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Flavor dependence of *CP* asymmetries and thermal leptogenesis with strong right-handed neutrino mass hierarchy

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We prove that taking correctly into account the lepton flavour dependence of the *CP* asymmetries and washout processes, it is possible to obtain successful thermal leptogenesis from the decays of the second right-handed neutrino. The asymmetries in the muon and tau-flavour channels are then not erased by the inverse decays of the lightest right-handed neutrino N_1. In this way, we reopen the possibility of "thermal leptogenesis" in models with a strong hierarchy in the right-handed Majorana masses that is typically the case in models with up-quark neutrino-Yukawa unification.

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The measure of nonvanishing neutrino masses and mixings and the use of the seesaw mechanism to explain the smallness of these masses has made leptogenesis [1] one of the most promising mechanisms to generate the observed value of the baryon asymmetry of the Universe. The seesaw mechanism requires the presence of heavy right-handed neutrinos with complex Yukawa couplings that can generate a lepton asymmetry through their out-of-equilibrium decays. This "lepton" asymmetry is then transformed into baryon asymmetry by the standard model sphaleron processes violating $(B + L)$ [2]. When the right-handed neutrinos are produced through scatterings (mainly inverse decays) in the thermal plasma this scenario is known as "thermal leptogenesis."

In the thermal leptogenesis mechanism, it is usually assumed that the lepton asymmetry is generated by the out-of-equilibrium decays of the lightest right-handed neutrino N_1. The asymmetry generated by the two heavier right-handed neutrinos N_2 and N_3 that we take as hierarchically heavier than N_1, is usually assumed to be either negligible or efficiently washed out by the N_1 inverse decays [3]. In these conditions only N_1 generates a baryon asymmetry. The final value of the baryon asymmetry normalized to the entropy density $\eta_{\rm B}$ is then related to the *CP* asymmetry ε_1 in the decays of the lightest right-handed neutrino as

$$\eta_{\rm B} = \frac{8}{23}\frac{n_{N_1}}{s}\chi_1\varepsilon_1 \simeq 3\times 10^{-4}\chi_1\varepsilon_1, \qquad (1)$$

where the factor $8/23$ is the fraction of the $B - L$ asymmetry converted into baryon asymmetry by sphalerons. n_{N_1}/s is the number density of right-handed neutrinos normalized to the entropy density that in equilibrium is approximately equal to $0.2/g_*$ with $g_* \simeq 106$ the number of propagating states in the standard model plasma. Finally χ_1 is the efficiency factor describing the fraction of the *CP* asymmetry that survives the washout by inverse decays and scattering processes. Now, for hierarchical right-handed neutrinos, it is possible to relate ε_1 to the mass of the lightest right-handed neutrino, the observed neutrino masses, and an effective leptogenesis phase as

$$\varepsilon_1 = \frac{3}{8\pi}\frac{M_1}{v_2^2}(m_3 - m_1)\sin\delta_{\rm L} \le \frac{3}{8\pi}\frac{M_1 m_{\rm atm}}{v_2^2} \simeq 2\times 10^{-6}\left(\frac{M_1}{10^{10}\ {\rm GeV}}\right)\left(\frac{m_{\rm atm}}{0.05\ {\rm eV}}\right), \qquad (2)$$

where the maximum value for this asymmetry corresponds to fully hierarchical neutrinos with $m_3 = m_{\rm atm} \simeq 0.05$ eV, $m_1 \simeq 0$, and a maximal phase $\delta_{\rm L}$.

The baryon asymmetry is precisely measured from the cosmic microwave background as $\eta_{\rm B}^{\rm CMB} = 0.9\times 10^{-10}$ [4]. Thus, using Eqs. (1) and (2) with $\chi_1 = 1$ we obtain a lower bound on the mass of N_1 [5]:

$$M_1 \ge M_1^{\rm min} = 1.5\times 10^9\ {\rm GeV}\left(\frac{\eta_{\rm B}^{\rm CMB}}{9\times 10^{-11}}\right)\left(\frac{0.05\ {\rm eV}}{m_{\rm atm}}\right). \qquad (3)$$

However this lower bound represents a serious problem for many flavour models, specially those incorporating some grand unification symmetry. In these models we usually expect that neutrino Yukawas are closely related to the up-quark Yukawa matrices. As show in the appendix of [6] (see also Ref. [7]) and explained below, in these models it is possible to obtain the mass of the lightest right-handed neutrino as

$$|M_1| \simeq \frac{y_1^2 v_2^2}{|W_{1k}^2 m_{\nu_k}|} \simeq \frac{m_u^2}{|W_{11}^2 m_1 + W_{12}^2 m_{\rm sol} + W_{13}^2 m_{\rm atm}|}, \qquad (4)$$

with W_{ij} the rotation from the basis of diagonal $\nu_{\rm L}$ masses to the basis where the (left-handed) neutrino-Yukawa matrix $Y_\nu Y_\nu^\dagger$ is diagonal, m_{ν_k} the left-handed neutrino masses that we take as hierarchical, and we have approximated $y_1 v_2 \simeq m_u$. From this expression it is straightforward to see that we can only obtain $M_1 \ge 10^9$ GeV if

$$\frac{m_u^2}{10^9\ {\rm GeV}} = 4\times 10^{-6}\ {\rm eV}\left(\frac{m_u^2}{(2\ {\rm MeV})^2}\right) \gtrsim |W_{11}^2 m_1 + W_{12}^2 m_{\rm sol} + W_{13}^2 m_{\rm atm}|. \qquad (5)$$

This implies that, barring accidental cancellations (or some specific textures as in [8]), having a large enough N_1 is only possible if we have simultaneously $m_1 \lesssim 4 \times 10^{-6}$ eV, $W_{12}^2 \lesssim 5 \times 10^{-4}$ (for $m_{\rm sol} \simeq 0.008$ eV), and $W_{13}^2 \lesssim 8 \times 10^{-5}$ (for $m_{\rm atm} \simeq 0.05$ eV). These conditions are usually not satisfied in most flavour models where the typical values for the lightest right-handed neutrino mass are close to 10^6–10^7 GeV with W_{12} of the order of mixing angle in solar neutrinos [9–11]. Therefore, apparently, these models cannot produce a sufficient baryon asymmetry through the thermal leptogenesis mechanism.

It is evident that this represents a very serious problem for all these flavour models, that are forced to abandon thermal leptogenesis and use other mechanisms to produce the observed baryon asymmetry [12]. In this paper we present a solution to this problem. We show that when the flavour dependence of the lepton asymmetries and the erasure processes are correctly taken into account, it is still possible to generate a large enough baryon asymmetry in the decays of the second right-handed neutrino N_2 that then survives the washout due to the inverse decays of N_1. The basic idea behind this statement is that the asymmetries generated by N_2 in the different flavour channels are washed out by different N_1 Yukawa interactions. In fact, the decays of the second right-handed neutrino N_2 create different asymmetries in the τ, μ, and e channels. As we will show these asymmetries do not mix through the thermal interactions with the plasma before the N_1 inverse decays become effective. Therefore, it is evident that the different flavour channels will only be erased by the corresponding $N_1 \to HL_i$ interaction and not by the sum over all lepton flavours as usually done [3,6,13–16]. Taking into account these effects, we prove here that in models with $Y_\nu \simeq Y_u$, the lepton asymmetry in the τ and μ channels is only mildly erased by N_1 decays. Some of these ideas were already present in the literature. For instance, the possibility of generating the asymmetry in N_2 decays was previously discussed by P. di Bari in the one flavour approximation [13]; an early discussion of flavour effects in leptogenesis, mainly in N_1 decays, can be found in [17], and flavour effects in the context of low scale leptogenesis models were recently studied in [18]. In this paper we show that combining these elements it is possible to obtain successful leptogenesis in realistic models of fermion masses with up-quark neutrino-Yukawa unification. Therefore, we reopen the possibility of "thermal leptogenesis" in these models.

In the following we give a general proof of this statement. Although this mechanism is equally valid in supersymmetric and nonsupersymmetric models we only discuss the nonsupersymmetric example. The supersymmetric case is completely analogous [19].

The general expression for leptonic Yukawa couplings and Majorana masses in a seesaw standard model is, $\mathcal{L}_{\rm Yuk} = L^T Y^e e_{\rm R}^c H_1 + L^T Y^\nu \nu_{\rm R}^c H_2 - \frac{1}{2} \nu_{\rm R}^{cT} \mathcal{M} \nu_{\rm R}^c$, with L_i, e_{Ri}, ν_{Ri}, $(i = e, \mu, \tau)$ the three generations of leptons and Y^e, Y^ν, and $\mathcal{M}$, 3×3 matrices. In the basis of diagonal Y^e and $\mathcal{M}$, the neutrino-Yukawa matrix is $Y^\nu = V_{\rm L}^\dagger \cdot D_{Y_\nu} \cdot V_{\rm R}$, where D_{Y_ν} is a diagonal mass of neutrino-Yukawa couplings, that in the class of models discussed here is $D_{Y_\nu} = {\rm diag}(y_1, y_2, y_3) \simeq {\rm diag}(m_u/v_2, m_c/v_2, m_t/v_2)$. The effective Majorana mass for the left-handed neutrinos is obtained through the seesaw mechanism as

$$m_\nu = v_2^2 Y^\nu \cdot \mathcal{M}^{-1} \cdot Y^{\nu T} = U \cdot D_{m_\nu} \cdot U^T, \qquad (6)$$

where the mixing matrix U is the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix that is sufficiently known from neutrino experiments, and $D_{m_\nu} = {\rm diag}(m_1, m_2, m_3)$ the three neutrino masses. At the moment, we know only two mass differences, $\Delta m_{\rm atm}^2 \simeq 2.7 \times 10^{-3}$ eV^2 and $\Delta m_{\rm atm}^2 \simeq 7.0 \times 10^{-5}$ eV^2. Although in the following we assume hierarchical spectrum with normal hierarchy, other cases can be easily treated in the same way.

Now we define the Hermitian matrix $\Lambda \equiv v_2^4 y_1^4 \mathcal{M}^{-1\dagger} \mathcal{M}^{-1}$ [6] and, in the basis of diagonal $Y^{\nu\dagger} Y^\nu$, from Eq. (6) we have

$$\frac{\Lambda}{v_2^4 y_1^4} = D_{Y_\nu}^{-1} V_{\rm L}^* m_\nu^\dagger V_{\rm L}^\dagger D_{Y_\nu}^{-2} V_{\rm L} m_\nu V_{\rm L}^T D_{Y_\nu}^{-1}$$
$$\equiv D_{Y_\nu}^{-1} \Delta^\dagger D_{Y_\nu}^{-2} \Delta D_{Y_\nu}^{-1}, \qquad (7)$$

with $\Delta = V_{\rm L} m_\nu V_{\rm L}^T \equiv W D_{m_\nu} W^T$. From here the matrix Λ can be written as

$$\Lambda_{ij} = (\vec{\lambda}_i^\dagger) \cdot (\vec{\lambda}_j) \equiv \sum_k (\vec{\lambda}_i^*)_k (\vec{\lambda}_j)_k, \qquad \vec{\lambda}_i = \frac{y_1}{y_i} \begin{pmatrix} \Delta_{1i} \\ \frac{y_1}{y_2} \Delta_{2i} \\ \frac{y_1}{y_3} \Delta_{3i} \end{pmatrix}. \qquad (8)$$

Given the strong hierarchy in the neutrino-Yukawa eigenvalues (similar to the hierarchy in the up-quark masses) we have in practice always that $|\vec{\lambda}_1|^2 \gg |\vec{\lambda}_2|^2, |\vec{\lambda}_3|^2$. For instance, to change this hierarchy in the simplified case where we take $W_{13} \simeq 0$ and $y_1/y_3 \simeq 0$ would require that $|\Delta_{11}| < |\Delta_{22}| y_1^2/y^2$. In these conditions it is very easy to understand that the largest eigenvalue of the matrix Λ, corresponding to $v_2^4 y_1^4/|M_1|^2$, will be given in very good approximation by the (1, 1) element of this matrix. Therefore the lightest right-handed Majorana neutrino mass is given by:

$$|M_1| \simeq \frac{y_1^2 v_2^2}{|\lambda_1|} \simeq \frac{y_1^2 v_2^2}{|\Delta_{11}|} = \frac{m_u^2}{|W_{1k}^2 m_{\nu_k}|}. \qquad (9)$$

From this expression we can estimate the mass of the lightest right-handed neutrino. If the observed neutrino mixings are not present in the neutrino-Yukawa matrices and come mainly from the seesaw mechanism itself we have that $V_{\rm L} \simeq \mathbb{1}$ and $W \simeq U_{\rm PMNS}$. Then we have that $|M_1| \simeq 2 m_u^2/m_{\rm sol} \simeq 1 \times 10^6$ GeV. Only in the case $W \simeq \mathbb{1}$ we can have $|M_1| \simeq m_u^2/m_1 = 10^8(4 \times 10^{-5}\ \mathrm{eV}/m_1)$ GeV. Therefore, typically N_1 is too light to generate a sizeable baryon asymmetry. In the same way the associated eigenvector to this eigenvalue is given by $\hat{\lambda}_1 \simeq \vec{\lambda}_1/\Delta_{11}$. Using this eigenvector we can also obtain the Yukawa couplings of N_1 in the basis of diagonal Y^e and $\mathcal{M}$ as

$$(Y_\nu)_{i1} = (V_\mathrm{L}^\dagger \cdot D_{Y_\nu}(\hat{\lambda}_1))_i = \frac{y_1}{\Delta_{11}} V_\mathrm{L}^\dagger \cdot \begin{pmatrix} \Delta_{11} \\ \Delta_{21} \\ \Delta_{31} \end{pmatrix} = \frac{y_1}{\Delta_{11}} \cdot \begin{pmatrix} U_{1k} m_{\nu_k} W_{1k} \\ U_{2k} m_{\nu_k} W_{1k} \\ U_{3k} m_{\nu_k} W_{1k} \end{pmatrix}, \tag{10}$$

where we used that $\sum_i (V_\mathrm{L})^*_{ij} \Delta_{i1} = \sum_k U_{jk} m_{\nu_k} W_{1k}$.

At this point we can already discuss the behavior of a lepton asymmetry created by the out-of-equilibrium decays of the next-to-lightest right-handed neutrino N_2. The asymmetry in the different flavour channels ($a = e, \mu, \tau$), assuming $M_1 \ll M_2 \ll M_3$, is given by [13]

$$\epsilon_2^a = \frac{1}{8\pi (Y^\dagger Y)_{22}} \left(\frac{3}{2} \frac{M_2}{M_3} \mathrm{Im}[Y^*_{a2} Y_{a3} Y^*_{k2} Y_{k3}] + \frac{M_1}{M_2} \mathrm{Im}[Y^*_{a2} Y_{a1} Y^*_{k2} Y_{k1}] \right). \tag{11}$$

Owing to the hierarchy in the neutrino-Yukawa matrices and the right-handed Majorana neutrinos, the first term in Eq. (11) usually dominates and therefore we take it to estimate the size of the asymmetry. Now we have that, $1/v_2^2\, \mathrm{Im}[Y^*_{a2} (m_\nu)_{ak} Y^*_{k2}] = 1/M_3\, \mathrm{Im}[Y^*_{a2} Y_{a3} Y^*_{k2} Y_{k3}] + 1/M_1\, \mathrm{Im}[Y^*_{a2} Y_{a1} Y^*_{k2} Y_{k1}]$. Therefore, the numerator in ϵ_2^a includes the M_3 contribution to the $(m_\nu)_{ak}$ element. To maximize the asymmetry we assume this element to be of order $m_{\nu_3} = m_\mathrm{atm}$ and define the corresponding *CP*-violating phase as δ_ν, then we have

$$\epsilon_2^a \simeq \frac{3}{8\pi} \frac{M_2 m_3}{v_2^2} \frac{|Y_{a2}|}{|Y_{32}|} \delta_\nu, \tag{12}$$

where we use that the hierarchy in the Yukawa elements is such that $Y_{3j} > Y_{2j}, Y_{1j}$. It is clear that asymmetry in the τ channel can be as large as $\epsilon_2^\tau = 3 \times 10^{-6}(M_2/10^{10}\ \mathrm{GeV})$. In the muon and electron channels the asymmetry will be further suppressed by $|Y_{22}|/|Y_{32}|$ and $|Y_{12}|/|Y_{32}|$, respectively, and thus we can expect a smaller asymmetry specially in the electron channel [20].

These asymmetries are partially converted, through sphaleron processes, to a $\Delta_a = \frac{1}{3}B - L_a$ asymmetry and, in part, washed away by N_2 inverse decays. The final lepton asymmetry after the decay of N_2 will be given by

$$\eta_{\Delta_a} \simeq \frac{n_{N_2}}{s} \epsilon_2^a \chi_2^a, \tag{13}$$

with n_{N_2}/s the number density of N_2 normalized to the entropy density [22]. The efficiency factor χ_2^a is the fraction of the produced asymmetry that survives after the end of N_2 decays and inverse decays and it is approximately given by

$$\chi_2^a \simeq \frac{\Gamma(N_2 \to l_a H)}{H|_{T \simeq M_2}} \simeq \frac{\tilde{m}_2^a}{m_*}, \tag{14}$$

with $m_2^a \equiv v_2^2 Y_{a2} Y^*_{a2}/M_2$ [14] and m_* is the equilibrium mass equal to $\simeq 1 \times 10^{-3}$ eV in the SM [14]. The final values of this efficiency factor are model dependent and therefore we keep this factor as a parameter in the following.

At temperatures slightly below M_2 we have three different asymmetries in the Δ_τ, Δ_μ, and Δ_e channels. These asymmetries remain basically unchanged from $T \sim M_2$ to $T \sim M_1$ and, in particular, the different flavours do not mix. In fact, at this temperature, both N_2 and N_3 right-handed neutrinos have already decayed. Therefore the only possible Yukawa interactions in the plasma are those with the right-handed charged leptons and with N_1. Taking into account that the $Y^e_{3,3}$ Yukawa interactions are in thermal equilibrium at temperatures below $\simeq 10^{14}$ GeV, it is more convenient to use the basis of diagonal charged lepton Yukawas to include these interactions in the thermal masses. Thus, only the Yukawa interactions with N_1 can change flavour, although N_1 are very rare in the plasma up to temperatures of the order of M_1 when they reach the equilibrium density [15,16]. Therefore, we have to discus the washout of the different asymmetries by the N_1 decays and inverse decays. First, it is easy to see that this washout will be exponential and not linear if the N_1 are thermally produced at temperatures close to M_1. From Ref. [15] and taking $\varepsilon_1 \simeq 0$, the Boltzman equation for this case would be [23]:

$$\frac{\partial \eta_{\Delta_a}}{\partial z} = -\frac{1}{4} z^2 e^{-z} \sqrt{1 + \frac{\pi}{2} z \frac{\tilde{m}_1^a}{m_*}} A_{ab} \eta_{\Delta_b}, \quad \text{with} \quad A(T \leq 10^9\ \mathrm{GeV}) \simeq \begin{pmatrix} -0.86 & 0.1 & 0.1 \\ 0.06 & -0.65 & 0.017 \\ 0.06 & 0.017 & -0.65 \end{pmatrix}, \tag{15}$$

with $\tilde{m}_1^a \equiv v_2^2 Y_{a1} Y^*_{a1}/M_1$ takes into account that a given asymmetry Δ_a can only be erased by the corresponding lepton flavour, $z = M_1/T$ and $m_* \simeq 1 \times 10^{-3}$ eV [24]. Now we can use Eqs. (9) and (10) and we see that

$$\begin{aligned}\tilde{m}_1^e &= |\sum_k U_{1k} m_{\nu_k} W_{1k}|^2/|\Delta_{11}| \simeq |m_2 W_{12}/\sqrt{2} + m_1 W_{11}/\sqrt{2}|^2/|\Delta_{11}|,\\ \tilde{m}_1^\mu &= |\sum_k U_{2k} m_{\nu_k} W_{1k}|^2/|\Delta_{11}| \simeq | - m_2 W_{12}/2 + m_1 W_{11}/2 + m_3 W_{13}/\sqrt{2}|^2/|\Delta_{11}|,\\ \tilde{m}_1^\tau &= |\sum_k U_{3k} m_{\nu_k} W_{1k}|^2/|\Delta_{11}| \simeq |m_2 W_{12}/2 - m_1 W_{11}/2 + m_3 W_{13}/\sqrt{2}|^2/|\Delta_{11}|,\end{aligned} \tag{16}$$

where $m_3 = m_{\rm atm}$ and $m_3 = m_{\rm sol}$ and we have used bimaximal mixings in the PMNS matrix to make an estimate. The final values of $\tilde{m}_1^a$ depend on the matrix W. In case the mixings in the neutrino-Yukawa matrix are small, we have that $W \simeq U$ and so

$$\begin{aligned}\tilde{m}_1^e &\simeq \left|\frac{m_2}{2} + \frac{m_1}{2}\right| \simeq \frac{m_{\rm sol}}{2},\\ \tilde{m}_1^\mu &\simeq \left|-\frac{m_2}{2\sqrt{2}} + \frac{m_1}{2\sqrt{2}}\right|^2 \Big/ \left|\frac{m_2}{2}\right| \simeq \frac{m_{\rm sol}}{4},\\ \tilde{m}_1^\tau &\simeq \left|\frac{m_2}{2\sqrt{2}} + \frac{m_1}{2\sqrt{2}}\right|^2 \Big/ \left|\frac{m_2}{2}\right| \simeq \frac{m_{\rm sol}}{4}.\end{aligned} \tag{17}$$

Using these values in the case $\eta^{\rm ini}_{\Delta_\tau} \simeq \eta^{\rm ini}_{\Delta_\mu} \gg \eta^{\rm ini}_{\Delta_e}$, Eq. (15) can be easily integrated (taking $\eta_{\Delta_\tau}(z) \simeq \eta_{\Delta_\mu}(z)$, $\eta_{\Delta_e}(z) \simeq 0$ and $A_{\tau\tau} + A_{\tau\mu} = 0.63$) and we obtain

$$\eta^{\rm fin}_{\Delta_{\mu,\tau}} = \eta^{\rm ini}_{\Delta_{\mu,\tau}} e^{-0.75 \tilde{m}_1^{\mu,\tau}/m_*}. \tag{18}$$

If now we take the central value from the solar mass difference as $m_{\rm sol} = 0.008$ eV, we have that the muon and tau asymmetries are only erased by a factor $\exp(-1.5) \simeq 0.22$ and hence they can survive the washout from N_1 decays. This should be compared with the washout that we would obtain in the single flavour approximation, that would be $\simeq e^{-0.75 \sum_a \tilde{m}_1^a/m_*} = e^{-0.75 m_{\rm sol}/m_*} = 2.4 \times 10^{-3}$. Now the final asymmetry in the τ channel would be

$$\begin{aligned}\eta^{\rm fin}_{\Delta_\tau} &\simeq 0.001 \chi_2^\tau \frac{3}{8\pi} \frac{M_2 m_3}{v_2^2} \delta_\nu e^{-0.75 m_{\rm sol}/(4 m_*)}\\ &\lesssim 6.5 \times 10^{-10} (M_2/10^{10}\ {\rm GeV}) \chi_2^\tau,\\ \eta_B &= \frac{28}{79}(\eta^{\rm fin}_{\Delta_\tau} + \eta^{\rm fin}_{\Delta_\mu}) \simeq \frac{56}{79} \eta^{\rm fin}_{\Delta_\tau},\end{aligned} \tag{19}$$

where we have used that $\eta_{\Delta_\tau} \simeq \eta_{\Delta_\mu}$ and $\eta_{\Delta_e} \simeq 0$. Thus in this case, it could be possible to generate a sufficient baryon asymmetry from the μ and τ channels even for efficiency factors χ_2^τ and χ_2^μ of order 0.15 and $M_2 \simeq 10^{10}$ GeV. For smaller efficiency factors we could still have successful leptogenesis with a somewhat heavier M_2. This situation is typically found for instance in models with non-Abelian flavour symmetries. In this paper we have only presented a simple estimate to show that this mechanism can generate a large enough baryon asymmetry. A full computation in a model with an $SU(3)$ flavour symmetry and spontaneous CP violation is now in progress [9,11,25].

From Eq. (16) and taking into account that $\Delta_{11} = \sum_k W_{1k}^2 m_{\nu_k}$ we can see the washout by N_1 can only be stronger than the result obtained in Eq. (17) if the contribution from the atmospheric neutrino mass is sizable (assuming that $m_1 \ll m_{\rm sol}$). This requires basically that $W_{13} m_3 \geq W_{12} m_2$. Clearly this possibility cannot be excluded, but taking into account that $m_2/m_3 \simeq 1/6$, $W = V_{\rm L} \cdot U$, and that $U_{12} \simeq 1/\sqrt{2}$ and $U_{13} \leq 0.2$ it requires large left-handed mixings in the neutrino-Yukawa matrix. This could be possible in models based in Abelian flavour symmetries or discrete symmetries, although it is clear the final result is highly model dependent and the different models must be studied in detail.

We have shown that taking correctly into account the lepton flavour dependence of the CP asymmetries and washout processes, it is possible to obtain successful thermal leptogenesis from the decays of the second right-handed neutrino. The asymmetries in the muon and tau-flavour channels are then not erased by the inverse decays of the lightest right-handed neutrino N_1. Therefore "thermal leptogenesis" is still viable in models with a strong hierarchy in the right-handed Majorana masses that is typically the case in models with up-quark neutrino-Yukawa unification.

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[23] Notice that, due to the exponential washout, order one factors in the A matrices can be important [17].
[24] In this equation we consider only the effects of decays and inverse decays. The effects of $\Delta L = 1$ and $\Delta L = 2$ scatterings do not change the results presented here for $\tilde{m}_1^a \geq 10^{-3}$ eV [15].
[25] O. Vives (work in progress).