Mechanism for baryogenesis via feebly interacting massive particles

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We present a simple mechanism which allows the simultaneous generation of the baryon asymmetry of the Universe along with its dark matter content. To this goal, we employ the out-of-equilibrium decays of heavy bath states into a feebly coupled dark matter particle and Standard Model charged fermions. These decays lead to dark matter production via the freeze-in mechanism and, assuming that they further violate CP , can generate a viable matter-antimatter asymmetry in the resonant regime. We illustrate this mechanism by studying a particular realization of this general scenario, where the role of the heavy bath particles is played by $SU(3)_c \times SU(2)_L$ -singlet vectorlike fermions with a nonzero hypercharge and dark matter is identified with a gauge-singlet real scalar field. We show that in the context of this simple model the cosmological constraints for the dark matter abundance and the baryon asymmetry are satisfied for masses of heavy vectorlike fermion states of a few TeV, potentially within reach of the high-luminosity run of the Large Hadron Collider. Dark matter, in turn, is predicted to be rather light, with a mass of a few keV.

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I. INTRODUCTION

Among the numerous open questions in contemporary high-energy physics, the origin of cosmic dark matter (DM) and that of the baryon asymmetry in the Universe occupy a pivotal place. Not only do they constitute two of the most striking and fundamental pieces of evidence for the existence of physics beyond the Standard Model (BSM) of particle physics, but they do so only once the latter is placed within a cosmological framework. Resolving either (let alone both) of these questions most likely requires particle physics to be viewed from a cosmological standpoint and, conversely, cosmology to be analyzed in terms of the behavior of the fundamental constituents of matter and their interactions. And, indeed, in doing so during the past decades model builders have not been short of ideas concerning the nature of dark matter and the mechanism through which matter came to dominate over antimatter in the Universe.

Among the various dark matter candidates that have been proposed we can mention axions [1–3], primordial black holes [4–6], a vast number of incarnations of weakly or

feebly interacting massive particles (WIMPs or FIMPs; for reviews, cf., e.g., Refs. [7,8]), gravitationally produced dark matter [9], and asymmetric dark matter (for reviews, cf., e.g., Refs. [10,11]), just to name a few. Similarly, the matter-antimatter asymmetry of the Universe has been explained in terms of different mechanisms of baryogenesis, like baryogenesis based on grand unified theories (GUTs) baryogenesis [12], electroweak baryogenesis [13,14], the Affleck-Dine mechanism [15], and leptogenesis [16]. Interestingly, during the past decade, there have also been several attempts to actually link the two questions, most notably in the contexts of WIMP baryogenesis [17–19] or asymmetric dark matter [20–22].

Recently, the main idea behind Akhmedov-Rubakov-Smirnov (ARS) leptogenesis [23], namely, that of a lepton number asymmetry being generated through CP -violating sterile neutrino oscillations, was exploited in Ref. [24], where the authors proposed that a baryon asymmetry could, instead, be also generated by augmenting the SM with exotic scalars and fermions directly coupling to quarks. The fermions, which were taken to be singlets under the SM gauge group can, moreover, play the role of viable dark matter candidates through the freeze-in mechanism [25,26], whereas their CP -violating oscillations, in the presence of electroweak sphaleron transitions, can generate the observed baryon asymmetry of the Universe.

In this paper, we propose a mechanism which borrows ideas both from Dirac leptogenesis [27–30] and from this scenario of “freeze-in baryogenesis.” As we will describe in

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detail in the following, we consider a heavy particle species which is charged under (parts of) the SM gauge group and which can decay through feeble interactions into SM fermions along with a neutral particle. The latter is our dark matter candidate, produced upon the decays of the heavy particle through the freeze-in mechanism, along the lines presented in Ref. [24]. However, in our case, the decays *themselves* violate CP , in a similar manner as in leptogenesis models. Then, in the presence of electroweak sphalerons, we will see that an asymmetry can be generated between SM fermions and antifermions. The relevant processes proceed through feeble couplings, preventing them from ever reaching equilibrium and, thus, satisfying the third Sakharov condition [31]. The first condition, namely, baryon number violation, is satisfied due to the active sphaleron processes in a way that resembles neither GUT baryogenesis (explicit B violation in decays) nor leptogenesis (explicit L violation in decays), although we will also comment on the possibility of direct B/L violation as well.

The paper is structured as follows: In Sec. II, we discuss the general features of our mechanism, namely, the way through which the required dark matter abundance and the baryon and lepton asymmetries can be generated, without adopting any concrete microscopic model. In Sec. III, we propose a simple model as a proof of concept that concrete incarnations of our mechanism can, indeed, be constructed. We compute the predicted dark matter abundance and baryon asymmetry, quantify the effects of processes that wash out the latter, and briefly comment on the phenomenological perspectives of our model, notably in relation to searches for long-lived particles (LLPs) at the Large Hadron Collider (LHC). Finally, in Sec. IV, we summarize our main findings and conclude. Some more technical aspects are left for the Appendix.

II. DARK MATTER AND BARYOGENESIS FROM FREEZE-IN: GENERAL FRAMEWORK

Before presenting a concrete realization of our take on the freeze-in baryogenesis idea, let us briefly summarize a few key notions that will be useful for the discussion that follows: frozen-in DM and how the freeze-in framework can enable us to satisfy the three Sakharov conditions that are necessary for successful baryogenesis. A concrete implementation of these ideas will be presented in detail in Sec. III.

A. Freeze-in DM abundance

The freeze-in mechanism for dark matter production relies on two basic premises:

- (i) The initial DM abundance is zero.
- (ii) Dark matter interacts only extremely weakly (“feebly”) with the Standard Model particles (along with any other particles that are in thermal equilibrium with them).

Under these assumptions, and further assuming that dark matter production takes place in a radiation-dominated Universe, dark matter never reaches thermal equilibrium with the plasma. Instead, it is produced through the out-of-equilibrium decays or annihilations of bath particles and all dark matter depletion processes, the rate of which typically scales as $\langle\sigma v\rangle \times n_{\text{DM}}^2$ (where σ is the dark matter annihilation cross section, v its velocity, and n_{DM} its number density), can be ignored.

B. Baryon asymmetry abundance Y_B

In general, the decays and/or annihilations that are responsible for dark matter production can also violate both the baryon number B and C/CP .¹ Then, as long as we stick to the freeze-in framework, these processes occur out of equilibrium with the thermal plasma, thus fulfilling all three Sakharov conditions.

Intuitively, and ignoring all washout processes, if we denote the measure of CP violation by ϵ_{CP} , we would expect the generated asymmetry in the SM fermion $Y_{\Delta f}$ to be connected to the dark matter abundance Y_{DM} through a relation of the type

$$Y_{\Delta f} \sim \epsilon_{CP} Y_{\text{DM}}. \quad (1)$$

In reality, this limit cannot be attained given that some amount of washout is almost inevitable, whereas concrete realizations typically require the introduction of additional particles and decay channels. In this respect, this relation may be viewed as an upper limit to the asymmetry that can be generated through decays that simultaneously produce dark matter.

In fact, in the following, when studying a concrete incarnation of our decay-induced freeze-in baryogenesis idea, we will see the following.

- (i) Since we will be starting with non-self-conjugate initial states F_i (Dirac fermions), CPT conservation and unitarity impose the existence of multiple decay channels of the F_i ’s for a nonvanishing CP asymmetry to be generated. To this goal, we will exploit possible decays of the heavy fermions into different Standard Model fermions (leptons), i.e., flavor effects.
- (ii) The freeze-in framework will necessitate extremely small values for the couplings involved in the decay process. The predicted CP violation, being an effect that arises from the interference of tree-level and one-loop processes, is, then, even further suppressed, which will lead us to consider resonantly enhanced configurations in self-energy-type diagrams [32–35].

¹Very similar remarks can be made if the decay violates, instead, the lepton number. This is also the option that we will adopt later in this work.

Therefore, at least two heavy fermions F_i must be added.

- (iii) The baryon and/or lepton number need not be violated by the decay processes. As proposed, e.g., in Ref. [27], a CP asymmetry can be translated to a baryon asymmetry by the electroweak sphalerons. If, additionally, the BSM heavy fermions are immune to the action of sphalerons, in the end a net asymmetry will be generated in the baryon and lepton sectors.

III. A CONCRETE REALIZATION

Let us now elaborate the previous considerations through a concrete, simple model. We extend the SM by two heavy vectorlike leptons F_i which are singlets under $SU(3)_c \times SU(2)_L$ but carry hypercharge and a real gauge-singlet scalar S which is our freeze-in DM candidate. We moreover impose a discrete $\mathbb{Z}_2$ symmetry on the Lagrangian, under which all exotic states are taken to be odd while the SM particles are even. Under these assumptions, the Lagrangian reads

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_S + \mathcal{L}_{SF}, \quad (2)$$

where $\mathcal{L}_{\text{SM}}$ is the SM Lagrangian,

$$\mathcal{L}_S = \partial_\mu S \partial^\mu S - \frac{\mu_S^2}{2} S^2 + \frac{\lambda_S}{4!} S^4 + \lambda_{Sh} S^2 (H^\dagger H) \quad (3)$$

describes interactions of dark matter with itself and with the Standard Model Higgs doublet, and

$$\begin{aligned} \mathcal{L}_{SF} = & \sum_i (\bar{F}_i (i\mathcal{D}) F_i - M_i \bar{F}_i F_i) \\ & - \sum_{\alpha, i} (\lambda_{ai} S \bar{F}_i P_R e_\alpha + \lambda_{ai}^* S \bar{e}_\alpha P_L F_i), \end{aligned} \quad (4)$$

where e_α are the right-handed SM charged leptons of flavor $\alpha = \{e, \mu, \tau\}$ and λ_{ai} denote the feeble couplings. The heavy fermions F_i are assumed to carry the same lepton number with the SM leptons, so their interactions are lepton number conserving. Note that, without loss of generality, we have neglected potential off-diagonal couplings among the F_i 's. For simplicity in what follows, we will also set the Higgs portal coupling λ_{Sh} to zero.

The tree-level proper decay width $\Gamma_{F_i \rightarrow e_\alpha S}|_0$ in the limit $M_i \gg m_{e_\alpha} + m_S$ is given by

$$\Gamma_{F_i \rightarrow e_\alpha S}|_0 \simeq \frac{|\lambda_{ai}|^2}{16\pi g_{F_i}} M_i, \quad (5)$$

where $g_{F_i} = 2$ are the internal degrees of freedom of species F_i . The equilibrium decay rate density $\gamma_{F_i \rightarrow e_\alpha S}$ evaluated in the decaying particle rest frame is [26]

$$\begin{aligned} \gamma_{F_i \rightarrow e_\alpha S} & \equiv \int d\Pi_{F_i} d\Pi_{e_\alpha} d\Pi_S (2\pi)^4 \\ & \times \delta^{(4)}(p_{F_i} - p_{e_\alpha} - p_S) f_{F_i}^{\text{eq}} |\mathcal{M}|_{F_i \rightarrow e_\alpha S}^2 \\ & = \frac{g_{F_i}}{2\pi^2} M_i^2 \Gamma_{F_i \rightarrow e_\alpha S} T K_1\left(\frac{M_i}{T}\right), \end{aligned} \quad (6)$$

where $d\Pi_k = d^3\mathbf{p}_k / (2\pi)^3 2E_k$ is the elementary Lorentz invariant phase space volume of species k , $|\mathcal{M}|^2$ denotes the squared matrix element summed, but not averaged, over the internal degrees of freedom of the initial states, K_1 is the modified Bessel function of the second kind of the order of one, and the distribution function of F_i approximately follows the Maxwell-Boltzmann distribution: $f_{F_i}^{\text{eq}} = e^{-E_{F_i}/T}$. The corresponding thermally averaged equilibrium decay rate density $\langle \gamma_{F_i \rightarrow e_\alpha S} \rangle$ reads

$$\begin{aligned} \langle \gamma_{F_i \rightarrow e_\alpha S} \rangle & \equiv \frac{\gamma_{F_i \rightarrow e_\alpha S}}{n_{F_i}^{\text{eq}}} = \frac{K_1(M_i/T)}{K_2(M_i/T)} \Gamma_{F_i \rightarrow e_\alpha S} \\ & \simeq \begin{cases} \frac{M_i}{2T} \Gamma_{F_i \rightarrow e_\alpha S}, & T \gg M_i, \\ \Gamma_{F_i \rightarrow e_\alpha S}, & T \ll M_i, \end{cases} \end{aligned} \quad (7)$$

where $n_{F_i}^{\text{eq}}$ is the equilibrium number density of F_i assuming zero chemical potential:

$$n_{F_i}^{\text{eq}} \equiv g_{F_i} \int \frac{d^3\mathbf{p}}{(2\pi)^3} f_{F_i}^{\text{eq}} = \frac{g_{F_i}}{2\pi^2} M_i^2 T K_2\left(\frac{M_i}{T}\right). \quad (8)$$

The time dilation factor $K_1(M_i/T)/K_2(M_i/T)$ implies that decays are inhibited at temperatures higher than the decaying state mass M_i .

The dominant $2 \leftrightarrow 2$ scattering processes modifying the abundance of S are those which involve the $U(1)_Y$ gauge boson as an external state, i.e., $F_i B \leftrightarrow e_\alpha S$, $F_i S \leftrightarrow e_\alpha B$, and $F_i \bar{e}_\alpha \leftrightarrow S B$, which will be henceforth referred to as gauge scatterings. The corresponding matrix elements depend on the product of a feeble and a gauge coupling, whereas all other scattering processes involve higher powers of feeble couplings and are therefore subleading. All relevant Feynman diagrams are shown in Fig. 1.

The equilibrium interaction rate density of a generic scattering process $ab \rightarrow cd$ is [36]

$$\begin{aligned} \gamma_{ab \rightarrow cd} & \equiv \int d\Pi_a d\Pi_b d\Pi_c d\Pi_d (2\pi)^4 \delta^{(4)}(p_a + p_b - p_c - p_d) \\ & \times f_a^{\text{eq}} f_b^{\text{eq}} |\mathcal{M}|_{ab \rightarrow cd}^2 \\ & = \frac{T}{512\pi^6} \int_{\tilde{s}_{\min}}^\infty d\tilde{s} \frac{1}{\sqrt{\tilde{s}}} |\mathbf{p}_{\text{in}}| |\mathbf{p}_{\text{fin}}| K_1\left(\frac{\sqrt{\tilde{s}}}{T}\right) \\ & \times \int d\Omega |\mathcal{M}|_{ab \rightarrow cd}^2, \end{aligned} \quad (9)$$

FIG. 1. Feynman diagrams that contribute to the generation of the baryon asymmetry.

where $\mathbf{p}_{\text{in}}$, $\mathbf{p}_{\text{fin}}$, and $\sqrt{s}$ are the initial and final momenta and the energy in the center-of-momentum frame, respectively, with $\tilde{s}_{\text{min}} = \max\{(m_a + m_b)^2, (m_c + m_d)^2\}$. In Fig. 2, we present the decay and gauge scattering rate densities involving F_1 and the right-handed electron e_e as external states in terms of the dimensionless parameter $z \equiv M_1/T$. For the feeble coupling, we use the value $\lambda_{e1} = 2.145 \times 10^{-8}$ and work at lowest order in perturbation theory and in the limit $M_1 \gg m_{e_\alpha} + m_S$. In order to study their effect on the Boltzmann equations, it is convenient to normalize them to the Hubble parameter H and to the number density of photons $n_\gamma = 2\zeta(3)T^3/\pi^2$, where ζ is the Riemann zeta function [37]. The integrals appearing in the various scattering rate densities have been computed with the Cuba numerical library [38].

FIG. 2. Tree-level decay and gauge scattering rate densities involving F_1 and e_e as external states, normalized to the Hubble parameter H and the equilibrium number density of photons n_γ . The feeble coupling of the interactions is $\lambda_{e1} \simeq 2.15 \times 10^{-8}$, and the densities have been evaluated at tree level and in the limit $M_1 \gg m_{e_\alpha} + m_S$.

The s-channel resonance that appears in the scattering $F_i B \leftrightarrow e_\alpha S$ has been regularized by the finite decay width of F_i . Also, the IR-induced resonances due to the exchange of massless SM leptons that appear in the u channel of $F_i B \leftrightarrow e_\alpha S$ and in the t channel of $F_i \bar{e}_\alpha \leftrightarrow S B$ have both been regulated by the thermal mass of the right-handed SM leptons $m_e^2(T) = g_Y^2 T^2/8$ (see [39] and references therein). Note that, as thermal effects are irrelevant at the temperature of a few TeV that is of interest to us, we include the thermal mass only when it acts as a regulator of the IR divergences. For a few examples in which finite temperature effects can become important, cf., e.g., Refs. [40–42].

A. Out-of-equilibrium decays

As we already mentioned, one of the crucial ingredients of our setup is that all processes involving dark matter (and, for that matter, CP violation) must never attain chemical equilibrium. In order to fulfill this condition, an upper bound on the magnitude of the freeze-in couplings can be obtained by requiring the total thermally averaged decay rate to be smaller than the Hubble expansion rate at the characteristic temperature $T = M_i$. Because of the feeble nature of the λ_{ai} couplings, successful baryogenesis requires a resonant enhancement of the asymmetry, which, in turn, implies that the masses of the two heavy fermions have to be very close to each other. Then, for $M_1 \simeq M_2$ the out-of-equilibrium condition reads

$$\left(\sum_{\alpha,i} \langle \gamma_{F_i \rightarrow e_\alpha S} \rangle \lesssim H \right) |_{T=M_1}. \quad (10)$$

The Hubble parameter is given by

$$H(T) \simeq \frac{1.66 \sqrt{g_{*p}}}{M_{\text{Pl}}} T^2, \quad (11)$$

where $M_{\text{Pl}} \simeq 1.22 \times 10^{19}$ GeV is the Planck mass and $g_{*p} \simeq 106.75$ are the effective degrees of freedom related to the energy density. At leading order and in the limit $M_i \gg m_{e_a} + m_s$, the out-of-equilibrium condition (10) reads ($g_{F_i} = 2$)

$$\sum_{\alpha,i} \Gamma_{F_i \rightarrow e_a S} \lesssim 2H|_{T=M_1} \Rightarrow \sum_{\alpha,i} |\lambda_{ai}|^2 \lesssim 2.83 \times 10^{-13} \left(\frac{M_1}{\text{TeV}} \right). \quad (12)$$

Thus, for heavy leptons at the TeV range, the couplings have to be smaller than $\sim 10^{-7}$ for the decays to proceed out of equilibrium.

B. Freeze-in DM abundance

The freeze-in DM abundance Y_S that is produced from decays and $2 \rightarrow 2$ scattering processes in our model follows the Boltzmann equation

$$\begin{aligned} s \frac{dY_S}{dt} = & \sum_{\alpha,i} \{F_i \leftrightarrow e_a S\} + \sum_{\alpha,i} \{F_i B \leftrightarrow e_a S\} \\ & - \sum_{\alpha,i} \{F_i S \leftrightarrow e_a B\} + \sum_{\alpha,i} \{F_i \bar{e}_a \leftrightarrow SB\} \\ & + 2 \sum_{i,j} \{F_i \bar{F}_j \leftrightarrow SS\} + 2 \sum_{\alpha,\beta} \{\bar{e}_\alpha e_\beta \leftrightarrow SS\}. \end{aligned} \quad (13)$$

In writing Eq. (13), we have used the notations

$$\{ab \leftrightarrow cd\} \equiv (ab \leftrightarrow cd) + (\bar{a} \bar{b} \leftrightarrow \bar{c} \bar{d}), \quad (14a)$$

$$[ab \leftrightarrow cd] \equiv (ab \leftrightarrow cd) - (\bar{a} \bar{b} \leftrightarrow \bar{c} \bar{d}), \quad (14b)$$

$$\begin{aligned} (ab \leftrightarrow cd) \equiv & \int d\Pi_a d\Pi_b d\Pi_c d\Pi_d (2\pi)^4 \delta^{(4)} [|\mathcal{M}|_{ab \rightarrow cd}^2 \\ & \times f_a f_b (1 \pm f_c)(1 \pm f_d) \\ & - |\mathcal{M}|_{cd \rightarrow ab}^2 f_c f_d (1 \pm f_a)(1 \pm f_b)], \end{aligned} \quad (14c)$$

where $\delta^{(4)}$ is an abbreviation for $\delta^{(4)}(p_a + p_b - p_c - p_d)$, f_k is the distribution function of species k , and s denotes the entropy density. We will make the following assumptions.

- (i) The initial DM abundance is zero. Combined with the feeble couplings, this allows us to ignore the inverse decays $e_a S \rightarrow F_i$, i.e., $f_S \simeq 0$.
- (ii) The DM production takes place during the radiation-dominated era. At this epoch, time and temperature are related by $\dot{T} \simeq -HT$, which is valid for $\partial g_{*p}/\partial T \simeq 0$. Using this relation, we may switch variables and write

$$s \frac{d}{dt} Y_S = s H z \frac{d}{dz} Y_S. \quad (15)$$

The entropy density during the radiation-dominated era is given by

$$s = \frac{2\pi^2}{45} g_{*s} T^3, \quad (16)$$

where g_{*s} are the effective degrees of freedom with respect to the entropy density.

- (iii) The distribution functions of the visible sector species obey the Maxwell-Boltzmann statistics; i.e., we neglect Bose-enhancement and Pauli-blocking factors. Hence, we can write, in general,

$$(ab \leftrightarrow cd) = \gamma_{ab \rightarrow cd} \frac{Y_a Y_b}{Y_a^{\text{eq}} Y_b^{\text{eq}}} - \gamma_{cd \rightarrow ab} \frac{Y_c Y_d}{Y_c^{\text{eq}} Y_d^{\text{eq}}}. \quad (17)$$

Let us first focus on the heavy lepton decays $F_i \rightarrow e_a S$ and the corresponding CP -conjugate processes, which provide the dominant contribution to DM production for $z \gtrsim 1$ [26]. Under the aforementioned assumptions, the Boltzmann equation for the DM abundance can be written as

$$\begin{aligned} s H z \frac{d}{dz} Y_S & \simeq \sum_{\alpha,i} \{F_i \leftrightarrow e_a S\} \\ & = 2 \sum_{\alpha,i} \gamma_{e_a S}^{F_i} \frac{Y_{F_i + \bar{F}_i}}{Y_{F_i + \bar{F}_i}^{\text{eq}}} + \mathcal{O}(\epsilon^2) \\ & = \frac{g_{F_1}}{\pi^2} M_1^2 T K_1 \left(\frac{M_1}{T} \right) \sum_{\alpha} \Gamma_{F_1 \rightarrow e_a S} \\ & \quad + \frac{g_{F_2}}{\pi^2} M_2^2 T K_1 \left(\frac{M_2}{T} \right) \sum_{\alpha} \Gamma_{F_2 \rightarrow e_a S}, \end{aligned} \quad (18)$$

where $\gamma_{e_a S}^{F_i}$ is the equilibrium decay rate density at tree level, ϵ denotes the CP asymmetry, and we have used $Y_{F_i + \bar{F}_i} \simeq Y_{F_i + \bar{F}_i}^{\text{eq}}$.² If we consider $M_1 \simeq M_2$ (resonant case) and substitute the decay width (5), the Hubble parameter (11), and the entropy density (16) in Eq. (18), then at tree level and in the limit $M_i \gg m_{e_a} + m_s$ the DM abundance simplifies to

$$Y_S(z) = \frac{45 M_{\text{Pl}}}{32 \times 1.66 \pi^5 g_{*s} \sqrt{g_{*p}}} \frac{\sum_{\alpha,i} |\lambda_{ai}|^2}{M_1} \int_{z_{RH}}^z dz' z'^3 K_1(z'), \quad (19)$$

where $z_{RH} \equiv M_1/T_{RH}$ and we have considered for simplicity that $\partial g_{*s}/\partial T \simeq 0$. In our analysis, T_{RH} will be set to 10^{12} GeV (although the precise value of T_{RH} does not affect our results, as long as $T_{RH} > M_i$). At the present day $T = T_0 \simeq 2.73$ K, so $T_0 \ll M_1 \ll T_{RH}$ and the integral contributes a factor of $3\pi/2$, yielding

²Gauge scatterings keep $F_i(\bar{F}_i)$ close to thermal equilibrium down to $z \sim 25$, when they eventually freeze out. Since the baryon asymmetry is generated prior to their freeze-out (when sphalerons become inactive), this is a valid approximation.

$$Y_S(z_0) = \frac{135M_{\text{Pl}}}{64 \times 1.66\pi^4 g_{*s} \sqrt{g_{*p}}} \frac{\sum_{a,i} |\lambda_{ai}|^2}{M_1}. \quad (20)$$

The experimentally observed DM abundance is

$$Y_S(z_0) = \frac{\Omega_{\text{DM}} \rho_c}{m_S s_0}, \quad (21)$$

where $\Omega_{\text{DM}} h^2 = 0.1200 \pm 0.0012$, $\rho_c \equiv 3H_0^2/8\pi G \simeq 10.537h^2 \text{ GeV } m^{-3}$ is the critical density, and $s_0 \simeq 2.9 \times 10^9 m^{-3}$ is the entropy density at the present day [43]. If we ignore dark matter production through scattering processes, the DM mass m_S can be related to the heavy lepton mass M_1 and to the feeble couplings as

$$m_S = \frac{64 \times 1.66\pi^4 g_{*s} \sqrt{g_{*p}} \Omega_{\text{DM}} \rho_c}{135M_{\text{Pl}} s_0} \frac{M_1}{\sum_{a,i} |\lambda_{ai}|^2}. \quad (22)$$

Lyman- α forest observations can be used in order to extract a lower bound on the DM mass, the exact value of which depends on the underlying mechanism of DM genesis. For DM candidates that freeze out, the current bound is $m_{\text{DM}} \gtrsim (1.9\text{--}5.3) \text{ keV}$ at 95% C.L. [44–46]. This limit can be mapped onto the case of freeze-in-produced DM yielding $m_S \gtrsim (4\text{--}16) \text{ keV}$ [47–50]. This, in turn, imposes an upper limit on the feeble couplings:

$$\sum_{a,i} |\lambda_{ai}|^2 \lesssim 7.55 \times 10^{-16} \left(\frac{M_1}{\text{TeV}} \right) \quad (23)$$

in order not to overclose the Universe, where we have used the more conservative bound $m_S \gtrsim 4 \text{ keV}$. Thus, we see that the Lyman- α forest sets a more severe constraint than the out-of-equilibrium condition of Eq. (10), forcing the feeble couplings to lie in the 10^{-8} range and below for heavy lepton masses at the TeV scale.

For a more rigorous treatment, one should take into account the impact of scattering processes, which can modify the predicted DM abundance at $z \lesssim 1$, an epoch when scatterings are dominant. Including such scattering processes will be essential for the calculation of the baryon asymmetry presented in the next section. We focus only on scattering processes which involve gauge bosons as external states and ignore the subleading ones [third row in Eq. (13)]. In this case, the Boltzmann equation takes the form

$$\begin{aligned} sH z \frac{dY_S}{dz} &= \sum_{a,i} \{F_i \leftrightarrow e_a S\} + \sum_{a,i} \{F_i B \leftrightarrow e_a S\} \\ &\quad - \sum_{a,i} \{F_i S \leftrightarrow e_a B\} + \sum_{a,i} \{F_i \bar{e}_a \leftrightarrow SB\} \\ &= 2 \sum_{a,i} (\gamma_{e_a S}^{F_i} + \gamma_{e_a S}^{F_i B} + \gamma_{SB}^{F_i \bar{e}_a} + \gamma_{e_a B}^{F_i S}) + \mathcal{O}(\epsilon^2), \end{aligned} \quad (24)$$

FIG. 3. DM abundance generated solely by decays, as well as by decays and scatterings for heavy lepton masses $M_1 \simeq M_2 = 1.2 \text{ TeV}$ and couplings $\sum_{a,i} |\lambda_{ai}|^2 \simeq 6.25 \times 10^{-16}$.

where we have made use of the freeze-in approximation $f_S \simeq 0$ and considered $Y_{F_i + \bar{F}_i} \simeq Y_{F_i + \bar{F}_i}^{\text{eq}}$. In Fig. 3, we fix the masses of the heavy leptons $M_1 \simeq M_2 = 1.2 \text{ TeV}$ and the couplings $\sum_{a,i} |\lambda_{ai}|^2 \simeq 6.25 \times 10^{-16}$, and we compare the predicted DM abundance as estimated if we take into account only decays of the heavy leptons (lower, green line) with the result obtained through a full-blown numerical solution of Eq. (24) (upper, yellow line). We observe that, as expected, the inclusion of the scattering processes does not drastically modify the predicted amount of dark matter in the Universe. Note that we have also cross-checked all of our results by implementing our model in FeynRules [51] and computing the predicted freeze-in DM abundance with CalcHEP and micrOMEGAs 5 [52,53]. Given our findings, we conclude that the analytic estimate of the DM mass given by Eq. (22) constitutes a reliable approximation. For the values of the physical parameters used in Fig. 3, the value of the dark matter mass in order to reproduce the observed DM abundance in the Universe is $m_S \simeq 5.81 \text{ keV}$, which is compatible with the Lyman- α forest bounds discussed previously.

C. CP asymmetry

The CP asymmetry generated through the decays of $F_{1,2}$ arises, at lowest order, due to the interference of the tree-level and one-loop self-energy Feynman diagrams (wave-part contribution) as shown in Figs. 1(a) and 1(b). It can be defined in terms of the decay widths as

$$\begin{aligned} \epsilon_{ai} &\equiv \frac{\Gamma(F_i \rightarrow e_a S) - \Gamma(\bar{F}_i \rightarrow \bar{e}_a S)}{\sum_a \Gamma(F_i \rightarrow e_a S) + \Gamma(\bar{F}_i \rightarrow \bar{e}_a S)} \\ &= \frac{\Gamma(F_i \rightarrow e_a S) - \Gamma(\bar{F}_i \rightarrow \bar{e}_a S)}{2\Gamma_i}, \end{aligned} \quad (25)$$

where Γ_i is the total decay width of F_i . If we separate the matrix element $\mathcal{M}|_k$, with k being the loop order

($k = 0, 1, \dots$), into a coupling constant part c_k (denoting a collection of coupling constants) and an amplitude part $\mathcal{A}_k$, i.e., $\mathcal{M}|_k \equiv \sum_{i=0}^k c_i \mathcal{A}_i$, we can rewrite the CP asymmetry as [54]

$$\epsilon_{ai} = -2 \frac{\text{Im}\{c_0 c_1^*\} \int d\Pi_{e_a} d\Pi_S (2\pi)^4 \delta^{(4)} \text{Im}\{\mathcal{A}_0 \mathcal{A}_1^*\}}{\sum_a |c_0|^2 \int d\Pi_{e_a} d\Pi_S (2\pi)^4 \delta^{(4)} |\mathcal{A}_0|^2}. \quad (26)$$

The imaginary part of the amplitudes is related to the discontinuity of the corresponding graph as $2i\text{Im}\{\mathcal{A}_0 \mathcal{A}_1^*\} = \text{Disc}\{\mathcal{A}_0 \mathcal{A}_1^*\}$, which can be computed using the Cutkosky cutting rules. For a nondegenerate F_i spectrum, $M_j - M_i \gg \Gamma_j$, we obtain [55]

$$\epsilon_{ai} = -\frac{1}{16\pi} \frac{1}{1-x_j} \frac{\text{Im}\{\lambda_{ai}^* \lambda_{aj} [\lambda^\dagger \lambda]_{ji}\}}{[\lambda^\dagger \lambda]_{ii}}, \quad (27)$$

where $x_j \equiv M_j^2/M_i^2$ and we have used the standard notation $[\lambda^\dagger \lambda]_{ii} \equiv |\lambda_{ei}|^2 + |\lambda_{\mu i}|^2 + |\lambda_{\tau i}|^2$ and $[\lambda^\dagger \lambda]_{ji} \equiv \lambda_{ej}^* \lambda_{ei} + \lambda_{\mu j}^* \lambda_{\mu i} + \lambda_{\tau j}^* \lambda_{\tau i}$. This is half the result obtained in standard leptogenesis [see, e.g., Eq. (5.13) in Ref. [54]], because we deal with Dirac fermions where only a charged lepton propagates in the loop, whereas Majorana neutrinos can decay to both a charged lepton or a light neutrino accompanied by the Higgs field.

From CPT invariance and unitarity, we know that the total decay width of a state and its CP conjugate are equal:

$$\sum_a \Gamma(F_i \rightarrow e_a S) = \sum_a \Gamma(\bar{F}_i \rightarrow \bar{e}_a S) \equiv \Gamma_i. \quad (28)$$

Hence, the CP asymmetry vanishes when summed over the flavors: $\epsilon_i \equiv \sum_a \epsilon_{ai} = 0$. This can be seen by summing Eq. (27) over flavors, a case in which the argument of the imaginary part becomes real and it vanishes identically:

$$\sum_a \text{Im}\{\lambda_{ai}^* \lambda_{aj} [\lambda^\dagger \lambda]_{ji}\} = \text{Im}\{[\lambda^\dagger \lambda]_{ij} [\lambda^\dagger \lambda]_{ji}\} = 0. \quad (29)$$

However, flavor effects *can* lead to a nonvanishing baryon or lepton asymmetry, since washout processes are flavor dependent and, therefore, the lepton asymmetries in each flavor are washed out in a different way [56–59].

Because of the feeble nature of the λ_{ai} couplings, a resonant enhancement of the CP asymmetry is needed in order to generate a sufficiently large baryon asymmetry. This occurs when the mass difference between the heavy leptons is of the order of their decay widths and is related to the wave-part contribution to the CP asymmetry [Fig. 1(b)]. The resulting CP asymmetry is given by

$$\epsilon_{ai} = -\frac{1}{16\pi} \frac{1-x_j}{(1-x_j)^2 + g_j^2} \frac{\text{Im}\{\lambda_{ai}^* \lambda_{aj} [\lambda^\dagger \lambda]_{ji}\}}{[\lambda^\dagger \lambda]_{ii}}, \quad (30)$$

where $g_j \equiv \Gamma_j/M_i$. In deriving Eq. (30), we have regularized the divergence for small x_j by applying the resummation procedure presented in Ref. [34], with the regulator given by $M_1 \Gamma_2$. Note that a more complete treatment of the resonance enhancement requires the employment of non-equilibrium techniques, which yield, in general, different regulators. Such an analysis has been performed, e.g., in Ref. [60], where the nearly degenerate states are considered to be out of equilibrium with the bath, in contrast to our model, where the F_i 's are kept close to equilibrium and the required deviation occurs for the DM states. Such an analysis is beyond the scope of this paper and is left for future work.

If we also express the coupling constants in terms of their magnitude and phase, $\lambda_{ai} = |\lambda_{ai}| e^{i\phi_{ai}}$, Eq. (30) can be rewritten as

$$\begin{aligned} \epsilon_{ai} &= -\frac{1}{16\pi} \frac{1-x_j}{(1-x_j)^2 + g_j^2} \frac{|\lambda_{ai}| |\lambda_{aj}|}{[\lambda^\dagger \lambda]_{ii}} \\ &\quad \times \sum_{\beta \neq \alpha} |\lambda_{\beta i}| |\lambda_{\beta j}| \sin(-\phi_{ai} + \phi_{aj} - \phi_{\beta j} + \phi_{\beta i}) \\ &\equiv -\frac{1}{16\pi} \frac{1-x_j}{(1-x_j)^2 + g_j^2} \frac{|\lambda_{ai}| |\lambda_{aj}|}{[\lambda^\dagger \lambda]_{ii}} \sum_{\beta \neq \alpha} |\lambda_{\beta i}| |\lambda_{\beta j}| p_{\alpha\beta}^{ij}, \end{aligned} \quad (31)$$

where $p_{\alpha\beta}^{ij} = -p_{\alpha\beta}^{ji} = -p_{\beta\alpha}^{ij}$. The resonance condition reads $M_2 - M_1 \sim \Gamma_2/2$ [34], and in this case the CP asymmetry can be maximally enhanced to

$$\epsilon_{a1}^{\text{res}} = \frac{g_{F_1}}{2} \frac{|\lambda_{a1}| |\lambda_{a2}|}{[\lambda^\dagger \lambda]_{11} [\lambda^\dagger \lambda]_{22}} \sum_{\beta \neq \alpha} |\lambda_{\beta 1}| |\lambda_{\beta 2}| p_{\alpha\beta}^{12}, \quad (32a)$$

$$\epsilon_{a2}^{\text{res}} = g_{F_2} \frac{|\lambda_{a1}| |\lambda_{a2}|}{[\lambda^\dagger \lambda]_{11}^2 + [\lambda^\dagger \lambda]_{22}^2} \sum_{\beta \neq \alpha} |\lambda_{\beta 1}| |\lambda_{\beta 2}| p_{\alpha\beta}^{12}. \quad (32b)$$

These are the expressions that we will be using in the numerical analysis that follows in order to compute ϵ_{ai} .

D. Baryon asymmetry

In the previous sections, we described how our model can satisfy two out of the three Sakharov conditions, namely, C/CP violation and out-of-equilibrium dynamics. The last condition to be fulfilled in order to generate a baryon asymmetry is the violation of the baryon or lepton number. In the case of Majorana heavy states, a conserved lepton number cannot be consistently assigned in the presence of interaction and mass terms in the Lagrangian, and, therefore, L is violated. This is not the case in our model, where the heavy states are of Dirac nature. In this case, we may rely on sphaleron departure from equilibrium during the epoch that the lepton asymmetry is generated, along the lines described in Ref. [30].

In the model we propose, the heavy leptons F_i carry the same lepton number as the SM leptons, and the total lepton asymmetry is $Y_L = Y_{L_{\text{SM}}} + Y_{L_F}$, where $Y_{L_{\text{SM}}} \equiv \sum_{\alpha} Y_{L_{\alpha}}$ and $Y_{L_F} = \sum_i Y_{L_{F_i}}$. All processes conserve the combination $Y_{B-L} \equiv Y_B - Y_{L_{\text{SM}}} - Y_{L_F}$, i.e., $dY_{B-L}/dz = 0$. We also assume that the Universe is initially totally symmetric, $Y_{B-L_{\text{SM}}}|_{z_{RH}} = Y_{L_F}|_{z_{RH}} = 0$, and, therefore, at any z it holds $Y_{B-L_{\text{SM}}} = Y_{L_F}$.

Since sphalerons are insensitive to the lepton asymmetry Y_{L_F} , as F_i are $SU(3)_c \times SU(2)_L$ singlets, they affect only the nonzero lepton asymmetry stored in the SM lepton sector $Y_{L_{\text{SM}}}$. In particular, they convert it to a baryon

asymmetry by imposing certain relations among the chemical potentials of the various species (see the Appendix). Once sphalerons depart from equilibrium, which occurs at $T_{sph} = 131.7 \pm 2.3$ GeV [61], the baryon and lepton numbers are separately conserved. When the heavy leptons decay away, the total baryon asymmetry, being proportional to $Y_{B-L_{\text{SM}}}$, vanishes. However, if sphalerons become inactive during the decay epoch of the heavy leptons, then the baryon asymmetry freezes at $Y_B \propto Y_{B-L_{\text{SM}}}|_{T_{sph}}$, which, in general, is not null.

Taking into consideration all the decay and $2 \rightarrow 2$ scattering processes, the full Boltzmann equations of the asymmetries read

$$\begin{aligned} -sHz \frac{dY_{\Delta F_i}}{dz} = & \sum_{\alpha} [F_i \leftrightarrow e_{\alpha} S] + \sum_{\alpha} [F_i B \leftrightarrow e_{\alpha} S] + \sum_{\alpha} [F_i S \leftrightarrow e_{\alpha} B] + \sum_{\alpha} [F_i \bar{e}_{\alpha} \leftrightarrow SB] \\ & + \sum_{\alpha, \beta, j} [F_i \bar{e}_{\alpha} \leftrightarrow \bar{F}_j e_{\beta}] + \sum_{\alpha, \beta} [F_i \bar{e}_{\alpha} \leftrightarrow F_j \bar{e}_{\beta}] + \sum_{\alpha, \beta} [F_i e_{\beta} \leftrightarrow F_j e_{\alpha}] \\ & + \sum_{\alpha, \beta} [F_i \bar{F}_{j \neq i} \leftrightarrow e_{\alpha} \bar{e}_{\beta}] + [F_i \bar{F}_{j \neq i} \leftrightarrow SS] + \sum_{\alpha, \beta, j} [F_i F_j \leftrightarrow e_{\alpha} e_{\beta}] + [F_i S \leftrightarrow F_{j \neq i} S], \end{aligned} \quad (33)$$

$$\begin{aligned} -sHz \frac{dY_{\Delta_{\alpha}}}{dz} = & \sum_i [F_i \leftrightarrow e_{\alpha} S] + \sum_i [F_i B \leftrightarrow e_{\alpha} S] + \sum_i [F_i S \leftrightarrow e_{\alpha} B] + \sum_i [F_i \bar{e}_{\alpha} \leftrightarrow SB] \\ & + \sum_{i, j, \beta} [F_i \bar{e}_{\alpha} \leftrightarrow \bar{F}_j e_{\beta}] + \sum_{i, j, \beta \neq \alpha} [F_i \bar{e}_{\alpha} \leftrightarrow F_j \bar{e}_{\beta}] + \sum_{i, j, \beta \neq \alpha} [F_i e_{\beta} \leftrightarrow F_j e_{\alpha}] \\ & + \sum_{i, j, \beta \neq \alpha} [F_i \bar{F}_j \leftrightarrow e_{\alpha} \bar{e}_{\beta}] + \sum_{\beta \neq \alpha} [\bar{e}_{\alpha} e_{\beta} \leftrightarrow SS] + \sum_{i, j, \beta} [F_i F_j \leftrightarrow e_{\alpha} e_{\beta}] + \sum_{\beta \neq \alpha} [e_{\beta} S \leftrightarrow e_{\alpha} S]', \end{aligned} \quad (34)$$

where $Y_{\Delta F_i} \equiv Y_{F_i} - Y_{\bar{F}_i}$ and $Y_{\Delta_{\alpha}} \equiv Y_B/3 - Y_{L_{\text{SM}\alpha}}$. The primed term indicates that one has to subtract the contribution due to the on-shell propagation of F_i , usually referred to as real intermediate state subtraction (RISS), which is already taken into account by the successive decays $e_{\alpha} S \leftrightarrow F_i \leftrightarrow e_{\alpha} S$.

The various terms in the Boltzmann equations can be expressed in terms of the CP asymmetry ϵ_{ai} , the tree-level rate densities, and the asymmetric abundances. As is typically done, we linearize in the SM chemical potentials [62]

$$\frac{Y_{e_{\alpha}(\bar{e}_{\alpha})}}{Y_{e_{\alpha}}^{\text{eq}}} \equiv e^{\pm \mu_{e_{\alpha}}/T} \simeq 1 \pm \frac{\mu_{e_{\alpha}}}{T} = 1 \pm \frac{Y_{\Delta e_{\alpha}}}{2Y_{\gamma}}. \quad (35)$$

All nongauge interactions are subleading in comparison to the gauge interactions and can be safely discarded. We include only the CP -violating part of the RISS term, which ensures that the source term (proportional to ϵ_{ai}) takes the correct form; i.e., it vanishes when all species are in chemical equilibrium. The Boltzmann equations can be rewritten as³

$$-sHz \frac{dY_{\Delta F_i}}{dz} = \sum_{\alpha} \left(y_{F_i} - \frac{Y_S}{Y_S^{\text{eq}}} y_{e_{\alpha}} \right) (\gamma_{e_{\alpha} S}^{F_i} + \gamma_{e_{\alpha} S}^{F_i B}) + \sum_{\alpha} \left(\frac{Y_S}{Y_S^{\text{eq}}} y_{F_i} - y_{e_{\alpha}} \right) \gamma_{e_{\alpha} B}^{F_i S} + \sum_{\alpha} (y_{F_i} - y_{e_{\alpha}}) \gamma_{SB}^{F_i \bar{e}_{\alpha}}, \quad (36)$$

³In deriving Eq. (24), we dropped dark matter annihilation processes altogether. In baryo- and leptogenesis, inverse reactions can be crucial and need to be taken into account. In order to do so in an efficient manner, in Eqs. (36) and (37) we will approximate $f_S \simeq f_S^{\text{eq}} \frac{Y_S}{Y_S^{\text{eq}}}$ —a relation which is rigorously applicable to bath particles. We have checked that—as expected, since we have restricted ourselves to regions of the parameter space in which the condition of Eq. (10) is satisfied—this is, indeed, a good approximation which leads to only a small ($\sim 2\%$) reduction of the predicted dark matter abundance for large values of z .

$$\begin{aligned}
-sHz \frac{dY_{\Delta\alpha}}{dz} = & 2 \sum_i \epsilon_{ai} \left[\left(1 - \frac{Y_S}{Y_S^{\text{eq}}} \right) \sum_{\rho} (\gamma_{e_{\rho}S}^{F_i} + \gamma_{e_{\rho}S}^{F_i B} + \gamma_{SB}^{F_i \bar{e}_{\rho}} - \gamma_{e_{\rho}B}^{F_i S}) \right] \\
& + \sum_i \left[y_{F_i} - \frac{Y_S}{Y_S^{\text{eq}}} \sum_{\beta} (B_{e_{\beta}S}^{F_i} y_{e_{\beta}}) \right] \gamma_{e_{\alpha}S}^{F_i} + \sum_i \left(y_{F_i} - \frac{Y_S}{Y_S^{\text{eq}}} y_{e_{\alpha}} \right) \gamma_{e_{\alpha}S}^{F_i B} \\
& + \sum_i \left(\frac{Y_S}{Y_S^{\text{eq}}} y_{F_i} - y_{e_{\alpha}} \right) \gamma_{e_{\alpha}B}^{F_i S} + \sum_i (y_{F_i} - y_{e_{\alpha}}) \gamma_{SB}^{F_i \bar{e}_{\alpha}}, \quad (37)
\end{aligned}$$

where $y_{F_i} \equiv Y_{\Delta F_i}/Y_{F_i}^{\text{eq}}$, $y_{e_{\alpha}} \equiv Y_{\Delta e_{\alpha}}/Y_{\gamma}$, $B_{e_{\beta}S}^{F_i}$ denotes the branching ratio of the decay $F_i \rightarrow e_{\beta}S$, and we have used $Y_{F_i+\bar{F}_i} \simeq Y_{F_i+\bar{F}_i}^{\text{eq}}$. In deriving the equations above, we have taken into account that the CP asymmetry in scattering processes stemming from self-energy diagrams [Fig. 1(b)] is always equal to the CP asymmetry from decays [63]. As described in Ref. [54], one must also include the contributions from the RISS $2 \rightarrow 3$ scatterings involving gauge bosons, in order to obtain the correct form for the scattering source terms. These processes can, however, be neglected as subleading with regards to their impact on washout. Note that the source term in the Boltzmann equation of the asymmetry in F_i vanishes as a consequence of CPT and unitarity, as explained in Sec. III C. Lastly, these equations are not totally independent, as they are related through the conservation of Y_{B-L} , resulting in $\sum_i Y_{\Delta F_i} = \sum_{\alpha} Y_{\Delta\alpha}$.

One can express the asymmetries in each SM flavor $y_{e_{\alpha}}$ in terms of $Y_{\Delta\alpha}$, solving a system of algebraic equations which relate the chemical potentials and abundances in equilibrium of the various species (see the Appendix). For the temperatures of interest we find

$$\begin{pmatrix} y_{e_e} \\ y_{e_{\mu}} \\ y_{e_{\tau}} \end{pmatrix} = -\frac{1}{711} \begin{bmatrix} 230 & -7 & -7 \\ -7 & 230 & -7 \\ -7 & -7 & 230 \end{bmatrix} \begin{pmatrix} Y_{\Delta e} \\ Y_{\Delta\mu} \\ Y_{\Delta\tau} \end{pmatrix} \frac{1}{Y_{\gamma}}. \quad (38)$$

In the same way, we have calculated the amount of the SM lepton asymmetry converted to baryon asymmetry by the sphaleron transitions to be

$$Y_B = \frac{22}{79} \sum_{\alpha} Y_{\Delta\alpha}. \quad (39)$$

The system of the five coupled Boltzmann equations of the asymmetries (36) and (37) has been solved numerically, using the DM abundance generated by decays and scattering processes of Eq. (24). We confirm that the proposed model can indeed generate the observed baryon asymmetry $Y_B = 8.71 \pm 0.06 \times 10^{-11}$ [43] in the resonant regime. As an illustration, we present in Fig. 4 an explicit example, using the following values of the parameters:

$$\begin{aligned}
M_1 &= 1.2 \text{ TeV}, & M_2 &= M_1 + \Gamma_2/2, \\
|\lambda_{e1}| &= 2.145 \times 10^{-8}, & |\lambda_{e2}| &= |\lambda_{\tau 2}| = 9 \times 10^{-10}, \\
|\lambda_{\mu 1}| &= |\lambda_{\tau 1}| = 9 \times 10^{-9}, & |\lambda_{\mu 2}| &= 8 \times 10^{-10}, \\
p_{e\mu}^{12} &= p_{e\tau}^{12} = p_{\mu\tau}^{12} = 1. \quad (40)
\end{aligned}$$

The corresponding resonant CP asymmetries ϵ_{ai} turn out to be $\epsilon_{e1} \simeq 2.1 \times 10^{-1}$, $\epsilon_{\mu 1} \simeq -5.7 \times 10^{-2}$, $\epsilon_{\tau 1} \simeq -1.53 \times 10^{-1}$ and $\epsilon_{e2} \simeq 1.53 \times 10^{-3}$, $\epsilon_{\mu 2} \simeq -4.2 \times 10^{-4}$, $\epsilon_{\tau 2} \simeq -1.11 \times 10^{-3}$. We have explicitly verified that Y_{B-L} is conserved for all values of z or, equivalently, that the relation $\sum_i Y_{\Delta F_i} = \sum_{\alpha} Y_{\Delta\alpha}$ holds. Similarly, the baryon asymmetry vanishes in the absence of flavor effects, i.e., when $|\lambda_{e1}| = |\lambda_{\mu 1}| = |\lambda_{\tau 1}|$ and $|\lambda_{e2}| = |\lambda_{\mu 2}| = |\lambda_{\tau 2}|$.

Note that, contrary to the standard leptogenesis scenarios, the SM flavor asymmetries $Y_{\Delta\alpha}$ are constantly generated, since the CP -violating processes always occur out of equilibrium. They eventually attain their final value as soon as the DM state freezes in at $z \sim 3-5$. On the other hand, as the heavy leptons decay, their asymmetries $Y_{\Delta F_i}$ decrease and eventually vanish at high z . The asymmetry in F_2 is many orders of magnitude smaller than the one in F_1 ,

FIG. 4. The generated DM abundance Y_S , the asymmetries $Y_{\Delta\alpha}$, $Y_{\Delta F_1}$, and the baryon asymmetry Y_B as a function of the dimensionless parameter $z = M_1/T$, for the heavy leptons masses and feeble couplings given in Eq. (40). The dashed line illustrates the evolution that Y_B would follow in the absence of sphaleron decoupling, while the vertical line denotes the temperature at which sphalerons depart from equilibrium.

FIG. 5. Combinations of the feeble couplings and the heavy lepton mass which can generate the observed baryon asymmetry, for the set of parameters depicted in Table I. The dashed lines depict representative DM masses $m_S = \{2, 4, 6, 16\}$ keV for which the observed DM abundance of the Universe can be reproduced. Masses below 4 keV are excluded from current Lyman- α forest observations (gray shaded area).

due to its smaller couplings $\lambda_{\alpha 2} \ll \lambda_{\alpha 1}$, and is not shown in Fig. 4. Once the F_i have decayed away, conservation of Y_{B-L} implies that also $Y_{B-L_{SM}} \rightarrow 0$. Thus, the model predicts an equal amount of baryon Y_B and SM lepton $Y_{L_{SM}}$ asymmetries left in the Universe [30].

Our numerical analysis reveals that at least one of the feeble couplings must have a magnitude larger than $\sim 10^{-8}$ in order to account for the observed baryon asymmetry. Our main results are summarized in Fig. 5, where we show the combinations of the feeble couplings and the heavy lepton mass for which a viable baryon asymmetry can be generated, for the two illustrative scenarios depicted in Table I.

In both cases, the CP asymmetry is resonantly enhanced; i.e., the mass of F_2 is $M_2 = M_1 + \Gamma_2/2$, and we consider $p_{e\mu}^{12} = p_{e\tau}^{12} = p_{\mu\tau}^{12} = 1$ for the phases of the feeble couplings. The λ_{e1} coupling is being treated as a free parameter and lies in the 10^{-8} ballpark.

The dark matter relic density constraint can be satisfied throughout the parameter ranges depicted in the figure, for an appropriate choice of the mass m_S . However, in the shaded region, the required dark matter mass turns out to be in conflict with the Lyman- α bound of Eq. (23).

We observe that the first scenario (upper, blue line) is excluded from Lyman- α constraints for all values of the heavy lepton mass M_1 . On the contrary, the second scenario

(lower, orange line) is able to account for both the baryon asymmetry and the DM abundance for a wide range of M_1 , $550 \text{ GeV} \lesssim M_1 \lesssim 2800 \text{ GeV}$. Note that, interestingly, we find that for any combination of parameter values the relic density constraint is satisfied for DM masses that do not exceed $\sim 6 \text{ keV}$. This implies that if Lyman- α constraints become stronger in the future, they may be able to fully exclude the model's viable parameter space.

E. Phenomenological aspects

In the previous paragraphs, we saw that our simple model, described by the Lagrangian of Eq. (2), can provide a common explanation for the observed dark matter content and baryon asymmetry of the Universe. As we showed, this can be achieved by assuming highly degenerate heavy fermions F_i with masses $M_i \sim 1 \text{ TeV}$ and Yukawa-like couplings of the order of $\lambda_{\alpha 1} \sim 10^{-8} - 10^{-7}$ and $\lambda_{\alpha 2} \sim 10^{-10}$ for F_1 and F_2 , respectively.

Our model is intended to serve mainly as a proof of concept for the freeze-in baryogenesis mechanism that we propose. In this spirit, a full-blown analysis of its phenomenological predictions goes well beyond the scope of the present work, especially since alternative constructions assuming, e.g., different quantum numbers for the various particles involved, can lead to wildly different phenomenological signatures. Still, despite its simplicity, this model does exhibit an interesting phenomenology, which we will briefly comment upon. Note that a simpler variant of this model involving a single vectorlike heavy fermion F was already studied, e.g., in Refs. [64,65], mainly from the viewpoint of its predicted dark matter abundance as well as its signatures in new physics searches at the LHC.

First, the heavy fermions can mediate lepton flavor-violating decays of the type $\ell_\beta \rightarrow \ell_\alpha \gamma$. In particular, given the constraint $\text{Br}(\mu \rightarrow e \gamma) < 4.2 \times 10^{-13}$ [66] and assuming that (as was the case throughout the analysis presented in the previous sections) $\lambda_{e1} \sim \lambda_{\mu 1}$, the analysis presented in Ref. [65] restricts the product of the Yukawa-like couplings of F_1 to the first two generation leptons as

$$\frac{\sqrt{\lambda_{e1}\lambda_{\mu 1}}}{M_1} \lesssim 3.6 \times 10^{-5} \text{ GeV}^{-1}, \quad (41)$$

which is easily satisfied in the cosmologically relevant part of our parameter space. In deriving this bound, we have neglected the interference between Feynman diagrams involving different species of heavy fermions running in the loop, since the couplings $\lambda_{\alpha 2} \ll \lambda_{\alpha 1}$. Besides, constraints stemming from measurements of the muon lifetime turn out to be subleading [65].

On the side of collider searches, let us first point out the fact that, given the feeble nature of the $\lambda_{\alpha i}$ couplings, the mean proper decay length of F_1 is of the order of a few centimeters, whereas that of F_2 is of the order of several meters. Given these lifetime ranges, both F_1 and F_2 can be

TABLE I. Values of the parameters used for the two scenarios shown in Fig. 5.

	$ \lambda_{\mu 1} $	$ \lambda_{\tau 1} $	$ \lambda_{e2} $	$ \lambda_{\mu 2} $	$ \lambda_{\tau 2} $
Scenario 1	10^{-9}	2×10^{-9}	9×10^{-10}	8×10^{-10}	9×10^{-10}
Scenario 2	9×10^{-9}	9×10^{-9}	9×10^{-10}	8×10^{-10}	9×10^{-10}

looked for in searches for LLPs at the LHC. In particular, as shown in Refs. [64,65], searches for displaced leptons accompanied by missing energy can target the production and decay of F_1 and, for the values of $\lambda_{\alpha 1}$ that we assume, already exclude F_1 masses up to ~ 400 – 500 GeV. F_2 , on the other hand, survives long enough to escape the detector and can be probed by searches for heavy stable charged particles, with masses smaller than ~ 500 – 600 GeV being already excluded. The corresponding searches at the high-luminosity run of the LHC will probe $F_{1,2}$ masses reaching up to ~ 800 GeV and ~ 1.5 TeV, respectively [64].

In summary, if our model is to simultaneously explain the dark matter abundance and the baryon asymmetry of the Universe while remaining consistent with Lyman- α bounds, then it predicts interesting signatures for LLP searches at the LHC. In conjunction with the fact that, as we showed, explaining the matter-antimatter asymmetry of the Universe tends to favor $F_{1,2}$ masses lying in the range of a few TeV, we can hope that a substantial part of the cosmologically favored part of the parameter space will be probed by the LHC within the next few years.

IV. CONCLUSIONS AND OUTLOOK

In this paper, we discussed a mechanism for the simultaneous generation of the dark matter density and the baryon asymmetry of the Universe. To this goal, we relied on the out-of-equilibrium decays of heavy bath particles into a (feebly coupled) dark matter state along with Standard Model charged fermions, which leads to dark matter production via the freeze-in mechanism. If, moreover, CP is violated by these same decay processes, a viable matter-antimatter asymmetry can also be generated either directly (if the decays also lead to baryon or lepton number violation) or through the interplay of the generated CP asymmetry with the SM electroweak sphalerons.

As a proof of concept, we employed a simple model in which the role of the heavy bath particles was played by two $SU(3)_c \times SU(2)_L$ -singlet vectorlike fermions with a nonzero hypercharge, and dark matter was identified with a gauge singlet real scalar field. We showed that, indeed, such a simple construction can lead to both dark matter production and successful baryogenesis, and we briefly discussed the phenomenological prospects of such a construction, particularly in relation to searches for long-lived particles at the LHC.

Our proposal draws inspiration from the idea presented in Ref. [24]: First, the dark matter content and the baryon asymmetry of the Universe are generated simultaneously through the freeze-in mechanism. Second, at least within the framework of our concrete incarnation of this proposal, in both cases the electroweak sphalerons are used in order to convert a CP asymmetry into a baryonic one. In our case, however, instead of considering oscillations—as in ARS leptogenesis—we rather relied on decays in order to generate this initial CP asymmetry—a process which is

somewhat reminiscent of GUT baryogenesis. Moreover, we presented a general thermodynamical treatment that can also be applied in cases in which the baryon or lepton number is violated already at the decay level.

There are several topics which we did not expand upon in this paper and which would merit much more detailed investigation. First, more elaborate models can certainly be developed in this context of freeze-in baryogenesis, potentially establishing connections with theoretically well-motivated UV completions of the Standard Model. As an example, we could mention the fact that we assumed rather *ad hoc* values for the couplings of the heavy bath particles to dark matter and to the SM fermions. More realistic flavor structures can be envisaged in quite a straightforward manner, whereas nothing forbids dark matter itself to be asymmetric, along the lines described in Ref. [20]. Second, albeit related to the previous point, our discussion of more phenomenological aspects has been fairly elementary. Concrete, well-motivated constructions can have important implications for flavor physics (e.g., flavor-violating decays of or contributions to electric or magnetic moments of SM fermions), collider physics (e.g., new resonances and/or final states in LHC searches, potentially involving long-lived particles [67]), or cosmological measurements (e.g., direct detection experiments or cosmic microwave background observations). Such considerations, which are typically fairly model dependent, are left for future work.

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APPENDIX: SPECTATOR PROCESSES

During the baryon and/or lepton number violation era, there are various processes which can modify the number densities of the particles species. They are called “spectator

processes,” because they affect the baryon or lepton number indirectly by distributing the asymmetry among all SM fermions and Higgs fields. These processes include the gauge interactions, the Yukawa interactions, and the electroweak and QCD nonperturbative sphaleron transitions. If these processes are in thermal equilibrium with the cosmic plasma, then their effect is to impose certain relations among the chemical potentials of the various particle species. Recall that the asymmetry in the particle and antiparticle equilibrium number densities of species k , denoted as $n_{\Delta k} \equiv n_k - n_{\bar{k}}$, is given in the ultrarelativistic limit $m_k \ll T$ and $\mu_k \ll T$ by

$$n_{\Delta k} \equiv n_k - n_{\bar{k}} = \begin{cases} g_k \mu_k \frac{T^2}{3}, & k: \text{boson}, \\ g_k \mu_k \frac{T^2}{6}, & k: \text{fermion}, \end{cases} \quad (\text{A1})$$

where μ_k is the chemical potential and g_k are the independent degrees of freedom of species k . The total baryon n_B and lepton n_L asymmetries can be written as

$$n_B = \sum_k B_k n_{\Delta k}, \quad n_L = \sum_k L_k n_{\Delta k}, \quad (\text{A2})$$

where B_k and L_k are the baryon and lepton numbers of species k , respectively.

In our case, the lepton asymmetry is mainly generated during $T_{\text{EW}} \leq T \ll 10^8$ GeV when all spectator processes are active, including the sphaleron processes. These violate the $n_{B+L} \equiv n_B + n_L$ asymmetry but conserve the orthogonal combination $n_{B-L} \equiv n_B - n_L$. A rough estimation is that sphalerons will erase the n_{B+L} asymmetry, resulting in

$$n_B = \frac{n_{B+L}}{2} + \frac{n_{B-L}}{2} \xrightarrow{\text{sphalerons}} n_B = \frac{n_{B-L}}{2}, \quad (\text{A3a})$$

$$n_L = \frac{n_{B+L}}{2} - \frac{n_{B-L}}{2} \xrightarrow{\text{sphalerons}} n_L = -n_B = -\frac{n_{B-L}}{2}. \quad (\text{A3b})$$

A more careful treatment shows that, in spite of the fact that sphalerons violate n_{B+L} , thermodynamic equilibrium requires a nonzero value, i.e., $n_{B+L} \neq 0$. To quantify the asymmetry distribution, one can always relate the n_{B-L} number density asymmetry with n_B and n_L as follows:

$$\begin{aligned} n_{B-L} &\equiv n_B - n_L = c n_{B-L} - (c-1) n_{B-L} \\ &\Rightarrow \begin{cases} n_B = c n_{B-L}, \\ n_L = (c-1) n_{B-L}, \end{cases} \end{aligned} \quad (\text{A4})$$

where $c \lesssim 1$ is the constant of proportionality. To evaluate it, we assign chemical potentials to the particle spectrum of the model. This consists of all the SM species with three families and one Higgs doublet, together with the two BSM vectorlike fermions F_i .

1. SM

Let us first consider only the SM spectrum and assign the following nonvanishing chemical potentials [68,69]:

$$\mu_{q_\alpha} \quad (\text{LH quark doublets}), \quad (\text{A5a})$$

$$\mu_{u_\alpha} \quad (\text{RH up-quark singlets}), \quad (\text{A5b})$$

$$\mu_{d_\alpha} \quad (\text{RH down-quark singlets}), \quad (\text{A5c})$$

$$\mu_{\ell_\alpha} \quad (\text{LH lepton doublets}), \quad (\text{A5d})$$

$$\mu_{e_\alpha} \quad (\text{RH charged lepton singlets}), \quad (\text{A5e})$$

$$\mu_H \quad (\text{Higgs doublet}, H = (h^+, h^0)^T), \quad (\text{A5f})$$

where $\alpha = \{1, 2, 3\}$ is a SM flavor index.⁴ The corresponding antiparticle states have opposite chemical potentials, e.g., $\mu_{\tilde{H}} = -\mu_H$, where $\tilde{H} = (h^{0*}, -h^-)^T$. Note that the electrically chargeless gauge bosons have vanishing chemical potentials, as they carry no conserved quantum number, while the chemical potential of the $W^\pm$ gauge bosons vanishes, because the third component of the weak isospin is zero at $T > T_{\text{EW}}$ [68]. Hence, the number density asymmetries of the various components can be written, according to Eq. (A1), as

$$n_{\Delta q_\alpha} \equiv n_{q_\alpha} - n_{\bar{q}_\alpha} = 6\mu_{q_\alpha} \frac{T^2}{6}, \quad (\text{A6a})$$

$$n_{\Delta u_\alpha} \equiv n_{u_\alpha} - n_{\bar{u}_\alpha} = 3\mu_{u_\alpha} \frac{T^2}{6}, \quad (\text{A6b})$$

$$n_{\Delta d_\alpha} \equiv n_{d_\alpha} - n_{\bar{d}_\alpha} = 3\mu_{d_\alpha} \frac{T^2}{6}, \quad (\text{A6c})$$

$$n_{\Delta \ell_\alpha} \equiv n_{\ell_\alpha} - n_{\bar{\ell}_\alpha} = 2\mu_{\ell_\alpha} \frac{T^2}{6}, \quad (\text{A6d})$$

$$n_{\Delta e_\alpha} \equiv n_{e_\alpha} - n_{\bar{e}_\alpha} = \mu_{e_\alpha} \frac{T^2}{6}, \quad (\text{A6e})$$

$$n_{\Delta H} \equiv n_H - n_{\tilde{H}} = 2\mu_H \frac{T^2}{3} = 4\mu_H \frac{T^2}{6}. \quad (\text{A6f})$$

The baryon and lepton number density asymmetries for each flavor $(n_{B_\alpha}, n_{L_\alpha})$, as well as the corresponding total ones (n_B, n_L) , can be written in terms of the chemical potentials in Eq. (A5) as follows:

⁴Note that in this Appendix we are introducing a slightly different notation with respect to the rest of the paper, denoting the SM right-handed leptons by $\{e_1, e_2, e_3\}$ instead of $\{e_e, e_\mu, e_\tau\}$.

$$n_{B_\alpha} \equiv \sum_k B_{k_\alpha} n_{\Delta k_\alpha} = \frac{1}{3} (n_{\Delta q_\alpha} + n_{\Delta u_\alpha} + n_{\Delta d_\alpha})$$

$$= (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6}, \quad (\text{A7a})$$

$$n_B \equiv \sum_{\alpha=1}^3 n_{B_\alpha} = \sum_{\alpha=1}^3 (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6}, \quad (\text{A7b})$$

$$n_{L_\alpha} \equiv \sum_k L_{k_\alpha} n_{\Delta k_\alpha} = n_{\Delta \ell_\alpha} + n_{\Delta e_\alpha} = (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6}, \quad (\text{A7c})$$

$$n_L \equiv \sum_{\alpha=1}^3 n_{L_\alpha} = \sum_{\alpha=1}^3 (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6}. \quad (\text{A7d})$$

These 16 SM chemical potentials are not totally independent; they are related through the processes that attain chemical equilibrium during the era of the asymmetry generation. At $T \ll 10^8$ these are

- (i) $SU(2)_L$ sphaleron transitions which induce an effective 12-fermion effective operator $\mathcal{O}_{(B-L)_{\text{SM}}} = \sum_{i=1}^{N_f} (q_i q_i q_i \ell_i)$, which implies

$$\sum_{\alpha=1}^3 (3\mu_{q_\alpha} + \mu_{\ell_\alpha}) = 0; \quad (\text{A8})$$

- (ii) $SU(3)_c$ sphaleron transitions which give rise to

$$\sum_{\alpha=1}^3 (2\mu_{q_\alpha} - \mu_{u_\alpha} - \mu_{d_\alpha}) = 0; \quad (\text{A9})$$

- (iii) Yukawa interactions which imply

$$\mu_{q_\alpha} - \mu_H - \mu_{d_\alpha} = 0, \quad (q_\alpha \tilde{H} \bar{d}_\alpha + \text{H.c.}), \quad (\text{A10a})$$

$$\mu_{q_\alpha} + \mu_H - \mu_{u_\alpha} = 0, \quad (q_\alpha H \bar{u}_\alpha + \text{H.c.}), \quad (\text{A10b})$$

$$\mu_{\ell_\alpha} - \mu_H - \mu_{e_\alpha} = 0, \quad (\ell_\alpha \tilde{H} \bar{e}_\alpha + \text{H.c.}). \quad (\text{A10c})$$

Out of the 11 constraints of Eqs. (A8)–(A10), only ten of them are linearly independent, as the QCD sphaleron-induced relation (A9) can be obtained by adding Eqs. (A10a) and (A10b). An additional independent constraint can be obtained from the hypercharge ($\Upsilon = Q - t_3$) neutrality condition, which implies

- (iv) hypercharge constraint

$$n_\Upsilon \equiv \sum_k \Upsilon_k n_{\Delta k} = 0$$

$$\Rightarrow \sum_{\alpha=1}^3 \left(\frac{1}{6} n_{\Delta q_\alpha} + \frac{2}{3} n_{\Delta u_\alpha} - \frac{1}{3} n_{\Delta d_\alpha} - \frac{1}{2} n_{\Delta \ell_\alpha} - n_{\Delta e_\alpha} \right)$$

$$+ \frac{1}{2} n_{\Delta H} = 0$$

$$\Rightarrow \sum_{\alpha=1}^3 (\mu_{q_\alpha} + 2\mu_{u_\alpha} - \mu_{d_\alpha} - \mu_{\ell_\alpha} - \mu_{e_\alpha}) + 2\mu_H = 0, \quad (\text{A11})$$

where Υ_k is the hypercharge of species k . Two more constraints are imposed by the equality of the baryon flavor asymmetries [54].

- (v) Baryon flavor asymmetry equality

$$n_{B_1} = n_{B_2} = n_{B_3}$$

$$\Rightarrow 2\mu_{q_3} + \mu_{u_3} + \mu_{d_3} = 2\mu_{q_2} + \mu_{u_2} + \mu_{d_2}$$

$$= 2\mu_{q_1} + \mu_{u_1} + \mu_{d_1}. \quad (\text{A12})$$

Hence, there are in total 13 independent constraints [Eqs. (A8), (A10)–(A12)], and, therefore, one can express the 16 SM chemical potentials (A5) in terms of three chemical potentials, that we chose to be $\{\mu_{\ell_\alpha}\}$. Solving the system of equations, one obtains

$$\mu_{q_\alpha} = -\frac{1}{9} (\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}), \quad (\text{A13a})$$

$$\mu_{u_\alpha} = \frac{5}{63} (\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}), \quad (\text{A13b})$$

$$\mu_{d_\alpha} = -\frac{19}{63} (\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}), \quad (\text{A13c})$$

$$\mu_{e_1} = \frac{1}{21} (17\mu_{\ell_1} - 4\mu_{\ell_2} - 4\mu_{\ell_3}), \quad (\text{A13d})$$

$$\mu_{e_2} = \frac{1}{21} (-4\mu_{\ell_1} + 17\mu_{\ell_2} - 4\mu_{\ell_3}), \quad (\text{A13e})$$

$$\mu_{e_3} = \frac{1}{21} (-4\mu_{\ell_1} - 4\mu_{\ell_2} + 17\mu_{\ell_3}), \quad (\text{A13f})$$

$$\mu_H = \frac{4}{21} (\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}). \quad (\text{A13g})$$

Now the baryon and lepton number density asymmetries can be written in terms of the chemical potentials μ_{ℓ_α} as

$$n_{B_\alpha} = (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6} = -\frac{4}{9}(\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}) \frac{T^2}{6}, \quad (\text{A14a})$$

$$n_B = 3n_{B_\alpha} = -\frac{4}{3}(\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}) \frac{T^2}{6}, \quad (\text{A14b})$$

$$n_{L_\alpha} = (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6}, \quad (\text{A14c})$$

$$n_L = \sum_{\alpha=1}^3 (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6} = \frac{17}{7}(\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}) \frac{T^2}{6}. \quad (\text{A14d})$$

The flavor asymmetry combination $n_{\Delta_\alpha} \equiv n_{B_\alpha} - n_{L_\alpha} = n_B/3 - n_{L_\alpha}$ can be written as

$$\begin{aligned} n_{\Delta_\alpha} &= (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha} - 2\mu_{\ell_\alpha} - \mu_{e_\alpha}) \frac{T^2}{6} \\ \Rightarrow \begin{pmatrix} n_{\Delta_1} \\ n_{\Delta_2} \\ n_{\Delta_3} \end{pmatrix} &= -\frac{1}{63} \begin{bmatrix} 205 & 16 & 16 \\ 16 & 205 & 16 \\ 16 & 16 & 205 \end{bmatrix} \begin{pmatrix} \mu_{\ell_1} \\ \mu_{\ell_2} \\ \mu_{\ell_3} \end{pmatrix} \frac{T^2}{6}. \end{aligned}$$

Inverting the expression above, one finds

$$\begin{pmatrix} \mu_{\ell_1} \\ \mu_{\ell_2} \\ \mu_{\ell_3} \end{pmatrix} = -\frac{1}{711} \begin{bmatrix} 221 & -16 & -16 \\ -16 & 221 & -16 \\ -16 & -16 & 221 \end{bmatrix} \begin{pmatrix} n_{\Delta_1} \\ n_{\Delta_2} \\ n_{\Delta_3} \end{pmatrix} \frac{6}{T^2}. \quad (\text{A15})$$

Using Eq. (A6), we can express the chemical potentials μ_{ℓ_α} in terms of the flavor asymmetries n_{Δ_α} and relate them to n_{Δ_α} as⁵

$$\begin{pmatrix} n_{\Delta_\ell_1} \\ n_{\Delta_\ell_2} \\ n_{\Delta_\ell_3} \end{pmatrix} = -\frac{2}{711} \begin{bmatrix} 221 & -16 & -16 \\ -16 & 221 & -16 \\ -16 & -16 & 221 \end{bmatrix} \begin{pmatrix} n_{\Delta_1} \\ n_{\Delta_2} \\ n_{\Delta_3} \end{pmatrix}. \quad (\text{A16})$$

The total $B - L$ asymmetry is

$$n_{B-L} \equiv n_B - n_L = -\frac{79}{21}(\mu_{\ell_1} + \mu_{\ell_2} + \mu_{\ell_3}) \frac{T^2}{6}. \quad (\text{A17})$$

Hence, in the SM and for $T_{sph} \ll T \ll 10^8$ GeV, the total baryon n_B and the n_L asymmetries are related to n_{B-L} by

$$n_B = \frac{28}{79} n_{B-L} = \frac{28}{79} \sum_{\alpha=1}^3 n_{\Delta_\alpha}, \quad (\text{A18})$$

⁵Note that our result agrees with Ref. [57], while the authors of Ref. [30] define n_{Δ_α} per gauge degree of freedom, and, therefore, their result is smaller by a factor of 2.

$$n_L = -\frac{51}{79} n_{B-L} = -\frac{51}{79} \sum_{\alpha=1}^3 n_{\Delta_\alpha}. \quad (\text{A19})$$

2. Our model

We extend the SM to include also two heavy vectorlike leptons F_1 and F_2 with chemical potentials μ_{F_1} and μ_{F_2} , respectively, and asymmetry

$$n_{\Delta F_i} = \mu_{F_i} \frac{T^2}{6}. \quad (\text{A20})$$

Since the BSM fermions are gauged under $U(1)_Y$, with $\Upsilon_{F_i} = \Upsilon_e = -1$, the hypercharge constraint (A11) is modified to

$$\begin{aligned} n_Y &\equiv \sum_k \Upsilon_k n_{\Delta k} = 0 \\ &\Rightarrow \sum_{\alpha=1}^3 \left(\frac{1}{6} n_{\Delta q_\alpha} + \frac{2}{3} n_{\Delta u_\alpha} - \frac{1}{3} n_{\Delta d_\alpha} - \frac{1}{2} n_{\Delta \ell_\alpha} - n_{\Delta e_\alpha} \right) \\ &\quad + \frac{1}{2} n_{\Delta H} - \sum_{i=1}^2 n_{\Delta F_i} = 0 \\ &\Rightarrow \sum_{\alpha=1}^3 (\mu_{q_\alpha} + 2\mu_{u_\alpha} - \mu_{d_\alpha} - \mu_{\ell_\alpha} - \mu_{e_\alpha}) + 2\mu_H \\ &\quad - (\mu_{F_1} + \mu_{F_2}) = 0. \end{aligned} \quad (\text{A21})$$

The theory conserves the total $n_{B-L} \equiv n_B - n_L = n_B - n_{L_{SM}} - n_{L_F}$ asymmetry, where

$$\begin{aligned} n_{B_\alpha} &\equiv \sum_k B_{k_\alpha} n_{\Delta k_\alpha} = \frac{1}{3} (n_{\Delta q_\alpha} + n_{\Delta u_\alpha} + n_{\Delta d_\alpha}) \\ &= (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6}, \end{aligned} \quad (\text{A22a})$$

$$\begin{aligned} n_B &\equiv \sum_{\alpha=1}^3 n_{B_\alpha} = \sum_{\alpha=1}^3 (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6} \\ &= 3(2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6}, \end{aligned} \quad (\text{A22b})$$

$$\begin{aligned} n_{L_\alpha} &\equiv \sum_{k:SM} L_{k_\alpha} n_{\Delta k_\alpha} = n_{\Delta \ell_\alpha} + n_{\Delta e_\alpha} \\ &= (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6}, \end{aligned} \quad (\text{A22c})$$

$$n_{L_{SM}} \equiv \sum_{\alpha=1}^3 n_{L_\alpha} = \sum_{\alpha=1}^3 (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6}, \quad (\text{A22d})$$

$$n_{L_F} \equiv \sum_{i=1}^2 n_{\Delta F_i} = \sum_{i=1}^2 \mu_{F_i} \frac{T^2}{6}, \quad (\text{A22e})$$

$$n_L \equiv n_{L_{\text{SM}}} + n_{L_F} = \left(\sum_{\alpha=1}^3 (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) + \sum_{i=1}^2 \mu_{F_i} \right) \frac{T^2}{6}. \quad (\text{A22f})$$

If we also assume that the early Universe is totally symmetric, $n_{B-L_{\text{SM}}}|_0 = 0$ and $n_F|_0 = 0$, then we obtain an additional constraint to $\mu_{F_1} + \mu_{F_2}$, that is,

$$\begin{aligned} n_B - n_{L_{\text{SM}}} - n_{L_F} &= \sum_{\alpha=1}^3 n_{\Delta_\alpha} - \sum_{i=1}^2 n_{\Delta_{F_i}} = 0 \\ \Rightarrow \sum_{\alpha=1}^3 (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha} - 2\mu_{\ell_\alpha} - \mu_{e_\alpha}) - (\mu_{F_1} + \mu_{F_2}) &= 0. \end{aligned} \quad (\text{A23})$$

Replacing this in Eq. (A21), we obtain the modified hypercharge constraint expressed in terms of the SM chemical potentials.

(vi) Modified hypercharge constraint

$$\sum_{\alpha=1}^3 (-\mu_{q_\alpha} + \mu_{u_\alpha} - 2\mu_{d_\alpha} + \mu_{\ell_\alpha}) + 2\mu_H = 0. \quad (\text{A24})$$

Now the 16 SM chemical potentials are constrained under the $SU(2)$ sphaleron (A8) and Yukawa (A10) processes (ten constraints), as well as under the baryon flavor asymmetry equality (A12) and the modified hypercharge (A24) conditions (three constraints). Solving the system of equations, we may express all chemical potentials in terms of μ_{e_α} :

$$\mu_{q_\alpha} = -\frac{11}{144}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}), \quad (\text{A25a})$$

$$\mu_{u_\alpha} = -\frac{13}{72}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}), \quad (\text{A25b})$$

$$\mu_{d_\alpha} = \frac{1}{36}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}), \quad (\text{A25c})$$

$$\mu_{\ell_1} = \frac{1}{48}(43\mu_{e_1} - 5\mu_{e_2} - 5\mu_{e_3}), \quad (\text{A25d})$$

$$\mu_{\ell_2} = \frac{1}{48}(-5\mu_{e_1} + 43\mu_{e_2} - 5\mu_{e_3}), \quad (\text{A25e})$$

$$\mu_{\ell_3} = \frac{1}{48}(-5\mu_{e_1} - 5\mu_{e_2} + 43\mu_{e_3}), \quad (\text{A25f})$$

$$\mu_H = -\frac{5}{48}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}). \quad (\text{A25g})$$

Now the baryon and lepton number density asymmetries can be written in terms of the chemical potentials μ_{e_α} as

$$n_{B_\alpha} = (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha}) \frac{T^2}{6} = -\frac{11}{36}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}) \frac{T^2}{6}, \quad (\text{A26a})$$

$$n_B = 3n_{B_\alpha} = -\frac{11}{12}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}) \frac{T^2}{6}, \quad (\text{A26b})$$

$$n_{L_\alpha} = (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6}, \quad (\text{A26c})$$

$$n_{L_{\text{SM}}} = \sum_{\alpha=1}^3 (2\mu_{\ell_\alpha} + \mu_{e_\alpha}) \frac{T^2}{6} = \frac{19}{8}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}) \frac{T^2}{6}, \quad (\text{A26d})$$

$$\begin{aligned} n_{L_F} &= (\mu_{F_1} + \mu_{F_2}) \frac{T^2}{6} \\ &= \sum_{\alpha=1}^3 (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha} - 2\mu_{\ell_\alpha} - \mu_{e_\alpha}) \frac{T^2}{6} \\ &= -\frac{79}{24}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}) \frac{T^2}{6} = n_B - n_{L_{\text{SM}}}, \end{aligned} \quad (\text{A26e})$$

$$n_L \equiv n_{L_{\text{SM}}} + n_{L_F} = -\frac{11}{12}(\mu_{e_1} + \mu_{e_2} + \mu_{e_3}) \frac{T^2}{6} = n_B. \quad (\text{A26f})$$

The flavor asymmetry combination $n_{\Delta_\alpha} \equiv n_{B_\alpha} - n_{L_\alpha} = n_B/3 - n_{L_\alpha}$ can be written as

$$\begin{aligned} n_{\Delta_\alpha} &= (2\mu_{q_\alpha} + \mu_{u_\alpha} + \mu_{d_\alpha} - 2\mu_{\ell_\alpha} - \mu_{e_\alpha}) \frac{T^2}{6} \\ \Rightarrow \begin{pmatrix} n_{\Delta_1} \\ n_{\Delta_2} \\ n_{\Delta_3} \end{pmatrix} &= -\frac{1}{72} \begin{bmatrix} 223 & 7 & 7 \\ 7 & 223 & 7 \\ 7 & 7 & 223 \end{bmatrix} \begin{pmatrix} \mu_{e_1} \\ \mu_{e_2} \\ \mu_{e_3} \end{pmatrix} \frac{T^2}{6}. \end{aligned}$$

Inverting the expression above, one finds

$$\begin{pmatrix} \mu_{e_1} \\ \mu_{e_2} \\ \mu_{e_3} \end{pmatrix} = -\frac{1}{711} \begin{bmatrix} 230 & -7 & -7 \\ -7 & 230 & -7 \\ -7 & -7 & 230 \end{bmatrix} \begin{pmatrix} n_{\Delta_1} \\ n_{\Delta_2} \\ n_{\Delta_3} \end{pmatrix} \frac{6}{T^2}. \quad (\text{A27})$$

Using Eq. (A6), we can express the chemical potentials μ_{e_α} in terms of the flavor asymmetries $n_{\Delta_{e_\alpha}}$ and relate them to n_{Δ_α} as

$$\begin{pmatrix} n_{\Delta_{e_1}} \\ n_{\Delta_{e_2}} \\ n_{\Delta_{e_3}} \end{pmatrix} = -\frac{1}{711} \begin{bmatrix} 230 & -7 & -7 \\ -7 & 230 & -7 \\ -7 & -7 & 230 \end{bmatrix} \begin{pmatrix} n_{\Delta_1} \\ n_{\Delta_2} \\ n_{\Delta_3} \end{pmatrix}. \quad (\text{A28})$$

Hence, in this model and for $T_{sph} \ll T \ll 10^8$ GeV, the total baryon n_B and the SM lepton $n_{L_{SM}}$ asymmetries are related to $n_{B-L_{SM}}$ by

$$n_B = \frac{22}{79} n_{B-L_{SM}} = \frac{22}{79} \sum_{\alpha=1}^3 n_{\Delta\alpha} \quad (A29)$$

$$n_{L_{SM}} = -\frac{57}{79} n_{B-L_{SM}} = -\frac{57}{79} \sum_{\alpha=1}^3 n_{\Delta\alpha}. \quad (A30)$$

This result agrees with that obtained in Ref. [24].

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