LEPTOGENESIS IN U(1) EXTENSIONS OF THE STANDARD MODEL*

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We present the ingredients of a computation needed for a semi-classical treatment for estimating the amount of baryogenesis via leptogenesis. Our main focus is on the computation of the finite temperature CP-asymmetry factor originating from Majorana fermion decays into a lepton and a scalar particle at next-to-leading order in perturbation theory. Such decays emerge naturally in U(1) extensions of the Standard Model, such as the superweak extension (SWSM). We emphasize the importance of all cuts in the one-loop corrections in order to obtain physically meaningful expressions and present benchmark predictions for the CP-asymmetry factor as a function of temperature in the SWSM.

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1. Introduction

The Standard Model (SM) of particle physics correctly describes the scattering processes at high-energy colliders [1] and is consistent with most of the low-energy measurements, too. There are, however, well-established observations in the cosmic and intensity frontiers that cannot be explained within the Standard Model [2]. In this contribution, we focus on the origin of baryon asymmetry, which is best measured in the Cosmic Microwave Background power spectrum and cannot be estimated correctly in the SM [3].

Evidence is accumulating that CP-violation in the lepton sector can be much larger than that in the quark sector [4], implying that leptogenesis may provide an explanation for the observed baryon asymmetry — called lepto-baryogenesis — which requires an epoch in the evolution of the Universe with Majorana-type fermions (such as heavy right-handed neutrinos, RHNs) with complex Yukawa couplings and sphaleron conversion [5] allowed. The RHNs may gain mass from a scalar vacuum expectation value (VEV) coupled to

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RHNs through the Majorana-type Yukawa term in the Lagrangian, while the sphaleron process becomes negligible when the sphaleron rate drops below the Hubble rate near the electroweak phase transition [6]. In this contribution, we study thermal leptogenesis [7] in such models beyond the SM (BSM).

2. Lepto-baryogenesis

Electroweak sphaleron processes violate baryon (B) and lepton (L) number, but conserve $B-L$ [8]. These processes redistribute asymmetries among fermions and are suppressed exponentially with decreasing temperature, but unsuppressed at $T \gtrsim 132$ GeV [6].

The standard way to estimate baryogenesis quantitatively is: *(i)* to generate a non-vanishing lepton asymmetry ΔL via lepton number violating processes (see Ref. [9] for a review), and then *(ii)* use the sphaleron conversion to generate a non-vanishing baryon asymmetry $\Delta B \propto \Delta L$. Here, we discuss how to estimate ΔL in BSM theories with RHNs.

The neutrinos in the SM are massless, hence the observation of non-vanishing neutrino masses provides a clear sign for BSM physics. Adding RHNs to the particle spectrum is a simple way to generate the light neutrino masses. The RHN decays can lead to a lepton asymmetry, which can be estimated either in non-equilibrium quantum field theory based on the Kadanoff–Baym equations [10], as done in [11, 12], or by a system of Boltzmann equations that are semi-classical in the sense that they can be derived from the former ones employing the quasi-particle approximation. In the latter, however, double counting of contributions appearing in both subsequent decays and scatterings has to be remedied by employing real intermediate state (RIS) subtraction [13], which leads to collision terms respecting unitarity. Strictly speaking, the Boltzmann approach assumes the system being

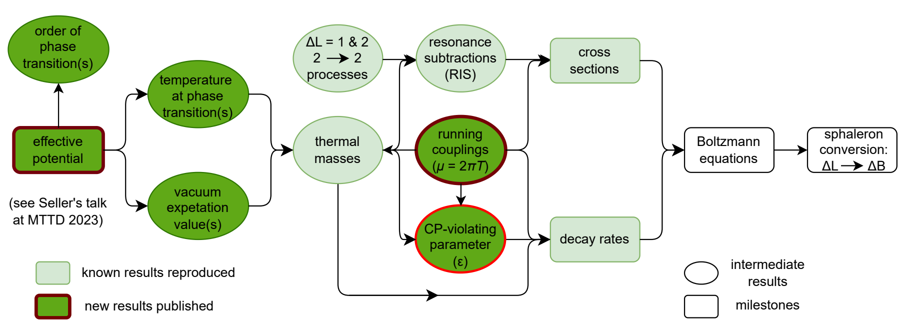

Fig. 1. Required ingredients for estimating the baryon asymmetry by the Boltzmann equations.

near statistical equilibrium. Yet, it is technically significantly simpler, being highly modular in the necessary ingredients as shown in Fig. 1, hence, it is better suited for exploratory estimates of measure of lepton asymmetry in BSM theories, as often done in the literature. (See Ref. [14] for a detailed calculation.)

3. Measure of lepton asymmetry in BSM theories with RHNs

The Boltzmann equation for the abundance of lepton asymmetry $\mathcal{Y}_{\Delta L} = \mathcal{Y}_\ell - \mathcal{Y}_{\bar\ell}$ (*i.e.* difference of the comoving charged lepton and anti-lepton number densities, with $\mathcal{Y}$ being equal to the ratio of number density to entropy density, $\mathcal{Y} = n/s$) reads as

$$\frac{\mathrm{d}\mathcal{Y}_{\Delta L}}{\mathrm{d}z} = \frac{1}{sHz}\left[\left(\epsilon\gamma_{\mathrm{D}} - \gamma_{ab\to N\ell}\frac{\mathcal{Y}_{\Delta L}}{\mathcal{Y}_\ell^{\mathrm{eq}}}\right)\left(\frac{\mathcal{Y}_N}{\mathcal{Y}_N^{\mathrm{eq}}} - 1\right) - W\mathcal{Y}_{\Delta L}\right] + \mathcal{O}\left(\epsilon^2\right) , \quad (1)$$

where $\mathcal{Y}_\ell^{\mathrm{eq}}$ is the equilibrium value of the lepton abundance (and similarly the quantities with index N for the RHNs). The dimensionless parameter $z = \Lambda/T$ is the inverse temperature measured in some arbitrarily chosen, unphysical energy scale Λ that does not affect physical predictions, $H(z)$ denotes the Hubble parameter at the time corresponding to the temperature Λ/z, and $s(z)$ represents the entropy density. The full rate equation and the relevant thermal rates $\gamma_{ab\to\mathrm{leptons}}(z)$ for the (model-dependent) processes $ab \to$ leptons will be precisely derived elsewhere. The neutrino density is subject to change due to decays and scattering processes. The asymmetry is generated by CP-violating decays of the sterile neutrinos, giving the term in the first line, and it is proportional to the CP-asymmetry factor ϵ, whose estimation is the focus of this presentation. We parametrize the CP-violating decays involving the RHNs using the CP-conserving tree-level decay rate γ_{D}. The term with W, which is a collection of terms emerging from the scattering processes, drives to equilibration, *i.e.* the washout of lepton asymmetry.

4. Estimating the CP-asymmetry factor

While the CP-asymmetry factor is still often used as a constant number obtained in $T = 0$ quantum field theory, it is clearly more appropriate to use $\epsilon(T)$ as a function of temperature. In Ref. [15], we presented a comprehensive study of the finite temperature CP-asymmetry factor needed in the semi-classical treatment of leptogenesis originating from Majorana fermion decays into a lepton and a scalar particle. The motivation for that study was to further improve the computation of Ref. [14] in two aspects. *(i)* To take into account the neglected two cuts[1] in the vertex diagram (out of the three in total). While those cuts are suppressed exponentially as $\exp(-m_N/T)$ for large mass m_N of the decaying heavy neutrino (as considered in Ref. [14]), the additional contribution of the neglected cuts (cuts 2 and 3 in Fig. 2) may be relevant for the low-scale leptogenesis when $m_N \approx T$ [16], which is the case in the superweak extension of the SM (SWSM) [17]. *(ii)* The expressions for the physically relevant thermal self-energy and vertex function in Ref. [14] are quadratic in statistical factors, but should be linear at one-loop accuracy, which had called for a correction. We mention, however, that the correct expression of $\epsilon(T)$ can be read off from the source term of lepton asymmetry obtained in the non-equilibrium formalism (*e.g.* the contribution of the self-energy and cut 1 can be extracted from Eq. (44) of [18]).

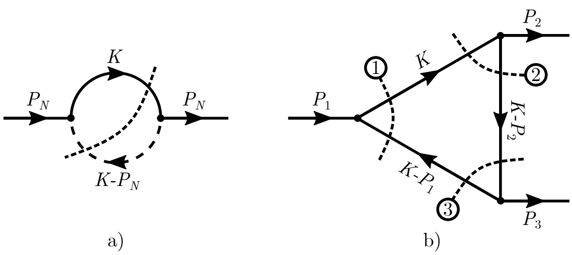

Fig. 2. One-loop cut diagrams contributing to the CP-asymmetry factor. (a) The cut neutrino self-energy (bubble) diagram with intermediate lepton and scalar fields in the loop. (b) The three cuts of the one-loop vertex correction (triangle) diagram. The particle types of the incoming and outgoing states as well as the internal lines, depend on the particular process considered, $N_i(P_1) \to \phi(P_2) + L(P_3)$ or $\phi(P_1) \to N_i(P_2) + L(P_3)$.

5. Superweak extension of the Standard Model

The BSM theories we have in mind involve the simplest extension with a $\mathrm{U}(1)_z$ gauge group — the SWSM being a particular example. U(1) extensions are the simplest ultraviolet complete extensions with sufficiently rich phenomenology to explain the known BSM phenomena. The neutral gauge boson belonging to the extra U(1) symmetry must be massive, otherwise it would mix with the photon. In order to acquire such a mass from spontaneous breaking of the U(1) symmetry by the vacuum, the existence of a second scalar is needed in addition to the SM scalar.

The superweak force is a minimal, anomaly-free U(1) extension of the Standard Model with such a complex scalar field and three families of RHNs. It is designed to explain the origin of: *(i)* neutrino masses and mixing matrix

[1] Using the optical theorem, from the cut diagrams we can derive the imaginary parts of the decay amplitudes, which contribute to the CP asymmetry factor [9].

elements, *(ii)* dark matter, *(iii)* cosmic inflation, *(iv)* stabilization of the electroweak vacuum, and *(v)* leptogenesis without altering the predictions of the SM beyond the experimental uncertainties of measurements at colliders.

In previous editions of this conference, we presented the model and some of its phenomenological consequences [19–22], which can be summarized as follows:

1. Dirac and Majorana neutrino mass terms are generated by the SSB of the scalar fields, providing the origin of neutrino masses and oscillations [24, 25].
2. The lightest new particle is a natural and viable candidate for WIMP dark matter if it is sufficiently stable [26].
3. The second scalar together with the established BEH field can stabilize the vacuum and be related to the accelerated expansion now and inflation in the early Universe [27, 28].
4. Diagonalization of neutrino mass terms leads to the PMNS matrix, which, in turn, can be the source of lepto-baryogenesis. This is being explored in ongoing research, and it is the subject of this contribution.

The smallness of the g_z coupling — as implied by the adjective "superweak" — in the SWSM ensures that the predictions do not conflict with the existing experimental results. In particular, the fruitless search for dark photons by the NA64 experiment [29] can be translated into a strong upper limit of $g_z \lesssim O(10^{-4})$. This small coupling also implies that the thermal mass of the RHN is typically smaller than its mass at $T = 0$, as shown in Fig. 3. In the shaded region at intermediate temperatures, the thermal masses are such

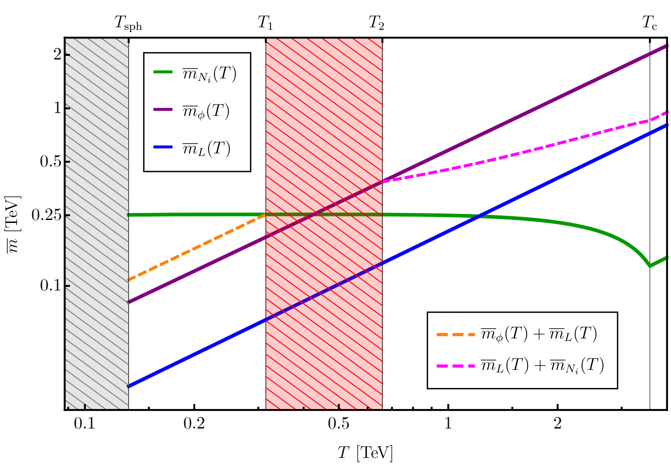

Fig. 3. Thermal masses for the lighter of the two heavy RHNs ($\bar{m}_{N_i}$), the lepton ($\bar{m}_L$), and the Brout–Englert–Higgs field ($\bar{m}_\phi$) in the SWSM at two specific values of the VEV ratio.

that decays involving sterile RHNs, leptons, and scalar fields are kinematically excluded. We indicate a low temperature cutoff at $T_{\rm sph} \simeq 132$ GeV, where the sphaleron processes freeze-out. In this example, the vacuum masses of the two heavy neutrinos are similar, $m_{N_j} = 1.1\, m_{N_i} = 275$ GeV, and the mass of the singlet scalar is $m_\chi = 650$ GeV with the singlet VEV being $w = 10\, v$, where $v \simeq 246.22$ GeV is the VEV of the SM scalar field. The value of the scalar mixing coupling is $\lambda = 0.1$. Above $T_{\rm pt}$, the scalar field is in the symmetric phase and $m_{N_i} \propto T$. (In the range shown, we neglect vacuum masses of the SM fields.)

The estimates of the critical temperatures of the superweak and electroweak phase transitions were performed in Ref. [30] and presented in the previous edition of this conference by Seller [22]. With the parameters in our examples, the phase transition occurs around $T_{\rm pt} \simeq 2.6$ TeV, as indicated by the kink in the mass of the RHN at this temperature.

6. CP-asymmetry factor in the SWSM

In general, the CP asymmetry factor for the $a \to b + c$ decay is given as the thermal average

$$\epsilon_{a\to b+c} = \frac{\displaystyle\int_{z_a}^{\infty} {\rm d}y_a\, f_{t(a)}(-y_a)\sqrt{y_a^2 - z_a^2}\int_{-1}^{1} {\rm d}x\, \epsilon_{\mathcal{M}}(y_a, x) f_{t(b)}(y_b) f_{t(c)}(y_c)}{\displaystyle\int_{z_a}^{\infty} {\rm d}y_a\, f_{t(a)}(-y_a)\sqrt{y_a^2 - z_a^2}\int_{-1}^{1} {\rm d}x\, f_{t(b)}(y_b) f_{t(c)}(y_c)}\,, \tag{2}$$

where $z_a = m_a/T$, $f_{\rm B/F}(y) = [\exp(y) \mp 1]^{-1}$ denote the statistical factors, with $t(p) = $ B(ose) or F(ermi), giving the statistics type of particle p, and $\epsilon_{\mathcal{M}}$ is the amplitude-level CP-asymmetry factor (see Refs. [14, 15]). The last is proportional to the imaginary part of the amplitude, and its form depends on the specific process and particle physics model. The details of its computation are presented explicitly in Refs. [15, 31].

In Fig. 4, we show the temperature-dependent CP-asymmetry factor computed in the SWSM assuming a 10% relative mass gap between the heavy RHNs as above (solid black) line. The red hatched area is excluded kinematically, where either $N_i \to \phi + L$ or $\phi \to N_i + L$ is forbidden (but other processes giving a non-zero CP violation can happen). Below it, the contribution due to the purely thermal cuts of the vertex function is also shown separately (green, dash-dotted). The temperatures T_i correspond to kinematic thresholds, where $m_{N_i}(T_1) = m_\phi(T_1) + m_L(T_1)$, or $m_\phi(T_2) = m_{N_i}(T_2) + m_L(T_2)$, or $m_\phi(T_3) = m_{N_j}(T_3) + m_L(T_3)$. In the gray hatched area, $T \lesssim T_{\rm sph}$. We see that the thermal CP-asymmetry factor deviates significantly from the $T = 0$ estimate (solid blue line obtained with vanishing masses), where

CP-violating processes occur. At $T > T_2$, where $\phi \to N_i + L$, the self-energy contribution dominates, yet the first and third cuts would only contribute at $T > T_3$. At high temperature where $w(T) = 0$, the $\epsilon_{\phi\to N_i+L}$ becomes roughly a constant (apart from a dependence on the temperature through running couplings with renormalization scale $\mu = 2\pi T$) as there are no other scales left in the problem.

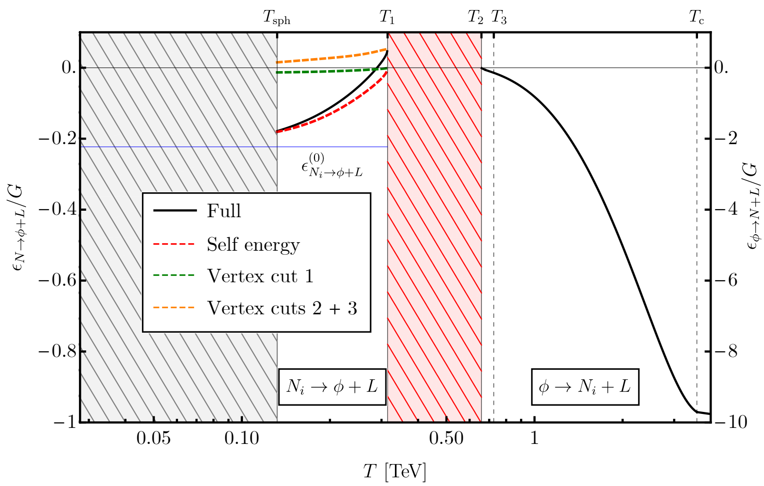

Fig. 4. CP-asymmetry factor of $N_i \to \phi + L$ and $\phi \to N_i + L$ decays. The values of the free parameters are the same as in Fig. 3.

7. Conclusions

In this contribution, we presented a computational scheme to estimate the measure of lepto-baryogenesis based on the Boltzmann equations. We focused on the proper evaluation of the CP-asymmetry factor, whose correct computation had been rather cumbersome to extract from the existing literature prior to the papers [15, 31]. We showed that the thermal effects are non-negligible within the context of the superweak extension of the Standard Model, whose parameter space can be partly constrained by comparing the generated baryon asymmetry to the experimentally measured value.

The estimates presented here depend most sensitively on the mass splits of the RHNs, as the self-energy contribution is proportional to the inverse of $m^2_{N_j} - m^2_{N-i}$, which is called resonance enhancement. The ratio of the scalar VEVs affects the prediction for ϵ by small, but non-negligible amount such that the smaller the VEV of the new scalar as compared to that of the SM scalar, the larger $|\epsilon|$. The effect of the varying scalar mass is surprisingly small, and that of the new gauge coupling g_z is negligible.

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REFERENCES

[1] https://cms-results.web.cern.ch/cms-results/public-results/publications/SMP

[2] Particle Data Group (R.L. Workman *et al.*), *Prog. Theor. Exp. Phys.* **2022**, 083C01 (2022).

[3] Planck Collaboration (N. Aghanim *et al.*), *Astron. Astrophys.* **641**, A1 (2020), arXiv:1807.06205 [astro-ph.CO].

[4] T2K Collaboration (K. Abe *et al.*), *Eur. Phys. J. C* **83**, 782 (2023), arXiv:2303.03222 [hep-ex].

[5] V.A. Kuzmin, V.A. Rubakov, M.E. Shaposhnikov, *Phys. Lett. B* **155**, 36 (1985).

[6] M. D'Onofrio, K. Rummukainen, A. Tranberg, *Phys. Rev. Lett.* **113**, 141602 (2014), arXiv:1404.3565 [hep-ph].

[7] J.A. Harvey, M.S. Turner, *Phys. Rev. D* **42**, 3344 (1990).

[8] M. Fukugita, T. Yanagida, *Phys. Lett. B* **174**, 45 (1986).

[9] B. Garbrecht, *Prog. Part. Nucl. Phys.* **110**, 103727 (2020), arXiv:1812.02651 [hep-ph].

[10] G. Baym, L.P. Kadanoff, «Quantum Statistical Mechanics», *Benjamin*, New York 1962.

[11] M. Beneke, B. Garbrecht, M. Herranen, P. Schwaller, *Nucl. Phys. B* **838**, 1 (2010), arXiv:1002.1326 [hep-ph].

[12] T. Frossard *et al.*, *Phys. Rev. D* **87**, 085009 (2013), arXiv:1211.2140 [hep-ph].

[13] E.W. Kolb, S. Wolfram, *Nucl. Phys. B* **172**, 224 (1980); *Erratum ibid.* **195**, 542 (1982).

[14] G.F. Giudice *et al.*, *Nucl. Phys. B* **685**, 89 (2004), arXiv:hep-ph/0310123.

[15] K. Seller, Z. Szép, Z. Trócsányi, *J. High Energy Phys.* **2025**, 034 (2025), arXiv:2409.07180 [hep-ph].

[16] B. Garbrecht, *Nucl. Phys. B* **847**, 350 (2011), arXiv:1011.3122 [hep-ph].

[17] Z. Trócsányi, *Symmetry* **12**, 107 (2020), arXiv:1812.11189 [hep-ph].

[18] S. Biondini *et al.*, *Int. J. Mod. Phys. A* **33**, 1842004 (2018), arXiv:1711.02864 [hep-ph].

[19] T.J. Kärkkäinen, Z. Trócsányi, *Acta Phys. Pol. B Proc. Suppl.* **15**, 2-A16 (2022), arXiv:2111.07789 [hep-ph].

[20] Z. Trocsanyi, *Acta Phys. Pol. B Proc. Suppl.* **17**, 2-A1 (2024), `arXiv:2311.08828 [hep-ph]`.

[21] Z. Péli, *Acta Phys. Pol. B Proc. Suppl.* **17**, 2-A16 (2024), `arXiv:2311.08203 [hep-ph]`.

[22] K. Seller, *Acta Phys. Pol. B Proc. Suppl.* **17**, 2-A17 (2024), `arXiv:2311.05347 [hep-ph]`.

[23] Z. Péli, Z. Trócsányi, *Phys. Rev. D* **108**, L031704 (2023), `arXiv:2305.11931 [hep-ph]`.

[24] S. Iwamoto, T.J. Kärkkäinen, Z. Péli, Z. Trócsányi, *Phys. Rev. D* **104**, 055042 (2021), `arXiv:2104.14571 [hep-ph]`.

[25] T.J. Kärkkäinen, Z. Trócsányi, *J. Phys. G: Nucl. Part. Phys.* **49**, 045004 (2022), `arXiv:2105.13360 [hep-ph]`.

[26] S. Iwamoto, K. Seller, Z. Trócsányi, *J. Cosmol. Astropart. Phys.* **2022**, 035 (2022), `arXiv:2104.11248 [hep-ph]`.

[27] Z. Péli, I. Nándori, Z. Trócsányi, *Phys. Rev. D* **101**, 063533 (2020), `arXiv:1911.07082 [hep-ph]`.

[28] Z. Péli, Z. Trócsányi, *Phys. Rev. D* **106**, 055045 (2022), `arXiv:2204.07100 [hep-ph]`.

[29] NA64 Collaboration (D. Banerjee *et al.*), *Phys. Rev. D* **97**, 072002 (2018), `arXiv:1710.00971 [hep-ex]`.

[30] K. Seller, Z. Szép, Z. Trócsanyi, *J. High Energy Phys.* **2023**, 096 (2023), `arXiv:2301.07961 [hep-ph]`.

[31] K. Seller, Z. Szép, Z. Trócsanyi, *Eur. Phys. J. C* **85**, 1295 (2025), `arXiv:2511.03445 [hep-ph]`.