Flavoured CP asymmetries for type II seesaw leptogenesis

R. González Felipe
Instituto Superior de Engenharia de Lisboa - ISEL, Rua Conselheiro Emídio Navarro, 1959-007 Lisboa, Portugal

Departamento de Física and Centro de Física Teórica de Partículas (CFTP), Instituto Superior Técnico, Universidade de Lisboa, Avenida Rovisco Pais, 1049-001 Lisboa, Portugal
ricardo.felipe@ist.utl.pt

F. R. Joaquim
Departamento de Física and Centro de Física Teórica de Partículas (CFTP), Instituto Superior Técnico, Universidade de Lisboa, Avenida Rovisco Pais, 1049-001 Lisboa, Portugal
filipe.joaquim@ist.utl.pt

H. Serôdio
Departament de Física Teòrica and IFIC, Universitat de València - CSIC, E-46100, Burjassot, Spain
hugo.serodio@ific.uv.es

A novel contribution to the leptonic CP asymmetries in type II seesaw leptogenesis scenarios is obtained for the cases in which flavour effects are relevant for the dynamics of leptogenesis. In the so-called flavoured leptogenesis regime, the interference between the tree-level amplitude of the scalar triplet decaying into two leptons and the one-loop wave-function correction with leptons in the loop, leads to a new nonvanishing CP asymmetry contribution. The latter conserves total lepton number but violates lepton flavour. Cases in which this novel contribution may be dominant in the generation of the baryon asymmetry are briefly discussed.

Keywords: Leptogenesis, neutrino physics, seesaw mechanism

1. Introduction

Leptogenesis[1] is perhaps the most appealing mechanism to explain the matter-antimatter asymmetry observed in the Universe. One of its remarkable features is the possibility of establishing a bridge between neutrino physics at high and low energies, through the well-known seesaw mechanism for neutrino mass generation. Several scenarios are conceivable in this context. Namely, the canonical ones are the type I[2,3,4,5,6], type II[7,8,9,10,11] and type III[12] seesaws, in which neutrino masses are mediated by the three-level exchange of heavy singlet fermions, $SU(2)$-triplet scalars and $SU(2)$-triplet fermions, respectively. Particularly economical is the type II seesaw scenario with one triplet, where the flavour pattern of the Yukawa couplings between the scalar triplets and the Standard Model (SM) doublets uniquely determines the flavour structure of the low-energy effective neutrino mass matrix. This feature is particularly interesting when the type II seesaw is embedded in a beyond-the-SM framework, where those Yukawa couplings trigger new sources of lepton flavour violation that may be relevant for processes like radiative charged-lepton decays[13]. This is indeed what happens in the supersymmetric type II seesaw where model-independent predictions can be made for the rates of lepton flavour violating decays in terms of the low-energy neutrino parameters[14,15,16,17,18,19].

To successfully implement leptogenesis in a minimal type II seesaw framework (without introducing heavy singlet fermions), at least two scalar triplets are needed[20].[a] The complex Yukawa couplings of the Higgs triplets to leptons, as well as their complex couplings to the standard model Higgs doublet, provide the necessary sources of CP violation for leptogenesis. In particular, the CP asymmetry in the decay of the scalar triplet into two leptons arises from the interference of the corresponding tree-level and one-loop amplitudes. A nonvanishing lepton asymmetry is then generated via the out-of-equilibrium decays of the triplet scalars in the early Universe, which is afterwards partially converted into a baryon asymmetry by nonperturbative sphaleron processes[22].

Departure from thermal equilibrium crucially depends on the expansion rate of the Universe. Since at very high temperatures ($T \gtrsim 10^{12}$ GeV) all charged lepton flavours are out of thermal equilibrium, their states are indistinguishable. Interactions involving the τ and μ Yukawa couplings enter in equilibrium at $T \lesssim 10^{12}$ GeV and $T \lesssim 10^{9}$ GeV, respectively. The corresponding lepton doublets are distinguishable mass eigenstates below these temperature scales and, therefore, their flavour effects should be properly taken into account in the leptogenesis dynamics. Such effects turn out to be relevant in type I seesaw leptogenesis scenarios[23,24,25,26]. In particular, in the flavoured regime the washout processes can be less significant than in the unflavoured one, and the low-energy leptonic phases affect directly the final asymmetry so that it is possible to have successful leptogenesis just from low-energy leptonic CP violation[27,28]. Also, the upper bound on each individual flavoured asymmetry is not suppressed when the absolute neutrino mass scale increases.

So far, flavour effects in type II seesaw leptogenesis have been only partially addressed[29,13,30]. The purpose of this work is to study the importance of these effects on the leptonic CP asymmetries generated in minimal type II seesaw scenarios where leptogenesis is implemented through the out-of-equilibrium decays of scalar triplets into two leptons. The required CP asymmetries in those decays are guaranteed by the interference of the tree-level and one-loop decay amplitudes, in the presence of complex couplings of the triplets with the SM Higgs and the leptons. It turns out that there is a novel contribution to the flavoured leptogenesis

[a]In the presence of only one scalar triplet, the CP asymmetry induced by the triplet decays is generated beyond the one-loop level and is therefore highly suppressed[21].

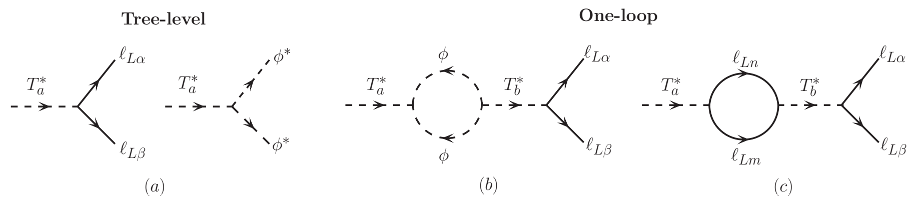

Fig. 1. Tree-level and one-loop Feynman diagrams contributing to the CP asymmetry in a type II seesaw framework.

asymmetries coming from the wave-function renormalization correction. We briefly discuss some cases in which this contribution may be dominant.

2. Type II seesaw leptogenesis

In the type II seesaw mechanism, neutrino masses are generated through the tree-level exchange of hypercharge $Y = 1$ scalar $\mathrm{SU}(2)_L$ triplets. As already pointed out, a single triplet is enough to explain the low-energy neutrino spectrum. However, in this case it is not possible to generate a leptonic CP asymmetry for leptogenesis, since all the interference terms vanish at one loop. Thus, we need to add at least another triplet to allow for the generation of a non-zero CP asymmetry[20]. Here, we will consider an extension of the SM in which n_L scalar triplets T_a are added. Following the usual SU(2) representation, each T_a can be written in terms of the corresponding charge eigenstates T_a^0, T_a^+ and T_a^{++} as

$$T_a = \begin{pmatrix} T_a^0 & -\frac{T_a^+}{\sqrt{2}} \\ -\frac{T_a^+}{\sqrt{2}} & T_a^{++} \end{pmatrix} . \tag{1}$$

The relevant Lagrangian is given by

$$\mathcal{L}_{II} = \mathcal{L}_{SM} + \mathrm{Tr}\left[(D_\mu T_a)^\dagger (D^\mu T_a)\right] - \left(\mathbf{Y}^{T_a}_{\alpha\beta}\overline{\ell_{L\alpha}} T_a^\dagger \ell^c_{L\beta} + \mathrm{h.c.}\right) - V_T \,, \tag{2}$$

where $\mathcal{L}_{SM}$ contains the SM terms and V_T accounts for the scalar potential terms involving the triplets,

$$\begin{aligned} V_T =& M_a^2\,\mathrm{Tr}\left(T_a^\dagger T_a\right) + \mu_a \tilde{\phi}^T T_a \tilde{\phi} + g_{ab}\phi^\dagger T_a^\dagger T_b \phi + h_{ab}\phi^\dagger\phi\,\mathrm{Tr}\left(T_a^\dagger T_b\right) \\ &+ \lambda'_{abcd}\mathrm{Tr}\left(T_a^\dagger T_b T_c^\dagger T_d\right) + \lambda_{abcd}\mathrm{Tr}\left(T_a^\dagger T_b\right)\mathrm{Tr}\left(T_c^\dagger T_d\right) + \mathrm{h.c.}\,. \end{aligned} \tag{3}$$

In the above equations, D_μ stands for the usual covariant derivative, $\phi = (\phi^+, \phi^0)^T$ is the SM Higgs doublet ($\tilde{\phi} = i\sigma_2\phi^*$), and $\ell = (\nu_L, l_L)^T$ is a SM lepton doublet.

Each T_a has two decay modes: $T_a^\dagger \to \ell_\alpha \ell_\beta$ and $T_a^\dagger \to \phi^*\phi^*$ (see Fig. 1a). For the first channel, neglecting the masses of the final states, the tree-level decay rates are

$$\Gamma\left(T_a^\dagger \to \ell_\alpha\ell_\beta\right) = \frac{M_a}{8\pi}\left|\mathbf{Y}^{T_a}_{\alpha\beta}\right|^2 c_{\alpha\beta}\,, \quad \text{with} \quad c_{\alpha\beta} = \begin{cases} 2-\delta_{\alpha\beta} & \text{for } T_a^0\,,\, T_a^{++} \\ 1 & \text{for } T_a^+ \end{cases} . \tag{4}$$

Summing over the final flavour states we get

$$\Gamma\left(T_a^{--} \to l^- l^-\right) = \Gamma\left(T_a^{0*} \to \nu\nu\right) = \frac{M_a}{8\pi} \sum_{\alpha,\beta\geq\alpha} \left|\mathbf{Y}^{T_a}_{\alpha\beta}\right|^2 c_{\alpha\beta} \,, \qquad \Gamma\left(T_a^{-} \to l^- \nu\right) = \frac{M_a}{8\pi} \sum_{\alpha,\beta} \left|\mathbf{Y}^{T_a}_{\alpha\beta}\right|^2 . \tag{5}$$

Note that in the case of T_a^{--} and T_a^0 the sum is ordered, while for T_a^- is not. Therefore, at tree level $\Gamma\left(T_a^{--} \to l^- l^-\right) = \Gamma\left(T_a^{0*} \to \nu\nu\right) = \Gamma\left(T_a^{-} \to l^- \nu\right)$. This is the result one would expect since T_a^{0*}, T_a^- and T_a^{--} belong to the same $SU(2)_L$ multiplet, which is not yet broken at the leptogenesis scale. For the decay channel into the Higgs scalars, we get

$$\Gamma\left(T_a^{--} \to \phi^-\phi^-\right) = \Gamma\left(T_a^{0*} \to \phi^{0*}\phi^{0*}\right) = \Gamma\left(T_a^{-} \to \phi^-\phi^{0*}\right) = \frac{|\mu_a|^2}{8\pi M_a} . \tag{6}$$

The total decay rate is then given by

$$\Gamma_{T_a} \equiv \Gamma\left(T_a^\dagger \to \ell\ell, \phi^*\phi^*\right) = \frac{M_a}{8\pi}\left[\mathrm{Tr}\left(\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_a}\right) + |\lambda_a|^2\right] , \tag{7}$$

with $\lambda_a = \mu_a/M_a$.

The leptonic CP asymmetries relevant for leptogenesis stem from the interference between the tree-level and one-loop T_a decay amplitudes. In our framework, there is no vertex correction contributing to the leptogenesis CP asymmetry at the one-loop level. The only diagrams contributing to the CP asymmetries are those coming from wave function renormalization, shown in Figs. 1b and 1c. Notice that the former is both total lepton number and lepton flavour violating, while the second one is only lepton flavour violating. The interference between the one-loop diagram (1b) and the corresponding tree-level one leads to the CP asymmetry[13]

$$\epsilon_a^{\alpha\beta}(\text{wave 1}) \simeq \frac{c_{\alpha\beta}}{2\pi} \frac{\sum_{b\neq a} g(z_b)\mathrm{Im}\left(\lambda_a^*\lambda_b \mathbf{Y}^{T_b}_{\alpha\beta}\mathbf{Y}^{T_a *}_{\alpha\beta}\right)}{\mathrm{Tr}\left(\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_a}\right) + |\lambda_a|^2} , \tag{8}$$

where $z_b \equiv M_b^2/M_a^2$ and

$$g(z_b) = \frac{\sqrt{z_b}\,(1-z_b)}{(z_b-1)^2 + (\Gamma_{T_b}/M_a)^2} . \tag{9}$$

We recall that in the type II seesaw the Yukawa couplings are directly related to the effective light neutrino mass matrix by

$$\mathbf{m}_\nu = \sum_a \mathbf{m}_\nu^a \,, \quad \mathbf{m}_\nu^a = 2\langle T_a^0\rangle^* \mathbf{Y}^{T_a} = 2\frac{\lambda_a v^2}{M_a}\mathbf{Y}^{T_a} , \tag{10}$$

with $v = \langle\phi^0\rangle = 174$ GeV. One can then rewrite the leptonic CP asymmetries in terms of these quantities. Indeed, using the branching ratio relations

$$\mathcal{B}_a^\ell\,\Gamma_{T_a} \equiv \frac{M_a}{8\pi}\mathrm{Tr}\left(\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_a}\right) , \quad \mathcal{B}_a^\phi\,\Gamma_{T_a} \equiv \frac{M_a}{8\pi}|\lambda_a|^2 , \tag{11}$$

and

$$\begin{aligned}\sqrt{\mathcal{B}^\ell_a\mathcal{B}^\phi_a}\,\Gamma_{T_a} &= \frac{M_a^2}{16\pi v^2}\left[\mathrm{Tr}(\mathbf{m}_\nu^{a\dagger}\mathbf{m}_\nu^a)\right]^{1/2}, \\ \mathrm{Im}\left(\lambda_a^*\lambda_b\mathbf{Y}^{T_b}_{\alpha\beta}\mathbf{Y}^{T_a*}_{\alpha\beta}\right) &= \frac{M_aM_b}{4v^4}\mathrm{Im}\left[\left(\mathbf{m}_\nu^b\right)_{\alpha\beta}\left(\mathbf{m}_\nu^{a*}\right)_{\alpha\beta}\right],\end{aligned} \tag{12}$$

Eq. (8) can be expressed in the form

$$\epsilon_a^{\alpha\beta}(\text{wave }1) \simeq \frac{c_{\alpha\beta}}{4\pi}\frac{M_a\sqrt{\mathcal{B}^\ell_a\mathcal{B}^\phi_a}}{v^2}\frac{\sum_{b\neq a} g(z_b)\mathrm{Im}\left[\left(\mathbf{m}_\nu^b\right)_{\alpha\beta}\left(\mathbf{m}_\nu^{a*}\right)_{\alpha\beta}\right]}{\left[\mathrm{Tr}(\mathbf{m}_\nu^{a\dagger}\mathbf{m}_\nu^a)\right]^{1/2}}\,. \tag{13}$$

For the second contribution to the wave function renormalization, coming from the interference between the one-loop diagram (1c) and the corresponding tree-level one, we obtain

$$\epsilon_a^{\alpha\beta}(\text{wave }2) \simeq \frac{c_{\alpha\beta}}{2\pi}\frac{\sum_{b\neq a} z_b^{-1/2} g(z_b)\,\mathrm{Im}\left[\mathrm{Tr}\left(\mathbf{Y}^{T_b\dagger}\mathbf{Y}^{T_a}\right)\mathbf{Y}^{T_b}_{\alpha\beta}\mathbf{Y}^{T_a*}_{\alpha\beta}\right]}{\mathrm{Tr}\left(\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_a}\right)+|\lambda_a|^2}\,. \tag{14}$$

Rewriting the imaginary part as

$$\mathrm{Im}\left[\mathrm{Tr}\left(\mathbf{Y}^{T_b\dagger}\mathbf{Y}^{T_a}\right)\mathbf{Y}^{T_b}_{\alpha\beta}\mathbf{Y}^{T_a*}_{\alpha\beta}\right] = \frac{M_a^2M_b^2}{16v^8|\lambda_a|^2|\lambda_b|^2}\mathrm{Im}\left[\mathrm{Tr}\left(\mathbf{m}_\nu^{b\dagger}\mathbf{m}_\nu^a\right)(\mathbf{m}_\nu^b)_{\alpha\beta}(\mathbf{m}_\nu^{a*})_{\alpha\beta}\right], \tag{15}$$

we get for the CP asymmetry

$$\begin{aligned}\epsilon_a^{\alpha\beta}(\text{wave }2) \simeq &\frac{c_{\alpha\beta}}{16\pi}\frac{M_a\sqrt{\mathcal{B}^\ell_a\mathcal{B}^\phi_a}}{v^6|\lambda_a|^2\left[\mathrm{Tr}\left(\mathbf{m}_\nu^{a\dagger}\mathbf{m}_\nu^a\right)\right]^{1/2}}\times \\ &\sum_{b\neq a}\frac{M_b^2\,g(z_b)\,\mathrm{Im}\left[\mathrm{Tr}\left(\mathbf{m}_\nu^{b\dagger}\mathbf{m}_\nu^a\right)(\mathbf{m}_\nu^b)_{\alpha\beta}(\mathbf{m}_\nu^{a*})_{\alpha\beta}\right]}{z_b^{1/2}|\lambda_b|^2},\end{aligned} \tag{16}$$

in terms of the various contributions to the neutrino mass matrix.

The above novel contribution to the CP asymmetries is only relevant within the flavoured leptogenesis regime. Indeed, summing over the final flavours we get

$$\epsilon_a(\text{wave }2) \propto \mathrm{Im}\left[\mathrm{Tr}\left(\mathbf{Y}^{T_b\dagger}\mathbf{Y}^{T_a}\right)\mathrm{Tr}\left(\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_b}\right)\right] = 0\,. \tag{17}$$

Therefore, the only contribution surviving in the unflavoured regime comes from the Higgs loop. The final unflavoured asymmetry is given by

$$\begin{aligned}\epsilon_a &\simeq \frac{1}{2\pi}\frac{\sum_{b\neq a} g(z_b)\mathrm{Im}\left[\lambda_a^*\lambda_b\mathrm{Tr}\left(\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_b}\right)\right]}{\mathrm{Tr}\left[\mathbf{Y}^{T_a\dagger}\mathbf{Y}^{T_a}\right]+|\lambda_a|^2} \\ &= \frac{1}{4\pi}\frac{M_a\sqrt{\mathcal{B}^\ell_a\mathcal{B}^\phi_a}}{v^2}\frac{\sum_{b\neq a} g(z_b)\mathrm{Im}\left[\mathrm{Tr}\left(\mathbf{m}_\nu^{a\dagger}\mathbf{m}_\nu^b\right)\right]}{\left[\mathrm{Tr}\left(\mathbf{m}_\nu^{a\dagger}\mathbf{m}_\nu^a\right)\right]^{1/2}}\,.\end{aligned} \tag{18}$$

We stress that in the regime where the final-state flavour discrimination is important, the new contribution $\epsilon_a^{\alpha\beta}(\text{wave}\,2)$ given in Eq. (14) can, in principle, dominate over $\epsilon_a^{\alpha\beta}(\text{wave}\,1)$. In a minimal scenario with two scalar triplets $T_{1,2}$, with $\mathbf{Y}^{T_1} \simeq \mathbf{Y}^{T_2} \sim y$ and $M_1 \ll M_2$, the condition for $\epsilon_a^{\alpha\beta}(\text{wave}\,2) \gg \epsilon_a^{\alpha\beta}(\text{wave}\,1)$ would roughly be

$$y^2 \gg \frac{M_2}{M_1}\lambda_1\lambda_2\,. \tag{19}$$

This can be easily achieved if the triplets couple strongly to leptons but very weakly to the SM Higgs doublet. Ultimately, if one of the triplets couples mainly to the leptons (thus, not giving any contribution to neutrino masses) then $\epsilon_a^{\alpha\beta}(\text{wave}\,1) \simeq 0$ and $\epsilon_a^{\alpha\beta}(\text{wave}\,2)$ is the only contribution for the CP asymmetries. To illustrate this, let us consider a simple example with two triplets $T_{1,2}$ of masses $M_{1,2}$. We assume

$$\mathbf{Y}^{T_1} \simeq \frac{M_1\mathbf{m}_\nu}{2\lambda_1 v^2}\,,\ \mathbf{Y}^{T_2} = \mathbf{K}\mathbf{Y}^{T_1}\mathbf{K}\,, \tag{20}$$

with $\mathbf{K} = \mathrm{diag}(e^{i\pi/2},1,1)$. The above approximation for $\mathbf{Y}^{T_1}$ is valid provided that $\lambda_1 M_2 \gg \lambda_2 M_1$. The effective neutrino mass matrix is constructed from $\mathbf{m}_\nu = \mathbf{U}^*\mathrm{diag}(m_1,m_2,m_3)\mathbf{U}^\dagger$, where $\mathbf{U}$ is the PMNS lepton mixing matrix in the standard PDG parametrization and m_i are the neutrino masses. We take the best-fit values from the latest global analysis of all neutrino oscillation data[31,32] and consider a hierarchical neutrino mass spectrum with $m_1 \simeq 0$. We also assume maximal Dirac-type CP violation, i.e. the phase $\delta = \pi/2$, and neglect any Majorana-type CP violation. As for the high-energy parameters, we choose $\lambda_1 = 10\lambda_2 = 5\times 10^{-6}$ and take $M_2 = 10M_1 = 10^{10}\,\mathrm{GeV}$ (to ensure that leptogenesis takes place within the flavoured regime). In this case, $\epsilon_1^{\alpha\beta}(\text{wave}\,1) \simeq 0$ (for all α, $\beta = e,\mu,\tau$), while

$$\epsilon_1^{\alpha\beta}(\text{wave}\,2) \simeq \begin{pmatrix} -0.02 & -5.62 & -6.92 \\ -5.62 & 5.84 & 4.59 \\ -6.92 & 4.59 & 10.08 \end{pmatrix} \times 10^{-7}\,, \tag{21}$$

which is sufficiently large to give a sizeable contribution to the baryon asymmetry.

3. Flavoured Boltzmann equations

The final asymmetry crucially depends on the efficiency of leptogenesis, which is dictated by the solution of the relevant Boltzmann equations in the flavoured regime. Before discussing these equations in detail, it is worth commenting on some general features that are present in the unflavoured regime. It has been shown[33] that, for unflavoured leptogenesis, the efficiency is maximal when either $\mathcal{B}_a^\ell \ll \mathcal{B}_a^\phi$ or $\mathcal{B}_a^\ell \gg \mathcal{B}_a^\phi$ (for a recent analysis see Ref. [34]). This can be easily understood if one recalls that, in the type II seesaw, lepton number is violated only if both decay channels of the scalar triplet (i.e. to two leptons and to two Higgs scalars) are present. Thus, even when the total decay rate Γ_{T_a} and the gauge scattering rates are much larger than the Hubble rate, if either the decay rate to leptons or to Higgs doublets is out of thermal equilibrium (i.e. $\mathcal{B}_a^\ell \ll 1$ or $\mathcal{B}_a^\phi \ll 1$), lepton number is not erased by the corresponding inverse decays, and there is no suppression of the leptogenesis efficiency. As we shall see below, some of these features remain valid in the flavoured regime.

We restrict our analysis to a minimal scenario of two scalar triplets $T_{1,2}$, with T_1 lighter that T_2, so that the effects of lepton-number violating processes due to T_2 can be safely neglected. In general, the lepton asymmetry produced in the T_2 decays will be washed out by the interactions of T_1. We denote by n_x the number density of the particle x, and define its comoving number density $Y_x = n_x/s$, where $s = (2\pi^2/45)\, g_* T^3$ is the total entropy density at leptogenesis temperatures ($g_* = 106.75$). We also define the comoving asymmetries $\Delta_x = Y_x - Y_{\bar{x}}$ and denote by $\Sigma_T = Y_{T_1} + Y_{\bar{T}_1}$ the total triplet density.

The relevant Boltzmann equations, which describe the evolution of the asymmetries Δ_T, Δ_ϕ and Δ_{ℓ_α} ($\alpha = e, \mu, \tau$) as functions of $z = M_1/T$, read as [b]

$$szH(z)\,\frac{d\Sigma_T}{dz} = -\left(\frac{\Sigma_T}{\Sigma_T^{eq}} - 1\right)\gamma_D - 2\left(\frac{\Sigma_T^2}{\Sigma_T^{eq^2}} - 1\right)\gamma_A\,, \tag{22a}$$

$$szH(z)\,\frac{d\Delta_T}{dz} = -\,\gamma_D\left(\frac{\Delta_T}{\Sigma_T^{eq}} + \sum_{\alpha,\beta}\mathcal{B}_1^{\alpha\beta}\frac{\Delta_{\ell_\alpha}}{Y_\ell^{eq}} - \mathcal{B}_1^\phi\frac{\Delta_\phi}{Y_\phi^{eq}}\right)\,, \tag{22b}$$

$$szH(z)\,\frac{d\Delta_\phi}{dz} = \sum_{\alpha,\beta} X_{\alpha\beta} - 2\mathcal{B}_1^\phi\gamma_D\left(\frac{\Delta_\phi}{Y_\phi^{eq}} - \frac{\Delta_T}{\Sigma_T^{eq}}\right)\,, \tag{22c}$$

$$szH(z)\,\frac{d\Delta_{\ell_\alpha}}{dz} = \sum_{\beta}\left[X_{\alpha\beta} - 2\mathcal{B}_1^{\alpha\beta}\gamma_D\left(\frac{\Delta_T}{\Sigma_T^{eq}} + \frac{\Delta_{\ell_\alpha} + \Delta_{\ell_\beta}}{2Y_\ell^{eq}}\right)\right]\,, \tag{22d}$$

where

$$\begin{aligned} X_{\alpha\beta} = &\left(\frac{\Sigma_T}{\Sigma_T^{eq}} - 1\right)\gamma_D\epsilon_1^{\alpha\beta} + 2\gamma_D\left(\mathcal{B}_1^\ell\,\epsilon_1^{\alpha\beta} - \mathcal{B}_1^{\alpha\beta}\,\epsilon_1\right) \\ &- \left(2\frac{\Delta_\phi}{Y_\phi^{eq}} + \frac{\Delta_{\ell_\alpha} + \Delta_{\ell_\beta}}{Y_\ell^{eq}}\right)\left(2\gamma'^{\bar{\phi}\bar{\phi}}_{\ell_\alpha\ell_\beta} + \gamma^{\bar{\ell}_\beta\bar{\phi}}_{\ell_\alpha\phi}\right). \end{aligned} \tag{23}$$

In the above equations,

$$H(z) = \frac{H_0(M_1)}{z^2}\,, \quad H_0(T) = \sqrt{\frac{4\pi^3}{45}g_*}\,\frac{T^2}{m_P} \tag{24}$$

is the Hubble constant at temperature T, $m_P = 1.22 \times 10^{19}$ GeV is the Planck mass, and the suffix eq denotes equilibrium values. We have

$$Y_T^{eq} = \frac{45\,g_T}{4\pi^4 g_*}z^2K_2(z)\,, \quad Y_\ell^{eq} = \frac{3}{4}\,\frac{45\,\zeta(3)g_\ell}{2\pi^4 g_*}\,, \quad Y_\phi^{eq} = \frac{45\,\zeta(3)g_\phi}{2\pi^4 g_*}\,, \tag{25}$$

[b] Since scalar triplets are not self-conjugated states, a triplet-antitriplet asymmetry Δ_T is generated and a Boltzmann equation for this asymmetry must be included.

where g_x are the degrees of freedom of the particle ($g_T = 1$ for each triplet component, $g_\ell = 2$ and $g_\phi = 2$), $\zeta(3) \simeq 1.202$ and $K_i(z)$ are the modified Bessel functions. Finally, the relevant interaction densities contributing to leptogenesis are the decays and inverse decays, given by the standard expression

$$\gamma_D = s\,\Gamma_{T_1}\,\Sigma_T^{eq}\,\frac{K_1(z)}{K_2(z)}\,, \tag{26}$$

gauge scattering processes of the triplets, approximately given by the s-wave contribution

$$\gamma_A = \frac{M_1 T^3 e^{-2M_1/T}}{64\pi^4}(9g^4 + 12g^2 g_Y^2 + 3g_Y^4)\left(1 + \frac{3T}{4M_1}\right), \tag{27}$$

and $\Delta L = 2$ scattering processes due to $\ell_\alpha \ell_\beta \leftrightarrow \bar{\phi}\bar{\phi}$ and $\ell_\alpha \phi \leftrightarrow \bar{\ell}_\beta \bar{\phi}$ generated in s-channel and t-channel, respectively. For these processes, the reaction densities are obtained as

$$\gamma = \frac{T}{64\pi^4}\int_0^\infty ds\, s^{1/2} K_1(\sqrt{s}/T)\,\hat{\sigma}(s), \tag{28}$$

where the reduced cross sections are

$$\hat{\sigma}_{\ell_\alpha \ell_\beta}^{\bar{\phi}\bar{\phi}} = \frac{3x|\lambda_1|^2|\mathbf{Y}_{\alpha\beta}^{T_1}|^2}{2\pi}\left[\frac{1}{(1-x)^2 + (\Gamma_{T_1}/M_1)^2}\right], \tag{29a}$$

$$\hat{\sigma}_{\ell_\alpha \phi}^{\bar{\ell}_\beta \bar{\phi}} = \frac{6|\lambda_1|^2|\mathbf{Y}_{\alpha\beta}^{T_1}|^2}{\pi}\left[-\frac{1}{1+x} + \frac{\ln(1+x)}{x}\right], \tag{29b}$$

and $x = s/M_1^2$. We recall that in the Boltzmann equations the term due to on-shell triplet exchange must be subtracted from $\gamma_{\ell_\alpha \ell_\beta}^{\bar{\phi}\bar{\phi}}$. This procedure leads to the subtracted reaction density

$$\gamma'_{\alpha\beta} = \gamma_{\alpha\beta} - \frac{1}{2}\,\mathcal{B}_1^{\alpha\beta}\mathcal{B}_1^{\phi}\gamma_D, \tag{30}$$

where $\mathcal{B}_1^{\alpha\beta} = \Gamma(T_1^\dagger \to \ell_\alpha \ell_\beta)/\Gamma_{T_1}$.

The Boltzmann equations given in Eqs. (22) have been obtained following the standard general procedure, as described e.g. in Ref. [35]. These equations contain the relevant processes contributing to triplet leptogenesis and share the same structure of those in the unflavoured regime[33]. The main difference resides in the fact that quantities that depend on the lepton flavours are now treated independently.

We integrate the system of equations (22) with the following initial conditions: $\Sigma_T(z \ll 1) = \Sigma_T^{eq}(z \ll 1)$, $\Delta_T(z \ll 1) = 0$, $\Delta_\phi(z \ll 1) = 0$, and $\Delta_{\ell_\alpha}(z \ll 1) = 0$. To demonstrate how a large leptogenesis efficiency may arise due to the novel CP asymmetry contribution given in Eq. (14), we consider the example case presented at the end of the previous section. In Fig. 2 we plot the evolution of the asymmetries Δ_T, Δ_ϕ, and Δ_{ℓ_α} in the three-flavoured regime, with the CP asymmetries given by the matrix (21), $M_2 = 10M_1 = 10^{10}$ GeV and $\lambda_1 = 10\lambda_2 = 5 \times 10^{-6}$. As can be seen from the figure, large leptonic asymmetries develop,

$$\Delta_{\ell e} \simeq -2.23 \times 10^{-7}, \quad \Delta_{\ell\mu} \simeq \Delta_{\ell\tau} \simeq 1.08 \times 10^{-7}, \tag{31}$$

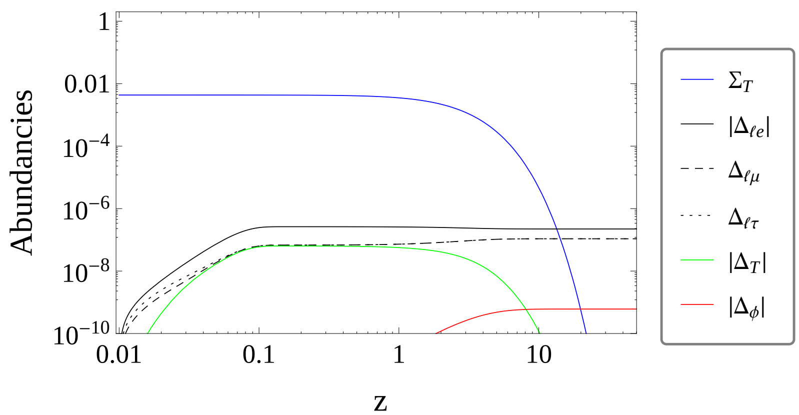

Fig. 2. Evolution of asymmetries for the example case in which the flavoured CP asymmetries are given by Eq. (21), $M_2 = 10M_1 = 10^{10}$ GeV and $\lambda_1 = 10\lambda_2 = 5 \times 10^{-6}$. The final baryon asymmetry is $\Delta_B \simeq 8 \times 10^{-8}$.

and a sizable baryon asymmetry can be generated. Indeed, at temperatures below 10^9 GeV, the final baryon asymmetry Δ_B can be estimated as

$$\Delta_B = 3\, c_{\rm sph} \sum_{\alpha,\beta} A^{-1}_{\alpha\beta}\, \Delta_{\ell_\beta}, \tag{32}$$

where $c_{\rm sph} = 28/79$ is the sphaleron conversion factor, and the matrix A is given by[25]

$$A = \begin{pmatrix} -151/179 & 20/179 & 20/179 \\ 25/358 & -344/537 & 14/537 \\ 25/358 & 14/537 & -344/537 \end{pmatrix}. \tag{33}$$

For the example under consideration, we obtain $\Delta_B \simeq 8 \times 10^{-8}$.

It is worthwhile to comment on the efficiency of leptogenesis in this case. Defining the flavoured efficiency factors,

$$\eta_\alpha = \frac{|\Delta_{\ell\alpha}|}{|\sum_\beta \epsilon_1^{\alpha\beta}|}, \tag{34}$$

we estimate $\eta_e \simeq 0.18$, $\eta_\mu \simeq 0.22$, $\eta_\tau \simeq 0.14$, i.e. the efficiency in all flavours is large. This stems from the fact that the decay of the scalar triplet into two Higgs doublets is strongly out of equilibrium, since $\mathcal{B}_1^\phi\, \Gamma_{T_1} \ll H$. Notice also that the values of the efficiency parameters for the various lepton flavours differ from each other. This is required to guarantee that the total lepton asymmetry does not vanish. The crucial point here is that the strength of the washout parameters is mainly controlled by the Yukawa couplings $\mathbf{Y}^{T_1}_{\alpha\beta}$, which are in general different for each lepton flavour. These couplings enter directly into the Boltzmann equations through the leptonic branching ratios $\mathcal{B}_1^{\alpha\beta}$ (cf. Eqs. (22b) and (22d)). Were all these

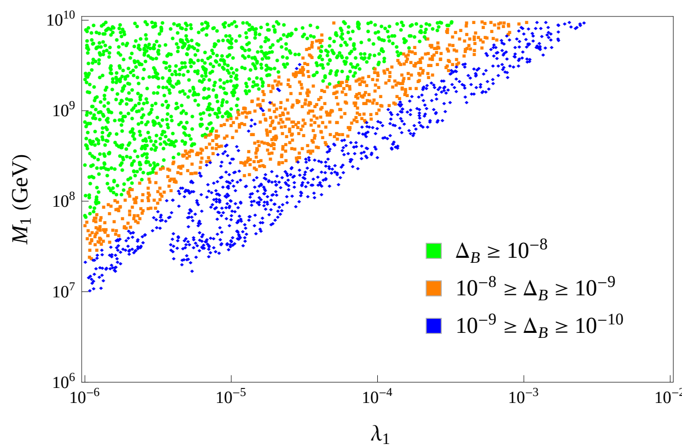

Fig. 3. Baryon asymmetry in the (λ_1, M_1) parameter space, for the example case with the flavoured CP asymmetries of Eq. (21), $M_2 = 10M_1$ and $\lambda_1 = 10\lambda_2$.

branching ratios equal, then the efficiency of leptogenesis would be the same in all flavours, leading to a vanishing total lepton asymmetry [c].

We have also randomly varied the triplet mass M_1 and the coupling λ_1, but keeping the relations $M_2 = 10M_1$ and $\lambda_1 = 10\lambda_2$. The results are presented in Fig. 3, which clearly shows that there is a large region of the parameter space where flavoured type II seesaw leptogenesis is efficient and leads to a viable baryon asymmetry, exclusively dominated by the CP asymmetry coming from the one-loop diagram (1c) in Fig. 1.

4. Conclusions

The class of scenarios where the total unflavoured lepton asymmetry is zero ($\epsilon_1 = 0$), while its individual flavour contributions are not, is known as purely flavoured leptogenesis (PFL)[36]. These scenarios have been studied in type-I seesaw leptogenesis, and it has been shown that, in order to get a sizable contribution from the Boltzmann equations, the various flavour projections have to be different. In the case where these projections are of the same order for each flavour, known as lepton flavour equilibration, PFL may become ineffective[37]. This is however a model-dependent conclusion which cannot be generalized to all PFL models[28,27,36,38]. In the context of type-II seesaw PFL, the relevance of lepton flavour equilibration and flavour-dependent washout processes through the study of the Boltzmann equations has not been addressed yet. The example case presented in this work is a first step in this direction.

[c]We remark that, by summing over the lepton flavours α in Eq. (22d), one does not recover the unflavoured Boltzmann equation for the total lepton asymmetry[33], unless all $\mathcal{B}_1^{\alpha\beta}$ are equal.

From a theoretical viewpoint, a simple way to suppress the contribution of diagram (b) to the leptonic CP asymmetry with respect to the one of diagram (c) (cf. Fig. 1) is by imposing some symmetry that forbids the trilinear terms $\mu_a \tilde{\phi}^T T_a \tilde{\phi}$. Clearly, in this case, no effective neutrino mass term can be generated. To implement the type II seesaw mechanism the symmetry should be softly broken. A well-known example is the soft breaking of lepton number L. For instance, considering a Lagrangian invariant under $U(1)_L$ with the symmetry transformations $\ell_L \to e^{i\alpha_L}\ell_L\,,\, e_R \to e^{i\alpha_L} e_R\,,\, T_a \to e^{-2i\alpha_L} T_a\,,\, \phi \to \phi\,$, only the lepton-number conserving loop diagram (c) is allowed. Neutrino masses can then be generated by an explicit soft breaking of the lepton number, or by the spontaneous breaking induced by the vacuum expectation values of additional scalar fields, $\mu_a \sim \langle \eta \rangle$. A drawback in the latter scenario is the appearance of pseudo Nambu-Goldstone bosons (Majorons) after the spontaneous breaking of the global lepton number, which are tightly constrained by experimental searches. This problem can be avoided by replacing the continuous symmetry by a discrete one; for example, by requiring the Lagrangian to be invariant under the Z_3 transformations $\{\ell_L, e_R, T_a\} \to e^{i2\pi/3}\{\ell_L, e_R, T_a\}$, $\phi \to \phi$.

To conclude, in this work we have studied the relevance of flavoured effects on the leptonic CP asymmetries generated in purely type II seesaw scenarios, where leptogenesis is implemented through the out-of-equilibrium decays of scalar triplets. We have found a novel contribution to the flavoured leptogenesis asymmetries coming from the wave-function renormalization correction, with leptons running inside the loops. This contribution conserves total lepton number but violates lepton flavour and, therefore, does not vanish in the flavoured leptogenesis regime at temperatures $T \lesssim 10^{12}$ GeV. We have also briefly discussed some possible scenarios in which such contribution may dominate the cosmological baryon asymmetry.

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