the CMB [49], and neutrino oscillations [37] depending on the m_ϕ values. However, their constraints are weaker than others in the parameter space we are interested in.

In figure 3, we show the allowed parameter regions for the second leptogenesis in the case of $m_\phi = 10^{-2}$ eV, which corresponds to $T_* \simeq 2.7$ TeV. The blue regions do not satisfy the three conditions required for the second leptogenesis. The red and yellow regions are excluded by theoretical constraints to avoid a dominant quartic coupling at $T = T_{mr}$ and to prevent the thermalization of ϕ, respectively. The experimental constraints are too weak to be shown in the figure. Consequently, the second leptogenesis is expected to occur in the blank regions of the figure. In numerical evaluations, we use the average density of the DM $\rho_{\mathrm{DM}} = 1.2 \times 10^{-6}$ GeV/cm^3, not the local density because we investigate phenomena in the early universe.

The allowed region changes with different values of m_ϕ. For example, the constraint $\Gamma_{\phi N_i \to \phi N_i} > H$ becomes stronger for larger m_ϕ, while the quartic coupling constraint becomes more stringent for smaller m_ϕ. Since we consider the scenario where the oscillation of ϕ begins before the sphaleron decoupling as explained in section 2, m_ϕ needs to be larger than 10^{-5} eV, which is derived from $T_* > T_{\mathrm{sph}} \simeq 100$ GeV. Thus, our relevant mass region is 10^{-5} eV $< m_\phi < 30$ eV. We have checked that the allowed region does not change significantly, and we can find a lot of parameter points to achieve the second leptogenesis in this mass region.

5 Quantitative result of the second leptogenesis

The asymmetry production in leptogenesis is evaluated by the density matrix equation including the flavor effect [73–75]. Here, we employ the formalism given in refs. [76–78] with the addition of the temperature dependence in the masses. The equation is given by

$$\frac{\mathrm{d}N_{N_i}}{\mathrm{d}z} = -D_i(N_{N_i} - N_{N_i}^{\mathrm{eq}}), \tag{5.1}$$

$$\begin{aligned}\frac{\mathrm{d}N_{\alpha\beta}}{\mathrm{d}z} &= \sum_i \left[\varepsilon_{\alpha\beta}^{(i)} D_i(N_{N_i} - N_{N_i}^{\mathrm{eq}}) - \frac{1}{2}W_i\{P_i, N\}_{\alpha\beta}\right] \\ &\quad - \frac{\Gamma_\tau}{Hz}[I_\tau, [I_\tau, N]]_{\alpha\beta} - \frac{\Gamma_\mu}{Hz}[I_\mu, [I_\mu, N]]_{\alpha\beta},\end{aligned} \tag{5.2}$$

where $z = M_{01}/T$, $i = 1, 2, 3$, and α, $\beta = e$, μ, τ. N_{N_i} and the diagonal terms $N_{\alpha\alpha}$ are the number of N_i and $B/3 - L_\alpha$, where L_α is the lepton number for each flavor, respectively, in a portion of the comoving volume that contains one photon at the era when N_i is relativistic and in thermal equilibrium [79, 80]. The off-diagonal terms $N_{\alpha\beta}$ $(\alpha \neq \beta)$ represent the coherence between the flavors. The number of the total $B - L$ is given by $N_{B-L} = \sum_\alpha N_{\alpha\alpha}$. The term D_i accounts for the decay and inverse decay of N_i. The washout effect in $B - L$ asymmetry is described by W_i. We consider the washout effect from the inverse decay and neglect one from other lepton-number-violating processes. D_i and W_i are given by [80]

$$D_i = \frac{(yy^\dagger)_{ii}}{8\pi Hz} M_i(T) \frac{K_1(M_i(T)/T)}{K_2(M_i(T)/T)}, \tag{5.3}$$

$$W_i = \frac{2}{3} D_i N_{N_i}^{\mathrm{eq}}, \tag{5.4}$$

where $N_{N_i}^{\mathrm{eq}} = (3/8)(M_i(T)/T)^2 K_2(M_i(T)/T)$ is the equilibrium value of N_{N_i}, and y is the Yukawa matrix for the interaction among the lepton doublet, the Higgs doublet, and the heavy neutrinos. K_n is a modified n-th Bessel function of the second kind. We note that the temperature dependence of the mass is included in these terms. P_i is the projection matrix constructed with the Yukawa matrix. Decoherence effects by the interchanges between the left-handed and right-handed leptons are described by the double commutator terms, where $\Gamma_{\mu(\tau)} \simeq 8 \times 10^{-3}(\sqrt{2}m_{\mu(\tau)}/v)^2 T$ [77] is the rate of the process involving $\mu(\tau)$, $I_\mu = \mathrm{diag}(0,1,0)$, and $I_\tau = \mathrm{diag}(0,0,1)$.

Since M_{0i} in our scenario is much lighter than the Davidson-Ibarra bound, $M_{0i} \gtrsim 10^8$ GeV [81], for the heavy neutrinos with hierarchical masses, we consider resonant leptogenesis, where the heavy neutrinos possess very close masses [82, 83]. The CP asymmetry parameter $\varepsilon^{(i)}_{\alpha\beta}$ is divided into contributions from vertex and self-energy diagrams [33]. The self-energy part $\varepsilon^{S(i)}_{\alpha\beta}$ is resonantly enhanced with degenerate masses and can dominate the CP asymmetry parameter. $\varepsilon^{S(i)}_{\alpha\beta}$ is given by

$$\begin{aligned}\varepsilon^{S(i)}_{\alpha\beta} = \frac{1}{16\pi(yy^\dagger)_{ii}} \sum_{j\neq i} \Big\{ & i\left[y^*_{i\alpha}y_{j\beta}(yy^\dagger)_{ji} - y_{i\beta}y^*_{j\alpha}(yy^\dagger)_{ij}\right]\frac{M_j}{M_i} \\ & + i\left[y^*_{i\alpha}y_{j\beta}(yy^\dagger)_{ij} - y_{i\beta}y^*_{j\alpha}(yy^\dagger)_{ji}\right]\Big\} \frac{(M_j^2 - M_i^2)M_i^2}{(M_j^2 - M_i^2)^2 + M_i^4\Gamma_j^2/M_j^2}, \end{aligned} \tag{5.5}$$

where M_j and Γ_j are the temperature-dependent mass and decay rate of N_j, respectively. Thus, $\varepsilon^{(i)}_{\alpha\beta}$ is significantly enhanced when the resonant condition, $|M_j - M_i| \simeq \Gamma_j/2$, is satisfied.

As a benchmark point to evaluate N_{N_1} and N_{B-L}, we assume the following input parameters: $m_\phi = 10^{-2}\,\mathrm{eV}$ ($T_* \simeq 2.7\,\mathrm{TeV}$), $M_{01} = 0.1\,\mathrm{GeV}$, and $M_{*1} = 2.4 \times 10^5\,\mathrm{GeV}$. The current masses of N_2 and N_3 are chosen to satisfy the resonant condition at the second leptogenesis, $\Delta M_{12} \equiv M_{02} - M_{01} = 0.5 \times 10^{-19}$ GeV and $\Delta M_{13} \equiv M_{03} - M_{01} = 4.0 \times 10^{-19}$ GeV. M_{*i} ($i = 2, 3$) are determined by imposing $M_{*i}/M_{*1} = M_{0i}/M_{01}$ by which the resonant condition is also satisfied at the first leptogenesis.

The Yukawa matrix y is set to satisfy the neutrino oscillation data with the normal ordering masses [65] by using the Casas-Ibarra parametrization [32], $y = \sqrt{2}\hat{M}_N^{1/2} R \hat{m}_\nu^{1/2} U^\dagger / v$, where $\hat{M}_N$ and $\hat{m}_\nu$ are the diagonal mass matrices of the heavy and active neutrinos, respectively, U is the PMNS matrix [84, 85], v is the vacuum expectation value of the Higgs, and R is a complex orthogonal matrix. We note that $\hat{M}_N$ is evaluated at the current temperature. Six parameters in y are undetermined by the neutrino oscillation data: the lightest neutrino mass m_{ν_1}, two Majorana phases α_1, α_2 in the PMNS matrix U, and three complex phases ω_1, ω_2, ω_3 in R. The notations of the Majorana phases, and the complex phases follow refs. [65] and [76], respectively. We assume the following values for them; $m_{\nu_1} = 0$ eV, $\alpha_1 = \alpha_2 = 0$, $\omega_1 = \omega_2 = 0$, and $\omega_3 = 0.2e^{-i\pi/4}$.

We consider thermal leptogenesis with the initial condition $N_{N_1} = N_{B-L} = 0$ at $z = 10^{-10}$ ($T = 10^9$ GeV). Figure 4 shows the behavior of N_{N_1} and N_{B-L} with the above inputs. The first leptogenesis occurs a little later than $T = M_{*1}$ ($z \simeq 10^6$) because of weak washout [80]. The produced lepton number is converted to the baryon number by the sphaleron process, and it is fixed after the sphaleron decoupling. For simplicity, we assume that the decoupling occurs instantaneously, and the final baryon number is determined by N_{B-L} at the temperature $T = 100$ GeV ($z = 10^3$) shown by the purple lines in figure 4.

The baryon-to-photon ratio can be obtained by $\eta_B = a_{\rm sph} N_{B-L}/f$ at the sphaleron decoupling, where $a_{\rm sph}$ and f are the sphaleron conversion rate [29, 30] and the photon dilution factor [79] due to the increase of the photon number by annihilation of particles from the first leptogenesis till the BBN, respectively. We use $a_{\rm sph} = 28/79$, the value for the SM plasma. Since the heavy neutrinos have nearly degenerate masses, we evaluate f

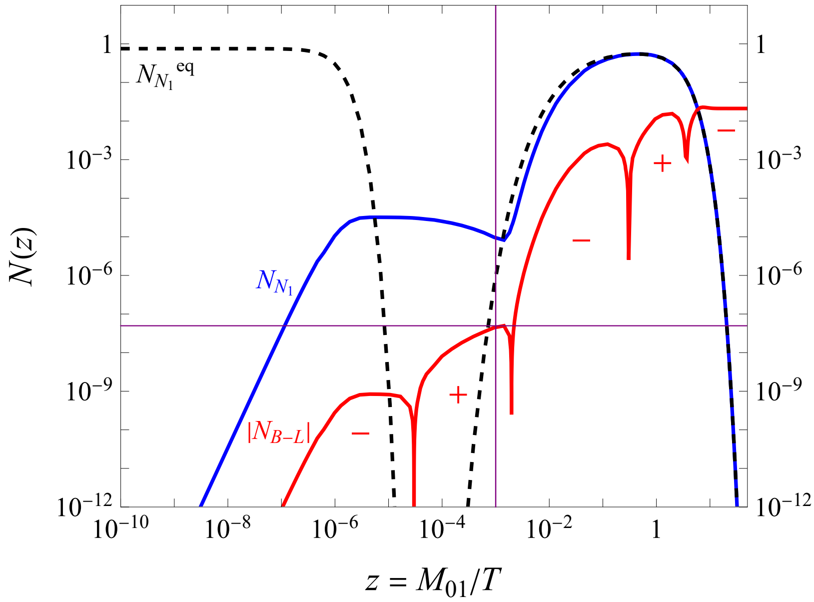

Figure 4. N_{N_1} and $|N_{B-L}|$ for resonant leptogenesis with the temperature-dependent masses of the neutrinos on the benchmark point described in the text. Here, $|N_{B-L}|$ is the sum of the $B-L$ asymmetry deposited in each flavor, and its sign is represented together. The baryon asymmetry is determined after the sphaleron decoupling ($T \simeq 100$ GeV corresponding to $z \simeq 10^{-3}$ with $M_{01} = 0.1$ GeV) with positive N_{B-L}, but the lepton asymmetry is fixed only after the second leptogenesis ($z \simeq 1$) with negative N_{B-L}. The baryogenesis fitting at the sphaleron decoupling ($z \simeq 10^{-3}$) is shown together (thin lines).

including the effect of N_1, N_2, and N_3, not only N_1. Then, we have $f = 1232/43$.[1] As a result, we obtain $\eta_B \simeq 6.14 \times 10^{-10}$, which is consistent within 1σ level with the current observations [1, 86]. Also, the positive N_{B-L} at the sphaleron decoupling in our benchmark point provides the correct sign for the baryon asymmetry.

The second leptogenesis commences shortly after $T = M_{01}$ (or $z = 1$), due to the effects of strong washout [80]. The lepton number generated during this period persists into the current universe, as all lepton-number-violating processes have ceased. In this scenario, the photon dilution factor f' is calculated based on the change in photon numbers from the second leptogenesis to the era of BBN, rather than from the first leptogenesis. This is because the N_1 species is relativistic and returns to thermal equilibrium before the onset of the second leptogenesis. We adopt $f' = 176/43$, where the thermal bath consists of $e^\pm$, ν_ℓ, γ, and N_i before the second leptogenesis.[2] Consequently, we derive a lepton asymmetry value of $\eta_L = N_{B-L}/f' \simeq 5.0 \times 10^{-3}$. The lepton asymmetry is flavor-universal due to neutrino oscillation, so we consider the summation of the $B-L$ asymmetry across all flavors, rather than focusing solely on the electron component. Since N_{B-L} is negative in the late epoch of our benchmark model, a positive lepton asymmetry is also guaranteed.

This demonstrates that the second leptogenesis can significantly amplify lepton asymmetry, increasing it by several orders of magnitude from the baryon asymmetry. In our analysis,

[1] With m generations of the heavy neutrinos, $f = 11(427 + 7m)/172$, which leads to the commonly used value $f = 2387/86$ in the case of $m = 1$ [79].

[2] With m generations of the heavy neutrinos, $f' = 11(43 + 7m)/172$.

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the benchmark point, which was not fully optimized, already indicates $\eta_L \sim 10^{-3}$. This is remarkably close to the EMPRESS data, which suggests $\eta_L \sim 10^{-2}$, and represents a substantial deviation from $\eta_B \sim 10^{-10}$. A more refined analysis or the addition of more Majorana neutrinos might aid in reconciling the slight discrepancy from the observed values.

6 Summary and outlook

It is notable that a significant deviation of the lepton asymmetry from the baryon asymmetry can be explained in a rather simple framework of the second leptogenesis. This scenario allows only larger lepton asymmetry than the baryon asymmetry, not the other way around, in accordance with the measurement. A more comprehensive study will follow in the subsequent work. In the future, there will be increased CMB data from the Simons Observatory [87] and CMB-S4 [88] that can either confirm or refute the discrepancy [23].

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