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Dirac leptogenesis from asymmetry wash-in via scatterings

Tomáš Blažek,[1,*] Julian Heeck,[2,†] Jan Heisig,[3,2,‡] Peter Maták,[1,§] and Viktor Zaujec[1,¶]

[1]Department of Theoretical Physics, Comenius University, Mlynská dolina, 84248 Bratislava, Slovak Republic

[2]Department of Physics, University of Virginia, Charlottesville, Virginia 22904-4714, USA

[3]Institute for Theoretical Particle Physics and Cosmology, RWTH Aachen University, D-52056 Aachen, Germany

Leptogenesis typically requires the introduction of heavy particles whose out-of-equilibrium decays are essential for generating a matter-antimatter asymmetry, according to one of Sakharov's conditions. We demonstrate that in Dirac leptogenesis, scatterings between the light degrees of freedom – Standard Model particles plus Dirac neutrinos – are sufficient to generate the asymmetry. Due to its vanishing source term in the Boltzmann equations, the asymmetry of right-handed neutrinos solely arises through wash-in processes. Sakharov's conditions are satisfied because the right-handed neutrino partners are out of equilibrium. Consequently, heavy degrees of freedom never needed to be produced in the early universe, allowing for a reheating temperature well below their mass scale. Considering a minimal leptoquark model, we discuss the viable parameter space along with the observational signature of an increased number of effective neutrinos in the early universe.

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I. INTRODUCTION

The discovery of nonzero neutrino masses through neutrino oscillations provides clear evidence of physics beyond the Standard Model (SM) and necessitates the introduction of new particles. Thus far, experiments have not resolved the fundamental property of neutrinos as either Dirac or Majorana particles, leaving two qualitatively different scenarios for their mass generation.

Majorana neutrinos, particularly when implemented in a seesaw mechanism, often offer suitable conditions for baryogenesis via leptogenesis [1, 2]. Essentially, the $\Delta L = 2$ interactions responsible for generating Majorana neutrino masses can induce a lepton asymmetry in the early universe, which can then be converted into a baryon asymmetry via sphalerons [3]. According to Sakharov's conditions [4], this process necessitates lepton number and CP violation, as well as out-of-equilibrium dynamics, typically achieved through the freeze-out or freeze-in of the heavy seesaw states. The ability to simultaneously account for neutrino masses and the baryon asymmetry is a significant attraction of Majorana neutrino models.

However, Dirac neutrinos also have the potential to generate a matter-antimatter asymmetry, as demonstrated in Ref. [5]. In this scenario, the smallness of the Dirac-neutrino mass term effectively decouples the right-handed neutrino partners ν_R from the rest of the SM plasma. Without ever violating the lepton number, it becomes possible to create an effective lepton asymmetry by hiding an opposite asymmetry in the decoupled ν_R sector [5]. Similar to the seesaw case, fulfilling Sakharov's out-of-equilibrium condition typically involves the decay of heavy particles [5, 6].

Interestingly, since the ν_R themselves are out of equilibrium, Dirac leptogenesis offers the possibility to satisfy Sakharov's conditions without actually involving the decay of heavy particles, X_i. As we show in this paper, it opens up a new variant of the mechanism in which the temperature of our universe never reached the heavy particle's mass scale while generating the asymmetry through freeze-in due to $2 \to 2$ scatterings of SM particles and ν_R. However, CPT invariance and unitarity constraints [7–10] introduce additional conditions on the viability of the simple Dirac-leptogenesis models outlined in Ref. [6]. In particular, flavor-blind models require the existence of at least three different final states in the scattering of ν_R for a non-vanishing asymmetry, as present in the case of X_i being a leptoquark, *i.e.* model c in [6]. While the source term for the ν_R asymmetry in the Boltzmann equation is still vanishing, an asymmetry can be produced via wash-in [11]. In fact, we show that this contribution is large enough to explain the observed baryon asymmetry, $Y_{\Delta B} \simeq 0.9 \times 10^{-10}$ [2, 12], rendering Dirac leptogenesis via scatterings a viable scenario. To the best of our knowledge, the underlying mechanism – a combination of wash-in and freeze-in of the nonequilibrium particle density – has not been considered in the literature before.

The remainder of the paper is structured as follows. In Sec. II, we introduce the leptoquark model used throughout the study. We compute the scattering asymmetry and discuss the Boltzmann equations in Secs. III and IV, respectively. The implications for the baryon asymmetry are discussed in Sec. V. We conclude in Sec. VI.

II. SIMPLE MODEL

The simple renormalizable Dirac-leptogenesis models considered in Ref. [6] can successfully explain the ob-

*E-mail: tomas.blazek@fmph.uniba.sk
†E-mail: heeck@virginia.edu; ORCID: 0000-0003-2653-5962.
‡E-mail: heisig@physik.rwth-aachen.de
§E-mail: peter.matak@fmph.uniba.sk
¶E-mail: viktor.zaujec@fmph.uniba.sk

served baryon asymmetry when assuming that the universe was hot enough to thermally produce a significant number of on-shell heavy mediator particles. However, here we are interested in the alternative scenario of a universe whose temperature never reached the mediators' mass scale in its entire (post-inflationary) history. In this case, on-shell mediator production from the thermal bath is highly suppressed by the high-energy tails of thermal distributions and is, hence, negligible. Accordingly, we could equally well do away with them altogether and describe their effects with effective operators, similar to Ref. [13]. However, we opt to keep the mediator particles as degrees of freedom in our theory to facilitate direct comparison with the mediator decay case from Ref. [6].

As a simple model that contains all ingredients for successful Dirac leptogenesis via scattering, we introduce several scalar particles $X_i \sim (\mathbf{3}, \mathbf{1}, -1/3)$, *i.e.* the same gauge quantum numbers as the right-handed down quark, or the S_1 leptoquark in the notation of Ref. [14]. The relevant Yukawa interactions take the form

$$\mathcal{L} = \bar{Q}^c F_i L \bar{X}_i + \bar{d}^c_R G_i \nu_R \bar{X}_i + \bar{u}^c_R K_i e_R \bar{X}_i + \text{H.c.}\,, \quad (1)$$

where we have suppressed gauge and flavor indices and assumed the X_i to be mass eigenstates with masses M_i. Focusing on a minimal model allowing for a successful leptogenesis, we only consider two copies of X, $i = 1, 2$, and take the coupling matrices to be proportional to the identity matrix in flavor space. To simplify our discussion, we have also imposed a baryon-number symmetry that forbids the di-quark couplings QQX and $u_R d_R X$, *cf.* case *c* among the renormalizable Dirac-leptogenesis models from Ref. [6]. Upon assigning $(B-L)(X_i) = -2/3$, the above Lagrangian is $B-L$ conserving. With these choices, the X_i do not mediate any processes violating baryon number, lepton number, or lepton flavor.

While $B-L$ is hence conserved over the entire history of our universe, ν_R number need not be, allowing sphalerons [3] – which are blind to the ν_R – to convert the matching asymmetry $Y_{\Delta\nu_R} = Y_{\Delta(B-L_{\text{SM}})}$ into a baryon asymmetry [3, 15]

$$Y_{\Delta B} = \frac{28}{79} Y_{\Delta(B-L_{\text{SM}})} = \frac{28}{79} Y_{\Delta\nu_R}\,. \quad (2)$$

The interactions in $\mathcal{L}$ indeed break ν_R number if G and either F or K are nonzero, seemingly allowing for the production of a CP asymmetry in ν_R. As we will show below, the necessary condition for a ν_R asymmetry is actually more subtle and requires the simultaneous presence of all three coupling matrices in $\mathcal{L}$.

III. SCATTERING CROSS SECTIONS AND ASYMMETRIES

Right-handed neutrino interactions enter the matter–antimatter asymmetry evolution through thermally averaged cross sections. These are computed using [16]

$$\begin{aligned}\langle\sigma v\rangle_{12\to34} &= \frac{g_1 g_2}{n_1^{\text{eq}} n_2^{\text{eq}}} \left(\prod_{a=1}^{4} \int \frac{\mathrm{d}^3 p_a}{(2\pi)^3 2E_a} \right) \mathrm{e}^{-\frac{E_1+E_2}{T}} \\ &\quad \times (2\pi)^4 \delta^{(4)}(p_1+p_2-p_3-p_4) |\bar{\mathcal{M}}|^2 \\ &= \frac{g_1 g_2 T}{32\pi^4 n_1^{\text{eq}} n_2^{\text{eq}}} \int \mathrm{d}s\, s^{3/2} \sigma(s) K_1\!\left(\frac{\sqrt{s}}{T}\right),\end{aligned} \quad (3)$$

where we employ the Maxwell–Boltzmann densities and neglect the Pauli blocking factors as an approximation.[1]

All fermions in our model are considered massless. Therefore, in equilibrium, their densities n_a^{eq} can be written as a product of the right-handed neutrino equilibrium density $n_{\nu_R}^{\text{eq}} = 3\zeta(3)T^3/(4\pi^2)$ and the number of degrees of freedom g_a.

The interactions are mediated by heavy X_i scalars. In the s-channel, the cross sections contain singularities for $s = M_i^2$. These are usually regularized by inserting a finite mediator width, leading to double counting the on-shell X_i production, which must be subtracted in the Boltzmann equation [7]. Alternatively, one may follow the procedures introduced in Refs. [17, 18]. Focusing on temperatures much lower than the mediator masses, we neglect the energy dependence of the propagators, disregarding the region in which the singularities occur. Within this approximation, the lowest-order cross sections of the t- and u-channel processes are equal to one-third of the corresponding s-channel total cross sections. We will use this relation to simplify our Boltzmann equations.

For reactions containing two massless particles in the initial and final states, the CPT symmetry allows us to define asymmetries of total cross sections as

$$\Delta\sigma_{12\to34} = \sigma_{12\to34} - \sigma_{34\to12} \quad (4)$$

obeying the unitarity constraints [7–10]

$$\sum_f \Delta\sigma_{12\to f} = 0\,, \quad (5)$$

where we sum over all possible final states. Relations of this type are essential to guarantee vanishing asymmetry in thermal equilibrium as required by Sakharov's conditions [4].

To see how unitarity and CPT symmetry constrain the asymmetries in our model, let us look at the $\nu_R d_R \to LQ$ cross section. The asymmetry follows from the interference of the tree and loop diagrams in Figs. 1a and 1b,

[1] If this is not the case and one wishes to keep the quantum statistics, it is necessary to include analogous statistical factors to cutting rules and asymmetry calculation. Otherwise, the CPT and unitarity constraints [7–10] and, consequently, the Sakharov conditions [4], will be violated.

respectively, and is proportional to

$$\sum_{i,j,k} \frac{\Im[\mathrm{tr}(G_j^\dagger G_i)\mathrm{tr}(F_i^\dagger F_k)\mathrm{tr}(K_k^\dagger K_j)]}{M_i^2 M_j^2 M_k^2}\,, \tag{6}$$

where we implicitly sum over flavors. The traces in the numerator of Eq. (6) can be viewed as 2×2 Hermitian matrices and can be parametrized as[2]

$$\mathrm{tr}(F_i^\dagger F_j) = \frac{M_i M_j}{T_{\mathrm{reh}}^2} \sum_a f_a (\sigma_a)_{ij}\,, \tag{7}$$

where $f_a = (f_0, \boldsymbol{f})$ are real dimensionless components of a Euclidean four-vector, $\sigma_a = (1, \boldsymbol{\sigma})$, and $\boldsymbol{\sigma}$ is the vector of Pauli matrices. In the flavor-blind limit employed here, $f_0 = |\boldsymbol{f}|$; furthermore, since $T_{\mathrm{reh}} \ll M_i$ by assumption, $|f_a| \ll 1$ in the perturbative-coupling regime. Analogously rewriting $\mathrm{tr}(G_i^\dagger G_j)$ and $\mathrm{tr}(K_i^\dagger K_j)$, thermal averaging yields the asymmetry

$$\Delta\langle\sigma_1 v\rangle = \frac{512}{\pi^2}\frac{T^4}{T_{\mathrm{reh}}^6}\frac{\boldsymbol{g}.(\boldsymbol{f}\times\boldsymbol{k})}{\zeta(3)^2} \equiv \frac{512}{\pi^2}\frac{T^4}{T_{\mathrm{reh}}^6}\frac{\varepsilon}{\zeta(3)^2}\,, \tag{8}$$

where the subscript 1 refers to the $\nu_R d_R \to LQ$ reaction; the cross sections of the other s-channel processes, $\nu_R d_R \to e_R u_R$ and $e_R u_R \to LQ$, will be further labeled by 2 and 3, respectively. Here, we introduced ε as a convenient dimensionless measure of CP violation, further discussed below.

From the diagrams in Figs. 1c and 1d we obtain the asymmetry in the $\nu_R d_R \to e_R u_R$ cross section proportional to $\boldsymbol{g}.(\boldsymbol{k}\times\boldsymbol{f}) = -\varepsilon$ and clearly

$$\Delta\langle\sigma_1 v\rangle + \Delta\langle\sigma_2 v\rangle = 0\,. \tag{9}$$

For the $\nu_R d_R$ initial state, the only two-particle final states allowed by the Lagrangian density are LQ, $e_R u_R$, and $\nu_R d_R$. The total-cross-section asymmetry of the elastic scattering must vanish by the CPT symmetry. Therefore, from the three possible final states mentioned in Sec. I, only two have nonzero asymmetries. These must be of opposite sign and cancel each other due to the unitarity constraints, as we observe in Eq. (9).[3]

Finally, the tree-level symmetric parts of the s-channel cross sections result in thermal averages

$$\langle\sigma_1 v\rangle = \frac{64}{3\pi}\frac{T^2}{T_{\mathrm{reh}}^4}\frac{f_0 g_0 + \boldsymbol{f}.\boldsymbol{g}}{\zeta(3)^2} \equiv \frac{64}{3\pi}\frac{T^2}{T_{\mathrm{reh}}^4}\frac{\alpha_1}{\zeta(3)^2}\,, \tag{10}$$

$$\langle\sigma_2 v\rangle = \frac{32}{3\pi}\frac{T^2}{T_{\mathrm{reh}}^4}\frac{k_0 g_0 + \boldsymbol{k}.\boldsymbol{g}}{\zeta(3)^2} \equiv \frac{32}{3\pi}\frac{T^2}{T_{\mathrm{reh}}^4}\frac{\alpha_2}{\zeta(3)^2}\,, \tag{11}$$

$$\langle\sigma_3 v\rangle = \frac{64}{3\pi}\frac{T^2}{T_{\mathrm{reh}}^4}\frac{f_0 k_0 + \boldsymbol{f}.\boldsymbol{k}}{\zeta(3)^2} \equiv \frac{64}{3\pi}\frac{T^2}{T_{\mathrm{reh}}^4}\frac{\alpha_3}{\zeta(3)^2}\,, \tag{12}$$

[2] A similar parametrization was used to study CP violation in models containing multiple Higgs doublets [19, 20].

[3] In more complex models, the CPT and unitarity constraints can be made explicit using the cyclic diagrams of Ref. [21] or holomorphic cutting rules [22–24].

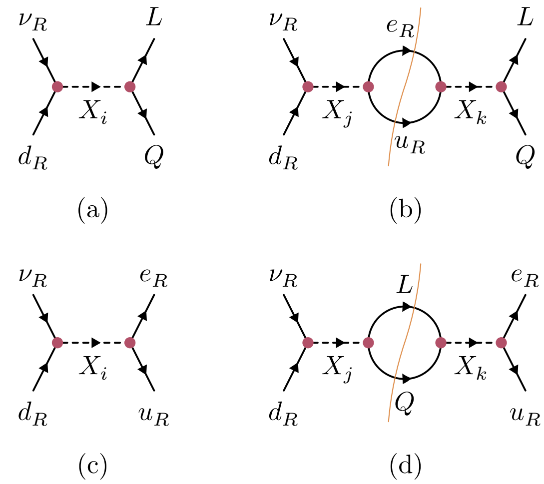


FIG. 1: Tree and loop diagrams contributing to the $\nu_R d_R \to LQ$ and $\nu_R d_R \to e_R u_R$ asymmetries. The cuts represent the imaginary parts of the loop diagrams.

where we introduced $\alpha_{1,2,3} \geq 0$ as useful measures of cross-section strength. In the flavor-blind limit, we have the inequality

$$|\varepsilon| \leq \sqrt{2\alpha_1\alpha_2\alpha_3} \tag{13}$$

and our small-asymmetry calculation requires

$$|\varepsilon| \ll \frac{\pi}{12}\alpha_1, \frac{\pi}{24}\alpha_2, \frac{\pi}{12}\alpha_3\,. \tag{14}$$

The α_i are essentially free parameters, allowing, in particular, for large hierarchies.

IV. BOLTZMANN EQUATIONS

To study particle densities and asymmetries, we follow the conventions of Ref. [6] and introduce

$$\Sigma_a = \frac{n_a + n_{\bar a}}{s}\,, \quad \Delta_a = \frac{n_a - n_{\bar a}}{s}\,, \tag{15}$$

where $s = h_* T^3\, 2\pi^2/45$ denotes the entropy density. These quantities are summed over all flavors for each particle species. Since the baryon-asymmetry generation involves sphalerons, we require the reheating temperature T_{reh} to be above their decoupling temperature of 130 GeV [25]. As an approximation, we assume the SM particles are in thermal equilibrium immediately after reheating and set the effective number of relativistic degrees of freedom entering the energy and entropy density to constant $g_* = h_* = 106.75$. This assumption is apparently unrealistic but serves as a benchmark scenario to showcase the relevant dynamics. A more realistic reheating phase is expected to introduce $\mathcal{O}(1)$ corrections but not to change the qualitative picture.

As the initial abundance of right-handed neutrinos is negligible by assumption, they are produced via $LQ \to$

$\nu_R d_R$, $e_R u_R \to \nu_R d_R$, and crossed reactions. The symmetric part of their density obeys the Boltzmann equation

$$\frac{\mathrm{d}\Sigma_{\nu_R}}{\mathrm{d}x} = -\frac{1}{x^4}\frac{\Gamma}{\mathcal{H}}\bigg|_{T_{\rm reh}}\left(\Sigma_{\nu_R} - \Sigma_{\nu_R}^{\rm eq}\right), \qquad (16)$$

where $x \equiv T_{\rm reh}/T$, $\mathcal{H} = \pi\sqrt{g_*/90}\,T^2/M_{\rm Pl}$ is the Hubble parameter with reduced Planck mass $M_{\rm Pl} \simeq 2.4 \times 10^{18}\,\mathrm{GeV}$, and

$$\Gamma = \frac{5}{18} s\Sigma_{\nu_R}^{\rm eq}(\langle\sigma_1 v\rangle + \langle\sigma_2 v\rangle) \qquad (17)$$

parametrizes the ν_R interaction rate. We solve Eq. (16) analytically and obtain

$$\Sigma_{\nu_R}(x) = \frac{405\zeta(3)}{4\pi^4 h_*}\left(1 - \exp\left[-\frac{\Gamma}{\mathcal{H}}\bigg|_{T_{\rm reh}}\frac{x^3-1}{3x^3}\right]\right), \quad (18)$$

which rises steeply and quickly flattens out for $x \gg 1$, illustrated in Fig. 2, reaching the equilibrium value for large interaction rates, $\Gamma/\mathcal{H}|_{T_{\rm reh}} \gg 1$, corresponding to

$$2\alpha_1 + \alpha_2 \gg \frac{3\sqrt{g_*}\pi^4 T_{\rm reh}\zeta(3)}{40\sqrt{10}M_{\rm Pl}} \simeq \frac{T_{\rm reh}}{8\times 10^{16}\,\mathrm{GeV}}\,. \quad (19)$$

Notice that the α_i themselves scale with $T_{\rm reh}^4$ times Lagrangian parameters.

The evolution of an asymmetry generally results from two competing parts. The source term contains asymmetries of thermally averaged cross sections, whereas the wash-out terms are proportional to particle density asymmetries Δ_a. The latter are subject to constraints implied by the symmetries and conservation laws of the model. In particular, the interactions in Eq. (1) separately conserve the lepton and baryon numbers, resulting in

$$\Delta_{\nu_R} = -\Delta_{e_R} - \Delta_L\,, \quad \Delta_{d_R} = -\Delta_{u_R} - \Delta_Q\,. \quad (20)$$

Moreover, there are three other global $U(1)$ symmetries: considering $\nu_R \to \mathrm{e}^{\mathrm{i}\alpha}\nu_R$ with $d_R \to \mathrm{e}^{-\mathrm{i}\alpha}d_R$ leaves the Lagrangian density invariant, and the same independently applies to e_R, u_R and Q, L field transformations. The conservation of the respective charges then implies

$$\Delta_{\nu_R} = \Delta_{d_R}\,, \quad \Delta_{e_R} = \Delta_{u_R}\,, \quad \Delta_L = \Delta_Q\,. \quad (21)$$

Therefore, it is sufficient to solve the Boltzmann equations for Δ_L and Δ_{e_R} asymmetries, while the others are fixed through Eqs. (20) and (21).

It has been argued in Ref. [26] that to calculate the source term of an asymmetry, correspondingly charged particles must be produced in contributing reactions, while in agreement with Sakharov's conditions [4], the initial states must be out of equilibrium. In our scenario, the only out-of-equilibrium particles present in the universe are right-handed neutrinos. Once they are in the initial state, for the model in Eq. (1), they cannot be produced in the same $2 \to 2$ reaction. The source term of the right-handed neutrino asymmetry thus vanishes, but, fortunately, it does not mean that no asymmetry in their density can be produced. To observe this, let us finally write the Boltzmann equations for the density asymmetries,

$$\begin{aligned}\frac{\mathrm{d}\Delta_L}{\mathrm{d}x} &= \frac{\Sigma_{\nu_R}^{\rm eq}}{3\mathcal{H}}\frac{\mathrm{d}s}{\mathrm{d}x}\bigg\{\frac{1}{12}\Delta\langle\sigma_1 v\rangle\left(\Sigma_{\nu_R}^{\rm eq} - \Sigma_{\nu_R}\right) \qquad (22)\\ &\quad + \frac{5}{27}\langle\sigma_3 v\rangle\left(\Delta_L - 2\Delta_{e_R}\right)\\ &\quad + \frac{4}{27}\langle\sigma_1 v\rangle\bigg[\left(\Delta_L - \frac{17}{8}\Delta_{\nu_R}\right)\\ &\quad + \frac{1}{4}\frac{\Sigma_{\nu_R}}{\Sigma_{\nu_R}^{\rm eq}}\left(\Delta_L - \frac{3}{2}\Delta_{\nu_R}\right)\bigg]\bigg\},\\ \frac{\mathrm{d}\Delta_{e_R}}{\mathrm{d}x} &= \frac{\Sigma_{\nu_R}^{\rm eq}}{3\mathcal{H}}\frac{\mathrm{d}s}{\mathrm{d}x}\bigg\{\frac{1}{12}\Delta\langle\sigma_2 v\rangle\left(\Sigma_{\nu_R}^{\rm eq} - \Sigma_{\nu_R}\right) \qquad (23)\\ &\quad + \frac{5}{27}\langle\sigma_3 v\rangle\left(2\Delta_{e_R} - \Delta_L\right)\\ &\quad + \frac{4}{27}\langle\sigma_2 v\rangle\bigg[\left(2\Delta_{e_R} - \frac{17}{8}\Delta_{\nu_R}\right)\\ &\quad + \frac{1}{4}\frac{\Sigma_{\nu_R}}{\Sigma_{\nu_R}^{\rm eq}}\left(2\Delta_{e_R} - \frac{3}{2}\Delta_{\nu_R}\right)\bigg]\bigg\}.\end{aligned}$$

The first lines on the right-hand sides of Eqs. (22) and (23) correspond to the source terms containing the cross section asymmetries. According to Eq. (9), they only differ in sign. The lines below contain the wash-out terms which are linear in the density asymmetries Δ_L, Δ_{e_R}, and $\Delta_{\nu_R} = -\Delta_L - \Delta_{e_R}$. Again, the contributions of $e_R u_R \to LQ$ and crossed reactions enter Eqs. (22) and (23) with opposite signs. If there were no other contributions, we would obtain $\Delta_L = -\Delta_{e_R}$ and vanishing Δ_{ν_R}. However, the remaining wash-out terms depend on different combinations of Δ_L and Δ_{e_R}. According to Eq. (20), this allows for a nonzero and, in fact, sufficiently large Δ_{ν_R} which is washed-in [11], even though its source term vanishes identically. The only exception occurs for $\alpha_1 = \alpha_2$, or $\langle\sigma_1 v\rangle = 2\langle\sigma_2 v\rangle$. In this case, we can add Eq. (22) to Eq. (23) and obtain $\mathrm{d}\Delta_{\nu_R}/\mathrm{d}x$ proportional to $-\Delta_{\nu_R}$, erasing any ν_R asymmetry from the beginning.

The Boltzmann equations only depend on the ratios $\alpha_i M_{\rm Pl}/(\sqrt{g_*}T_{\rm reh})$ and $\varepsilon M_{\rm Pl}/(h_*\sqrt{g_*}T_{\rm reh})$, so the dependence on $T_{\rm reh}$ (and $M_{\rm Pl}$, g_*, and h_*) can be completely scaled away. Furthermore, the Boltzmann equations have been derived under the approximation of *small* asymmetries, which immediately implies that all asymmetries are directly proportional to $\varepsilon M_{\rm Pl}/(h_*\sqrt{g_*}T_{\rm reh})$, in particular $\Delta_{\nu_R} \propto \varepsilon M_{\rm Pl}/(h_*\sqrt{g_*}T_{\rm reh})$. This effectively leaves the three $\alpha_i M_{\rm Pl}/(\sqrt{g_*}T_{\rm reh})$ as free parameters in the Boltzmann equations.

Numerical solutions for $T_{\rm reh} = 10^8\,\mathrm{GeV}$ and various coupling values are shown in Fig. 2. In each of the three cases, fast ν_R production can be observed at the very beginning. In Fig. 2a, the solution corresponds to weak right-handed neutrino interactions with the SM bath. Their density never reaches the equilibrium value and

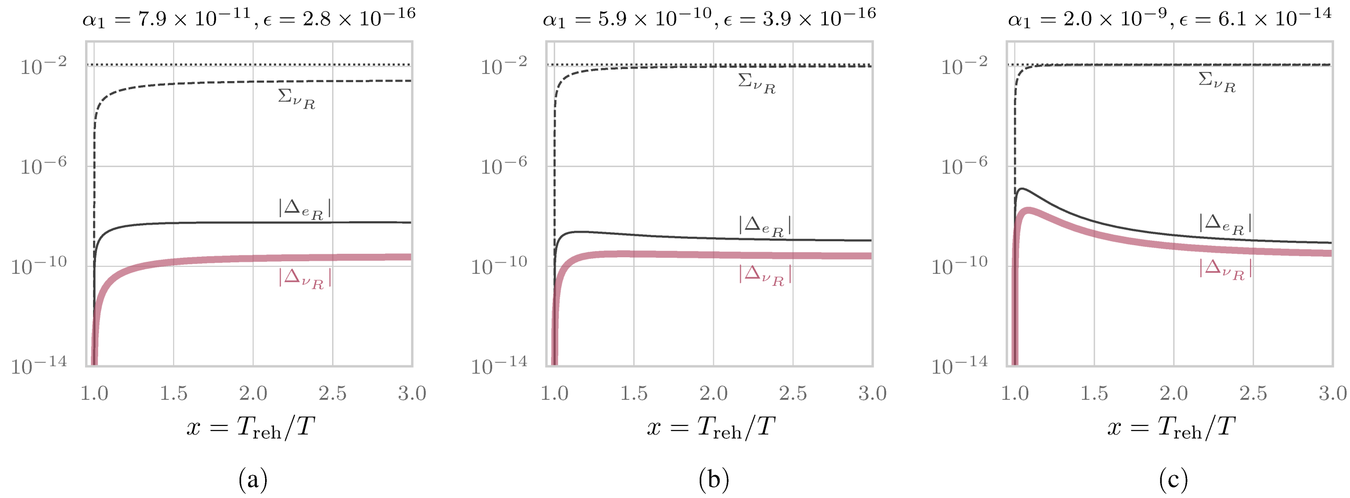


FIG. 2: Evolution of $\Sigma_{\nu_R}(x)$ (black dashed line), $\Delta_{e_R}(x)$ (black solid line), and $\Delta_{\nu_R}(x)$ (thick red line) for multiple parameter values and $T_{\rm reh} = 10^8$ GeV. The horizontal dotted line corresponds to $\Sigma_{\nu_R}^{\rm eq} \simeq 10^{-2}$. In all three cases, the ratio of the wash-out cross sections was fixed by $\alpha_2 = 2\alpha_1$ and $\alpha_3 = \alpha_1$, and ε was chosen to obtain the observed baryon asymmetry.

freezes in at $x \simeq 1.5$. The same happens to $|\Delta_L|$ and $|\Delta_{e_R}|$, which are roughly the same, while Δ_{ν_R} is suppressed by three orders of magnitude as it only comes from their tiny difference.

In Figs. 2b and 2c, at stronger couplings, both source and wash-out cross sections are enhanced, and initially produced asymmetries are larger. However, when right-handed neutrinos reach equilibrium, the source term vanishes, and the asymmetries are partially washed out.

V. BARYON ASYMMETRY

While the differential equations for Δ_L and Δ_{e_R} are not analytically solvable in general, we can obtain an analytic approximation in the limit of negligible Σ_{ν_R}, *i.e.* in the freeze-in regime:

$$\begin{aligned}\Delta_{\nu_R}(\infty) \simeq{}& \frac{23328\,\zeta(3)(\alpha_1-\alpha_2)(\Gamma/\mathcal{H})|_{T_{\rm reh}}\,\varepsilon}{\pi^5 h_*(2\alpha_1+\alpha_2)(50\alpha_1+\alpha_2-60\alpha_3)} \\ &\times\left[\xi\left(\frac{50\alpha_1+17\alpha_2}{30\alpha_1+15\alpha_2}\left.\frac{\Gamma}{\mathcal{H}}\right|_{T_{\rm reh}}\right)\right. \\ &\left.-\xi\left(\frac{16\alpha_2+60\alpha_3}{30\alpha_1+15\alpha_2}\left.\frac{\Gamma}{\mathcal{H}}\right|_{T_{\rm reh}}\right)\right],\end{aligned} \tag{24}$$

where $\xi(y)$ is proportional to the incomplete Gamma function $\Gamma(5/3,0,y)$:

$$\xi(y) \equiv \frac{\Gamma(5/3,0,y)}{y^{5/3}} = \frac{5}{3} - \frac{3y}{8} - \frac{3y^2}{22} + \mathcal{O}(y^3)\,. \tag{25}$$

Notice that $\Delta_{\nu_R}(\infty)$ is well-behaved in the limit $50\alpha_1 + \alpha_2 - 60\alpha_3 \to 0$. At leading order in small α_i, *i.e.* for small couplings in Eqs. (10)–(12), the expression simplifies considerably:

$$\Delta_{\nu_R}(x\to\infty) \simeq 0.297\,\frac{M_{\rm Pl}(\alpha_1-\alpha_2)}{\sqrt{g_*}T_{\rm reh}}\frac{M_{\rm Pl}}{\sqrt{g_*}h_*T_{\rm reh}}\varepsilon \tag{26}$$

$$\sim \frac{9}{g_*h_*}\frac{M_{\rm Pl}^2}{T_{\rm reh}^2}\,\mathcal{O}\left(\frac{T_{\rm reh}^{10}}{M^{10}}F^{10}\right), \tag{27}$$

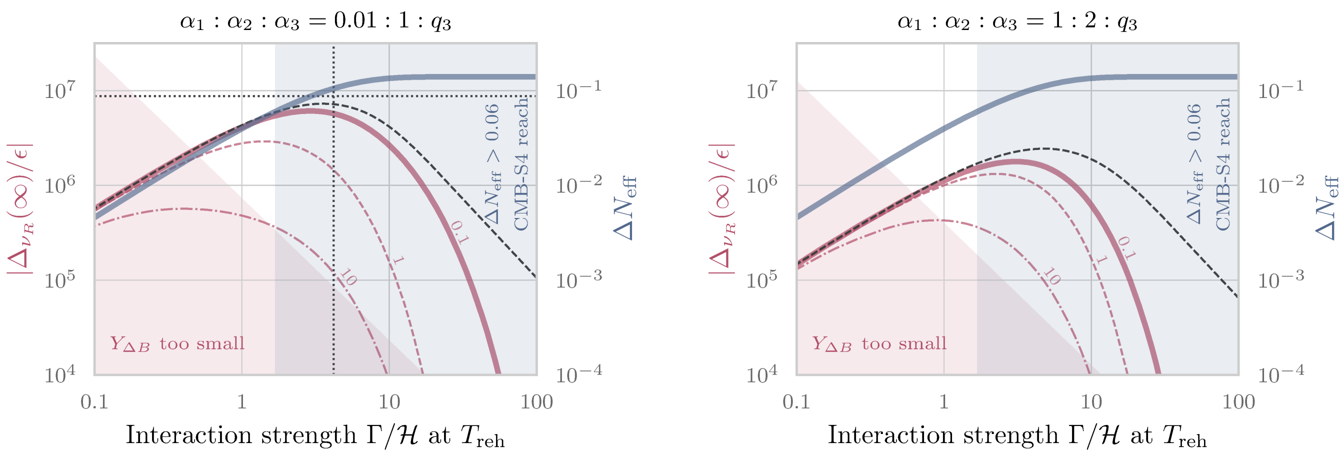
where in the second line we have assumed both X masses to be of similar order and all Yukawa couplings of order F. Despite the suppression by ten powers of the small quantities $T_{\rm reh}/M$ and F, the above can easily give the required $\frac{79}{28}\,0.9\times10^{-10}$ by choosing a low reheating temperature $T_{\rm reh} \ll M_{\rm Pl}$. A *lower* bound on $T_{\rm reh}$ of order 10^4 GeV can be obtained from the inequality (13), which can only be mildly relaxed in the freeze-out regime. For $T_{\rm reh}$ near this lower bound, we need $\mathcal{O}(M/F) \sim 3\times10^7$ GeV to explain the baryon asymmetry, so the X masses can easily be orders of magnitude above the reheating temperature even for perturbatively small couplings F. Notice that in this freeze-in regime, α_3 plays no role since the underlying reaction $e_R u_R \to LQ$ does not involve ν_R; numerically, we find that increasing α_3 for fixed ε always suppresses the final asymmetry.

Equation (26) shows an increase in ν_R asymmetry for increasing $\alpha_{1,2}$, typical for the freeze-in regime. Once the interactions become strong enough to thermalize the ν_R, Eq. (19), the asymmetry will of course start to exponentially decrease. Still in the limit of weak ν_R interactions, we can approximate the *maximal* ν_R asymmetry using Eq. (24) as

$$\Delta_{\nu_R}^{\rm max}(x\to\infty) \simeq 0.4\,\frac{M_{\rm Pl}}{\sqrt{g_*}h_*T_{\rm reh}}\varepsilon\,, \tag{28}$$

which can be achieved for

$$\alpha_{1,3} \ll \alpha_2 \simeq 3.8\,\frac{\sqrt{g_*}T_{\rm reh}}{M_{\rm Pl}} \tag{29}$$



(a) (b)

FIG. 3: Numerically computed values of $|\Delta_{\nu_R}(\infty)/\varepsilon|$ (red contours labeled by q_3) plotted with $\Delta N_{\rm eff}$ (grey line) for $T_{\rm reh} = 10^8$ GeV. The ratio of the wash-out cross sections was fixed as indicated in the upper part of each panel. The black dashed line shows the freeze-in approximation from Eq. (24). In Fig. 3a, the black dotted lines correspond to the maximal asymmetry from Eq. (28) and the respective α_2 value from Eq. (29). The red-shaded region is excluded by Eq. (13) for $q_3 = 0.1$; for higher q_3 values, this constraint will be slightly relaxed. The gray-shaded region corresponds to $\Delta N_{\rm eff} > 0.06$, which can be reached by CMB-S4 experiments.

or, almost, for values

$$\alpha_{2,3} \ll \alpha_1 \gg 0.8\, \frac{\sqrt{g_*} T_{\rm reh}}{M_{\rm Pl}}\,. \tag{30}$$

Even though the position and value of the $\Delta^{\rm max}_{\nu_R}(x \to \infty)$ lie strictly speaking outside of the freeze-in regime used in our analytical approximation, they serve as useful guides.

For fixed ratios of the three α_i, we show the final ν_R asymmetry as a function of the ν_R interaction rate Γ from Eq. (17) over Hubble at reheating in Fig. 3. We can clearly see the features derived above: for small interaction rates, the linear freeze-in approximation from Eq. (26) works very well, although deep into that regime the inequality from Eq. (13) makes it impossible to obtain the measured baryon asymmetry. At higher reheating temperatures, the constraint of Eq. (14) also becomes relevant. For $\Gamma/\mathcal{H}|_{T_{\rm reh}} \gg 1$ the asymmetry decreases exponentially, resulting in a maximum around $\Gamma/\mathcal{H}|_{T_{\rm reh}} \sim \mathcal{O}(1)$, depending on α_3.

For large ν_R interaction rates, $\Gamma/\mathcal{H}|_{T_{\rm reh}} \gg 1$, the ν_R density Σ_{ν_R} starts to saturate the equilibrium value. Since the ν_R are decoupled from the SM bath, these essentially massless fermions contribute to the radiation density in the early universe, usually parametrized via the effective number of neutrinos, $N_{\rm eff}$ (see *e.g.* Ref. [6]):

$$\begin{aligned}\Delta N_{\rm eff} &\simeq 0.14 \left(\frac{106.75}{g_*}\right)^{4/3} \frac{\Sigma_{\nu_R}(\infty)}{\Sigma^{\rm eq}_{\nu_R}} \\ &\simeq 0.14 \left(\frac{106.75}{g_*}\right)^{4/3} \left(1 - \exp\left[-\frac{\Gamma}{3\mathcal{H}}\bigg|_{T_{\rm reh}}\right]\right).\end{aligned} \tag{31}$$

We show $\Delta N_{\rm eff}$ together with the ν_R asymmetry in Fig. 3. Currently, even three fully thermalized ν_R that decouple above the electroweak scale are within the experimental uncertainty [12], but upcoming CMB-S4 results will be sensitive to $\Delta N_{\rm eff}$ down to 0.06 [27, 28] and hence probe the region

$$2\alpha_1 + \alpha_2 > \frac{T_{\rm reh}}{1.5 \times 10^{17}\,\text{GeV}}\,. \tag{32}$$

If a larger value than the SM value $N_{\rm eff} \simeq 3$ is found, it could be interpreted as the contribution of the Dirac partners ν_R. On the other hand, even SM-compatible observations of $N_{\rm eff}$ at CMB-S4 cannot exclude Dirac leptogenesis, since the freeze-in regime remains as a finite region of parameter space where the correct baryon asymmetry can be achieved without excessive $\Delta N_{\rm eff}$ (see Fig. 3).

Overall, Dirac leptogenesis via scattering is an efficient mechanism to explain the matter–antimatter asymmetry of our universe in models with Dirac neutrinos, without ever violating $B - L$. In the flavor-blind – effectively one-generational – case studied above, this extends the viable parameter space of all Dirac-leptogenesis models of Ref. [6] that contain at least three different Yukawa matrices, *i.e.* models *b* and *c*, into the region $T_{\rm reh} \ll M_X$. Within our approximations, the models with only two Yukawa matrices would yield a vanishing asymmetry [29]. However, the introduction of flavor effects could change this conclusion and allow for all Dirac leptogenesis models to work in the low-reheating regime. A detailed discussion of such flavor effects is left for future work.

VI. CONCLUSION

Neutrino oscillations have proven neutrinos to be massive particles, but without any indication of whether they are Dirac or Majorana particles. While we anxiously await an experimental resolution of this question through studies of neutrinoless double-beta decay, it behooves us to investigate possible impacts of *Dirac neutrinos*. Compared to Majorana neutrinos, the connection between neutrino mass and matter–antimatter asymmetry has hardly been explored, even though Dirac leptogenesis [5] beautifully employs the neutrino-mass smallness to hide a lepton asymmetry in the decoupled ν_R bath, letting sphalerons create a baryon asymmetry in a $B-L$-conserving universe. Similar to standard leptogenesis, this is easily accomplished through the out-of-equilibrium decays of new heavy particles [5, 6].

In this paper, we have demonstrated that the non-thermalization of the ν_R is sufficient to satisfy Sakharov's conditions *without* the need for heavy mediator decays. This makes it possible to obtain the observed baryon asymmetry purely from scatterings without ever producing the mediators. This is relevant for inflationary scenarios with small reheating temperatures compared to the mediator masses.

In the flavor-blind limit employed here, a non-vanishing asymmetry requires the existence of at least three different scattering states in the considered $2 \to 2$ processes, which we exemplified by considering heavy leptoquarks mediating the interactions $\nu_R d_R \leftrightarrow LQ$, $\nu_R d_R \leftrightarrow e_R u_R$, and $e_R u_R \leftrightarrow LQ$. While the source term for the Δ_{ν_R} vanishes in the Boltzmann equations, the existence of wash-in processes does enable the ν_R asymmetry generation.

We show that sufficiently large asymmetries can be achieved in a wide range of parameter space with a maximal value reached for intermediate interaction strength that lie between the freeze-in regime (very weak couplings) and the semi-thermalized regime (larger couplings). The viable window is limited towards smaller couplings as the ν_R abundance becomes too small to provide large Δ_{ν_R} and towards larger couplings by the diminishing deviation from thermal equilibrium of ν_R. The scenario can be probed by an enhanced N_{eff}: while evading current constraints, CMB-S4 experiments are expected to test the semi-thermalized regime as well as the maximum asymmetry scenario.

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