

Leptogenesis with GeV-Scale Right-Handed Neutrinos

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Abstract. We review the relation between leptogenesis and the discrete symmetry of charge-parity conjugation. The requirement of respecting the theorem of combined charge-parity-time reversal invariance at the level of the kinetic equations describing leptogenesis in the early Universe poses an interesting challenge that may most efficiently be addressed by the use of closed-time-path techniques. A byproduct of these methods is an accurate and unified description of leptogenesis from oscillations of right-handed neutrinos that applies to the regime of ultraheavy as well as GeV-scale RHNs. For the latter scenario, we discuss the mechanism on the example of the oscillatory and overdamped parametric regimes and very briefly comment on the prospect of experimental tests.

1. Introduction

The 2016 edition of the DISCRETE symposium in Warsaw takes place fifty years after Andrei D. Sakharov famously noted that the cosmic matter-antimatter asymmetry can be explained through the violation of discrete symmetries in conjunction with non-equilibrium thermodynamics [1]. The since sought-after, but yet unknown mechanism of creating the asymmetry has been called *baryogenesis*, and necessary ingredients for this process are often paraphrased as the celebrated *Sakharov conditions*:

- baryon number B violation,
- charge C and charge-parity CP violation,
- interactions out of thermal equilibrium.

Baryon-number violation is necessary to create $B \neq 0$ starting from initial conditions with $B \equiv 0$, C violation for the Universe to treat particles different from antiparticles and CP violation to avoid unwanted cancellations between fermionic degrees of freedom of left and right chirality.

The out-of-equilibrium condition is perhaps the most defining for the field of baryogenesis, and it is useful to look at it from different angles. To name two of these, one may note that an equilibrium state with an equal number of particles and antiparticles maximizes the entropy, while a state with more particles than antiparticles has a smaller entropy. The reduction of entropy due to a baryon asymmetry must therefore be compensated or overcompensated by an increase in entropy carried by other degrees of freedom, which can be achieved by a thermodynamically irreversible change of the state variables, *e.g.* through the expansion and



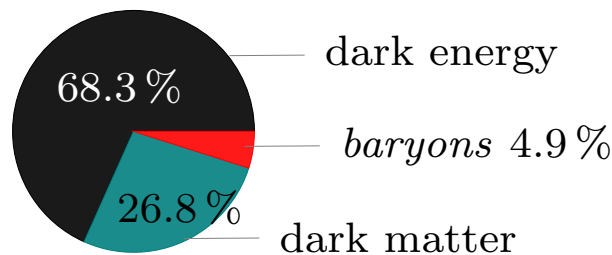


Figure 1. Contributions to the energy density of the Universe.

cooling of the Universe. Alternatively, one may argue that due to the combined invariance of quantum field theory under charge, parity and time reversal (CPT), an equilibrium state (which is time-reversal invariant) cannot support a non-vanishing (*i.e.* C and CP violating) baryon number.

By the year 1966, a number of theoretical and experimental insights have had come together setting the initial point for the problem of baryogenesis: the formulation of the Dirac equation [2] together with the discovery of the positron [3], the CPT -invariance theorem [4, 5], the observation of parity [6] and CP violation [7] in weak interactions and finally the detection of the cosmic microwave background (CMB) as the smoking gun of the hot big bang [8]. Modern surveys of the CMB also give the best measurement of the asymmetry [9] (*cf.* Figure 1), which is in excellent agreement with big bang nucleosynthesis (BBN) [10]. The accurate understanding that has been achieved of the CMB and of BBN is a stunning success of applying particle physics well beyond laboratory experiments to the early Universe, and it motivates to tackle unresolved problems such as dark matter and baryogenesis.

2. Non-Equilibrium and Neutrinos

The non-equilibrium conditions may be realized in a large number of ways. To avoid reliance on peculiar initial conditions, there is a certain appeal to scenarios where the Universe is initially very close to equilibrium and a more pronounced deviation occurs subsequently during the expansion process. According to these rules of the game, baryogenesis from first-order phase transitions, such as electroweak baryogenesis, and baryogenesis from out-of-equilibrium decays of non-relativistic particles, such as ultraheavy sterile neutrinos, are of principal interest. We also find it interesting to relax above rule a little and to allow for a deviation of equilibrium due to very weakly coupled degrees of freedom that are not occupied initially and only populated at a later time due to subsequent equilibration. An example for this latter category is leptogenesis from GeV-scale sterile neutrinos. In this section, we discuss the parametric hints that point to an involvement of sterile neutrinos in baryogenesis.

Key observables that have recently been undergoing rapid experimental improvement are the neutrino mass and mixing parameters that can be inferred from a combination of the solar neutrino and atmospheric neutrino fluxes as well as from oscillations of reactor neutrinos and in muon neutrino beams. We denote by m_i the mass of a light neutrino, where $i = 1, 2, 3$, $\Delta m_{ij}^2 = (m_i - m_j)^2$ and $m_1 < m_2 < m_3$ for the so-called normal hierarchy (NH), whereas $m_3 < m_1 < m_2$ for inverted hierarchy (IH). For the present discussion, we quote the mass differences $\Delta m_{21}^2 = 7.50 \times 10^{-5} \text{eV}^2$, $\Delta m_{31}^2 = 2.457 \times 10^{-3} \text{eV}^2$ for NH, (the best fit values for IH only differ by very little) summarized in Ref. [12] as well as the upper cosmological bound on the sum of the masses [9] $\sum_i m_i < 0.23 \text{eV}$.

These mass scales are way below those within the Standard Model (SM) and call for an explanation, where perhaps the most popular is the type I seesaw mechanism. It introduces n_N

hypothetical right-handed neutrinos (RHNs) N with a Majorana mass matrix M (which we can take to be diagonal through field redefinitions). The RHNs couple to SM leptons $\ell = (\nu, e_L)$ and Higgs doublets ϕ , where $\langle |\phi| \rangle = v = 174 \text{ GeV}$:

$$\mathcal{L}_{\text{SM}} \rightarrow \mathcal{L}_{\text{SM}} + \frac{1}{2} \bar{N}_i^c (i \not{\partial} - M_{ij}) N_j - Y_{ia}^* \bar{\ell}_a \phi^\dagger N_i - Y_{ia} \bar{N}_i \phi \ell_a, \quad (1)$$

with $a = e, \mu, \tau$. The RHNs and the SM leptons mix through the Yukawa couplings Y , and their masses can be arranged in terms of a $(3 + n_N) \times (3 + n_N)$ matrix:

$$\frac{1}{2} (\bar{\nu} \bar{N}^c) \underbrace{\begin{pmatrix} 0 & m_D \\ m_D & M \end{pmatrix}}_{=: \mathcal{M}} \begin{pmatrix} \nu^c \\ N \end{pmatrix}, \quad (2)$$

where $m_D = Y^\dagger v$. Assuming, $\|M\| \gg \|m\|$, it is easy to block diagonalize

$$\begin{pmatrix} m & 0 \\ 0 & M \end{pmatrix} = U \mathcal{M} U^\dagger, \quad U = \begin{pmatrix} 1 - \theta \theta^\dagger & \theta \\ -\theta^\dagger & 1 - \theta^\dagger \theta \end{pmatrix}, \quad \theta = Y^\dagger v M^{-1}. \quad (3)$$

This leads to a Majorana mass matrix for *light* neutrinos $m = v^2 Y^\dagger M^{-1} Y^*$ (with the eigenvalues m_i), diagonalized by U_ν , as well as one for *heavy* neutrinos $M_N = M + \frac{1}{2} (\theta^\dagger \theta M + M \theta^T \theta^*) + \mathcal{O}(\theta^3)$, diagonalized by U_N , where

$$\nu^{\text{light}} = U_\nu^\dagger \left(\left(\mathbb{1} - \frac{1}{2} \theta \theta^\dagger \right) \nu - \theta N^c \right), \quad (4a)$$

$$\nu^{\text{heavy}} = U_N^\dagger \left(\left(\mathbb{1} - \frac{1}{2} \theta^T \theta^* \right) N - \theta^T \nu^c \right). \quad (4b)$$

The RHNs are their own antiparticles and can decay either $N_i \rightarrow \ell \phi$ or $N_i \rightarrow \bar{\ell} \phi^\dagger$ and therefore violate lepton number L .

3. Out-of Equilibrium Decays, Asymmetries and Real Intermediate States

Now due to the approximate conformal invariance, a gas of thermalized and relativistic (*i.e.* of mass much below the temperature T) particles remains in equilibrium as the Universe expands. Another way of stating this is that the distribution function of a non-interacting massless particle in an expanding box remains of the Bose-Einstein or Fermi-Dirac form as the volume is changed, while this is not true for a massive species. Therefore, a deviation from thermal equilibrium will occur when the mass of a particular species becomes appreciable compared to the temperature. As we discuss below, ideally, the temperature when equilibration occurs coincides with the temperature where the particle becomes non-relativistic.

While in Sakharov's seminal work, the heavy particle responsible for the asymmetry bears the name "Maximon", in the arena of leptogenesis, the RHN is the gladiator [11]. The simplest meaningful network of kinetic equations describing leptogenesis may be stated as

$$\frac{dn_{Ni}}{dt} + 3Hn_{Ni} = \Gamma_i(n_{Ni} - n_{Ni}^{\text{eq}}), \quad (5a)$$

$$\frac{dn_L}{dt} + 3Hn_L = \varepsilon \Gamma_i(n_{Ni} - n_{Ni}^{\text{eq}}) - Wn_L, \quad (5b)$$

where

$$\Gamma_i = \frac{|Y_i|^2}{8\pi} M_{ii}, \quad W = \frac{\Gamma_i}{4} \left(\frac{M_{ii}}{T} \right)^{\frac{3}{2}} e^{-\frac{M_{ii}}{T}}$$

are the decay rate of RHN N_i and the washout rate, respectively. Further, we have introduced here the Hubble expansion rate H , the number density n_X of a species X (L denoting leptons) and the decay asymmetry

$$\varepsilon = (\Gamma_{Ni \rightarrow \ell H} - \Gamma_{Ni \rightarrow \bar{\ell} H^*}) / (\Gamma_{Ni \rightarrow \ell H} + \Gamma_{Ni \rightarrow \bar{\ell} H^*}). \quad (6)$$

(Γ_i is here taken to be the tree-level, CP -symmetric decay rate, while $\Gamma_{Ni \rightarrow \ell H}$ and $\Gamma_{Ni \rightarrow \bar{\ell} H^*}$ include CP violating contributions that typically arise at next-to-leading order.)

The coincidence between the equilibration temperature and the transition to the non-relativistic regime occurs when $\Gamma_i \sim H$ for $T \sim M_{ii}$, *i.e.* if $[YY^\dagger]_{ii}/(8\pi) \sim T^2/m_{\text{Pl}} \sim M_{ii}^2/m_{\text{Pl}}$. For smaller interaction strengths Γ_i , the processes that produce the lepton asymmetry n_L in Eq. (5b) will be suppressed, whereas for larger interactions, the system will be maintained close to equilibrium due to the right-hand side of Eq. (5a), what suppresses the outcome, such that above coincidence corresponds to ideal circumstances for creating a large lepton asymmetry. For simplicity, we now drop the flavour indices, such that the Majorana mass for light neutrinos reads $m \sim Y^2 v^2/M$. Using this to substitute the heavy neutrino mass M in the previous relations, we obtain that

$$m \sim 8\pi v^2/m_{\text{Pl}} \sim 0.1 \text{ meV}. \quad (7)$$

This is somewhat below the observed mass scale that we have quoted in Section 2 but yet remarkably close given the fact that the masses of the known fermions span over many orders of magnitude. Moreover, ignoring the flavour structure, this hints to Γ being moderately larger than H for $M \sim T$, implying that the RHNs are moderately close to equilibrium when they become non-relativistic. This situation is referred to as *strong washout*. To summarize this point, intriguingly, the neutrino mass scale points to the possibility that the RHNs of the type-I seesaw mechanism may play a key role in baryogenesis. Strong washout has the extra appeal that the outcome of leptogenesis is largely independent of the initial conditions, and there are fairly accurate analytic methods of solving Eqs.(10) in that regime [13, 14].

The decay asymmetry ε can be calculated by standard means from interfering tree-level with loop amplitudes for the process $N_i \rightarrow \ell H$ as well as for its CP conjugate. These interferences can be represented by the blue cuts in Figure 2. As stated above, we denote the CP symmetric tree-level result for the decay rate by Γ_i . Then,

$$\Gamma_{Ni \rightarrow \ell \phi} = \Gamma_i(1 + \varepsilon) = \Gamma_{\bar{\ell} \phi^* \rightarrow Ni}, \quad (8)$$

where the latter equality is a direct consequence of the CPT theorem. Likewise, the charge-parity conjugate processes are related through conjugation of the couplings,

$$\Gamma_{Ni \rightarrow \bar{\ell} \phi^*} = \Gamma_i(1 - \varepsilon) = \Gamma_{\ell \phi \rightarrow Ni}. \quad (9)$$

These relations suggest that if, for definiteness, $\varepsilon > 0$, then processes generating a positive asymmetry, *i.e.* those that create leptons or destroy antileptons, are faster than those that generate a negative asymmetry. Substituting this without further ado as rates into Boltzmann equations, this would neither be consistent with Eq. (5b) (where a process creating leptons of a definite charge and its inverse have the same rate) nor with the CPT theorem, that bars the creation of an asymmetry in equilibrium.

This difficulty is often addressed by including lepton number violating $2 \leftrightarrow 2$ processes, *i.e.* $\ell H \leftrightarrow \bar{\ell} H^*$ and by subtracting from these *partly* the contributions from the on-shell exchange of RHNs, namely precisely in a manner that complies with the CPT theorem, arguing that these contributions are already included within the $1 \leftrightarrow 2$ processes [15]. This procedure is known as real intermediate state (RIS) subtraction.

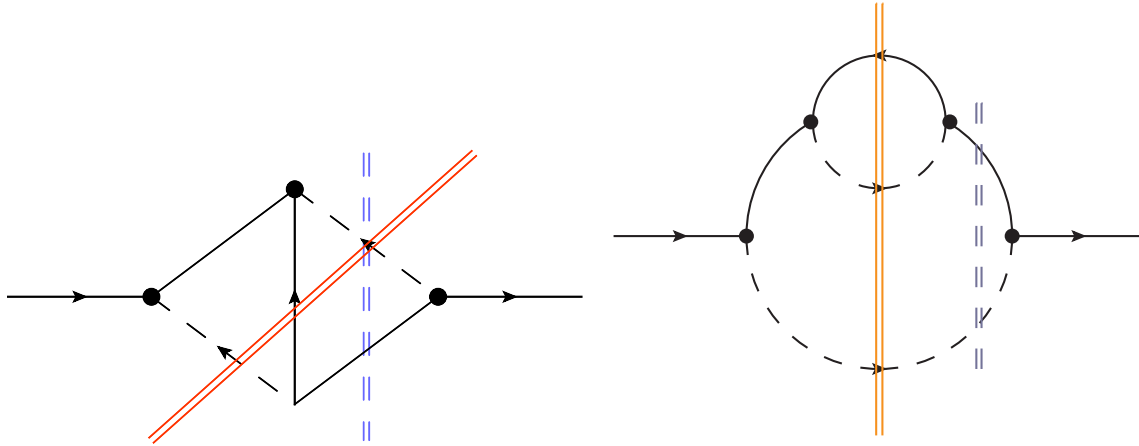


Figure 2. One loop contributions to the asymmetry in decays and inverse decays. RHNs are represented by black solid lines without arrows, leptons by solid lines with arrows and Higgs bosons by dashed lines with arrows. The double lines represent cuts, where the dashed one leads to interferences between tree and loop $1 \leftrightarrow 2$ processes, while the solid ones to interferences between scattering processes. In those contributions relevant for generating a CP asymmetry, scattering proceeds via an on-shell RHN.

Instead of substituting S -matrix elements into Boltzmann equations, a complementary and somewhat more systematic and direct way of addressing this problem is to directly solve for the time evolution of the operators of interest, *e.g.* the lepton charge density. The unitary time evolution follows from the Hamiltonian but a perturbative expansion becomes rather involved quickly. The task is greatly simplified in a functional approach, where the perturbation expansion can be organized within simple Feynman diagrammatic rules. This calculus is known as the Schwinger-Keldysh or Closed-Time-Path (CTP) formalism [16, 17]. The key difference to scattering problems is that for S -matrix elements, one does calculations in terms of time-ordered Green functions and in particular propagators, while in the CTP framework, diagrams are constructed from a more general basis of propagators (for example time-ordered Feynman propagators and Wightman functions without time ordering) that contain both spectral and statistical information.

An introduction to the derivation of kinetic equations from the CTP formalism in relativistic field theory is given in Ref. [18]. The kinetic equations descend from Schwinger-Dyson equations, and the various scattering processes appear in the self-energy terms. By construction, these processes are counted *inclusively*, such that the CPT theorem is respected. The propagators pertaining to different boundary conditions that appear in the CTP approach can be associated with certain cuts, that in turn lead to the CP -violating contribution. This is illustrated in Figure 2: Besides the interferences of tree and loop amplitudes, we encounter also interferences between $2 \leftrightarrow 2$ scattering amplitudes. The CP -violating contributions from the latter interferences come precisely from the region of momentum integration where one of the RHNs propagates on shell. This is shown for leptogenesis in Refs. [19, 20], and for a model with scalar particles in Refs. [21, 22]. The CP cuts that appear in fully inclusive calculations on soft leptogenesis, where substantial cancellations occur, are discussed in Ref. [23], where also numerous calculational details on handling CP asymmetries in the CTP approach are presented.

4. Resonant Leptogenesis and from Oscillations of GeV-Scale Sterile Neutrinos

Now suppose that the width $k^0\Gamma_i$ is larger or of the same order than the squared mass splitting $(M_{ii} - M_{jj})^2$, where k^0 is the typical energy of the RHN, *i.e.* $k^0 \approx M_{ii}$ in the non-relativistic case whereas $k^0 \sim T$ in the relativistic case. For the mass splitting to be comparable to the width, a mass degeneracy is therefore required in the non-relativistic regime, what leads to the typical scenarios of resonant leptogenesis [24, 25, 26, 27]. (Note here that even when $(M_{ii} - M_{jj})^2 \ll M_{ii}^2$ it is still possible that $(M_{ii} - M_{jj})^2 \gg k^0\Gamma_i$ such that no resummation of the RHN propagators is necessary). In the relativistic regime, a mass degeneracy is not necessary. This can typically lead to successful leptogenesis with RHNs around the GeV scale [28, 29]. (For specific configurations of the flavoured couplings in the seesaw mechanism, also scenarios with RHNs above the electroweak scale are viable [30].)

Clearly, when the squared mass splitting is smaller than the width, the expansion of the RHN propagator through the insertion of a wave function correction, as indicated in Figure 2, cannot be performed. Instead, the wave function corrections need to be resummed, what is automatically achieved in the CTP formulation of the Schwinger-Dyson equations. To illustrate this point explicitly, we note the kinetic equation for the deviation $\delta f_{Nh}(\mathbf{k})$ of the RHN distribution from equilibrium [32, 33],

$$\delta f'_{Nh} + \frac{a^2(\eta)}{2k^0} i[M^2, \delta f_{Nh}] + f_N^{\text{eq}'} = -2 \left\{ \text{Re}[Y^* Y^t] \frac{k \cdot \hat{\Sigma}_N^A}{k^0} - i h \text{Im}[Y^* Y^t] \frac{\tilde{k} \cdot \hat{\Sigma}_N^A}{k^0}, \delta f_{Nh} \right\}. \quad (10)$$

This follows from the Schwinger-Dyson equations (or, more specifically to the CTP framework, the Kadanoff-Baym equations) from integrating out k^0 (*i.e.* taking the zeroth moment) and taking an appropriate trace over the spinor structure. Furthermore, in above equation, $\hat{\Sigma}_N^A$ is the reduced (*i.e.* stripped of Yukawa couplings that appear explicitly in the equation) spectral (cut part) of the self energy, $a(\eta)$, η are the scale factor of the expanding Universe and the conformal time ($\tau \equiv d/d\eta$), h denotes helicity, and $\tilde{k} = (|\mathbf{k}|, |k^0|\mathbf{k}/|\mathbf{k}|)$. We have suppressed the flavour indices of δf_{Nijh} and of the heavy neutrino mass matrix M_{ij} as well as on the Yukawa couplings Y_{ia} . This equation implicitly resums all wave-function insertions of $\hat{\Sigma}_N^A$. The off-diagonal components of δf_{Nijh} are crucial because these also appear in the kinetic equation for the left-handed SM leptons, where they lead to the creation of the CP asymmetry.

For non-relativistic neutrinos, one often encounters the situation that one can neglect the term $\delta f'_{Nh}$ because all eigenvalues of this ordinary and linear differential equation that follow from the commutator and anticommutator terms are larger than the rate of change $f_N^{\text{eq}'} / f_N^{\text{eq}}$ due to the Hubble expansion. (f_N^{eq} denotes the equilibrium Fermi-Dirac distribution of the RHNs.) In this case (which can include situations where $(M_{ii} - M_{jj})^2 \ll k^0\Gamma_i$), one can derive a quasi-static solution for $\delta f'_{Nh}$ that is more general than the well-known results but that agrees with these when $(M_{ii} - M_{jj})^2 \gg k^0\Gamma_i$ [33, 34]. In particular, these solutions apply throughout the strong washout regime. In case there are eigenvalues that are smaller than $f_N^{\text{eq}'} / f_N^{\text{eq}}$, which may happen outside the strong washout regime, it is necessary to track the full time-dependence of δf_{Nh} . In summary, making use of Eq. (10), it is possible to make accurate predictions for resonant leptogenesis for all ratios between $(M_{ii} - M_{jj})^2$ and $k^0\Gamma_i$. (We disagree with Ref. [40] that argues that there are additional contributions from the “mixing” of RHNs, that would result in a factor two enhancement of the asymmetry for typical parameter points within the strong washout regime.)

Now for GeV-scale RHNs, that are relativistic at the relevant temperatures above the electroweak phase transition (where baryon-minus-lepton number violating sphaleron transitions freeze out) the production rate of the individual RHNs may or may not be below the Hubble rate. While there exist different suitable analytic approximations for various regions in parameter space, the situation is more involved compared to the strong washout regime in standard resonant

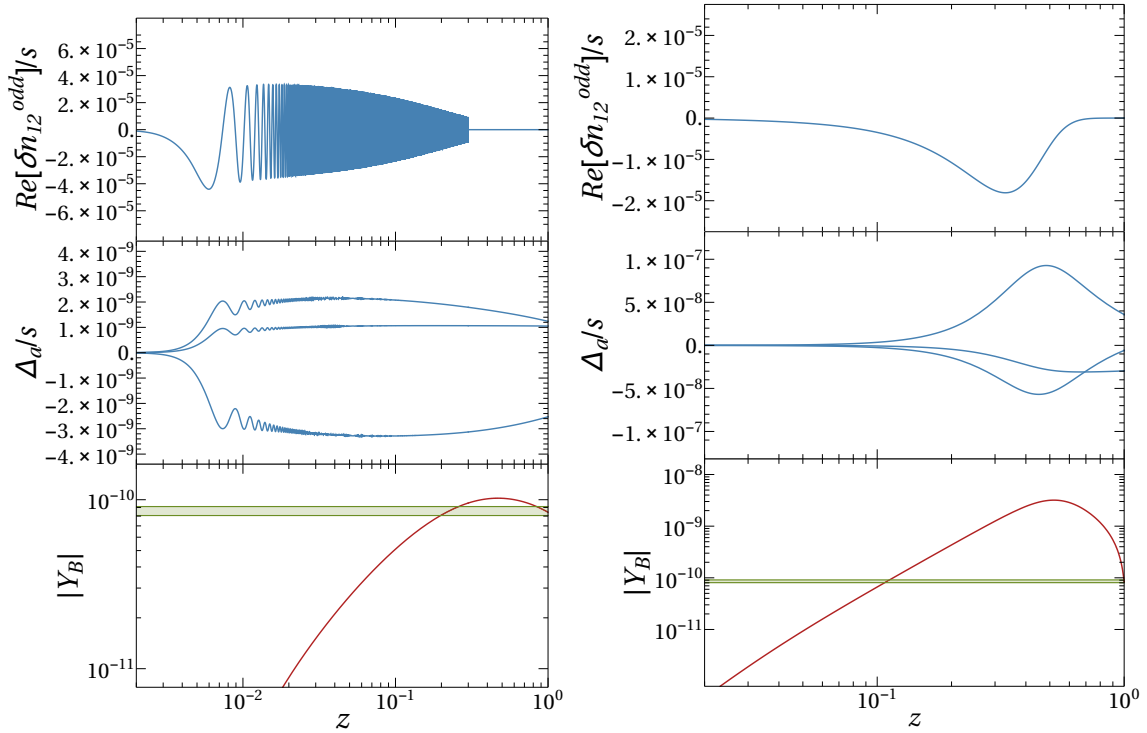


Figure 3. Examples for the oscillatory (left) and the overdamped regime of leptogenesis from GeV-scale RHNs. The helicity-odd number density of RHNs (*i.e.* the phase-space integral of $\delta f_{N+\frac{1}{2}} - \delta f_{N-\frac{1}{2}}$) is denoted by δn , $\Delta_a = B/3 - L_a$ (B being the baryon number density and L_a the number density of SM leptons of flavour $a = e, \mu, \tau$), $Y_B = B/s$, s the entropy density and $z = T_{EW}/T$ with T_{EW} the temperature of the electroweak phase transition. The green band indicates the range for the baryon asymmetry that is in agreement with current observations.

leptogenesis. To illustrate the different parametric regions, we present in Figure 3 the *oscillatory* and the *overdamped* regimes [35].

In the oscillatory regime, $|M_{ii} - M_{jj}| \gg T_{EW}$ and $(M_{ii} - M_{jj})^2 \gg TT_i$, such that many oscillations can take place before the electroweak phase transition. The Yukawa couplings are small compared to the overdamped regime, such that the baryon asymmetry only suffers a small amount of washout prior to the electroweak phase transition. In the overdamped regime $(M_{ii} - M_{jj})^2 \gg TT_i$ does not hold (for at least one of the N_i), *i.e.* there is a smaller mass splitting and larger Yukawa couplings. Interestingly, this does not necessarily have to correspond to fine tuning but can also be motivated by an approximate conservation of lepton number [36, 37, 38, 39]. We can observe that larger Yukawa couplings also lead to a stronger washout of the asymmetry toward the electroweak phase transition.

What can be seen from both panels of Figure 3 is that at early times, the baryon asymmetry Y_B is very small. This is because the asymmetry generated initially is purely flavoured, such that summing the asymmetries over the active lepton flavours yields a vanishing result. Only at a later time toward the electroweak phase transition, an appreciable amount of asymmetry is transferred to the RHNs (where it takes the form of a helicity asymmetry), such that a non-vanishing net lepton number is present in the SM leptons. Sphaleron processes then lead to a non-vanishing baryon asymmetry Y_B .

As direct and indirect signals of RHNs are proportional to these quantities, the active-sterile mixing angles θ , Eq. (3), given present and anticipated constraints from the PMNS matrix and

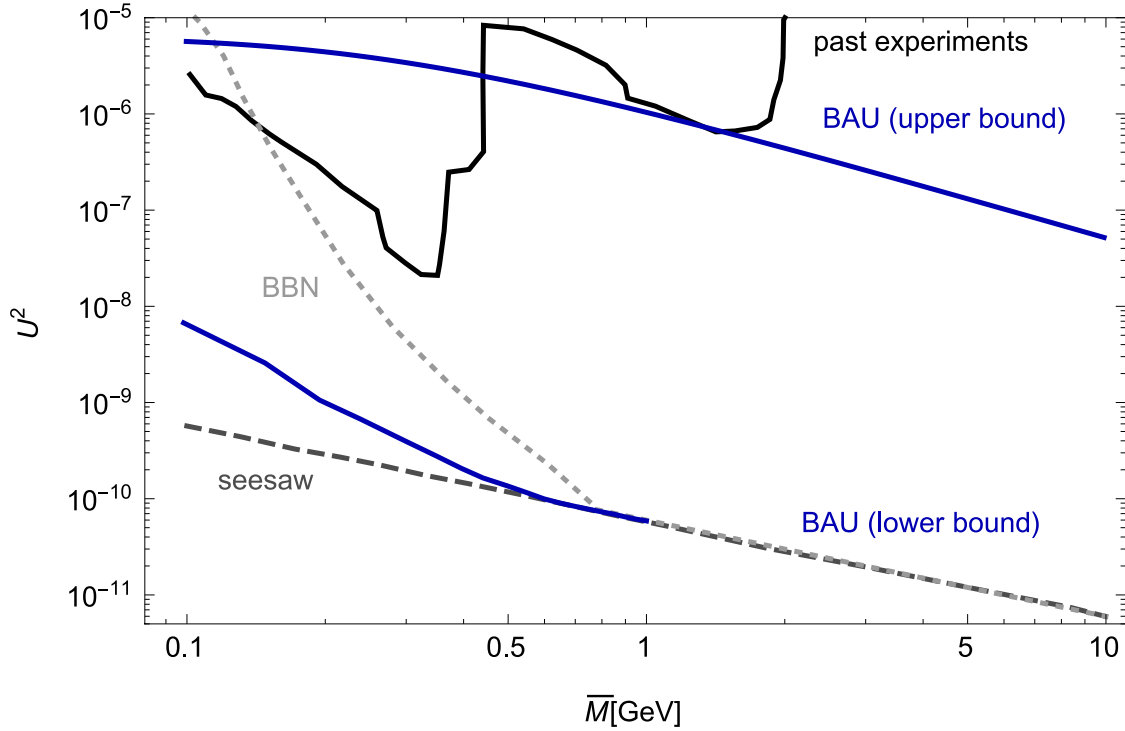


Figure 4. Between the two solid blue lines, there is the allowed range of the total active sterile mixing U^2 over the average mass $(M_1 + M_2)/2$ in a scenario with $n_N = 2$ RHNs, assuming normal hierarchy. There is a lower bound on the mixing from the requirement of explaining the light neutrino masses through the seesaw mechanism, and an additional one from big bang nucleosynthesis (BBN).

the light neutrino masses, are of phenomenological relevance. Recent work has discussed this topic, also concerning the different flavour effects in some detail [35, 41, 42, 43, 44]. For brevity, we here just present in Figure 4 the flavour-summed quantity $U^2 = \text{tr} \theta^\dagger \theta$ that characterizes the total active-sterile mixing. Along the upper limit, leptogenesis is viable in the overdamped regime that allows for large mixing angles whereas the lower bound occurs at the transition between both regimes, where the baryon asymmetry turns out maximal for a given mixing angle (through tuning the mass degeneracy).

It can be anticipated that the experimental bounds will improve considerably in the future. Relevant in the lower mass range are NA62 and the proposed SHiP facility, whereas heavier RHNs may be particularly well accessible at future lepton colliders (see *e.g.* [45, 46, 47, 48]).

5. Conclusions

Closed-time-path techniques are the most powerful tool to meet the requirements of the *CPT* theorem when formulating kinetic equations for the dynamical emergence of the matter-antimatter asymmetry in the early Universe. In particular, they lead to a unified formulation of *CP* violation from mixing of RHNs throughout the entire parameter space of the type-I seesaw mechanism. This includes resonant leptogenesis from ultraheavy RHNs as well as leptogenesis from GeV-scale RHNs. In the latter case, it is possible that RHNs responsible for leptogenesis are accessible by present and proposed experiments. As this requires large active-sterile mixing, leptogenesis would then likely proceed in a regime of overdamped oscillations.

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