

Leptogenesis via multiscalar coherent evolution in a supersymmetric neutrino seesaw model

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A novel scenario of leptogenesis is investigated in the supersymmetric neutrino seesaw model. The right-handed sneutrino $\tilde{N}$ and the ϕ field in the $\tilde{L}H_u$ direction of the slepton and Higgs doublets together start coherent evolution after inflation with a right-handed neutrino mass M_N smaller than the Hubble parameter of inflation. Then, after some period the motion of $\tilde{N}$ and ϕ is drastically changed by the cross coupling $M_N h_\nu \tilde{N}^* \phi \phi$ from the $M_N NN$ and $h_\nu NLH_u$ terms, and the significant asymmetries of $\tilde{N}$ and $\tilde{L}$ are generated. The $\tilde{L}$ asymmetry is fixed later by the thermal effect as the lepton number asymmetry for baryogenesis, while the $\tilde{N}$ asymmetry disappears through the decays $\tilde{N} \rightarrow \tilde{L}\tilde{H}_u, \tilde{L}H_u$ with almost the same rate but opposite final lepton numbers.

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Baryogenesis via leptogenesis is considered one of the promising scenarios to explain the baryon number asymmetry in the Universe [1]. Leptogenesis is interesting particularly from the point of view that it may be related to the neutrino mass generation. In the supersymmetric standard model, as investigated fully so far, leptogenesis may be realized via the Affleck-Dine mechanism [2,3] in the $\tilde{L}H_u$ flat direction of the slepton and Higgs doublets, $\tilde{L}$ and H_u , requiring a very small mass of the lightest ordinary neutrino [4]. It is also possible to realize the leptogenesis on the flat manifold of $\tilde{L}-H_u-H_d$, where the restriction on the lightest neutrino mass may be considerably moderated [5]. In this article, we investigate another novel scenario of leptogenesis in the supersymmetric seesaw model for neutrino masses [6]. The lepton number asymmetry is indeed generated via the coherent evolution of the multiscalar fields, the right-handed sneutrino $\tilde{N}$ and the ϕ field in the $\tilde{L}H_u$ direction. The potential terms provided with the supersymmetric neutrino seesaw and also the thermal effect [7] play important roles for leptogenesis in the respective epochs. The leptogenesis is completed when the generated lepton number asymmetry is fixed to some significant value by the thermal effect at the scale much higher than the gravitino mass $m_{3/2} \sim 10^3$ GeV. Hence, this scenario is not restricted by the low-energy electroweak physics.

In the present scenario of leptogenesis, it is supposed that some of the masses M_N 's of the right-handed neutrinos N 's (antineutrinos strictly) are smaller than the Hubble parameter $H_{\text{inf}} \sim 10^{13}$ GeV during inflation. Specifically, we describe the generation of lepton number asymmetry by considering the simple case that

$$M_{N_1} < H_{\text{inf}} \quad (1)$$

for one N_1 while $M_{N_2}, M_{N_3} > H_{\text{inf}}$ for the others N_2, N_3 . The lepton doublets are arranged with unitary transformation so

that only L_1 has the Yukawa coupling with N_1 . Then, the right-handed sneutrino $\tilde{N}_1$ and the ϕ field in $\tilde{L}_1 H_u$ together start coherent evolution with large initial field values after the inflation in the manner of Affleck and Dine. The motions of $\tilde{N}_1$ and ϕ are linked through the superpotential term $h_\nu N_1 L_1 H_u$. ($\tilde{N}_2 = \tilde{N}_3 = 0$ due to the large masses $M_{N_2}, M_{N_3} > H_{\text{inf}}$.) Henceforth the generation indices are suppressed by considering only the one generation for leptogenesis, and the relevant scalar fields are specified as

$$\tilde{N}, \tilde{L} = \begin{pmatrix} \phi/\sqrt{2} \\ 0 \end{pmatrix}, \quad H_u = \begin{pmatrix} 0 \\ \phi/\sqrt{2} \end{pmatrix}. \quad (2)$$

If some of the M_N 's are smaller than H_{inf} in general, the coherent evolution after the inflation may be much more multidimensional involving the $\tilde{N}$'s, $\tilde{L}$'s and H_u . The leptogenesis scenario is essentially valid even in such cases, where the main source for asymmetry generation is the cross coupling $M_N h_\nu \tilde{N}^* \phi \phi$ from the $M_N NN$ and $h_\nu NLH_u$ terms of the supersymmetric neutrino seesaw model.

The relevant superpotential is given by

$$W = \frac{M_N}{2} NN + \frac{e^{i\delta_N}}{4M} NNNN + h_\nu NLH_u. \quad (3)$$

The $NNNN$ term may originate in the physics of Planck scale, and its phase factor $e^{i\delta_N}$ is included here with real M_N . The NNN term is discarded for simplicity by requiring the R parity. The $LH_u LH_u$ term is not considered either, since it does not provide a significant effect if the Yukawa coupling $h_\nu > 0$ is not extremely small. As seen later in Eqs. (6) and (7), the Yukawa coupling $h_\nu \sim 3 \times 10^{-3}$ is relevant for the present scenario of leptogenesis starting at the large scale $\sim 10^{15}$ GeV. Then, its value at the electroweak scale M_W is evaluated as $h_\nu(M_W) \sim 10^{-3}$ by considering the renormalization group effects mainly provided by the top quark loop for the H_u field. The ordinary neutrino mass via seesaw mechanism is roughly estimated as

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$$m_\nu \sim 10^{-4} \text{ eV} \left(\frac{h_\nu(M_W)}{10^{-3}} \right)^2 \left(\frac{10^{11} \text{ GeV}}{M_N} \right) \quad (4)$$

depending on $h_\nu(M_W)$ and M_N . (The neutrino mixing is present in general with the matrix form of h_ν .) Hence, this neutrino relevant for leptogenesis should be identified with the lightest one, being compatible with the data on the atmospheric and solar neutrino experiments [8,9]. It is interesting that the lightest neutrino mass is expected to be $m_\nu \sim 10^{-4}$ eV for the present leptogenesis with $\tilde{N}$ and ϕ , while $m_\nu \lesssim 10^{-8}$ eV is required for the conventional Affleck-Dine leptogenesis in the $\tilde{L}H_u$ flat direction [4].

The scalar potential is given with W in Eq. (3) as

$$\begin{aligned} V = & -c_N H^2 |\tilde{N}|^2 - c_\phi H^2 |\phi|^2 + \left| M_N \tilde{N} + \frac{e^{i\delta_N}}{M} \tilde{N} \tilde{N} \tilde{N} + \frac{h_\nu}{2} \phi \phi \right|^2 \\ & + h_\nu^2 |\tilde{N}|^2 |\phi|^2 + H \left(b_N \frac{M_N}{2} \tilde{N} \tilde{N} + a_N \frac{e^{i\delta_N}}{4M} \tilde{N} \tilde{N} \tilde{N} + \text{H.c.} \right) \\ & + H \left(a_h \frac{h_\nu}{2} \tilde{N} \phi \phi + \text{H.c.} \right) + V_{\text{th}}(\phi). \end{aligned} \quad (5)$$

Here the soft supersymmetry breaking terms are induced by the expansion of the universe with the Hubble parameter H . The thermal terms [7] are also included in $V_{\text{th}}(\phi)$. The D^2 terms are vanishing for the ϕ field. The evolution of the scalar fields is governed by the equations of motion with this potential V and the redshift of H .

(i) Inflation epoch: $H = H_{\text{inf}}$

The scalar fields settle into one of the minima $(\tilde{N}_0, \phi_0)$ of V during the inflation with $H = H_{\text{inf}}$, which are determined as

$$|\tilde{N}_0| \sim |\phi_0| \sim 3 \times 10^{15} \text{ GeV} \left(\frac{H_{\text{inf}} M}{10^{13} \text{ GeV} \times 10^{18} \text{ GeV}} \right)^{1/2} \quad (6)$$

for the Yukawa coupling

$$h_\nu \sim 3 \times 10^{-3} \left(\frac{H_{\text{inf}}/M}{10^{13} \text{ GeV}/10^{18} \text{ GeV}} \right)^{1/2}. \quad (7)$$

(We henceforth take $H_{\text{inf}} = 10^{13}$ GeV typically.) We consider for definiteness the case with this range of h_ν , though it does not require a fine-tuning. If $h_\nu < (H_{\text{inf}}/M)^{1/2}$, $|\tilde{N}_0|$, and $|\phi_0|$ take larger values. If $h_\nu > (H_{\text{inf}}/M)^{1/2}$, on the other hand, $\phi_0 = 0$ may be obtained due to the $h_\nu^2 |\tilde{N}|^2 |\phi|^2$ term. The leptogenesis can be realized even in these cases with some modifications of scenario, which will be described elsewhere.

(ii) Oscillation epoch: $H_{\text{inf}} > H > H_{\text{tr}}$

After the inflation the Hubble parameter decreases as $H = (2/3)t^{-1}$ in the matter dominated universe, and the multi-scalar coherent evolution of $\tilde{N}$ and ϕ starts with the initial

condition $(\tilde{N}, \phi) = (\tilde{N}_0, \phi_0)$ at $t = t_0 \sim H_{\text{inf}}^{-1}$, as given in Eq. (6). The higher order potential terms suppressed by M are soon reduced by redshift, and the quartic couplings $h_\nu^2 |\phi|^4$ and $h_\nu^2 |\tilde{N}|^2 |\phi|^2$ dominate in this epoch with h_ν as given in Eq. (7). Then, driven by these quartic couplings, the scalar fields oscillate in magnitude with scaling by redshift as

$$|\tilde{N}| \sim |\phi| \sim (H_{\text{inf}} M)^{1/2} (H/H_{\text{inf}})^{2/3} \propto H^{2/3}. \quad (8)$$

The field phases, however, remain almost constant except for the vicinities of $\tilde{N} = 0$ and $\phi = 0$, and the significant asymmetries of $\tilde{N}$ and $\tilde{L}$ do not appear in this epoch.

(iii) Transition epoch: $H_{\text{tr}} \gtrsim H > H_{\text{th}}$

Since $\tilde{N}$ and ϕ decrease with H as given in Eq. (8), the mass term $M_N^2 |\tilde{N}|^2$ and $M_N h_\nu$ cross coupling $M_N h_\nu \tilde{N}^* \phi \phi$ become comparable to the quartic couplings $h_\nu^2 |\phi|^4$ and $h_\nu^2 |\tilde{N}|^2 |\phi|^2$ with $|\tilde{N}| \sim |\phi| \sim M_N/h_\nu$ and the Hubble parameter

$$H_{\text{tr}} \sim 10^{10} \text{ GeV} \left(\frac{M_N}{10^{11} \text{ GeV}} \right)^{3/2}, \quad (9)$$

where $h_\nu \sim (H_{\text{inf}}/M)^{1/2}$ with $H_{\text{inf}} = 10^{13}$ GeV is taken from Eq. (7). The thermal mass term should also be considered at $H \sim H_{\text{tr}}$, which is given by $(yT_p)^2 |\phi|^2$ with relevant coupling y under the condition $y|\phi| < T_p$ [7]. The temperature T_p of the dilute plasma of inflaton decay products is given in terms of the reheating temperature T_R of the universe after the inflaton decay is completed:

$$T_p \sim (T_R^2 H M_P)^{1/4}, \quad (10)$$

where $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. The thermal mass is constrained at $H \sim H_{\text{tr}}$ as $yT_p < T_p^2/|\phi| \sim h_\nu T_p^2/M_N$ for $y|\phi| < T_p$ with $|\tilde{N}| \sim |\phi| \sim M_N/h_\nu$. Hence, the thermal mass term $(yT_p)^2 |\phi|^2$ is smaller than the $M_N^2 |\tilde{N}|^2$ and $M_N h_\nu \tilde{N}^* \phi \phi$ terms at $H \sim H_{\text{tr}}$ for the right-handed neutrino mass

$$M_N \gtrsim 10^{10} \text{ GeV} \left(\frac{h_\nu}{3 \times 10^{-3}} \right)^{4/5} \left(\frac{T_R}{10^9 \text{ GeV}} \right)^{4/5}. \quad (11)$$

In this situation, the $M_N^2 |\tilde{N}|^2$ and $M_N h_\nu \tilde{N}^* \phi \phi$ terms as well as the $h_\nu^2 |\phi|^4$ term dominate for $H \lesssim H_{\text{tr}}$, so that the motion of $\tilde{N}$ and ϕ is changed drastically. Specifically, the $\tilde{N}$ field oscillates mainly driven by the mass term $M_N^2 |\tilde{N}|^2$ with $|\tilde{N}| \propto H$. The motion of ϕ follows after $\tilde{N}$ toward the new stable configuration with $(h_\nu/2) \phi \phi \approx -M_N \tilde{N}$ so as to make $|F_N|^2 \sim |F_\phi|^2 \ll M_N^2 |\tilde{N}|^2$ in V , where $F_N \approx M_N \tilde{N} + (h_\nu/2) \phi \phi$ and $F_\phi = h_\nu \tilde{N} \phi$. Consequently, the scalar fields decrease roughly as

$$|\tilde{N}| \sim (M_N/h_\nu)(H/H_{\text{tr}}) \propto H, \quad (12)$$

$$|\phi| \sim (M_N/h_\nu)(H/H_{\text{tr}})^{1/2} \propto H^{1/2}, \quad (13)$$

with $|F_N| \sim |F_\phi| \propto H^{3/2}$. Through this drastic change in the multiscalar coherent evolution, the significant asymmetries of $\tilde{N}$ and $\tilde{L}$ appear, which is really seen in the rate equations

$$\begin{aligned} \frac{d}{dt} \left(\frac{n_{\tilde{N}}}{H^2} \right) \simeq & -\frac{2}{H^2} \text{Im}[b_N M_N H \tilde{N} \tilde{N}] \\ & -\frac{2}{H^2} \text{Im} \left[M_N \tilde{N} F_N^* + \frac{a_h h_\nu}{2} H \tilde{N} \phi \phi \right], \end{aligned} \quad (14)$$

$$\frac{d}{dt} \left(\frac{n_{\tilde{L}}}{H^2} \right) \simeq -\frac{2}{H^2} \text{Im} \left[\frac{h_\nu}{2} \phi \phi F_N^* + \frac{a_h h_\nu}{2} H \tilde{N} \phi \phi \right], \quad (15)$$

with $n_{\tilde{L}} = n_{H_u} = n_\phi/2$. The main sources are scaled as $\text{Im}[(h_\nu/2)\phi\phi F_N^*]/H^2 \simeq -\text{Im}[M_N \tilde{N} F_N^*]/H^2 \propto H^{5/2}/H^2$ with $|F_N| \propto H^{3/2}$, and hence the asymmetries $n_{\tilde{N}}$ and $n_{\tilde{L}}$ oscillate rapidly by the exchange $\tilde{N} \leftrightarrow \tilde{L}$. The sum $n_{\tilde{N}} + n_{\tilde{L}}$, however, varies rather moderately with the remaining sources $\propto H^3$, since the main sources are cancelled as $\text{Im}[F_N F_N^*] = 0$ with $F_N \simeq M_N \tilde{N} + (h_\nu/2)\phi\phi$.

(iv) Completion epoch: $H_{\text{th}} \gtrsim H \gg m_{3/2}$

After the transition epoch continues for some period, the thermal log term [7] eventually becomes significant on the evolution of ϕ . It is mainly provided as

$$a_{\text{th}} \alpha_s^2 T_p^4 \ln(|\phi|^2/T_p^2) \quad (16)$$

($a_{\text{th}} = 9/8$) through the modification of $\text{SU}(3)_C$ coupling due to the decoupling of top quark from the plasma with large mass $h_t |\phi|/\sqrt{2} > T_p$. This thermal log term acts as the effective mass term for the ϕ field giving $(a_{\text{th}} \alpha_s^2 T_p^4/|\phi|^2)\phi \propto H^{1/2}$ in $\partial V/\partial \phi^*$. It dominates over the term $F_N h_\nu \phi^* \propto H^2$ in $\partial V/\partial \phi^*$ ($|F_N| \sim |F_\phi| = h_\nu |\tilde{N}| |\phi|$) with the Hubble parameter

$$\begin{aligned} H_{\text{th}} \sim 10^7 \text{ GeV} \left(\frac{h_\nu}{3 \times 10^{-3}} \right)^{4/3} \left(\frac{M_N}{10^{11} \text{ GeV}} \right)^{-1/6} \\ \times \left(\frac{T_R}{10^9 \text{ GeV}} \right)^{4/3}, \end{aligned} \quad (17)$$

where Eqs. (10), (13), and (12) are considered. Then, the rotation of the ϕ field phase is accelerated by this thermal log term with the change of field scaling

$$|\phi| \propto H^{1/2} \rightarrow H^{3/2} \quad (18)$$

while keeping $|\tilde{N}| \propto H$. After a while the top quark enters the plasma at $H \sim 0.1 H_{\text{th}}$ with $|\phi| \sim T_p$ ($h_t \sim 1$). Then, the thermal mass term $T_p^2 |\phi|^2$ instead becomes important, and the ϕ field decreases as $|\phi| \propto H^{7/8}$. In the preceding epoch, the significant exchange of asymmetries, $n_{\tilde{N}} \leftrightarrow n_{\tilde{L}}$, took place through the $\tilde{N}$ - ϕ couplings, as seen in Eqs. (14) and (15).

These couplings are actually turned off in this epoch with the rapid decrease of $|\phi| \propto H^{3/2}$ and $H^{7/8}$ later due to the thermal terms, and the ϕ and $\tilde{N}$ evolve almost independently.

In this way, the $\tilde{L}$ asymmetry is fixed to some significant value as the lepton number asymmetry for $t > H_{\text{th}}^{-1}$,

$$n_{\tilde{L}} \simeq n_L \equiv \epsilon_L [(3/2)H^2 M], \quad (19)$$

since the $\tilde{L}$ violating sources in Eq. (15) decrease fast enough as H^4/H^2 and $H^{11/4}/H^2$ later with rapidly varying phase of $\tilde{N}^* \phi \phi$. This concludes that the thermal effect plays the positive role for the completion of leptogenesis, which is in salient contrast to the conventional Affleck-Dine mechanism where the thermal effect rather suppresses the asymmetry seriously. The resultant lepton-to-entropy ratio after the reheating is estimated with $s \simeq 3H_R^2 M_p^2/T_R$ as

$$\frac{n_L}{s} \sim 10^{-10} \left(\frac{\epsilon_L}{1} \right) \left(\frac{M}{10^{18} \text{ GeV}} \right) \left(\frac{T_R}{10^9 \text{ GeV}} \right). \quad (20)$$

Here the reheating temperature is restricted as $T_R \lesssim 10^8 - 10^{10} \text{ GeV}$ to avoid the gravitino problem [10–12]. The lepton number asymmetry is converted to the baryon number asymmetry through the electroweak anomalous effect as $n_B/s = -(8/23)n_L/s$ [13]. Hence, the sufficient baryon-to-entropy ratio can be provided for the nucleosynthesis with $\eta = (2.6-6.2) \times 10^{-10}$ [14].

The motion of $\tilde{N}$ after the decoupling from ϕ for $H < H_{\text{th}} \ll M_N$ is determined by the $M_N^2 |\tilde{N}|^2$ and $b_N H M_N \tilde{N} \tilde{N}$ terms, and the analytic solution is obtained in a good approximation with $H \ll M_N$ for the two eigenmodes $\eta_R(t)$ and $\eta_I(t)$ in $\tilde{N}(t)$ as

$$\eta_{R,I}(t) \simeq \overline{\eta_{R,I}} \cos[M_N t + \sigma_{R,I}(|b_N|/3) \ln t + \delta_{R,I}] \quad (21)$$

with $\sigma_R = +1$, $\sigma_I = -1$, and

$$\tilde{N} \equiv H(M/M_N)^{1/2} (b_N/|b_N|)^{-1/2} (\eta_R + i \eta_I). \quad (22)$$

The parameters $\overline{\eta_{R,I}}$ and $\delta_{R,I}$ are determined as the result of $\tilde{N}$ motion from $t = t_0 \sim H_{\text{inf}}^{-1}$ through $t > H_{\text{th}}^{-1} \gg M_N^{-1}$. The $\tilde{N}$ asymmetry is evaluated with Eqs. (21) and (22) as

$$n_{\tilde{N}}(t) \simeq -2H^2 M \overline{\eta_R} \overline{\eta_I} \cos[(2|b_N|/3) \ln t + \delta_R - \delta_I], \quad (23)$$

where the rapid oscillations of η_R and η_I with $M_N t$ in Eq. (21) are canceled.

This $\tilde{N}$ asymmetry oscillates slowly in $\ln t$ for some while due to the b_N term, as seen in Eq. (23). Then, the incoherent decays of $\tilde{N}$ become significant with the dominant modes

$$\tilde{N} \rightarrow \bar{L} \bar{H}_u [L = -1], \quad \tilde{L} H_u [L = +1], \quad (24)$$

where the decay products are ultra-relativistic with $h_\nu |\tilde{N}| \ll M_N/2$. The motion of $\tilde{N}$ is significantly decelerated by these $\tilde{N}$ decays at $H \sim \Gamma_{\tilde{N}} \simeq (h_\nu^2/4\pi) M_N$ ($\sim 10^5 \text{ GeV}$ nu-

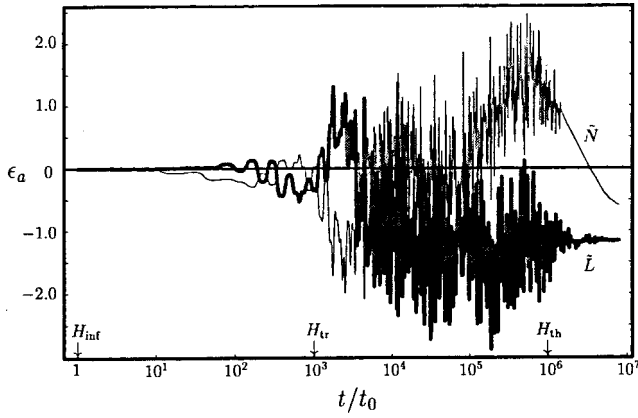


FIG. 1. Typical time variations of the asymmetries $\epsilon_{\tilde{N}}(t)$ (thin) and $\epsilon_{\tilde{L}}(t) \approx \epsilon_L(t)$ (bold) are depicted.

merically), so that it is linked again to ϕ , tracking the instantaneous minimum of V as $\tilde{N} \approx -(h_\nu/2M_N)\phi\phi$ with $|\tilde{N}| \propto |\phi|^2 \propto H^{7/4}$ in magnitude and $d\theta_{\tilde{N}}/dt \approx 2d\theta_\phi/dt \propto H^{1/4}$ in phase. Then, the $\tilde{N}$ asymmetry remaining after the transition epoch diminishes rapidly through the decays as $n_{\tilde{N}} = 2(d\theta_{\tilde{N}}/dt)|\tilde{N}|^2 \propto H^{15/4}$. It is the essential point that the decay modes (24) have almost the same rate $\Gamma_{\tilde{N}}/2$ but the opposite final lepton numbers $L = \pm 1$. This means that the $\tilde{N}$ asymmetry does not leave any significant lepton number asymmetry.

The equations of motion for $\tilde{N}$ and ϕ are solved by numerical calculations to confirm the present scenario of leptogenesis. The typical time variations of $n_{\tilde{N}}(t)$ and $n_{\tilde{L}}(t) \approx n_L(t)$ are depicted in Fig. 1 in terms of the asymmetry fractions $\epsilon_a \equiv n_a / [(3/2)H^2 M]$. Here the model parameters are taken for example as $H_{\text{inf}} = 10^{13}$ GeV, $M = 5 \times 10^{18}$ GeV, $M_N = 10^{11}$ GeV, $e^{i\delta_N} = e^{i(3/10)\pi}$, $h_\nu = 3 \times 10^{-3}$,

$c_N = 1.2$, $c_\phi = 0.8$, $b_N = 1.3e^{i(2/3)\pi}$, $a_N = 1.5e^{i(5/4)\pi}$, $a_h = 0.8e^{i(1/4)\pi}$, $T_R = 10^9$ GeV. The relevant scales, H_{inf} , H_{tr} , and H_{th} , are marked together specifying the respective epochs. We really observe the expected changes of the asymmetries through $H_{\text{inf}} \rightarrow H_{\text{tr}} \rightarrow H_{\text{th}}$, resulting in the desired lepton number asymmetry $\epsilon_L \sim 1$. Particularly, the variations of $\epsilon_{\tilde{L}}$ and $\epsilon_{\tilde{N}}$ are separated for $t > H_{\text{th}}^{-1}$; $\epsilon_{\tilde{L}}$ is fixed to some significant value while $\epsilon_{\tilde{N}}$ oscillates slowly in $\ln t$ as given in Eq. (23). It is also checked that the sum $n_{\tilde{N}} + n_{\tilde{L}}$ varies rather moderately in the transition epoch, while the respective asymmetries oscillate rapidly.

In summary, we have investigated the leptogenesis via multiscalar coherent evolution in the supersymmetric seesaw model. The right-handed sneutrino $\tilde{N}$ and the ϕ field in $\tilde{L}H_u$ of the slepton and Higgs doublets together start coherent evolution after the inflation with M_N smaller than H_{inf} . Then, after some period the motion of $\tilde{N}$ and ϕ is drastically changed by the cross coupling $M_N h_\nu \tilde{N}^* \phi \phi$, and the significant asymmetries of $\tilde{N}$ and $\tilde{L}$ are generated. The $\tilde{L}$ asymmetry is fixed later by the thermal effect as the lepton number asymmetry n_L . The $\tilde{N}$ asymmetry, on the other hand, disappears through the incoherent decays $\tilde{N} \rightarrow \bar{L}\tilde{H}_u, \tilde{L}H_u$ with almost the same rate but opposite final lepton numbers. The sufficient amount of n_L for baryogenesis can be obtained with the lightest neutrino mass $m_\nu \lesssim 10^{-3}$ eV given by the seesaw mechanism with the right-handed neutrino mass $M_N \sim 10^{10} - 10^{14}$ GeV $\lesssim H_{\text{inf}}$.

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