

Triplet Higgs field assisted leptogenesis from axion oscillation after inflation

Sasmita Mishra^{*}

*Department of Physics and Astronomy, National Institute of Technology Rourkela,
Sundargarh, Odisha, India, 769 008*



(Received 17 June 2025; accepted 17 October 2025; published 19 November 2025)

Leptogenesis via axion oscillation after inflation is an alternate mechanism of thermal leptogenesis. In this mechanism, the requirement of the existence of a lepton number (L) violating process in equilibrium to drive the lepton number requires the temperature of leptogenesis to be at $\sim 10^{13}$ GeV. Triplet scalars, due to their interaction with gauge bosons, make a suitable candidate to prevail in the thermal bath via gauge scattering at such a high energy. Also, owing to its interaction with Standard Model (SM) leptons and the Higgs scalar, it can mediate the $\Delta L = 2$ process. Moreover, just one triplet is enough to serve the purpose as opposed to thermal leptogenesis where at least one more triplet scalar/right-handed neutrino is required to generate a sizable CP violation. In this work, we study a model where the SM gauge group is extended with $U(1)_{PQ}$, with the addition of one Higgs doublet, one scalar triplet, and one complex scalar singlet. The presence of a complex scalar singlet decouples the Peccei-Quinn (PQ) symmetry breaking from the electroweak scale. It also provides a common source of an axionlike particle and seesaw scale. The scalar triplet offers a common link between leptogenesis via axion oscillation and neutrino mass. Further, with the presence of one triplet scalar, the lepton flavor violation process is directly determined from low-energy neutrino oscillation data.

DOI: [10.1103/y6f5-4yxh](https://doi.org/10.1103/y6f5-4yxh)

I. INTRODUCTION

The origin of the baryon asymmetry of the Universe (BAU) after an early epoch of cosmological inflation continues to be one of the most important questions of physics beyond the Standard Model (SM). The observations from cosmic microwave background (CMB) radiation and big bang nucleosynthesis give a measurement of the asymmetry [1] $Y_B = \frac{(n_B - n_{\bar{B}})}{s}$ as 8.7×10^{-11} , where n_B , $n_{\bar{B}}$, and $s = \frac{2\pi^2}{45} g_{*s} T^3$ denote the number densities of baryons, antibaryons, and entropy density of the Universe, respectively. Baryogenesis via leptogenesis [2] has become a well-known scenario to explain the matter-antimatter asymmetry of the Universe based on Sakharov's conditions [3]. The conditions provide the necessary ingredients for dynamical generation of BAU, viz, baryon (lepton) number $[B(L)]$ violation, C and CP violation, and out-of-equilibrium dynamics in the processes involving baryons. The well-known seesaw framework could explain tiny nonzero neutrino mass and the consequent leptogenesis leading to adequate baryon asymmetry. One such scenario is the SM

augmented with one scalar triplet where through a type-II seesaw mechanism [4–7], light neutrino masses can naturally arise from Weinberg's dimension-five operator. The unique feature of type-II seesaw-induced light neutrino mass is that the presence of only one Yukawa coupling matrix can completely be determined from the low-energy values of neutrino masses and mixing angles up to an overall factor and renormalization group equation (RGE) evolution effects [8]. With one triplet added to the SM, although it is possible to theoretically explain the light neutrino masses and mixing satisfying experimental data, thermal leptogenesis is not viable. It was shown in Ref. [9] that baryogenesis via heavy Higgs leptogenesis requires its mass $\geq 10^{13}$ GeV with the amount of CP violation $\sim 10^{-7}$. However, to generate CP violation, an $\sim 10^{-7}$ one triplet scalar is inadequate, which leads to the inclusion of another right-handed neutrino or triplet scalar. As a result, the opportunities to connect high- and low-scale dynamics is lost in a model. In order to preserve the intriguing connection between high- and low-scale dynamics, the authors in Refs. [10,11] explored a minimal extension of the Standard Model by introducing a single scalar triplet. This framework can simultaneously account for the baryon asymmetry of the Universe (BAU) and give rise to a Starobinsky-like inflationary epoch [12]. In this work, we show an alternate mechanism, where without losing the direct link between neutrino mass and flavor structure, leptogenesis can be achieved with the help of one triplet scalar using axion oscillation after inflation [13].

^{*}Contact author: mishras@nitrkl.ac.in

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

The quantum-chromodynamics (QCD) axion and/or axionlike-particles (ALPs) appear from a spontaneously broken global $U(1)_{\text{PQ}}$ symmetry as a pseudo-Nambu-Goldstone boson, the phase of a complex singlet scalar field. The simplest realizations of the Peccei-Quinn (PQ) mechanism is in the Weinberg-Wilczek (WW) model [14,15]. The WW model was soon ruled out by laboratory searches as the axion decay constant f_a is of the order of the electroweak scale $v_{\text{ew}} \simeq 246$ GeV not being sufficiently suppressed. This led to models in which the PQ symmetry breaking is decoupled from the electroweak scale via the introduction of an SM singlet scalar field σ acquiring a vacuum expectation value (VEV) f_a . In all string theories, axions and ALPs appear naturally as Kaluza-Klein zero modes of antisymmetric tensor fields when the extra six dimensions are compactified. Determined by the anomaly cancellation conditions, the couplings of these fields result in the anomalous CP violating couplings in a low-energy effective Lagrangian. Thus, axions and ALPs are a potential bridge connecting the low-energy effective theories like SM and high-energy theories like string theory, theories of extra dimensions, UV complete theories like grand unified theories (GUT), and more. ALPs have wide applications in theoretical physics. Some of them remain relatively light and play an important role in baryogenesis, the role of the inflaton, dark matter, and dark energy (for a review, see Ref. [16]). In this work, we explore one such model that contains ALP and additional Higgs scalars in the SM to address the issue of neutrino mass and baryogenesis. We intend to study the common origin of neutrino mass and baryogenesis through spontaneously broken PQ symmetry in a way that the PQ symmetry breaking scale and the seesaw scale arises from a common origin. A related study of the common origin of resolving the strong CP problem and the seesaw mechanism explaining the smallness of neutrino masses with spontaneously broken PQ symmetry was studied in Ref. [17].

Sakharov's conditions for attaining successful baryogenesis are based on CPT invariance. In an attempt by Cohen and Kaplan in their scenario of spontaneous baryogenesis [18], the authors have shown that when CPT is spontaneously broken, baryon asymmetry is generated in thermal equilibrium. In an attempt in a similar direction, it was shown in Ref. [19] that the PQ symmetry may be explicitly broken in the early Universe, inducing the rotation of the PQ-charged scalar field. The asymmetry of the PQ charge arising from rotation can be converted to baryon asymmetry via QCD and electroweak sphaleron transitions. In Ref. [20], baryogenesis was shown from the cosmological evolution of a CP -breaking modulus scalar either during the big bang or around the end of inflation with a condition that the modulus scalar evolves in time, while $\Delta L = 2$ rates related to neutrino masses go out of equilibrium at a decoupling temperature $T_{\text{dec}} \sim 10^{11}$ GeV.

In this work, we pursue the idea of the possibility of baryogenesis via leptogenesis by spontaneously breaking

CPT invariance by introducing an ALP, which couples derivatively to the fermion current j_μ in the Lagrangian. The derivative coupling gives rise to an effective chemical potential associated with the fermion number, leading to a relative shift in the energy levels of fermions and anti-fermions. This amounts to creating the difference between B and L satisfying one of Sakharov's conditions. After the end of inflation, the ALP starts to oscillate coherently with nonzero velocity, which temporarily induces an effective chemical potential for the fermion number. For baryogenesis, there must be an external source of baryon or lepton number violation in equilibrium. The reheating temperature in this scenario that turns out to be $\simeq 10^{13}$ GeV makes it a high-energy-scale phenomenon, making it an alternate mechanism of thermal leptogenesis. An inert Higgs doublet assisted lepton number violating operator with a much lower reheating temperature was studied in Ref. [21] via axion oscillation. In this work, we propose that the presence of a single triplet scalar is suitable enough to assist the $\Delta L = 2$ process and generate light neutrino mass via a type-II seesaw mechanism, making it a unified picture. While the involvement of gauge coupling in the case of triplet scalars ensures its presence in the thermal bath in the early Universe, the PQ symmetry breaking is decoupled from the electroweak scale via the introduction of an SM singlet scalar field σ acquiring a VEV f_a .

The paper is organized as follows: In Sec. II, we lay down the model where the SM is extended with $U(1)_{\text{PQ}}$ symmetry with the addition of one Higgs doublet, one complex scalar singlet, and one triplet scalar. In Sec. III, we discuss the triplet-scalar-assisted leptogenesis via axion oscillation. In Sec. IV, we discuss the lepton flavor violation by connecting the branching ratios with a light neutrino mass matrix through triplet Yukawa coupling. The conclusions are given in Sec. V.

II. THE MODEL

Although the presence and abundance of axions are linked with an effective low-energy consequence of string theory, the theory also provides strong evidence that exact global symmetries cannot survive in contact with quantum gravity. Also, obtaining an automatic $U(1)_{\text{PQ}}$ in GUT or connecting the axion and grand unification is a long-standing problem [22]. This is similar to the so-called PQ quality problem [23]: Why is $U(1)_{\text{PQ}}$ a good symmetry of UV physics? Addressing this issue, the simultaneous presence of the horizontal flavor group $SU(3)_f$ and vertical $SO(10)$ leads to an automatic $U(1)_{\text{PQ}}$ if the $SO(10)$ Higgs representations are properly chosen was shown in Ref. [23] with a global $SU(3)_f$ and local $SU(3)_f$ in Ref. [22]. Some related proposals without GUTs were carried out in Refs. [24–26].

In this section, we present a model motivated by the idea of the connection of massive neutrinos, often related to a high-energy-scale seesaw mechanism, with the presence of

spontaneously broken PQ symmetry [27]. For UV completion of the theory, we choose Peccei and Quinn symmetry through $U(1)_{\text{PQ}}$ symmetry. The gauge group of the model is the SM gauge group extended by $U(1)_{\text{PQ}}$ with the addition of a complex scalar field σ and a Higgs doublet. The addition of σ decouples the PQ symmetry breaking from the electroweak scale. The particle content of the model is given in Table I (Sec. II) with the corresponding charges. The determination of PQ charges is summarized in Appendix A. TABLE I. The particle content of the model.

Field	Spin	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_{\text{PQ}}$
q_L	$\frac{1}{2}$	3	2	$+\frac{1}{6}$	0
u_R	$\frac{1}{2}$	3	1	$+\frac{2}{3}$	X_u
d_R	$\frac{1}{2}$	3	1	$-\frac{1}{3}$	X_d
l_L	$\frac{1}{2}$	1	2	$-\frac{1}{2}$	X_l
e_R	$\frac{1}{2}$	1	1	-1	X_e
H_u	0	1	2	$-\frac{1}{2}$	$-X_u$
H_d	0	1	2	$+\frac{1}{2}$	$-X_d$
Δ	0	1	3	+1	X_Δ
σ	0	1	1	0	X_σ

The Lagrangian for the fermion sector is given by

$$-\mathcal{L} = Y_u \bar{q}_L u_R H_u + Y_d \bar{q}_L d_R H_d + Y_e \bar{l}_L e_R H_d + \frac{1}{2} Y_\Delta l_L^T C i \tau_2 \Delta l_L + \text{H.c.}, \quad (1)$$

where the triplet scalar Δ can be represented as

$$\Delta \equiv \frac{\vec{\tau} \cdot \vec{\Delta}}{\sqrt{2}} = \begin{pmatrix} \frac{\Delta^+}{\sqrt{2}} & \Delta^{++} \\ \Delta^0 & -\frac{\Delta^+}{\sqrt{2}} \end{pmatrix}, \quad (2)$$

where $\vec{\Delta} = (\Delta_1, \Delta_2, \Delta_3)$ are $SU(2)_L$ components of the scalar triplet. The scalar potential is given by

$$V = V(\sigma) + V(\Delta) + V(H_u, H_d) + V(\sigma, H_u, H_d) + V(\sigma, \Delta) + V(\Delta, H_u, H_d) + V(\sigma, \Delta, H_u, H_d), \quad (3)$$

where the individual terms are given as

$$V(\sigma) = \lambda \left(\sigma^* \sigma - \frac{f_a^2}{2} \right)^2,$$

$$V(\Delta) = \mu_\Delta^2 \text{Tr}(\Delta^\dagger \Delta) + \lambda_{\Delta 4} \text{Tr}(\Delta^\dagger \Delta) \text{Tr}(\Delta^\dagger \Delta) + \lambda_9 \text{Tr}(\Delta^\dagger \Delta)^2,$$

$$V(H_u, H_d) = -\mu_1^2 |H_u|^2 + \lambda_1 |H_u|^4 - \mu_2^2 |H_d|^2 + \lambda_2 |H_d|^4 + \lambda_{12} |H_u|^2 |H_d|^2 + \lambda_4 |H_u^\dagger H_d|^2,$$

$$V(\sigma, H_u, H_d) = \lambda_{13} |\sigma|^2 |H_u|^2 + \lambda_{23} |\sigma|^2 |H_d|^2 + (\lambda_5 \sigma^2 \tilde{H}_u^\dagger H_d + \text{H.c.}),$$

$$V(\sigma, \Delta) = \lambda_{\Delta 3} |\sigma|^2 \text{Tr}(\Delta^\dagger \Delta),$$

$$V(\Delta, H_u, H_d) = \text{Tr}(\Delta^\dagger \Delta) [\mu_{\Delta 1} |H_u|^2 + \mu_{\Delta 2} |H_d|^2] + \lambda_7 H_u^\dagger \Delta \Delta^\dagger H_u + \lambda_8 H_d^\dagger \Delta \Delta^\dagger H_d + \lambda_9 (\Delta^\dagger \Delta)^2,$$

$$V(\sigma, \Delta, H_u, H_d) = \lambda_6 \sigma H_u^\dagger \Delta^\dagger H_d + \text{H.c.},$$

where $\tilde{H}_u = i \tau_2 H_u^*$. The simultaneous presence of Y_Δ and λ_6 ensures explicit breaking of the lepton number.

The vacuum configurations of the PQ-charged complex singlet is given by

$$\sigma = \frac{1}{\sqrt{2}} (S + f_a) e^{i\phi_a/f_a}. \quad (4)$$

At the first stage of symmetry breaking, the radial component of the scalar singlet σ acquires a VEV f_σ thereby breaking the PQ symmetry. The phase of the complex scalar field $\theta = \phi_a/f_a$, where the pseudo-Nambu-Goldstone boson ϕ_a is identified as the ALP. The massive radial mode is settled to the minimum and ignored below the scale f_a . Assuming that the PQ symmetry is broken before the end of inflation, the initial value of the ALP $\phi_{a0} = \theta_0 f_a$, where the value of the angle θ_0 takes a random value in the range $\theta_0 \in [0, 2\pi)$. So, the initial value of ϕ_{a0} results in it being constant during the superhorizon scale.

The potential of the ALP can originate from instanton effects in a strongly coupled hidden sector characterized by a dynamical scale Λ_H and is given by

$$V(\theta) = \Lambda_H^4 (1 - \cos \theta). \quad (5)$$

The potential can be approximated to $V(\theta) \simeq \frac{1}{2} m_a^2 \phi_a^2$, where the ALP mass $m_a \simeq \Lambda_H^2/f_a$. After the PQ symmetry breaking, the potential can be written as

$$V_1 = V(\theta) + V_{H_u, H_d} + \lambda_{13} f_a^2 |H_u|^2 + \lambda_{23} f_a^2 |H_d|^2 + \text{Tr}(\Delta^\dagger \Delta) [\mu_\Delta^2 + \mu_{\Delta 1} |H_u|^2 + \mu_{\Delta 2} |H_d|^2 + \lambda_{\Delta 3} f_a^2 + \lambda_{\Delta 4} \text{Tr}(\Delta^\dagger \Delta)] + \lambda_7 H_u^\dagger \Delta \Delta^\dagger H_u + \lambda_8 H_d^\dagger \Delta \Delta^\dagger H_d + \lambda_9 (\Delta^\dagger \Delta)^2 + (\lambda_5 f_a \tilde{H}_u^\dagger H_d + \lambda_6 f_a H_u^\dagger \Delta^\dagger H_d + \text{H.c.}). \quad (6)$$

After the first stage of $U(1)_{\text{PQ}}$ symmetry breaking, when the electroweak and triplet scalar VEVs are not switched on ($\langle H_{u,d} \rangle = 0$), the triplet scalar Δ acquires a bare mass μ_Δ as well as a mass induced by the VEV of the complex singlet scalar charged under $U(1)_{\text{PQ}}$. The model now can be treated as a two-Higgs-doublet model extended with a real triplet scalar; see Ref. [28]. In this case, the scalar spectrum can be calculated and is summarized in Appendix B. With electroweak and triplet scalar VEVs switched on, the scalar spectrum has been calculated in Ref. [27].

A. Neutrino mass

After the first stage of symmetry breaking, at a later epoch the remaining scalar fields attain the vacuum configuration given by

$$\begin{aligned} H_u &\supset \frac{v_u + h_u}{\sqrt{2}} \begin{pmatrix} 1 \\ 0 \end{pmatrix}, & H_d &\supset \frac{v_d + h_d}{\sqrt{2}} \begin{pmatrix} 0 \\ 1 \end{pmatrix}, \\ \Delta &\supset \frac{v_\Delta + \delta^0}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ 1 & 0 \end{pmatrix}, \end{aligned} \quad (7)$$

where v_u , v_d , and v_Δ denote the corresponding VEVs. The light neutrino mass matrix from Eq. (1) can be written as

$$m_\nu = Y_\Delta v_\Delta. \quad (8)$$

Using the VEVs of H_u , H_d , and Δ from Eq. (7), the part of the potential dependent on v_Δ is given as

$$\begin{aligned} \langle V_\Delta \rangle &= (\mu_\Delta^2 + \lambda_{\Delta 3} f_a^2 + \lambda_{\Delta 1} v_u^2 + (\lambda_{\Delta 2} + \lambda_8) v_d^2) v_\Delta^2 \\ &\quad + 2\lambda_6 f_a v_u v_d v_\Delta + \mathcal{O}(v_\Delta^4) + \text{terms independent of } v_\Delta. \end{aligned}$$

For a hierarchy $f_a \gg v_{u,d} \gg v_\Delta$, the minimization of the scalar potential is given by

$$2M_\Delta^2 v_\Delta + 2\lambda_6 f_a v_u v_d \approx 0, \quad (9)$$

$$M_\Delta^2 = \mu_\Delta^2 + \lambda_{\Delta 3} f_a^2 + \lambda_{\Delta 1} v_u^2 + (\lambda_{\Delta 2} + \lambda_{\Delta 8}) v_d^2. \quad (10)$$

Writing $\mu = \lambda_6 f_a$ and $v_u v_d \approx v_{\text{ew}}^2$, Eq. (8) can be written as

$$m_\nu = Y_\Delta v_\Delta \approx \frac{Y_\Delta \mu v_{\text{ew}}^2}{M_\Delta^2}, \quad v_\Delta \approx \frac{\lambda_6 f_a v_u v_d}{M_\Delta^2}. \quad (11)$$

The VEV v_Δ has the following range of allowed values:

$$\mathcal{O}(1) \text{ GeV} > v_\Delta \gtrsim 0.05 \text{ eV}, \quad (12)$$

where the upper bound is derived from the ρ -parameter constraints determined from precision measurements [29], and the lower limit ensures the generation of the observed neutrino masses while also requiring perturbative Yukawa couplings. In the limit $\mu \approx M_\Delta$ and $v_{\text{ew}} \sim \mathcal{O}(100) \text{ GeV}$, the mass of the triplet scalar is constrained as

$$10^4 \text{ GeV} \leq M_\Delta \leq 10^{13} \text{ GeV}. \quad (13)$$

The above bound is obtained by considering that there should be at least one neutrino mass at least of the order of $m_\nu \sim 0.05 \text{ eV}$. By diagonalizing the mass matrix $m_\nu = Y_\Delta v_\Delta$ with a diagonalizing matrix U_L ,

$$U_L^T m_\nu U_L = m_{\text{Diag}}, \quad (14)$$

the eigenvalues of the light neutrino masses m_i , $i = 1, 2, 3$ can be obtained. In the present work, we have taken U_L as the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix U_{PMNS} times a diagonal matrix involving Majorana phases $\text{Diag}(1, e^{i\alpha}, e^{i\beta})$. The structure of the U_{PMNS} matrix can be formed using three rotation matrices with corresponding rotating angles θ_{ij} : $i, j = 1, 2, 3$ and a Dirac CP violating phase δ_{CP} . The PMNS matrix for the light neutrino sector is

$$U_{\text{PMNS}} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta_{CP}} \\ -s_{12}c_{23} - c_{12}s_{13}s_{23}e^{i\delta_{CP}} & c_{12}c_{23} - s_{12}s_{13}s_{23}e^{i\delta_{CP}} & c_{13}s_{23} \\ s_{12}s_{23} - c_{12}s_{13}c_{23}e^{i\delta_{CP}} & -c_{12}s_{23} - s_{12}s_{13}c_{23}e^{i\delta_{CP}} & c_{13}c_{23} \end{pmatrix}, \quad (15)$$

where $c_{ij} = \cos(\theta_{ij})$, $s_{ij} = \sin(\theta_{ij})$.

III. LEPTOGENESIS FROM AXION OSCILLATION

The axion is the angular direction of the complex scalar field σ . At the first stage of symmetry breaking with the radial component of the scalar singlet σ , a VEV develops, f_a , and breaks the PQ symmetry. The massive radial mode quickly settles to its minimum and can be integrated out

below the scale f_a . Assuming that the PQ symmetry is broken before the end of inflation, the initial value of the axion at the end of inflation $\phi_{a0} = \theta_0 f_a$, where the value of the angle θ_0 takes a random value in the range $\theta_0 \in [0, 2\pi)$. So, the initial value of ϕ_{a0} results in being constant during the superhorizon scale. The evolution of the axion field in the early Universe is governed by the classical field equation

$$\ddot{\phi}_a + 3H\dot{\phi}_a = -\partial_{\phi_a} V(\theta), \quad (16)$$

where $V(\theta)$ is given in Eq. (5). The axion mass $m_a \simeq \Lambda_H^2/f_a$ can originate from instanton effects in a strongly coupled hidden sector characterized by a dynamical scale Λ_H . For compatibility, m_a should be smaller than the scale of inflation H_{inf} , $m_a \lesssim H_{\text{inf}}$. After the end of inflation, the axion field remains at rest until the Hubble parameter drops to $H \simeq m_a$. Once the axion field starts to move, the effective Lagrangian is given by

$$\mathcal{L} \supset \frac{1}{f_a} \partial_\mu \phi_a j^\mu; \quad j^\mu = \sum_f \bar{\psi}_f \gamma^\mu \psi_f, \quad (17)$$

where ψ_f is the SM lepton (l_i) and quark (q_i) doublets and right-handed singlets of different flavors. Integrating Eq. (17) by parts,

$$\mathcal{L} \supset -\frac{1}{f_a} \phi_a \partial_\mu j^\mu \rightarrow \frac{\phi_a}{f_a} \frac{N_g}{8\pi^2} (g_2^2 W_{\alpha\beta} \tilde{W}^{\alpha\beta} - g_1^2 B_{\alpha\beta} \tilde{B}^{\alpha\beta}), \quad (18)$$

where N_g is the number of SM fermion generations. The electroweak field strength tensors are represented as $W_{\mu\nu}$ and $B_{\mu\nu}$. The gauge couplings of the electroweak sector are represented as g_1 and g_2 . In Eq. (18), the anomaly equation has been used. Considering the axion field is spatially homogeneous $\phi_a = \phi_a(t)$, the coupling presented in Eq. (18) transforms into an effective chemical potential μ_{eff} for the fermion number. The latter can be identified by performing an integration by parts in Eq. (18),

$$\mathcal{L} \supset \frac{\dot{\phi}_a}{f_a} j^0 = \mu_{\text{eff}} n_F, \quad (19)$$

where $\mu_{\text{eff}} = \dot{\phi}_a/f_a$ and $j^0 = n_F = n_f - n_{\bar{f}}$. The nonzero effective chemical potential μ_{eff} creates a shift in the energy levels of particles and antiparticles. In thermal equilibrium, the minimum of the free energy is obtained for a nonzero fermion-antifermion asymmetry n_F ,

$$n_{f,\bar{f}}^{\text{eq}} \sim T^3 \left(1 \pm \frac{\mu_{\text{eff}}}{T} \right), \quad n_F^{\text{eq}} = n_f^{\text{eq}} - n_{\bar{f}}^{\text{eq}} \sim \mu_{\text{eff}} T^2. \quad (20)$$

The equilibrium number density of the equilibrium number density can be calculated taking the quantum-statistical effects for counting relativistic degrees of freedom in the classical Boltzmann approximations, so the necessary condition for generating a lepton asymmetry is satisfied [13]. The mechanism does not depend on CP violating phases or the heavy neutrino mass spectrum, unlike thermal leptogenesis. Lepton number violation is mediated by the $\Delta L = 2$ scattering processes. The scalar triplets generate the $LL \leftrightarrow \bar{H}\bar{H}$ density rate γ_{T_s} by s -channel exchange and the $LH \leftrightarrow \bar{L}\bar{H}$ density rate γ_{T_t} by t -channel exchange. Here, H is a generic Higgs scalar representing H_u/H_d . The process $LL \leftrightarrow HH$ is in thermal equilibrium

$$\Gamma_L|_{T=M_\Delta} = n_l^{\text{eq}} \sigma_{\text{eff}} > H|_{T=M_\Delta}, \quad n_l^{\text{eq}} = \frac{2}{\pi^2} T^3, \quad (21)$$

where $\sigma_{\text{eff}} = \langle \sigma_{\Delta L=2} v \rangle$ is the thermally averaged cross section of the $\Delta L = 2$ process, and

$$H|_{T=M_\Delta} = \sqrt{\frac{\pi^2 g_*}{90}} \frac{M_\Delta^2}{M_{\text{Pl}}}. \quad (22)$$

Identifying

$$\Gamma_L|_{T=M_\Delta} = n_l^{\text{eq}} \langle \sigma_{\Delta L=2} v \rangle \approx \frac{Y_\Delta^2 \mu^2}{M_\Delta}, \quad (23)$$

and using the expression for light neutrino mass in Eq. (11) $m_\nu \simeq Y_\Delta \frac{|\mu| v_{\text{ew}}^2}{2M_\Delta^2}$, the limit on the triplet Higgs mass results in

$$M_\Delta > 10^{12} \text{ GeV}. \quad (24)$$

The other lepton number violating processes $LL \leftrightarrow \Delta$ and $HH \leftrightarrow \Delta$ must not coexist, otherwise the lepton number asymmetry will be rapidly washed out. If both processes are in equilibrium, the solutions corresponding to chemical potentials are $\mu_L = \mu_{u_L} = \mu_H = 0$. To maintain the lepton asymmetry in the plasma, the process $LL \leftrightarrow \Delta$ must be efficient, while the process $HH \leftrightarrow \Delta$ is out of equilibrium. The latter condition can be represented as

$$\Gamma_{\text{ID}}(HH \leftrightarrow \Delta)|_{T=M_\Delta} < H|_{T=M_\Delta}, \quad (25)$$

where

$$\Gamma_{\text{ID}}(HH \leftrightarrow \Delta)|_{T=M_\Delta} \approx \Gamma_D(\Delta \rightarrow HH) \simeq \frac{\mu^2}{32\pi M_\Delta}. \quad (26)$$

Using $v_\Delta \sim \frac{\mu v_{\text{ew}}^2}{2M_\Delta^2}$, the condition in Eq. (25) on v_Δ translates to

$$v_\Delta \lesssim \mathcal{O}(1) \text{ GeV}. \quad (27)$$

In Refs. [10,11], it was shown that the combination of the triplet Higgs field and the SM Higgs field can simultaneously give rise to a Starobinsky-like inflationary epoch [12]. Efforts to use an axion with a shift symmetry to ensure a prolonged slow-roll background evolution was studied in [30]. Primordial inflation is produced by a combination of σ (nonminimally coupled to the scalar curvature), and the SM Higgs boson was explored in [31]. In the framework of natural inflation [32], Nambu-Goldstone fields like the axion can drive the inflation, but the results from Planck 2018 + BICEP/Keck (BK15) Polarization + BAO data strongly disfavor the framework [33]. The alternate mechanism of the axion as an inflaton candidate has been studied in Refs. [34,35]. A model with chaotic inflation where the

role of the inflaton is played by the Higgs triplet in a type-II seesaw mechanism for generating the small masses of left-handed neutrinos was studied in Ref. [36]. In this work, we assume, during inflation, that the axion remains a spectator that is responsible for creating a chemical potential required for leptogenesis, and the inflation is driven by another inflaton field.

The requirement that the inflaton must not decay before the end of inflation implies

$$\Gamma_\phi \leq H_{\text{inf}}, \quad (28)$$

where Γ_ϕ is the decay width of the inflaton. Planck measurements of the CMB anisotropies, combining information from the temperature and polarization maps and the lensing reconstruction, put a constraint on the scale of inflation as reference $\frac{H_{\text{inf}}}{M_{\text{Pl}}} < 2.5 \times 10^{-5}$ (95% CL) [1]. Baryon asymmetry produced during reheating, in our case, is dependent on the initial axion field value ϕ_{a0} . The axion field acquires quantum fluctuations during inflation by an amount $H_{\text{inf}}/2\pi$. The baryonic isocurvature fluctuation puts a constraint on $H_{\text{inf}}/2\pi\phi_{a0} \ll 5 \times 10^{-5}$ [37]. The constraint results in $H_{\text{inf}} \leq 6 \times 10^{11} \text{ GeV} (\frac{f_a}{10^{15} \text{ GeV}})$. We take into account these bounds on H_{inf} in our numerical analysis for estimating the BAU from the leptogenesis via axion oscillation in the next subsection.

A. Numerical analysis and results

The successful implementation of triplet assisted leptogenesis requires that the triplet scalar mediated $\Delta L = 2$ interactions must be in thermal equilibrium before the onset of axion oscillation

$$\Gamma_L \gg H \gtrsim m_a, \quad (29)$$

where

$$\Gamma_L = n_l^{\text{eq}} \sigma_{\text{eff}}, \quad n_l^{\text{eq}} = \frac{2}{\pi^2} T^3, \quad (30)$$

where the equilibrium number density n_l^{eq} has been calculated by taking the quantum-statistical effects for counting relativistic degrees of freedom. Now using the conditions given in Eq. (30), the rough temperature of the leptogenesis and axion mass scale can be obtained as

$$T \sim T_L = \frac{\sqrt{g_*}}{\pi \sigma_{\text{eff}} M_{\text{Pl}}} \sim 10^{13} \text{ GeV},$$

$$m_a \sim \sigma_{\text{eff}} T_L^3 \sim 10^8 \text{ GeV}, \quad (31)$$

where

$$\sigma_{\text{eff}} \approx \frac{3}{32\pi} \frac{\bar{m}^2}{v_{\text{ew}}^4} \simeq 1 \times 10^{-31} \text{ GeV}^{-2},$$

$$\bar{m}^2 = \sum_{i=1}^3 \bar{m}_i^2 \approx \Delta m_{\text{atm}}^2 \simeq 2.4 \times 10^{-3} \text{ eV}^2. \quad (32)$$

The dependence of σ_{eff} on light neutrino mass can be understood from the calculation of Feynman diagrams representing $\Delta L = 1$ and $\Delta L = 2$ processes using the Feynman rules for the dimension-five operator $\kappa \propto m_\nu/v_{\text{ew}}^2$ [38,39]. The solution to the temperature is obtained by using the energy density due to radiation

$$\rho_R = \frac{\pi^2 g_*}{3} T^4. \quad (33)$$

After the end of inflation, the temperature rises to the maximal value

$$T_{\text{max}} = 5 \times 10^{13} \text{ GeV} \left(\frac{\Gamma_\phi}{10^9 \text{ GeV}} \right)^{1/4} \left(\frac{H_{\text{inf}}}{10^{11} \text{ GeV}} \right)^{1/2}. \quad (34)$$

After that, due to inflaton oscillation, the energy density scales as matter, and hence the temperature decreases. During reheating, the temperature drops as $T \propto R^{-3/8}$ until the radiation comes to dominate the energy density of the Universe at time $t = t_{\text{rh}} \simeq \Gamma_\phi^{-1}$ when $\rho_R = \rho_\phi$ (ρ_ϕ is the energy density of the inflaton) where the reheating temperature is ($T_{\text{rh}} \simeq 0.3 \sqrt{\Gamma_\phi M_{\text{Pl}}}$),

$$T_{\text{rh}} \simeq 2 \times 10^{13} \text{ GeV} \left(\frac{\Gamma_\phi}{10^9 \text{ GeV}} \right)^{1/2}. \quad (35)$$

To get the final lepton asymmetry and hence the BAU, the following equations must be solved simultaneously:

$$\ddot{\phi}_a + 3H\dot{\phi}_a = -\frac{\partial V(\phi_a)}{\partial \phi_a}, \quad (36)$$

$$\dot{\rho}_\phi + 3H\rho_\phi = -\Gamma_\phi \rho_\phi, \quad (37)$$

$$\dot{\rho}_R + 4H\rho_R = \Gamma_\phi \rho_\phi, \quad (38)$$

$$H^2 = \frac{\rho_{\text{tot}}}{3M_{\text{Pl}}^2}, \quad (39)$$

$$\rho_{\text{tot}} \approx \rho_\phi + \rho_R, \quad (40)$$

$$\dot{n}_L + 3Hn_L = -\Gamma_L(n_L - n_L^{\text{eq}}). \quad (41)$$

Taking the quantum-statistical effects for counting the relativistic degrees of freedom, the equilibrium number density in the classical Boltzmann approximations is

$$n_L^{\text{eq}} = \frac{4}{\pi^2} \mu_{\text{eff}} T^2. \quad (42)$$

The lepton asymmetry is partly converted to baryon asymmetry by the electroweak sphaleron processes. The final BAU of η_B can be calculated as

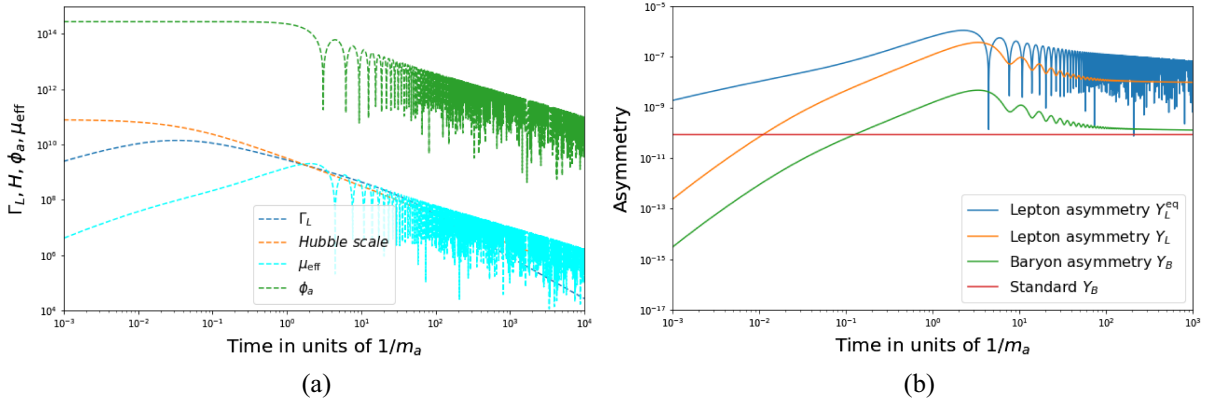


FIG. 1. The plots represent the evolution of various quantities of interest as a function of the time in units of $1/m_a$. (a) The evolution of H , Γ_L , ϕ_a , and μ_{eff} , showing the evolution of the Hubble parameter H , the rate of $\Delta L = 2$ process Γ_L , and the evolution of the ALP ϕ_a and μ_{eff} . (b) Evolution of lepton and baryon asymmetry, showing the evolution of the equilibrium value of the lepton asymmetry $Y_L^{\text{eq}} = n_L^{\text{eq}}/s$, the lepton asymmetry $Y_L = n_L/s$, and the final baryon asymmetry $Y_B = n_B/s$. The red line is the reference line for the observed value of $Y_B = 8.75 \times 10^{-11}$. The numerical evolution of the quantities has been done for the benchmark values: $f_a = 3 \times 10^{14}$ GeV, $m_a = 5 \times 10^9$ GeV, $\Gamma_\phi = 3 \times 10^8$ GeV, and $H_{\text{inf}} = 8 \times 10^{10}$ GeV.

$$\eta_B = c_{\text{sph}} \eta_L, \quad \eta_L = \frac{n_L}{s}, \quad (43)$$

where s is the entropy density of the Universe and $c_{\text{sph}} = \frac{28}{79}$. The plots in Fig. 1 represent the evolution of various quantities of interest as a function of the time in units of $1/m_a$. The numerical evolution of the quantities has been done by solving Boltzmann equations (36)–(41) for the benchmark values $f_a = 3 \times 10^{14}$ GeV, $m_a = 5 \times 10^9$ GeV, $\Gamma_\phi = 3 \times 10^8$ GeV, and $H_{\text{inf}} = 8 \times 10^{10}$ GeV. In Fig. 1(a), the evolution of the Hubble parameter H , the rate of the $\Delta L = 2$ process Γ_L , the evolution of the ALP ϕ_a , and μ_{eff} are shown in orange, blue, red, and cyan, respectively. Since the PQ symmetry is broken before the end of inflation, after N e -folds during inflation, the ALP field ϕ_a attenuates as $\phi_{a,\text{end}} \simeq \phi_{a,\text{ini}} e^{-\lambda N}$, where $\lambda = \frac{3}{2} \left(1 - \sqrt{1 - \frac{4m_a^2}{9H_{\text{inf}}^2}} \right)$. So, depending upon the number of e -folds $N \sim 50$ – 60 , and using the values of H_{inf} and m_a , at the end of inflation, the ALP field value ends up being $\phi_{a,\text{end}} \sim 0.9\phi_{a,\text{ini}}$. After the end of inflation, the analytical solution of the ALP field equation in Eq. (36) in a radiation-dominated Universe is given by

$$\phi_a(t) \simeq A \phi_{a,\text{ini}} (m_a t)^{-1/4} J_{1/4}(m_a t), \quad (44)$$

where A is a constant that can be fixed from initial conditions, and $J_{1/4}$ is the Bessel function of the first kind. For $m_a t \ll 1$, the ALP field is effectively frozen by Hubble friction and $\phi_a(t) \simeq \text{constant}$. For $m_a t \gg 1$, the ALP field behaves as $\phi_a(t) \simeq t^{-3/4} \cos(m_a t + \delta)$. δ is a phase factor; i.e., it oscillates with an amplitude decreasing as $\phi_a(t) \sim t^{-3/4}$. The numerical evolution of $\phi_a(t)$ is shown in green in Fig. 1(a). Once $\phi_a(t)$ is known, $\dot{\phi}_a(t)$ can be calculated to find $\mu_{\text{eff}} = \dot{\phi}_a/\phi_a$, and is shown in cyan in Fig. 1(a).

In Fig. 1(b), the evolution of the equilibrium value of lepton asymmetry $Y_L^{\text{eq}} = n_L^{\text{eq}}/s$, lepton asymmetry $Y_L = n_L/s$, and final baryon asymmetry $Y_B = n_B/s$ are shown. The red line is the reference line for the observed value of $Y_B = 8.75 \times 10^{-11}$. The initial value of ALP at the end of inflation and its velocity are taken to be $\phi_a \simeq 0.9f_a$ and $\dot{\phi}_a = 0$. Since the equilibrium lepton number density Y_L^{eq} depends on $\mu_{\text{eff}} = \dot{\phi}_a/\phi_a$, it tracks the evolution of $\dot{\phi}_a$. Once Γ_L , the $\Delta L = 2$ triplet scalar mediated scattering catches up being in equilibrium the lepton asymmetry Y_L , and hence, the baryon asymmetry Y_B starts to grow, and as the Γ_L process goes out of equilibrium, the lepton asymmetry freezes. Finally, due to the sphaleron process, the lepton asymmetry gets converted to baryon asymmetry $Y_B \sim 8.7 \times 10^{11}$.

In our numerical analysis, we find the final lepton asymmetry is highly sensitive to Γ_L through its dependence on the thermally averaged cross section $\sigma_{\text{eff}}/\langle \sigma_{\Delta L=2} v \rangle$. As per Eq. (32), the latter is determined partly from the light neutrino mass m_ν . In the model considered in the present work, the light neutrino mass matrix is generated through the type-II seesaw mechanism. The direct link between m_ν and the triplet Yukawa coupling Y_Δ is now directly linked with the high-energy-scale phenomenon of leptogenesis. In the next section, we discuss the low-energy-scale connection of Y_Δ in the lepton flavor violation processes.

IV. LEPTON FLAVOR VIOLATION

The generation of light neutrino masses and mixing generated through the seesaw mechanism can induce lepton flavor violation (LFV) effects in a charged lepton sector like $l_i \rightarrow l_j \gamma$. The LFV effects are suppressed in models where the SM is extended with heavy particles responsible for the seesaw mechanism. The processes can be sizeable in

the minimal supersymmetric standard model embedded in GUTs where they can be enhanced through the one-loop exchange of sleptons (and gauginos). The LFV effects manifest from the Yukawa sector of the superpotential and the RGE effects that introduce off-diagonal elements of the slepton mass matrix. The supersymmetric version of the model considered in this work and the LFV effects can be compared with the model as studied in Ref. [8].

The existence of LFV effects in the slepton sector can result in enhanced LFV processes in the charged lepton sector. For example, the radiative decays $l_i \rightarrow l_j + \gamma$ can be estimated as

$$\text{BR}(l_i \rightarrow l_j + \gamma) \simeq \frac{\alpha^3}{G_F^2} \frac{|(m_L^2)_{ji}|^2}{m_S^8} \tan^2 \beta \text{BR}(l_j \rightarrow l_i \nu_j \bar{\nu}_i), \quad (45)$$

where α and G_F denote the fine structure and Fermi constant, respectively, and m_S is a typical supersymmetry mass particle inside the loops. It is useful to work with the quantities

$$R_{\tau\mu} \equiv \frac{\text{BR}(\tau \rightarrow \mu\gamma)}{\text{BR}(\mu \rightarrow e\gamma)} \simeq \left| \frac{(m_L^2)_{\tau\mu}}{(m_L^2)_{\mu e}} \right|^2 \frac{\text{BR}(\tau \rightarrow \mu\nu_\tau \bar{\nu}_\mu)}{\text{BR}(\mu \rightarrow e\nu_\mu \bar{\nu}_e)}, \quad (46)$$

$$R_{\tau e} \equiv \frac{\text{BR}(\tau \rightarrow e\gamma)}{\text{BR}(\mu \rightarrow e\gamma)} \simeq \left| \frac{(m_L^2)_{\tau e}}{(m_L^2)_{\mu e}} \right|^2 \frac{\text{BR}(\tau \rightarrow e\nu_\tau \bar{\nu}_e)}{\text{BR}(\mu \rightarrow e\nu_\mu \bar{\nu}_e)}. \quad (47)$$

Between the GUT scale M_{GUT} and the seesaw scale M_Δ the RGE induced slepton mass matrix inherits the LFV effect as

$$(m_L^2)_{\tau e} \simeq -\frac{9m_0^2 + 3a_0^2}{8\pi^2} (Y_\Delta^\dagger Y_\Delta)_{ij} \ln \frac{M_{\text{GUT}}}{M_\Delta}, \quad (48)$$

where m_0 and A_0 are the universal soft mass for scalars and universal trilinear coupling defined at the GUT scale. The striking feature of the type-II seesaw mechanism allows us to replace $(Y_\Delta^\dagger Y_\Delta)_{ij}$ by $(m_\nu^\dagger m_\nu)_{ij}$, and hence the link between LFV effects and low-energy neutrino parameters becomes direct. Using the ratios $\frac{\text{BR}(\tau \rightarrow \mu\nu_\tau \bar{\nu}_\mu)}{\text{BR}(\mu \rightarrow e\nu_\mu \bar{\nu}_e)} = 0.1737$ and $\frac{\text{BR}(\tau \rightarrow e\nu_\tau \bar{\nu}_e)}{\text{BR}(\mu \rightarrow e\nu_\mu \bar{\nu}_e)} = 0.1784$, the quantities of interest now can be written as

$$R_{\tau\mu} \simeq 0.1737 \times \left| \frac{(m_\nu^\dagger m_\nu)_{\tau\mu}}{(m_\nu^\dagger m_\nu)_{\mu e}} \right|^2, \quad (49)$$

$$R_{\tau e} \simeq 0.1784 \times \left| \frac{(m_\nu^\dagger m_\nu)_{\tau e}}{(m_\nu^\dagger m_\nu)_{\mu e}} \right|^2. \quad (50)$$

V. CONCLUSION

In this work, we have studied a model where the presence of a single triplet scalar accounts for light neutrino

mass and baryogenesis via leptogenesis. In contrast to the standard thermal leptogenesis, the initial lepton asymmetry is generated from the CPT violating effect from the ALP which is sourced from the spontaneous breaking of $U(1)_{\text{PQ}}$. The scale of PQ symmetry breaking is decoupled from the electroweak symmetry breaking at a scale $\sim 10^{14}$ GeV. The PQ symmetry is broken before the end of inflation, and the initial field value of ALP is uniform. The classical evolution of the ALP induces a chemical potential for the fermion number due to the derivative coupling of the ALP with the electroweak field strength tensors. The requirement that there be an external source of baryon or lepton number violation process in equilibrium is met by the presence of a single triplet scalar is suitable enough to assist the $\Delta L = 2$ process and generate light neutrino mass via the type-II seesaw mechanism. This requires the reheat temperature be $\simeq 10^{13}$ GeV, making it a high-energy-scale phenomenon. Unlike the case of thermal leptogenesis, where one needs at least two scalar triplets or one each of a scalar triplet and right-handed singlet, a single triplet is enough to generate light neutrino mass and baryogenesis via leptogenesis. Moreover, the $U(1)_{\text{PQ}}$ extended SM provides a common link between neutrino mass generation and leptogenesis. We find the production of final lepton asymmetry is highly sensitive to light neutrino mass m_ν . The proportionality relation between m_ν and the triplet Yukawa coupling Y_Δ makes it now directly linked with the high-energy-scale phenomenon of leptogenesis. With this setup again determining that the flavor structure of the model from the low-energy-scale data is preserved, the model provides an excellent link between low-energy- and high-energy-scale phenomena. Cosmology allows us to verify physical processes occurring at such high mass scales. For example, Ref. [40] has shown that if a coupling exists between the SM Higgs field and the ALP arising from the breaking of a global $U(1)$ symmetry, it can induce a strong first-order phase transition, thereby generating stochastic gravitational waves that could be detectable by current and future gravitational-wave observatories.

ACKNOWLEDGMENTS

S. M. thanks the anonymous referee for their critical comments and suggestions.

DATA AVAILABILITY

No data were created or analyzed in this study.

APPENDIX A: PQ CHARGES

The invariance of the Lagrangian and the scalar potential requires the constraints on the PQ charges as

$$-X_l + X - e - X_d = 0, \quad (\text{A1})$$

$$2X_l + X_\Delta = 0, \quad (\text{A2})$$

$$2X_\Sigma - X_u - X_d = 0, \quad (\text{A3})$$

$$X_\sigma + X_u - X_\Delta - X_d = 0. \quad (\text{A4})$$

The above equations can be solved in terms of X_u and X_d to get

$$\begin{aligned} X_l &= -\frac{3X_u}{4} + \frac{X-d}{4}, & X_e &= -\frac{3X_u}{4} + \frac{5X_d}{4}, \\ X_\Delta &= \frac{3X_u}{2} - \frac{X_d}{2}, & X_\sigma &= \frac{X_u}{2} + \frac{X_d}{2}. \end{aligned} \quad (\text{A5})$$

The orthogonality of the hypercharge and axion currents requires

$$X_u^2 v_u^2 = X_d^2 v_d^2. \quad (\text{A6})$$

Using the conditions in Eqs. (A5) and (A6), all PQ charges can be determined up to an overall renormalization. By choosing the renormalization condition

$$X_\sigma = 1, \quad (\text{A7})$$

and defining $x \equiv v_u/v_d$, the PQ charges are given by

$$\begin{aligned} X_u &= \frac{2}{x^2 + 1}, & X-d &= \frac{2x^2}{x^2 + 1}, & X_l &= \frac{x^2 - 3}{2(x^2 + 1)}, \\ X_e &= \frac{5x^2 - 3}{2(x^2 + 1)}, & X_\Delta &= \frac{3 - x^2}{x^2 + 1}. \end{aligned} \quad (\text{A8})$$

APPENDIX B: SCALAR SPECTRUM

With the electroweak VEVs switched off,

$$\frac{\partial V}{\partial f_a} = 2f_a(2\lambda_3 f_a^2 - \mu_3^2). \quad (\text{B1})$$

Using Eq. (B1) in expanding the scalar potential up to second order in the dynamical fields and with the unbroken SM symmetry, the scalar spectrum is given by:
a real scalar SM singlet σ^0 ,

$$M_{\sigma^0}^2 = 4\lambda_3 f_a^2, \quad (\text{B2})$$

a real pseudoscalar SM singlet as a part of η_σ^0

$$M_{\eta_\sigma^0}^2 = 0, \quad (\text{B3})$$

corresponding to the zero-mass mode of the PQ-breaking field corresponding to the axion,
a complex triplet with mass

$$M_\Delta^2 = \lambda_{\Delta 3} f_a^2 + \mu_\Delta^2, \quad (\text{B4})$$

and the complex doublets H_u and H_d with the mass matrix in the basis (H_u^*, H_u) ,

$$M_H^2 = \begin{pmatrix} \lambda_{13} f_a^2 - \mu_1^2 & \lambda_5 f_a^2 \\ \lambda_5^2 f_a^2 & \lambda_{23} f_a^2 - \mu_2^2 \end{pmatrix}. \quad (\text{B5})$$

Diagonalizing the above mass matrix, the corresponding eigenvalues are given as

$$\begin{aligned} 2M_{u,d}^2 &= (\lambda_{13} + \lambda_{23})f_a^2 - \mu_1^2 - \mu_2^2 \\ &\pm \sqrt{((\lambda_{13} - \lambda_{23})f_a^2 - \mu_1^2 + \mu_2^2)^2 + 4\lambda_5^2 f_a^4}. \end{aligned} \quad (\text{B6})$$

-
- [1] N. Aghanim *et al.*, Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020); **652**, C4(E) (2021).
 - [2] M. Fukugita and T. Yanagida, Baryogenesis without grand unification, *Phys. Lett. B* **174**, 45 (1986).
 - [3] A. D. Sakharov, Violation of CP invariance, C asymmetry, and baryon asymmetry of the universe, *Pis'ma Zh. Eksp. Teor. Fiz.* **5**, 32 (1967).
 - [4] George Lazarides, Q. Shafi, and C. Wetterich, Proton lifetime and fermion masses in an $SO(10)$ model, *Nucl. Phys. B* **181**, 287 (1981).
 - [5] M. Magg and C. Wetterich, Neutrino mass problem and gauge hierarchy, *Phys. Lett.* **94B**, 61 (1980).
 - [6] Rabindra N. Mohapatra and Goran Senjanovic, Neutrino masses and mixings in gauge models with spontaneous parity violation, *Phys. Rev. D* **23**, 165 (1981).
 - [7] J. Schechter and J. W. F. Valle, Neutrino masses in $SU(2) \times U(1)$ theories, *Phys. Rev. D* **22**, 2227 (1980).
 - [8] Anna Rossi, Supersymmetric seesaw without singlet neutrinos: Neutrino masses and lepton flavor violation, *Phys. Rev. D* **66**, 075003 (2002).
 - [9] James N. Fry, Keith A. Olive, and Michael S. Turner, Higgs bosons and the evolution of baryon asymmetries, *Phys. Rev. D* **22**, 2977 (1980).
 - [10] Neil D. Barrie, Chengcheng Han, and Hitoshi Murayama, Affleck-Dine leptogenesis from Higgs inflation, *Phys. Rev. Lett.* **128**, 141801 (2022).
 - [11] Neil D. Barrie, Chengcheng Han, and Hitoshi Murayama, Type II Seesaw leptogenesis, *J. High Energy Phys.* **05** (2022) 160.
 - [12] Alexei A. Starobinsky, A new type of isotropic cosmological models without singularity, *Phys. Lett.* **91B**, 99 (1980).

- [13] Alexander Kusenko, Kai Schmitz, and Tsutomu T. Yanagida, Leptogenesis via axion oscillations after inflation, *Phys. Rev. Lett.* **115**, 011302 (2015).
- [14] Steven Weinberg, A new light boson?, *Phys. Rev. Lett.* **40**, 223 (1978).
- [15] Frank Wilczek, Problem of strong P and T invariance in the presence of instantons, *Phys. Rev. Lett.* **40**, 279 (1978).
- [16] Luca Di Luzio, Maurizio Giannotti, Enrico Nardi, and Luca Visinelli, The landscape of QCD axion models, *Phys. Rep.* **870**, 1 (2020).
- [17] Y. H. Ahn and Eung Jin Chun, Minimal models for axion and neutrino, *Phys. Lett. B* **752**, 333 (2016).
- [18] Andrew G. Cohen and David B. Kaplan, Spontaneous baryogenesis, *Nucl. Phys.* **B308**, 913 (1988).
- [19] Raymond T. Co and Keisuke Harigaya, Axiogenesis, *Phys. Rev. Lett.* **124**, 111602 (2020).
- [20] Mateusz Duch, Alessandro Strumia, and Arsenii Titov, Baryogenesis from cosmological CP breaking, *arXiv*: 2504.03506.
- [21] Arghyajit Datta, Soumen Kumar Manna, and Arunansu Sil, Spontaneous leptogenesis with sub-GeV axionlike particles, *Phys. Rev. D* **110**, 095035 (2024).
- [22] Luca Di Luzio, Accidental SO(10) axion from gauged flavour, *J. High Energy Phys.* **11** (2020) 074.
- [23] Howard M. Georgi, Lawrence J. Hall, and Mark B. Wise, Grand unified models with an automatic Peccei-Quinn symmetry, *Nucl. Phys.* **B192**, 409 (1981).
- [24] Darwin Chang, Palash B. Pal, and Goran Senjanovic, Axions from chiral family symmetry, *Phys. Lett.* **153B**, 407 (1985).
- [25] Joerg Jaeckel and Andreas Ringwald, The low-energy frontier of particle physics, *Annu. Rev. Nucl. Part. Sci.* **60**, 405 (2010).
- [26] Palash B. Pal, The strong CP question in $SU(3)(C) \times SU(3)(L) \times U(1)(N)$ models, *Phys. Rev. D* **52**, 1659 (1995).
- [27] Stefano Bertolini, Luca Di Luzio, Helena Kolečová, and Michal Malinský, Massive neutrinos and invisible axion minimally connected, *Phys. Rev. D* **91**, 055014 (2015).
- [28] B. Ait-Ouazghour and M. Chabab, The Higgs potential in 2HDM extended with a real triplet scalar: A roadmap, *Int. J. Mod. Phys. A* **36**, 2150131 (2021).
- [29] Shinya Kanemura and Kei Yagyu, Radiative corrections to electroweak parameters in the Higgs triplet model and implication with the recent Higgs boson searches, *Phys. Rev. D* **85**, 115009 (2012).
- [30] Enrico Pajer and Marco Peloso, A review of axion inflation in the era of Planck, *Classical Quantum Gravity* **30**, 214002 (2013).
- [31] Guillermo Ballesteros, Javier Redondo, Andreas Ringwald, and Carlos Tamarit, Unifying inflation with the axion, dark matter, baryogenesis and the seesaw mechanism, *Phys. Rev. Lett.* **118**, 071802 (2017).
- [32] Katherine Freese, Joshua A. Frieman, and Angela V. Olinto, Natural inflation with pseudo—Nambu-Goldstone bosons, *Phys. Rev. Lett.* **65**, 3233 (1990).
- [33] Nina K. Stein and William H. Kinney, Natural inflation after Planck 2018, *J. Cosmol. Astropart. Phys.* **01** (2022) 022.
- [34] Guido D’Amico, Nemanja Kaloper, and Albion Lawrence, Monodromy inflation in the strong coupling regime of the effective field theory, *Phys. Rev. Lett.* **121**, 091301 (2018).
- [35] Jihn E. Kim, Hans Peter Nilles, and Marco Peloso, Completing natural inflation, *J. Cosmol. Astropart. Phys.* **01** (2005) 005.
- [36] Chian-Shu Chen and Chia-Min Lin, Type II Seesaw Higgs triplet as the inflaton for chaotic inflation and leptogenesis, *Phys. Lett. B* **695**, 9 (2011).
- [37] Keisuke Harigaya, Ayuki Kamada, Masahiro Kawasaki, Kyohei Mukaida, and Masaki Yamada, Affleck-Dine baryogenesis and dark matter production after high-scale inflation, *Phys. Rev. D* **90**, 043510 (2014).
- [38] Yosef Nir, Introduction to leptogenesis, in *6th Rencontres du Vietnam: Challenges in Particle Astrophysics* (2007), 2, *arXiv*:hep-ph/0702199.
- [39] Alessandro Strumia, Baryogenesis via leptogenesis, in *Les Houches Summer School on Theoretical Physics: Session 84: Particle Physics Beyond the Standard Model* (2006), p. 655, *arXiv*:hep-ph/0608347.
- [40] P. S. Bhupal Dev, Francesc Ferrer, Yiyang Zhang, and Yongchao Zhang, Gravitational waves from first-order phase transition in a simple axion-like particle model, *J. Cosmol. Astropart. Phys.* **11** (2019) 006.