CP violating decays in leptogenesis scenarios

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Abstract

We compute the CP violation in the decays of heavy electroweak singlet neutrinos, arising from both the one—loop vertex corrections and the wave function mixing. We extend the computation to the supersymmetric version of the model and discuss the implications for the generation of a lepton number asymmetry by the out of equilibrium decay of the heavy (s)neutrinos in the early Universe, to be reprocessed later in the observed baryon excess by anomalous electroweak processes.

A very attractive scenario, proposed by Fukugita and Yanagida [1], for the generation of the baryon asymmetry of the Universe, is based on the production of a lepton asymmetry by the out of equilibrium decays of heavy ($m \gg \text{TeV}$) electroweak singlet neutrinos. Their decays into light leptons and Higgs bosons can violate CP if the Yukawa couplings involved have unremovable phases, and can then lead to the production of an excess of antileptons over leptons in the final state. The lepton asymmetry so produced is then partially converted into a baryon asymmetry by anomalous electroweak processes [2], which are in equilibrium at temperatures larger than the electroweak phase transition one, and as a result, the amount of baryons present in the Universe, $n_B/n_{\gamma} \simeq 5 \times 10^{-10}$, can be accounted for [3, 4].

To postulate the existence of right–handed 'sterile' neutrinos constitutes one of the simplest and most economical extensions of the standard model, being strongly motivated by Grand Unified theories (GUT) such as SO(10). It also allows to implement the see–saw mechanism which naturally accounts for non–zero, but still very small, light neutrino masses. Moreover, the values of the light neutrino masses suggested by the MSW solution to the solar neutrino problem, the atmospheric neutrino anomaly and the hot dark matter scenarios [5], point towards an intermediate scale, $M \sim 10^9$ – 10^{13} GeV, for the right–handed neutrino masses. This further encourages the consideration of scenarios where the heavy neutrino decays act as sources for the generation of the baryon asymmetry of the Universe, since for these large masses it becomes easier to achieve the out of equilibrium conditions required for the efficient generation of an asymmetry, and also because temperatures comparable to the intermediate scale are more likely to be produced as a result of the reheating process at the end of inflation than the temperatures required in the conventional baryogenesis GUT scenarios.

In addition to the non–supersymmetric version originally considered by Fukugita and Yanagida, the extension of this scenario to the supersymmetric case has been studied by Campbell, Davidson and Olive [6], and also a related scenario where the asymmetry is produced in the decay of heavy scalar neutrinos produced non–thermally by the coherent oscillations of the scalar field at the end of inflation has been discussed by Murayama et al. [7]. A key ingredient for all these scenarios is the one–loop CP violating asymmetry

involved in the heavy (s)neutrino decay, and the reconsideration of this quantity will be the main issue of this work. As will be discussed, there are contributions to the asymmetry that have not been included previously, and we also want to discuss the different results present in the literature on this subject.

Let us start with the non–supersymmetric version of the model, with a Lagrangian given by

$$\mathcal{L} = -\lambda_{ij} \epsilon_{\alpha\beta} \overline{N_j} P_L \ell_i^{\alpha} H^{\beta} + h.c. , \qquad (1)$$

where $\ell_i^T = (\nu_i \ l_i^-)$ and $H^T = (H^+ \ H^0)$ are the lepton and Higgs doublets $(\epsilon_{\alpha\beta} = -\epsilon_{\beta\alpha},$ with $\epsilon_{12} = +1)$, and we are taking the N_j to be the Majorana mass eigenstate fields, of mass M_j . Since the scale of the heavy neutrinos is so much larger than the electroweak scale, it is reasonable to work directly in the symmetric phase where all particles except the heavy Majorana neutrinos N_j are massless, and all components of the neutral and charged complex Higgs fields are physical.

The basic quantity we are interested in is

$$\epsilon_{\ell}^{N} \equiv \frac{\Gamma_{N\ell} - \Gamma_{N\bar{\ell}}}{\Gamma_{N\ell} + \Gamma_{N\bar{\ell}}},\tag{2}$$

where $\Gamma_{N\ell} \equiv \sum_{\alpha,\beta} \Gamma(N \to \ell^{\alpha} H^{\beta})$ and $\Gamma_{N\bar{\ell}} \equiv \sum_{\alpha,\beta} \Gamma(N \to \bar{\ell}^{\alpha} H^{\beta\dagger})$ are the N decay rates (in the N rest frame), summed over the neutral and charged leptons (and Higgs fields) which appear as final states in the N decays¹. Note that at tree level

$$\Gamma_{N_i\ell} = \Gamma_{N_i\bar{\ell}} = \frac{(\lambda^{\dagger}\lambda)_{ii}}{16\pi} M_i. \tag{3}$$

The asymmetry ϵ_{ℓ}^{N} arises from the interference of the one-loop diagrams depicted in Fig. 1 with the tree level coupling. The vertex correction in Fig. 1a is the one that is usually considered, but it has however been pointed out [8, 9, 10] that the wave function piece in Fig. 1b also contributes to the asymmetry, in an amount which is typically comparable to the vertex contribution. We have computed these CP violating asymmetries and we obtain for the vertex contribution to ϵ_{ℓ}^{N}

¹For the Majorana light neutrinos, one should think of ℓ and $\bar{\ell}$ as being the left and right–handed helicities of ν .

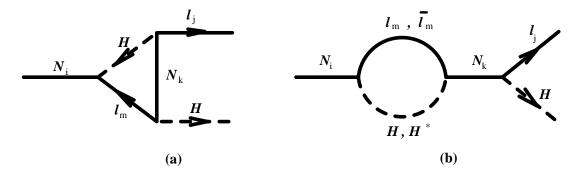


Figure 1: Diagrams contributing to the vertex (Fig. 1a) and wave function (Fig. 1b) CP violation in the heavy singlet neutrino decay.

$$\epsilon_{\ell}^{N_i}(vertex) = \frac{1}{8\pi} \sum_{k} f(y_k) \mathcal{I}_{ki} , \qquad (4)$$

where $y_k \equiv M_k^2/M_i^2$, $f(x) = \sqrt{x} (1 - (1+x)\ln[(1+x)/x])$, and we have defined

$$\mathcal{I}_{ki} \equiv \frac{\operatorname{Im}\left[\left(\lambda^{\dagger}\lambda\right)_{ki}^{2}\right]}{\left(\lambda^{\dagger}\lambda\right)_{ii}}.$$
(5)

For the wave function piece we find

$$\epsilon_{\ell_j}^{N_i}(wave) = -\frac{1}{8\pi} \sum_{k \neq i} \frac{M_i}{M_k^2 - M_i^2} \frac{\operatorname{Im}\left\{ \left| M_k \left(\lambda^{\dagger} \lambda \right)_{ki} + M_i \left(\lambda^{\dagger} \lambda \right)_{ik} \right| \lambda_{jk}^* \lambda_{ji} \right\}}{(\lambda^{\dagger} \lambda)_{ii}}.$$
 (6)

A useful check of the result can be obtained using the unitarity relation

$$|T_{fi}|^2 - |T_{if}|^2 \simeq -2 \text{ Im} \left[\sum_n T_{ni} T_{nf}^* T_{if} \right],$$

where we are retaining only the leading order terms; i and f are the initial and final states, and n are the possible intermediate states connecting them. T_{fn} are the transition amplitudes between the intermediate and final states, which in the present case arise from both the s and t-channel N_k exchanges, which correspond to wave and vertex contributions respectively.

The result in Eq. (4) coincides with the one recently obtained in ref. [4], differing by factors of 2 and 4 from those in ref. [11] and [3] respectively². The wave function piece,

²The normalization of the ϵ parameters are sometimes different (and not always explicit).

which in this scenario was only considered previously by Liu and Segrè [10], is a factor of two larger than their result due to the fact that their computation applies to the case in which the scalar field is real and only the neutral lepton runs in the loop. Adding the charged lepton loop contribution the result is doubled. In Eq. (6) we do not include in the sum over the flavour of the intermediate neutrinos the case k = i, since in general states degenerate with the initial one do not contribute to the CP violating asymmetry [12]. We also assumed that $|M_k - M_i| \gg \Gamma_{N_k}$ in the computation. Notice that the contribution from k = i in the vertex piece, being proportional to $\text{Im}\left[(\lambda^{\dagger}\lambda)_{ii}^2\right] = 0$, also vanishes. Hence, both sums in Eqs. (4) and (6) may be restricted to $k \neq i$. In Eq. (6) we have not summed over the final lepton flavour, but after summing over it one gets

$$\epsilon_{\ell}^{N_i}(wave) = -\frac{1}{8\pi} \sum_{k \neq i} \frac{M_i M_k}{M_k^2 - M_i^2} \mathcal{I}_{ki},$$
(7)

so that the piece proportional to M_i in the square brackets in Eq. (6) actually gives a null contribution to the total lepton number asymmetry (although it may generate an asymmetry in the individual leptonic numbers L_j). In ref. [10] it was pointed out that there should also be in the vertex contribution, in addition to the result in Eq. (4), a term like the one just discussed appearing in Eq. (6). This term appears however only when a single real scalar Higgs field is allowed to run in the loop, but it can be shown that the new contributions arising from the two real scalars in the decomposition of the standard model Higgs fields $H^{\alpha} = (H_1^{\alpha} + iH_2^{\alpha})/\sqrt{2}$ actually cancel each other, leaving only the term proportional to \mathcal{I}_{ki} present in Eq. (4).

In the particular case in which a hierarchy among the right-handed neutrino masses is considered, i.e. for $y_k \gg 1$, we have that the wave function contribution becomes twice as large as the vertex one, and hence the total asymmetry produced in the decay of the lightest heavy neutrino³, N_1 , becomes

$$\epsilon_{\ell}^{N_1} \equiv \epsilon_{\ell}^{N_1}(vertex) + \epsilon_{\ell}^{N_1}(wave) \simeq -\frac{3}{16\pi} \sum_{k \neq 1} \frac{1}{\sqrt{y_k}} \mathcal{I}_{k1}. \tag{8}$$

The inclusion of the wave function piece then increases by a factor of three the CP violating asymmetry in the case of large mass hierarchies [10].

³In the case of large hierarchies, the asymmetries produced by the heavier neutrino decays are usually erased before the lightest one decays.

If the departure from equilibrium in N_1 decay is large, the lepton number asymmetry produced, per unit entropy, is

$$\frac{n_L}{s} \simeq \frac{\epsilon_\ell^{N_1}}{s} \frac{g_N T^3}{\pi^2},\tag{9}$$

where $g_N = 2$ are the spin degrees of freedom of the Majorana neutrino N_1 , so that their number density (assuming Maxwell–Boltzmann statistics) before they decay is $g_N T^3/\pi^2$ if we assume that they were in chemical equilibrium before becoming non–relativistic (see below). Using that $s = (2/45)g_*\pi^2T^3$, where $g_* = g_{bosons} + (7/8)g_{fermions}$ is the effective number of relativistic degrees of freedom contributing to the entropy, ($g_* = 106.75$ in the standard non–supersymmetric version of the model), we get

$$\frac{n_L}{s} \simeq 4 \times 10^{-3} \epsilon_\ell^{N_1}. \tag{10}$$

This lepton asymmetry will then be reprocessed by anomalous electroweak processes, leading to a baryon asymmetry [13]

$$n_B = -\left(\frac{24 + 4N_H}{66 + 13N_H}\right) n_L \simeq -\frac{28}{79} n_L. \tag{11}$$

Here N_H is the number of Higgs doublets ($N_H = 1$ in the standard scenario considered now, while $N_H = 2$ in the supersymmetric version to be discussed below if the scalar Higgs bosons are assumed to be lighter than the electroweak symmetry breaking scale). Combining Eqs. (10) and (11) and assuming that the Universe expansion is adiabatic, the present baryon asymmetry which results is

$$\frac{n_B}{s} \simeq -1.5 \times 10^{-3} \epsilon_\ell^{N_1}.$$
 (12)

Values of $\epsilon_\ell^N \simeq -5 \times 10^{-8}$ are then required to account for the value $n_B/s \simeq 0.6-1 \times 10^{-10}$ inferred from the successful theory of primordial nucleosynthesis [14]. In the case in which there is a hierarchy in the heavy neutrino masses, the Yukawa parameters need then to satisfy

$$\sum_{k \neq 1} \frac{M_1}{M_k} \, \mathcal{I}_{k1} \simeq 0.9 \times 10^{-6}. \tag{13}$$

Let us also briefly mention that the lepton asymmetry produced by N decays is smaller than the value in Eq. (9) if the departure from equilibrium is not large during the decay epoch, as is the case if the decay rate $\Gamma_1 \equiv \Gamma_{N_1\ell} + \Gamma_{N_1\bar{\ell}}$ is comparable or larger than the expansion rate of the Universe when N_1 becomes non relativistic, i.e. $\Gamma_1 \geq H(T = M_1)$, with H(T) being Hubble's constant at temperature T. We also want to emphasize that a problem of these scenarios is that the Yukawa interactions are not effective in establishing an equilibrium population of heavy neutrinos. To see this, note that before N_1 becomes non-relativistic, i.e. for $x \equiv M_1/T \ll 1$, the thermally averaged rate scales as T^{-1} , since $\langle \Gamma_1 \rangle = K_1(x)/K_2(x)\Gamma_1 \simeq (x/2)\Gamma_1 \propto M_1^2/T$, where $K_{1,2}$ are Bessel functions [3]. Taking into account that $H \propto T^2/M_{Planck}$ and supposing that $\Gamma_1 \ll H(T=M_1)$ (so that the departure from equilibrium during the decay is significant), one concludes that $\langle \Gamma_1 \rangle \ll H$ holds true also for all temperatures larger than M_1 . Hence, the inverse decays are unable to establish chemical equilibrium for N_1 at any temperature $T > M_1$. The pair production of N_1 from Higgs boson or light lepton scattering may do somewhat better than inverse decays at high temperatures, since the relevant rate scales as $\langle \sigma v \rangle \propto T$, but it is anyhow insufficient to achieve chemical equilibrium for N_1 because the rates need to be quite suppressed at $T \simeq M_1$ in order not to erase any lepton asymmetry generated during the decay, and consequently are also suppressed with respect to the expansion rate at higher temperatures. Hence, additional interactions of the heavy neutrinos are required for them to have a thermal distribution prior to the decay, so that the leptogenesis scenario can be successful. This point was recently remarked in ref. [4], where it was shown that gauge interactions which are naturally present in GUT scenarios (new Z' or $SU(2)_R$ gauge bosons) can easily do the job of producing a thermal population of heavy neutrinos at $T > M_1$.

Turning now to the supersymmetric version of this scenario, the interactions of the heavy (s)neutrino field can be derived from the superpotential

$$W = \frac{1}{2}M_i N_i N_i + \lambda_{ij} \epsilon_{\alpha\beta} L_i^{\alpha} H^{\beta} N_j.$$
 (14)

Supersymmetry breaking terms are of no relevance for the mechanism of lepton number generation and we are then left with the following trilinear couplings in the Lagrangian, in terms of four component spinors,

$$\mathcal{L} = -\lambda_{ij}\epsilon_{\alpha\beta} \left\{ M_j \tilde{N}_j^* \tilde{L}_i^{\alpha} H^{\beta} + \overline{N_j} P_L \ell_i^{\alpha} H^{\beta} + \overline{(\tilde{h}^{\beta})^c} P_L \ell_i^{\alpha} \tilde{N}_j + \overline{(\tilde{h}^{\beta})^c} P_L N_j \tilde{L}_i^{\alpha} \right\} + h.c. \quad (15)$$

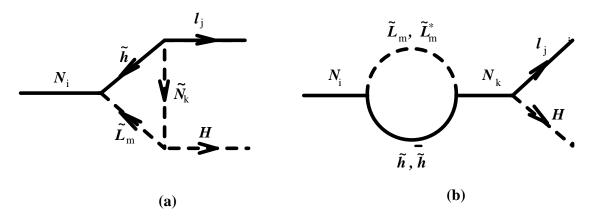


Figure 2: Supersymmetric contribution to the diagrams leading to CP violation in the $N \to \ell H$ decay.

From these couplings we obtain the tree level relations

$$\Gamma_{N_i\ell} + \Gamma_{N_i\bar{\ell}} = \Gamma_{N_i\tilde{L}} + \Gamma_{N_i\tilde{L}^*} = \Gamma_{\tilde{N}^*\ell} = \Gamma_{\tilde{N}\tilde{L}} = \frac{(\lambda^{\dagger}\lambda)_{ii}}{8\pi} M_i.$$
 (16)

There are now new diagrams, like the ones in Figure 2, contributing to the generation of a leptonic asymmetry. We will denote the corresponding asymmetry parameters in the supersymmetric case with a tilde, so that for instance $\tilde{\epsilon}_{\ell}^{N}(vertex)$ will receive contributions from the 'standard' diagram in Figure 1a as well as from the one in Figure 2a which involves superpartners in the loop⁴. We also define the slepton asymmetry associated to N decays as

$$\tilde{\epsilon}_{\tilde{L}}^{N} \equiv \frac{\Gamma_{N\tilde{L}} - \Gamma_{N\tilde{L}^{*}}}{\Gamma_{N\tilde{L}} + \Gamma_{N\tilde{L}^{*}}},\tag{17}$$

which arises from diagrams similar to those in Figures 1 and 2 but with $\tilde{L}\tilde{h}$ in the final state. We similarly define the asymmetry parameters associated to sneutrino decays as

$$\tilde{\epsilon}_{\ell}^{\tilde{N}^*} \equiv \frac{\Gamma_{\tilde{N}^*\ell} - \Gamma_{\tilde{N}\bar{\ell}}}{\Gamma_{\tilde{N}^*\ell} + \Gamma_{\tilde{N}\bar{\ell}}} \quad , \quad \tilde{\epsilon}_{\tilde{L}}^{\tilde{N}} \equiv \frac{\Gamma_{\tilde{N}\tilde{L}} - \Gamma_{\tilde{N}^*\tilde{L}^*}}{\Gamma_{\tilde{N}\tilde{L}} + \Gamma_{\tilde{N}^*\tilde{L}^*}}. \tag{18}$$

After a direct computation of these quantities we obtain

$$\tilde{\epsilon}_{\ell}^{N_i}(vertex) = -\frac{1}{8\pi} \sum_{k} g(y_k) \mathcal{I}_{ki} , \qquad (19)$$

⁴For later convenience, we define $\tilde{\epsilon}_{\ell}^{N}$ as in Eq. (2), rather than normalizing it to the total N decay rate which also includes the slepton final states.

where $g(x) = \sqrt{x} \ln[(1+x)/x]$, while

$$\tilde{\epsilon}_{\ell}^{N}(wave) = 2\epsilon_{\ell}^{N}(wave). \tag{20}$$

For the remaining channels we get similar results

$$\tilde{\epsilon}_{\tilde{L}}^{N}(vertex) = \tilde{\epsilon}_{\ell}^{\tilde{N}^{*}}(vertex) = \tilde{\epsilon}_{\tilde{L}}^{\tilde{N}}(vertex) = \tilde{\epsilon}_{\ell}^{N}(vertex), \tag{21}$$

and also the wave function contributions are equal

$$\tilde{\epsilon}_{\tilde{L}}^{N}(wave) = \tilde{\epsilon}_{\ell}^{\tilde{N}^{*}}(wave) = \tilde{\epsilon}_{\tilde{L}}^{\tilde{N}}(wave) = \tilde{\epsilon}_{\ell}^{N}(wave). \tag{22}$$

The vertex pieces were computed previously in ref. [6] and our results are proportional. The sneutrino decay asymmetry was also considered in ref. [7] in the limit in which all heavy (s)neutrinos are degenerate and, specialised to that case, our results are in agreement (except for the overall sign). The wave function contributions were not considered previously and, as in the non–supersymmetric case, they are non–negligible. Let us also note that there are additional one loop diagrams involving the four–scalar couplings appearing in the F–terms of the scalar potential. Although helpful to cure the divergences in the loops, they do not contribute to the asymmetry generation.

For the study of the decays of thermal populations of heavy neutrinos and sneutrinos, it is convenient to introduce the total asymmetry defined as

$$\tilde{\epsilon}_i \equiv \tilde{\epsilon}_{\ell}^{N_i} + \tilde{\epsilon}_{\tilde{L}}^{N_i} + \tilde{\epsilon}_{\ell}^{\tilde{N}_i^*} + \tilde{\epsilon}_{\tilde{L}}^{\tilde{N}_i} = 4\tilde{\epsilon}_{\ell}^N, \tag{23}$$

resulting in

$$\tilde{\epsilon}_i = -\frac{1}{2\pi} \sum_{k \neq i} \left[g(y_k) + \frac{2\sqrt{y_k}}{y_k - 1} \right] \mathcal{I}_{ki}. \tag{24}$$

In particular, in the case of hierarchical masses $(y_k \gg 1)$, we find that again the wave contribution becomes twice as large as the vertex one, giving

$$\tilde{\epsilon}_1 \simeq -\frac{3}{2\pi} \sum_{k \neq 1} \frac{1}{\sqrt{y_k}} \, \mathcal{I}_{k1}. \tag{25}$$

With the definition in Eq. (23), the resulting lepton asymmetry is⁵

$$\frac{n_L}{s} = \frac{\tilde{\epsilon}_1}{s} \frac{T^3}{\pi^2} \simeq 1 \times 10^{-3} \tilde{\epsilon}_1, \tag{26}$$

where we have used that the effective number of degrees of freedom in the supersymmetric scenario is approximately doubled, i.e. $g_*^{SUSY} \simeq 2 \ g_*^{SM}$. Hence, to account for the baryon asymmetry of the Universe we need now

$$\sum_{k \neq 1} \frac{M_1}{M_k} \, \mathcal{I}_{k1} \simeq 0.7 \times 10^{-6}. \tag{27}$$

Comparing this with Eq. (13), we see that the required parameters are similar in the supersymmetric and standard scenarios.

To summarize, we have computed all the contributions to the CP violating asymmetries arising at one-loop in the decays of heavy (s)neutrinos, both in the standard non-supersymmetric Fukugita-Yanagida scenario and in its supersymmetric version. We have discussed the different results present in the literature and showed that the contribution from wave function mixing is relevant in the computation of the CP violating asymmetries. The baryon number generated in both scenarios was also obtained.

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⁵Had we normalised $\epsilon_{\ell,\tilde{L}}^N$ to the total N decay rate, their contribution to n_L would have to be multiplied by $g_N = 2$ as in Eq. (9).

References

- [1] M. Fukugita and T. Yanagida, Phys. Lett. **B174** (1986) 45.
- [2] N. S. Manton, Phys. Rev. **D28** (1983) 2212. V. A. Kuzmin, V. A. Rubakov and M.
 E. Shaposhnikov, Nucl. Phys. **B155** (1985) 36.
- [3] M. A. Luty, Phys. Rev. **D45** (1992) 455.
- [4] M. Plümacher, hep-ph/9604229.
- [5] G. Gelmini and E. Roulet, Rep. Prog. Phys. **58** (1995) 1207.
- [6] B. A. Campbell, S. Davidson and K. A. Olive, Nucl. Phys. **B399** (1993) 111.
- [7] H. Murayama et al., Phys. Rev. Lett. **70** (1993) 1912.
- [8] A. Yu. Ignatiev, V. A. Kuzmin and M. E. Shaposhnikov, JETP Lett. 30 (1979) 688.
- [9] F. J. Botella and J. Roldan, Phys. Rev. **D44** (1991) 966.
- [10] J. Liu and G. Segrè, Phys. Rev. **D48** (1993) 4609.
- [11] M. Fukugita and T. Yanagida, *Physics of Neutrinos*, in: *Physics and Astrophysics of Neutrinos*, Eds. M. Fukugita and A. Suzuki, (Springer-Verlag, Tokyo, 1994).
- [12] J. Liu and G. Segrè, Phys. Rev. **D49** (1994) 1342.
- [13] J. A. Harvey and M. S. Turner, Phys. Rev. **D42** (1990) 3344.
- [14] C. J. Copi, D. N. Schramm and M. S. Turner, Science **267** (1995) 192.