

Radiative seesaw model with baryon number violation and upper limit on neutron-antineutron transition time

Rabindra N. Mohapatra¹ and Nobuchika Okada²

¹*Maryland Center for Fundamental Physics and Department of Physics, University of Maryland, College Park, Maryland 20742, USA*

²*Department of Physics, University of Alabama, Tuscaloosa, Alabama 35487, USA*



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The minimal scotogenic model where small neutrino masses arise via radiative seesaw, is known to provide a unified framework for neutrino mass and origin of matter via leptogenesis. However if the inflation reheat temperature of the universe is below the sphaleron decoupling temperature, then leptogenesis fails and one way to understand the origin of matter would be to add an effective interaction involving the right-handed neutrino (RHN) N of the form $\frac{1}{\Lambda^2} N u_R d_R d_R$. This model can lead to observable neutron-antineutron ($n - \bar{n}$) oscillation if N is a Majorana fermion. We show that if RHNs are produced nonthermally, we can get a cosmological upper limit on the transition time $\tau_{n-\bar{n}}$, which is within the reach of the planned European spallation source high-intensity baryon extraction and measurement/neutron-antineutron experiment. Proton stability is guaranteed by the Z_2 invariance, which prevents the appearance of the Dirac mass term for the neutrino.

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I. INTRODUCTION

The origin of matter antimatter asymmetry in the universe is a fundamental mystery of cosmology. Particle physics in the past three decades has opened up several pathways to resolving this mystery, thereby providing a guide to different new physics scenarios beyond the standard model (BSM). It is therefore important to isolate predictions of these ideas and test them. We focus on one such direction in this paper.

A popular mechanism to understand the origin of matter is leptogenesis, which involves the decays of the hypothetical heavy Majorana right-handed neutrino (RHN) [1], used to understand small neutrino masses via the type I seesaw mechanism [2–6]. The decay process responsible is $N \rightarrow \ell + \phi$, where ℓ and ϕ denote, respectively, the left-handed lepton doublet of the standard model and the scalar doublet responsible for electroweak symmetry breaking. The interesting point about this mechanism is that it uses [1,7–10] the same interactions that play the key role in understanding the small neutrino masses.

Within the leptogenesis framework, there are two possibilities: the first possibility is where the RHN is generated by the thermal processes in the early Universe. This will

happen if the reheat temperature of the Universe after inflation is above the RHN masses. On the other hand, if the reheat temperature is lower than the RHN mass (i.e., $M_{N_a} > T_R$), the RHNs can only be produced via inflaton decay, leading to a nonthermal picture for leptogenesis (see, for example, Refs. [11–14]).

In this paper, we will consider the nonthermal route in a popular neutrino mass model known as the scotogenic model, extended to include baryon number (B) violation, and study its implications.

The scotogenic model for neutrino mass generation extends the standard model by adding three right-handed neutrinos and an extra scalar doublet (η), which is prevented from having a vacuum expectation value (VEV) [15–18]. The neutrino masses then arise as a one-loop effect [15,16] with combined suppression by the loop factor and the RHN masses. This helps to lower the seesaw scale, possibly making it accessible to colliders. One can also apply this model to implement thermal leptogenesis using the $N \rightarrow \ell + \eta$ decay as has been discussed in the literature (see, for instance, Refs. [19–24]).

The cosmology of this model depends, however, on whether the reheat temperature T_R after inflation is higher or lower than the RHN masses. In the latter case, in particular, if $T_R \leq T_{\text{sph}} \leq M_N$, where T_{sph} is the sphaleron decoupling temperature, neither thermal nor nonthermal leptogenesis can work to explain the origin of matter. To understand baryogenesis, therefore, we add an effective B-violating interaction of the form $\frac{1}{\Lambda^2} N u_R d_R d_R$, as has already been discussed in the literature [25–27]. If the

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RHNs are used in the usual type I seesaw way to get neutrino masses [2–6], then this extended model leads to rapid proton decay induced by the Dirac mass term in the seesaw matrix and is therefore not phenomenologically viable. On the other hand, if one uses the radiative seesaw (or scotogenic seesaw), this problem is avoided [25], since the Dirac mass term connecting N and ν is forbidden by a Z_2 symmetry. Furthermore, the new interaction $\frac{1}{\Lambda^2} N u_R d_R d_R$, generates the B-violating process of neutron-antineutron oscillation [6,28,29], which is observable for 10 TeV–100 TeV scale for M_N and Λ . This will make the model testable. It is therefore important to discuss the cosmological consequences of this model for baryogenesis.

We find that the model has implications for neutron-antineutron oscillation transition time, $\tau_{n-\bar{n}}$, arising from the fact that the lightest member of the second Higgs doublet η in the model is unstable. Making sure that its decay does not affect big bang nucleosynthesis (BBN) and also does not dilute the baryon asymmetry leads to upper limits on $\tau_{n-\bar{n}}$, keeping it in accessible range for the upcoming high-intensity baryon extraction and measurement (HIBEAM)/neutron-antineutron (NNBAR) experiment planned at European spallation source (ESS). To the best of our knowledge, this is the second model where $n - \bar{n}$ oscillation is directly responsible for the origin of matter, with the model leading to an upper bound on the $n - \bar{n}$ oscillation time (the other model [30] being the postsphaleron baryogenesis model with $S \rightarrow u_R d_R d_R u_R d_R d_R$ decay [31], where S is a real scalar field.). We then provide a high scale UV completion of the model and point out how proton decay is forbidden by the same Z_2 symmetry that leads to scotogenic neutrino mass generation.

The paper is organized as follows: in Sec. II, we present the effective Lagrangian of the model with baryon number violating RHN coupling; in Sec. III, we discuss the cosmology of the model, addressing questions such as baryogenesis; in Sec. IV, we discuss cosmological limits on $\tau_{n-\bar{n}}$. In Sec. V, we discuss constraints on the model from baryon number washout. In Sec. VI, we present a renormalizable UV complete version of the model. In Sec. VII, we add some comments on the model and conclude.

II. THE SCOTOGENIC LAGRANGIAN WITH B -VIOLATION

The scotogenic model for neutrino mass [15,16] is an extension of the standard model, which adds to it three heavy Majorana neutrinos N_i ($i = 1, 2, 3$) (called RHN below) and a new $SU(2)_L$ doublet, η having no VEV and interacting only with the singlet neutrinos N as $Y_D \bar{\ell} \eta N$. This results from a new Z_2 symmetry which allows only this new interaction and Majorana mass terms for the N as well as the extended two Higgs potential discussed below. Since the η doublet does not have a VEV, the model does not have a Dirac neutrino mass term and hence, does not give the usual

tree-level seesaw. Neutrino mass arises at the one-loop level. The original model has no B-violation.

The Z_2 invariant Higgs potential of the model is given by

$$V(\phi, \eta) = -\mu^2 \phi^\dagger \phi + m^2 \eta^\dagger \eta + \frac{\lambda_1}{2} (\phi^\dagger \phi)^2 + \lambda_2 (\eta^\dagger \eta)^2 + \lambda_3 (\phi^\dagger \phi) (\eta^\dagger \eta) + \lambda_4 (\phi^\dagger \eta) (\eta^\dagger \phi) + \frac{1}{2} \lambda_5 [(\phi^\dagger \eta) (\phi^\dagger \eta) + \text{H.c.}], \quad (1)$$

where ϕ is the SM Higgs doublet. To this theory, we add a new real gauge singlet field Φ to play the role of inflaton, which couples to two RHNs and a new coupling of N to three right-handed quarks so that the relevant effective Lagrangian for the model in the RHN sector looks as follows:

$$\mathcal{L}_I = Y_D^{ij} \bar{\ell}_i \eta N_j - M_i \bar{N}_i^c N_i + \frac{\kappa_i}{\Lambda^2} N_i u_R d_R d_R + \lambda_0 \Phi \bar{N}_i^c N_i + \text{H.c.}, \quad (2)$$

where ℓ is the left-handed lepton doublet of SM, and κ_i are complex phases. We will spell out the detailed action of the Z_2 symmetry that makes the above theory natural when we present the UV-complete Lagrangian in a subsequent section.

Successful inflation can be implemented via the following nonminimal coupling of Φ to gravity:

$$\mathcal{L}_g = -\frac{1}{2} (M_P^2 + 2\xi \Phi^2) R. \quad (3)$$

A suitable choice of the coupling constant $\xi > 0$ leads to a successful inflation scenario (see, for example, Ref. [32]).

Using the fact that the only field that has a VEV is the SM Higgs doublet ϕ , the masses of the physical scalars are given by [18]

$$\begin{aligned} m_h^2 &= \lambda_1 v^2, \\ m_{\eta^+}^2 &= m^2 + \frac{1}{2} \lambda_3 v^2, \\ m_{\eta_R^0}^2 &= m^2 + \frac{1}{2} (\lambda_3 + \lambda_4 + \lambda_5) v^2, \\ m_{\eta_I^0}^2 &= m^2 + \frac{1}{2} (\lambda_3 + \lambda_4 - \lambda_5) v^2, \end{aligned} \quad (4)$$

where $\langle \phi^0 \rangle = \frac{v}{\sqrt{2}}$, and $\eta_{R,I}^0$ are real and imaginary parts of the η^0 field. In what follows, we will denote $\eta_R \equiv \eta$ and choose $\lambda_5 < 0$ so that it is the lightest Z_2 -odd particle in the model, as noted above.

A. Phenomenology of the new physical η scalars in the model

The implications of the η -fields for LHC searches have been discussed, and various constraints have been derived on $\lambda_{3,4,5}$ [33–35], which parametrize their masses and

couplings. They depend sensitively on the relative mass differences among the component fields in the η doublet. The charged scalar $\eta^\pm$ can be produced at the Large Electron-Positron Collider (LEP) via the Drell-Yan process, $e^+e^- \rightarrow \eta^+\eta^-$, mediated by the photon and the Z boson. From the LEP 2 results, one obtains the lower bound $m_{\eta^\pm} > 70$ GeV [36]. The charged scalar $\eta^\pm$ can also be produced at the LHC. Once produced, it decays as $\eta^\pm \rightarrow W^\pm + \eta$. This production and decay topology are analogous to chargino ($\tilde{\chi}^\pm$) pair production followed by $\tilde{\chi}^\pm \rightarrow W^\pm + \tilde{\chi}_1^0$ in the minimal supersymmetric standard model, where $\tilde{\chi}_1^0$ is the lightest neutralino. Although the production cross section for an $\eta^\pm$ pair is smaller than that for chargino pair production by a factor of a few, it is expected that the LHC sparticle search results impose similar constraints on the inert doublet model. For example, the ATLAS Collaboration has reported sparticle search results using up to 140 fb^{-1} of integrated luminosity at $\sqrt{s} = 13$ TeV [37]. If the mass difference between the chargino and the neutralino is large, the lower bound on the chargino mass is found to be around 800 GeV. At the LHC, associated production of $\eta_I \eta^\pm$ can also occur through a W -boson mediated process, followed by the decays $\eta_I \rightarrow \eta + Z$ and $\eta^\pm \rightarrow W^\pm + \eta$. This process is analogous to $\tilde{\chi}_2^0 \tilde{\chi}^\pm$ production, followed by $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 + Z$ and $\tilde{\chi}^\pm \rightarrow W^\pm + \tilde{\chi}_1^0$, where $\tilde{\chi}_2^0$ is the second-lightest neutralino. Searches for three leptons plus missing transverse energy arising from leptonic decays of the W and Z bosons yield even stronger constraints: assuming the lightest neutralino is much lighter, the lower bound on the chargino mass reaches approximately 1 TeV [37]. These severe constraints, however, do not apply when the lightest neutralino is nearly degenerate with the chargino. For example, for $m_{\tilde{\chi}^\pm} - m_{\tilde{\chi}_1^0} \simeq 1$ GeV, only the LEP bound is relevant [37].

Note that when perturbativity of the quartic couplings is imposed up to the inflationary or Planck scale, their magnitudes are typically bounded by $|\lambda_i| \lesssim 0.3$. See, for example, Ref. [38] for a detailed renormalization group analysis. For our choice of parameters, we have explicitly checked that all running couplings λ_i remain in the perturbative regime. Moreover, the running quartic coupling λ_1 of the SM Higgs doublet remains positive for $|\lambda_{3,4,5}| \gtrsim 0.1$, ensuring the stability of the electroweak vacuum.

Furthermore, we demand that the masses of the η fields be less than a TeV to avoid destabilizing the SM Higgs boson mass, since the one-loop contribution to the Higgs boson mass in the theory is given by

$$\Delta m_\phi^2 \sim \frac{\lambda_{3,4}}{16\pi^2} m_\eta^2. \quad (5)$$

This is the naturalness bound on η -masses.

To close this section, we note that the lightest η particle is unstable, unlike in the original scotogenic model, and decays to $\eta \rightarrow u d \bar{\nu}$ with a decay width

$$\Gamma_\eta \simeq 1.3 \times 10^{-7} \frac{Y_D^2 m_\eta^7}{M_N^6 \beta^4}, \quad (6)$$

where $\beta = \Lambda/m_\eta$, and clearly within our assumption of effective B-violating Lagrangian $\beta \gg 1$. The η field could be fairly long lived with lifetime close to one second or less. This will have consequences for the $n - \bar{n}$ transition, as we discuss later.

III. COSMOLOGY OF THE MODEL

The cosmology of the model is as follows: for values of Φ of order or higher than the Planck mass, M_P , the potential $V(\Phi)$ in the Einstein frame is almost a constant as a function of Φ , giving rise to inflation. As Φ rolls down, the inflaton potential $V(\Phi)$ is no more a constant and inflation ends [39]. After this, the inflation field decays causing reheating of the Universe. We arrange parameters such that the inflaton decay width $\Gamma_{\Phi \rightarrow NN} \ll \Gamma_N$ so that once N is created from the decay of inflaton, it immediately decays and reheats the Universe. Thus production of N is a nonthermal process. The number density of N subsequent to Φ decay, n_N , is then equal to twice n_Φ , the number density of the inflaton.

A. Failure of nonthermal leptogenesis

Baryogenesis in this model can take place either via $N \rightarrow \ell + \eta$ decay in combination with sphalerons or via $N \rightarrow 3q$ decay. If, however, the reheat temperature T_R after inflation is below the sphaleron decoupling temperature $T_{\text{sph}} \simeq 130$ GeV, which we assume here, leptogenesis does not work and matter-antimatter asymmetry can only arise from the $N \rightarrow 3q$ decay in combination with CP violation.

B. Baryogenesis via $N \rightarrow u_R d_R d_R$ decay

The baryon number generated by N decay is given by

$$n_B = n_N \text{Br}(N \rightarrow 3q) \epsilon_B = 2 \frac{\rho_\Phi}{m_\Phi} \text{Br}(N \rightarrow 3q) \epsilon_B, \quad (7)$$

where ϵ_B is the baryon asymmetry parameter that arises from tree and one-loop interference of diagrams for $N \rightarrow 3q$ decay, and we have used $n_N = 2n_\Phi = \frac{\rho_\Phi}{m_\Phi}$ with the inflaton mass m_Φ and ρ_Φ being the inflaton energy density. Assuming η to be out of equilibrium, the radiation energy density created by N decay at the time of inflaton decay is estimated as

$$\rho_{\text{rad}} = \rho_\Phi \left(1 - \frac{1}{2} \text{Br}(N \rightarrow \eta \ell) \right) = \frac{\pi^2}{30} g_* T_R^4, \quad (8)$$

where $g_* \simeq 100$ is the effective degrees of freedom of relativistic particles in thermal plasma. Using $\rho_{\text{rad}}/s = 3/4T_R$ for the thermal plasma, we estimate the baryon asymmetry generated by the $N \rightarrow 3q$ decay as

$$\frac{n_B}{s} \simeq \frac{3T_R}{2m_\Phi} \frac{\text{Br}(N \rightarrow 3q)}{1 - \frac{1}{2}\text{Br}(N \rightarrow \eta\ell)} \epsilon_B. \quad (9)$$

Similarly, we can derive an expression for $Y_\eta \equiv \frac{n_\eta}{s}$, as follows:

$$Y_\eta(t = \tau_\Phi) = \frac{3T_R}{2m_\Phi} \frac{\text{Br}(N \rightarrow \eta\ell)}{1 - \frac{1}{2}\text{Br}(N \rightarrow \eta\ell)}. \quad (10)$$

Combining Eqs. (9) and (10), we can express Y_η in terms of the parameters of the theory:

$$Y_\eta \simeq \frac{10^{-10}}{\epsilon_B} \frac{\text{Br}(N \rightarrow \ell + \eta)}{\text{Br}(N \rightarrow 3q)}, \quad (11)$$

where we have used the observed value, $n_B/s \simeq 10^{-10}$. In the above equation, the asymmetry parameter ϵ_B is given by [25,26]

$$\epsilon_B \simeq \frac{\text{Im}(\kappa_1^{2*}\kappa_2^2)}{3072\pi^3|\kappa_1|^2} \left(\frac{M_1}{\Lambda}\right)^4 \frac{M_1 M_2}{M_2^2 - M_1^2}. \quad (12)$$

where we have denoted RHN masses for different generations $M_{1,2,3}$ with at least two of them being quasidegenerate to facilitate an adequate value for ϵ_B .

The partial decay widths of $N \rightarrow \ell + \eta$ and $N \rightarrow 3q$ are, respectively, estimated as

$$\Gamma_{N \rightarrow \ell + \eta} \simeq \frac{Y_D^2}{8\pi} M_N, \quad \Gamma_{N \rightarrow 3q} \simeq \frac{M_N^5}{192\pi^3 \Lambda^4}. \quad (13)$$

Using $\frac{\text{Br}(N \rightarrow \eta\ell)}{\text{Br}(N \rightarrow 3q)} = \frac{\Gamma_{N \rightarrow \ell + \eta}}{\Gamma_{N \rightarrow 3q}}$, we get

$$Y_\eta \simeq \frac{10^{-10}}{\epsilon_B} 24\pi^2 Y_D^2 \beta^4. \quad (14)$$

This gives us an expression for the energy density of η field when it decays ($t = \tau_\eta$) as

$$\rho_\eta(t = \tau_\eta) = m_\eta Y_\eta s(t = \tau_\eta), \quad (15)$$

where τ_η is the lifetime of η , and $s(t = \tau_\eta)$ is the entropy density of the Universe at $t = \tau_\eta$.

IV. THEORETICAL LIMITS ON $\tau_{n-\bar{n}}$ IN THE MODEL

To convert the above results into a bound on $n - \bar{n}$ transition time, we note that $n - \bar{n}$ oscillation in the model arises from the exchange of Majorana N between two

$u_R d_R d_R$ vertices. The resulting $\Delta B = 2$ interaction has the strength

$$\mathcal{L}_{\Delta B=2} = \frac{1}{\Lambda^4 M_N} u_R d_R d_R u_R d_R d_R + \text{H.c.} \quad (16)$$

Clearly, this is directly connected to the Majorana mass of the neutrinos. After hadron dressing, we find the $\tau_{n-\bar{n}}$ transition time to be

$$\tau_{n-\bar{n}}[s] \simeq \frac{M_N \Lambda^4}{|\mathcal{M}|} \times 6.6 \times 10^{-25} \simeq 10^{-20} M_N^5 \beta^4, \quad (17)$$

where M_N is given in units of GeV, and $\mathcal{M} = -6.5 \times 10^{-5} \text{ GeV}^6$ from a lattice QCD calculation [40]. Similar hadron dressing factors $|\mathcal{M}|$ also arise in early bag model calculations [41]. Proton decay is forbidden in the model since η -field does not have a VEV.

A. Constraints from the decay time of η

Note that the new neutral scalar field η has interactions mediated by W, Z bosons which keep it in equilibrium until quite low temperature. The processes of concern here are $\eta + \eta \rightarrow WW, ZZ, hh$ which deplete its density, whereas the inverse processes get it into thermal equilibrium. The temperature at which η fully decouples from the cosmic soup can be determined by demanding that

$$n_\eta \langle \sigma(\eta\eta \rightarrow SM) v_{\text{rel}} \rangle \simeq \frac{n_\eta(T_d)}{4\pi m_\eta^2} = \sqrt{\frac{\pi^2 g_*}{90}} \frac{T_d^2}{M_P}. \quad (18)$$

Taking $n_\eta(T_d) \simeq \left(\frac{m_\eta T_d}{2\pi}\right)^{3/2} e^{-m_\eta/T_d}$, $g_* \simeq 100$ and $M_P \simeq 2.4 \times 10^{18} \text{ GeV}$, we find T_d and plot it as a function of m_η in Fig. 1. For η to be nonthermal, we must have $T_d > T_R$. The naturalness bound of $m_\eta < \text{TeV}$ then implies from Fig. 1 that the reheat temperature of the Universe in our model must be at or below 10 GeV.

Due to its large mass, η cloud carries a lot of energy, which is released when it decays. This affects the big bang cosmology. If η decays after BBN, the energetic decay products may destroy light nuclei successfully synthesized in the BBN epoch. We therefore demand that $\tau_\eta \gtrsim 1 \text{ sec}$, or equivalently, the temperature of the thermal plasma T_{decay} when η decays should be above a few MeV.

For $m_\eta < M_N$, η decays to $\nu + 3q$ mediated by the Majorana RHN, and this decay width is estimated as

$$\Gamma_\eta \simeq PS^{(4)} \frac{Y_D^2 m_\eta^7}{M_N^2 \Lambda^4} \simeq 1.3 \times 10^{-7} \times \frac{Y_D^2 m_\eta^7}{M_N^6 \beta^4}, \quad (19)$$

where we estimate the four body phase space factor $PS^{(4)} = \frac{1}{2(4\pi)^5 \Gamma(4) \Gamma(3)} \simeq 1.3 \times 10^{-7}$. Using $\Gamma_\eta = H(T_{\text{decay}})$, we find

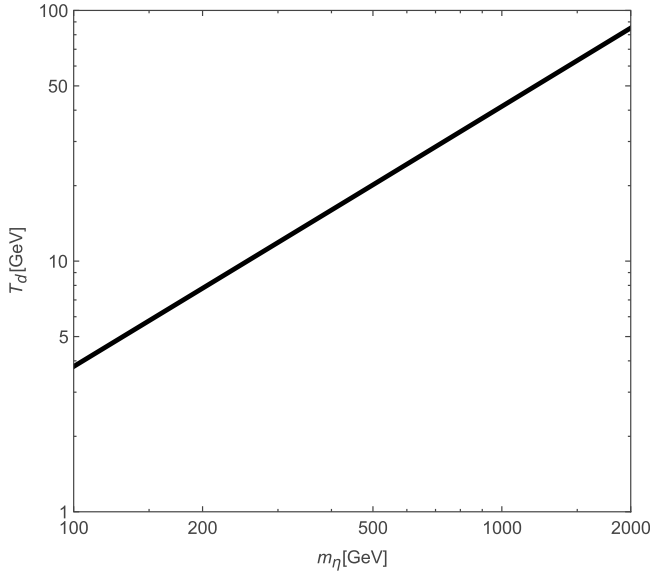


FIG. 1. Profile of the T_d as a function of η mass. We find that $m_\eta/T_d \simeq 20\text{--}27$.

$$T_{\text{decay}} = \left(1.3 \times 10^{-7} \times \sqrt{\frac{90}{\pi^2 g_*}} \frac{Y_D^2 m_\eta^7}{M_N^6 \beta^4} M_P \right)^{1/2}. \quad (20)$$

We use an approximate expression for one-loop m_ν in this model, as follows:

$$m_\nu \simeq \frac{|\lambda_5| Y_D^2 v^2}{16\pi^2 M_N}, \quad (21)$$

with $m_\nu \simeq 5 \times 10^{-11}$ GeV for a typical neutrino mass scale from the neutrino oscillation data to replace $Y_D^2 = 2.6 \times 10^{-13} \frac{M_N}{|\lambda_5|}$. Here, M_N is given in units of GeV. Using the fact that $\tau_{n-\bar{n}} \simeq 10^{-20}(\text{sec.}) M_N^5 \beta^4$ in this model, we get

$$\tau_{n-\bar{n}}[\text{s}] = \simeq 2.4 \times 10^{-22} \frac{m_\eta^7}{|\lambda_5| T_{\text{decay}}^2}, \quad (22)$$

where m_η and T_{decay} are given in units of GeV.

The limits on $\tau_{n-\bar{n}}$ for plausible parameter ranges of the model are shown in Fig. 2. We see that for most of the parameter range of the model, the predicted upper limits are within the accessible range of the planned HIBEAM/NNBAR experiment at ESS. We also note from Fig. 2 that higher values of $|\lambda_5|$ lead to lower $\tau_{n-\bar{n}}$ that are in conflict with current experimental limits. Thus, η lifetime considerations prefer lower values of $|\lambda_5|$.

B. Limits from dilution effects on n_B from the decay of η

Since η is a long-lived particle, it may dominate the energy density of the Universe when it decays. If so, the

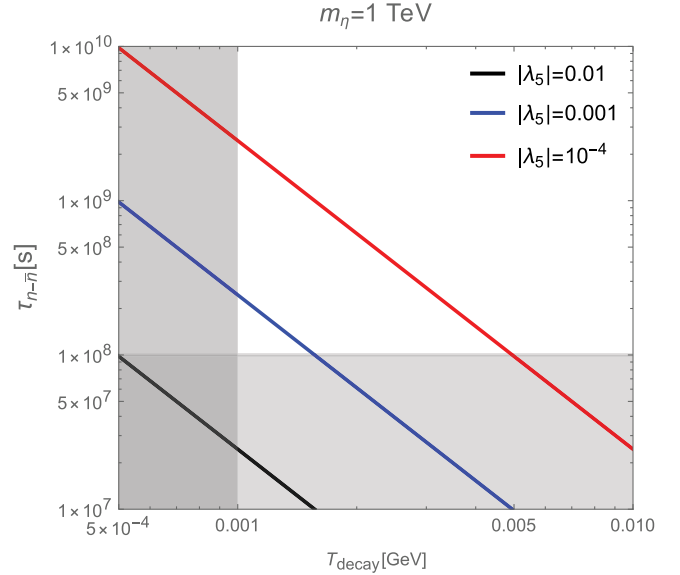


FIG. 2. Predicted free neutron-antineutron oscillation time $\tau_{n-\bar{n}}$ is shown as a function of T_{decay} for several values of $|\lambda_5|$. The gray-shaded region is excluded by the current experimental limit, $\tau_{n-\bar{n}} > 10^8$ sec. [42,43] and $T_{\text{decay}} > 1$ MeV from BBN.

baryon asymmetry created before the decay is diluted. To avoid the dilution, we must therefore demand that $\rho_\eta(t = \tau_\eta) < \rho_{\text{rad}}(t = \tau_\eta)$. This condition translates to $m_\eta Y_\eta S(T_{\text{decay}}) < \rho_{\text{rad}}(T_{\text{decay}})$, or equivalently,

$$m_\eta Y_\eta < \frac{3}{4} T_{\text{decay}}. \quad (23)$$

Combining this with Y_η from Eq. (14), T_{decay} from Eq. (20), and $Y_D^2 = 2.6 \times 10^{-13} \frac{M_N}{|\lambda_5|}$, the condition translates to

$$\frac{m_\eta Y_\eta}{3/4 T_{\text{decay}}} = 5.2 \times 10^{-20} \frac{\beta^6}{\epsilon_B} \sqrt{\frac{M_N^7}{m_\eta^5 |\lambda_5|}} < 1. \quad (24)$$

Applying this condition to Eq. (17), we get an upper bound on $\tau_{n-\bar{n}}$, as follows:

$$\tau_{n-\bar{n}}[\text{s}] < \frac{3.5 \times 10^7 m_\eta^{25/7} \epsilon_B^{10/7} |\lambda_5|^{5/7}}{\beta^{32/7}}, \quad (25)$$

where m_η is given in units of GeV.

Noting the naturalness bound of $m_\eta \lesssim 1$ TeV, $\beta \equiv \frac{\Lambda}{M_N} \gg 1$, $\epsilon_B < 1$, and $|\lambda_5| \ll 1$ as suggested by Fig. 2, we plot the theoretical upper limits on $\tau_{n-\bar{n}}$ in Figs. 3–5. We find that most of the upper limits are below 10^{11} sec for a large range of parameters $\frac{\Lambda}{M_N}$, ϵ_B and $|\lambda_5|$ in the model. We see that the predicted upper limits on $\tau_{n-\bar{n}}$ from dilution constraint are complementary to the decay constraints. Combining the results in Figs. 2 and 4, we see that our model allows only a limited range of values for $\tau_{n-\bar{n}}$, all of

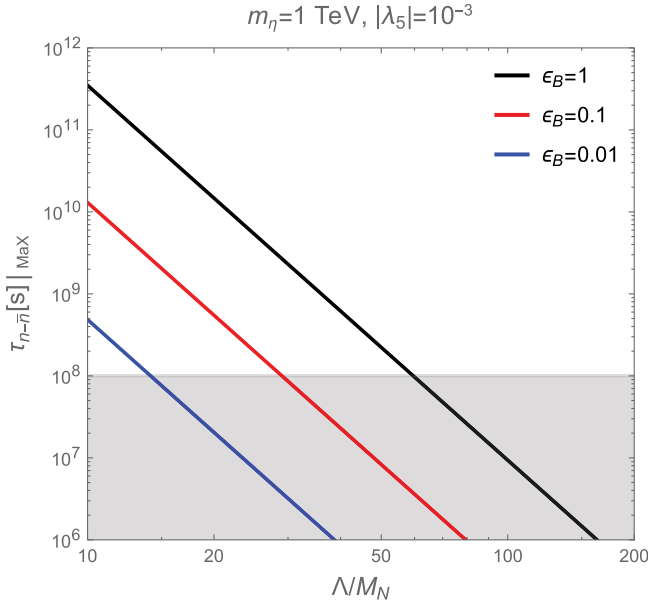


FIG. 3. Theoretical upper limit on $\tau_{n-\bar{n}}$ as a function of $\frac{\Lambda}{M_N}$ for $|\lambda_5| = 10^{-3}$ and various values of ϵ_B .

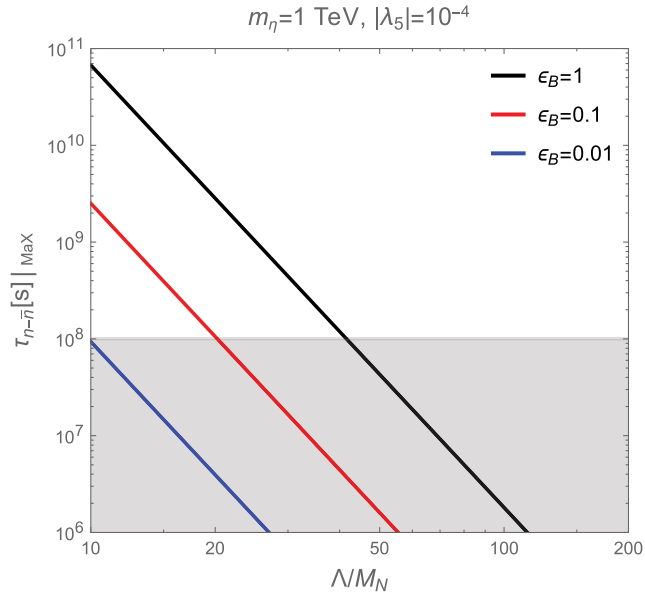


FIG. 4. Theoretical upper limit on $\tau_{n-\bar{n}}$ as a function of $\frac{\Lambda}{M_N}$ for $|\lambda_5| = 10^{-4}$ and various values of ϵ_B .

which is accessible to the planned ESS HIBEAM/NNBAR experiment [45]. A negative result in this experiment, therefore, will rule out our model for baryogenesis. This is the main result of our paper.

V. AVOIDING BARYON ASYMMETRY WASH-OUT

We have to make sure that there is no wash-out process that can reduce the baryon asymmetry in the model.

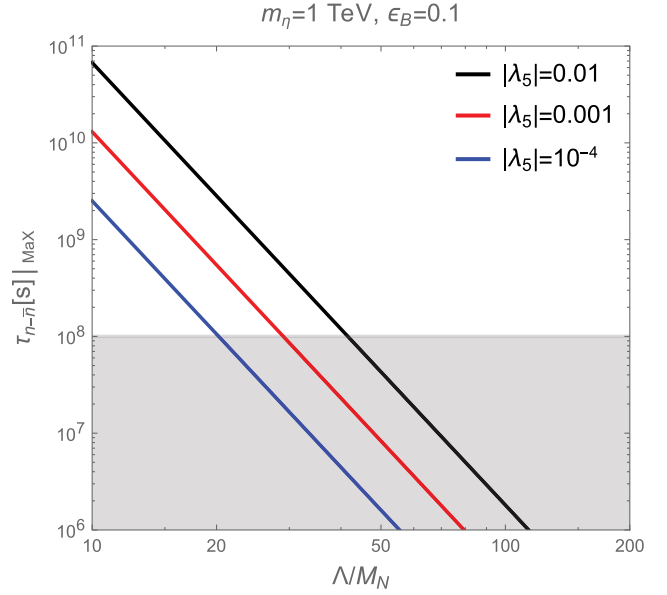


FIG. 5. Theoretical upper limit on $\tau_{n-\bar{n}}$ as a function of $\frac{\Lambda}{M_N}$ for $\epsilon_B = 0.1$ for different values of $|\lambda_5| = 10^{-4}$.

One process that can wash out the baryon asymmetry is $qq \rightarrow \bar{q}\bar{q}\bar{q}\bar{q}$ generated by the $\Delta B = 2$ interaction, $\frac{1}{\Lambda^4 M_N} u_R d_R d_R u_R d_R d_R$. Its rate must be such that its decoupling temperature T'_d is higher than the reheating temperature T_R . To calculate T'_d , we set $n_q(T'_d)\sigma(qq \rightarrow 4\bar{q})v_{\text{rel}} \simeq \frac{T_d'^2}{M_P}$. Taking $n_q(T'_d) \simeq 4 \times 3 \times \frac{3\zeta(3)}{4\pi^2} T_d'^3$ and $\sigma(qq \rightarrow 4\bar{q}) \simeq PS^{(4)} \frac{T_d'^8}{m_N^2 \Lambda^8}$, where we estimate the four body phase space factor $PS^{(4)} = \frac{1}{2(4\pi)^5 \Gamma(4) \Gamma(3)} \simeq 1.3 \times 10^{-7}$, we get T'_d as

$$T'_d \simeq 5.8 \times \left(\frac{M_N}{M_P}\right)^{1/9} \beta^{8/9} M_N. \quad (26)$$

From our results shown in Figs. 1–4, we can see $10^8 < \tau_{n-\bar{n}}[\text{s}] \lesssim 10^{11}$ and $\beta \lesssim 50$. Using Eq. (17), we find

$$4.0 \times 10^5 < M_N[\text{GeV}] \beta^{4/5} \lesssim 1.6 \times 10^6, \quad (27)$$

which means $M_N[\text{GeV}] = 10^4 - 10^5$ for $10 < \beta \lesssim 50$. Hence, we find $T'_d = \mathcal{O}(m_N)$. To avoid the wash-out, the reheating temperature T_R must be less than T'_d . Our choice $T_R < T_{\text{sph}} \simeq 130$ GeV satisfies this condition. (For a discussion of wash-out in generic neutron oscillation models, see Ref. [44]).

VI. UV-COMPLETE MODEL FOR THE EFFECTIVE B-VIOLATING INTERACTION

In this section, we present a high scale renormalizable model [25] which in the low energy effective theory leads to the Lagrangian in Eq. (2). We work within the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ and extend its particle

content by the addition of three RHNs (N_i), an extra doublet (η) and an $SU(2)_L$ -singlet, color-triplet scalar (Δ) with hypercharge $Y = +2/3$ and a complex gauge singlet Φ (inflaton). We denote the quark and lepton doublets of the SM by $Q_L^T = (u_L, d_L)$ and $L^T = (\nu_L, e_L)$, and singlets by u_R, d_R, e_R ; the SM Higgs doublet field is denoted by ϕ . We impose an additional Z_2 -symmetry under which Q_L, u_R, d_R, N_a , and η fields are taken to be odd, whereas L, e_R, ϕ, Δ , and Φ fields are taken to be even. The gauge- and Z_2 -invariant Yukawa interaction Lagrangian involving the leptons and the new fields Δ, Φ , and N in the model is given by

$$\mathcal{L} = \Phi \bar{N}^c N + \bar{Q}_L \phi u_R + \bar{Q}_L \tilde{\phi} d_R + \bar{L} \tilde{\phi} e_R + \bar{L} \eta N + \text{H.c.}, \quad (28)$$

where, for simplicity, we have suppressed the coupling constants as well as flavor indices.

To this, if we add the coupling and mass terms of the form

$$\mathcal{L}_Y = \lambda_1 \Delta u_R d_R + \lambda_2 N \Delta^* d_R + \text{H.c.}, \quad (29)$$

the low energy effective Lagrangian for $M_N, m_\eta \ll M_\Delta$ has the additional terms

$$\mathcal{L}_{\text{eff}}^N = \frac{\kappa_i}{\Lambda^2} (N_i u^c d^c d^c) + Y_D \bar{L} \eta N + \text{H.c.} \quad (30)$$

where $\Lambda^2 = M_\Delta^2 / (\lambda_1 \lambda_2)$. It then becomes immediately clear that combination of Eqs. (29) and (30) leads to an operator of type $u_R d_R d_R \nu$, when η acquires a VEV. This leads to proton decay $p \rightarrow K^+ \bar{\nu}, \pi^+ \bar{\nu}$. It is therefore essential that

η does not acquire a VEV in this model and also $m_\eta > m_p$, as is the case for our model.

VII. CONCLUSION

In conclusion, we first show that if the heavy right-handed neutrinos in the minimal scotogenic model for neutrino mass are produced in the decay of the inflaton field and $T_R \leq 10$ GeV, then usual leptogenesis scenario does not work to produce simultaneously the baryon asymmetry. In order to cure this problem, we introduce a dimension-six effective baryon-violating term involving the RHN so that its B-violating decay can produce the desired amount of baryon-antibaryon asymmetry in the Universe. An experimental signature of this model is an observable baryon number-violating process of neutron-antineutron oscillation. The model has a long-lived light neutral scalar, whose decay should not affect big bang nucleosynthesis and should not dilute the produced baryon asymmetry. We find that our model allows only a limited range of values for $\tau_{n-\bar{n}}$, all of which is accessible to the planned ESS HIBEAM/NNBAR experiment. A negative result in this experiment will therefore rule out our model for baryogenesis. We also present a UV-complete renormalizable high scale version of this model.

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DATA AVAILABILITY

No data were created or analyzed in this study.

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