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Throwing away antimatter via neutrino oscillations during the reheating era

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Abstract. The simplest possibility to explain the baryon asymmetry of the Universe is to assume that radiation is created asymmetrically between baryons and anti-baryons after the inflation. We propose a new mechanism of this kind where CP-violating flavor oscillations of left-handed leptons in the reheating era distribute the lepton asymmetries partially into the right-handed neutrinos while net asymmetry is not created. The asymmetry stored in the right-handed neutrinos is later washed out by the lepton number violating decays, and it ends up with the net lepton asymmetry in the Standard Model particles, which is converted into the baryon asymmetry by the sphaleron process. This scenario works for a range of masses of the right-handed neutrinos while no fine-tuning among the masses is required. The reheating temperature of the Universe can be as low as $\mathcal{O}(10)$ TeV if we assume that the decays of inflatons in the perturbative regime are responsible for the reheating. For the case of the reheating via the dissipation effects, the reheating temperature can be as low as $\mathcal{O}(100)$ GeV.

Keywords: baryon asymmetry, leptogenesis, particle physics - cosmology connection, neutrino masses from cosmology

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Contents

1 Introduction

Missing antimatter is one of the mysteries in the history of the Universe. The baryon asymmetry cannot be the initial condition in the inflationary cosmology while the thermal history within the Standard Model of particle physics seems to fail to explain it. It is plausible that the neutrino masses may be something to do with this mystery since the Majorana masses of neutrinos together with the sphaleron process provides us with a new source of the baryon number violation as well as CP violation [1].

It is well-known that three conditions need to be satisfied for the creation of the baryon asymmetry after inflation [2]: baryon number violation, CP violation, and a stage with out-of-equilibrium. Recently it has been shown that all the three conditions can be satisfied at the very beginning of the Universe in the Standard Model with a dimension five operator to generate the neutrino Majorana masses [3, 4]. (See also the reheating era baryogenesis [3, 5, 6], and non-thermal leptogenesis [7–9].) In these scenarios, the right-handed neutrinos are not necessary, and thus the effective theory to describe the phenomena is the same as the one for the low energy experiments, such as the neutrino oscillation experiments as well as the neutrinoless double beta decays. This provides us with tight connections between the baryon asymmetry of the Universe and the low energy experiments.

The key fact is that flavor oscillations of active neutrinos during the reheating era provide CP violation just as in the neutrino oscillation phenomena we observe today [3]. It was shown that only with adding the higher dimensional Majorana mass term, $LLHH$, to the renormalizable Standard Model, the baryon asymmetry can be generated during the thermalization process [4]. The oscillations are induced due to the misalignment of the eigenbasis of the effective mass matrices, governed by the matter effects, and that of the $LLHH$ interactions. The observed baryon asymmetry is shown to be explained if the reheating temperature is higher than 10^8 GeV. The scenario works for whatever mechanism for the generation of neutrino Majorana masses at sufficiently high renormalization scale.

The baryogenesis with two or three generations of the right-handed neutrinos have been studied widely in connection with the generation of the neutrino masses by the seesaw mechanism [10–14]. (See also refs. [15, 16].) Thermal leptogenesis [1] assumes the thermal bath of the Standard Model including the right-handed neutrinos as the initial condition, and the lepton asymmetry is produced by the out-of-equilibrium decays of right-handed neutrinos, which requires $T_R \gtrsim 10^8$ GeV. (See refs. [17, 18] for reviews.) It has been shown that the reheating temperature can be lower if one tunes the difference between the right-handed neutrino masses so that the resonant effects take place [19, 20]. The right-handed neutrinos can be much lighter than $\mathcal{O}(100)\,\mathrm{GeV}$ when the flavor oscillations of the right-handed neutrinos get important while some tuning in the mass spectrum is necessary [21, 22]. In this scenario, the abundance of the right-handed neutrinos, which is assumed to be zero at the beginning of the Universe, are generated through the scattering of the left-handed leptons. The lepton asymmetries originated from the oscillation among right-handed neutrinos are stored separately into the left-handed and right-handed neutrino sectors. In the case of the neutrinos with the Dirac masses, one can also consider the possibility that the lepton asymmetry is stored in the right-handed neutrinos [23].

In this paper, we consider flavor oscillations of active neutrinos during the reheating era in the seesaw model [10–14]. Through the Yukawa interactions between the lepton doublets and the right-handed neutrinos, the CP-violating flavor oscillations distribute lepton asymmetries into the left-handed and right-handed neutrinos, while total lepton asymmetry is conserved. The lepton asymmetry stored in the left-handed leptons is, in turn, converted into the baryon asymmetry by the sphaleron process. If the right-handed neutrinos never come into the thermal equilibrium until the sphaleron process shuts off at $T \sim 100\,\mathrm{GeV}$, the created baryon asymmetry remains today. In the scenario where the reheating is caused by the perturbative decay of the inflaton, the reheating temperature should satisfy $T_R \gtrsim 7\,\mathrm{TeV}$ for the successful baryogenesis. We also discuss the possibility that the reheating is due to the dissipation effect. In that case the right-handed neutrino can be lighter than $100\,\mathrm{GeV}$ and the reheating temperature can be as low as $T_R \sim 100\,\mathrm{GeV}$.

The new mechanism does not require a fine-tuning of the mass degeneracy. Since the density matrices of the initial left-handed neutrinos are not in the thermal ones in the reheating era, the asymmetry via oscillation is produced at the leading order in the perturbation of the neutrino Yukawa couplings. As a result, large enough baryon asymmetry can be produced. The mechanism works with a single right-handed neutrino, and thus no tuning among masses of right-handed neutrinos is necessary.

This paper is organized as follows. The main idea is shown in the section 2 by assuming the inflaton decays in the perturbative regime. In section 3, we discuss the case with reheating via the dissipation processes. The last section is devoted to conclusions and discussion.

2 Active neutrino oscillation for baryogenesis

We introduce a singlet fermion to the Standard Model gauge group, N, which is one of three right-handed neutrinos responsible for the seesaw mechanism. For a while, we ignore its mass. The Lagrangian is given as

$$\mathcal{L} \supset -y_{Ni}\tilde{H}^*\bar{N}\hat{P}_L L_i, \tag{2.1}$$

where y_{Ni} are the Yukawa coupling constants, L_i $(i = e, \mu, \tau)$ is the lepton doublet field, $\hat{P}_L$ is a left-handed projection operator and $\tilde{H} \equiv i\sigma_2 H$ is the Higgs field. We restrict ourselves in the case for $|y_{N_i}| \ll \mathcal{O}(1)$. Here we take the basis that y_{Ni} is real by making the phase rotation of L_e, L_μ and N without loss of generality.

2.1 Inflaton decay in the perturbative regime

We introduce an inflaton field, ϕ, which once dominates over the Universe. The mass is m_ϕ. Let us first assume that the reheating of the Universe proceeds via the ϕ perturbative decays for simplicity. We suppose that the decay has some branching fraction B to the active neutrinos,

$$\phi \to L_\phi + X, \quad \bar{L}_\phi + \bar{X}. \tag{2.2}$$

Here X denotes arbitrary final states, that is supposed not to contribute to the baryon asymmetry, and $B \leq 1$. The final lepton state, L_ϕ, is in general a linear combination of L_e, L_μ, and L_τ. Through the dominant decay channels of ϕ the Universe is reheated to the temperature, $T = T_R \simeq (g_* \pi^2/90)^{-1/4}\sqrt{\Gamma M_P}$, with the total decay width Γ, the effective relative degrees of freedom $g_* \simeq 106.75$, and the reduced Planck mass $M_P \simeq 2.4 \times 10^{18}\,\text{GeV}$.

At the moment of the inflaton decay $t = t_R \equiv 1/\Gamma$, there are two components in the Universe. One is the thermal plasma which is generated at $t < t_R$.[1] The thermal distribution is characterized by the temperature T_R, which should satisfy

$$T_R \leq m_\phi \tag{2.3}$$

for the regime of the perturbative decay. Another component is the direct decay product at $t = t_R$ which includes the active leptons, L_ϕ. These leptons are generally out of equilibrium. For instance, if we consider a two-body decay to the lepton, the component includes monochromatic modes of the leptons with energy around $m_\phi/2$. The lepton will be thermalized promptly due to the interaction with the thermal plasma.

In the following, we will discuss the leptogenesis via the active lepton/neutrino oscillations during this rapid thermalization process and show that this scenario works with low reheating temperature if there is a sufficient amount of CP violation. The lepton asymmetries are divided into two sectors: the active neutrino sector, Δ_{vis}, and N sector, Δ_N, while the total asymmetry is zero, i.e. $\Delta_{\text{vis}} + \Delta_N = 0$, due to the conservation of the lepton number once we ignore the mass term of N. After the thermalization of left-handed leptons only Δ_{vis} is important and can be converted into the baryon asymmetry by the sphaleron process. If N is not thermalized until the temperature drops to $T < T_{\text{sph}} \sim 100\,\text{GeV}$, where the sphaleron process freezes out, Δ_N would not be transferred back into the visible sector. As a result, the produced baryon asymmetry is maintained until today.

The asymmetry Δ_N is produced in the following way. At $t = t_R$, the lepton of momentum $\mathbf{p}$ produced by an inflaton decay, L_ϕ, is represented as a quantum state

$$|L_\phi, t_R\rangle, \tag{2.4}$$

which evolves as

$$|L_\phi, t\rangle = \sum_i c_i \exp\left[-i\int_{t_R}^{t} E_i dt'\right] |i\rangle \tag{2.5}$$

where $|i\rangle$ is the flavor eigenstate of the left-handed leptons, $i = e, \mu, \tau$, with momentum p which is around m_ϕ. We have defined

$$c_i \equiv \langle i | L_\phi, t_R\rangle. \tag{2.6}$$

[1]The Boltzmann equation for the inflaton density, ρ_ϕ, is $d\rho_\phi/dt + 3H\rho_\phi = -\Gamma\rho_\phi$. At $t < t_R$, plasma (including the leptons) with energy density $\rho_r \sim \Gamma t \rho_\phi$ is produced. It is most abundant at around one Hubble time, $t \sim 1/H_{\text{inf}}$ with H_{inf} being the Hubble parameter during inflation, after the inflation. The plasma is thermalized immediately by fast interactions among the Standard Model particles. At the last period of the reheating, which we focus with $t \sim 1/\Gamma$, one finds $\rho_r \sim \rho_\phi$ i.e. the preexisting plasma is as abundant as the inflaton. Therefore, the preexisting plasma exists in general. The baryon asymmetry is also produced at $t \ll t_R$ by the mechanism described here but the amount is suppressed due to the dilution from copiously produced plasma. This is the reason we focus the last period. The $\mathcal{O}(1)$ uncertainty to the baryon asymmetry induced by this sudden decay approximation can be absorbed into the definition of T_R.

Here c_τ can be taken to be real by the field redefinition of L_τ without loss of generality, but c_e and c_μ are in general complex numbers. Notice that the state $|L_\Phi, t_R\rangle$ is not necessary to be the same as the flavor state $|i\rangle$, and in general it is in the superposition of different flavor states. Hereafter, we assume the general case. The flavor oscillation phenomena happen through thermal potentials, which are created by the preexisting thermal plasma.[2] For $|\mathbf{p}| \gtrsim T$, the dispersion relation becomes flavor dependent such as [24]

$$E_i \simeq y_i^2 \frac{T^2}{16|\mathbf{p}|} + \cdots, \quad (i = e, \mu, \tau), \tag{2.7}$$

where y_i are the Yukawa coupling constant for the charged leptons, $y_i = m_i/\langle H\rangle$. We assumed $y_\tau \gg y_N$, and '$\cdots$' contains the flavor-blind terms irrelevant for the flavor oscillation.

The thermal plasma plays two important roles. One is to induce the thermal potential for the flavor oscillation as just discussed. The other is that it prevents the flavor oscillation from lasting too long. The oscillation is terminated when the leptons annihilate with the plasma. The free propagation time scale, $t_{\rm MFP}$, is given approximately as the inverse of the thermalization rate,

$$t_{\rm MFP}(\mathbf{p}) \simeq \Gamma_{\rm LPM}^{-1} \simeq \left(\alpha_2^2 T \sqrt{\frac{T}{|\mathbf{p}|}}\right)^{-1}. \tag{2.8}$$

where we have taken into account the Landau-Pomeranchuk-Migdal (LPM) effects [25, 26] for estimating the energy loss process important for the thermalization. The inelastic scattering rate via a t-channel gauge boson exchange is naïvely $\mathcal{O}(\alpha_2^2 T)$. However, at the quantum level, one must take into account the coherent multiple gauge boson emissions, when an energetic lepton is injected into the medium. This effect leads to the suppression factor $\sqrt{T/|\mathbf{p}|}$.

The leptons from the inflaton decays lose the energy and settle down to a state with $|\mathbf{p}| \sim T$ after traveling in the plasma for a typical time scale of thermalization

$$t_{\rm th} \equiv t_{\rm MFP}(m_\phi). \tag{2.9}$$

The scattering via gauge interactions does not touch the flavor and the flavor oscillation continues even after the scattering while the pair annihilation of the leptons via gauge interactions terminates the oscillation. Notice that when the energy $|\mathbf{p}|$ is larger than the temperature, the annihilation rate

$$\Gamma_{\rm pair} \sim \alpha_2^2 \frac{T^2}{|\mathbf{p}|} \tag{2.10}$$

is smaller than the energy loss rate $\Gamma_{\rm LPM}$ for given $\mathbf{p}$. The pair annihilation happens most effectively after the energy of the lepton drops down to $|\mathbf{p}| \sim T = T_R$. The flavor oscillation is also the most effective for $|\mathbf{p}| \sim T$. The time scale of the pair annihilation after the energy loss is given as

$$\Delta t_{\rm pair} \equiv \Gamma_{\rm pair}^{-1}\Big|_{|\mathbf{p}|\sim T} \sim (\alpha_2^2 T)^{-1}, \tag{2.11}$$

[2]The coherent forward scatterings with negligible momentum exchanges, $|\Delta\mathbf{p}| \sim 0$, induce the thermal potentials. Since the exchanged momentum $|\Delta\mathbf{p}|$ is much smaller than T, the scattering takes place coherently with multiple particles in the plasma. Thus we can safely use the thermal field theory approach to estimate the thermal potentials. In particular, the thermal potential can be obtained from the two point function of the lepton fields in the thermal environment. We thank Markus Luty for useful discussion on this point.

which is even shorter than the time scale of the thermalization, $t_{\rm th}$. Therefore, the quantum state of the leptons shortly after the time scale, $t_{\rm th}$, is given by

$$|L_\phi, t_R + t_{\rm th}\rangle \simeq \sum_i c_i \exp\left[-i\frac{y_i^2}{16\alpha_2^2} + \dots\right]|i\rangle. \tag{2.12}$$

The integration in eq. (2.5) is approximated by $E_i \Delta t_{\rm pair}$ evaluated at $|\mathbf{p}| = T$. The evolution of each flavor component differs by a phase, and for τ the difference is,

$$\frac{y_\tau^2}{16\alpha_2^2} \sim 0.005. \tag{2.13}$$

This can be the origin of the baryon asymmetry $\mathcal{O}(10^{-10})$. The effects are not suppressed by a ratio of the neutrino masses or charged lepton masses to the energy scale of the problem, m_ϕ or T_R. The matter effects in the finite temperature plasma make it possible to induce the large quantum oscillation phenomenon. We emphasize here that even though the oscillation is stopped by the time scale of the pair annihilations, the density matrices in the flavor space are still not collapsed into the flavor eigenbasis until the Yukawa interactions get important.

After the evolution of the quantum state, the flavor is "observed" by the flavor dependent interaction with the thermal plasma. At this stage, the lepton state is identified as one of the flavors, e, μ or τ by the Yukawa interactions of the charged leptons. As a rare process, however, the flavor can be "observed" by the neutrino Yukawa interaction in eq. (2.1). The observation through the Yukawa interaction of eq. (2.1) happens at the probability of

$$\eta \sim \frac{\sum |c_i|^2 \sigma_{\nu_i t_L \to N t_R}}{\sum |c_i|^2 \sigma_{\nu_i t_L \to \tau_R b_R}} \sim \frac{|y_N|^2}{y_\tau^2}, \tag{2.14}$$

where we have defined $|y_N|^2 \equiv \sum_i |y_{N_i}|^2$. The probability is normalized by the process with the largest cross section, i.e., the scattering via y_τ.

As in the ordinary neutrino oscillation this rare process can have CP asymmetry since there are strong phases (CP-even phases) from the oscillation and the CP-odd phases in the new interactions including the inflaton couplings. This is because we cannot remove all of the CP phases from the field redefinition as we have performed. The CP asymmetry in the probability is given by

$$P_{L_\phi \to L_N} - P_{\overline{L}_\phi \to \overline{L}_N} = \eta \left(|\langle L_N | L_\phi, t_R + t_{\rm th}\rangle|^2 - |\langle \overline{L_N} | \overline{L}_\phi, t_R + t_{\rm th}\rangle|^2\right). \tag{2.15}$$

Here we have defined the state $\langle L_N|$ as the eigenstate in the interaction basis of eq. (2.1) which satisfies $\langle L_N | i\rangle = y_{Ni}/|y_N|$. Thus,

$$\langle L_N | L_\phi, t_R + t_{\rm th}\rangle \simeq \sum_i c_i \exp\left[i\frac{y_i^2}{16\alpha_2^2} + \dots\right]\frac{y_{Ni}}{|y_N|}. \tag{2.16}$$

The probability is estimated as

$$P_{L_\phi \to L_N} - P_{\overline{L}_\phi \to \overline{L}_N} \simeq \sum_{i>j} 4\Im[c_i c_j^*] \sin\left(\frac{y_i^2 - y_j^2}{16\alpha_2^2}\right)\frac{y_{N_i} y_{N_j}}{y_\tau^2} \sim \frac{c_\tau y_{N_\tau} \sum_{i=e,\mu} \Im[c_i^* y_{N_i}]}{4\alpha_2^2}. \tag{2.17}$$

The leptonic asymmetry in N is produced with this probability for each leptons generated by the inflaton decays.

Since the inflaton decays provide the leptons in terms of the number density divided by the entropy density as $\sim 3BT_R/4m_\phi$, Δ_N to entropy density is given as

$$\begin{aligned}\frac{\Delta_N}{s} &\simeq \frac{3}{4}\frac{T_R}{m_\phi}B \times \left(P_{L_\phi \to L_N} - P_{\overline{L}_\phi \to \overline{L}_N}\right)\\ &\simeq \frac{3}{4}\frac{T_R}{m_\phi}B\xi_{CP}\frac{|y_N|^2}{4\alpha_2^2}\\ &= 10^{-10}\frac{T_R}{m_\phi}B\xi_{CP}\left(\frac{|y_N|}{10^{-6}}\right)^2. \end{aligned} \tag{2.18}$$

Here we have defined

$$\xi_{CP} \equiv \frac{c_\tau y_{N_\tau}\sum_{i=e,\mu}\Im[c_i^*]y_{N_i}}{|y_N|^2}.$$

The CP asymmetry parameter $\xi_{CP} = \mathcal{O}(1)$ unless one tunes the CP phases, flavor-structure of the left-handed lepton coupling to inflaton or N. In fact neutrino oscillation measurements have revealed the violation of flavor symmetry in the left-handed lepton sector, and also suggested the maximal CP-violation [27]. We stress here that this value of generated asymmetry does not depend on T_R once T_R/m_ϕ is fixed. This implies that the reheating temperature has no restriction in generating Δ_N.

Since $\Delta_N = -\Delta_{\rm vis}$, the non-zero Δ_N means that there exists

$$\frac{\Delta_{\rm vis}}{s} \sim -10^{-10}\frac{T_R}{m_\phi}B\xi_{CP}\left(\frac{|y_N|}{10^{-6}}\right)^2. \tag{2.19}$$

This is transferred into the baryon asymmetry via the sphaleron process. The required value of the lepton asymmetry converted from the measured baryon asymmetry of the universe [28–30] is

$$\left(\frac{\Delta_{\rm vis}}{s}\right)^{\rm required} = -(2.45 \pm 0.01) \times 10^{-10}. \tag{2.20}$$

Comparing with eq. (2.18), we see that enough amount of baryon asymmetry can be generated just after the reheating. The question is whether this asymmetry remains until today.

Let us consider the condition for preserving the baryon asymmetry until today. Obviously, Δ_N should not be transferred back to the visible sector via eq. (2.1) until the sphaleron process becomes inefficient at the temperature lower than $T_{\rm sph}$. Otherwise, the sphaleron would washout the baryon asymmetry. Therefore, N should not be thermalized until $T = T_{\rm sph}$. The interaction rate of relativistic N with the thermal plasma is given by

$$\Gamma_N^{\rm th} \simeq \gamma_N |y_N|^2 T \tag{2.21}$$

where $\gamma_N \simeq 0.01$ is the numerical result from refs. [31–33] which includes $2 \leftrightarrow 2$ and $1 \leftrightarrow 2$ processes as well as the LPM effect. By comparing $\Gamma_N^{\rm th}$ with the Hubble parameter at the radiation dominant era, $H \simeq \sqrt{g_* \pi^2 T^4/90M_P^2}$, one obtains the temperature that N is thermalized

$$T \lesssim T_{\rm th}^N \simeq 7\,\text{TeV}\left(\frac{|y_N|}{10^{-6}}\right)^2. \tag{2.22}$$

The thermalization of N can be avoided if we take into account its Majorana mass parameter, M_N, as

$$\delta\mathcal{L} = -\frac{M_N}{2}\bar{N}^c N \tag{2.23}$$

which satisfies

$$M_N \gtrsim T_{\rm th}^N. \tag{2.24}$$

In this case, before the thermalization occurs N becomes non-relativistic so that the asymmetry, Δ_N, is washed-out while the produced baryon asymmetry corresponding to $\Delta_{\rm vis}$ untouched. On the other hand, $T_R \gtrsim M_N$ is necessary for our discussion, i.e. the "observation" of active states with producing N has to be valid kinematically. One arrives at the condition for our scenario in inflaton perturbative decay

$$T_R \gtrsim M \gtrsim T_{\rm th}^N \simeq 7\,{\rm TeV}\left(\frac{|y_N|}{10^{-6}}\right)^2. \tag{2.25}$$

This condition predicts specific patterns of the both active and right-handed neutrino masses and the relating phenomena, as we shall see soon.

2.2 Implications on neutrino physics

Since N is the right-handed neutrino, it gives a mass of active neutrino through the type-I seesaw mechanism,

$$\delta m_\nu = \frac{|y_N|^2 \langle H\rangle^2}{M_N}. \tag{2.26}$$

Here $\langle H\rangle \simeq 174\,{\rm GeV}$ is the Higgs vacuum expectation value. Substituting the condition for the baryogenesis eq. (2.25) to eq. (2.26) one can estimate the active neutrino mass as

$$\delta m_\nu \lesssim \frac{|y_N|^2 \langle H\rangle^2}{T_{\rm th}^N} \simeq 4\times 10^{-3}\,{\rm eV} \quad ({\rm for}\ M_N \gtrsim T_{\rm th}^N). \tag{2.27}$$

Compared to the neutrino mass scales, $\sqrt{\Delta m^2_{\rm sol}} \simeq 9\times 10^{-3}$ eV and $\sqrt{\Delta m^2_{\rm atm}} \simeq 5\times 10^{-2}$ eV, the N particle which is responsible for baryogenesis can significantly contribute to the active neutrino masses only for the lightest or the second lightest ones. Two other right-handed neutrinos need to explain the rest of the neutrino masses.

Based on the above discussion, the baryogenesis scenario predicts the active neutrinos in either normal hierarchy (NH) or inverted hierarchy (IH), i.e. not degenerated. The sum of the active neutrino masses is determined for each mass hierarchy, $\sum_I m_{\nu_I} \simeq 0.06\ (0.10)\,{\rm eV}$ for the NH (IH)[3] with $I = 1, 2, 3$ denoting the generation of active neutrinos in the mass basis. The sum of masses has been constrained by the observations of cosmic microwave background (CMB) and baryonic acoustic oscillation given as $\sum_I m_{\nu_I} < 0.12\,{\rm eV}$ [35]. The value is consistent with eq. (2.27). The future observations should improve the upper bound so that the scenario can be tested.

The prediction on the lightest neutrino mass in eq. (2.27) impacts on neutrinoless double beta decay. Its decay rate is characterized by the effective neutrino mass $m_{\rm eff}$, whose definition is $m_{\rm eff} = \sum_I m_{\nu_I}[U_{\rm PMNS}]^2_{eI}$. Since the lightest neutrino mass is at most $|\delta m_\nu|$, we find

$$|m_{\rm eff}| \lesssim 7\times 10^{-3}\ {\rm eV}\ \ \text{for NH}\ \ \text{and}\ \ 0.01\ {\rm eV} \lesssim |m_{\rm eff}| \lesssim 0.05\ {\rm eV}\ \ \text{for IH}. \tag{2.28}$$

The masses of other two right-handed neutrinos are restricted since they should not wash out the lepton asymmetry $\Delta_{\rm vis}$. Here for simplicity we restrict ourselves in the case

[3]For the estimation of the total neutrino mass we use results in a global analysis of neutrino oscillation measurements [34].

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that the reheating temperature is so low, e.g. $T_R \lesssim 100\,\text{TeV}$, that all the interaction rates via Standard Model Yukawa couplings are faster than the expansion rate.[4] Under this most dangerous circumstance for the wash out, one can obtain four possible mass patterns to evade the wash out:

$$\text{Case 1}: \quad M_{2,3} \ll 100\,\text{GeV}, \tag{2.29}$$

$$\text{Case 2}: \quad T_R \ll M_{2,3}, \tag{2.30}$$

$$\text{Case 3}: \quad M_2 \ll 100\,\text{GeV} \ll T_R \ll M_3, \tag{2.31}$$

$$\text{Case 4}: \quad M_3 \ll 100\,\text{GeV} \ll T_R \ll M_2. \tag{2.32}$$

The Case 1 says that the masses of $N_{2,3}$ that violates the lepton number are almost negligible at $T > T_{\text{sph}}$. Case 2 is the possibility that $N_{2,3}$ are so heavy that the thermal production are kinematically suppressed, and that they are hardly thermalized. Case 3 or 4 is the composition of Cases 1 and 2. Apart from Case 2 there exists the right-handed neutrino with the mass below the electroweak scale.

Such a light particle may be probed experimentally. The right-handed neutrinos in the mass range of $\mathcal{O}(1\text{–}10)$ GeV will be searched for in future beam-dump and collider experiments e.g. ref. [36]. Moreover such a particle can impact on the neutrinoless double beta decay as an additional intermediate state and thus may be tested indirectly [37].[5] The contribution behaves as M_α^{-2} because of a suppression of nuclear matrix element when the mass is larger than its typical momentum exchange $\sqrt{\langle p^2 \rangle} \sim 200\,\text{MeV}$.

2.3 Numerical estimation

Here we perform a numerical simulation to confirm the discussion on the asymmetry separation. See ref. [4] for the detail analysis. We focus on the two components of the density matrices:

$$(\rho_{\mathbf{k}})_{ij} = \int_{|\mathbf{p}|\sim|\mathbf{k}|} \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{\rho_{ij}(\mathbf{p},t)}{s}, \tag{2.33}$$

$$(\delta\rho_T)_{ij} = \int_{|\mathbf{p}|\sim T} \frac{d^3\mathbf{p}}{(2\pi)^3} \left(\frac{\rho_{ij}(\mathbf{p})}{s} - \frac{\rho_{ij}^{\text{eq}}(\mathbf{p})}{s} \right), \tag{2.34}$$

and those for anti-leptons. Here s is the entropy density. The first component, $\rho_{\mathbf{k}}$, represents the energetic leptons produced by the ϕ decay with initial typical momentum, $|\mathbf{k}| = m_\phi$. The second component, $\delta\rho_T$, represents leptons that deviate from the thermal distribution with the typical momentum $|\mathbf{p}| \sim T$. Here $T \simeq T_R$ is the temperature, and $\rho_{ij}^{\text{eq}} = \delta_{ij}/(e^{|\mathbf{p}|/T} + 1)$ represents the density matrix in the thermal equilibrium, which denotes the preexisting thermal plasma. We did not write down the equation for the right-handed neutrino since $\Delta_N = -\Delta_{\text{vis}}$ is guaranteed and we will estimate Δ_{vis}.

[4]If this assumption is removed, some of the charged lepton Yukawa coupling can be neglected when the washout is effective. There can be flavor-dependent lepton symmetry, and thus some component of Δ_{vis} can not be washed out. In this case, there can be mass patterns where M_2 or M_3 is below T_R. The extension is straightforward.

[5]If two degenerate right-handed neutrinos affect the process simultaneously the contribution to the decay rate is always destructive [38].

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The time evolutions of the matrices can be obtained by solving the kinetic equations, which are derived from first principle with approximations [39]. The equations are given as

$$i\frac{d\rho_{\mathbf{k}}}{dt} = [\Omega_{\mathbf{k}}, \rho_{\mathbf{k}}] - \frac{i}{2}\{\Gamma^d_{\mathbf{k}}, \rho_{\mathbf{k}}\}, \tag{2.35}$$

$$i\frac{d\delta\rho_T}{dt} = [\Omega_T, \delta\rho_T] - \frac{i}{2}\{\Gamma^d_T, \delta\rho_T\} + i\delta\Gamma^p_T, \tag{2.36}$$

where $\Omega_{\mathbf{k}} = E_i(|k|)\delta_{ij}$ and $\Omega_T = E_i(T)\delta_{ij}$. The destruction and production rates for leptons are given by

$$\left(\Gamma^d_{\mathbf{k}}\right)_{ij} \simeq C\alpha_2^2 T\sqrt{\frac{T}{|\mathbf{k}|}}\delta_{ij} \tag{2.37}$$

$$\left(\Gamma^d_T\right)_{ij} \simeq C'\alpha_2^2 T\delta_{ij} + \gamma_L T(\delta_{i\tau}\delta_{\tau j}y_\tau^2 + \delta_{i\mu}\delta_{\mu j}y_\mu^2 + \delta_{ie}\delta_{ej}y_e^2) + \gamma_N T y_{N_i} y_{N_j}, \tag{2.38}$$

$$\left(\delta\Gamma^p_T\right)_{ij} \simeq C\alpha_2^2 T\sqrt{\frac{T}{|\mathbf{k}|}}\,(\rho_{\mathbf{k}})_{ij} - C'\alpha_2^2 T\,(\delta\bar{\rho}_T)_{ij}\,. \tag{2.39}$$

The equations for the anti-leptons are obtained by replacing ρ with $\bar{\rho}$ everywhere and reversing the sign of Ω's. In the actual numerical computation, the kinetic equations of right-handed leptons and the red-shift of momenta are taken into account [4].

Now let us briefly explain the terms in the rates. The terms with the coefficient C, C' corresponds to the $\Gamma_{\rm LPM}, \Gamma_{\rm pair}$ respectively. We put a parameter $C, C' = \mathcal{O}(1)$ to take into account the theoretical uncertainty in the LPM effects and in the energy distributions of the inflaton decay product. The terms with γ_L, γ_N describe the scattering, and (inverse) decay via the Yukawa interactions of charged lepton and eq. (2.1), respectively, which are important for "observing" the flavor. Here, we take

$$\gamma_L = \gamma_N$$

because of the same kinematics and gauge structure if we neglect g_Y. This term with γ_N divides the lepton asymmetries into the two sectors. The numerical result of γ_N can be found from [31–33]. We have neglected the scattering via Yukawa interactions for the high energy mode since it is much slower than the energy loss process.

We have approximated that the Higgs bosons are in thermal distributions, which is justified as follows. If the asymmetry of Higgs produced by the CP-violating effect is transferred into the right-handed neutrinos, the contribution is suppressed by $\mathcal{O}(y_N^4)$. On the other hand, most left-handed leptons produced from the interaction with Higgs bosons are in flavor eigenstates, and that the flavor oscillations are suppressed. Thus we can safely neglect the out-of-equilibrium effects of the Higgs bosons.

Now we are ready to solve the kinetic equations. The initial conditions of the density matrices for our scenario are as follows

$$\rho_{\mathbf{k}}|_{t=t_R} = \bar{\rho}_{\mathbf{k}}|_{t=t_R} = \frac{3}{4}\frac{T_R}{m_\phi}Bc_i^* c_j, \quad \delta\rho_T|_{t=t_R} = \delta\bar{\rho}_T|_{t=t_R} = 0. \tag{2.40}$$

We have assumed the absence of the deviation from the thermal equilibrium for the preexisting thermal plasma at $t = t_R$. We take $\rho_{\mathbf{k}}$ corresponding to eq. (2.6).

In figure 1, we show the asymmetry $\Delta_L/(sB)$ by varying the Hubble time $(t - t_R)/t_R$ with $T_R = 10^7$ GeV, $m_\phi = T_R$, and $\gamma_N = \gamma_L = 0.01$. The vertical purple dashed and blue solid

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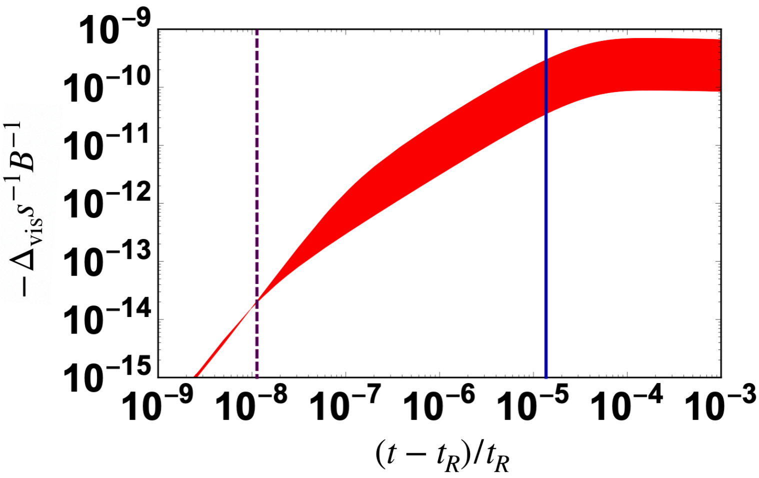

Figure 1. The time evolution of the lepton asymmetry with $B = 1, m_\phi = T_R = 10^7\,\mathrm{GeV}$ and $y_N = 10^{-6}(1, 1, 1)$, where $c_i = \frac{1}{\sqrt{3}}(\exp i, \exp 2i, 1)$ is taken as an example. The red band corresponds to the variance of $C = C'$ between 1/3 and 3. The blue solid and purple dashed lines are the time scale of $t_{\rm th}$ and $1/(\gamma_L y_\tau^2 T_R)$, respectively, at which the pair annihilation and the scattering with the tau Yukawa coupling are important.

lines represent the time scale of $t_{\rm th} = t_{\rm MFP}(m_\phi)$ and that the density matrices are collapsed into the flavor eigenbasis due to the charged τ-Yukawa interaction, $(\gamma_L y_\tau^2 T)^{-1}$. One finds that the asymmetry production lasts much shorter than $t_R \simeq 1/H|_{t=t_R}$, and the behavior around the timescales are consistent with the discussions in the previous section. We have also checked that the amount of asymmetry does not change much by changing the reheating temperature.

3 Case with reheating via dissipation processes

Since the mechanism discussed previously is tied to the inflaton sector, for certain reheating dynamics, the asymmetry can be enhanced significantly. In what follows, we consider the reheating scenario where $T_R/m_\phi \gg \mathcal{O}(1)$. For the inflaton perturbative decays, it was a thermal blocking effect that prevents the T_R from going beyond m_ϕ. This is because the decays of ϕ to the daughter particles with thermal mass $\sim T_R \gtrsim m_\phi$ are kinematically forbidden. Thus even if we increase the couplings among inflaton and particles, we can have at most $T_R \sim m_\phi$ from the perturbative decays. However, then a dissipation effect becomes important to increase the temperature.

When the thermal blocking effect is important, a thermal dissipation effect is also important [40–45]. If the dissipation effect is efficient, the reheating proceeds through the scatterings among the inflaton condensate and preexisting thermal plasma.[6] The process can be represented as

$$\phi + Y \to X + L_\phi. \tag{3.1}$$

[6] For simplicity, we assume that the parametric resonance effect is not important. This is justified in the following ALP inflation model because the ALP has derivative couplings.

For instance, one may take Y = a Higgs boson, X = a right-handed charged anti-lepton. This process is kinematically allowed even if m_ϕ is much smaller than the thermal masses of particles. This implies the parameter region of

$$T_R \gg m_\phi \tag{3.2}$$

is also possible. In particular, the inflaton condensate loses energy of around $m_\phi \ll T_R$ per one scattering. This means when the energy of the inflaton condensate all becomes the radiation, the scatterings take place $3T_R/4m_\phi \gg 1$ times in a unit volume. Therefore, in this scenario $3T_R/4m_\phi$ leptons carrying momenta

$$|\mathbf{p}| \sim T_R$$

are produced. They are out of equilibrium and thermalized after undergoing the flavor oscillation as in the previous part. As a result the same formula of eq. (2.18) is expected, but with $T_R/m_\phi \gg 1$. In particular, when $T_R/m_\phi \gtrsim 10^2$ the baryon asymmetry can be explained with

$$|y_N| \lesssim 10^{-7}. \tag{3.3}$$

Thus,

$$T_{\rm th}^N \lesssim T_{\rm sph} \tag{3.4}$$

can be satisfied. This means that Δ_N is not transferred back to the visible sector until the sphaleron freezes out. The baryon asymmetry remains until today. Consequently, our scenario works for the reheating temperature satisfying

$$T_R \gtrsim T_{\rm sph} \sim 100\,\text{GeV}. \tag{3.5}$$

In this case, N can be identified with any of the three right-handed neutrinos. We notice that the dissipation effect is important only when the thermal plasma from the perturbative decays of ϕ is abundant. This forbids $m_\phi \to 0$, which suppressed the perturbative decay rates and thus the plasma. The question is how large T_R/m_ϕ can be in concrete inflation models.

Let us discuss this possibility in more detail by introducing ALP inflation models [46–48] where the inflaton, ϕ, is an axion-like particle (ALP). The inflation scale can be as high as $\Lambda_{\rm inf} \simeq 100\,\text{GeV}$–$100\,\text{TeV}$.[7] The ALP (effective) mass[8] $m_\phi \sim \Lambda_{\rm inf}^2/f$ and decay constant, f, have typical relation fixed by the CMB normalization of the primordial density perturbation:

$$m_\phi \sim 10^{-6} f. \tag{3.6}$$

The flavor oscillation occurs by introducing flavor-dependent couplings of ϕ to leptons responsible for the reheating. The couplings are given as

$$\delta\mathcal{L} \simeq \frac{i\phi}{f}\sum_{ij} c_{ij} y_j H^* \bar{e}_i \hat{P}_L L_i, \tag{3.7}$$

[7]The QCD axion window can be opened and the moduli problem can be alleviated due to the low-scale inflation if inflation lasts long enough and if no mixing between the inflaton and the axion [49–51]. If there is a mixing which shifts the axion phase by π, the QCD axion can be set on the hilltop and thus a heavier QCD axion dark matter than usual is also possible [52].

[8]For the ALP miracle scenario [46, 47] m_ϕ should be identified as the effective mass of the inflaton. The inflaton mass at the vacuum, on the other hand, is highly suppressed due to an upside-down symmetry so that ϕ is long-lived.

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where c_{ij} are dimensionless constants related with derivative ϕ-lepton couplings since ϕ is an ALP. The reheating through the kind of couplings is shown to be successful in refs. [47, 48] if the ϕ coupling to τ is large enough. This coupling also contributes to the dispersion relations and the scattering rates of the leptons, while the contributions, which are suppressed by $\mathcal{O}\left((y_i T_R/f)^2\right)$, are negligible. Thus eq. (2.18) holds with $|\mathbf{p}| \sim T_R \sim 10^{2-3} m_\phi$, $B = 1$, and c_i to be an eigenvector of c_{ij}. The reheating occurs instantaneously, i.e. after inflation the energy density of the inflaton promptly becomes the radiation, $T_R \sim \Lambda_{\rm inf}$, if $f \lesssim 10^8\,\text{GeV}$. As a result, the scenarios predict

$$T_R \sim 10^{2-3} m_\phi. \tag{3.8}$$

The phenomenological implications are as follows. The flavor mixing for the flavor oscillation leads to the process $\tau \to \mu/e + \phi$ if kinematically allowed. This process can be searched for in Belle II experiment [47]. $|y_N| \sim 10^{-8}$–10^{-7} is predicted from eqs. (2.18), (2.20) and (3.8). The right-handed neutrino mass is

$$M_N \simeq 6\,\text{GeV} \left(\frac{|y_N|}{10^{-7}}\right)^2 \left(\frac{0.05\,\text{eV}}{\delta m_\nu}\right). \tag{3.9}$$

The beam-dump and collider tests, as well as the enhancement of the neutrinoless double beta decay rate, are interesting as discussed previously. The dark matter candidate may be the inflaton itself if the inflaton potential has an upside-down symmetry [46, 47]. The dark matter can be searched for in the IAXO experiment [53–55]. In this case the lightest right-handed neutrino may be the candidate as well.

4 Discussion and conclusions

We have discussed a mechanism of baryogenesis through the CP violation in the flavor oscillation in the reheating era. Since a part of the lepton asymmetry is distributed to the right-handed neutrinos through the Yukawa interactions, one obtains asymmetry in the Standard Model sector, which is converted into the baryon asymmetry by the sphalerons. The scenario works for the reheating temperature of the Universe greater than $\mathcal{O}(10)$ TeV or $\mathcal{O}(100)$ GeV depending on the dynamics of inflaton. The baryogenesis mechanism can be compatible with various new physics models that favor low reheating temperature, such as supersymmetric model avoiding gravitino problem, or require low reheating temperature due to the low cutoff scale, e.g. relaxion models, models with large extra-dimensions, composite Higgs models.

We note that the mother particle producing the left-handed leptons may not be the inflaton, but the moduli, or heavy fermions that once dominate the Universe. Even in those cases, the mechanism works. One can also apply the mechanism to the asymmetric dark matter scenario [56–59]. By assuming that the matter couples to the dark matter with baryon (lepton) number preserving interaction, the matter from the inflaton decay undergoes the flavor oscillation caused by the misalignment of the oscillation and interaction basis. The matter-antimatter asymmetry is distributed to the dark matter due to the CP-violating oscillation.

We have discussed the dissipation effect which enhances the produced lepton asymmetry. This effect can also enhance the asymmetry production of the scenario in ref. [4]. For $T_R/m_\phi \gtrsim \mathcal{O}(10^3)$, $T_R \lesssim 10^7\,\text{GeV}$ is possible to explain the baryon asymmetry of the Universe.

Lastly, let us briefly mention the inflaton decays to right-handed neutrinos via e.g. $\kappa_{ij}\phi \bar{N}_i N_j$. The produced right-handed neutrinos can undergo flavor oscillations if their Majorana masses are parametrically degenerated. The flavor oscillation can also induce CP-violation. We expect larger effect on the asymmetry production than that in refs. [21, 22], because κ_{ij} need not to be flavor-diagonal and the produced N_i can be abundant. In the scenario, one can have the sterile neutrino dark matter produced from the direct inflaton decays. We, however, did not consider such inflaton couplings or the parametrically degenerated right-handed neutrino masses for simplicity. We will come back to the possibility in future.

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