More Viable Parameter Space for Leptogenesis

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Abstract

Lepton flavour asymmetries generated at the onset of the oscillations of sterile neutrinos with masses above the electroweak scale can be large enough to partly survive washout and to explain the baryon asymmetry of the Universe. This opens up new regions of parameter space, where Leptogenesis is viable within the type-I seesaw framework. In particular, we find it possible that the sterile neutrinos masses are substantially below $10^9$ GeV, while not being degenerate. However, the required reheat temperature that is determined by the begin of the oscillations lies some orders of magnitude above the sterile neutrino mass-scale.

1 Introduction

While the type-I seesaw mechanism is a very plausible way of generating neutrino masses [1], it may turn out to evade experimental test. Augmenting the Standard Model (SM) by Majorana masses for the active neutrinos introduces nine new parameters (three masses, and in the PMNS matrix, three mixing angles, one Dirac and two Majorana phases), but there are nine extra numbers describing the masses and the couplings of the RHNs. Quite generically, the RHNs decouple from the SM either due to their high masses or their tiny couplings, such that these extra parameters remain practically unobservable.

Therefore, Leptogenesis [2] models based on the seesaw mechanism may likewise escape from direct tests. Nonetheless, the requirement of successful Leptogenesis is very useful in order to gain constraints on the parameter space
of the seesaw model, in particular when embedded in more general beyond
the SM (BSM) frameworks. A very well known example is the lower bound
on the masses of the RHNs, $2 \times 10^9$ GeV \cite{4,3}, which consequently implies
a lower bound of $3 \times 10^9$ GeV on the reheat temperature of the Universe
(assuming a vanishing initial abundance of the RHNs).

Numerous loopholes to this bound arise when considering extensions of
the SM beyond the seesaw scenario (we refer here always to the type-I vari-
ant), see \textit{e.g.} Ref. \cite{5}. Notable are therefore the few possibilities that allow for
smaller RHN masses and lower reheat temperatures, but that yet rely on the
minimal seesaw mechanism, thus avoiding a further proliferation of free pa-
rameters. The most well-known of these is resonant Leptogenesis, relying on
the enhancement of the lepton-number violating charge-parity ($CP$) asymme-
try due to the mixing of nearly mass-degenerate RHNs \cite{6,7,8,9,10,11,12}.

Here, we consider Leptogenesis from a source term that conserves lepton
number but violates the individual flavours. Its presence was first pointed
out in Ref. \cite{13}, and it was subsequently studied in a more accurate manner
in Refs. \cite{14,15,16}. We refer to it as the ARS scenario after the names of
the authors of Ref. \cite{13}. It is based on the usual seesaw model given by the
Lagrangian

$$
\mathcal{L} = \frac{1}{2} \bar{N} (i \not{\!D} - M) N + \bar{\ell} i \not{\!D} \ell + \bar{R} i \not{\!D} R + \partial_{\mu} \phi \partial^{\mu} \phi \\
- \bar{\ell} Y^\dagger \tilde{\phi} - \bar{N} Y^\dagger \tilde{\phi} \tilde{\phi} - \phi \not{\!D} \phi - \phi \not{\!D} \phi \not{\!D} R,
$$

where $\phi$ is the Higgs doublet, $\tilde{\phi} = (\epsilon \phi)^\dagger$ and $\epsilon$ is the SU(2)$_L$-invariant antisym-
metric tensor. The RHNs are given by the Majorana spinors $N_i$, for which we
assume three flavours, $i = 1, 2, 3$, and the left- and right-handed leptons of
the SM by $\ell_a$ and $R_a$, where $a = e, \mu, \tau \equiv 1, 2, 3$. We choose the usual flavour
bases for the RHNs where their mass matrix is $M = \text{diag}(M_1, M_2, M_3)$ and
for the SM leptons such that $h = \text{diag}(h_e, h_\mu, h_\tau)$. For definiteness, we take
$M_1 < M_2 < M_3$.

In the ARS scenario, oscillations of the RHNs are the dominant source of
flavour asymmetries (that conserve total lepton number) in the SM leptons $\ell$
at temperatures $T \gg M_i$. The RHN masses are typically taken to be at the
GeV-scale or below \cite{13,14,15}, such that the couplings $Y$ are suppressed
enough (note the seesaw relation $m_\nu = Y^t M^{-1} Y v^2 / 2$, with $m_\nu$ the mass
matrix of the SM neutrinos and $v = 246$ GeV) that the asymmetries are only
weakly washed out prior to the electroweak phase transition (EWPT) at the
The temperature $T_{EW} \approx 140$ GeV, where baryon-number ($B$) violating sphaleron processes freeze out. Some washout is essential though, because in order to end up with a net lepton asymmetry, the individual flavour asymmetries must experience different suppression. In Ref. [15], it is demonstrated that in spite of the low mass scale of the RHNs, no degeneracy is necessary in order to account for the observed baryon asymmetry of the Universe (BAU).

Here, we propose that ARS-type Leptogenesis is also viable with RHNs above the electroweak scale, more precisely in the mass range from $10^4$ GeV–$10^8$ GeV and mass ratios of order one. For the subsequent considerations, it is useful to recall the equilibrium neutrino mass $m_\star = Y_\nu^2 v^2/(2M_i)$, where $\Gamma_\star = Y_\nu^2 M_i/(8\pi)$ and $H\big|_{T=M_i} = \Gamma_\star$ ($H$ being the Hubble rate), such that $m_\star \approx 1.1 \times 10^{12}$ GeV. In the present scenario, the flavour asymmetries are produced at temperatures above the masses $M_i$ for all three $N_i$. Given that the observed neutrino mass-differences are larger than $m_\star$ and the large mixing angles, the asymmetries within each of the $\ell_{e,\mu,\tau}$ experience strong washout, defined through the condition $\sum_i Y_{ai}^\dagger Y_{ai} M_i/(8\pi) \geq H$ (i.e. the washout rate for $\ell_a$ exceeds the Hubble rate) at some instance during the expansion of the Universe. Therefore, the conversion of the flavour asymmetries into a baryon asymmetry is less efficient than in the usual ARS scenario. However, the mass differences of the SM neutrinos are not too far above $m_\star$, such that there should be parametric regions where the washout of the asymmetries is not prohibitively strong. Moreover, it turns out that the stronger washout can be made up for by larger initial asymmetries. In Section 2 we outline the way of calculating the BAU in the proposed model and in Section 3 we identify a viable point in parameter space and discuss some parametric dependences. We conclude in Section 4.

2 Lepton Asymmetry and Washout

Oscillations of the RHNs and $CP$ violation.— We now consider the dynamics of the seesaw model when $M_i \gg T_{EW}$ and for a high reheat temperature of the SM particles. While the Universe expands, the Higgs particle and leptons are maintained very close to thermal equilibrium through gauge interactions, while the RHNs may be far from equilibrium over a wide temperature range, because they are only interacting through the Yukawa couplings $Y$. As a plausible initial condition, we assume that the abundance of the RHNs vanishes. The RHNs are then produced from interactions with
leptons and Higgs bosons. For $T \gg M_i$, the fastest processes involve the radiation of an extra gauge boson (2 ↔ 2 processes) \[17, 18\], while for $T \ll M_N i$, inverse decays (2 ↔ 1 processes) dominate.

Now, provided the matrix $Y Y^\dagger$ is non-diagonal, the RHNs emerge as superpositions of their mass eigenstates and oscillate. As a result, they can decay into leptons $\ell$ and Higgs-bosons $\phi$ in an asymmetric manner \[13, 14\]. It is therefore interesting to note the temperature, at which the first oscillation is completed for a typical RHN with momentum of order of the temperature $T$. In order to quantify this, we define $a_R = m_{Pl} \sqrt{45/(4\pi^3 g_*)}$, where $m_{Pl}$ is the Planck mass and $g_*$ the number of relativistic degrees of freedom of the SM at high temperatures. For the scale factor in the radiation-dominated Universe, we take $a(\eta) = a_R \eta$, where $\eta$ is the conformal time. By this choice, $T = 1/\eta$. Besides, we define $z = T_{ref}/T = \eta T_{ref}$. Note that the choice $T_{ref} = M_1$ leads to the usual definition for the parameter $z$ in Leptogenesis calculations \[3\]. More than one full oscillation is completed, when

$$\int_0^\eta \frac{|M_i^2 - M_j^2|}{2T} a(\eta)d\eta = \frac{\Delta M^2 a_R \eta^3}{6} \gtrsim 2\pi$$

$$\Leftrightarrow z \gtrsim z_{osc}^{ij} = \left[12\pi T_{ref}^3 (a_R |M_i^2 - M_j^2|)^{\frac{1}{2}} \right]^{\frac{1}{3}},$$

what consequently defines $T_{osc}^{ij} = T_{ref}/z_{osc}^{ij}$ and $T_{osc}^{13} = T_{osc}^{13}$, as the highest of these oscillation temperatures. Below $T_{osc}^{ij}$, the oscillations can be averaged, what leads to a result in agreement with the usual perturbative calculations for the decay asymmetry of the RHNs \[6, 7, 8, 9, 10, 11\]. We recall at this point that for the conventional lepton-number violating contributions, the asymmetry is proportional to $M_i M_j/|M_i - M_j|^2$, while the purely flavoured source in the ARS scenario is enhanced by the factor $T^2/|M_i - M_j|^2$ \[15\]. The absence of the chirality-flipping insertions of Majorana masses thus explains why the ARS-type source yields much larger asymmetries at temperatures $T \gg M_i$ than the conventional source.

The charge density $q_{\alpha}$ in the lepton flavour $\alpha$ is dominantly produced at
temperatures around $T = T_{\text{osc}}$ and is given by \cite{14, 15}

$$
\frac{q_{\ell a}}{s} \approx -\frac{1}{g_w} \sum_{j \neq i} \frac{1}{y_{ai}^* y_{cj}^* y_{ja} - y_{ai} y_{cj} y_{ja}} \frac{\text{sign}(M^2_{ii} - M^2_{jj})}{M^2_{ii} - M^2_{jj}}
$$

\times \left( \frac{M^2_{Pl}}{|M^2_{ii} - M^2_{jj}|} \right)^{\frac{2}{3}} \times 8.4 \times 10^{-5} \gamma_{av}^2 ,

(3)

where $s$ is the entropy density. Below the temperature $T_{osc}$, the $q_{\ell a}/s$ are approximately conserved initially until washout processes become important. We take here a numerical factor that is slightly smaller compared to the one of Ref. \cite{15}, due to a more conservative estimate of $T_{osc}$, see relations (2). Note that Eq. (3) corresponds to a conservative estimate that may receive corrections of order one \cite{15}. An accurate leading order calculation appears viable and will be tackled in the future. One can immediately verify that Eq. (3) implies that $\sum a q_{\ell a} = 0$, because these asymmetries are purely flavoured. Note that differently from Ref. \cite{15}, we define $q_X$ as the charge density within one component of the SU(2) multiplet, hence the factor $1/g_w$ with $g_w = 2$. We take the value $\gamma_{av} = 3 \times 10^{-3}$, following Ref. \cite{15}. One may take into account a temperature dependence due to the running coupling constants \cite{17, 20}, which we ignore here in view of the accuracy of other approximations used in the present calculation of the freeze-out asymmetry.

In addition, we estimate the temperature $T_{res}$, above which the largest of the widths of the RHNs exceeds the mass splitting, such that Eq. (3) should not be applied, as

$$
T_{res} = \sqrt{(M^2_{ii} - M^2_{jj})/(\gamma_{av} \text{max}[YY^\dagger]|_{ii})}.
$$

(4)

**Washout of left-handed leptons.**— A net lepton asymmetry is then present at the time of sphaleron freeze-out during the EWPT, provided the individual lepton flavours are affected differently by the washout. For non-relativistic RHNs the washout rate of $\ell_a$ is

$$
W_{a}^{\text{NR}} = \sum_i Y_{ai} Y_{ia} \left( \frac{M_i}{T_{\text{ref}}} \right)^{\frac{1}{2}} \frac{3}{T_{\text{ref}}^{2\frac{3}{2} \pi^2 \gamma_{av}^2}} e^{-\frac{2M_i}{T_{\text{ref}}}} \approx \frac{3}{T_{\text{ref}}^{2\frac{3}{2} \pi^2 \gamma_{av}^2}} = \tilde{W}_{a}^{\text{NR}} \sum_i Y_{ai} Y_{ia} ,
$$

(5)
while in the non-relativistic regime, the rate of equilibration of $\ell_a$ with the RHNs is

$$W^\text{rel}_a = \sum_i Y^\dagger_{ai} Y_{ia} \gamma_{av} T = \bar{W}^\text{rel}_a \sum_i Y^\dagger_{ai} Y_{ia}. \quad (6)$$

One should note that the latter rate is lepton-number conserving, as it does not rely on the insertion of Majorana mass terms, which only enter at sub-leading order in the relativistic regime. Below, we consider a region in parameter space where $|Y_{ie}| \ll |Y_{i\mu}, Y_{i\tau}|$, such that only $q_{le}$ is not nullified by washout, and any asymmetry that is transferred from $q_{le}$ into relativistic RHNs will scatter or decay into $\ell_{\mu,\tau}$. We therefore take for the effective washout rate of $\ell_e$

$$W_e = \sum_i Y_{ai}^\dagger Y_{ia} \left[ \vartheta(T - M_i) \max(\bar{W}^{\text{NR}}_a, \bar{W}^\text{rel}_a) + \vartheta(M_i - T) \bar{W}^{\text{NR}}_a \right]. \quad (7)$$

This accounts for the relativistic $2 \leftrightarrow 2$ scatterings at high temperatures, for the $1 \leftrightarrow 2$ decays and inverse decays when $M_i \gtrsim T$ and eventually for the freeze-out of the reactions due to Maxwell suppression when $M_i \gg T$.

**Complete and partial chemical equilibrium of right-handed SM leptons.**— The Yukawa couplings $h_a$ mediate chemical equilibration between the $\ell_a$ and R$_a$, at a rate that is given by $h_a^2 \gamma_{fl} T$ with $\gamma_{fl} = 5 \times 10^{-3}$ [19, 20]. Full equilibration implies $q_{le} = q_{Ra}$. Again, we ignore the temperature dependence of $\gamma_{fl}$, which is present because of running coupling strengths. The redistribution of the asymmetries is quantitatively relevant, because it reduces the washout rates, as only the $\ell_a$ couple to the RHNs. While for $\tau$ and $\mu$, washout can be assumed to be complete no matter whether chemical equilibrium with R$_{\mu,\tau}$ is established or not, the $e$ leptons equilibrate around $T = 3.1 \times 10^4 \text{GeV}$ (defined by the condition $h^2 \gamma_{fl} T = H$), at the bottom of the relevant mass range. We therefore need to account for a partial equilibration of R$_e$ with $\ell_e$, which is done through Eqs. (8) below. Other spectator effects may be included following Ref. [21], what would impact the results at the 10% to 20% level, which we ignore here.

**Final baryon asymmetry.**— Putting together above details, equations that can be used to determine the BAU for the present scenario are simply
given by \((\text{cf. Ref. } [19])\):

\[
\frac{dq_{\ell e}}{dz} = -\gamma \frac{a_R}{T_{\text{ref}}} h_e^2 (q_{\ell e} - q_{Re}) - W_e q_{\ell e}, \quad (8a)
\]

\[
\frac{dq_{Re}}{dz} = -\gamma \frac{a_R}{T_{\text{ref}}} h_e^2 (q_{Re} - q_{\ell e}). \quad (8b)
\]

As initial conditions at \(z = 0\), we take \(q_{Re} = 0\) and \(q_{\ell e}\) obtained from Eq. (3).

The solution evaluated at \(z = T_{\text{ref}}/T_{\text{EW}}\) yields the e-lepton charge density \(L_e = g_w q_{\ell e} + q_{Re}\) at the EWPT. Due to the strong washout in the \(\mu, \tau\) flavours, there is no asymmetry in these sectors, and we obtain for the BAU

\[
B \approx -(28/79) L_e \quad [22],
\]

which is to be compared with the observed value \(B_{\text{obs}}/s = 8.8 \times 10^{-11}\) \([23, 24]\).

### 3 Example Scenario

On the case of a specific parametric reference point, we demonstrate now that the proposed scenario can explain the BAU in a phenomenologically viable way. We use the parametrisation of the neutrino Yukawa couplings from Ref. \([25]\) and follow the notation of Ref. \([15]\). For the mixing angles of the PMNS matrix and the mass differences of the SM neutrinos, we take the best fit values from Ref. \([26]\). We assume a normal mass hierarchy with the lightest neutrino mass \(m_1 = 2.5\) meV. As for the RHNs, we take \(1 : 2 : 3\) for the mass ratios \(M_1 : M_2 : M_3\). For the remaining angles, we choose

\[
\begin{align*}
\delta &= 0.2, \\
\omega_{23} &= 0.6 + 1.4i, \\
\alpha_1 &= 0, \\
\omega_{13} &= 0.1 - 1.5i, \\
\alpha_2 &= 2.6, \\
\omega_{12} &= -1.9 - 1.0i.
\end{align*}
\]

We vary \(M_1\), keeping the mass ratios of the RHNs fixed. In Figure 1 we show the ratio of the final BAU to the observed value as a function of \(M_1\). It turns out that \(B/B_{\text{obs}} = 1\) for \(M_1 \approx 2.5 \times 10^5\) GeV and \(T_{\text{osc}} \approx 3.0 \times 10^9\) GeV. This value for \(T_{\text{osc}}\) can be interpreted as the minimum reheat temperature for the scenario studied here, and it happens to coincide with the corresponding value for standard Leptogenesis \([3]\). The increase of the asymmetry with \(M_1\) can be easily understood in terms of the initial asymmetry \((3)\) and the seesaw relation. The deviation from a simple power law, visible as a wiggle, is due to the partial equilibration of the \(R_e\), described by Eqs. \((8)\). The
mass range in which our approximations are viable is given by the condition $T_{\text{osc}} < T_{\text{res}}$, and from Figure 2 we see that we should choose $M_1$ to be below about $10^{11}$ GeV.

Within the studied range for $M_1$, we find that $L_{\ell e} |_{T=T_{\text{EW}}} / L_{\ell e} |_{T=T_{\text{osc}}}$ varies between 0.25 and 0.39, indicating that for the reference point, the $|Y_{ie}|$ take small values such that the washout of the $\ell_e$ is as small as possible. Moreover, the moderately large imaginary parts of the mixing angles imply some tuning that may be estimated as $\cosh(\omega_{23})^2 \cosh(\omega_{13})^2 \cosh(\omega_{12})^2 \approx 60$, what implies that the $Y_{ia}$ are generically about a factor eight larger than expected from the standard seesaw relation without a particular alignment. This amount of tuning can also be explicitly verified by observing cancellations of individual
terms that add up to the largest of the light neutrino masses. Therefore, finding viable mixing angles for small $M_i$ requires parametric tweaks. However, from Figure 1 it is clear that the situation becomes largely relaxed when allowing for larger $M_i$.

4 Conclusions

We have studied the ARS mechanism for RHNs above the electroweak scale, what opens up a substantial region of parameter space for Leptogenesis based on the type-I seesaw model. Non-degenerate RHNs with masses substantially below $10^9$ GeV may account for the BAU, what has not been found to be viable before and allows for new cosmologically consistent BSM scenarios. One should note however that the reheat temperature given by $T_{osc}$ is by orders of magnitude above the mass scale of the RHNs [cf. Eq. (2)], unless assuming a mass degeneracy. Going to lower masses and reheat temperatures appears to require parametric tuning, either of the Yukawa couplings or the RHN masses, that should then be taken to be close to degenerate. (In turn, the source term that we discuss here generically adds asymmetries to scenarios of resonant Leptogenesis at low temperature scales.) Another option may be to extend the particle content of the model in such a way that the production of RHNs at the temperature $T_{osc}$ becomes more efficient. It is interesting to note that the loop correction to the Higgs mass-square at the parameter point with the lowest viable RHN masses considered here ($M_1 = 2.5 \times 10^5$ GeV) yields $\sum_i [YY^\dagger]_{ii} M_i^2/(16\pi^2) \approx 200$ GeV$^2$, such that the electroweak scale is not necessarily destabilised by the RHNs, a desirable feature also emphasised in Ref. [27]. Future work on this scenario should encompass a more systematic analysis of the viable parameter space involving three RHNs as well as a more accurate determination of the relevant $CP$-violating rates.

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