PHYSICAL REVIEW D **96,** 015015 (2017)

Calculable cosmological *CP* violation and resonant leptogenesis

Avtandil Achelashvili[*] and Zurab Tavartkiladze[†]

Center for Elementary Particle Physics, ITP, Ilia State University, 0162 Tbilisi, Georgia

(Received 29 November 2016; published 17 July 2017)

Within the extension of MSSM by two right-handed neutrinos, with tree level mass degeneracy, we address the issue of leptogenesis. Investigating the quantum corrections in detail, we show that the lepton asymmetry is induced at the one-loop level and a decisive role is played by the tau lepton Yukawa coupling. On a concrete and predictive neutrino model, which enables us to predict the *CP* violating δ phase and relate it to the cosmological *CP* asymmetry, we demonstrate that the needed amount of the baryon asymmetry is generated via the resonant leptogenesis.

DOI: 10.1103/PhysRevD.96.015015

I. INTRODUCTION

The simplest extension of the standard model (SM), required for accommodation of the atmospheric and solar neutrino data [1], is the inclusion of the SM singlet right-handed neutrinos (RHN). The latter, having the Majorana mass, can generate neutrino masses via the seesaw mechanism. It is remarkable that this simple construction also offers an elegant way for generating the baryon asymmetry of the Universe through thermal leptogenesis [2] (for reviews see [3–5]). In order to reduce the number of parameters entering the *CP* asymmetry, the minimalistic approach with texture zeros has been put forward in Ref. [6]. This approach enables one to relate the *CP* violating phase δ (appearing in the neutrino oscillations) with the cosmological *CP* asymmetry [6–15]. The setup looks especially attractive with two (or more) quasidegenerate RHNs [9–12,15] because, besides the further reduction of the model parameter number, it offers the possibility for resonant leptogenesis [16–18] (for recent discussions on resonant leptogenesis see [19–22]).

With two degenerate RHNs, in [11] all possible one texture zero 3×2 Dirac-type Yukawa couplings have been investigated. As it turns out, because of a very limited number of parameters, these types of models are either disfavored by the current data [1] or do not generate enough of the baryon asymmetry. In order to circumvent this obstacle, in a recent work [15] the setup with two degenerate RHNs and two texture zero 3×2 Dirac-type Yukawa couplings augmented with a single $\Delta L = 2$ lepton number violating the $d = 5$ operator has been investigated. All textures, within such a setup, giving experimentally viable neutrino mass matrices have been studied in great detail. As it turned out [15], some of them together with a successful neutrino sector give interesting predictions and allow one to calculate the cosmological *CP* phase in terms of the neutrino *CP* phase δ.

Encouraged by these findings, in this paper we aim to investigate such a construction in detail from the viewpoint of the leptogenesis. Thus, we start our studies with the minimal SUSY standard model augmented with two RHNs, which at high energy scales are strictly degenerate in mass. The degeneracy is lifted by the renormalization. As we show, taking into account the charged lepton Yukawa couplings in the renormalization procedure (where, in a regime of RHN masses $\lesssim 10^7$ GeV,[1] the decisive role is played by the tau lepton's Yukawa coupling), the nonzero cosmological lepton asymmetry emerges at the one-loop level. Moreover, the sufficient baryogenesis is realized even with RHN masses near the TeV scale and also with low values of the MSSM parameter $\tan\beta(\sim 1)$. As we have mentioned, to make the scenario viable, in Ref. [15] we have included the single $\Delta L = 2, d = 5$ operator, which we adopt also in this paper. The inclusion of such terms does not alter renormalization group (RG) studies, and the results mentioned above are robust. For demonstrative purposes we pick up one of the viable models of [15]. That is the concrete neutrino texture zero mass matrix (referred to as the texture P_1), which emerges via the integration of two (quasi)degenerate RHNs and the single $\Delta L = 2,\ d = 5$ operator. The model's predictive power allows one to compute the cosmological *CP* phase in terms of observed neutrino parameters and *CP* phases (not measured yet, but predicted by the model).

Note that an approach, similar to the one we pursue in this paper, could also work within a non-SUSY framework (i.e., within the SM augmented with two degenerate RHNs). However, since for a solution to the gauge hierarchy problem the supersymmetry appears to be a well motivated (and perhaps the best so far) framework, we choose to perform our investigations within the SUSY setup.

The paper is organized as follows. In Sec. II, we first describe our setup and then, proving the emergence of the

[*]avtandil.achelashvili.1@iliauni.edu.ge
[†]zurab.tavartkiladze@gmail.com

[1]These mass values that we consider within our studies avoid the relic gravitino problem [23,24].

cosmological CP violation via charged lepton Yukawas at the one-loop level, give a detailed calculation of CP violation relevant for the leptogenesis. In Sec. III we present the neutrino scenario (discussed in Ref. [15] together with other scenarios), with the prediction of the CP phase δ and its relation with the cosmological CP violation. On this scenario we demonstrate that leptonic asymmetry, induced at quantum level (and computed in Sec. II), leads to desirable baryon asymmetry via resonant leptogenesis. Then we present one example of the renormalizable UV completion of our model and prove the robustness of all obtained results. Appendix A includes details and various aspects of the RG studies. In Appendix B we investigate the effects of the scalar components of the RHN superfields in the net baryon asymmetry.

II. TWO QUASIDEGENERATE RHN AND COSMOLOGICAL CP

In this section, we first describe our setup and then give a detailed calculation of CP violation relevant for the leptogenesis.

Our framework is the MSSM augmented with two right-handed neutrinos N_1 and N_2. This extension is enough to build a consistent neutrino sector accommodating the neutrino data [1] and also to have a successful leptogenesis scenario. The relevant lepton superpotential couplings, which we are starting with, are given by

$$W_{\rm lept} = l^T Y_e^{\rm diag} e^c h_d + l^T Y_\nu N h_u - \frac{1}{2} N^T M_N N, \quad (1)$$

where h_d and h_u are down- and up-type MSSM Higgs doublet superfields, respectively, and $l^T = (l_1, l_2, l_3)$, $e^{cT} = (e_1^c, e_2^c, e_3^c)$, $N^T = (N_1, N_2)$. We work in a basis in which the charged lepton Yukawa matrix is diagonal and real,

$$Y_e^{\rm diag} = {\rm Diag}(\lambda_e, \lambda_\mu, \lambda_\tau). \quad (2)$$

Moreover, we assume that the RHN mass matrix M_N is strictly degenerate at high scale. For the latter we take the GUT scale $M_G \simeq 2\times 10^{16}$ GeV.[2] Therefore, we assume

$$\text{at } \mu = M_G\colon\ M_N = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} M(M_G). \quad (3)$$

This form of M_N is crucial for our studies. Although it is interesting and worthwhile to study, we do not attempt here to justify the form of M_N (and of the textures considered below) by symmetries. Our approach here is rather phenomenological aiming to investigate possibilities, outcomes, and implications of the textures we consider. Since (3) at a tree level leads to the mass degeneracy of the RHNs, it has interesting implications for resonant leptogenesis [9–11] and also, as we will see below, for building predictive neutrino scenarios [11,15].

For the leptogenesis scenario two necessary conditions need to be satisfied. First of all, at the scale $\mu = M_{N_{1,2}}$ the degeneracy between the masses of N_1 and N_2 has to be lifted. And, at the same scale, the neutrino Yukawa matrix $\hat{Y}_\nu$—written in the mass eigenstate basis of M_N—must be such that ${\rm Im}[(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{12}]^2 \neq 0$. [These can be seen from Eq. (40) with a demand $\epsilon_{1,2} \neq 0$.] Below we show that both these are realized by radiative corrections and the needed effect already arises at the one-loop level, with a dominant contribution due to the Y_e Yukawa couplings (in particular from λ_τ) in the RG.

A. Loop induced cosmological CP violation

Radiative corrections are crucial for the cosmological CP violation. We will start with radiative corrections to the M_N matrix. RG effects cause lifting of the mass degeneracy and, as we will see, are important also for the phase misalignment (explained below).

At the GUT scale, the M_N has an off-diagonal form with $(M_N)_{11} = (M_N)_{22} = 0$ [see Eq. (3)]. However, at low energies, RG corrections generate these entries. Thus, we parametrize the matrix M_N at scale μ as

$$M_N(\mu) = \begin{pmatrix} \delta_N^{(1)}(\mu) & 1 \\ 1 & \delta_N^{(2)}(\mu) \end{pmatrix} M(\mu). \quad (4)$$

While all entries of the matrix M_N run, for our studies the ratios $\frac{(M_N)_{11}}{(M_N)_{12}} = \delta_N^{(1)}$ and $\frac{(M_N)_{22}}{(M_N)_{12}} = \delta_N^{(2)}$ will be relevant (for which we will write and solve RG equations below). That is why we have written M_N in the form given in Eq. (4). With $|\delta_N^{(1,2)}| \ll 1$, the M (at scale $\mu = M$) will determine the masses of RHNs M_1 and M_2, while $\delta_N^{(1,2)}$ will be responsible for their splitting and for complexity in M_N (the phase of the overall factor M does not contribute to the physical CP). As it will turn out (see below),

$$\delta_N^{(1)} = (\delta_N^{(2)})^* \equiv -\delta_N. \quad (5)$$

Therefore, M_N is diagonalized by the transformation

$$U_N^T M_N U_N = M_N^{\rm Diag} = {\rm Diag}(M_1, M_2),$$
$$\text{with } U_N = P_N O_N P_N',$$
$$M_1 = |M|(1-|\delta_N|), \qquad M_2 = |M|(1+|\delta_N|), \quad (6)$$

where

[2] Degeneracy of M_N can be guaranteed by some symmetry at high energies. For concreteness, we assume this energy interval to be $\geq M_G$ (although the degeneracy at lower energies can be considered as well).

$$P_N = \mathrm{Diag}(e^{-i\eta/2}, e^{i\eta/2}),$$
$$O_N = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 & -1 \\ 1 & 1 \end{pmatrix},$$
$$P'_N = \mathrm{Diag}(e^{-i\phi_M/2}, ie^{-i\phi_M/2}),$$
$$\text{with} \quad \eta = \mathrm{Arg}(\delta_N),$$
$$\phi_M = \mathrm{Arg}(M). \tag{7}$$

In the N's mass eigenstate basis, the Dirac-type neutrino Yukawa matrix will be $\hat{Y}_\nu = Y_\nu U_N$. In the CP asymmetries, the components $(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{21}$ and $(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{12}$ appear [see Eq. (40)]. From (6) and (7) we have

$$[(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{21}]^2 = -[(O_N^T P_N^* Y_\nu^\dagger Y_\nu P_N O_N)_{21}]^2,$$
$$[(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{12}]^2 = -[(O_N^T P_N^* Y_\nu^\dagger Y_\nu P_N O_N)_{12}]^2. \tag{8}$$

Therefore, we see that the CP violating part should come from the combination $P_N^* Y_\nu^\dagger Y_\nu P_N$, which in a matrix form is

$$P_N^* Y_\nu^\dagger Y_\nu P_N = \begin{pmatrix} (Y_\nu^\dagger Y_\nu)_{11} & |(Y_\nu^\dagger Y_\nu)_{12}| e^{i(\eta-\eta')} \\ |(Y_\nu^\dagger Y_\nu)_{21}| e^{i(\eta'-\eta)} & (Y_\nu^\dagger Y_\nu)_{22} \end{pmatrix},$$
$$\text{with} \quad \eta' = \mathrm{Arg}[(Y_\nu^\dagger Y_\nu)_{21}]. \tag{9}$$

We see that the $\eta' - \eta$ difference (mismatch) will govern the CP asymmetric decays of the RHNs. Without including the charged lepton Yukawa couplings in the RG effects we will have $\eta' \simeq \eta$ with a high accuracy. It was shown in Ref. [21] that by ignoring Y_e Yukawas no CP asymmetry emerges at $\mathcal{O}(Y_\nu^4)$ order and nonzero contributions start only from $\mathcal{O}(Y_\nu^6)$ terms [22]. Such corrections are extremely suppressed for $Y_\nu \lesssim 1/50$. Since in our consideration we are interested in cases with $M_{1,2} \lesssim 10^7$ GeV giving $|(Y_\nu)_{ij}| < 7 \times 10^{-4}$ (well fixed from the neutrino sector and the desired value of the baryon asymmetry), these effects (i.e., order $\sim Y_\nu^6$ corrections) will not have any relevance. In Ref. [11] in the RG of M_N the effect of Y_e, coming from two-loop corrections, was taken into account and showed that sufficient CP violation can emerge. Below we show that including Y_e in the Y_ν's one-loop RG, will induce a sufficient amount of CP violation. This mainly happens via λ_τ Yukawa coupling. Thus, below we give a detailed investigation of λ_τ's effect.

Using M_N's RG given in Eq. (A3) (of Appendix A 1), for $\delta_N^{(1,2)}$, which are the ratios $\frac{(M_N)_{11}}{(M_N)_{12}}$ and $\frac{(M_N)_{22}}{(M_N)_{12}}$ [see parametrization in Eq. (4)], we can derive the following RG equations:

$$16\pi^2 \frac{d}{dt}\delta_N^{(1)} = 4(Y_\nu^\dagger Y_\nu)_{21} + 2\delta_N^{(1)}[(Y_\nu^\dagger Y_\nu)_{11} - (Y_\nu^\dagger Y_\nu)_{22}]$$
$$- 2(\delta_N^{(1)})^2 (Y_\nu^\dagger Y_\nu)_{12} - 2\delta_N^{(1)}\delta_N^{(2)}(Y_\nu^\dagger Y_\nu)_{21}$$
$$- \frac{1}{4\pi^2}(Y_\nu^\dagger Y_e Y_e^\dagger Y_\nu)_{21} + \cdots, \tag{10}$$

$$16\pi^2 \frac{d}{dt}\delta_N^{(2)} = 4(Y_\nu^\dagger Y_\nu)_{12} + 2\delta_N^{(2)}[(Y_\nu^\dagger Y_\nu)_{22} - (Y_\nu^\dagger Y_\nu)_{11}]$$
$$- 2(\delta_N^{(2)})^2 (Y_\nu^\dagger Y_\nu)_{21} - 2\delta_N^{(1)}\delta_N^{(2)}(Y_\nu^\dagger Y_\nu)_{12}$$
$$- \frac{1}{4\pi^2}(Y_\nu^\dagger Y_e Y_e^\dagger Y_\nu)_{12} + \cdots, \tag{11}$$

where in the second lines of (10) and (11) are given two-loop corrections depending on Y_e. Dots there stand for higher order irrelevant terms. From two-loop corrections we keep only Y_e dependent terms. Remaining contributions are not relevant for us.[3] From (10) and (11) we see that dominant contributions come from the first terms of the right-hand side (RHS) and from those given in the second rows. Other terms give contributions of order $\mathcal{O}(Y_\nu^4)$ or higher and thus will be ignored. At this approximation we have

$$\delta_N^{(1)}(t) \simeq \delta_N^{(2)*}(t) \equiv -\delta_N(t)$$
$$\simeq -\frac{1}{4\pi^2}\int_t^{t_G} dt \left(Y_\nu^\dagger \left(\mathbf{1} - \frac{1}{16\pi^2} Y_e Y_e^\dagger \right) Y_\nu \right)_{21}, \tag{12}$$

where $t = \ln \mu$, $t_G = \ln M_G$, and we have used the boundary conditions at the GUT scale $\delta_N^{(1)}(t_G) = \delta_N^{(2)}(t_G) = 0$. For evaluation of the integral in (12) we need to know the scale dependence of Y_ν and Y_e. This is found in Appendix A 1 by solving the RG equations for Y_ν and Y_e. Using Eqs. (A5) and (A6), the integral of the matrix appearing in (12) can be written as

$$\int_{t_M}^{t_G} Y_\nu^\dagger \left(\mathbf{1} - \frac{1}{16\pi^2} Y_e Y_e^\dagger \right) Y_\nu dt$$
$$\simeq \bar{\kappa}(M) Y_{\nu G}^\dagger \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & \bar{r}_\tau(M) \end{pmatrix} Y_{\nu G}, \tag{13}$$

where

$$\bar{r}_\tau(M) = \frac{\int_{t_M}^{t_G} \kappa(t) r_\tau(t)(1 - \frac{\lambda_\tau^2}{16\pi^2}) dt}{\int_{t_M}^{t_G} \kappa(t) dt}, \quad \bar{\kappa}(M) = \int_{t_M}^{t_G} \kappa(t) dt, \tag{14}$$

[3] Omitted terms are either strongly suppressed or do not give any significant contribution neither to the CP violation nor to the RHN mass splittings.

$$r_\tau(\mu) = \eta_\tau^2(\mu), \qquad \kappa(\mu) = \eta_t^6(\mu)\eta_{g\nu}^2(\mu), \tag{15}$$

and we have ignored $\lambda_{e,\mu}$ Yukawa couplings. For the definition of η factors see Eq. (A6). The $Y_{\nu G}$ denote the corresponding Yukawa matrix at scale $\mu = M_G$. On the other hand, we have

$$(Y_\nu^\dagger Y_\nu)|_{\mu=M} \simeq \kappa(M) Y_{\nu G}^\dagger \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & r_\tau(M) \end{pmatrix} Y_{\nu G}. \tag{16}$$

(Derivations are given in Appendix A 1.)

Comparing (13) with (16) we see that the difference in these matrix structures (besides overall flavor universal RG factors) are in the RG factors $r_\tau(M)$ and $\bar{r}_\tau(M)$. Without the λ_τ Yukawa coupling these factors are equal, and there is no mismatch between the phases η and η' [defined in Eqs. (7) and (9)] of these matrices. Nonzero $\eta' - \eta$ will be due to the deviation, which we parametrize as

$$\xi = \frac{\bar{r}_\tau(M)}{r_\tau(M)} - 1. \tag{17}$$

This value can be computed numerically by evaluation of the appropriate RG factors. However, it is useful to have an approximate expression for ξ, which is given by

$$\xi \simeq \left[\frac{\lambda_\tau^2(M)}{16\pi^2}\ln\frac{M_G}{M} + \frac{1}{3}\frac{\lambda_\tau^2(M)}{(16\pi^2)^2}[3\lambda_t^2 + 6\lambda_b^2 + 10\lambda_\tau^2 - (2c_e^a + c_\nu^a)g_a^2]_{\mu=M}\left(\ln\frac{M_G}{M}\right)^2\right]_{1-\text{loop}} - \left[\frac{\lambda_\tau^2(M)}{16\pi^2}\right]_{2-\text{loop}}, \tag{18}$$

where one- and two-loop contributions are indicated. The derivation of this expression is given in Appendix A 1. As we see, nonzero ξ is induced already at the one-loop level [without the two-loop correction of $\frac{\lambda_\tau^2}{16\pi^2}$ in Eq. (14)]. However, the inclusion of the two-loop correction can contribute to ξ by the amount of ~3%–5% (because of the $\ln\frac{M_G}{M}$ factor suppression), and we have included it.

Now we are ready to write down quantities that have direct relevance for the leptogenesis. From (12), with definitions introduced above and by obtained relations, we have

$$|\delta_N(M)|e^{i\eta} = \frac{1}{4\pi^2}\frac{\bar{\kappa}(M)}{\kappa(M)} \times [|(Y_\nu^\dagger Y_\nu)_{21}|e^{i\eta'} + \xi|(Y_\nu)_{31}(Y_\nu)_{32}|e^{i(\phi_{31}-\phi_{32})}]_{\mu=M}, \tag{19}$$

where ϕ_{31} and ϕ_{32} are phases of the matrix elements $(Y_\nu)_{31}$ and $(Y_\nu)_{32}$, respectively, at scale $\mu = M$. Equation (19) shows well that in the limit $\xi \to 0$, we have $\eta = \eta'$, while the mismatch of these two phases are due to $\xi \neq 0$. With $\xi \ll 1$, from (19) we derive

$$\eta - \eta' \simeq \xi\frac{|(Y_\nu)_{31}(Y_\nu)_{32}|}{|(Y_\nu^\dagger Y_\nu)_{21}|}\sin(\phi_{31} - \phi_{32} - \eta'). \tag{20}$$

We stress that the one-loop renormalization of the Y_ν matrix plays the leading role in the generation of ξ, i.e., in the CP violation.[4] [This is also demonstrated by Eq. (18).]

The value of $|\delta_N(M)|$, which characterizes the mass splitting between the RHNs, can be computed taking the absolute value of both sides of (19),

$$|\delta_N(M)| = \frac{\kappa_N}{4\pi^2}|(Y_\nu^\dagger Y_\nu)_{21} + \xi(Y_\nu)_{31}(Y_\nu^*)_{32}|_{\mu=M}\ln\frac{M_G}{M}, \quad \text{with} \quad \kappa_N = \frac{\bar{\kappa}(M)}{\kappa(M)\ln\frac{M_G}{M}}. \tag{21}$$

These expressions can be used upon the calculation of the leptogenesis, which we will do in the next section for one concrete model of the neutrino mass matrix.

III. PREDICTIVE NEUTRINO TEXTURE AND BARYON ASYMMETRY

In this section we apply obtained results within the setup of the couplings (1) augmented by single $\Delta L = 2$, $d = 5$ operator. As was shown in [15], this could lead to the successful and predictive neutrino sectors. With the addition of this $d = 5$ operator, the results obtained above can remain intact. We consider one neutrino scenario that allows us to predict the CP phase δ and relate it with the cosmological CP violation leading to desirable baryon asymmetry via resonant leptogenesis. First we discuss the neutrino sector and then turn to the investigation of the leptogenesis. At the end, we present one possible renormalizable UV completion (giving rise to the $\Delta L = 2, d = 5$ operator that we utilize) maintaining all obtained results.

A. P_1 neutrino texture: Relating δ and cosmological CP

In the work of Ref. [15], within the setup of two (quasi) degenerate RHNs were studied neutrino mass matrices that emerged from two zero 3×2 Yukawa textures in combination with one $d = 5$ entry. In this way, all experimentally viable neutrino mass matrices have been investigated, which also predicted CP violation and gave promise for successful leptogenesis. Here, for concreteness we consider one scenario of the neutrino mass matrix—called in [15]

[4]Note that since RG equations for M_N and Y_ν in the non-SUSY case have similar structures (besides some group-theoretical factors), the ξ would be generated also within the non-SUSY setup.

the P_1-type texture—and show that it admits having calculable *CP* violation.

Thus, we consider the Yukawa matrix with the form

$$Y_\nu = \begin{pmatrix} 0 & 0 \\ a_2 e^{i\alpha_2} & b_2 e^{i\beta_2} \\ a_3 e^{i\alpha_3} & b_3 e^{i\beta_3} \end{pmatrix} = \begin{pmatrix} e^{ix} & 0 & 0 \\ 0 & e^{iy} & 0 \\ 0 & 0 & e^{iz} \end{pmatrix} \begin{pmatrix} 0 & 0 \\ a_2 & b_2 \\ a_3 & b_3 e^{i\phi} \end{pmatrix} \begin{pmatrix} e^{i\omega} & 0 \\ 0 & e^{i\rho} \end{pmatrix}, \tag{22}$$

with

$$\omega = \alpha_2 - \beta_2 + \rho, \qquad y = \beta_2 - \rho,$$
$$z = \alpha_3 - \alpha_2 + \beta_2 - \rho, \qquad \phi = \alpha_2 - \alpha_3 + \beta_3 - \beta_2, \tag{23}$$

where, only one phase ϕ will be relevant for the cosmological *CP* asymmetry. The phases x, y, z can be removed by proper phase redefinitions of the states l_i and e_i^c. Using this and the form of M_N, given in Eq. (3), via the seesaw formula we get the following contribution to the neutrino mass matrix:

$$M_\nu^{ss} = -\langle h_u^0 \rangle^2 Y_\nu M_N^{-1} Y_\nu^T. \tag{24}$$

Besides this, we include the d = 5 operator

$$\frac{\tilde{d}_5 e^{ix_5}}{M_*} l_1 l_2 h_u h_u, \tag{25}$$

where M_* and $\tilde{d}_5$ are some cutoff scale and dimensionless coupling, respectively. With proper phase redefinitions of l_i states, without loss of any generality, both of these can be taken real and the phase x_5 selected as $x_5 = \omega + \rho - \arg(M)$. The origin of the operator (25) and consistency of our construction will be discussed in Sec. III C. Taking into account these and Eq. (24), the neutrino mass matrix at scale M will have the form

$$M_\nu(M) = -\begin{pmatrix} 0 & d_5 & 0 \\ d_5 & 2a_2 b_2 & a_3 b_2 + a_2 b_3 e^{i\phi} \\ 0 & a_3 b_2 + a_2 b_3 e^{i\phi} & 2a_3 b_3 e^{i\phi} \end{pmatrix} \times \frac{v_u^2(M)}{M \cdot e^{-i(\omega+\rho)}}, \quad \text{with} \quad d_5 = \tilde{d}_5 \frac{|M|}{M_*}, \tag{26}$$

where in M_N we have ignored (1,1) and (2,2) elements, which are induced at one-loop level and are so suppressed that they have no impact on light neutrino masses and mixings. By the renormalization (discussed in Appendix A 2) for the neutrino mass matrix at scale M_Z we obtain

$$M_\nu(M_Z) = \begin{pmatrix} 0 & d_5 & 0 \\ d_5 & 2a_2 b_2 & (a_3 b_2 + a_2 b_3 e^{i\phi}) r_{\nu 3} \\ 0 & (a_3 b_2 + a_2 b_3 e^{i\phi}) r_{\nu 3} & 2a_3 b_3 e^{i\phi} r_{\nu 3}^2 \end{pmatrix} \bar{m}, \quad \text{with} \quad \bar{m} = -\frac{r_{\bar{m}} v_u^2(M_Z)}{M \cdot e^{-i(\omega+\rho)}}, \tag{27}$$

where the couplings a_i, b_i, d_5, and phases appearing in (27) are given at scale M. The RG factors $r_{\nu 3}$ and $r_{\bar{m}}$ are given in Eqs. (A17) and (A18), respectively. The neutrino mass matrix (27) is of the P_1 type investigated in details in [15].

Noting that we are working in a basis of a diagonal charged lepton mass matrix, the neutrino mass matrix can be related to the lepton mixing matrix U by

$$M_\nu = P U^* P' M_\nu^{\rm diag} U^+ P, \tag{28}$$

where $M_\nu^{\rm diag} = (m_1, m_2, m_3)$ and the phase matrices and U are

$$P = {\rm Diag}(e^{i\omega_1}, e^{i\omega_2}, e^{i\omega_3}),$$
$$P' = {\rm Diag}(1, e^{i\rho_1}, e^{i\rho_2}), \tag{29}$$

$$U = \begin{pmatrix} c_{13}c_{12} & c_{13}s_{12} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}s_{13}c_{12}e^{i\delta} & c_{23}c_{12} - s_{23}s_{13}s_{12}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}s_{13}c_{12}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{13}s_{12}e^{i\delta} & c_{23}c_{13} \end{pmatrix}. \tag{30}$$

As was discussed in detail in [15], the texture (27) allows only a normal neutrino mass hierarchy. Using the conditions $M_\nu^{(1,1)} = M_\nu^{(1,3)} = 0$ in Eq. (28), we obtain the following predictions:

$$m_3^2 = \frac{\Delta m_{atm}^2 + \Delta m_{\rm sol}^2 c_{12}^2}{1 - s_{13}^2 {\rm cot}_{23}^2 (1 + t_{13}^2)^2 - t_{13}^4}, \qquad \cos\rho_1 = \frac{m_3^2 t_{13}^4 - m_1^2 c_{12}^4 - m_2^2 s_{12}^4}{2 m_1 m_2 c_{12}^2 s_{12}^2}, \tag{31}$$

TABLE I. Results from the P_1-type texture of Eq. (27). Masses are given in eVs.

δ	ρ_1	ρ_2	Works with
± 0.378	± 3.036	± 2.696	NH, $\sin^2\theta_{23} = 0.49$ and best fit values [of Eq. (33)] for remaining oscillation parameters, $(m_1, m_2, m_3) = (0.00613, 0.0106, 0.0499)$, $m_{\beta\beta} = 0$

$$\delta = \arg[m_1 c_{12}^2 + m_2 s_{12}^2 e^{i\rho_1}] - \arg[m_1 - m_2 e^{i\rho_1}],$$
$$\rho_2 = \pm\pi - \arg[m_1 c_{12}^2 + m_2 s_{12}^2 e^{i\rho_1}] + 2\arg[m_1 - m_2 e^{i\rho_1}], \tag{32}$$

where by definition $\Delta m_{\rm atm}^2 = m_3^2 - m_2^2$ and $\Delta m_{\rm sol}^2 = m_2^2 - m_1^2$. With the inputs

$$\sin^2\theta_{12} = 0.304, \quad \sin^2\theta_{23} = 0.49, \quad \sin^2\theta_{13} = 0.0218,$$
$$\Delta m_{\rm atm}^2 = 0.002382\ {\rm eV}^2, \quad \Delta m_{\rm sol}^2 = 7.5\times 10^{-5}\ {\rm eV}^2, \tag{33}$$

we obtain the values

$$m_1 = 0.00613\ {\rm eV}, \quad m_2 = 0.0106\ {\rm eV}, \quad m_3 = 0.0499\ {\rm eV},$$
$$\rho_1 = \pm 3.036, \quad \delta = \pm 0.378, \quad \rho_2 = \pm 2.696. \tag{34}$$

Notice that besides $\sin^2\theta_{23}$ all inputs of Eq. (33) are taken to be the best fit values [1]. The results are summarized in Table I.

At the same time, from (28) we have the relations

$$2a_2 b_2 \bar{m} = e^{2i\omega_2}\mathcal{A}_{22},$$
$$2a_3 b_3 e^{i\phi}\bar{m} r_{\nu 3}^2 = e^{2i\omega_3}\mathcal{A}_{33},$$
$$(a_3 b_2 + a_2 b_3 e^{i\phi})\bar{m} r_{\nu 3} = e^{i(\omega_2+\omega_3)}\mathcal{A}_{23}, \tag{35}$$

with

$$\mathcal{A}_{ij} = U_{i1}^* U_{j1}^* m_1 + U_{i2}^* U_{j2}^* m_2 e^{i\rho_1} + U_{i3}^* U_{j3}^* m_3 e^{i\rho_2}. \tag{36}$$

Note that from the neutrino sector all $\mathcal{A}_{ij}$ numbers are determined with the help of zero entries in the matrix of Eq. (27). With the help of the phases appearing in (22), without loss of generality we can take $a_i, b_i > 0$. With this, from Eqs. (35) we can express $|\bar{m}|$ and the couplings a_3, $b_{2,3}$ in terms of a_2 and $|M|$ as follows:

$$|\bar{m}| = \frac{v_u^2(M_Z)}{|M|} r_{\bar{m}},$$
$$a_3 = \frac{a_2}{r_{\nu 3}}\left|\frac{1}{\mathcal{A}_{22}}\left(\mathcal{A}_{23} \pm \sqrt{\mathcal{A}_{23}^2 - \mathcal{A}_{22}\mathcal{A}_{33}}\right)\right|,$$
$$b_2 = \frac{1}{a_2}\frac{|\mathcal{A}_{22}|}{2|\bar{m}|}, \qquad b_3 = \frac{1}{a_3}\frac{|\mathcal{A}_{33}|}{2|\bar{m}| r_{\nu 3}^2}. \tag{37}$$

Also, for the phase ϕ we get the following prediction:

$$\phi = {\rm Arg}\left[\left(\frac{\mathcal{A}_{23}}{\sqrt{\mathcal{A}_{22}\mathcal{A}_{33}}} \mp \sqrt{\frac{\mathcal{A}_{23}^2}{\mathcal{A}_{22}\mathcal{A}_{33}} - 1}\right)^2\right]. \tag{38}$$

Notice that there is a pair of solutions. When for the a_3's expression in Eq. (37) we are taking the $+$ sign, in Eq. (38) we should take the sign $-$, and vice versa.

From these, using results given in Table I, we find the numerical value of ϕ,

$$\text{for } \delta = +0.378\colon \phi_+ = +1.287, \ \phi_- = -1.287,$$
$$\text{for } \delta = -0.378\colon \phi_+ = -1.287, \ \phi_- = +1.287, \tag{39}$$

where ϕ's subscripts correspond to the signs taken in (38). These and the relations of (37) will be used upon calculation of the baryon asymmetry, which we do in the next subsection.

B. Resonant leptogenesis

The CP asymmetries ϵ_1 and ϵ_2 generated by out-of-equilibrium decays of the quasidegenerate fermionic components of N_1 and N_2 states, respectively, are given by [17,18][5]

$$\epsilon_1 = \frac{{\rm Im}[(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{21}]^2}{(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{11}(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{22}}\frac{(M_2^2 - M_1^2)M_1\Gamma_2}{(M_2^2 - M_1^2)^2 + M_1^2\Gamma_2^2},$$
$$\epsilon_2 = \epsilon_1(1 \leftrightarrow 2). \tag{40}$$

Here M_1, M_2 (with $M_2 > M_1$) are the mass eigenvalues of the RHN mass matrix. These masses, within our scenario, are given in (6) with the splitting parameter given in Eq. (21). The decay widths of fermionic RHNs are given by $\Gamma_i = \frac{M_i}{4\pi}(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{ii}$. Moreover, the imaginary part of $[(\hat{Y}_\nu^\dagger \hat{Y}_\nu)_{21}]^2$ will be computed with the help of (8) and (9) with the relevant phase given in Eq. (20). Using general expressions (20) and (21) for the neutrino model discussed in the previous subsection, we get

[5] In Appendix B we investigate the contribution to the baryon asymmetry via decays of the scalar components of the RHN superfields. As we show, these effects are less than 3%.

TABLE II. Cases with different values of $m_t(m_t)$ and M_S.

	Case $(\mathbf{I}_-)$	Case $(\mathbf{I})$	Case $(\mathbf{II}_-)$	Case $(\mathbf{II})$
$m_t(m_t)$	162.77 GeV	163.48 GeV	162.77 GeV	163.48 GeV
M_S	10^3 GeV	10^3 GeV	2×10^3 GeV	2×10^3 GeV

$$\eta-\eta'\simeq -\xi\frac{\frac{a_2b_2}{a_3b_3}\sin\phi}{(\frac{a_2b_2}{a_3b_3}+\cos\phi)^2+\sin^2\phi},$$
$$|\delta_N(M)|=\frac{\kappa_N}{4\pi^2}|a_2b_2+a_3b_3(1+\xi)e^{i\phi}|\ln\frac{M_G}{M}. \quad (41)$$

With these, since we know the possible values of the phase ϕ [see Eq. (39)], and with the help of the relations (37) we can compute $\epsilon_{1,2}$ in terms of $|M|$ and a_2. Recalling that the lepton asymmetry is converted to the baryon asymmetry via sphaleron processes [25], with the relation $\frac{n_b}{s}\simeq -1.48\times 10^{-3}(\kappa_f{}^{(1)}\epsilon_1+\kappa_f{}^{(2)}\epsilon_2)$ we can compute the baryon asymmetry. For the efficiency factors $\kappa_f{}^{(1,2)}$ we will use the extrapolating expressions [3] [see Eq. (40) in Ref. [3]], with $\kappa_f{}^{(1)}$ and $\kappa_f{}^{(2)}$ depending on the mass scales $\tilde{m}_1=\frac{v_u^2(M)}{M_1}(\hat{Y}_\nu^\dagger\hat{Y}_\nu)_{11}$ and $\tilde{m}_2=\frac{v_u^2(M)}{M_2}(\hat{Y}_\nu^\dagger\hat{Y}_\nu)_{22}$, respectively.

Within our studies we will consider the RHN masses $\simeq|M|\lesssim 10^7$ GeV. With this, we will not have the relic gravitino problem [23,24]. For the simplicity, we consider all SUSY particle masses to be equal to $M_S<|M|$, with M_S identified with the SUSY scale, below which we have just SM. As it turns out, via the RG factors, the asymmetry also depends on the top quark mass. Therefore, we will consider cases given in Table II, where cases of low top quark masses by $1-\sigma$ deviation are included [i.e., cases $(\mathbf{I}_-)$ and $(\mathbf{II}_-)$]. It is remarkable that the observed baryon asymmetry

$$\left(\frac{n_b}{s}\right)_{\rm exp}=(8.65\pm 0.085)\times 10^{-11} \quad (42)$$

(the recent value reported according to WMAP and Planck [26]) can be obtained even for low values of the MSSM parameter $\tan\beta=\frac{v_u}{v_d}$ (defined at the SUSY scale $\mu=M_S$). This, for different cases and different values of M, is demonstrated in Table III. For the calculations we have used the RG factors found by numerical computations. The details of this procedure, appropriate boundary and matching conditions, are given in Appendix A 3.

While Table III deals with cases of the low $\tan\beta$, in plots of Fig. 1 we show baryon asymmetries as functions of a_2 (the logs of these values for convenience) for different values of the parameters M_S, M, $\tan\beta$, and the phases ϕ of Eq. (39). We see that needed baryon asymmetry is obtained for a wide range of phenomenologically interesting values of parameters. With the values of a_2 giving the needed values of the baryon asymmetry, we have also calculated [via relations of Eq. (37)] the values of a_3, $b_{2,3}$, which also turned out to be suppressed, i.e., a_3, $b_{2,3}\lesssim a_2$.

C. Renormalizable UV completion and consistency check

Upon building the neutrino mass matrix (26), together with seesaw contribution (24) (emerged via integration of $N_{1,2}$ states) we have used the $d=5$ operator (25). Here we present one renormalizable completion of the model, which gives the latter operator. Also we check the whole construction and show what conditions should be satisfied in order to have a fully consistent model without affecting obtained results.

For building a fully renormalizable model, we introduce two additional RHN states $\mathcal{N}$ and $\bar{\mathcal{N}}$ with the following superpotential couplings:

$$\lambda l_1\mathcal{N}h_u+\bar{\lambda}l_2\bar{\mathcal{N}}h_u-M_*\mathcal{N}\bar{\mathcal{N}}. \quad (43)$$

TABLE III. Baryon asymmetry for various values of M and for the minimal (allowed) value of $\tan\beta$. The values of $(\frac{n_b}{s})_{\rm max}$ given here are obtained for all cases of Eq. (39), but for different values of a_i, b_j. [For phase sign choices see (38), (39), and comments after these equations.]

Case	M[GeV]	$\tan\beta$	$r_{\nu 3}$	$r_{\tilde{m}}$	r_{v_u}	κ_N	$10^5\times\xi$	$10^{11}\times(\frac{n_b}{s})_{\rm max}$
$(\mathbf{I}_-)$	3×10^3	1.63	$\simeq 1$	0.8861	0.9713	1.230	5.678	8.573
$(\mathbf{I.1})$	3×10^3	1.636	$\simeq 1$	0.8849	0.9709	1.242	5.729	8.565
$(\mathbf{I.2})$	10^4	1.665	$\simeq 1$	0.8343	0.953	1.211	5.490	8.564
$(\mathbf{I.3})$	10^5	1.72	$\simeq 1$	0.7530	0.9218	1.1596	5.0317	8.559
$(\mathbf{I.4})$	10^6	1.775	$\simeq 1$	0.6883	0.8944	1.118	4.574	8.557
$(\mathbf{I.5})$	10^7	1.831	$\simeq 1$	0.6369	0.8703	1.0834	4.118	8.565
$(\mathbf{II}_-)$	6×10^3	1.608	$\simeq 1$	0.8685	0.9677	1.197	5.462	8.557
$(\mathbf{II.1})$	6×10^3	1.615	$\simeq 1$	0.8670	0.9673	1.206	5.515	8.564
$(\mathbf{II.2})$	10^4	1.627	$\simeq 1$	0.8468	0.9600	1.195	5.416	8.563
$(\mathbf{II.3})$	10^5	1.681	$\simeq 1$	0.7671	0.9295	1.147	4.968	8.557
$(\mathbf{II.4})$	10^6	1.736	$\simeq 1$	0.7034	0.9027	1.108	4.523	8.565
$(\mathbf{II.5})$	10^7	1.79	$\simeq 1$	0.6524	0.8790	1.076	4.072	8.564

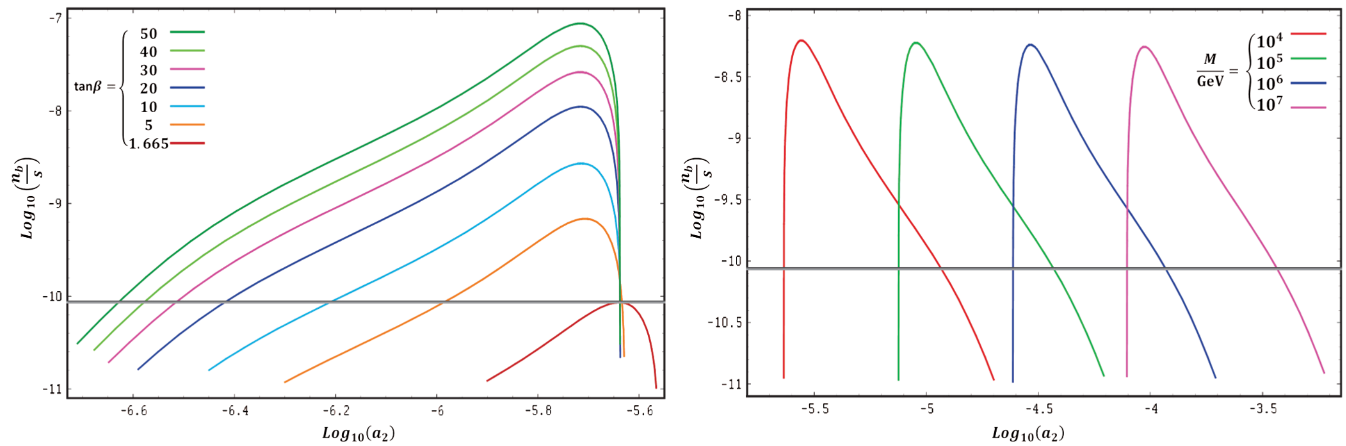

FIG. 1. Left: Curves for case (I) (see Table II), with $M = 10^4$ GeV, $\delta = +0.378$, $\phi = \phi_+ = +1.287$, and with different values of $\tan\beta$. Right: Curves for case (II) (see Table II), with $\tan\beta = 15$ GeV, $\delta = -0.378$, $\phi = \phi_+ = -1.287$, and with different values of M. Gray horizontal bands correspond to the experimental value of the baryon asymmetry within the $1-\sigma$ range given in Eq. (42).

With these and the couplings of (1)–(3), (22), after removing the phases x, y, x, ω, ρ in Y_ν (by proper redefinition of the fields) without loss of generality $\bar\lambda$ and M_* can be taken real and $\arg(\lambda) = \arg(\bar m)$. Thus, the full (i.e., "extended") Yukawa and RHN matrices will be

$$Y_\nu^{\rm ext} = \begin{matrix} & \begin{matrix} N_1 & N_2 & \mathcal{N} & \bar{\mathcal{N}} \end{matrix} \\ \begin{matrix} l_1 \\ l_2 \\ l_3 \end{matrix} & \begin{pmatrix} 0 & 0 & \lambda & 0 \\ a_2 & b_2 & 0 & \bar\lambda \\ a_3 & b_3 e^{i\phi} & 0 & 0 \end{pmatrix} \end{matrix}, \qquad M_N^{\rm ext} = \begin{matrix} & \begin{matrix} N_1 & N_2 & \mathcal{N} & \bar{\mathcal{N}} \end{matrix} \\ \begin{matrix} N_1 \\ N_2 \\ \mathcal{N} \\ \bar{\mathcal{N}} \end{matrix} & \begin{pmatrix} 0 & M & 0 & 0 \\ M & 0 & 0 & 0 \\ 0 & 0 & 0 & M_* \\ 0 & 0 & M_* & 0 \end{pmatrix} \end{matrix} \tag{44}$$

With these forms, integration of heavy RHN states leads to the neutrino mass matrix

$$M_\nu = -v_u^2 Y_\nu^{\rm ext} (M_N^{\rm ext})^{-1} (Y_\nu^{\rm ext})^T, \tag{45}$$

which, as desired, indeed has the form of (26) with

$$d_5 = |\lambda| \bar\lambda \frac{|M|}{M_*}. \tag{46}$$

Furthermore, one should make sure that via loops the couplings λ and $\bar\lambda$ instead of zeros in the textures of Eq. (44) do not induce entries which would affect and/or spoil the results of the neutrino sector and leptogenesis. To check this, one can apply one-loop RGs for the neutrino Yukawas and RHN masses. Namely, in Eqs. (A2) and (3) with the replacements $Y_\nu \to Y_\nu^{\rm ext}$, $M_N \to M_N^{\rm ext}$ we can estimate the one-loop contributions due to the λ, $\bar\lambda$ couplings.[6] Since the structure of $Y_\nu^{\rm ext}$ may be altered only by the second term on the RHS of (A2), we will calculate only the contribution due to this type of entry. By the same reason, for the $M_N^{\rm ext}$'s correction, we will focus only on the first term (and on its transpose) on the RHS of Eq. (A3). Doing so, with an assumption $M_* > |M|$, at scale $\mu = M_*$ we obtain

$$\delta Y_\nu^{\rm ext} \approx -\tfrac{3}{16\pi^2} \begin{pmatrix} 0 & 0 & \lambda|\lambda|^2 & 0 \\ a_2\bar\lambda^2 & b_2\bar\lambda^2 & 0 & \bar\lambda(a_2^2+b_2^2+\bar\lambda^2) \\ \times & \times & 0 & \bar\lambda(a_2 a_3 + b_2 b_3 e^{i\phi}) \end{pmatrix} \ln\tfrac{M_G}{M_*},$$

$$\delta M_N^{\rm ext} \approx -\tfrac{1}{8\pi^2} \begin{pmatrix} \times & \times & a_2\bar\lambda M_* & b_2\bar\lambda M \\ \times & \times & b_2\bar\lambda M_* & a_2\bar\lambda M \\ a_2\bar\lambda M_* & b_2\bar\lambda M_* & 0 & (|\lambda|^2+\bar\lambda^2)M_* \\ b_2\bar\lambda M & a_2\bar\lambda M & (|\lambda|^2+\bar\lambda^2)M_* & 0 \end{pmatrix} \ln\tfrac{M_G}{M_*}, \tag{47}$$

[6] Since (as we have seen) the couplings a_i, b_i are small, their corrections in the RG of $Y_\nu^{\rm ext}$ do not harm anything.

where we have taken into account that at scale $\mu = M_G$ the couplings $Y_\nu^{\rm ext}$, $M_N^{\rm ext}$ have forms given in Eq. (44). In (47) $\times$ stands for the corrections which do not depend on λ and/or $\bar{\lambda}$. Comparing (47) with (44) we see that the structure of $Y_\nu^{\rm ext}$ is not changed and $\delta Y_\nu^{\rm ext}$ can be negligible for $\lambda, \bar{\lambda} \lesssim \lambda_\tau/10$. In fact, from the neutrino sector, we have

$$d_5|\bar{m}| = |\mathcal{A}_{12}| \simeq 1.07 \times 10^{-11}\ \text{GeV} \tag{48}$$

[see Eqs. (27) and (36) for definitions]. With this, on the other hand, we have

$$d_5 \approx 4.15 \times 10^{-12}\left(\frac{M}{10^4\ \text{GeV}}\right)\left(\frac{1}{\sin\beta}\right)^2\left(\frac{0.85}{r_{\bar{m}}}\right). \tag{49}$$

With this and $M_* = (3-10)M$, the (46) can be satisfied by the selection

$$|\lambda| \approx \bar{\lambda} = \left(d_5\frac{M_*}{|M|}\right)^{1/2} \simeq (3.5-6.4)\times 10^{-6}\left(\frac{M}{10^4\ \text{GeV}}\right)^{1/2}\left(\frac{1}{\sin\beta}\right)\left(\frac{0.85}{r_{\bar{m}}}\right)^{1/2}. \tag{50}$$

This in turn gives

$$\text{for}\ \ M \lesssim 10^7\ \text{GeV},\quad \tan\beta > 1.6 \Rightarrow |\lambda| \approx \bar{\lambda} < 3\times 10^{-4}. \tag{51}$$

We checked and made sure that, for such small values of λ, $\bar{\lambda}$, the corrections $\delta Y_\nu^{\rm ext}$ and $\delta M_N^{\rm ext}$ are affecting neither the neutrino sector nor the leptogenesis. We have also checked that two-loop corrections are very suppressed too and can be safely ignored. The selection $M_* = (3-10)M$ is convenient because the states $\mathcal{N}$, $\bar{\mathcal{N}}$ (having the mass M_*) decouple earlier than the states $N_{1,2}$ and will not contribute to the leptogenesis process. With all these we conclude that the results obtained in previous subsections stay robust.

Closing this section, we comment (as was also noted in Sec. II), that throughout our studies we have not attempted to explain and justify texture zeros by symmetries. Our approach here was to consider such textures that give a predictive and consistent scenario allowing us to calculate cosmological *CP* violation. The forms of the matrices in Eqs. (3), (22), and/or (44) with specific coupling selections are such that their structures and the model's predictive power (as was demonstrated) are not ruined by radiative corrections. For our purposes this was already satisfactory. A more fundamental explanation should be pursued elsewhere.

ACKNOWLEDGMENTS

We thank A. Pilaftsis for discussions and helpful comments. The work is partially supported by Shota Rustaveli National Science Foundation (Contract No. DI/12/6-200/13). Z. T. thanks CETUP* (Center for Theoretical Underground Physics and Related Areas) for its hospitality and partial support during the 2016 Summer Program. Z. T. also would like to express a special thanks to the Mainz Institute for Theoretical Physics (MITP) for its hospitality and support.

APPENDIX A: RENORMALIZATION GROUP STUDIES

1. Running of Y_ν, Y_e and M_N matrices and approximation for ξ

RG equations for the charged lepton and neutrino Dirac Yukawa matrices, appearing in the superpotential of Eq. (1), at one-loop order have the forms [27,28]

$$16\pi^2\frac{d}{dt}Y_e = 3Y_eY_e^\dagger Y_e + Y_\nu Y_\nu^\dagger Y_e + Y_e[\text{tr}(3Y_d^\dagger Y_d + Y_e^\dagger Y_e) - c_e^a g_a^2],\quad c_e^a = \left(\frac{9}{5}, 3, 0\right), \tag{A1}$$

$$16\pi^2\frac{d}{dt}Y_\nu = Y_eY_e^\dagger Y_\nu + 3Y_\nu Y_\nu^\dagger Y_\nu + Y_\nu[\text{tr}(3Y_u^\dagger Y_u + Y_\nu^\dagger Y_\nu) - c_\nu^a g_a^2],\quad c_\nu^a = \left(\frac{3}{5}, 3, 0\right). \tag{A2}$$

$g_a = (g_1, g_2, g_3)$ denote gauge couplings of $U(1)_Y$, $SU(2)_w$, and $SU(3)_c$ gauge groups, respectively. Their one-loop RG have forms $16\pi^2\frac{d}{dt}g_a = b_a g_a^3$, with $b_a = (\frac{33}{5}, 1, -3)$, where the hypercharge of $U(1)_Y$ is taken in $SU(5)$ normalization.

The RG for the RHN mass matrix at the two-loop level has the form [28]

$$16\pi^2\frac{d}{dt}M_N = 2M_NY_\nu^\dagger Y_\nu - \frac{1}{8\pi^2}M_N[Y_\nu^\dagger Y_eY_e^\dagger Y_\nu + Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu + Y_\nu^\dagger Y_\nu\text{tr}(3Y_u^\dagger Y_u + Y_\nu^\dagger Y_\nu)] + \frac{1}{8\pi^2}M_NY_\nu^\dagger Y_\nu\left(\frac{3}{5}g_1^2 + 3g_2^2\right) + (\text{transpose}). \tag{A3}$$

Let us start with renormalization of the Y_ν's matrix elements. Ignoring in Eq. (A2) the $\mathcal{O}(Y_\nu^3)$ order entries (which are very small because within our studies $|(Y_\nu)_{ij}| \lesssim 10^{-4}$), and from charged fermion Yukawas keeping λ_τ and λ_t, we will have

$$16\pi^2 \frac{d}{dt}\ln(Y_\nu)_{ij} \simeq \delta_{i3}\lambda_\tau^2 + 3\lambda_t^2 - c_\nu^a g_a^2. \tag{A4}$$

This gives the solution

$$(Y_\nu)_{ij}(\mu) = (Y_{\nu G})_{ij}(\eta_\tau(\mu))^{\delta_{i3}}\eta_t^3(\mu)\eta_{g\nu}(\mu), \tag{A5}$$

where $Y_{\nu G}$ denotes the Yukawa matrix at scale M_G and the scale dependent RG factors are given by

$$\eta_{t,b,\tau}(\mu) = \exp\left(-\frac{1}{16\pi^2}\int_t^{t_G}\lambda_{t,b,\tau}^2(t')dt'\right),$$
$$\eta_a(\mu) = \exp\left(\frac{1}{16\pi^2}\int_t^{t_G} g_a^2(t')dt'\right),$$
$$\eta_{g\nu}(\mu) = \exp\left(\frac{1}{16\pi^2}\int_t^{t_G} c_\nu^a g_a^2(t')dt'\right) = \eta_1^{3/5}(\mu)\eta_2^3(\mu),$$
$$\text{with}\quad t = \ln\mu, \qquad t' = \ln\mu', \qquad t_G = \ln M_G. \tag{A6}$$

From these, for the combination $Y_\nu^\dagger Y_\nu$ at scale $\mu = M$ we get the expression given in Eq. (16).

On the other hand, for the RHN mass splitting and for the phase mismatch [depending on ξ defined in Eq. (17)], the integrals/factors of Eqs. (13), (14), (15), and (16) will be relevant. For obtaining approximate analytical results [for the expression of $\frac{\bar{r}_\tau(M)}{r_\tau(M)}$] we will use expansions. Namely, we introduce the notation

$$\mathcal{K} = \kappa r_\tau\left(1 - \frac{\lambda_\tau^2}{16\pi^2}\right) \tag{A7}$$

and make a Taylor expansion of $\mathcal{K}(t)$ and $\kappa(t)$ near the point $t = t_M$, in powers of $(t - t_M)$. As it turns out, this will allow us to calculate $\xi = \frac{\bar{r}_\tau(M)}{r_\tau(M)} - 1$ in powers of $\frac{\lambda_\tau^2}{16\pi^2}$ (and possibly in powers of other couplings appearing in higher degrees—together with appropriate $\frac{1}{16\pi^2}$ factors). We have

$$\mathcal{K}(t) = \mathcal{K}(t_M) + \mathcal{K}'(t_M)(t - t_M) + \frac{1}{2}\mathcal{K}''(t_M)(t - t_M)^2 + \cdots,$$
$$\kappa(t) = \kappa(t_M) + \kappa'(t_M)(t - t_M) + \frac{1}{2}\kappa''(t_M)(t - t_M)^2 + \cdots, \tag{A8}$$

where primes denote derivatives with respect to t. Plugging these in Eq. (14) and performing integration we will get

$$\bar{r}_\tau(M) = \frac{\mathcal{K}(t_M)}{\kappa(t_M)}\left(1 + \frac{1}{2}\frac{\mathcal{K}'(t_M)}{\mathcal{K}(t_M)}(t_G - t_M) + \frac{1}{6}\frac{\mathcal{K}''(t_M)}{\mathcal{K}(t_M)}(t_G - t_M)^2 + \cdots\right)$$
$$\times\left(1 + \frac{1}{2}\frac{\kappa'(t_M)}{\kappa(t_M)}(t_G - t_M) + \frac{1}{6}\frac{\kappa''(t_M)}{\kappa(t_M)}(t_G - t_M)^2 + \cdots\right)^{-1}. \tag{A9}$$

Using in (A9) expression (A7) for $\mathcal{K}$ and keeping in expansion terms up to $(t - t_M)^2$, we get

$$\frac{\bar{r}_\tau(M)}{r_\tau(M)} - 1 \simeq \frac{1}{2}\frac{r_\tau'}{r_\tau}\bigg|_{t=t_M}(t_G - t_M) + \frac{1}{6}\left(\frac{r_\tau''}{r_\tau} + \frac{1}{2}\frac{\kappa' r_\tau'}{\kappa r_\tau}\right)_{t=t_M}(t_G - t_M)^2 - \frac{\lambda_\tau^2(M)}{16\pi^2}. \tag{A10}$$

As we see, the flavor universal RG factor κ drops out at first order of $(t_G - t_M)$. The last term in Eq. (A10) is due to the two-loop correction in the RG of M_N [in particular the $M_N Y_\nu^\dagger Y_e Y_e^\dagger Y_\nu$ term of the RHS of Eq. (A3)]. The remaining terms are due to one-loop corrections, proving that cosmological CP violation emerges already at the one-loop level.

Using in (A10) expressions for the scale factors given in Eqs. (A6) and (15), using the RG for λ_τ [easily obtained from Eq. (A1)], and keeping terms up to the order of $\frac{1}{(16\pi^2)^2}$, we obtain the expression for ξ given in Eq. (18).

2. Neutrino mass matrix renormalization

In the energy interval $M_S \le \mu < M$ (where M_S is the SUSY scale) the RG for the neutrino mass matrix is [28,29]

$$M_S \le \mu < M\colon\ 16\pi^2\frac{d}{dt}M_\nu = Y_e Y_e^\dagger M_\nu + M_\nu Y_e^* Y_e^T + M_\nu[6\mathrm{tr}(Y_u^\dagger Y_u) - 2c_\nu^a g_a^2]. \tag{A11}$$

Below the M_S scale, effectively we have SM and the RG is [29]

$$\mu < M_S:\ 16\pi^2 \frac{d}{dt} M_\nu = \frac{1}{2} Y_e Y_e^\dagger M_\nu + \frac{1}{2} M_\nu Y_e^* Y_e^T + M_\nu [\mathrm{tr}(6Y_u^\dagger Y_u + 6Y_d^\dagger Y_d + 2Y_e^\dagger Y_e) - 3g_2^2 + 4\lambda], \tag{A12}$$

where λ is the SM Higgs self-coupling [emerging from the self-interaction term $\lambda(H^\dagger H)^2$ of the SM Higgs doublet H]. We will also need the RG evaluation of the VEVs v_u and v, which in appropriate energy intervals are given by [30–33]

$$\mu > M_S:\ 16\pi^2 \frac{d}{dt} v_u = v_u \left(-3\lambda_t^2 + \frac{1}{4} c_\nu^a g_a^2\right), \tag{A13}$$

$$\mu < M_S:\ 16\pi^2 \frac{d}{dt} v = v \left(-3\lambda_t^2 - 3\lambda_b^2 - \lambda_\tau^2 + \frac{3}{4} c_\nu^a g_a^2\right). \tag{A14}$$

At scale M, after decoupling of the RHN states, the neutrino mass matrix is formed with the form

$$M_\nu^{ij}(M) = \begin{pmatrix} \times & \times & \times \\ \times & \times & \times \\ \times & \times & \times \end{pmatrix} \frac{v_u^2(M)}{M}, \tag{A15}$$

where $\times$ stands for entries depending on Yukawa couplings. After renormalization, keeping λ_τ, λ_t, and g_a in the above RGs, for the neutrino mass matrix at scale M_Z we obtain

$$M_\nu^{ij}(M_Z) = \begin{pmatrix} \times & \times & (\times)\cdot r_{\nu 3} \\ \times & \times & (\times)\cdot r_{\nu 3} \\ (\times)\cdot r_{\nu 3} & (\times)\cdot r_{\nu 3} & (\times)\cdot r_{\nu 3}^2 \end{pmatrix} \bar{m},$$
$$\text{with}\quad \bar{m} = \frac{v^2(M_Z) s_\beta^2}{M} r_{\bar{m}}, \tag{A16}$$

where $\times$ denotes entries determined at scale M corresponding to those in (A15), and RG factors are given by

$$r_{\nu 3} = \left(\frac{\eta_\tau(t_Z)}{\eta_\tau(t_{M_S})}\right)^{1/2} \left(\frac{\eta_\tau(t_{M_S})}{\eta_\tau(t_M)}\right), \tag{A17}$$

$$r_{\bar{m}} = \eta_\lambda^4 \left(\frac{\eta_t(t_{m_t})}{\eta_t(t_M)}\right)^{12} \left(\frac{\eta_b(t_Z)}{\eta_b(t_{M_S})}\right)^{12} \left(\frac{\eta_\tau(t_Z)}{\eta_\tau(t_{M_S})}\right)^4 \left(\frac{\eta_2(t_Z)}{\eta_2(t_M)}\right)^{\frac{15}{2}} \times \left(\frac{\eta_1^{3/5}(t_Z)\eta_1^{2/5}(t_{M_S})}{\eta_1(t_M)}\right)^{\frac{3}{2}}, \tag{A18}$$

where

$$\eta_\lambda = \exp\left(-\frac{1}{16\pi^2}\int_{t_{m_h}}^{t_{M_S}} \lambda(t)dt\right), \tag{A19}$$

and remaining η factors are defined in Eq. (A6).

We will also need the RG factor relating the VEV $v_u(M)$ to the $v(M_Z)$. Using Eqs. (A13) and (A14) we obtain

$$r_{v_u} = \frac{v_u(M)}{v(M_Z)s_\beta} = \left(\frac{\eta_t(t_{m_t})}{\eta_t(t_M)}\right)^3 \left(\frac{\eta_b(t_Z)}{\eta_b(t_{M_S})}\right)^3 \left(\frac{\eta_\tau(t_Z)}{\eta_\tau(t_{M_S})}\right) \times \left(\frac{\eta_2^3(t_Z)\eta_2^{-2}(t_{M_S})}{\eta_2(t_M)}\right)^{\frac{3}{4}} \left(\frac{\eta_1^3(t_Z)\eta_1^{-2}(t_{M_S})}{\eta_1(t_M)}\right)^{\frac{3}{20}}. \tag{A20}$$

3. Boundary and matching conditions

For finding the RG factors, appearing in the baryon asymmetry, we numerically solve renormalization group equations from the scale M_Z up to the $M_G \simeq 2\times 10^{16}$ GeV scale. For simplicity, for all SUSY particle masses we take the common mass scale M_S. Thus, in the energy interval $M_Z \le \mu < M_S$, the Standard Model RGs for $\overline{\mathrm{MS}}$ coupling constants are used. However, in the interval $M_S \le \mu \le M_G$, since we are dealing with the SUSY, the RGs for the $\overline{\mathrm{DR}}$ couplings are applied. Below we give boundary and matching conditions for the gauge couplings $g_{1,2,3}$, for Yukawa constant $\lambda_{t,b,\tau}$ and for the Higgs self-coupling λ.

a. Gauge couplings

We choose our inputs for the $\overline{\mathrm{MS}}$ gauge couplings at scale M_Z as follows:

$$\alpha_1^{-1}(M_Z) = \frac{3}{5} c_w^2 \alpha_{em}^{-1}(M_Z) + \frac{3}{5} c_w^2 \frac{8}{9\pi} \ln\frac{m_t}{M_Z},$$
$$\alpha_2^{-1}(M_Z) = s_w^2 \alpha_{em}^{-1}(M_Z) + s_w^2 \frac{8}{9\pi} \ln\frac{m_t}{M_Z},$$
$$\alpha_3^{-1}(M_Z) = \alpha_s^{-1}(M_Z) + \frac{1}{3\pi} \ln\frac{m_t}{M_Z}, \tag{A21}$$

where logarithmic terms $\ln\frac{m_t}{M_Z}$ are due to the top quark threshold correction [32,34]. Taking $\alpha_s(M_Z) = 0.1185$, $\alpha_{em}^{-1}(M_Z) = 127.934$, and $s_w^2 = 0.2313$, from (A21) we obtain

$$\alpha_1^{-1}(M_Z) = 59.0057 + \frac{8c_w^2}{15\pi} \ln\frac{m_t}{M_Z},$$
$$\alpha_2^{-1}(M_Z) = 29.5911 + \frac{8s_w^2}{9\pi} \ln\frac{m_t}{M_Z},$$
$$\alpha_3^{-1}(M_Z) = 8.4388 + \frac{1}{3\pi} \ln\frac{m_t}{M_Z}. \tag{A22}$$

With these inputs we run $g_{1,2,3}$ via the two-loop RGs from M_Z up to the scale M_S.

At scale $\mu = M_S$ we use the matching conditions between $\bar{\mathrm{DR}} - \overline{\mathrm{MS}}$ gauge couplings [35,36],

$$\text{at } \mu = M_S: \ \frac{1}{\alpha_1^{\bar{\mathrm{DR}}}} = \frac{1}{\alpha_1^{\overline{\mathrm{MS}}}}, \qquad \frac{1}{\alpha_2^{\bar{\mathrm{DR}}}} = \frac{1}{\alpha_2^{\overline{\mathrm{MS}}}} - \frac{1}{6\pi}, \\ \frac{1}{\alpha_3^{\bar{\mathrm{DR}}}} = \frac{1}{\alpha_3^{\overline{\mathrm{MS}}}} - \frac{1}{4\pi}. \tag{A23}$$

Above the scale M_S we apply two-loop SUSY RG equations in the $\bar{\mathrm{DR}}$ scheme [27].

b. Yukawa couplings and λ

At the scale M_S all SUSY states decouple, and we are left with the Standard Model with one Higgs doublet. Thus, the third family Yukawa couplings and the self-coupling are determined as

$$\lambda_t(m_t) = \frac{m_t(m_t)}{v(m_t)}, \qquad \lambda_b(M_Z) = \frac{2.89\ \mathrm{GeV}}{v(M_Z)}, \\ \lambda_\tau(M_Z) = \frac{1.746\ \mathrm{GeV}}{v(M_Z)}, \\ \lambda(m_h) = \frac{1}{4}\left(\frac{m_h}{v(m_h)}\right)^2, \quad \text{with} \quad v(M_Z) = 174.1\ \mathrm{GeV}, \\ m_h = 125.15\ \mathrm{GeV}, \tag{A24}$$

where $m_t(m_t)$ is the top quark running mass related to the pole mass as

$$m_t(m_t) = p_t M_t^{\mathrm{pole}}. \tag{A25}$$

The factor p_t is $p_t \simeq 1/1.0603$ [37], while the recent measured value of the top's pole mass is [38]

$$M_t^{\mathrm{pole}} = (173.34 \pm 0.76)\ \mathrm{GeV}. \tag{A26}$$

We take the values of (A24) as boundary conditions for solving two-loop RG equations [32,39] for $\lambda_{t,b,\tau}$ and λ from the M_Z scale up to the scale M_S.

Above the M_S scale, we have MSSM states including two doublets h_u and h_d, which couple with up-type quarks and down-type quarks and charged leptons, respectively. Thus, the third family Yukawa couplings at M_S are $\approx \lambda_t(M_S)/s_\beta, \lambda_b(M_S)/c_\beta$ and $\lambda_\tau(M_S)/c_\beta$, with $s_\beta \equiv \sin\beta$, $c_\beta \equiv \cos\beta$. Above the scale M_S we apply two-loop SUSY RG equations in the $\bar{\mathrm{DR}}$ scheme [27]. Thus, at $\mu = M_S$ we use the matching conditions between $\bar{\mathrm{DR}} - \overline{\mathrm{MS}}$ couplings,

$$\text{at } \mu = M_S: \ \lambda_t^{\bar{\mathrm{DR}}} \simeq \frac{\lambda_t^{\overline{\mathrm{MS}}}}{s_\beta}\left[1 + \frac{1}{16\pi^2}\left(\frac{g_1^2}{120} + \frac{3g_2^2}{8} - \frac{4g_3^2}{3}\right)\right], \\ \lambda_b^{\bar{\mathrm{DR}}} \simeq \frac{\lambda_b^{\overline{\mathrm{MS}}}}{c_\beta}\left[1 + \frac{1}{16\pi^2}\left(\frac{13g_1^2}{120} + \frac{3g_2^2}{8} - \frac{4g_3^2}{3}\right)\right], \\ \lambda_\tau^{\bar{\mathrm{DR}}} \simeq \frac{\lambda_\tau^{\overline{\mathrm{MS}}}}{c_\beta}\left[1 + \frac{1}{16\pi^2}\left(-\frac{9g_1^2}{40} + \frac{3g_2^2}{8}\right)\right], \tag{A27}$$

where expressions in brackets on the RHS of the relations are due to the $\bar{\mathrm{DR}} - \overline{\mathrm{MS}}$ conversions [36]. With Eq. (A27)'s matchings we run corresponding couplings from the scale M_S up to the M_G scale. Throughout the paper, above the mass scale M_S without using the superscript $\bar{\mathrm{DR}}$ we assume the couplings determined in this scheme.

APPENDIX B: CONTRIBUTION TO THE BARYON ASYMMETRY FROM $\tilde{N}$ DECAYS

The impact of the decays of the right-handed sneutrinos—the scalar partners of the RHNs—was estimated in [11] for specific textures. Here we give a more detailed investigation and give results for the neutrino model discussed in Sec. III A.

We will need to derive masses of the RH sneutrinos and their couplings to the components of the superfields l and h_u. For this purpose, we should include the soft breaking terms

$$V_{SB}^\nu = \tilde{l}^T A_\nu \tilde{N} h_u - \frac{1}{2}\tilde{N}^T B_N \tilde{N} + \mathrm{H.c.} + \tilde{l}^\dagger m_{\tilde{l}}^2 \tilde{l} + \tilde{N}^\dagger m_{\tilde{N}}^2 \tilde{N}, \tag{B1}$$

which, together with the superpotential couplings, will be relevant. As it turns out, A_ν and B_N couplings will be relevant. Therefore, first we will study their renormalization. After this, we investigate masses of the physical RH sneutrinos and their couplings to the lepton superfield components. These, at the end, will be used for the calculation of the contribution in the baryon asymmetry via the RH sneutrino decay processes.

1. Renormalization of soft A_ν and B_N terms

From general expressions of Ref. [27] we can derive RGs for A_ν and B_N, which at the one-loop level have the forms

$$16\pi^2 \frac{d}{dt}A_\nu = Y_e Y_e^\dagger A_\nu + 2\hat{A}_e Y_e^\dagger Y_\nu + 5Y_\nu Y_\nu^\dagger A_\nu \\ + A_\nu[\mathrm{tr}(3Y_u^\dagger Y_u + Y_\nu^\dagger Y_\nu) + 4Y_\nu^\dagger Y_\nu - c_\nu^a g_a^2] \\ + 2Y_\nu[\mathrm{tr}(3Y_u^\dagger \hat{A}_u + Y_\nu^\dagger A_\nu) + c_\nu^a g_a^2 M_{\tilde{V}_a}], \tag{B2}$$

$$16\pi^2 \frac{d}{dt}B_N = 2B_N Y_\nu^\dagger Y_\nu + 2Y_\nu^T Y_\nu^* B_N + 4M_N Y_\nu^\dagger A_\nu \\ + 4A_\nu^T Y_\nu^* M_N. \tag{B3}$$

Note that, applying these expressions for the third generation states, we can get expressions of [40] [see Eqs. (17) and (55) of this reference, which uses slightly different definitions for the couplings]. These results are also compatible with those given in [41] (with replacements $Y \to Y^T$, $A \to A^T$).

We parametrize the matrix B_N as

$$B_N = (M_N)_{12} m_B \begin{pmatrix} \delta_{BN}^{(1)} & 1 \\ 1 & \delta_{BN}^{(2)} \end{pmatrix}, \qquad \text{(B4)}$$

where all entries $(M_N)_{12}, m_B, \delta_{BN}^{(1,2)}$ run and their RGs can be derived from the RG equations given above. For the matrix A_ν, let us use the parametrization

$$A_\nu = m_A a_\nu, \qquad \text{(B5)}$$

where m_A is a constant and the elements of the a_ν matrix run. The matrix $\hat{A}_e$ is

$$\hat{A}_e = \mathrm{Diag}(A_e, A_\mu, A_\tau) \qquad \text{(B6)}$$

(similar to the structure of the Y_e Yukawa matrix). We will use the following boundary conditions:

$$\begin{aligned} &\text{at } \mu = M_G:\ a_\nu = Y_\nu, \qquad \delta_{BN}^{(1)} = \delta_{BN}^{(2)} = 0, \\ &\hat{A}_e = m_A \mathrm{Diag}(\lambda_e, \lambda_\mu, \lambda_\tau), \\ &\hat{A}_u = m_A Y_{uG}, \qquad \hat{A}_d = m_A Y_{dG}, \end{aligned} \qquad \text{(B7)}$$

which assume proportionality (alignment) of the soft SUSY breaking terms with the corresponding superpotential couplings.

With (B4), (B5), using (B3) we have

$$\begin{aligned} 16\pi^2 \frac{d}{dt}\delta_{BN}^{(1)} &\simeq 4(Y_\nu^\dagger Y_\nu)_{21} + 8\frac{m_A}{m_B}(Y_\nu^\dagger a_\nu)_{21}, \\ 16\pi^2 \frac{d}{dt}\delta_{BN}^{(2)} &\simeq 4(Y_\nu^\dagger Y_\nu)_{12} + 8\frac{m_A}{m_B}(Y_\nu^\dagger a_\nu)_{12}. \end{aligned} \qquad \text{(B8)}$$

Because of RG effects, the alignment between Y_ν and a_ν (which holds at the GUT scale) is violated. In particular,

$$16\pi^2 \frac{d}{dt}\left(\frac{(a_\nu)_{ij}}{(Y_\nu)_{ij}}\right) \simeq 2\delta_{i3}\frac{\lambda_\tau A_\tau}{m_A} + \frac{2}{m_A}(3\lambda_t A_t + c_\nu^a g_a^2 M_{\tilde{V}_a}), \qquad \text{(B9)}$$

where on the RHS we kept third family couplings, gauge couplings, and gaugino masses. From this we derive

$$\begin{aligned} a_\nu &\simeq \begin{pmatrix} 1+\epsilon_0 & 0 & 0 \\ 0 & 1+\epsilon_0 & 0 \\ 0 & 0 & 1+\epsilon_0+\epsilon \end{pmatrix} Y_\nu \\ &\quad \text{with} \quad \epsilon_0 = -\frac{1}{8\pi^2 m_A}\int_t^{t_G} dt(3\lambda_t A_t + c_\nu^a g_a^2 M_{\tilde{V}_a}), \\ \epsilon &= -\frac{1}{8\pi^2 m_A}\int_t^{t_G} dt \lambda_\tau A_\tau. \end{aligned} \qquad \text{(B10)}$$

Using (B10) in Eqs. (B8) and (B4) we obtain[7]

$$\begin{aligned} &\text{at } \mu = M: B_N = m_B M \begin{pmatrix} -\alpha\delta_N(1+\bar{\epsilon}_1) & 1 \\ 1 & -\alpha\delta_N^*(1+\bar{\epsilon}_2) \end{pmatrix}, \\ &\alpha = 1 + 2\frac{m_A}{m_B} \end{aligned} \qquad \text{(B11)}$$

and

$$\begin{aligned} \bar{\epsilon}_1 &= \frac{1}{4\pi^2 \alpha\delta_N}\int_{t_M}^{t_G} dt \left(Y_\nu^\dagger\left(\frac{\alpha}{16\pi^2}Y_e Y_e^\dagger + 2\frac{m_A}{m_B}\hat{\epsilon}\right)Y_\nu\right)_{21}, \\ \bar{\epsilon}_2 &= \frac{1}{4\pi^2 \alpha\delta_N^*}\int_{t_M}^{t_G} dt \left(Y_\nu^\dagger\left(\frac{\alpha^*}{16\pi^2}Y_e Y_e^\dagger + 2\frac{m_A^*}{m_B^*}\hat{\epsilon}^*\right)Y_\nu\right)_{21}^*, \\ &\quad \text{with} \quad \hat{\epsilon} = \mathrm{Diag}(\epsilon_0, \epsilon_0, \epsilon_0 + \epsilon). \end{aligned} \qquad \text{(B12)}$$

The form of B_N given in Eq. (B11) will be needed to construct the sneutrino mass matrix, which we will do below.

2. Sneutrino mass matrix and its diagonalization

For calculating scalar RHN masses, from (B1) we keep only the B_N term. We also include the mass^2 term $\tilde{N}^\dagger M_N^\dagger M_N \tilde{N}$ coming from the superpotential. Therefore, we consider the following quadratic potential:

$$V_{\tilde{N}}^{(2)} = \tilde{N}^\dagger M_N^\dagger M_N \tilde{N} - \left(\frac{1}{2}\tilde{N}^T B_N \tilde{N} + \mathrm{H.c.}\right). \qquad \text{(B13)}$$

With the transformation of the N superfields $N = U_N N'$ [according to Eq. (6), the U_N diagonalizes the fermionic RHN mass matrix], we obtain

$$V_{\tilde{N}}^{(2)} = \tilde{N}'^\dagger (M_N^{\mathrm{Diag}})^2 \tilde{N}' - \left(\frac{1}{2}\tilde{N}'^T U_N^T B_N U_N \tilde{N}' + \mathrm{H.c.}\right). \qquad \text{(B14)}$$

On the other hand, from (B11) we have

[7] Since in the β functions we are ignoring Y_ν couplings (due to their smallness), for all practical purposes the m_B can be treated as a constant.

$$U_N^T B_N U_N = m_B|M|\begin{pmatrix} 1-\tilde{\alpha}|\delta_N| & \frac{i}{2}\alpha|\delta_N|(\bar{\epsilon}_1-\bar{\epsilon}_2) \\ \frac{i}{2}\alpha|\delta_N|(\bar{\epsilon}_1-\bar{\epsilon}_2) & 1+\tilde{\alpha}|\delta_N| \end{pmatrix}, \quad \text{with} \quad \tilde{\alpha} = \alpha\left(1+\frac{\bar{\epsilon}_1+\bar{\epsilon}_2}{2}\right). \tag{B15}$$

With further phase redefinition

$$\tilde{N}' = \tilde{P}_1\tilde{N}'', \qquad \tilde{P}_1 = \mathrm{Diag}(e^{-i\tilde{\omega}_1/2}, e^{-i\tilde{\omega}_2/2}), \quad \text{with} \quad \tilde{\omega}_{1,2} = \mathrm{Arg}[m_B(1\mp\tilde{\alpha}|\delta_N|)], \tag{B16}$$

and by going to the real scalar components

$$\tilde{N}_1'' = \frac{1}{\sqrt{2}}(\tilde{N}_1^R + i\tilde{N}_1^I), \qquad \tilde{N}_2'' = \frac{1}{\sqrt{2}}(\tilde{N}_2^R + i\tilde{N}_2^I), \tag{B17}$$

we will have

$$\begin{aligned} &-\left(\frac{1}{2}\tilde{N}'^T U_N^T B_N U_N \tilde{N}' + \mathrm{H.c.}\right) = -\frac{|Mm_B|}{2}|1-\tilde{\alpha}|\delta_N||((\tilde{N}_1^R)^2 - (\tilde{N}_1^I)^2) \\ &\quad -\frac{|Mm_B|}{2}|1+\tilde{\alpha}|\delta_N||((\tilde{N}_2^R)^2 - (\tilde{N}_2^I)^2) - |M|\mathrm{Re}(m_B\delta_\epsilon)(\tilde{N}_1^R\tilde{N}_2^R - \tilde{N}_1^I\tilde{N}_2^I) + |M|\mathrm{Im}(m_B\delta_\epsilon)(\tilde{N}_1^I\tilde{N}_2^R + \tilde{N}_1^R\tilde{N}_2^I) \\ &\quad \text{with} \quad \delta_\epsilon = i\alpha|\delta_N|\frac{\bar{\epsilon}_1-\bar{\epsilon}_2}{2}e^{-i(\tilde{\omega}_1+\tilde{\omega}_2)/2}. \end{aligned} \tag{B18}$$

From (B14) and (B18) we obtain the mass^2 terms,

$$V_{\tilde{N}}^{(2)} = \frac{1}{2}\tilde{n}^{0T}M_{\tilde{n}}^2\tilde{n}^0, \quad \text{with} \quad \tilde{n}^{0T} = (\tilde{N}_1^R, \tilde{N}_1^I, \tilde{N}_2^R, \tilde{N}_2^I) \tag{B19}$$

and

$$M_{\tilde{n}}^2 = \begin{pmatrix} (\tilde{M}_1^0)^2 & 0 & -|M|\mathrm{Re}(m_B\delta_\epsilon) & |M|\mathrm{Im}(m_B\delta_\epsilon) \\ 0 & (\tilde{M}_2^0)^2 & |M|\mathrm{Im}(m_B\delta_\epsilon) & |M|\mathrm{Re}(m_B\delta_\epsilon) \\ -|M|\mathrm{Re}(m_B\delta_\epsilon) & |M|\mathrm{Im}(m_B\delta_\epsilon) & (\tilde{M}_3^0)^2 & 0 \\ |M|\mathrm{Im}(m_B\delta_\epsilon) & |M|\mathrm{Re}(m_B\delta_\epsilon) & 0 & (\tilde{M}_4^0)^2 \end{pmatrix}, \tag{B20}$$

where

$$\begin{aligned} &(\tilde{M}_1^0)^2 = |M|^2(1-|\delta_N|)^2 - |m_BM||1-\tilde{\alpha}|\delta_N||, \qquad (\tilde{M}_2^0)^2 = |M|^2(1-|\delta_N|)^2 + |m_BM||1-\tilde{\alpha}|\delta_N||, \\ &(\tilde{M}_3^0)^2 = |M|^2(1+|\delta_N|)^2 - |m_BM||1+\tilde{\alpha}|\delta_N||, \qquad (\tilde{M}_4^0)^2 = |M|^2(1+|\delta_N|)^2 + |m_BM||1+\tilde{\alpha}|\delta_N||. \end{aligned} \tag{B21}$$

The coupling of $\tilde{n}^0$ states with the fermions emerges from the F term of the superpotential $l^TY_\nu Nh_u$. Following the transformations, indicated above, we will have

$$\begin{aligned} (l^TY_\nu Nh_u)_F &\to \tilde{h}_u l^T Y_\nu \tilde{N} \\ &= e^{-i\tilde{\omega}_2/2}\tilde{h}_u l^T Y_\nu U_N(\rho_u e^{i(\tilde{\omega}_2-\tilde{\omega}_1)/2}, \rho_d)\tilde{n}^0, \\ &\quad \text{with} \quad \rho_u = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 & i \\ 0 & 0 \end{pmatrix}, \\ \rho_d &= \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & 0 \\ 1 & i \end{pmatrix}. \end{aligned} \tag{B22}$$

Performing the diagonalization of the matrix (B20) by the transformation $V_{\tilde{n}}^T M_{\tilde{n}}^2 V_{\tilde{n}} = (M_{\tilde{n}}^{\mathrm{Diag}})^2$, $\tilde{n}^0 = V_{\tilde{n}}\tilde{n}$, the fermion coupling with the scalar $\tilde{n}$ eigenstates will be

$$\begin{aligned} &\tilde{h}_u l^T Y_F \tilde{n} \quad \text{with} \quad Y_F = Y_\nu \tilde{V}^0 V_{\tilde{n}}, \\ &\tilde{V}^0 = U_N(\rho_u e^{-i\tilde{\omega}_1/2}, \rho_d e^{-i\tilde{\omega}_2/2}). \end{aligned} \tag{B23}$$

The coupling with the slepton $\tilde{l}$ is derived from the interaction term $h_u\tilde{l}^T(Y_\nu M_N^*\tilde{N}^* - A_\nu\tilde{N})$. Going from $\tilde{N}$ to the $\tilde{n}$ states, we obtain

$$h_u \tilde{l}^T Y_B \tilde{n} \quad \text{with} \quad Y_B = (Y_\nu M_N^* \tilde{V}^{0*} - A_\nu \tilde{V}^0) V_{\tilde{n}}. \tag{B24}$$

For given values of M, m_B, and m_A, with the help of Eqs. (B20), (B23), and (B24), we will have coupling matrices Y_F, Y_B and all other quantities needed for calculation of the baryon asymmetry created via the decays of the $\tilde{n}_{1,2,3,4}$ states.

3. Asymmetry via $\tilde{n}$ decays

Now we are ready to discuss the contribution to the net baryon asymmetry from the out of equilibrium resonant decays of the right-handed sneutrinos (RHSN). As we have seen, with SUSY breaking terms, the masses of RHSN's differ from their fermionic partners' masses. Thus we have mass-eigenstate RHSN's $\tilde{n}_{i=1,2,3,4}$ with masses $\tilde{M}_{i=1,2,3,4}$, respectively. With the SUSY scale M_S smaller (at least by a factor of 3) than the scale M, the states $\tilde{n}_i$ remain nearly degenerate.

For the resonant $\tilde{n}$ decays we will apply the resummed effective amplitude technique [17]. Effective amplitudes for the real $\tilde{n}_i$ decay, say into the lepton l_α ($\alpha = 1$, 2, 3) and antilepton $\bar{l}_\alpha$, respectively, are given by [17]

$$\hat{S}_{\alpha i} = S_{\alpha i} - \sum_j S_{\alpha j} \frac{\Pi_{ji}(\tilde{M}_i)(1-\delta_{ij})}{\tilde{M}_i^2 - \tilde{M}_j^2 + \Pi_{jj}(\tilde{M}_i)},$$
$$\hat{\bar{S}}_{\alpha i} = S^*_{\alpha i} - \sum_j S^*_{\alpha j} \frac{\Pi_{ji}(\tilde{M}_i)(1-\delta_{ij})}{\tilde{M}_i^2 - \tilde{M}_j^2 + \Pi_{jj}(\tilde{M}_i)}, \tag{B25}$$

where $S_{\alpha i}$ is a tree level amplitude and Π_{ij} is a two point Green function's (polarization operator of $\tilde{n}_i - \tilde{n}_j$) absorptive part. The *CP* asymmetry is then given by

$$\epsilon_i^{sc} = \frac{\sum_\alpha (|\hat{S}_{\alpha i}|^2 - |\hat{\bar{S}}_{\alpha i}|^2)}{\sum_\alpha (|\hat{S}_{\alpha i}|^2 + |\hat{\bar{S}}_{\alpha i}|^2)}. \tag{B26}$$

With Y_F and Y_B given by Eq. (B23) and (B24) we can calculate the polarization diagram's (with external legs $\tilde{n}_i$ and $\tilde{n}_j$) absorptive part Π_{ij}, which at the one-loop level is given by

$$\Pi_{ij}(p) = \frac{i}{8\pi}(p^2 Y_F^\dagger Y_F + p^2 Y_F^T Y_F^* + Y_B^\dagger Y_B + Y_B^T Y_B^*)_{ij}, \tag{B27}$$

where p denotes the external momentum in the diagram and upon evaluation of (B26), for Π we should use (B27) with $p = \tilde{M}_i$.

In an unbroken SUSY limit, neglecting finite temperature effects ($T \to 0$), the $\tilde{N}$ decay does not produce lepton asymmetry due to the following reason. The decays of $\tilde{N}$ in the fermion and scalar channels are, respectively, $\tilde{N} \to l\tilde{h}_u$ and $\tilde{N} \to \tilde{l}^* h_u^*$. Since the rates of these processes are the same due to SUSY (at $T = 0$), the lepton asymmetries created from these decays cancel each other. With $T \neq 0$, the cancellation does not take place and one has

$$\tilde{\epsilon}_i = \epsilon_i(\tilde{n}_i \to l\tilde{h}_u)\Delta_{BF}, \tag{B28}$$

with a temperature dependent factor Δ_{BF} given in [42].[8] Therefore, we just need to compute $\epsilon_i(\tilde{n}_i \to l\tilde{h}_u)$, which is the asymmetry created by $\tilde{n}_i$ decays in two fermions. Thus, in (B25) we take $S_{\alpha i} = (Y_F)_{\alpha i}$ and calculate $\epsilon_i(\tilde{n}_i \to l\tilde{h}_u)$ with (B26). The baryon asymmetry created from the lepton asymmetry due to $\tilde{n}$ decays is

$$\frac{\tilde{n}_b}{s} \simeq -8.46 \times 10^{-4} \sum_{i=1}^4 \frac{\tilde{\epsilon}_i}{\Delta_{BF}} \eta_i$$
$$= -8.46 \times 10^{-4} \sum_{i=1}^4 \epsilon_i(\tilde{n}_i \to l\tilde{h}_u)\eta_i, \tag{B29}$$

where an effective number of degrees of freedom (including two RHN superfields) $g_* = 228.75$ was used. η_i are efficiency factors that depend on $\tilde{m}_i \simeq \frac{(v \sin\beta)^2}{M} 2(Y_F^\dagger Y_F)_{ii}$ and take into account temperature effects by integrating the Boltzmann equations [42].

In Table IV we give results for the neutrino model discussed in Sec. III A. These are obtained for the SUSY particle masses $= M_S$ and for the different values of pairs (m_A, m_B) (see also the caption of Table IV). Upon the

TABLE IV. Values of $\frac{\Delta n_b}{s} = \frac{\tilde{n}_b}{s}$—contributions to the baryon asymmetry via decays of the right-handed sneutrinos for cases given in Table III [i.e., for values of a_2 giving $(\frac{n_b}{s})_{\max}$ given in Table III]. These values correspond to the phases $\delta = -0.378$ and $\phi_+ = -1.287$.

	$(m_A, m_B) = (100i, 500)$ GeV		$(m_A, m_B) = (500, 1000)$ GeV	
Case	$10^4 \times a_2$	$10^{11} \times \frac{\tilde{n}_b}{s}$	$10^4 \times a_2$	$10^{11} \times \frac{\tilde{n}_b}{s}$
(**I**$_-$)	0.016	0.25	0.016	0.24
(**I**.**1**)	0.0159785	0.25	0.0159785	0.25
(**I**.**2**)	0.0299	0.24	0.0299	0.24
(**I**.**3**)	0.0987	0.24	0.0987	0.24
(**I**.**4**)	0.3237	0.24	0.3237	0.24
(**I**.**5**)	1.05655	0.23	1.05655	0.23
(**II**$_-$)	0.0229	0.25	0.0229	0.24
(**II**.**1**)	0.0229	0.25	0.0229	0.24
(**II**.**2**)	0.02986	0.24	0.02986	0.24
(**II**.**3**)	0.09835	0.24	0.09835	0.24
(**II**.**4**)	0.322	0.24	0.322	0.24
(**II**.**5**)	1.05	0.23	1.05	0.23

[8]This expression is valid with alignment $A_\nu = m_A Y_\nu$, which we are assuming at the GUT scale, and thus Eq. (B28) can be well applicable for our estimates.

calculations, with obtained values of $\tilde{m}_i$, according to Ref. [42] we picked up the corresponding values of η_i and used them in (B29). From Table IV we see that a contribution to the net baryon asymmetry from the RHS decays is suppressed $\frac{\tilde{n}_b}{n_b} < 3 \times 10^{-2}$, i.e., is less than 3%. From Table IV we also see that the complexity of m_A practically does not change the results. This happens because the m_A in the Y_B coupling matrix appears in front of the Y_ν [see Eq. (B24)], which is strongly suppressed. From the structure of (B20), one can also make sure that the complexity of m_B will not affect the results. We have checked this by varying the phases of m_B. For instance, for case (I.1) and $m_A = 500$ GeV, $m_B = 1000 \times \{e^{i\pi/3}, e^{i\pi/2}, e^{2i\pi/3}\}$ GeV we have obtained $\frac{\tilde{n}_b}{s} \simeq 0.24 \times 10^{-11}$. Suppression of $\frac{\tilde{n}_b}{s}$ will always happen for the value of $|m_B|$ in the range of a few 100 GeV to a few TeV, because the mass degeneracy of $\tilde{n}_i$ states is lifted in such a way that resonant enhancement of $\frac{\tilde{n}_b}{s}$ does not happen. (This is unlike the case of soft leptogenesis [42], which requires $|m_B| \lesssim 10$ MeV. Considering the latter value unnatural, we did not pursue such possibilities within our studies.)

[1] M. C. Gonzalez-Garcia, M. Maltoni, and T. Schwetz, J. High Energy Phys. 11 (2014) 052.
[2] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).
[3] G. F. Giudice, A. Notari, M. Raidal, A. Riotto, and A. Strumia, Nucl. Phys. **B685**, 89 (2004).
[4] W. Buchmuller, P. Di Bari, and M. Plumacher, Ann. Phys. (Amsterdam) **315**, 305 (2005).
[5] S. Davidson, E. Nardi, and Y. Nir, Phys. Rep. **466**, 105 (2008).
[6] P. H. Frampton, S. L. Glashow, and T. Yanagida, Phys. Lett. B **548**, 119 (2002).
[7] A. Ibarra and G. G. Ross, Phys. Lett. B **591**, 285 (2004); S. Pascoli, S. T. Petcov, and A. Riotto, Nucl. Phys. **B774**, 1 (2007).
[8] Q. Shafi and Z. Tavartkiladze, Nucl. Phys. **B772**, 133 (2007).
[9] G. C. Branco, A. J. Buras, S. Jager, S. Uhlig, and A. Weiler, J. High Energy Phys. 09 (2007) 004.
[10] K. S. Babu, A. G. Bachri, and Z. Tavartkiladze, Int. J. Mod. Phys. A **23**, 1679 (2008).
[11] K. S. Babu, Y. Meng, and Z. Tavartkiladze, arXiv: 0812.4419.
[12] A. Meroni, E. Molinaro, and S. T. Petcov, Phys. Lett. B **710**, 435 (2012).
[13] K. Harigaya, M. Ibe, and T. T. Yanagida, Phys. Rev. D **86**, 013002 (2012).
[14] S. F. Ge, H. J. He, and F. R. Yin, J. Cosmol. Astropart. Phys. 05 (2010) 017.
[15] A. Achelashvili and Z. Tavartkiladze, Int. J. Mod. Phys. A **31**, 1650077 (2016).
[16] M. Flanz, E. A. Paschos, U. Sarkar, and J. Weiss, Phys. Lett. B **389**, 693 (1996).
[17] A. Pilaftsis, Phys. Rev. D **56**, 5431 (1997).
[18] A. Pilaftsis and T. E. J. Underwood, Nucl. Phys. **B692**, 303 (2004).
[19] S. Blanchet and P. Di Bari, New J. Phys. **14**, 125012 (2012).
[20] P. S. B. Dev, P. Millington, A. Pilaftsis, and D. Teresi, Nucl. Phys. **B886**, 569 (2014); **B891**, 128 (2015).
[21] P. S. B. Dev, P. Millington, A. Pilaftsis, and D. Teresi, Nucl. Phys. **B897**, 749 (2015).
[22] A. Pilaftsis and D. Teresi, Phys. Rev. D **92**, 085016 (2015).
[23] M. Y. Khlopov and A. D. Linde, Phys. Lett. **138B**, 265 (1984); J. R. Ellis, D. V. Nanopoulos, and S. Sarkar, Nucl. Phys. **B259**, 175 (1985).
[24] S. Davidson and A. Ibarra, Phys. Lett. B **535**, 25 (2002); K. Kohri, T. Moroi, and A. Yotsuyanagi, Phys. Rev. D **73**, 123511 (2006).
[25] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).
[26] P. A. R. Ade *et al.* (Planck Collaboration), Astron. Astrophys. **594**, A13 (2016).
[27] S. P. Martin and M. T. Vaughn, Phys. Rev. D **50**, 2282 (1994).
[28] S. Antusch and M. Ratz, J. High Energy Phys. 07 (2002) 059.
[29] P. H. Chankowski and Z. Pluciennik, Phys. Lett. B **316**, 312 (1993).
[30] W. Grimus, Lett. Nuovo Cimento **27**, 169 (1980).
[31] B. Pendleton and G. G. Ross, Phys. Lett. **98B**, 291 (1981).
[32] H. Arason, D. J. Castano, B. Keszthelyi, S. Mikaelian, E. J. Piard, P. Ramond, and B. D. Wright, Phys. Rev. D **46**, 3945 (1992).
[33] C. R. Das and M. K. Parida, Eur. Phys. J. C **20**, 121 (2001), and references therein.
[34] L. J. Hall, Nucl. Phys. **B178**, 75 (1981).
[35] I. Antoniadis, C. Kounnas, and K. Tamvakis, Phys. Lett. **119B**, 377 (1982); Y. Yamada, Phys. Lett. B **316**, 109 (1993).
[36] S. P. Martin and M. T. Vaughn, Phys. Lett. B **318**, 331 (1993).
[37] K. Melnikov and T. v. Ritbergen, Phys. Lett. B **482**, 99 (2000); P. Marquard, A. V. Smirnov, V. A. Smirnov, and M. Steinhauser, Phys. Rev. Lett. **114**, 142002 (2015); A. L. Kataev and V. S. Molokoedov, Eur. Phys. J. Plus **131**, 271 (2016).
[38] ATLAS, CDF, CMS, D0 Collaborations, arXiv:1403.4427.
[39] M. E. Machacek and M. T. Vaughn, Nucl. Phys. **B236**, 221 (1984).
[40] H. Baer, C. Balazs, M. Brhlik, P. Mercadante, X. Tata, and Y. Wang, Phys. Rev. D **64**, 015002 (2001).
[41] A. Ibarra and C. Simonetto, J. High Energy Phys. 04 (2008) 102.
[42] G. D'Ambrosio, G. F. Giudice, and M. Raidal, Phys. Lett. B **575**, 75 (2003).