Effects of Reheating on Charged Lepton Yukawa Equilibration and Leptogenesis

Arghyajit Datta,[1, 2, 3, *] Rishav Roshan,[4, †] and Arunansu Sil[3, ‡]

[1]*Department of Physics, Kyungpook National University, Daegu 41566, Republic of Korea*
[2]*Center for Precision Neutrino Research, Chonnam National University, Gwangju 61186, Republic of Korea*
[3]*Department of Physics, Indian Institute of Technology Guwahati, Assam 781039, India*
[4]*School of Physics and Astronomy, University of Southampton,*
Southampton, SO17 1BJ, U.K.

We show that the process of non-instantaneous reheating during the post-inflationary period can have a sizable impact on the charged lepton Yukawa equilibration temperature in the early Universe. This suggests relooking the effects of lepton flavors in the leptogenesis scenario where the production and decay of right-handed neutrinos take place within this prolonged era of reheating. We find this observation has the potential to shift the flavor regime(s) of leptogenesis compared to the standard thermal scenario.

The observed dominance of matter over antimatter in our Universe [1] requires an explanation where such an asymmetry is expected to be generated in a post-inflationary phase as otherwise, any pre-inflationary asymmetry, if present, would be erased due to the exponential expansion of the Universe during inflation. Among many possibilities, production of lepton asymmetry resulting from out-of-equilibrium decay of heavy right-handed neutrinos (RHN) N into Standard Model (SM) lepton ($\ell_L^T = (\nu_L, e_L)$) and Higgs (H) doublets, called *leptogenesis* [2–4], remains an attractive one. The obvious reason behind it lies in its close proximity with neutrino mass (m_ν) generation mechanism via type-I seesaw [5–12] where the Lagrangian is given by,

$$-\mathcal{L} = \overline{\ell}_{L_\alpha}(Y_\nu)_{\alpha i}\tilde{H}N_i + \frac{1}{2}\overline{N_i^c}(M_R)_i N_i + h.c., \quad (1)$$

(in the charged lepton diagonal basis) with $\alpha = e, \mu, \tau$ and $i = 1, 2, 3$. Such a lepton asymmetry then reprocesses itself in terms of baryon asymmetry by $B + L$ violating sphaleron transitions [13–15] (while conserving $B - L$) which remains in equilibrium from a temperature $T_s \sim 10^{12}$ GeV down to electroweak (EW) scale $T_{EW} \sim 100$ GeV [14, 16–18] in a radiation dominated Universe.

In the standard thermal leptogenesis scenario, the lepton or the $B - L$ asymmetry is generated during the radiation-dominated era which is indicative of the underlying assumption that the reheating temperature ($T_{\rm RH}$) after inflation is greater than the mass of the decaying (lightest) RHN, say M_1. At this end, to study the evolution of this $B - L$ asymmetry, effect of different SM interactions in terms of their respective equilibration temperatures (T^*) becomes very much relevant. This is due to the fact that a particle species (say X with degrees of freedom g_X) involved in thermal equilibrium at a given temperature T experiences chemical equilibrium conditions leading to an inter-connection between chemical potentials (μ_X) of different species while other constraints followed from conservation laws are also enforced as outlined in the Supplemental Material. In terms of particle (antiparticle) number density n_X ($n_{\bar{X}}$) to entropy (s)

ratio defined by[19] $Y_{\Delta_X} = n_{\Delta_X}/s = (n_X - n_{\bar{X}})/s \equiv g_X \frac{T^3}{6s}(\mu_X/T)$, the lepton asymmetry $Y_L \equiv Y_{\Delta_L}$ generated from N decay becomes [20]

$$Y_L = \frac{n_{\Delta L}}{s} = \sum_\alpha (2Y_{\Delta_{\ell_\alpha}} + Y_{\Delta_{e_{R_\alpha}}}), \quad (2)$$

where $Y_{\Delta_{e_{R_\alpha}}}$ is the asymmetry of the right handed charged leptons, if any. This Y_L in turn directly determine the $B - L$ asymmetry, $Y_{B-L} = -Y_L$ when $B = 0$.

Note that as long as the charged lepton Yukawa (Y_α) interactions do not enter equilibrium, Y_L is confined within the asymmetry in the lepton doublets only as $Y_{\Delta_{e_{R_\alpha}}} = 0$. The temperature at which a specific e_{R_α} equilibrates with corresponding lepton doublet, known as the equilibration temperature (ET), can be determined naively by comparing the Higgs boson decay (and inverse decay) rate $\langle \Gamma_\alpha \rangle$ with that of the Hubble expansion $\mathcal{H}$ [21, 22]. For example, the τ_R equilibration temperature in a radiation dominated (RD) Universe is obtained by solving $\langle \Gamma_\tau \rangle = \mathcal{H}(T)$ or equivalently

$$\frac{\pi Y_\tau^2}{192\zeta(3)}\frac{m_h^2(T)}{T} = 0.066 g_*^{1/2}\frac{T^2}{M_P}, \quad (3)$$

where $\mathcal{H}(T) = 0.066 g_*^{1/2}\frac{T^2}{M_P}$ is used (M_P is the reduced Planck mass), g_* represents effective number of relativistic degrees of freedom, and $m_h(T) \simeq 0.6T$ [23–25] is the thermal mass of the Higgs. The τ_R-ET $T^*_{0(\tau)}$ is found to be 5×10^{11} GeV as seen from the intersection point of solid blue line (standard) and $\langle \Gamma_\tau \rangle/\mathcal{H} = 1$ dotted horizontal line of Fig. 1. For μ_R and e_R, such temperatures $T^*_{0(\mu)}$ and $T^*_{0(e)}$ are found to be 10^9 GeV and 5×10^4 GeV respectively. Though a more refined analysis in the context of thermal field theory [26, 27] may slightly alter these numbers, such naive estimates of ETs remain sufficient for our purpose.

This finding plays a pivotal role behind the loss of lepton flavor coherence (fully or partially) as observed in flavor leptogenesis [20, 28–32]. Depending on which charged lepton Yukawa interaction(s) is(are) in equilibrium at a

given temperature T, the lepton doublet state(s) ℓ_{L_α} produced from the RHN-decay finds corresponding e_{R_α} in the bath to interact. As a result, the quantum coherence of the lepton doublet states gets broken, which in turn, dictates the number of orthogonal flavor alignments in that particular temperature regime.

A basic assumption behind such standard evaluation of ET is that the entire epoch happens below the reheating temperature, or in other words, a RD phase is prevalent. Though this sounds reasonable, at the same time, we can also recollect that reheating might not be an instantaneous process [33]. In its simplest version, the process of reheating begins in a matter-dominated phase where the oscillatory inflaton field (ϕ) decays perturbatively. The thermalization of these light decay products helps the Universe to attain quickly a maximum temperature $T_{\rm Max}$, subsequent to which the temperature falls to $T_{\rm RH}$ marking the commencement of RD era thereafter. During this period (*i.e.* $T_{\rm RH} < T < T_{\rm Max}$), temperature varies less sharply as compared to standard scaling $T \propto a^{-1}$ (a is the Friedmann–Robertson–Walker scale factor) and the Universe also expands faster than in the RD era at a given temperature.

In this letter, we show that such non-instantaneous reheating epoch exhibits a shift in ET as $\mathcal{H}$ would differ from that of a purely RD era, a situation which has not been discussed elsewhere to the best of our knowledge. In view of recent finding [34] that $T_{\rm Max}$ can be much larger than $T_{\rm RH}$, this study can be even more relevant. We further elaborate here for the first time[1] the impact of such shift in ET (due to the unusual relation between T and a) on the study of flavor leptogenesis too where the production and decay of RHN take place within $T_{\rm Max}$ and $T_{\rm RH}$. We begin the discussion on reheating below to estimate the modified temperature regimes where the right-handed charged leptons enter equilibrium, first without introducing any RHNs, and later we incorporate RHNs so as to discuss their influence on flavored leptogenesis.

As the era of reheating starts following the inflationary epoch, the discussion of reheating is somewhat connected to the form of inflaton potential near its minimum. To be specific, following [34] we consider a generic power-law form of the inflaton potential near the minimum, $V(\phi) = \lambda \frac{|\phi|^n}{M_P^{n-4}}$ the origin of which can be traced back to T-attractor models in no-scale supergravity [39]. For simplicity, we confine ourselves to a $n = 2$ case resembling the Starobinsky type of inflationary potential [40–42] in the large field limit. Here we skip details of inflation (see [43]) except the information that the magnitude of the parameter λ is found to

be $\sim 2 \times 10^{-11}$ so as to be consistent with inflationary observables like spectral index and tensor-to-scalar ratio within 95% C.L. of the Planck+BICEP2/Keck (PBK) constraints [44]. This leads to the effective mass of the inflaton as $m_\phi = (\partial_\phi^2 V(\phi))^{1/2} \simeq 1.5 \times 10^{13}$ GeV.
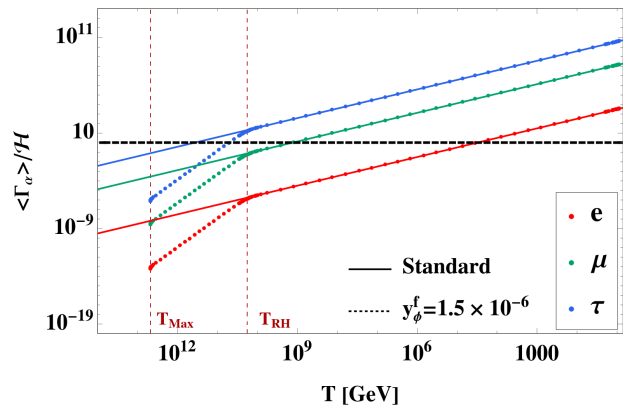

FIG. 1: Variation of $\langle \Gamma_\alpha \rangle / \mathcal{H}$ *w.r.t.* T for standard (solid lines) and modified (dotted) scenarios.

We introduce an effective interaction between the inflaton and any specific pair of SM fermion (f) anti-fermion fields as $y_\phi^f \phi \bar{f} f$ which will dictate the decay of the inflaton to radiation (composed of massless SM fermion fields in the early Universe). The origin of such effective interaction can be traced back to dimension-5 operator involving SM Yukawa interactions such as $\frac{\alpha_\ell}{\Lambda} \phi \bar{\ell}_L H E_R$, $\frac{\alpha_u}{\Lambda} \phi \bar{q}_L \tilde{H} u_R$, $\frac{\alpha_d}{\Lambda} \phi \bar{q}_L H d_R$, where Λ and α_f correspond to a cut-off scale and coupling constant(s) respectively. This would initiate a decay of inflaton to Higgs and a pair of fermion anti-fermion fields where a subsequent fast decay of the Higgs into another fermion anti-fermion pair (via SM Yukawa interactions) analogously represents the decay of the inflaton into fermion anti-fermion pair. Considering all such possible final state fermion anti-fermion pairs (denoted by g_f number of possibilities) couple uniformly to ϕ, the total decay width of ϕ can be represented by $\Gamma_\phi \simeq g_f^2 \frac{|y_\phi^f|^2}{8\pi} m_\phi$. Furthermore, y_ϕ^f is kept sufficiently small (quantified later) so as to exclude the possibility of parametric excitation of fermions [45] and ensure the perturbative decay of ϕ. The role of reheating is therefore limited here in changing the thermal history of the Universe to clearly exhibit the impact on charged lepton Yukawa ET.

The energy densities of the inflaton field (ρ_ϕ) and radiation (ρ_R) satisfy the evolution equations[2] in terms of

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[1] Earlier works on flavor leptogenesis during reheating [31, 35, 36] or within different modified cosmological scenarios[37, 38] did not incorporate such shift(s) in ET(s).

[2] Such scale factor dependance is explained in the Supplemental Material.

a:

$$\frac{d(\rho_\phi a^3)}{da} = -\frac{\Gamma_\phi}{\mathcal{H}}\rho_\phi a^2; \text{ and } \frac{d\left(\rho_R a^4\right)}{da} = \frac{a^3}{\mathcal{H}}\Gamma_\phi \rho_\phi, \quad (4)$$

where $\mathcal{H}^2 = (\dot{a}/a)^2 = (\rho_\phi + \rho_R)/3M_P^2$. At this moment, it is pertinent to define the temperature $T(a)$ connected to the ρ_R by the relation

$$T(a) = \left[\frac{30}{g_*(a)\pi^2}\right]^{1/4}\rho_R^{1/4}(a). \quad (5)$$

T can therefore be estimated as a solution to the coupled Eqs. (4) whose variation with the relative scale factor (a_{end} is the scale factor at the end of inflation, considered as a mere reference point) is presented in the top panel of Fig. 2 with $y_\phi^f = 1.5 \times 10^{-6}$ (with $g_f = 30$ kept throughout). Note that during the onset of reheating, $\mathcal{H}$ has the sole contribution from ρ_ϕ. However, as soon as the light decay products of ϕ thermalizes, ρ_R gets growing sharply and correspondingly a maximum temperature T_{Max} is attained. After this, ρ_R decreases slowly (as ϕ being decayed, ρ_ϕ diminishes) and finally a stage is reached where ρ_R eventually becomes equal to ρ_ϕ and starts dominating thereafter. This boundary is marked as the reheating temperature T_{RH}.
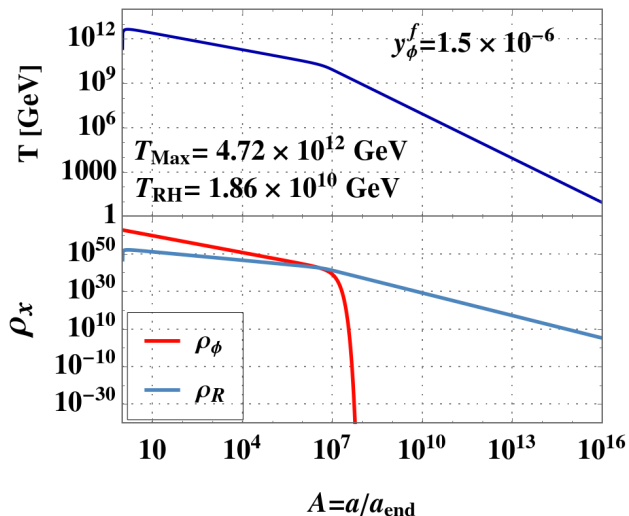

FIG. 2: Top (bottom) panel indicates evolutions of T (various energy densities) as function of the relative scale factor.

We now infer that provided the reheating temperature T_{RH} is smaller than the standard charged lepton Yukawa ET, $T_{0(\alpha)}^*$, for a specific flavor α, the epoch in which this equilibration (by equating $\langle\Gamma_\alpha\rangle$ with $\mathcal{H}$) happens may actually experience the era between T_{Max} and T_{RH}. The relatively faster expansion of the Universe in this era will obviously affect the $r.h.s$ of Eq. (3) so as to replace it with $\mathcal{H}$ which now depends on both ρ_ϕ and ρ_R. Using the ρ_ϕ

and ρ_R information obtained from Fig. 2 as input for $\mathcal{H}$, the modified $\langle\Gamma_\alpha\rangle/\mathcal{H}$ is plotted against T in Fig. 1 indicated by the dotted lines. As expected, prior to entering into the RD phase, it depicts a non-standard behavior (slope is different within T_{Max} and T_{RH}) as shown in Fig. 1, separately for three flavors of charged leptons. This is a new observation in our work, hitherto unexplored in the literature.

We note that Y_τ comes to equilibrium around $T_\tau^* \sim 5 \times 10^{10}$ GeV (compared to $T_{0(\tau)}^* \simeq 5 \times 10^{11}$ GeV in standard scenario) while μ_R and e_R attain equilibriums at their standard temperatures $T_{0(\mu)}^*$ and $T_{0(e)}^*$ respectively in this case. This is because with the specific choice of $y_\phi^f = 1.5 \times 10^{-6}$, T_{RH} is found to be 1.86×10^{10} GeV which falls below $T_{0(\tau)}^*$. Hence the non-standard effect on Y_μ and Y_e interactions do not appear in this particular example. For further smaller y_ϕ^f, T_{RH} may fall below $T_{0(\mu)}^*$ and correspondingly we would get a modified Y_μ ET. On the other hand, for sufficiently large y_ϕ^f, the temperature regime between T_{Max} and T_{RH} is getting shortened and hence such departure in evaluating the charged lepton ET, even if it is present, becomes less prominent. In the Supplemental Material, we elaborate on it exploring a range of y_ϕ^f.

Equipped with this basic understanding behind such a non-standard situation in case of prolonged reheating, we now turn our attention to the realm of flavor leptogenesis. The lepton asymmetry being generated from the out-of-equilibrium decay of the lightest RHN N_1, its mass M_1 is at the center of interest. This is because, in standard thermal leptogenesis, N_1 starts to decay at a temperature $T \sim M_1$ when the Universe is RD and N_1 is present in equilibrium within the thermal bath. Contrary to this, in case of non-instantaneous reheating under study, in general there can be three possibilities: (a) $M_1 < T_{\mathrm{RH}}$, (b) $T_{\mathrm{RH}} < M_1 < T_{\mathrm{Max}}$ and (c) $M_1 > T_{\mathrm{Max}}$. While case (a) is a scenario close to standard thermal leptogenesis (and we refrain from discussing it here), case (c) stands for a situation where N_1 can not be thermally produced, discussed separately in another work [46]. Here we focus on the case (b), where thermal production of N_1 as well as its decay seem to be a possibility within this reheating epoch.

To start, we need to introduce the energy density associated to N_1 (ρ_{N_1}) in the discussion as T_{Max} being greater than M_1, N_1 can be produced from the inverse decay[3]. The modified Hubble parameter is given by

$$\mathcal{H}^2 = \frac{\rho_\phi + \rho_R + \rho_{N_1}}{3M_P^2}. \quad (6)$$

Alongside, Eq. (4) for ρ_R would now have an additional term $\frac{a^3}{H}\langle\Gamma_{N_1}\rangle(\rho_{N_1} - \rho_{N_1}^{\mathrm{eq}})$ on the $r.h.s.$, the presence of

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[3] For simplicity, we do not consider production via scattering.

which is to represent the dilution of ρ_R due to production of N_1 from it (via inverse decays). The decay rate of N_1 is given by $\langle \Gamma_{N_1} \rangle$ [47] and ρ_{N_1} satisfies the other equation

$$\frac{d(\rho_{N_1} a^3)}{da} = -\frac{\langle \Gamma_{N_1} \rangle a^2}{\mathcal{H}}(\rho_{N_1} - \rho_{N_1}^{\text{eq}}). \tag{7}$$

Regarding initial conditions in order to solve the coupled equations for ρ_ϕ, ρ_R and ρ_{N_1}, we employ $\rho_{N_1} = \rho_R = 0$ whereas energy density at the end of inflation is taken to be $\rho_{\phi_{end}} = 3V(\phi_{end})/2$. The value of ϕ_{end} for the class of models under consideration is $0.78M_P$ [34]. All the three components of the energy densities are now plotted against the relative scale factor in Fig. 3. For this plot, we keep the parameter $y_\phi^f = 1.5 \times 10^{-6}$ same as in Fig. 2 and consider a specific $M_1 = 10^{11}$ GeV, lies in between T_{Max} and T_{RH} á la case (b). We consider $M_{2(3)}$ as 100 times heavier than $M_{1(2)}$ for discussion and hence thermal generation of $N_{2(3)}$ is not possible.

The neutrino Yukawa Y_ν appearing in $\langle \Gamma_{N_1} \rangle$ is evaluated with the help of Casas-Ibarra formalism [48],

$$Y_\nu = -i\frac{\sqrt{2}}{v}U D_{\sqrt{m}} \mathbf{R} D_{\sqrt{M}}, \tag{8}$$

where U is the PMNS [49–51] mixing matrix diagonalizing $m_\nu = -Y_\nu M_R^{-1} Y_\nu^T v^2/2$ ($v = 246$ GeV is the EW vev), $D_m(D_M)$ is the diagonal active neutrino (RHN) mass matrix and $\mathbf{R}$ is a complex orthogonal matrix of the form as in [52] parametrized by complex mixing angle θ_R. The values of $\text{Re}[\theta_R], \text{Im}[\theta_R]$ are so chosen as to produce correct baryon asymmetry via leptogenesis and also to keep Y_ν entries perturbative. We incorporate the best fit values of mixing angles and mass-squared differences [50] to define U and light neutrino mass eigenvalues (with $m_1 = 0$).

As can be seen from Fig. 3, N_1 gradually increases, however, always remain subdominant compared to ρ_R. This is because it also decays out of equilibrium within the same period. As a result, the reheating temperature T_{RH} is essentially unchanged even in presence of this new component, ρ_{N_1}. From Fig. 1, we recall that Y_τ enters equilibrium below $T_\tau^* = 5 \times 10^{10}$ GeV in this modified scenario. Therefore, when N_1 starts to decay, i.e. around $T \lesssim M_1(= 10^{11}$ GeV), none of the charged

lepton Yukawa couplings reaches equilibrium. Hence, the flavor coherence is maintained and we encounter an unflavored regime of leptogenesis. Note that for the same M_1 value, in standard thermal leptogenesis happening in a RD era with $T_{0(\tau)}^* = 5 \times 10^{11}$ GeV being larger than M_1, Y_τ-interaction would already be in equilibrium breaking flavor coherence. In that case, contrary to our realization, lepton asymmetries would be produced in two orthogonal directions: τ and κ, κ forming the subspace of a coherent superposition of e and μ lepton flavors [20, 28–31] leading to flavor leptogenesis. Such a shift in flavor regime (from two flavor to unflavored) of leptogenesis is an important outcome of our study.
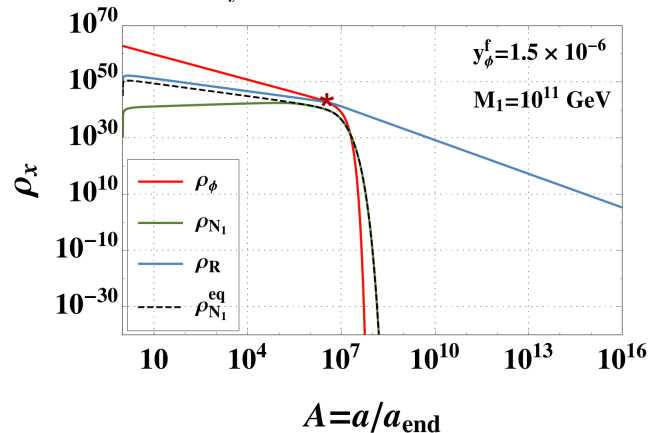

FIG. 3: Evolution of various components of total energy density are displayed. The $*$ indicates the T_{RH}.

Realizing that the case with $M_1 = 10^{11}$ GeV (and $y_\phi^f = 1.5 \times 10^{-6}$) corresponds to unflavored leptogenesis scenario during this extended reheating period, we proceed to evaluate the $B-L$ asymmetry using the following Boltzmann equation [47],

$$\frac{d(n_\Delta a^3)}{da} = -\frac{\langle \Gamma_{N_1} \rangle a^2}{\mathcal{H}} \left[\frac{\varepsilon_\ell}{M_1}(\rho_{N_1} - \rho_{N_1}^{\text{eq}}) + \frac{n_{N_1}^{\text{eq}}}{2n_\ell^{\text{eq}}} n_\Delta \right], \tag{9}$$

with $n_\Delta = n_{B-L}$. However, in a more general case where a shift of regimes of thermal leptogenesis still leads to a flavored one, the corresponding equation would be

$$\frac{d(n_{\Delta_i} a^3)}{da} = -\frac{\langle \Gamma_{N_1} \rangle a^2}{\mathcal{H}} \left[\frac{\varepsilon_{\ell_i}}{M_1}(\rho_{N_1} - \rho_{N_1}^{\text{eq}}) + \frac{1}{2}K_i^0 \sum_j (C_{ij}^\ell + C_j^H)\frac{n_{N_1}^{\text{eq}}}{n_\ell^{\text{eq}}} n_{\Delta_j} \right]. \tag{10}$$

Here $K_i^0 = (Y_\nu^*)_{\alpha 1}(Y_\nu)_{\alpha 1}/(Y_\nu^\dagger Y_\nu)_{11}$ is flavor projector [30] and C^ℓ, C^H matrices connect the asymmetries in lep-

ton and Higgs to asymmetries in $\Delta_i = B/3 - L_i$ (in terms of $n_{\Delta_{\ell_\tau}}$ and $n_{\Delta_{\ell_\kappa}}$ here)[30]. The CP asymmetry $\varepsilon_{\ell_\alpha}$ (for

unflavored case, $\varepsilon_\ell = \sum_\alpha \varepsilon_{\ell_\alpha}$) involved is obtained from the decay of N_1 to a specific flavor ℓ_α and estimated using standard expression [30, 53]. The final baryon asymmetry Y_B is related to n_{Δ_i} by [54]: $Y_B = \frac{28}{79} \sum_\alpha n_{\Delta_\alpha}/s$.
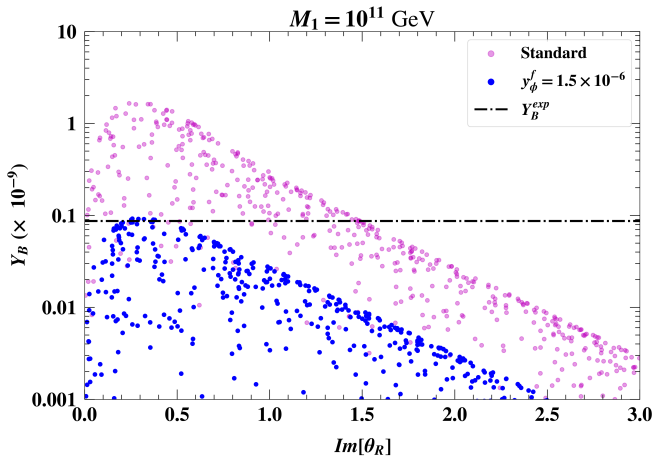

FIG. 4: Y_B in case (a) non-instantaneous reheating (blue) and (b) standard RD phase (magenta): variation with $\mathrm{Im}[\theta_R]$ is displayed.

In Fig. 4, we compare the parameter space of Y_B against $\mathrm{Im}[\theta_R]$, for the choices of M_1 and y_ϕ^f discussed above, obtained in usual thermal leptogenesis (fall in two flavored regimes, shown in magenta) with that in case of prolonged reheating (unflavored regime, shown in blue). Here we scan over the range of θ_R as: $\mathrm{Re}[\theta_R]$ (0 - 2π) and $\mathrm{Im}[\theta_R]$ (0-3). We observe a clear shift in the asymmetry by almost an order of magnitude, signifying the effect of non-instantaneous reheating. Furthermore, in terms of the correct amount of baryon asymmetry, the situation becomes more restrictive as evident from the presence of a few number of blue points over magenta ones across the correct Y_B line (horizontal black dashed). At this point, it is pertinent to inquire what happens if we choose a different set of M_1, y_ϕ^f, within case (b). To answer it, we include a scan for $y_\phi^f - M_1$ in the Supplemental Material which helps make a clear distinction with the standard scenario in terms of finding the effect of the shift in the ET on thermal leptogenesis. Though we restrict ourselves in effective interaction between inflaton and SM fermions only, inflaton-boson couplings such as $\sigma\phi^2 H^\dagger H$ and $\mu\phi H^\dagger H$ may also be present. However, in order to keep decays of ϕ via such interactions in perturbative regime (to avoid rapid production of particles via preheating [55]), the associated couplings should satisfy: $\sigma < 5.5\times10^{-12}, \mu < 5.2\times10^{-12} M_P$ [56]. It turns out that with such σ and μ values, our analysis with $y_\phi^f \phi \bar{f} f$ alone would not be impacted by the presence of inflaton-boson coupling(s).

We have shown that the effect of a prolonged reheating period, starting from the maximum temperature of the Universe following inflation to the reheating, affects the equilibration of the charged lepton Yukawa interactions significantly. Due to the faster expansion of the Universe in this reheating period, the charged lepton Yukawa interactions enter equilibrium in a delayed fashion. This new observation depends on the effective coupling of the inflaton field with SM fields. With low reheating temperature (i.e. with a smaller effective coupling of the inflaton field with SM fields), the era of reheating becomes longer and hence the ET(s) can be significantly smaller. Effects of this observation should also be applicable to a wider SM interactions and open a possibility for further investigation in different directions. As one such application, we study here its effect on leptogenesis where the lightest RHN belonging to the type-I seesaw is being created and decayed out of equilibrium within T_{Max} and T_{RH}. As expected, we find that the delayed entry of the charged lepton Yukawas in equilibrium can shift the flavor regimes of leptogenesis. This is a new observation and warrants more detailed studies. Alongside, there could be several directions that may open up. For example, in a follow-up study of this work, a new regime of flavor leptogenesis is found [46] in the context of non-thermal leptogenesis during extended reheating. Finally, we expect that any low scale leptogenesis scenario with low reheating temperature would be greatly impacted by our finding and can also be tested in current and near future experiments.

* arghyad053@gmail.com
† rishav.roshan@gmail.com
‡ asil@iitg.ac.in

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Effects of Reheating on Charged Lepton Yukawa Equilibration and Leptogenesis

Supplemental Material

Arghyajit Datta, Rishav Roshan and Arunansu Sil

This Supplemental Material is organized as follows. In Sec. I, we discuss the equilibrium conditions for the SM interactions responsible for the production of lepton asymmetry. In Sec. II, we analyze the dependence of T_{Max}, T_{RH} as well as the equilibration temperatures of the charged lepton Yukawa interactions on the effective individual inflaton-SM fermions coupling y_ϕ^f. In Sec. III, we provide the numerical analysis required to generate the baryon asymmetry in this modified setup and finally present the allowed parameter space in $y_\phi^f - M_1$ plane which is consistent with the observed baryon asymmetry.

I. THE EQUILIBRIUM CONDITIONS

The Boltzmann equations (BE) describing the evolution of lepton asymmetry (Y_L) produced in leptogenesis are in general connected with different particle species and their asymmetries in a sense that various SM interactions, if found to be in chemical equilibrium, develop relations among number density asymmetries between different species. Such number density asymmetry $(n_X - n_{\bar{X}})$ for a particle species X having number of degrees of freedom g_X is connected to the respective chemical potential (μ_X) via the relation [19]

$$n_X - n_{\bar{X}} = \frac{g_X T^2}{6} \begin{cases} \mu_X & \text{for fermions,} \\ 2\mu_X & \text{for bosons,} \end{cases} \tag{1}$$

as obtained from the respective distribution function. A particular SM interaction, whether or not in thermal equilibrium ($\equiv$ in chemical equilibrium) is determined by comparing the respective interaction rate (Γ_{Int}) with Hubble expansion rate of the Universe ($\mathcal{H}$) at any temperature T. If $\Gamma_{\mathrm{Int}} \gtrsim \mathcal{H}$, the corresponding interaction rate is fast enough to establish an equilibrium situation such that the particles involved satisfy an algebraic relation among their chemical potentials. For example, if the charged lepton Yukawa interaction $(Y_\alpha \bar{\ell}_{L_\alpha} H e_{R_\alpha})$ of a particular flavor α satisfies such condition, it would maintain the relation

$$\mu_{e_{R_\alpha}} - \mu_{\ell_{L_\alpha}} + \mu_H = 0. \tag{2}$$

Such inter-relations among the chemical potentials are important to quantify the lepton asymmetry we are interested in, as explained below.

The lepton asymmetry in terms of number densities of leptons and anti-leptons, normalized to the entropy density (s), is defined by

$$Y_L(\equiv Y_{\Delta_L}) = \frac{n_{\Delta_L}}{s} = \sum_\alpha \frac{1}{s} \Big[\big(n_{e_{L_\alpha}} - n_{\bar{e}_{L_\alpha}} \big) + \big(n_{\nu_{L_\alpha}} - n_{\bar{\nu}_{L_\alpha}} \big) + \big(n_{e_{R_\alpha}} - n_{\bar{e}_{R_\alpha}} \big) \Big], \tag{3}$$

where both the $SU(2)_L$ doublet and singlet leptons are included for general purpose. Using Eq. (1) and the fact that the two members of a $SU(2)_L$ multiplet will have same chemical potential, *i.e.* $\mu_\ell \equiv \mu_{e_L} = \mu_{\nu_L}$ resulting from the vanishing chemical potential of the electroweak gauge bosons above the electroweak scale, when the corresponding gauge interaction is in equilibrium, the total lepton number asymmetry can be expressed as:

$$Y_L = \frac{T^2}{6s} \sum_\alpha \big(2\mu_{\ell_\alpha} + \mu_{e_{R_\alpha}} \big), \tag{4}$$

$$\equiv \sum_\alpha \Big(2Y_{\Delta_{\ell_\alpha}} + Y_{\Delta_{e_{R_\alpha}}} \Big). \tag{5}$$

Now to determine whether few or all of these chemical potentials $\mu_{\ell_\alpha}, \mu_{e_{R_\alpha}}$, are non-zero and the relationships these chemical potentials might carry with other non-zero chemical potentials, one needs to identify the temperature regime $(T \sim M_1)$ where the lepton asymmetry is being generated from the decay of the lightest right handed neutrino (RHN) of mass M_1. Then, following conditions on the chemical potentials for different SM fields arise as a result of corresponding SM interactions, attaining thermal equilibrium with the decrease in the temperature of the Universe (in a radiation dominated era) should be considered.

- For temperature above 10^{13} GeV, interactions mediated by the SM gauge couplings and the top quark Yukawa coupling comes into equilibrium. Consequently, the following conditions on the chemical potentials of the participating particles arise: (i) The chemical potential of the gauge bosons (including gluons, W^i_μ and B_μ) are zero, $i.e.$,

$$\mu_g = \mu_{W^i} = \mu_B = 0, \tag{6}$$

  with $i = 1, 2, 3$. As a result, two members of a $SU(2)_L$ doublet reach chemical equilibrium and have same chemical potential: $\mu_Q \equiv \mu_{u_L} = \mu_{d_L}, \quad \mu_\ell \equiv \mu_{e_L} = \mu_{\nu_L}, \quad \mu_H \equiv \mu_{H^+} = \mu_{H^0}$. This is already incorporated in Eq. (4). (ii) Furthermore, equilibration of top quark Yukawa interaction yields,

$$\mu_t = \mu_{Q_3} + \mu_H. \tag{7}$$

- A temperature independent constraint follows from the hypercharge neutrality condition, leading to

$$\sum_\alpha (\mu_{Q_\alpha} + 2\mu_{u_\alpha} - \mu_{d_\alpha} - \mu_{\ell_\alpha} - \mu_{e_{R_\alpha}}) + 2\mu_H = 0, \tag{8}$$

  where u_α, d_α represents the $SU(2)_L$ singlet quarks of α flavor respectively.

- Around $T \sim 10^{13}$ GeV, (i) strong sphalerons come to equilibrium [14, 57] leading to:

$$\sum_\alpha (2\mu_{Q_\alpha} - \mu_{u_\alpha} - \mu_{d_\alpha}) = 0. \tag{9}$$

  (ii) On the other hand, equilibrium temperature of electroweak sphalerons is found to be smaller than that of strong sphaleron equilibrium temperature [14, 16–18] by an order of magnitude. Once in equilibrium, this enforces:

$$\sum_\alpha (3\mu_{Q_\alpha} + \mu_{\ell_\alpha}) = 0. \tag{10}$$

  (iii) At this stage, all the charged lepton Yukawa interactions are out of equilibrium and hence, the respective chemical potentials of the $SU(2)_L$ singlet right handed charged leptons can be set to zero ($i.e. \sum_\alpha Y_{\Delta_{e_{R_\alpha}}} = 0$). Then the lepton asymmetry, if produced from N_1 decay in this regime, along lepton doublet direction only contribute to the total lepton asymmetry and becomes,

$$Y_L = \sum_\alpha Y_{\Delta_{\ell_\alpha}}. \tag{11}$$

- However, when the temperature further falls down to $T < 10^{13}$ GeV, charged lepton (and remaining quark Yukawa interactions) Yukawa interactions start entering equilibrium one after other. As the temperature drops, interactions with larger Yukawa coupling attains thermal equilibrium faster. Once they are in equilibrium, the following relations among chemical potentials:

$$\mu_{e_{R_\alpha}} - \mu_{\ell_\alpha} + \mu_H = 0 \quad \text{for charged leptons,} \tag{12}$$
$$\mu_{u_\alpha} - \mu_{Q_\alpha} - \mu_H = 0 \quad \text{for up-type quarks,} \tag{13}$$
$$\mu_{d_\alpha} - \mu_{Q_\alpha} + \mu_H = 0 \quad \text{for down-type quarks} \tag{14}$$

  can be set. It is found that while Y_τ interaction enters equilibrium at $T = 5 \times 10^{11}$ GeV, that of Y_μ and Y_e interactions are at 10^9 GeV and 5×10^4 GeV respectively. In this case, depending upon which charged lepton Yukawa interactions are in equilibrium, following Eq. (5), contributions from $Y_{\Delta_{e_{R_\alpha}}}$ should also be incorporated in evaluating Y_L.

Finally, to quantify the lepton asymmetry contribution coming from right handed leptons or from the lepton doublets, one needs take into account the chemical potential relations (and other constraints) from Eq. (8)-(14), which eventually leads to the C^ℓ, C^H matrices mentioned in Eq[10] of the main text (see[20, 28, 30] for details). Note that the above discussion, particularly the temperature regimes where different SM interactions enter in equilibrium, is based on the assumption that the Universe undergoes a radiation dominated era during such decouplings. However, as evaluated in the text, the regime $T_{\text{Max}} > T(\sim M_1) > T_{\text{RH}}$ turns out to be an important one where the charged lepton Yukawa interactions are greatly affected and the respective equilibration conditions are to be applied with care.

II. ESTIMATING $T_{\text{Max}}, T_{\text{RH}}$ AND EQUILIBRATION TEMPERATURES AS FUNCTION OF y_ϕ^f

In this section, we evaluate the quantities important for studying reheating. In our work, the process of reheating proceeds perturbatively through the effective inflaton-SM fermion interaction $y_\phi^f \phi \bar{f} f$. As discussed in the main paper, this effective individual coupling y_ϕ^f plays a crucial role in understanding the entire reheating process by providing estimates for T_{Max} and T_{RH}. The respective equilibration temperatures (ET) of different charged lepton Yukawa interactions depend crucially on the temperature profile, dictated by y_ϕ^f, during this extended reheating period which in turn affects the flavor effect on leptogenesis. This makes y_ϕ^f an important parameter in the present setup. Though in the text, we consider a specific choice of y_ϕ^f for the demonstration purpose, it is indeed intriguing to have analytic expressions of the quantities like $T_{\text{Max}}, T_{\text{RH}}$ and ET as function of y_ϕ^f, for general purpose. This will also remain helpful for scanning the parameter space of the work as investigated in the next section.

A. Estimating T_{Max} and T_{RH}

To understand the reheating period, we first study the evolutions of inflaton and radiation density via Eq. 4 of the text as below,

$$\frac{d\rho_\phi}{dt} + 3\left(\frac{3n}{n+2}\right)\mathcal{H}\rho_\phi = -\Gamma_\phi \rho_\phi, \tag{15}$$

$$\frac{d\rho_R}{dt} + 4\mathcal{H}\rho_R = \Gamma_\phi \rho_\phi, \tag{16}$$

where the y_ϕ^f involvement comes through the decay width $\Gamma_\phi = g_f^2 \frac{y_\phi^{f\,2}}{8\pi} m_\phi$ with g_f representing the effective degrees of freedom (considering different generations of leptons and quarks of different colors) and n corresponds to the power involved in the inflationary potential. Here m_ϕ is the effective mass of the inflaton and in the adiabatic approximation, it can be estimated as:

$$m_\phi = \partial_\phi^2 V(\phi) = \lambda^{2/n} n(n-1) \rho_\phi^{\frac{n-2}{n}} M_P^{\frac{2(4-n)}{n}}. \tag{17}$$

In the early Universe just after inflation, the decay rate of the inflaton can be safely assumed to be smaller than the expansion rate of the Universe. This helps to solve Eq. (15) analytically (by neglecting the right-hand side of the equation) and is given by

$$\rho_\phi(a) \simeq \rho_{\text{end}} \left(\frac{a}{a_{\text{end}}}\right)^{-6n/(n+2)}, \tag{18}$$

where ρ_{end} represents the energy density of the inflaton at the end of the inflation and can be estimated by setting the slow roll parameter $\epsilon = 1$ which leads to $\rho_{\text{end}} = \frac{3}{2} V(\phi_{\text{end}})$[43] (for details, see T-attractor inflation models [34, 39]). The radiation energy density ρ_ϕ can also be estimated by solving the Eq. (16) after plugging Eqs. (17) and (18) into it. This gives,

$$\rho_R \simeq \frac{g_f^2 y_\phi^{f\,2}}{8\pi} \rho_{\text{end}}^{\frac{n-1}{n}} \sqrt{n(n-1)} \left(\frac{n+2}{14-2n}\right) \lambda^{1/n} M_P^{4/n} \left(\frac{a_{\text{end}}}{a}\right)^4 \left\{ \left(\frac{a}{a_{\text{end}}}\right)^{\frac{14-2n}{n+2}} - 1 \right\}. \tag{19}$$

Note that, irrespective of the choice of n, the radiation energy density depends on the effective coupling as $\rho_R \propto y_\phi^{f\,2}$. As a result, increasing y_ϕ^f increases the energy density of the radiation. Now, assuming instantaneous equilibration of the decay products of inflaton, the temperature of the Universe can be easily obtained from ρ_R using

$$T = \left(\frac{30\rho_R}{\pi^2 g_*}\right)^{\frac{1}{4}}, \tag{20}$$

where g_* is the effective relativistic degrees of freedom. Hence, just after inflation, the temperature of the Universe increases to a maximum value T_{Max} that can be obtained by maximizing T of Eq. (20) against the scale factor a with
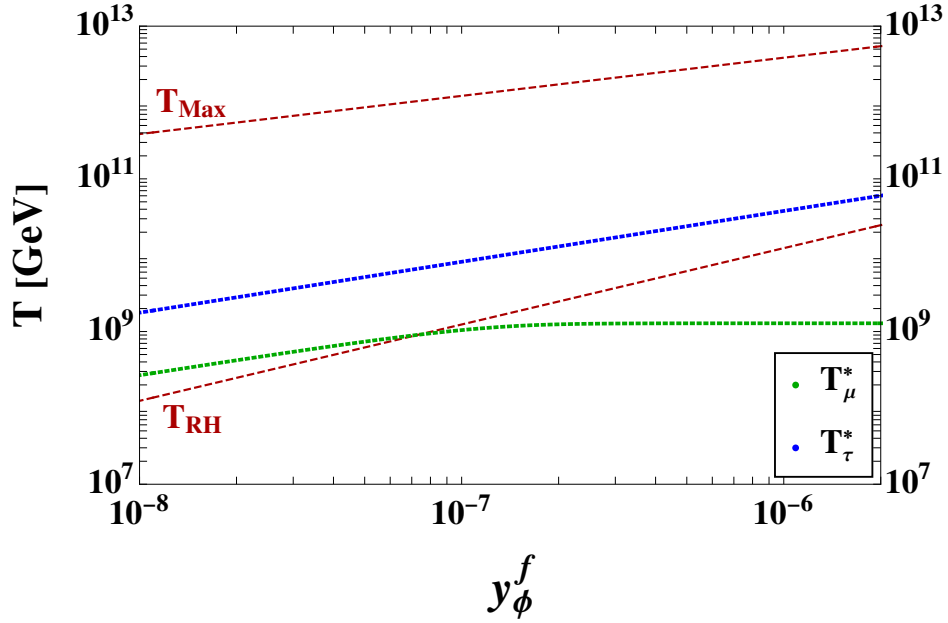
FIG. 1: Variation of T_{Max}, T_{RH} and ETs with respect to y_ϕ^f: The upper (lower) brown dashed line represents the variation of T_{Max} (T_{RH}) while T_τ^*, T_μ^*, are represented by blue and green dotted lines respectively.

the help of Eq. (19),

$$T_{\text{Max}} = \left[\frac{15 g_f^2 y_\phi^{f\,2}}{16 g_* \pi^3} \rho_{\text{end}}^{\frac{n-1}{n}} \sqrt{n(n-1)} \left(\frac{3n-3}{2n-4} \right)^{\frac{3(n-1)}{7-n}} M_P^{4/n} \right]^{1/4}. \tag{21}$$

Once T_{Max} is reached, the temperature profile starts falling as $T \propto a^{-\frac{3n-3}{2n+4}}$ and its continues to fall till an equality between matter and radiation energy density is reached. This point, known as the T_{RH}, can be estimated by comparing $\rho_\phi = \rho_R$. Using Eqs. (18) and (19), one obtains

$$T_{\text{RH}} = \left[(g_f y_\phi^f)^{2n} \frac{15 \{3n(n-1)\}^{n/2} \lambda}{2^{4n-1} \pi^{2+n} g_*} \left(\frac{n+2}{7-n} \right)^n M_P^4 \right]^{1/4}. \tag{22}$$

Following Eq.(21) and (22), one finds that $T_{\text{Max}} \propto y_\phi^{f\,1/2}$ while $T_{\text{RH}} \propto y_\phi^f$ (for $n = 2$). Such dependence becomes explicit in Fig. 1 where we plot T_{Max} (upper dashed brown line) and T_{RH} (bottom dashed brown line) as function of y_ϕ^f. We restrict ourselves with the choice of $y_\phi^f < 3.3 \times 10^{-6}$ as beyond this value, explosive production of fermions via preheating may take place, which we intend to avoid in the present work.

B. Equilibrium Temperature of e_{R_α} as function of y_ϕ^f

Next, we aim to identify the dependence of equilibration temperature of different charged lepton Yukawa interaction on y_ϕ^f. In order to estimate that, one may equate

$$\langle \Gamma_\alpha \rangle = \mathcal{H}, \tag{23}$$

with $\langle \Gamma_\alpha \rangle = \frac{\pi Y_\alpha^2}{192 \zeta(3)} \frac{m_h^2(T)}{T}$ associated to a specific Yukawa interaction Y_α. During the extended reheating era defined by $T_{\text{Max}} > T > T_{\text{RH}}$, due to the ρ_ϕ dominance, the Hubble can be approximated effectively by $\mathcal{H}^2 \simeq \frac{\rho_\phi}{3M_p^2}$. Thereafter, using $m_h(T) \simeq 0.6T$ in $\langle \Gamma_\alpha \rangle$ and Eq. (18) into Eq. (23), we write,

$$\frac{\pi Y_\alpha^2}{192 \zeta(3)} \frac{(0.6T)^2}{T} \simeq \sqrt{\frac{\rho_{\text{end}}}{3M_p^2} \left(\frac{a}{a_{\text{end}}} \right)^{-6n/(n+2)}}. \tag{24}$$

Further, in Eq. (24) one can replace the $a/a_{\rm end}$ dependence in terms of T using Eq. (19) and (20) so as to obtain an analytical expression of ET of respective charged lepton Yukawa interaction. At a much later stage (in the limit $a \gg a_{\rm end}$), the modified ET associated to Y_α in interaction T_α^* can then be written as,

$$T_\alpha^* = \mathbb{C}^{\frac{n-1}{n+1}} Y_\alpha^{\frac{2(n-1)}{n+1}} (g_f y_\phi^f)^{\frac{n}{n+1}}, \tag{25}$$

where $\mathbb{C}$ is a quantity independent of y_ϕ^f, given by

$$\mathbb{C} = 5\sqrt{3} \times 10^{-3} \lambda^{\frac{1}{2(n-1)}} M_p^{\frac{n+1}{n-1}} \left\{3n(n-1)\right\}^{\frac{n}{4(n-1)}} \left(\frac{n+2}{14-2n}\right)^{\frac{n}{2(n-1)}} \left(\frac{30}{8\pi^3 g_*}\right)^{\frac{n}{2(n-1)}}. \tag{26}$$

Such analytical estimate of the ETs as a function of y_ϕ^f during the extended period of reheating ($T_\alpha^* \propto y_\phi^{f\,2/3}$ with $n=2$) via Eq. (25) is in a good agreement with the numerical results we had in the main paper. The ETs for τ_R and μ_R are plotted in Fig.1 as a function of y_ϕ^f. Note that, e_R does not equilibrate above the $T_{\rm RH}$ obtained, within the specified range of y_ϕ^f considered here which means $T_e^* = T_{0(e)}^*$ beyond this point (in the radiation dominated era). As expected, both T_τ^* (blue dotted line) and T_μ^* (green dotted line) increase with the increase in y_ϕ^f during this extended reheating period. At some point with a typical value of y_ϕ^f, T_μ^* coincides with $T_{\rm RH}$ and thereafter T_μ^* saturates to its value $T_{0\mu}^*$ as in the radiation dominated era (indicated by the constant line). However, the Y_τ interaction continues to have modified ET throughout the entire range of y_ϕ^f considered.

At this stage, one may wonder what happens to the other SM interactions like the top, bottom, charm Yukawa interactions, EW sphaleron interactions, etc. In this regard, the point is to note that the source of the modified ET related to any interaction is related to $\mathcal{H}$ and hence all these other interactions should also be affected similarly. So overall a scaling effect $w.r.t.$ the charged lepton Yukawa ET should be observed. As the strength of these interactions is more compared to Y_τ, it is natural to expect these interactions to be in equilibrium by the time T_τ^* has reached.

III. LEPTON ASYMMETRY AS FUNCTION OF M_1 AND y_ϕ^f

With the above understanding on the $T_{\rm Max}$ and $T_{\rm RH}$ as a function of y_ϕ^f, it is evident that the parameter space for the lightest RHN of mass M_1 for case (b) of the main paper ($T_{\rm Max} > M_1 > T_{\rm RH}$) is limited within the region bounded by the $T_{\rm Max}$ and $T_{\rm RH}$ lines of Fig. 1. As a result, for the analysis of the lepton asymmetry generation from the decay of RHN N_1, we focus primarily on this range of M_1, indicated by the yellow shaded region of Fig. 2 corresponding to $y_\phi^f < 3.3 \times 10^{-6}$. Now, the lightest RHN mass M_1 being smaller than the temperature of the bath for the period: $T_{\rm Max} - T(\sim M_1)$ (for a specific y_ϕ^f), it should be allowed to be produced via inverse decay. Furthermore, for the rest of the reheating period, $i.e.$ for: $T(\sim M_1) - T_{\rm RH}$, its decay to the SM bath particles also contributes to the radiation. Hence, in addition to ρ_ϕ and ρ_R, contribution of ρ_{N_1} needs to be incorporated to study the evolution of different components and final lepton asymmetry. To estimate such evolutions, we involve a set of coupled Boltzmann equations as,

$$\frac{d\rho_\phi}{dt} + 3\mathcal{H}\rho_\phi = -\Gamma_\phi \rho_\phi, \tag{27a}$$

$$\frac{d\rho_R}{dt} + 4\mathcal{H}\rho_R = \Gamma_\phi \rho_\phi + \langle\Gamma_{N_1}\rangle(\rho_{N_1} - \rho_{N_1}^{\rm eq}), \tag{27b}$$

$$\frac{d\rho_{N_1}}{dt} + 3\mathcal{H}\rho_{N_1} = -\langle\Gamma_{N_1}\rangle(\rho_{N_1} - \rho_{N_1}^{\rm eq}), \tag{27c}$$

$$\frac{dn_\Delta}{dt} + 3\mathcal{H}n_\Delta = -\langle\Gamma_{N_1}\rangle\left[\frac{\varepsilon_\ell}{M_1}(\rho_{N_1} - \rho_{N_1}^{\rm eq}) + \frac{n_{N_1}^{\rm eq}}{2n_\ell^{\rm eq}}n_\Delta\right], \tag{27d}$$

where we consider the case with $n=2$, $\mathcal{H} = \sqrt{\frac{\rho_\phi + \rho_R + \rho_{N_1}}{3M_P^2}}$ and $n_\Delta = n_{B-L} \equiv -n_{\Delta_L}$ as defined in Eq. 3. For solution purpose, it is found to be convenient to replace the t-dependance in terms of scale factor a using the relation $\frac{d}{da} \equiv \frac{1}{a\mathcal{H}}\frac{d}{dt}$. The transformed equations are mentioned in the main paper.

At this stage, we recollect that apart from M_1 and y_ϕ^f, the lepton asymmetry, generated as a result of N_1 decay, is also controlled by the neutrino Yukawa coupling matrix Y_ν, via Γ_{N_1} in the BEs above. Introducing CI-parametrisation
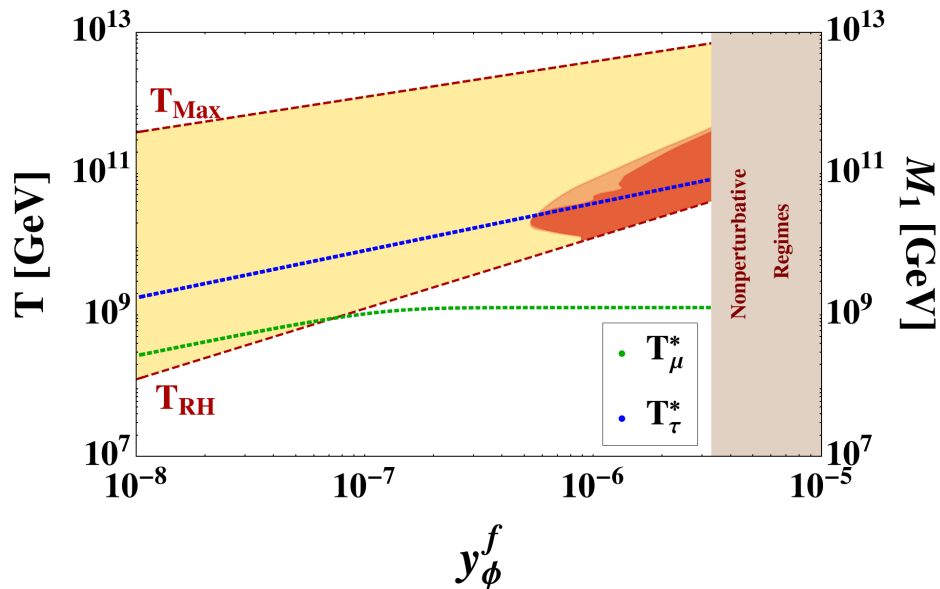
FIG. 2: Dark brown patch represents the allowed value of RHN mass M_1 and coupling y_ϕ^f which can produce the correct baryon asymmetry for the scenario with modified flavor regimes. On the contrary, light brown patch represents the allowed region when change in flavor effect was not taken into account. Here $y_\phi^f > 3.3 \times 10^{-5}$ indicates the nonperturbative regime.

makes Y_ν a function of M_1, $\text{Re}[\theta_R]$ and $\text{Im}[\theta_R]$, provided the heavier RHN masses carry a specific ratio with M_1 together with the neutrino oscillation parameters fixed from experimental data (see the main text for the details). The inflation related parameters such as m_ϕ and hence Γ_ϕ are so evaluated as to satisfy the inflationary predictions in terms of spectral index n_s and tensor to scalar ratio r, as stated in the main paper. With this understanding, we now proceed below to discuss the modification expected on flavor leptogenesis as a result of shift(s) in the ET(s) during extended reheating period and thereafter provide a guideline for scanning the parameter space that could correctly produce the observed baryon asymmetry.

While solving the above BEs to evaluate the lepton asymmetry, the effect of individual lepton flavor interactions enters in the picture if one of the charged lepton Yukawa interactions, say Y_τ, becomes comparable to the neutrino Yukawa interaction. As the RHN N_1 stays out of equilibrium during its decay for lepton asymmetry production, checking of equilibration condition for tau Yukawa interaction is sufficient to identify whether these interactions are strong enough. In such a situation, lepton (anti-lepton) states produced from the RHN decay start interacting with right-handed tau leptons and as a result, asymmetry produced in the tau flavor direction gets enhanced or diluted differently compared to other flavor directions. This can impact the total lepton asymmetry. Hence, one needs to take into account the time evolution of the lepton asymmetry along the tau flavor state and a coherent superposition state of muon and electron separately which is the core of flavor leptogenesis. On the other hand, at a lower temperature, when μ-Yukawa interaction becomes dominant over the expansion rate of the Universe, the quantum coherence of the lepton states produced from the RHN decay breaks down completely and one needs to look for the evolution of the lepton asymmetry along e, μ and τ separately. The BEs that describe such scenarios is given by (the Eq. [10] of the main paper),

$$\frac{dn_{\Delta_\alpha}}{dt} + 3\mathcal{H}n_{\Delta_\alpha} = -\langle\Gamma_{N_1}\rangle \left[\frac{\varepsilon_{\ell_\alpha}}{M_1}(\rho_{N_1} - \rho_{N_1}^{\text{eq}}) + \frac{1}{2}K_\alpha^0 \sum_\beta (C_{\alpha\beta}^\ell + C_\beta^H)\frac{n_{N_1}^{\text{eq}}}{n_\ell^{\text{eq}}}n_{\Delta_\beta} \right], \tag{28}$$

where $\alpha = e, \kappa$ when only Y_τ is in thermal equilibrium, while $\alpha = e, \mu, \tau$ when both Y_τ and Y_μ become non negligible. Here, $\varepsilon_{\ell_\alpha}$ denotes the CP asymmetry produced in the relevant lepton flavor directions, while K_α^0 represents the ratio between tree level decay rate of lightest RHN to relevant lepton flavor α and total decay rate of N_1. The asymmetry produced in the Higgs and lepton doublet sector is converted to $B/3 - L_\alpha$ number asymmetry via C^ℓ and C^H matrices.

We have already noticed that the parameters controlling the lepton asymmetry production are essentially: y_ϕ^f, M_1, $\text{Re}[\theta_R]$ and $\text{Im}[\theta_R]$. Among these, y_ϕ^f is bounded from above by 3.3×10^{-6} as stated before while M_1 is constraint

to be between $T_{\rm Max}$ and $T_{\rm RH}$ (for a specific y_ϕ^f) as indicated by the yellow patch of Fig. 2 for case (b) (the primary focus of the main paper). Keeping these in mind, we solve the BEs (Eq. (27a)-(27c) along with Eq. (27d) or Eq. (28) depending on the flavor regimes) for $y_\phi^f < 3.3 \times 10^{-6}$ and M_1 being scanned across the yellow region of Fig. 2 while allow ${\rm Re}[\theta_R]$ and ${\rm Im}[\theta_R]$ to vary arbitrarily with the sole constraint that such values must not make any element of Y_ν non-perturbative *i.e.* $(Y_\nu)_{ij} < \mathcal{O}(1)$.

It is found that only a stipulated parameter space, as indicated by the dark brown patch in Fig. 2, satisfies the correct amount of final baryon asymmetry. The information on the RHN neutrino mass M_1 is indicated in the right side Y-axis. Note that, this parameter space is inclusive of the shift in the equilibration temperature(s) of the charged lepton Yukawa interaction(s) that we observe in this work. To make a comparison, we include the light brown patch as well which corresponds to the parameter space producing correct amount of final baryon asymmetry production, however, without taking into account the shift in the ETs during this extended era of reheating. We notice that the parameter space is modified and more constrained when the impact of shift in ETs is correctly embedded in flavor leptogenesis scenario. In the main paper, we discuss such impact on lepton asymmetry production, due to modified (smaller compared to standard case) ETs, with the choice of a specific benchmark point due to the limited space. Here however, we scan the entire parameter space which in fact provides a clear confirmation about the importance of re-evaluating the charged lepton Yukawa equilibration during the reheating phase when lepton asymmetry production is simultaneously taking place.

We further notice that the blue dashed line, representative of the ET of tau leptons T_τ^*, goes through the dark brown patch and marks a demarkation between the unflavored and two flavored regimes of thermal leptogenesis. Note that for M_1 values falling above this line, N_1 decay effectively contributes to lepton asymmetry production while none of the charged lepton Yukawa interaction enters in equilibrium and as a result this case correspond to unflavored regime of leptogenesis. Contrary to this, for M_1 values lying below this line, N_1 decay starts after the Y_τ interaction equilibrates so as to consider τ and κ (a linear combination of e and μ) flavor directions separately as for such M_1, the T_μ^* remains below the $T_{\rm RH}$. Hence, these (M_1 below the blue dashed line) points belong to two flavor regime of thermal leptogenesis. Looking at this plot, we conclude that successful leptogenesis can be achieved with a minimum value of $T \sim M_1 = 1.3 \times 10^{10}$ GeV and $T_{\rm RH} = 1.1 \times 10^{10}$ GeV. At this stage, it is also pertinent to mention that the energy density of the RHN always remains subdominant to the radiation energy density. This is due to the fact that the RHN N_1 exists in the Universe for a very short period of time as it has to be produced and decay in between the era of $T_{\rm max}$ and $T_{\rm RH}$. As a result of which, the reheating temperature remains effectively unchanged even in the presence of RHN. This behavior is also depicted in Fig. 3 of the main text. A deviation of such a situation is discussed in a follow-up paper [46].