$\mathrm{SU}(2)_R$ and its axion in cosmology: A common origin for inflation, cold sterile neutrinos, and baryogenesis

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We introduce an axion-inflation model embedded in the left-right symmetric extension of the Standard Model in which W_R is coupled to the axion. This model merges three milestones of modern cosmology, i.e., inflation, cold dark matter, and baryon asymmetry. Thus, it can naturally explain the observed coincidences among cosmological parameters, i.e., $\eta_{\mathsf{B}} \approx P_\zeta$ and $\Omega_{\mathrm{DM}} \simeq 5\Omega_{\mathsf{B}}$. The source of asymmetry is spontaneous CP violation in the physics of inflation, and the lightest right-handed neutrino is the cold dark matter candidate with mass $m_{N_1} \sim 1$ GeV. The introduced mechanism does not rely on the largeness of the unconstrained CP-violating phases in the neutrino sector or fine-tuned masses for the heaviest right-handed neutrinos. It has two unknown fundamental scales, i.e., scale of inflation $\Lambda_{\mathrm{inf}} = \sqrt{HM_{\mathrm{Pl}}}$ and left-right symmetry breaking Λ_F. Sufficient matter asymmetry demands that $\Lambda_{\mathrm{inf}} \approx \Lambda_F$. Baryon asymmetry and dark matter today are remnants of a pure quantum effect (chiral anomaly) in inflation, which, thanks to flavor effects, has been memorized by cosmic evolution.

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The two pillars of the postinflationary scenarios of leptogenesis are (i) CP asymmetric decay of massive right-handed neutrinos (RHNs) after reheating, and (ii) washout processes to enhance the efficiency and eliminate the preexisting asymmetry to avoid theoretical uncertainties [1]. In addition, the lightest sterile neutrino may account for the dark matter (DM) [2]. The source of CP asymmetry is the CP-violating phases in the neutrino sector, unconstrained by the current data, which are assumed to be large enough. Moreover, low-scale leptogenesis mostly requires a fine-tuning of parameters, e.g., highly degenerate RHN masses such as $m_{N_3} \simeq m_{N_2}$ [3]. Also, in inflation models that produce matter symmetry, flavor effects make it difficult for the preexisting asymmetry to be washed out by the RH neutrino decays [4] (see also Fig. S3 of the Supplemental Material [5]).

This paper introduces a new framework for simultaneous baryogenesis and dark matter production within general relativity (GR), which avoids the above issues. The source of asymmetry is spontaneous CP violation by a W_R gauge field coupled to the inflaton that produces leptons and baryons in inflation. In this scenario, baryon asymmetry and DM are remnants of the same effect in inflation. Thus, it can naturally explain the observed coincidences among cosmological parameters, i.e., $\eta_{\mathsf{B}} \approx P_\zeta$ and $\Omega_{\mathrm{DM}} \simeq 5\Omega_{\mathsf{B}}$.

Early Universe physics is a subject that seeks answers for fundamental questions linking very high energy physics (immensely small scales) with extremely large cosmological scales. The Standard Model (SM) of particle physics, which has been highly successful in formulating fundamental particles at low energy scales, is greatly incomplete when it meets cosmology and astrophysics. The most glaring shortcomings of the SM are (I) the neutrino mass, (II) baryon asymmetry of the Universe (BAU), and (III) the particle nature of dark matter. Considering cosmic inflation as the leading paradigm for early Universe [6], we should add (IV) the particle nature of the inflaton field to this list. SM as a theory for particle physics also faces a number of issues, i.e., (i) the Higgs vacuum stability problem, (ii) accidental $\mathsf{B}-\mathsf{L}$ global symmetry, and (iii) *ad hoc* parity violation at the electroweak scale (EW).

Axion fields are suitable candidates for inflaton and are naturally coupled to gauge fields. As first discovered by the author, non-Abelian gauge fields may survive inflation and contribute to its physics while respecting the cosmological symmetries [7]. This discovery introduced a new class of inflation models accompanied by SU(2) gauge fields with an immensely rich phenomenology. The minimal realization of this idea is an SU(2) gauge field in GR coupled to a generic axion inflaton [7–9], so-called SU(2)-axion inflation. For a review on gauge fields in inflation, see [10]. The novel features shared by these models are (1) spontaneous P and CP violation satisfying all Sakharov conditions in inflation [11], (2) particle production by the gauge field in

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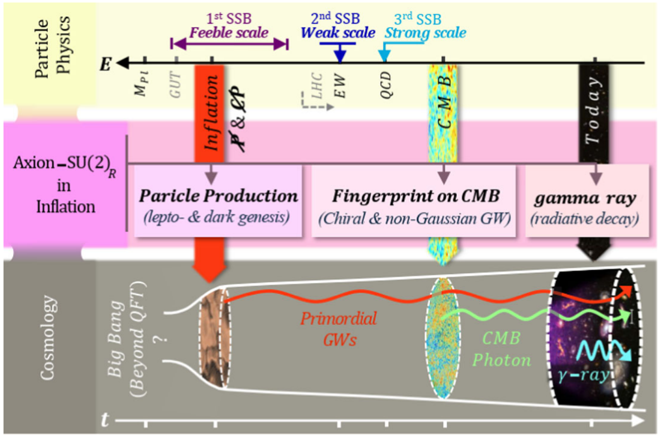

FIG. 1. $\mathrm{SU}(2)_R$-axion inflation: a natural common origin for inflation, fermionic dark matter, and matter asymmetry and its observational signatures.

inflation through the Schwinger effect [12–14] and the chiral anomaly [14], (3) naturally warm inflation [15], (4) and the prediction that a chiral and non-Gaussian gravitational wave (GW) background [10,16,17] will be detectable by future cosmic microwave background missions and laser interferometers [18]. Therefore, this inflation setup merges three open issues of the SM and cosmology, i.e., inflation, DM, and BAU. It gains additional value due to its unique observable signature on the GW background induced by GW-SU(2) field interactions. The connection of this SU(2) field to the SM, however, was still missing, and in this work we attempt to fill the gap.

The aim of this paper is to embed the SU(2)-axion inflation setting in gauge extensions of the SM and study its phenomenological and cosmological consequences. The most suitable beyond the SM theories are supersymmetry, grand unified theory (GUT), and left-right symmetric models (LRSM) of the weak interactions. We restrict this work to the most minimal realization, i.e., the LRSM. Originally proposed to explain P violation in low energy processes [19], the LRSM predicted massive neutrinos years before experimentation. Among its appealing features are natural $\mathsf{B-L}$ symmetry [20], entailed seesaw mechanisms [21], and use as a solution to vacuum stability problem at high scales [22].

In this paper, we assume that the gauge field in the SU(2)-axion inflation models is W_R in the LRSM. In a comparison with the minimal LRSM, here we have an axion φ, which is coupled to the $\mathrm{SU}(2)_R$ and drives cosmic inflation. We call this particle physics model for inflation *$\mathrm{SU}(2)_R$-axion inflation*. This new framework can simultaneously provide plausible explanations for the previously mentioned phenomena (I)–(IV) and (i)–(iii) (See Fig. 1). A more detailed analysis is presented in a followup work [23]. This paper can be a starting point for a further, more involved analysis of the rich and multifaceted phenomenology of these gauge extensions of the SM in inflation physics.

I. $\mathrm{SU}(2)_R$-AXION INFLATION MODEL

The minimal gauge group that implements the hypothesis of left-right symmetry is $\mathcal{G} \equiv \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{\mathsf{B-L}}$ (suppressing color). The theory includes three gauge fields $W_{L,R}$ and $A_{\mathsf{B-L}}$, which are associated with $\mathrm{SU}(2)_{L,R}$ and $\mathrm{U}(1)_{\mathsf{B-L}}$, respectively. The fermionic content is the SM quarks and leptons extended by three RHNs as

$$q_{iL,R} = \begin{pmatrix} \boldsymbol{u}_i \\ \boldsymbol{d}_i \end{pmatrix}_{L,R} \qquad \text{and} \qquad l_{iL,R} = \begin{pmatrix} \boldsymbol{\nu}_i \\ \boldsymbol{l}_i \end{pmatrix}_{L,R}, \tag{1}$$

where $\boldsymbol{\nu}_{iR}$ are three RHNs interacting by $\mathrm{SU}(2)_R$. It is accompanied by an extended Higgs sector consists of a Higgs bidoublet $\boldsymbol{\Phi}$ and $\mathrm{SU}(2)_{L,R}$ triplets $\boldsymbol{\Delta}_{L,R}$. The spontaneous symmetry breaking (SSB) structure of the LRSM is

$$\mathcal{G} \underset{\text{1st SSB}}{\overset{T<\Lambda_F}{\longrightarrow}} \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \underset{\text{2nd SSS}}{\overset{T<\Lambda_W}{\longrightarrow}} \mathrm{U}(1)_{em}.$$

Below the scale Λ_F, the first SSB happens, which breaks the LR symmetry and gives a vacuum expectation value (VEV) to the $\mathrm{SU}(2)_R$ triplet, i.e., $\langle \boldsymbol{\Delta}_R \rangle \neq 0$. At this point, $W_R^{\pm}, Z_R$, and $\mathbf{N}_i \equiv \boldsymbol{\nu}_i + \boldsymbol{\nu}_i^c$ become massive. Next, when the temperature gets below the EW scale, $T < \Lambda_W$, the Higgs bidoublet acquires a VEV, i.e., $\langle \boldsymbol{\Phi} \rangle \neq 0$, and the second SSB occurs, which gives the Dirac mass to the SM particles, active neutrinos included [21].

Now we add the inflaton, i.e., the axion field, which is coupled to the $\boldsymbol{W}_R$. As a concrete example we consider

$$S_{\rm inf} = \int d^4x \sqrt{-g} \left[-\frac{1}{2} (\partial_\mu \varphi)^2 - V(\varphi) + \mathcal{L}_{W_R} \right], \tag{2}$$

$$\mathcal{L}_{W_R} = -\frac{1}{2} \mathrm{Tr}[\boldsymbol{W}_{\mu\nu} \boldsymbol{W}^{\mu\nu}]_R - \frac{\lambda\varphi}{f} \mathrm{Tr}[\boldsymbol{W}_{\mu\nu} \tilde{\boldsymbol{W}}^{\mu\nu}]_R, \tag{3}$$

where $\boldsymbol{W}_{R\mu\nu}$ is the strength tensor of $\boldsymbol{W}_R$, $\tilde{\boldsymbol{W}}_R^{\mu\nu} \equiv \frac{1}{2} \frac{\epsilon^{\mu\nu\lambda\sigma}}{\sqrt{-g}} \boldsymbol{W}_{R\lambda\sigma}$, $f \lesssim 10^{-1} M_{\rm Pl}$, and $\lambda \lesssim 1$. For the sake of generality, we assume that $V(\varphi)$ is an arbitrary axion potential that is flat enough to support the slow-roll inflation, for instance, an axion monodromy inspired potential form [24]. This $\mathrm{SU}(2)$-axion inflation model and its cosmic perturbations were studied in [9]. $\mathrm{SU}(2)_R$-axion inflation has two unknown fundamental scales, i.e., the scale of inflation $\Lambda_{\rm inf} = \sqrt{M_{\rm Pl} H}$, and LR symmetry breaking Λ_F. Moreover, $\boldsymbol{W}_R$ may or may not have a VEV in inflation. Thus, we can distinguish among four different types of scenarios.

	$\Lambda_{\rm inf} > \Lambda_F$	$\Lambda_{\rm inf} < \Lambda_F$
$\langle \boldsymbol{W}_R \rangle = 0$	**I**	**II**
$\langle \boldsymbol{W}_R \rangle \neq 0$	$\mathbf{I_v}$	$\mathbf{II_v}$
Mass in Inflation	$m_{W_R} = 0$ $m_{N_i} = 0$	$m_{W_R} \neq 0$ $m_{N_i} \neq 0$

Scenarios **I** and $\mathbf{I}_v$ describe the case $\Lambda_{\rm inf} > \Lambda_F$, while **II** and $\mathbf{II}_v$ apply when $\Lambda_{\rm inf} < \Lambda_F$. The v subscript denotes systems in which the $\mathrm{SU}(2)_R$ acquires a VEV in inflation. In this work, we focus on scenarios **I** and **II**, $\langle \boldsymbol{W}_R\rangle = 0$, and leave cases $\mathbf{I}_v$ and $\mathbf{II}_v$ for future work. The RH fermions are coupled to the $\boldsymbol{W}_R$ field and its axion as

$$\mathcal{L}_{\Psi_{JR}} \supset \bar{\Psi}_{JR}\left(\boldsymbol{\sigma}^\mu[iD_\mu + g_R \boldsymbol{W}_{R\mu}] - \frac{\tilde{\lambda}\dot{\varphi}}{f}\right)\Psi_{JR}, \qquad (4)$$

where $\tilde{\lambda}$ is a constant, D_μ is the spinor covariant derivative, and RH fermions are collectively shown as

$$\Psi_{JR} = \{q_{iR}, l_{iR}\}, \quad \text{where } (J = 1, \ldots, 6). \qquad (5)$$

II. PARTICLE PRODUCTION IN INFLATION

Because of conformal symmetry, the gauge fields associated with the $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_{\mathsf{B-L}}$ group, as well as all the left-handed fermions, are exponentially decaying in inflation. However, the $\boldsymbol{W}_R$ associated with $\mathrm{SU}(2)_R$ is coupled to the inflaton and sourced by it. Subsequently, the generated $\mathrm{SU}(2)_R$ gauge field produces RH fermions. Note that the axion cannot create Weyl fermions [25] and that they are produced merely by W_R. The main particle physical consequences of this setup as the inflation physics are as follows: (i) P and C are maximally broken by the chiral nature of the $\mathrm{SU}(2)_R$ interaction with the axion, (ii) CP, B, and L are all violated by the nonperturbative effects of the $\boldsymbol{W}_R$, i.e., chiral (Adler-Bell-Jackiw), anomaly [26], (iii) $\mathsf{B-L}$ is conserved (violated) in scenario type **I** (type **II**), while $\mathsf{B-L}_{\rm SM}$ is violated in both scenarios, and (iv) out of thermal equilibrium conditions hold during inflation. Thus, all the Sakharov conditions required for BAU [27] are satisfied in inflation. The field equation of $\boldsymbol{W}_R$ is

$$(\partial_\mu - i g_R \boldsymbol{W}_{R\mu})\left[\boldsymbol{W}_R^{\mu\nu} + \frac{\lambda\varphi}{f}\tilde{\boldsymbol{W}}_R^{\mu\nu}\right] - m_{W_R}^2 \boldsymbol{W}_R^\nu = 0, \quad (6)$$

where g_R is the gauge coupling of $\boldsymbol{W}_R$. A massless gauge field (type **I**) with momentum $\mathbf{k}$ has two (transverse) polarization states specified by the polarization vectors $e^{\pm}(\mathbf{k})$ where $\mathbf{k}.e^{\pm}(\mathbf{k}) = 0$. The massive gauge field (type **II**) has an extra (longitudinal) mode with polarization vector $e^3(\mathbf{k}) = \mathbf{k}/k$ and its zero component coupled to it given by the constraint equation. Interestingly, the longitudinal mode and the zero component are decoupled from the axion and decay in inflation. Thus, we neglect them and refer the interested reader to [23] for detailed calculations. Now we define the following slowly increasing parameters:

$$\xi \equiv \frac{\lambda\dot{\varphi}}{2fH} \quad \text{and} \quad \tilde{\xi} \equiv \frac{\tilde{\lambda}}{\lambda}\xi. \qquad (7)$$

Imposing the Bunch-Davies vacuum, the transverse modes $f_\pm^a$ associated with $e^{\pm}(\mathbf{k})$ polarization states are

$$f_\pm^a(\mathbf{k},\tau) = \frac{e^{i\kappa_\pm \pi/2}}{(2\pi)^{\frac{3}{2}}\sqrt{2k}} W_{\kappa_\pm,\mu}(2ik\tau), \qquad (8)$$

where τ is the conformal time, $W_{\kappa_\pm,\mu}$ is the W-Whittaker function, $\kappa_\pm = \mp i\xi$, and $\mu^2 = \frac{1}{4} - \frac{m_{W_R}^2}{H^2}$. The energy density of W_R is

$$\langle \rho_{W_R}\rangle \simeq \left(\frac{H}{M_{\rm Pl}}\right)^2 \bar{\rho}\mathcal{T}(\xi, m_{W_R}), \qquad (9)$$

where $\bar{\rho} = 3M_{\rm Pl}^2 H^2$ and $\mathcal{T}(\xi, m_{W_R})$ is a function of ξ and m_{W_R}, which are shown in Eq. (S1) of the Supplemental Material [5]. For $\xi > m_{W_R}$, $\mathcal{T}$ increases (decreases) exponentially with the increase of ξ (m_{W_R}) as $\mathcal{T}(\xi, m_{W_R}) \propto \frac{1}{(2\pi)^2} e^{2(\xi - |\mu|)\pi}$, while for $\xi < m_{W_R}$, it has power-law behavior and softly increases with the increase of m_{W_R} (see Fig. S1 of the Supplemental Material [5]). During slow roll, ξ is an almost constant (gradually increasing) parameter. As a result, the axion slowly injects more and more energy into the gauge field sector, and inflation is warm.

The generated gauge boson field produces RH leptons and baryons in inflation. The anomalies of baryon and lepton currents are given, respectively, as

$$\nabla_\mu J_{\mathsf{B}}^{\mu R} = -\frac{3g_R^2}{16\pi^2}\mathrm{Tr}[\boldsymbol{W}^{\mu\nu}\tilde{\boldsymbol{W}}_{\mu\nu}]_R, \qquad (10)$$

$$\nabla_\mu J_{\mathsf{L}}^{\mu R} = -\frac{3g_R^2}{16\pi^2}\mathrm{Tr}[\boldsymbol{W}^{\mu\nu}\tilde{\boldsymbol{W}}_{\mu\nu}]_R + 2im_{N_i}\bar{\boldsymbol{\nu}}_{iR}\boldsymbol{\nu}_{iR}. \qquad (11)$$

However, the B- and L-violating interactions of the left-handed fermions remain negligible in inflation. The total lepton number is related to the SM one as

$$n_{\mathsf{L}} = n_{\underline{\mathsf{L}}} + \sum_i n_{N_i}, \qquad (\underline{\mathsf{L}} \equiv \mathsf{L}_{SM}), \qquad (12)$$

where n_{N_i} are the sterile neutrino lepton numbers. Using Eq. (8) in Eq. (10), we find that the baryon and lepton numbers are, respectively,

$$n_{\mathsf{B}} \simeq -g_R^2 H^3 \mathcal{K}(\xi, m_{W_R}), \qquad (13)$$

$$n_{\mathsf{L}} \simeq -H^3\left[g_R^2\mathcal{K}(\xi, m_{W_R}) + \sum_i \frac{\tilde{\xi}}{\pi}\left(\frac{m_{N_i}}{H}\right)^2 \mathcal{D}(\tilde{\xi}, m_{N_i})\right], \quad (14)$$

where $\mathcal{K}(\xi, m_{W_R})$ is the contribution of the chiral anomaly (a pure quantum effect) and $\mathcal{D}(\tilde{\xi}, m_{N_i})$ is the contribution of the mass term of the RHNs (in type **II** scenarios).

The explicit forms of these prefactors are presented in Eqs. (S2) and (S3) of the Supplemental Material [5], and their plots are shown in Fig. S2 of the Supplemental Material [5]. The prefactor $\mathcal{K}(\xi, m_{W_R})$ increases (decreases) with the increase of ξ (m_{W_R}) and for $\xi > M_{W_R}$ as

$$\mathcal{K}(\xi, m_{W_R}) \propto \frac{1}{(2\pi)^4} e^{2\xi\pi}. \tag{15}$$

The prefactor $\mathcal{D}(\tilde{\xi}, m_{N_i})$ is of order 1 and is symmetric with respect to $\tilde{\xi}$. The net $\mathsf{B} - \underline{\mathsf{L}}$ asymmetry ($\underline{\mathsf{L}} \equiv \mathsf{L}_{\rm SM}$) created by inflation is

$$n^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}} = \sum_i n^{\rm inf}_{N_i} \simeq -\alpha^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}} H^3,$$

$$\alpha^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}} \equiv \left[\frac{g_R^2}{2}\mathcal{K}(\xi, m_{W_R}) + \sum_i \frac{\tilde{\xi}}{\pi}\left(\frac{m_{N_i}}{H}\right)^2 \mathcal{D}(\tilde{\xi}, m_{N_i})\right]. \tag{16}$$

III. EVOLUTION AFTER REHEATING

The study of postinflationary evolution requires us to specify our parameter space further. For the sake of concreteness, we restrict the current analysis by assuming the following conditions: (C1) a hierarchical mass spectrum for the RH neutrinos (as implied by the neutrino oscillations) given as

$$m_{N_3} \gtrsim 10^{12}\ \text{GeV} \gg m_{N_2} \gtrsim 10^9\ \text{GeV} \gg m_{N_1}, \tag{17}$$

where $\mathbf{N}_1$ is much lighter than the EW scale and with feeble Yukawa interactions, i.e., a DM candidate. (C2) The *CP*-violating phases in the neutrino sector, unconstrained by the current data, are not enough to create the observed BAU. (C3) The postinflationary generation of RHNs with W_R interactions via freeze-out and freeze-in is negligible compared to their preexisting relics.

A. Memory effect and remnant asymmetries

Spectator effects have important impacts in the final values of B, $\underline{\mathsf{L}}$, and L_{N_1} (DM relic density). These effects are discussed in Sec. S2 of the Supplemental Material [5]. The final B and $\underline{\mathsf{L}}$ ($\equiv \mathsf{L}_{\rm SM}$) are (See Fig. 2)

$$n_{\mathsf{B}}(a) = 0.12\, n^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}}\left(\frac{a_{\rm inf}}{a}\right)^3 \quad \text{and} \quad n_{\underline{\mathsf{L}}}(a) = -\frac{7}{4} n_{\mathsf{B}}(a).$$

$\mathbf{N}_{3,2}$ decay to lighter particles, while $\mathbf{N}_1$ freezes out soon after inflation. Because of its feeble Yukawa interactions, it can account for the DM with a relic number density as

$$n_{N_1}(a) = -\frac{g_R^2 H^3}{6}\mathcal{K}(\xi, m_{W_R})\left(\frac{a_{\rm inf}}{a}\right)^3. \tag{18}$$

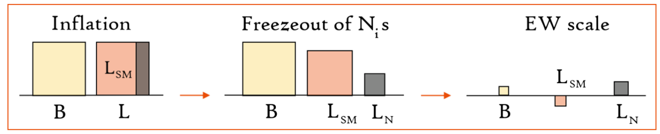

FIG. 2. Evolution of B, $\underline{\mathsf{L}}$, and $\mathbf{N}_1$ at three stages: (left panel) inflation, (middle panel) freeze-out of N_i's, and (right panel) EW scale.

B. Photon number density

Reheating starts at some point after the end of inflation. Here, we consider the phenomenological reheating model $\rho_{\rm reh} = \varepsilon(\frac{a_{\rm inf}}{a_{\rm reh}})^4 \rho_{\rm inf}$ in which ε is the efficiency of the reheating process and relates $\rho_{\rm reh}$ to the energy density at the end of inflation, $\rho_{\rm inf}$ (see Sec. S3 of the Supplemental Material [5]). Thus, the photon number density today ($g_{\rm eff,0} = \frac{43}{11}$) is

$$n_{\gamma,0} = \frac{6\sqrt{3}\zeta(3)}{\pi^2}\left(\frac{\varepsilon}{g_{\rm eff,0}}\right)^{\frac{3}{4}} (HM_{\rm Pl})^{\frac{3}{2}} \left(\frac{a_{\rm inf}}{a_0}\right)^3. \tag{19}$$

Note that, owing to condition (C3), the entropy injection by the decay of heavier RHNs after reheating is negligible.

C. Demands imposed by (C3)

Negligible freeze-out and freeze-in production of RHNs by W_R interactions requires

$$\frac{H}{M_{\rm Pl}} \lesssim \mathcal{A} \times 10^{-2} \varepsilon^{-\frac{1}{2}} \left(\frac{a_{\rm reh}}{a_{\rm inf}}\right)^2 \left(\frac{1}{g_R}\frac{m_{W_R}}{M_{\rm Pl}}\right)^{\frac{8}{3}}, \tag{20}$$

where $\mathcal{A} = \frac{1}{2}(\frac{g_{\rm eff}}{10^2})^{\frac{1}{2}}$ is of the order 1. (C3) imposes an upper bound on the scale of inflation. For details about the freeze-in production in this setup, see [23].

IV. BARYON-TO-PHOTON RATIO

Today we have

$$\eta_{\mathsf{B},0} \equiv \frac{n_{\mathsf{B},0}}{n_{\gamma,0}} \simeq -\frac{1}{3}\frac{\alpha^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}}}{\varepsilon^{\frac{3}{4}}}\left(\frac{H}{M_{\rm Pl}}\right)^{\frac{3}{2}}. \tag{21}$$

This is directly related to the amplitude of the primordial curvature power spectrum $P_\zeta(k) = P_\zeta(k_0)(\frac{k}{k_0})^{n_s-1}$ as

$$\eta_{\mathsf{B},0} \simeq -\frac{2(2\pi)^2}{3}\left(\frac{\epsilon\alpha^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}}}{\varepsilon^{\frac{3}{4}}}\right)\left(\frac{M_{\rm Pl}}{H}\right)^{\frac{1}{2}} P_\zeta(k_0), \tag{22}$$

where ϵ is the slow-roll parameter, $\eta_B \simeq 6 \times 10^{-10}$, and $P_\zeta(k_0) \approx 2 \times 10^{-9}$ [28]. The scale of inflation, then, is

$$\frac{H}{M_{\rm Pl}} \simeq 10^{-6} \frac{\epsilon^{\frac{1}{2}}}{(\alpha^{\rm inf}_{\mathsf{B}-\underline{\mathsf{L}}})^{\frac{2}{3}}}. \tag{23}$$

Combining Eqs. (20) and (23), we find the allowed parameter space. Interestingly, it demands that $\Lambda_F \approx \Lambda_{\rm inf}$, i.e., the LR SSB should coincide with the geometrical transition that ends inflation. More precisely, we need $HM_{\rm Pl} \approx g_R^{-2} m_{W_R}^2$. Moreover, the values of these parameters are within the natural range of parameters in GUT theories [23]. For instance, with a $\xi \in (2,4)$ and $\Delta N \equiv \ln(\frac{a_{\rm reh}}{a_{\rm inf}}) \gtrsim 2$, we find that $\frac{1}{g_R} m_{W_R} \sim 10^{-4} M_{\rm Pl}$ and $H \sim 10^{-8} M_{\rm Pl}$ (see Sec. S3 of the Supplemental Material [5]).

V. COLD DARK MATTER RELIC DENSITY

The number density of $\mathbf{N}_1$ neutrino today is $n_{N_1,0} \simeq 2.8 n_{\rm B,0}$. If it makes all the DM today, its mass is

$$m_{N_1} \approx 1.8 m_P \simeq 1.7 \text{ GeV}, \tag{24}$$

where m_P is the proton mass. Since the production mechanism is independent of the active-sterile mixing angles, $\mathbf{N}_1$ can have a lifetime much larger than the age of the Universe. Nevertheless, via its loop-mediated radiative decay channel, it can decay to gamma-ray photons with energy $E_\gamma \approx m_{N_1}/2$. Demanding that $\mathbf{N}_1$ is stable over the lifetime of the Universe gives $\theta_1 < 10^{-13}$. In this framework, the generation mechanism of $\mathbf{N}_1$ is independent of its Yukawa mixing with active neutrinos, and θ_1 can be any number that satisfies the above upper bound. Thus, it may provide observable effects to be probed by gamma-ray telescopes.

VI. DISCUSSION

This paper introduced the SU$(2)_R$-axion inflation model embedded in the left-right symmetric extension of the SM. This is a new framework for simultaneous baryogenesis and dark genesis. One of the most well-studied leptogenesis scenarios with new gauge interactions is the LRSM [29,30]. Let us explore the differences between previous studies and the current proposal. The scenarios discussed thus far in the literature rely on unconstrained CP-violating phases in the neutrino sector. (a) A relic abundance of RHNs is generated after reheating by W_R interactions via freeze-out or freeze-in mechanisms. (b) The asymmetric decay of RHNs, then, creates matter asymmetry. As an alternative mechanism, the current proposal set parameters such that the aforementioned phenomena (a) and (b) are negligible. The source of asymmetry is spontaneous CP violation in the physics of inflation, and the lightest right-handed neutrino is the cold dark matter candidate. Relic abundances of SM leptons, baryons, and RHNs are generated by the chiral anomaly of W_R in inflation. Sufficient asymmetry does not require fine-tuned masses for the heaviest right-handed neutrinos, but it demands that $\Lambda_{\rm inf} \approx \Lambda_F$. Therefore, this new framework relates the scale of SU$(2)_R \times$ U$(1)_{\rm B-L}$ breaking to the end of inflation and prefers scales above 10^{10} GeV. Interestingly, this is in the range suggested by the non-supersymmetric SO(10) GUT with an intermediate left-right symmetry scale. Owing to the common origin of inflation, cold dark matter, and baryon asymmetry, this framework can naturally explain the observed coincidences among cosmological parameters, i.e., $\eta_{\rm B} \approx P_\zeta$ and $\Omega_{\rm DM} \simeq 5\Omega_{\rm B}$ with $m_{N_1} \sim 1$ GeV.

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