Limits on $T_{\text{reh}}$ for thermal leptogenesis with hierarchical neutrino masses

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Abstract

We make a simple observation that if one of the right-chiral neutrinos is very heavy or its Yukawa couplings to the standard lepton doublets are negligible, so that it effectively decouples from the seesaw mechanism, the prediction for the baryon asymmetry of the Universe resulting from leptogenesis depends, apart from the masses $M_1$ and $M_2$ of the remaining two right-chiral neutrinos, only on the element $\tilde{Y}_{\nu_{22}}$ of the neutrino Yukawa coupling. For $M_2 \gtrsim 10M_1$ the lower bound on $M_1$ and also on $T_{\text{reh}}$, resulting from the requirement of 'successful leptogenesis' is then significantly increased compared to the one computed recently by Buchmüller et al. in the most general case. Within the framework of thermal leptogenesis, the only way to lower this limit is then to allow for sufficiently small mass difference $M_2 - M_1$. 
Accumulated over the past years data on solar and atmospheric neutrino oscillations are consistent with the minimal see-saw mechanism of neutrino mass generation. Barring the LSND anomaly, all the remaining data can be explained by the mixing of three light active neutrino species with differences of the masses squared equal to:

$$\Delta m_{\text{sol}}^2 \approx 7 \times 10^{-5} \text{eV}^2 \quad \Delta m_{\text{atm}}^2 \approx 2.5 \times 10^{-3} \text{eV}^2$$

and the so-called bi-large pattern of the unitary mixing matrix $U$. In principle, the differences (1) can result from different patterns of neutrino masses $m_{\nu_i}$ satisfying the WMAP bound \(\sum m_{\nu_i} < 0.7 \text{ eV}\), but the most natural in the context of the see-saw mechanism operating at some scale $M$ close to the GUT scale $\sim 2 \times 10^{16} \text{ GeV}$ seems to be the hierarchical pattern $m_{\nu_3} \approx \sqrt{\Delta m_{\text{atm}}^2}$, $m_{\nu_2} \approx \sqrt{\Delta m_{\text{sol}}^2}$, for which the mixing pattern is stable with respect to radiative corrections. In this mechanism, whose main virtue is that it naturally explains why neutrinos are so light, three heavy Majorana neutrinos $\nu^c_A (A = 1, 2, 3)$ couple to the standard $SU(2)$ lepton doublets through the Lagrangian term $L = \epsilon_{ij} H \nu^c_A Y^A l_j B + \text{H.c.}$ Below the mass scale of the lightest right-chiral neutrino the effective light neutrino mass matrix takes the form

$$m_{\nu} = -\frac{v^2}{2} Y^T \cdot M^{-1} \cdot Y_{\nu} = \frac{v^2}{2M_1} U_{\nu}^\dagger \cdot \text{diag}(\xi_1^2, \xi_2^2, \xi_3^2) \cdot U_{\nu}^\dagger,$$

where $M = M_1 \text{diag}(1, x_2^{-2}, x_3^{-2})$ is the right-chiral Majorana mass matrix, $v = \sqrt{2} \langle H^0 \rangle$ and $U_{\nu}$ diagonalizes $m_{\nu}$.

Another attractive feature of the see-saw mechanism is the possibility of explaining the baryon number $B$ asymmetry of the Universe by the out of equilibrium lepton number $L$ violating decays of the right-chiral neutrinos. The net lepton number generated by these decays at temperatures $T \lesssim M_1$ is subsequently converted into the baryon number by the sphaleron mediated $B + L$ violating transitions. Different theoretical models of the see-saw mechanism that aim at reproducing the experimentally measured neutrino properties in a natural way can be therefore further constrained by the requirement of 'successful leptogenesis'.

There are two classes of theoretical models of the see-saw mechanism which can give bi-large mixing for hierarchical light neutrino mass spectrum $m_{\nu_3} \approx \sqrt{\Delta m_{\text{atm}}^2}$, $m_{\nu_2} \approx \sqrt{\Delta m_{\text{sol}}^2}$, without too much artificial fine tuning of the parameters. These are the so-called 'lopsided' models in which large neutrino mixing arises from the charged lepton sector and models in which it is due to some special textures of the neutrino Yukawa matrix $Y_{\nu}$ in the mass eigenstate basis of the charged lepton. In both types of models the see-saw mechanism is usually dominated by contributions of one or at most two right-chiral neutrinos, what most naturally happens if their masses exhibit some hierarchy $1 = x_1 > x_2 \gg x_3$. One considers also minimal see-saw models with exactly two right-chiral neutrinos.

The purpose of this note is to investigate leptogenesis in a general class of theoretical see-saw models, in which the contribution to the see-saw formula of the third right-chiral neutrino is negligible. We will be working in the basis, in which the charged lepton Yukawa couplings are diagonal and therefore $U = U_{\nu}$. As it is well-known, the neutrino Yukawa coupling

\footnote{This bound relies on some theoretical assumption and can be questioned.}
matrix $Y_\nu$, which enters the formulae for the CP asymmetry parameters $\epsilon_i$ [13] and various reaction cross sections, cannot be uniquely reconstructed from the low energy quantities $m_{\nu_i}$ and $U_\nu$ and the Majorana masses $M_A$. Rather, one has [12, 14, 15, 16]

$$\tilde{Y}_{\nu}^{AB} \equiv (Y_\nu U_\nu)^{AB} = ix_A^{-1}\Omega_{AB}\xi_B \tag{3}$$

(no summation over $A$ and $B$), where the complex orthogonal matrix $\Omega$ accounts for the six-parameter ambiguity in translating the neutrino masses into the neutrino Yukawa couplings $\tilde{Y}_{\nu}$. We first point out that if one of the right-chiral neutrinos effectively decouples (either because it is very heavy or its couplings $\tilde{Y}_{\nu}^{3A}$ are negligible) the prediction for the baryon asymmetry of the Universe generated via leptogenesis depends only on $M_1, x_2$ and $\tilde{Y}_{\nu}^{22}$, that is, the six-parameter ambiguity encoded in $\Omega$ reduces to the dependence on two parameters only, which can be chosen as the modulus and the phase $\varphi$ of $\tilde{Y}_{\nu}^{22}$. For $M_2 \gg M_1$ ($x_2 \ll 1$), when the lepton number generated by the out of equilibrium decays of the second heavy neutrino is completely washed out by the time when the lightest right-chiral neutrinos decay, or simply the reheat temperature $T_{\text{reh}}$ is too low for the second heavy neutrino to be produced thermally, the leptogenesis depends only on $M_1$ and the product $x_2 \tilde{Y}_{\nu}^{22}$ which can be translated into $\varphi$ and the parameter $\tilde{m}_1$ defined in [17]. Compared to the analysis of such a scenario performed in [15] for the thermal production of heavy neutrinos, the assumption about the decoupling of one right-chiral neutrino has the following consequences. Firstly, for each point of the $(M_1, \tilde{m}_1, \varphi)$ parameter space the relevant CP asymmetry can be unambiguously computed; its maximal value is smaller by about 20% compared to the value $\epsilon_1^{\text{max}}$ found in [14] and adopted in [15]. Secondly, and more importantly, the range of possible $\tilde{m}_1$ values is substantially reduced: $\tilde{m}_1$ is always greater than $m_{\nu_2}$ (and not $m_{\nu_1}$ as in [15]). This significantly increases the lower limit on $M_1$ and, consequently, on the reheat temperature $T_{\text{reh}}$ for which the leptogenesis mechanism dominated by the decays of the lightest right-chiral neutrino can reproduce the observed baryon asymmetry of the Universe $Y_B \approx (0.8 - 1) \times 10^{-10}$ [1]. In this letter, for technical simplicity we integrate the Boltzmann equations on the non-supersymmetric scenario [17], but we believe that our results qualitatively hold in the supersymmetric case as well. The lower limit $T_{\text{reh}} \gtrsim 2 \times 10^{12}$ GeV (assuming $T_{\text{reh}} \gtrsim 10 M_1$), which we find in the strict decoupling limit, may then lead to too abundant gravitino production, incompatible with the standard nucleosynthesis. The only way to lower $T_{\text{reh}}$ is to consider almost degenerate two lighter right-chiral neutrinos as advocated in [19]. The simplification occurring in the decoupling limit allows us to extend the analysis also to this case. In particular, we investigate how fast the necessary reheat temperature $T_{\text{reh}}$ decreases with increasing degeneracy of the two lighter right-chiral neutrinos. At the end, we briefly discuss the conditions under which the results described here should apply.

For fixed $\xi_A$ (i.e. for fixed see-saw masses of the left-chiral neutrinos), $M_1, x_2$ and $x_3$ the formula (2) constitutes a set of $2(3 + 6/2) = 12$ real equations allowing to solve for six complex elements of $\tilde{Y}_{\nu}^{AB}$ in terms of the remaining three. This six-parameter ambiguity is encoded in the complex orthogonal matrix $\Omega$ in eq. (3). If the third right-chiral neutrino is decoupled or simply absent there are only six complex Yukawa couplings and, at first sight, the formula (2) should determine all of them unambiguously. However, it is easy to see that the resulting set
of equations
\[ \xi_A^2 \delta^{AB} + \sum_C x_C^2 \bar{Y}_\nu^C A \bar{Y}_\nu^{CB} = 0 , \quad A, B = 1, 2, 3 \] (4)
with \( C = 1, 2 \) is self-consistent only if \( m_{\nu_1} \propto \xi_1^2 = 0 \). In other words, the absence or decoupling of one right-chiral neutrino automatically ensures \( \det(m_\nu) = 0 \), that is \( m_{\nu_1} = 0 \) (with only one right-chiral neutrino two eigenvalues of \( m_\nu \) would vanish) and, hence, hierarchical or inversely hierarchical spectrum of the left-chiral neutrinos. For \( m_{\nu_1} = 0 \) not all equations (4) are independent; as a result there is still a two-parameter freedom which can be parametrized by \( \bar{Y}_{\nu_2}^{22} \).

Since for one right-chiral neutrino (no summation over \( C \) in eq. (4)) all Yukawa couplings are unambiguously determined by the formula (2), it follows, that for three right-chiral neutrinos a convenient parametrization of the ambiguity is given by the \( \bar{Y}_{\nu_2}^{22}, \bar{Y}_{\nu_2}^{32} \) and \( \bar{Y}_{\nu_2}^{33} \), as it allows to take the limit in which heavy right-chiral neutrinos are successively decoupled. From the above it is clear that the description of leptogenesis simplifies enormously in the decoupling limit for which the first condition is the sufficiently small value of \( m_{\nu_1} = 0 \) and, hence, the spectrum of left-chiral neutrinos close to hierarchical or inversely hierarchical.

For \( x_3 \bar{Y}_{\nu_2}^{32} = 0 \) and \( x_3 \bar{Y}_{\nu_2}^{33} = 0 \) the exact solution of the equations (4) in terms of the matrix \( \Omega \) reads
\[ \Omega = \begin{pmatrix} 0 & \pm iz & \pm i \sqrt{-1 - z^2} \\ 0 & -iz & i \sqrt{-1 - z^2} \\ 1 & 0 & 0 \end{pmatrix} , \quad \text{where} \quad z = (x_2 / \xi_2) \bar{Y}_{\nu_2}^{22} . \] (5)

It is obvious that all non-vanishing quantities: \( \bar{Y}_{\nu_2}^{12}, \bar{Y}_{\nu_2}^{13} \) and \( x_2 \bar{Y}_{\nu_2}^{23} \) depend then only on \( x_2 \bar{Y}_{\nu_2}^{22} \). The decoupling of the third right-chiral neutrino is a somewhat stronger requirement as it corresponds to \( x_3 \bar{Y}_{\nu_2}^{3A} \to 0 \) for all \( A = 1, 2, 3 \). In agreement with our previous discussion this requires one massless left-chiral neutrino because, as follows from eqs. (3) [5]
\[ x_2 \bar{Y}_{\nu_2}^{31} = i \xi_1 . \] (6)

While the results for the baryon asymmetry described below depend only on the particular form [5] of the matrix \( \Omega \), it is clear that the most natural way to realize the limit \( x_3 \bar{Y}_{\nu_2}^{32} = 0 \) and \( x_3 \bar{Y}_{\nu_2}^{33} = 0 \) is to have \( x_3 \to 0 \) and therefore also one nearly massless left-chiral neutrino.

Consider now the scenario with \( x_2 \lesssim 0.3 \), in which the decays of the second right-chiral neutrino do not contribute to the generated lepton asymmetry. In this limit only the asymmetry parameter \( \epsilon_1 \) is relevant and can be approximated by\(^2\) [4]
\[ \epsilon_1 \approx -\frac{3}{16\pi} \frac{2M_1}{v^2} \sum_A \text{Im} (\Omega_{1A}^2) m_{\nu_A}^2 . \] (7)

For \( \Omega \) given in [3] it is easy to find that
\[ \epsilon_1 \approx \begin{cases} \frac{-(3/16\pi)(\xi_3 / \xi_2)^4 |x_2 \bar{Y}_{\nu_2}^{22}|^2 \sin(2\varphi)}{\text{for} \ |z| \ll \xi_2 / \xi_3} \\ \frac{-(3/16\pi)(\xi_3^2 - \xi_2^2) \sin(2\varphi)}{\text{for} \ |z| \gg 1} \end{cases} , \] (8)
\(^2\)This is for the non-supersymmetric case; in this limit supersymmetric \( \epsilon_1 \) is twice as big [13].
Figure 1: $|\epsilon_1|$ and $\tilde{m}_1$ as a function of $|\tilde{Y}^{22}|$ for $M_1 = 10^{10}$ GeV, $x_2 = 0.1$ and strictly hierarchical light neutrino masses. Solid, dashed and dotted lines correspond to $\varphi = \arg(\tilde{Y}^{22}) = \pi/4$, $\pi/8$ and $\pi/12$, respectively. $\tilde{m}_1$ is insensitive to $\varphi$. The dash-dotted line shows $|\epsilon_1|_{\text{max}}$ from [14].

where $\varphi = \arg(\tilde{Y}^{22})$. For hierarchical masses of the left-chiral neutrinos this is illustrated in fig. [1a] where we plot $|\epsilon_1|$ as a function of $\tilde{Y}^{22}$ for fixed $M_1$ and $x_2$. The maximal value $|\epsilon_1|_{\text{max}} = (3/16\pi)(2M_1/v^2)(m_{\nu_3} - m_{\nu_2})$ attained in the decoupling limit is about 20% lower than the upper limit $|\epsilon_1|_{\text{max}} = (3/16\pi)(2M_1/v^2)\sqrt{\Delta m_{\text{atm}}^2}$ in the general hierarchical case derived in [14] and used in [15]. In the inversely hierarchical case, which is also compatible with the decoupling limit, $|\epsilon_1|$ is generically smaller, as then $m_{\nu_3} - m_{\nu_2} \lesssim (\Delta m_{\text{sol}}^2/2\sqrt{\Delta m_{\text{atm}}^2}) \ll \sqrt{\Delta m_{\text{atm}}^2}$.

Efficient generation of the lepton number asymmetry requires not only a sufficiently large value of $|\epsilon_1|$, but also that the lightest right-chiral neutrinos decay out of equilibrium. This requirement can be expressed as the condition that their decay rate $\Gamma_1$ must be smaller than the Hubble parameter $H$ at the temperature $T \sim M_1$ [20, 17]. This in turn translates into the following condition [14]:

$$\tilde{m}_1 \equiv \sum_A |\Omega^{2}_{1A}|m_{\nu_A} \lesssim 5 \times 10^{-3} \text{ eV}.$$  \hspace{1cm} (9)

For $\tilde{m}_1$ of this order and larger the wash-out processes are very efficient and suppress the resulting lepton asymmetry [17, 15]. In the decoupling limit the condition (9) can hardly be satisfied, because, as follows from the form (5) of the matrix $\Omega$, one has $\tilde{m}_1 > m_{\nu_2} \gtrsim 8 \times 10^{-3}$ eV in the hierarchical case (for inverse hierarchy $\tilde{m}_1 \gtrsim 0.05$ eV). Moreover, the smallest possible values of $\tilde{m}_1$ are obtained for $|z| \ll 1$ where $\epsilon_1$ is suppressed (see fig. [1]). As a result, in order to reproduce the observed baryon asymmetry through the leptogenesis the mass $M_1$ of the lightest right-chiral neutrino has to be about two orders of magnitude bigger than its lower limit found.
Figure 2: Lower limits on $M_1$ for $x_2 = 0.01$, 0.1, 0.3 and 0.95 (solid, dashed, dotted and dash-dotted lines, respectively) as functions of $|\tilde{Y}_{22}^{\nu}|$ (panel a) and $\tilde{m}_1$ (panel b).

in [15] for general see-saw models leading to the hierarchical spectrum of light neutrinos. This is illustrated in fig. 2 where the lower limits on $M_1$ (resulting from the requirement that the computed $B - L$ abundance $Y_{B-L}$ is not smaller than $2.8 \times 10^{-10}$) are plotted for $x_2 = 0.01$ (solid line), $x_2 = 0.1$ (dashed line), $x_2 = 0.3$ (dotted line) and $x_2 = 0.95$ (dash-dotted line) as functions of $|\tilde{Y}_{22}^{\nu}|$ (panel a) and $\tilde{m}_1$ (panel b). In agreement with our discussion, for $x \lesssim 0.3$, when the lepton number asymmetry arises entirely from the lightest neutrino decays, the limits on $M_1$ depend only on $x_2 |\tilde{Y}_{22}^{\nu}|$ that is on $\tilde{m}_1$ (the three curves corresponding to $x_2 = 0.01$, 0.1 and 0.3 merge into a single one when plotted against $\tilde{m}_1$) and not on $x_2$ and $|\tilde{Y}_{22}^{\nu}|$ separately. It is worth stressing that in the range of $\tilde{m}_1$ values accessible in the considered limit, the prediction for the baryon asymmetry of the Universe is almost insensitive to the initial abundance of the first right-chiral neutrinos [17, 15]. One should also note that our results are consistent and generalize those obtained in ref. [11] in the specific see-saw model with exactly two right-chiral neutrinos.

A non-trivial dependence on $x_2$ and $\tilde{Y}_{22}^{\nu}$ separately can arise only through direct contribution of the second heavy neutrino to the generated lepton asymmetry. This is shown in fig. 3a, where we plot the lower limit on $M_1$ for the case of degenerate two right-chiral neutrinos. The dependence of this lower limit on the degeneracy of the two right-chiral neutrinos is shown in fig. 3b. In agreement with the observation made in ref. [19], the degeneracy of two neutrinos

\footnote{For $|\tilde{Y}_{22}^{\nu}| \sim 0.1$ the curve corresponding to $x = 0.9999$ underestimates slightly the lower limit on $M_1$ because there the applicability condition $\Gamma_1 \ll M_2 - M_1$ for the formulae for $\epsilon_{1,2}$ is not satisfied and one should use more general expressions derived in [18].}
allows to significantly lower the reheat temperature. Linear dependence of $M_1^{\text{min}}$ on the degeneracy parameter $\delta = 1 - M_1/M_2 = 1 - x_2^2$ seen in fig. 3b for $\delta < 0.05$ is related to the fact that for degenerate two right-chiral neutrinos both, $\epsilon_1$ and $\epsilon_2$, are proportional to $M_1/\delta$ \[16\]. From fig. 3 it follows that in this framework lowering $M_1$ below $10^9 - 10^8$ GeV can be achieved if $\delta \lesssim 10^{-3} - 10^{-4}$. The stability of such a degenerate mass pattern is however an open question. Another possibility is to consider a non-thermal production of right-chiral neutrinos and — in the supersymmetric scenario — sneutrinos \[21\]. A considerably lower reheat temperature is also sufficient in the very interesting scenario proposed in \[22\] and elaborated recently in \[23\], in which the $B - L$ asymmetry is produced directly by the decays of right-chiral sneutrinos playing the role of the inflaton.

The results described above apply for $x_3 \tilde{Y}_{\nu}^{32} \to 0$, $x_3 \tilde{Y}_{\nu}^{33} \to 0$. Solving explicitly the set of equations (4) with $C = 1, 2, 3$ in terms of $\tilde{Y}_{\nu}^{22}$, $\tilde{Y}_{\nu}^{32}$ and $\tilde{Y}_{\nu}^{33}$ it is easy to see (we have also checked it numerically) that as long as

$$x_3|\tilde{Y}_{\nu}^{3A}| < \mathcal{O}(0.1) \min \left\{x_2|\tilde{Y}_{\nu}^{22}|, \xi_A\right\}, \quad \text{for} \quad A = 2, 3 \quad (10)$$

the resulting values of $\epsilon_1$ and $\tilde{m}_1$ and, hence, the predictions for the $B - L$ abundance are almost unchanged. In concrete theoretical models of the see-saw mechanism these conditions can be satisfied either if the independent couplings $\tilde{Y}_{\nu}^{32}$ and $\tilde{Y}_{\nu}^{33}$ are sufficiently weak (because $\tilde{Y}_{\nu}^{3A} = \sum_B \mathbf{Y}_{\nu}^{3B} \mathbf{U}^{B,A}$ and taking into account the bi-large structure of the matrix $\mathbf{U}$, this condition is automatically satisfied if the couplings $\mathbf{Y}_{\nu}^{3A}$ are small) or, more naturally, for the third right-chiral neutrino sufficiently heavy (typically in such models $M_3 \gtrsim 10^{16}$ GeV).
In the regime \( x_2 \lesssim 0.3 \) when only decays of the first right-chiral neutrino contribute to the \( B - L \) asymmetry, the observed baryon asymmetry can be generated for \( M_1 \) as low as \( \lesssim 10^9 \text{ GeV} \) if the conditions (10) are strongly violated. To see this explicitly, we note that this is possible only if \( \tilde{m}_1 \) can be made substantially smaller than \( m_{\nu_2} \) keeping at the same time not too small \( |\epsilon_1| \). From the formulae (7), (9) and the orthogonality relation \( \sum_B \Omega^2_{1B} = 1 \) it follows that this requires \( \Omega^2_{11} \approx 1 \) and, for \( A = 2, 3, \) \( \Omega^2_{1A} \sim O(m_{\nu_1}/m_{\nu_A}) \) with nonnegligible imaginary part. The orthogonality relation applied to the first column of \( \Omega \) then implies that \( \Omega^2_{31} \) is at most of order \( O(m_{\nu_1}/m_{\nu_3}) \).

(11) such a suppression of the element \( \Omega^2_{31} \) requires \( \tilde{Y}^2_{32} \) and/or \( \tilde{Y}^2_{33} \) violating the conditions (10). At least one of the couplings \( (\tilde{Y}^A_{\nu})^2 \) must be then of order

\[
(\tilde{Y}^A_{\nu})^2 \approx 3.3 \left( \frac{m_{\nu_A}}{\text{eV}} \right) \left( \frac{M_3}{10^{14} \text{ GeV}} \right) \quad A = 2 \text{ or } A = 3,
\]

(12) that is, it approaches the nonperturbative regime for \( M_3 \gtrsim 10^{16} \text{ GeV} \).

It is also interesting to note that in the decoupling limit the requirement of perturbativity of the neutrino Yukawa couplings imposes some additional constraints on the neutrino masses. From the estimate \( \xi^2_A = 3.3 \times 10^{-4}(m_{\nu_A}/\text{eV})(M_1/10^{10} \text{ GeV}) \) and (6) one gets:

\[
(\tilde{Y}^{A\nu})^2 \approx 3.3 \left( \frac{m_{\nu_1}}{\text{meV}} \right) \left( \frac{M_3}{10^{17} \text{ GeV}} \right)
\]

(13) Therefore, for \( m_{\nu_1} \approx 10^{-3} \text{ eV} \) the mass of the heaviest right-chiral neutrino cannot exceed \( 10^{17} \text{ GeV} \). Similarly, for \( M_2 \gg M_1 \), that is for \( x_2 \ll 1 \), when \( |\Omega_{23}| \approx 1 \), the requirement of perturbativity of \( \tilde{Y}_{23} \) leads to the bound \( M_2 \lesssim 10^{15} \text{ GeV} \).

In conclusion, if one of the right-chiral neutrinos effectively decouples from the see-saw mechanism the resulting spectrum of the left-chiral light neutrinos must be hierarchical (or inversely hierarchical). Then, if the baryon number asymmetry of the Universe is to be generated by the decays of the lightest right-chiral neutrino only, its mass has to be higher than \( \sim 2 \times 10^{11} \text{ GeV} \) irrespectively of its initial abundance. Therefore, if the decaying neutrinos are produced thermally the required reheat temperature \( T_{\text{reh}} \) has to be \( \gtrsim 10^{12} \text{ GeV} \). The decoupling limit is practically realized in a broad class of theoretical see-saw models invoked to reconcile the hierarchical spectrum of light neutrinos with their bi-large mixing pattern. Supersymmetric versions of such models face therefore the serious problem of gravitino overproduction if the right-chiral neutrinos are produced thermally.

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