Leptogenesis via Axion Oscillations after Inflation

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Once a light axionlike scalar field couples to the electroweak gauge bosons, its classical motion during reheating induces an effective chemical potential for the fermion number. In the presence of rapid lepton number (L)-violating processes in the plasma, such a chemical potential provides a favorable opportunity for baryogenesis via leptogenesis. We are able to demonstrate that L violation due to the exchange of heavy Majorana neutrinos is sufficient for a successful realization of this idea. Our mechanism represents a novel and minimal alternative to thermal leptogenesis, which turns out to be insensitive to the masses and CP-violating phases in the heavy neutrino sector. It is consistent with heavy neutrino masses close to the scale of grand unification and, quite complementary to thermal leptogenesis, requires the reheating temperature to be at least of order 10^{12} GeV.

The Peccei-Quinn (PQ) solution to the strong CP problem [1] has led to the prediction of a light scalar field called the axion [2], which appears generically in a number of models [3,4], and which has a characteristic coupling to the gauge fields of the form $f_a^{-1} a F \tilde{F}$, where f_a is the scale of PQ symmetry breaking. However, the motivation for considering axionic fields extends well beyond the context of the strong CP problem. Axions are ubiquitous in string theory, where at least one such field is generically associated with the Green-Schwarz mechanism of anomaly cancellation [5] and the scale f_a of which lies a few orders of magnitude below the Planck scale [6], but multiple other axions can also appear. While the model-independent axion is coupled to all gauge groups with a universal coupling strength, the additional axionic fields can couple to different groups with couplings that depend on both the gauge group and the particle content of the model [7]. The masses of these model-dependent axions can be different as they arise from their couplings to different anomalous groups. We will focus, in particular, on the axion (or linear combination of axions) that has a coupling to the electroweak $SU(2)$ gauge fields.

During inflation, light scalar fields develop large expectation values [8]. The relaxation of the axion field to the minimum of its effective potential begins once the Hubble rate becomes comparable to the axion mass. While the field $a(t)$ is slowly rolling towards its origin, its coupling to the $SU(2)$ gauge fields and, via the anomaly, to the fermionic current $j^\mu = \bar{\psi} \gamma^\mu \psi$ induces an effective CPT-violating term $a(t) F \tilde{F} \propto [\partial_t a(t)] j^0$, which serves as a chemical potential for fields carrying nonzero baryon or lepton number. (This illustrates that, while the axion could equally couple to the hypercharge gauge boson, a coupling to the standard model gluons would, by contrast, not provide a sufficient basis for leptogenesis.) Then, in the presence of rapid lepton

number-violating processes in the plasma—for example, due to the exchange of virtual heavy right-handed neutrinos—the conditions for successful leptogenesis are satisfied. Here, one important detail is that lepton number violation needs to occur at a fast rate, $\Gamma_L \gg \dot{a}/a$, so that the axion field acts as an adiabatic background during leptogenesis. It is therefore essential that the generation of the lepton asymmetry be driven by an external source, i.e., by scatterings with heavy neutrinos in the intermediate state in our case. If the rate of lepton number violation was instead tied to the change in the axion field value, we would, by contrast, not be able to interpret the axion velocity as an effective chemical potential [9]. As we will see, a consequence of the requirement of a fast rate Γ_L is, in particular, that we will have to adopt large values for the reheating temperature, $T_{\rm rh} \gtrsim 10^{12}$ GeV.

A similar scenario was discussed in connection with a flat direction that carries no baryon or lepton number [10]. Our treatment of the asymmetry is different, and we obtain very different results. Our scenario is also similar to leptogenesis via Higgs relaxation [11], where the Higgs coupling to $F\tilde{F}$ is assumed to arise from a higher-dimensional operator, unlike in the present scenario, where the required coupling appears generically for any axion coupled to the electroweak gauge fields. One can also draw an analogy to models of spontaneous baryogenesis [12], in particular to realizations of spontaneous baryogenesis during the electroweak phase transition [13]. Here, the Higgs field in the expanding bubble wall generates an effective chemical potential for the fermions. Our scenario is different in that the "wall" is represented by an axionic field moving in the timelike direction uniformly in space, unlike the bubble wall moving in a spacelike direction. Finally, we note that, in the context of electroweak baryogenesis, also the QCD axion may be employed to generate an effective chemical potential [14]. Successful baryogenesis then requires that the electroweak

phase transition be delayed below the GeV scale, for instance, due to some additional Higgs-dilaton coupling [15].

Provided that the axion a is to be identified with the pseudo-Nambu-Goldstone boson of a spontaneously broken $U(1)$ symmetry with a compact global topology, its initial value at the end of inflation is $a_0 = f_a \theta_0$, where the angle θ_0 takes a random value in the range $\theta_0 \in [0, 2\pi)$. Assuming that the PQ symmetry is broken sufficiently early before the end of inflation (and is not restored during reheating), this initial value ends up being constant on superhorizon scales. For definiteness, we set $a_0 = f_a$ and treat f_a as a free parameter in the following. Anticipating that the final baryon asymmetry will depend on a_0, we require that the baryonic isocurvature perturbations induced by the quantum fluctuations of the axion field during inflation be smaller than the observational upper limit. This implies a constraint on the Hubble rate during inflation: $H_{\text{inf}}/(2\pi)/a_0 \lesssim 10^{-5}$ [16], or

$$H_{\text{inf}} \lesssim 6 \times 10^{11} \text{ GeV} \left(\frac{f_a}{10^{15} \text{ GeV}} \right). \qquad (1)$$

The evolution of the homogeneous axion field in its effective potential V_{eff} around the origin is described by

$$\ddot{a} + 3H\dot{a} = -\partial_a V_{\text{eff}}, \qquad V_{\text{eff}} \approx \frac{1}{2} m_a^2 a^2, \qquad (2)$$

where we have neglected the backreaction of lepton number generation on the evolution of the axion field. m_a denotes the axion mass, which we assume to arise via dimensional transmutation, i.e., from an additional coupling of the axion to the gauge fields of some strongly coupled hidden sector. Given a dynamical scale Λ_H in this hidden sector, the axion mass is then of $\mathcal{O}(\Lambda_H^2/f_a)$. For consistency, we require m_a to be smaller than H_{inf}, the Hubble rate at the end of inflation:

$$m_a \lesssim H_{\text{inf}}. \qquad (3)$$

When inflation is over, the axion field remains practically at rest until the Hubble parameter drops to $H_{\text{osc}} = m_a$. Once the axion field is in motion, the effective Lagrangian contains the term

$$\mathcal{L}_{\text{eff}} \supset \frac{g_2^2}{32\pi^2} \frac{a(t)}{f_a} F\tilde{F} = -\frac{a(t)}{N_f f_a} \partial_\mu (\bar{\psi} \gamma^\mu \psi) \qquad (4)$$

$$= \frac{\partial_t a(t)}{N_f f_a} (\bar{\psi} \gamma^0 \psi) + \cdots = \mu_{\text{eff}} j^0 + \cdots, \qquad (5)$$

with g_2 being the $SU(2)$ gauge coupling and $N_f = 3$ the number of fermion generations in the standard model, where we have used the anomaly equation in Eq. (4), and integration by parts in Eq. (5). In the

following, we will absorb N_f in our definition of f_a and simply determine the effective chemical potential as $\mu_{\text{eff}} = \dot{a}/f_a$.

Now the necessary conditions for generating a lepton asymmetry are satisfied. A nonzero effective chemical potential shifts the energy levels of particles as compared to antiparticles. If the lepton number is not conserved, the minimum of the free energy in the plasma is reached for a different number density of leptons than for antileptons, i.e., for $n_L \equiv n_\ell - n_{\bar{\ell}} \neq 0$. Instead, if the lepton number violation is very rapid, the minimum of the free energy is obtained for an equilibrium number density of (In the following, we shall approximate all number and energy densities by their corresponding expressions in the classical Boltzmann approximation. We will only take care of quantum-statistical effects when counting relativistic degrees of freedom.)

$$n_L^{\text{eq}} = \frac{4}{\pi^2} \mu_{\text{eff}} T^2. \qquad (6)$$

Lepton number violation is mediated by the exchange of heavy neutrinos. In contrast to thermal leptogenesis [17], we will assume all heavy right-handed neutrino masses to be close to the scale of grand unification (GUT), $M_i \sim \mathcal{O}(10^{-1} \cdots 1)\Lambda_{\text{GUT}} \sim 10^{15} \cdots 10^{16}$ GeV, so that the heavy neutrinos are never produced thermally, i.e., $T \ll M_i$ at all times. This assumption serves the purpose to separate the mechanism under study from the contributions from ordinary thermal leptogenesis. In the expanding universe, the evolution of the L number density n_L is then described by the Boltzmann equation

$$\dot{n}_L + 3Hn_L \simeq -\Gamma_L (n_L - n_L^{\text{eq}}), \qquad \Gamma_L = 4n_\ell^{\text{eq}} \sigma_{\text{eff}}, \quad (7)$$

where $n_\ell^{\text{eq}} = 2/\pi^2 T^3$ and with $\sigma_{\text{eff}} \equiv \langle \sigma_{\Delta L=2} v \rangle$ denoting the thermally averaged cross section of two-to-two scattering processes with heavy neutrinos in the intermediate state that violate lepton number by two units,

$$\Delta L = 2: \ \ell_i \ell_j \leftrightarrow HH, \ \ell_i H \leftrightarrow \bar{\ell}_j \bar{H},$$
$$\ell_i^T = (\nu_i \quad e_i),$$
$$H^T = (h_+ \quad h_0), \quad i, j = 1, 2, 3. \qquad (8)$$

We note that the term proportional to n_L^{eq} now acts as a novel production term for the lepton asymmetry, as long as the axion field is in motion. For center-of-mass energies much smaller than the heavy neutrino mass scale, $\sqrt{s} \ll M_i$, the effective cross section σ_{eff} is practically fixed by the experimental data on the light neutrino sector [18], assuming the seesaw mass matrix [19]:

$$\sigma_{\text{eff}} \approx \frac{3}{32\pi} \frac{\bar{m}^2}{v_{\text{ew}}^4} \simeq 1 \times 10^{-31} \text{ GeV}^{-2}, \ \bar{m}^2 = \sum_{i=1}^{3} m_i^2, \quad (9)$$

where $v_{ew} \simeq 174$ GeV and where we have assumed that the sum of the light neutrino masses squared is of the same order of magnitude as the atmospheric neutrino mass difference, $\Delta m_{atm}^2 \simeq 2.4 \times 10^{-3}$ eV^2 [20].

For $a_0 \ll M_{Pl}$, and as long as $H \gg m_a$, i.e., prior to the onset of the axion oscillations, the axion energy density ρ_a is much smaller than the total energy density $\rho_{tot} = \rho_\varphi + \rho_R + \rho_a \approx \rho_\varphi + \rho_R$, where ρ_φ and ρ_R are the energy densities of the inflaton and of radiation. Reheating is described by a system of equations:

$$\dot{\rho}_\varphi + 3H\rho_\varphi = -\Gamma_\varphi \rho_\varphi, \qquad \dot{\rho}_R + 4H\rho_R = +\Gamma_\varphi \rho_\varphi, \qquad (10)$$

$$H^2 \equiv (\dot{R}/R)^2 = \frac{\rho_{tot}}{3M_{Pl}^2}, \qquad \rho_{tot} \approx (\rho_\varphi + \rho_R), \qquad (11)$$

where Γ_φ is the inflaton decay rate. The inflaton must not decay before the end of inflation, which implies

$$\Gamma_\varphi \lesssim H_{inf}. \qquad (12)$$

The rough temperature scale of leptogenesis as well as the axion mass scale in our scenario are determined by the requirement that the heavy neutrino-mediated $\Delta L = 2$ interactions must be in thermal equilibrium before the onset of axion oscillations, $\Gamma_L \gg H \gtrsim m_a$, which yields $T \sim T_L = g_*^{1/2}/(\pi\sigma_{eff}M_{Pl}) \sim 10^{13}$ GeV and $m_a \sim \sigma_{eff}T_L^3 \sim 10^8$ GeV. Upon closer examination, the solution for the temperature, $T^4 \equiv \pi^2/3/g_*\rho_R$, according to Eqs. (10) and (11) shows the following characteristic behavior: within roughly one Hubble time after the end of inflation, T quickly rises to its maximal value,

$$T_{max} \simeq 5 \times 10^{13} \text{ GeV} \left(\frac{\Gamma_\varphi}{10^9 \text{ GeV}}\right)^{1/4} \left(\frac{H_{inf}}{10^{11} \text{ GeV}}\right)^{1/2}, \qquad (13)$$

after which the temperature decreases because the energy density is dominated by the inflaton oscillations (which scale as matter). During reheating, the temperature drops as $T \propto R^{-3/8}$ until radiation comes to dominate at time $t = t_{rh} \simeq \Gamma_\varphi^{-1}$, when $\rho_R = \rho_\varphi$, and the reheating temperature is

$$T_{rh} \simeq 2 \times 10^{13} \text{ GeV} \left(\frac{\Gamma_\varphi}{10^9 \text{ GeV}}\right)^{1/2}. \qquad (14)$$

After the end of reheating, i.e., for $t > t_{rh}$, the expansion is then driven by relativistic radiation and the temperature simply decreases adiabatically, $T \propto R^{-1}$. In the case of a large axion decay constant, this phase of radiation domination, however, does not last all the way to the time of primordial nucleosynthesis. Instead, the axion comes to dominate the total energy density at some time prior to its
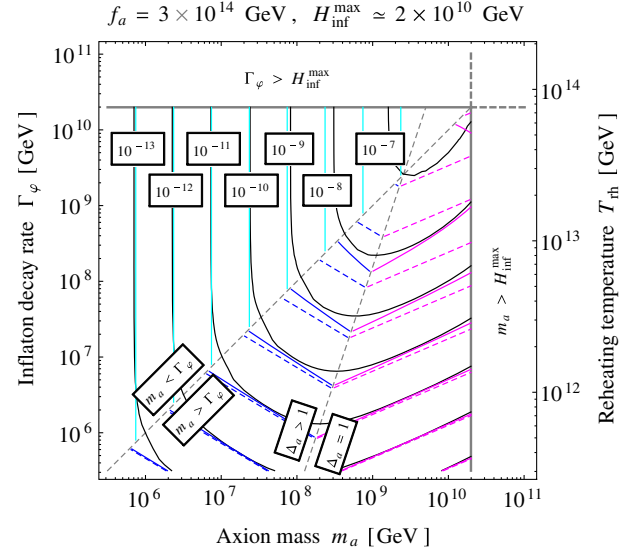

FIG. 1 (color online). Contour plot of the final baryon asymmetry η_B^0 as a function of the axion mass m_a and the inflaton decay rate Γ_φ for an axion decay constant of $f_a = 3 \times 10^{14}$ GeV. The black (bent) contours represent the full numerical result, while the colorful (straight) contours depict our analytical estimate according to Eqs. (18) and (19). In the lower part of the plot ($m_a > \Gamma_\varphi$), the effect of washout is illustrated by the difference between the dashed ($\kappa = 0$) and solid ($\kappa \neq 0$) lines.

decay, which marks the beginning of yet another stage of matter domination. The decay of the axion into relativistic gauge bosons and the corresponding renewed transition to radiation domination then represent a second installment of reheating, which can be described by the same set of equations as the primary reheating process, cf. Eqs. (10) and (11). With the axion decay rate

$$\Gamma_a \simeq \frac{\alpha^2}{64\pi^3}\frac{m_a^3}{f_a^2}, \qquad \alpha = \frac{g_2^2}{4\pi}, \qquad (15)$$

and using Eq. (14), we find for the *secondary reheating temperature* or axion decay temperature

$$T_{dec} \simeq 1 \times 10^4 \text{ GeV} \left(\frac{m_a}{10^9 \text{ GeV}}\right)^{3/2} \left(\frac{10^{15} \text{ GeV}}{f_a}\right). \qquad (16)$$

This temperature should be at least of $\mathcal{O}(10)$ MeV [21], which imposes a lower bound on m_a:

$$m_a \gtrsim 8 \times 10^4 \text{ GeV} \left(\frac{f_a}{10^{15} \text{ GeV}}\right)^{2/3}. \qquad (17)$$

The five differential equations in Eqs. (2), (7), (10), and (11) allow one to compute the present value of the baryon asymmetry (i.e., the baryon-to-photon ratio),

$$\eta_B^0 \equiv \frac{n_B^0}{n_\gamma^0} = c_{\rm sph} \frac{g_{*,s}^0}{g_*} \eta_L^a \simeq 0.013 \eta_L^a, \qquad (18)$$

where the sphaleron factor $c_{\rm sph}$ accounts for the conversion of the lepton asymmetry into baryon asymmetry by sphalerons. Here, $g_{*,s}^0$ and g_* denote the effective numbers of degrees of freedom contributing to the entropy density in the present epoch and during reheating, respectively. In the standard model $c_{\rm sph} = 28/79$, $g_{*,s}^0 = 43/11$, and $g_* = 427/4$. Last but not least, η_L^a in Eq. (18) stands for the final lepton asymmetry after the decay of the axion around $t \simeq \Gamma_a^{-1}$.

We determine η_L^a by solving the five differential equations in Eqs. (2), (7), (10), and (11) numerically. We also present approximate analytical solutions, which will be discussed in detail in an upcoming publication. It is convenient to parametrize η_L^a as follows:

$$\eta_L^a = C \Delta_a^{-1} \Delta_\varphi^{-1} \eta_L^{\rm max} e^{-\kappa}. \qquad (19)$$

The approximate analytical results agree with the numerical results, as shown in Fig. 1. In the following, we shall present analytical expressions for the individual factors on the right-hand side of Eq. (19). $\eta_L^{\rm max}$ denotes the all-time maximum value of the lepton asymmetry, which is reached around the time when the axion oscillations set it, i.e., at $t \sim t_{\rm osc} \simeq m_a^{-1}$. Note that this time does not necessarily coincide with the time when $\Gamma_L \simeq H$. Integrating the Boltzmann equation for the lepton asymmetry up to $t \sim t_{\rm osc}$, one approximately finds

$$\eta_L^{\rm max} \simeq \frac{\sigma_{\rm eff}}{g_*^{1/2}} \frac{a_0}{f_a} m_a M_{\rm Pl} \min\{1, (\Gamma_\varphi/m_a)^{1/2}\}, \qquad (20)$$

which is suppressed with respect to the would-be equilibrium lepton asymmetry, $\eta_L^{\rm eq} = n_L^{\rm eq}/n_\gamma^{\rm eq}$, evaluated at the same time by a factor $n_L/n_L^{\rm eq} \simeq T/T_L \min\{1, (\Gamma_\varphi/m_a)^{1/2}\}$. Remarkably enough, the maximal lepton asymmetry is rather insensitive to the axion decay constant; it only depends on the ratio a_0/f_a, which is expected to be $\mathcal{O}(1)$. Furthermore, $\eta_L^{\rm max}$ turns out to be directly proportional to the effective cross section $\sigma_{\rm eff}$. For $a_0 = f_a$ and given the value of $\sigma_{\rm eff}$ in Eq. (9), $\eta_L^{\rm max}$ is, hence, typically much larger than the observed value, $\eta_B^{\rm obs} \simeq 6 \times 10^{-10}$ [22],

$$\eta_L^{\rm max} \simeq 2 \times 10^{-5} \left(\frac{m_a}{10^9 \ {\rm GeV}}\right)^p \left(\frac{\Gamma_\varphi}{10^9 \ {\rm GeV}}\right)^q, \qquad (21)$$

where the powers p and q are given as in Eq. (20).

The two Δ factors in Eq. (19) account for the entropy production in inflaton and axion decays during reheating and at late times, respectively. We approximately have

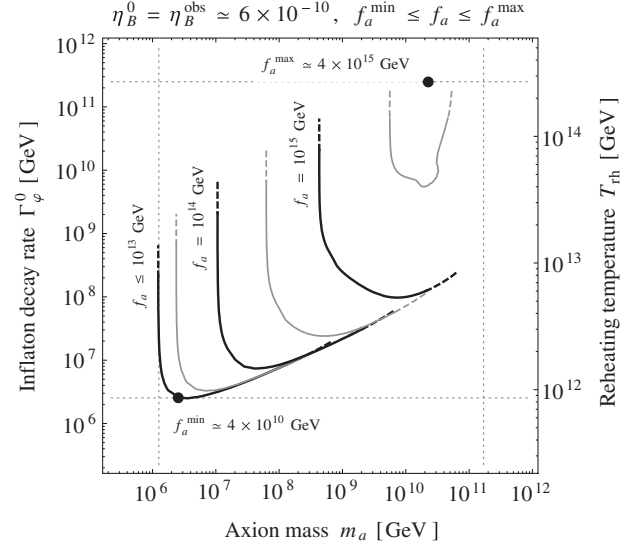$$\Delta_\varphi \simeq \max\{1, \Delta_\varphi'\}, \qquad \Delta_a \simeq \max\{1, \Delta_a'\}, \qquad (22)$$

FIG. 2. Contour lines for successful leptogenesis ($\eta_B^0 = \eta_B^{\rm obs}$) in the $m_a - \Gamma_\varphi$ plane for different values of the axion decay constant f_a. The dashed segments along the individual contours mark the regions where either m_a or Γ_φ become comparable to the maximally allowed Hubble rate $H_{\rm inf}^{\rm max}$, cf. Eq. (1). For $f_a \lesssim 10^{13}$ GeV, entropy production in axion decays ceases to affect η_B^0, cf. Eq. (23), which is reflected in the contour lines being no longer sensitive to changes in f_a. The lower bounds on m_a and f_a then directly follow from our constraints in Eqs. (1), (3), and (12). At the same time, the regime of large f_a, and hence the upper bounds on m_a and f_a, require a careful numerical analysis due to the strong impact of washout.

where $\Delta_\varphi' \simeq (m_a/\Gamma_\varphi)^{5/4}$ and with Δ_a' being given as

$$\Delta_a' \simeq \frac{2\pi^2}{\alpha} \frac{f_a a_0^2}{m_a M_{\rm Pl}^2} \min\{1, (\Gamma_\varphi/m_a)^{1/2}\}. \qquad (23)$$

In the region of parameter space in which we are able to successfully reproduce $\eta_B^{\rm obs}$, entropy production in axion decays begins to play a role for f_a values around 3×10^{13} GeV, cf. Fig. 2. For smaller values of f_a, we always have $\Delta_a = 1$ in the entire parameter region of interest.

The factor $e^{-\kappa}$ in Eq. (19) accounts for the washout of $\eta_L^{\rm max}$ during reheating due to the $\Delta L = 2$ washout processes, cf. the term proportional to $-\sigma_{\rm eff} n_L$ in Eq. (7). In case the axion begins to oscillate before the end of reheating, i.e., for $m_a \gtrsim \Gamma_\varphi$, one can estimate

$$\kappa \sim \frac{T_{\rm rh}}{T_L} \simeq 1 \left(\frac{T_{\rm rh}}{10^{13} \ {\rm GeV}}\right). \qquad (24)$$

For $m_a \lesssim \Gamma_\varphi$ on the other hand, washout is always negligible, so that κ can be safely set to $\kappa = 0$. A more careful treatment of the effect of washout on the final baryon asymmetry in our scenario is left for future work.

Finally, the factor C in Eq. (19) is a numerical fudge factor, which can, in principle, be estimated analytically, but which, in practice, is best determined by fitting η_L^a in Eq. (19) to the outcome of our numerical analysis. Specifically, we find $C \simeq 1.5$ for $m_a \lesssim \Gamma_\varphi$ and $C \simeq 2.2$ for $m_a \gtrsim \Gamma_\varphi$. The fact that these values are both of $\mathcal{O}(1)$ confirms the accuracy of our analytical estimate.

Altogether, the parameter dependence of the final lepton asymmetry in Eq. (19) can be summarized as follows (here, we neglect the effect of washout and set $\kappa \to 0$),

$$\eta_L^a \propto \frac{1}{T_L} \begin{cases} m_a^{-3/4}\Gamma_\varphi^{7/4}a_0 f_a^{-1} & ; m_a \gtrsim \Gamma_\varphi, \Delta_a = 1 \\ m_a a_0 f_a^{-1} & ; m_a \lesssim \Gamma_\varphi, \Delta_a = 1 \\ m_a^{3/4}\Gamma_\varphi^{5/4}M_{\rm Pl}^2 a_0^{-1} f_a^{-2} & ; m_a \gtrsim \Gamma_\varphi, \Delta_a > 1 \\ m_a^2 M_{\rm Pl}^2 a_0^{-1} f_a^{-2} & ; m_a \lesssim \Gamma_\varphi, \Delta_a > 1 \end{cases}.$$
$$(25)$$

Let us now determine the range of parameters that admit the correct value of the baryon asymmetry in view of the constraints in Eqs. (1), (3), (12), and (17), cf. Fig. 2, which shows the contour lines of successful leptogenesis for different values of f_a. The range of allowed values spans 5 orders of magnitude,

$$4 \times 10^{10} \text{ GeV} \lesssim f_a \lesssim 4 \times 10^{15} \text{ GeV}. \qquad (26)$$

For smaller values of f_a, it is not possible to generate a sufficiently large baron asymmetry, while keeping the baryonic isocurvature perturbations small enough. For larger values of f_a, the dilution of the asymmetry during the late-time decay of the axion is too strong. Varying f_a within the interval in Eq. (26), we then find that m_a, Γ_φ, and $T_{\rm rh}$ can take values within the following ranges:

$$1 \times 10^6 \text{ GeV} \lesssim m_a \lesssim 2 \times 10^{11} \text{ GeV}, \qquad (27)$$

$$3 \times 10^6 \text{ GeV} \lesssim \Gamma_\varphi \lesssim 3 \times 10^{11} \text{ GeV}, \qquad (28)$$

$$9 \times 10^{11} \text{ GeV} \lesssim T_{\rm rh} \lesssim 3 \times 10^{14} \text{ GeV}. \qquad (29)$$

These ranges of parameters are consistent with models of dynamical axions, as well as string axion models.

Finally, let us conclude. As we have been able to show, lepton number violation due to the exchange of heavy Majorana neutrinos, in combination with the effective chemical potential generated by a slowly rolling axionlike scalar field, is sufficient for a successful realization of baryogenesis via leptogenesis. In this scenario, the baryon asymmetry does neither depend on the concrete heavy neutrino mass spectrum nor on the amount of CP violation in the light and heavy neutrino sectors. In particular, it is consistent with (almost) degenerate heavy neutrino masses close to the GUT scale. Hence, while thermal leptogenesis assumes some of the heavy neutrino Yukawa couplings to be much smaller than $\mathcal{O}(1)$, our mechanism equally applies in the case of Yukawa couplings of $\mathcal{O}(10^{-1} \cdots 1)$. Furthermore, as can be seen from Eqs. (9) and (25), the baryon asymmetry increases with the light neutrino masses m_i. Thus, while thermal leptogenesis imposes an upper bound on the neutrino mass scale, $\bar{m} \lesssim 0.2$ eV, to avoid too strong washout [23], our scenario works for all experimentally allowed light neutrino masses. This bound will soon be probed by a multitude of terrestrial experiments [24] as well as in cosmological and astrophysical observations [25]. The axion-driven leptogenesis mechanism presented in this Letter therefore appears to be an attractive alternative to the conventional scenario of thermal leptogenesis. On the other hand, further work is clearly needed to embed it into a more complete model. Here, we expect our mechanism to yield a number of nontrivial model-building constraints due to some of its peculiar features: (i) the requirement of a high reheating temperature, (ii) the fact that f_a may potentially lie in the vicinity of the decay constant of the QCD axion, (iii) the possibility of baryonic isocurvature perturbations at a detectable level, etc. Such constraints may easily allow for a possibility to test our mechanism in the near future and assess whether it has indeed the potential to serve as a viable explanation for the baryon asymmetry of the Universe.

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