

Thermal leptogenesis in the Babu-Nandi-Tavartkiladze model of neutrino mass

P. S. Bhupal Dev^{1,2,*}, Srubabati Goswami^{3,†}, Debashis Pachhar^{3,4,‡} and Drona Vatsyayan^{5,6,§}

¹*Department of Physics and McDonnell Center for the Space Sciences,
Washington University, St. Louis, Missouri 63130, USA*

²*PRISMA⁺⁺ Cluster of Excellence and Mainz Institute for Theoretical Physics,
Johannes Gutenberg-Universität Mainz, 55099 Mainz, Germany*

³*Theoretical Physics Division, Physical Research Laboratory, Navrangpura, Ahmedabad 380009, India*

⁴*Indian Institute of Technology Kanpur, Kalyanpur, Kanpur Nagar, Uttar Pradesh 208016, India*

⁵*Department of Physics, Carleton University, Ottawa, Ontario K1S 5B6, Canada*

⁶*Arthur B. McDonald Canadian Astroparticle Physics Research Institute,
64 Bader Lane, Queen's University, Kingston, Ontario K7L 3N6, Canada*



(Received 18 August 2026; accepted 3 September 2026; published 1 October 2026)

We investigate neutrino mass and thermal leptogenesis in the Babu-Nandi-Tavartkiladze model featuring a scalar quadruplet (Φ) and a pair of vectorlike fermion triplets (Σ). In this framework, neutrino masses are generated via an effective dimension-7 operator $LLHH(H^\dagger H)/\Lambda^3$ at the tree level and via the dimension-5 operator $LLHH/\Lambda$ at the one-loop level. It naturally accommodates sub-eV neutrino masses even if the new physics scale Λ is $\mathcal{O}(\text{TeV})$, thus making the model a compelling target for experimental searches. We explore the viability of thermal leptogenesis in this model, which is distinct from the canonical seesaw-based leptogenesis due to the presence of vector-like fermions. We find that leptogenesis is viable for $M_\Sigma \gtrsim 10^7 \text{ GeV}$ for a hierarchical spectrum of fermion triplets. However, in the quasidegenerate regime, resonant enhancement of the CP asymmetry lowers this scale down to $\mathcal{O}(\text{TeV})$, reconciling successful leptogenesis with the originally motivated TeV-scale phenomenology and testability of the model at colliders.

DOI: [10.1103/PhysRevD.114.075001](https://doi.org/10.1103/PhysRevD.114.075001)

I. INTRODUCTION

Neutrino oscillation experiments over the past three decades have firmly established that at least two of the three active neutrino mass eigenstates must have small but nonzero eigenvalues. Since neutrinos are predicted to be exactly massless within the Standard Model (SM), these observations provide a concrete laboratory evidence for beyond the Standard Model (BSM) physics. A particularly elegant explanation of small neutrino masses is provided by the effective dimension-5 Weinberg operator $LLHH/\Lambda$ [1] (where L and H are the SM lepton and Higgs doublets, respectively, and Λ is the new physics scale), which generates Majorana neutrino masses after electroweak

symmetry breaking (EWSB). In ultraviolet-complete theories, this operator is commonly realized at the tree level through type-I [2–5], type-II [6–11], or type-III [12] seesaw mechanism involving either heavy fermion singlets, $SU(2)_L$ scalar triplets, or $SU(2)_L$ fermion triplets, respectively.

The Weinberg operator violates lepton number by two units. In the canonical seesaw mechanisms, the scale of lepton number violation (LNV) is typically the same as the mass of the heavy particles responsible for the smallness of neutrino masses. According to 't Hooft naturalness arguments [13], the natural value of the seesaw scale is $\sim 10^{14} \text{ GeV}$, close to the Grand Unified Theory scale. Such high scales associated with canonical dimension-5 seesaw scenarios often make direct experimental tests of LNV rather difficult.¹ Several mechanisms have been proposed, e.g., variations of the singlet seesaw with additional fermions [16–19], and radiative seesaw [20–29], which can naturally lower the seesaw scale to the TeV

*Contact author: bdev@wustl.edu

†Contact author: sruaba@prl.res.in

‡Contact author: debashisp@iitk.ac.in

§Contact author: drona@physics.carleton.ca

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.

¹Lowering the seesaw scale in the minimal scenario requires either tiny Yukawa couplings or approximate conservation of lepton number [14,15].

range, and hence, enhance its experimental testability [30,31]. Another appealing possibility is to generate neutrino masses through higher-dimensional operators [32–45]. In such scenarios, the tree-level neutrino mass is typically suppressed by factors of the form $1/\Lambda^{d-4}$, where $d \geq 7$ denotes the operator dimension. As a result of this additional suppression associated with higher-dimensional operators, the scale of new physics responsible for neutrino mass can naturally be lowered to the TeV range without requiring extremely small Yukawa couplings or fine-tunings. These operators may arise either at the tree level or radiatively at the loop level, and generally necessitate an extended particle spectrum involving new fermions and scalar fields in nontrivial representations of $SU(2)_L$. Since the associated new physics scale can lie around a few TeV, these models offer rich phenomenological prospects, including possible signatures at the Large Hadron Collider (LHC).

Here, we focus on a particular model realization by Babu, Nandi and Tavarukiladze (BNT) based on a dimension-7 operator involving scalar quadruplets and vectorlike fermion triplets [34]. The implications of this model for colliders have been studied in Refs. [45–50], for muon $g-2$ in Ref. [51] and the vacuum structure of the model was recently studied in Ref. [52]. Note that, in the BNT model the dimension-5 Weinberg operator can also be generated at one loop, and whether the dominant contribution to the neutrino mass comes from the dimension-7 operator or the dimension-5 operator depends on the masses of the new particles [46].

Apart from the origin of neutrino mass, another open problem in particle physics and cosmology is the origin of the observed baryon asymmetry of the Universe (BAU). An elegant framework to address this issue is the mechanism of leptogenesis [53], in which a primordial lepton asymmetry generated in the early Universe is partially converted into the observed baryon asymmetry through electroweak sphaleron processes [54]. What makes this option very attractive is that leptogenesis can occur naturally in seesaw models that inherently incorporate LNV through the Weinberg operator.² Conventional leptogenesis is usually realized at very high energy scales $\gtrsim \mathcal{O}(10^9)$ GeV [57,58], often associated with canonical seesaw models [59,60], making them difficult to test experimentally.³ This has

²An alternative is Dirac leptogenesis without violating global lepton number [55,56], provided there is a mechanism to source an initial asymmetry between left- and right-handed neutrinos (and their corresponding antiparticles). The electroweak sphalerons partially convert the left-handed lepton asymmetry into a baryon asymmetry, while the tiny Dirac Yukawa couplings prevent the right-handed neutrinos from reaching thermal equilibrium.

³Including flavor effects can lower the leptogenesis scale to $\sim 10^6$ GeV, but requires certain degree of cancellation between the tree and one-loop level contribution to the light neutrino mass matrix [61].

motivated growing interest in low-scale leptogenesis frameworks [62–64], where the origin of neutrino mass and the BAU can be connected to new physics accessible at present or future experiments. To this extent, it is natural to explore if leptogenesis can be realized in the neutrino mass models of higher-dimensional operators involving new fermion and scalar $\mathbf{n}$ triplets of $SU(2)_L$ (with $\mathbf{n} \geq 3$), and whether it is viable at the TeV scale.

To this effect, we study the possibility of thermal leptogenesis in the BNT model. The presence of vectorlike fermions can have interesting consequences for leptogenesis, which can be contrasted against the Majorana fermion singlet/triplet in type-I [59]/type-III [65] seesaw leptogenesis. Similarly, the presence of a new scalar multiplet in the model can be contrasted with type-II seesaw leptogenesis [66–68]. Building upon the features of leptogenesis with canonical seesaw, we will focus on the decays of heavy vectorlike triplets to realize thermal leptogenesis. Unlike the canonical seesaw, the Yukawa couplings in the BNT model are not necessarily suppressed by neutrino masses. The suppression is rather attributed to the small vacuum expectation value (VEV) of the quadruplet (v_Φ), which in turn is induced from the SM Higgs VEV via a quartic term in the scalar potential. Therefore, the analog of the Davidson-Ibarra (DI) bound [57] for vanilla leptogenesis does not apply here, hence allowing for a lower leptogenesis scale. We find that the scale of leptogenesis can be as low as 3.5×10^7 GeV in the hierarchical case. This can be contrasted against the Majorana fermion singlet (triplet) in type-I (III) leptogenesis, where vanilla leptogenesis occurs only above 10^9 GeV [57] (3×10^{10} GeV [65]).

In addition to the hierarchical case, we also consider the possibility of low-scale resonant leptogenesis considering a quasidegenerate spectrum of the vectorlike triplet fermions in the BNT model. As expected, the scale of leptogenesis in this case can be lowered to the TeV scale. We find a lower bound of 1.7 TeV on the triplet fermion mass, very close to the theoretical lower bound of 1.6 TeV which applies to both type-II and type-III seesaw leptogenesis [69].

A unique feature of the BNT model is that LNV can originate either from the scalar potential, when the scalar quadruplet carries a lepton number of -2 , or from the Yukawa interaction if the scalar quadruplet is assigned a zero lepton number. In the first case, the decay process of the vector-like fermions itself is not lepton number violating, which is similar to the Dirac leptogenesis scenario. However, because the vectorlike fermion can decay to two different scalars, a net lepton asymmetry can nevertheless be generated if the couplings are different. Note that this mechanism differs from the conventional Dirac leptogenesis because of the presence of an explicit lepton number violating term in the Lagrangian. Conversely, when the scalar does not carry the lepton number, the decay process

itself violates the lepton number, as in the canonical seesaw leptogenesis.

The rest of the article is organized as follows: In Sec. II, we describe the basic features of the BNT model including its scalar and fermion sectors. In Sec. III, we describe the neutrino mass generation in the BNT model and outline the adopted Yukawa parameterization procedure. We discuss the CP asymmetry generation and the Boltzmann equations relevant for thermal leptogenesis for a hierarchical spectrum of heavy fermion triplets in Sec. IV and present our results in Sec. V. We then discuss the possibility of low-scale leptogenesis and the relevant constraints in Sec. VI. We finally conclude in Sec. VII. The decay and scattering rates are given in Appendix A and the sphaleron conversion factor specific to this model is calculated in Appendix B.

II. THE BNT MODEL

In addition to the SM particles, the BNT model [34] contains a scalar quadruplet $\Phi \sim (\mathbf{1}, \mathbf{4}, 3/2)$ and 2 generations of a pair of vectorlike fermions $\Sigma_{L,R} \sim (\mathbf{1}, \mathbf{3}, 1)$, with

$$\Phi(\mathbf{1}, \mathbf{4}, 3/2) = \{\Phi^{+++}, \Phi^{++}, \Phi^+, \Phi^0\}, \quad (1)$$

$$\Sigma_{L,R}(\mathbf{1}, \mathbf{3}, 1) = \{\Sigma_{L,R}^{++}, \Sigma_{L,R}^+, \Sigma_{L,R}^0\}. \quad (2)$$

The particles and their charge assignments are given in Table I. We also show the global lepton number (L) assignments for the fields.

Note that $\Phi = \Phi_{ijk}$ ($i, j, k = 1, 2$) is a totally symmetric tensor of rank 3, whose components are given by

$$\begin{aligned} \Phi_{111} &= \Phi^{+++}, & \Phi_{112,121,211} &= \frac{\Phi^{++}}{\sqrt{3}}, \\ \Phi_{122,212,221} &= \frac{\Phi^+}{\sqrt{3}}, & \Phi_{222} &= \Phi^0. \end{aligned} \quad (3)$$

Similarly, $\Sigma = \Sigma_{ij}$ is a totally symmetric tensor of rank 2, and its components are

$$\Sigma_{11} = \Sigma^{++}, \quad \Sigma_{22} = \Sigma^0, \quad \Sigma_{12,21} = \frac{\Sigma^+}{\sqrt{2}}. \quad (4)$$

TABLE I. Particle content of the BNT model and their charge assignments under the SM gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$. We also show the lepton number assignments: Σ can be assigned either $L = -1$ or $+1$, and similarly Φ can carry $L = -2$ or 0 .

Field	$SU(3)_C \times SU(2)_L \times U(1)_Y$	L
L	$(\mathbf{1}, \mathbf{2}, -1/2)$	$+1$
H	$(\mathbf{1}, \mathbf{2}, 1/2)$	0
Σ	$(\mathbf{1}, \mathbf{3}, 1)$	$-1/+1$
Φ	$(\mathbf{1}, \mathbf{4}, 3/2)$	$-2/0$

The relevant Lagrangian for the model reads as

$$\mathcal{L} \supset \mathcal{L}_{\text{kin}} + \mathcal{L}_{\text{yuk}} - \mathcal{V}(H, \Phi), \quad (5)$$

where, suppressing the $SU(2)$ indices, the kinetic and Yukawa terms are, respectively, given by

$$\begin{aligned} \mathcal{L}_{\text{kin}} &= (D^\mu H)^\dagger (D_\mu H) + (D^\mu \Phi)^\dagger (D_\mu \Phi) \\ &\quad + \bar{\Sigma}_{L,Rj} i \gamma^\mu D_\mu \Sigma_{L,Rj}, \end{aligned} \quad (6)$$

$$\begin{aligned} -\mathcal{L}_{\text{yuk}} &= M_{\Sigma ij} \bar{\Sigma}_{Li} \Sigma_{Rj} + Y_{H\alpha i} (\bar{L}_\alpha)^c \epsilon \Sigma_{Li} H^* \\ &\quad + Y_{\Phi\alpha i} \bar{L}_\alpha \epsilon \Phi^* \Sigma_{Ri} + \text{H.c.}, \end{aligned} \quad (7)$$

where ϵ is the $SU(2)$ antisymmetric tensor, $D_\mu \equiv \partial_\mu - igT_a W_\mu^a - ig'(Y/2)B_\mu$ is the covariant derivative leading to gauge interactions of the BSM multiplets with the $SU(2)_L$ and $U(1)_Y$ gauge bosons W_μ and B_μ , respectively, and T_a generically stand for the $SU(2)$ generators in the isospin 1/2 (doublet), 1 (triplet) and 3/2 (quadruplet) representations for H , Σ and Φ fields, respectively.

The Yukawa couplings Y_H and Y_Φ are 3×2 complex matrices with $\alpha, \beta = e, \mu, \tau$, and $i, j = 1, 2$, each characterized by 12 real parameters,

$$Y_\Phi = \begin{pmatrix} Y_{e1}^\Phi & Y_{e2}^\Phi \\ Y_{\mu1}^\Phi & Y_{\mu2}^\Phi \\ Y_{\tau1}^\Phi & Y_{\tau2}^\Phi \end{pmatrix}, \quad Y_H = \begin{pmatrix} Y_{e1}^H & Y_{e2}^H \\ Y_{\mu1}^H & Y_{\mu2}^H \\ Y_{\tau1}^H & Y_{\tau2}^H \end{pmatrix}. \quad (8)$$

The Dirac mass term for the triplet, M_Σ , is taken to be real and diagonal, without loss of generality, so $M_\Sigma = \text{Diag}(M_{\Sigma_1}, M_{\Sigma_2})$.

The scalar potential of the model can be written as

$$\begin{aligned} \mathcal{V}(H, \Phi) &= -\mu_H^2 (H^\dagger H) + \lambda_1 (H^\dagger H)^2 + \mu_\Phi^2 (\Phi^\dagger \Phi) + \lambda_2 (\Phi^\dagger \Phi)^2 \\ &\quad + \lambda_3 (H^\dagger H) (\Phi^\dagger \Phi) + \lambda_4 (H^\dagger \tau_a H) (\Phi^\dagger T_a \Phi) \\ &\quad + (\lambda_5 H^3 \Phi^* + \text{H.c.}), \end{aligned} \quad (9)$$

where we take all the parameters to be real. Here $\tau_a = \sigma_a/2$ (with σ_a being the standard Pauli matrices) and T_a stand for the $SU(2)$ generators in the isospin 1/2 and 3/2 representations, respectively.

Upon integrating out the Σ field in Eq. (7), a Weinberg-like dimension-5 operator is obtained,

$$\mathcal{L}_5 = -\frac{(Y_\Phi Y_H^\dagger + Y_H^* Y_\Phi^T) \bar{L} L^c H \Phi}{M_\Sigma} + \text{H.c.}, \quad (10)$$

where we have suppressed the generation indices. It can be seen that once the scalars take a VEV, a Majorana mass is generated for the neutrinos, $m_\nu \propto v_\Phi v_H$.

Before discussing the scalar and fermion sectors of the model in detail, we briefly comment on the source of LNV

in the model—a crucial ingredient for leptogenesis. Similar to the case of type-II seesaw, LNV in the model can be attributed to either the Yukawa terms or to the scalar potential, depending on the L charge assigned to Φ . First, let us consider the case where Φ carries a charge $L = -2$ units while $\Sigma_{L,R}$ is assigned -1 . In this case, it can be seen that the λ_5 term in Eq. (9) softly breaks the L number, while $\mathcal{L}_{\text{yuk}}$ preserves it [34]. In the limit, $\lambda_5 \rightarrow 0$, the symmetry of the scalar potential is restored, therefore, its value can be naturally small in the 't Hooft sense [13]. Alternatively, if Φ has $L = 0$ charge, and $\Sigma_{L,R}$ carry a unit lepton number, then either the Y_H or Y_Φ term in $\mathcal{L}_{\text{yuk}}$ violates L number by two units, and must be suppressed [46]. Since lepton number is violated by two units (owing to the Weinberg operator) in both the scenarios, one can adopt either assignment, as shown in Table I, to realize leptogenesis. We discuss the implications later.

A. Scalar sector

After EWSB, the Higgs VEV, $v_H \simeq 246.2$ GeV induces a nonzero VEV for the neutral component of Φ via the λ_5 term,

$$v_\Phi = -\frac{\lambda_5 v_H^3}{2M_{\Phi^0}^2} \simeq 0.74 \text{ GeV} \left(\frac{\lambda_5}{0.1} \right) \left(\frac{1 \text{ TeV}}{M_{\Phi^0}} \right)^2. \quad (11)$$

Since a scalar quadruplet obtaining a VEV breaks the custodial symmetry, this induced VEV is further constrained by modifications to the ρ parameter: $\rho \simeq (1 - 6v_\Phi^2/v_H^2)$. Using the experimental value of the ρ parameter [70], the VEV is constrained to be much smaller than the SM Higgs VEV,

$$v_\Phi < 2.6 \text{ GeV}, \quad (12)$$

which implies that $\lambda_5 \lesssim 0.35$ for $M_\Phi \sim \mathcal{O}(1)$ TeV.

Further constraints on the scalar quadruplet can be placed from its production and decays at the LHC. The effective operator in Eq. (10) leads to an interesting phenomenological signal, namely the decays of the doubly charged component $\Phi^{\pm\pm}$ into a pair of same-sign leptons ($\Phi^{\pm\pm} \rightarrow \ell^\pm \ell^\pm$), after EWSB. The strength of this interaction is proportional to the active neutrino mass matrix, analogous to the case of Type-II seesaw, and goes like v_Φ^{-1} . These decays to same-sign dileptons compete against decays to gauge bosons $\Phi^{\pm\pm} \rightarrow W^\pm W^\pm$, which are proportional to v_Φ . Therefore, depending on the value of v_Φ , the dominant decay channel is either leptonic ($v_\Phi \lesssim \mathcal{O}(100 \text{ keV})$) or bosonic ($v_\Phi \gtrsim \mathcal{O}(100 \text{ keV})$).⁴ Current searches for the decays of doubly charged scalars to

⁴There are also cascade decays of $\Phi^{\pm\pm}$ to other components of the quadruplet (and vice versa), and further sequential decays to leptons or gauge bosons. See Refs. [45,48] for further discussion.

multilepton final states [71,72] or to a pair of gauge bosons [73,74] at the LHC lead to the following bounds [45]:

$$M_{\Phi^{\pm\pm}} > \begin{cases} 860 \text{ GeV}, & \text{Br}(\ell^\pm \ell^\pm) = 100\%, \\ 260 \text{ GeV}, & \text{Br}(W^\pm W^\pm) = 100\%. \end{cases} \quad (13)$$

Since the mass splitting between the different components of the quadruplet is restricted to be $\Delta M \lesssim 50$ GeV from electroweak precision data [48], the lower limits mentioned above should generically apply to all components of Φ . To be safe from these constraints in general, we consider $M_\Phi \gtrsim 1$ TeV in our numerical analysis.

B. Fermion sector

The model contains three neutral left-handed fermion fields $\nu_L, \Sigma_L^0, \Sigma_R^{0,c}$, with the following mass matrix:

$$-\mathcal{L}_{\nu\Sigma^0} = \frac{1}{2} \begin{pmatrix} \overline{\nu_L^c} & \overline{\Sigma_L^{0,c}} & \overline{\Sigma_R^0} \end{pmatrix} \begin{pmatrix} 0 & \frac{v_H}{\sqrt{2}} Y_H & \frac{v_\Phi}{\sqrt{2}} Y_\Phi^* \\ \frac{v_H}{\sqrt{2}} Y_H^T & 0 & M_\Sigma \\ \frac{v_\Phi}{\sqrt{2}} Y_\Phi^\dagger & M_\Sigma & 0 \end{pmatrix} \begin{pmatrix} \nu_L \\ \Sigma_L^0 \\ \Sigma_R^{0,c} \end{pmatrix} + \text{H.c.} \quad (14)$$

If only one generation of vectorlike fermion is introduced, this is a 5×5 matrix with rank 4, leading to two active neutrino masses, which can be contrasted with the case of Type-III seesaw where at least two generations of fermion triplets are required for two nonzero neutrino masses. While a single Σ generation is sufficient to explain two nonzero active neutrino masses, the generation of CP asymmetry via Σ decay requires at least two generations of Σ , as we discuss below. Therefore, for our case, the resulting neutrino mass matrix is 7×7 which has rank 7 and leads to three nonzero light neutrino masses.

Similarly, the singly charged component of Σ mixes with the SM charged leptons. From Eq. (2), we see that Σ_L^+ is positively charged with left chirality and Σ_R^+ is positively charged with right chirality. Similarly, $(\Sigma_R^+)^c$ is negatively charged with left chirality, and $(\Sigma_L^+)^c$ is negatively charged with right chirality. Thus the Dirac mass matrix in the charged fermion sector connecting the left- and right-chiral fields is given by

$$-\mathcal{L}_{\ell\Sigma} = \begin{pmatrix} \overline{\ell_L^-} & \overline{(\Sigma_R^+)^c} \end{pmatrix} \begin{pmatrix} m_\ell & v_H Y_H^*/2 \\ 0 & M_\Sigma \end{pmatrix} \begin{pmatrix} \ell_R^- \\ (\Sigma_L^+)^c \end{pmatrix} + \text{H.c.} \quad (15)$$

The mass eigenstates can be found by diagonalizing the above mass matrix.

Both CMS and ATLAS collaborations at the LHC have put stringent limits on the mass of the triplet fermions by analyzing final states including production of up to four light leptons, and up to three hadronically decaying

τ leptons [75,76], leading to $M_\Sigma > 1065$ GeV. Therefore, the mixing between the SM leptons and triplets is expected to be suppressed by m_ℓ/M_Σ , and the weak states almost coincide with the mass eigenstates for the charged leptons.

Since we are interested in the decays of Σ for leptogenesis, we work with the following hierarchy: $M_\Sigma > M_\Phi \gg M_H$, which leads to the following decay channels: $\Sigma \rightarrow L\Phi, \bar{L}H$ and $\bar{\Sigma} \rightarrow \bar{L}\Phi^*, LH^*$. The total decay widths and branching ratios can be found in Appendix A.

III. NEUTRINO MASS

The effective operator in Eq. (10) generates tree-level masses for the light neutrinos after EWSB as follows:

$$\begin{aligned} m_\nu^{\text{tree}} &= -v_H v_\Phi (Y_\Phi M_\Sigma^{-1} Y_H^\dagger + Y_H^* M_\Sigma^{-1} Y_\Phi^T) \\ &= -\frac{\lambda_5 v_H^4}{2M_\Phi^2} (Y_\Phi M_\Sigma^{-1} Y_H^\dagger + Y_H^* M_\Sigma^{-1} Y_\Phi^T), \end{aligned} \quad (16)$$

where the second line is obtained after integrating Φ out [cf. Eq. (11)], resulting in a dimension-7 neutrino mass; see Fig. 1 left panel. Notice the additional suppression from λ_5 and M_Φ^2 , compared to the type-I seesaw where $m_\nu \simeq -v_H^2 Y M_N^{-1} Y^\dagger$.

Moreover, the same λ_5 term responsible for the induced VEV of Φ also allows us to write the dimension-5 Weinberg operator at one-loop level, see Fig. 1 right panel, with the loop contribution given by [77]

$$m_\nu^{\text{loop}} = \frac{\lambda_5 v_H^2}{4\pi^2} \frac{(Y_\Phi Y_H^\dagger + Y_H^* Y_\Phi^T) M_\Sigma}{M_\Phi^2 - M_H^2} f(M_\Phi, M_H, M_\Sigma), \quad (17)$$

where the loop function is given by

$$f(x, y, z) = \frac{x^2}{x^2 - z^2} \ln\left(\frac{x^2}{z^2}\right) - \frac{y^2}{y^2 - z^2} \ln\left(\frac{y^2}{z^2}\right). \quad (18)$$

The total active neutrino mass is given by the sum

$$m_\nu = m_\nu^{\text{tree}} + m_\nu^{\text{loop}}. \quad (19)$$

It was shown in Ref. [46] that the loop contribution can dominate over the tree-level one in a wide region of the

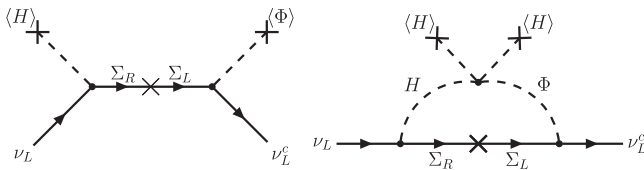


FIG. 1. Feynman diagrams for the tree and one-loop contributions to neutrino mass in the BNT model.

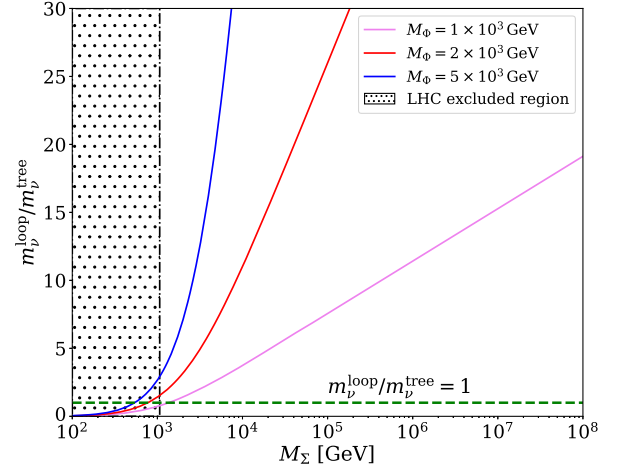


FIG. 2. Comparison of the loop (dimension-5) and tree (dimension-7) contributions to neutrino mass in the BNT model with respect to M_Σ for different values of M_Φ . The green dashed line represents $m_\nu^{\text{loop}}/m_\nu^{\text{tree}} = 1$, whereas the vertical shaded region is excluded from LHC searches for triplet fermions.

parameter space. In Fig. 2, we compare the loop and tree contributions to neutrino mass with respect to M_Σ for fixed values of M_Φ . It can be seen that for the allowed range of Σ masses above the TeV scale, the dimension-5 operator-induced loop contribution is dominant over the dimension-7 operator-induced tree contribution and the loop dominance increases with the higher values of M_Φ .

Next, we discuss the parameterization of the Yukawa matrices of Eq. (8) to satisfy the constraints from the low-energy neutrino sector, as the Yukawa values play an important role in determining the CP asymmetry produced in Σ decays.

Owing to two sets of distinct Yukawa matrices Y_Φ and Y_H in the BNT model, as opposed to a single Yukawa matrix in canonical seesaw, we adopt the generalized Casas-Ibarra (GCI) parametrization approach outlined in Ref. [78]. The neutrino masses in the model following the hierarchy $M_\Sigma > M_\Phi \gg M_H$ can be written as

$$m_\nu = \mathcal{Y}^T \mathcal{M} \mathcal{Y}, \quad (20)$$

where $\mathcal{Y} = (Y_\Phi^T, Y_H^T)^T$ is a 4×3 complex matrix [cf. Eq. (8)], and

$$\mathcal{M} \simeq \frac{\lambda_5 v_H^2}{4\pi^2 M_\Phi^2} \begin{pmatrix} 0 & 0 & \frac{f'_1}{M_{\Sigma_1}} & 0 \\ 0 & 0 & 0 & \frac{f'_2}{M_{\Sigma_2}} \\ \frac{f'_1}{M_{\Sigma_1}} & 0 & 0 & 0 \\ 0 & \frac{f'_2}{M_{\Sigma_2}} & 0 & 0 \end{pmatrix}, \quad (21)$$

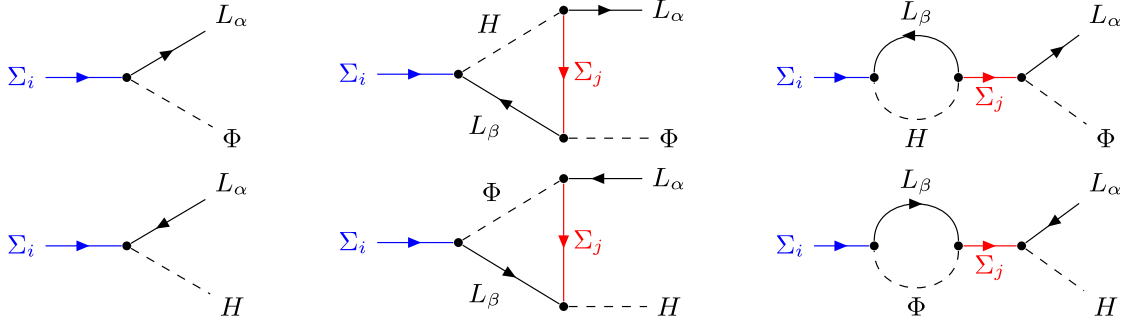


FIG. 3. Tree and one-loop diagrams for the decays of Σ to $L\Phi$ and $\bar{L}H$ that produce the CP asymmetry in the BNT model.

where

$$f'_i = M_\Phi^2 \ln(M_{\Sigma_i}^2/M_\Phi^2) - M_H^2 \ln(M_{\Sigma_i}^2/M_H^2), \quad (22)$$

where we have neglected the tree-level contribution to neutrino masses as they are subdominant (cf. Fig. 2). The Yukawa couplings can then be written as [78,79]

$$\mathcal{Y} = V^\dagger D_{\mathcal{M}}^{-1/2} \mathcal{R} D_m^{1/2} U^\dagger, \quad (23)$$

where V is the unitary matrix that diagonalizes $\mathcal{M}$ such that $\mathcal{M} = V^T D_{\mathcal{M}} V$, and is given by

$$V = \frac{1}{\sqrt{2}} \begin{pmatrix} i & 0 & -i & 0 \\ 0 & i & 0 & -i \\ 1 & 0 & 1 & 0 \\ 0 & 1 & 0 & 1 \end{pmatrix}, \quad (24)$$

U is the PMNS matrix diagonalizing the light neutrino mass matrix: $D_m = U^T m_\nu U$, and $\mathcal{R}$ is a complex semi-orthogonal 4×3 matrix

$$\mathcal{R} = \mathcal{R}'_{4 \times 4} \cdot \begin{pmatrix} \mathbb{I}_{3 \times 3} \\ 0_{1 \times 3} \end{pmatrix}, \quad (25)$$

where $\mathcal{R}' \in SO(4)$ and is characterized by six complex angles θ_{ij} .⁵

IV. LEPTOGENESIS

With the assumed hierarchy $M_\Sigma > M_\Phi \gg M_H$, the heavy fermion triplets can decay out-of-equilibrium to the SM leptons along with either the SM Higgs H or the quadruplet scalar Φ (and their conjugates) via the complex Yukawa interactions Y_H and Y_Φ , respectively. We define the yields

$$Y_i \equiv \frac{n_i}{s}, \quad Y_{\Delta i} \equiv Y_i - Y_{\bar{i}}, \quad (26)$$

⁵Some of these angles can be set to 0, for example, depending upon any imposed flavor symmetries or whether the lightest active neutrino remains massless. For instance, we find that fixing $\theta_{23}, \theta_{24} = 0$ enhances the CP asymmetry.

where n_i are the corresponding number densities, and $s = (2\pi^2/45)g_*T^3$ is the comoving entropy density of the Universe, with g_* being the effective relativistic degrees of freedom and T being the temperature.

If Φ carries a lepton number, these decays themselves do not violate lepton number. A nonvanishing CP asymmetry (ε_Σ) is nevertheless produced in their decays due to the interference of the tree and one-loop level diagrams shown in Fig. 3, leading to a net asymmetry in the lepton doublets ($Y_{\Delta L}$), as well as in the scalars ($Y_{\Delta\Phi}, Y_{\Delta H}$). This in turn induces an asymmetry among the triplets ($Y_{\Delta\Sigma}$). Because $M_\Phi \gg M_H$, the lepton number stored in the Φ sector is eventually washed out later via the $\Phi \rightarrow HHH$ decay sourced by the $\lambda_5 H^3 \Phi^*$ term in the scalar potential. On the other hand, when Φ does not carry a lepton number, either the decay channel $\Sigma \rightarrow \bar{L}H$ (and $\Sigma \rightarrow LH^*$) or $\Sigma \rightarrow L\Phi$ (and $\Sigma \rightarrow \bar{L}\Phi^*$) violates LN, while the decays to Φ conserve it. In either case, the $SU(2)_L$ sphalerons reprocess the asymmetry in the SM leptons to a baryon asymmetry

$$Y_{\Delta B} = 3c_{\text{sph}} Y_{\Delta L}, \quad (27)$$

where $c_{\text{sph}} = 0.461$ is the sphaleron conversion factor derived in Appendix B, and the factor of 3 comes from the three components of Σ . The per component lepton asymmetry can be expressed as

$$Y_{\Delta L} \simeq \varepsilon_\Sigma \eta Y_{\Sigma}^{\text{eq}}(T \gg M_\Sigma) = \varepsilon_\Sigma \eta \frac{135\zeta(3)}{2\pi^4 g_*}, \quad (28)$$

in terms of the CP asymmetry ε_Σ and efficiency parameter η such that $0 < \varepsilon_\Sigma, \eta \lesssim 1$. The η parameter incorporates the effects of washouts, and must be determined via the solution of the Boltzmann equations, with $\eta = 1$ implying maximal efficiency (negligible washouts). The equilibrium yield $Y_{\Sigma}^{\text{eq}} = n_{\Sigma}^{\text{eq}}/s$ is the total number of triplets ($n_\Sigma + n_{\bar{\Sigma}}$) per unit entropy, and $g_* = 114.75$ is the total number of relativistic degrees of freedom at high temperatures, obtained after accounting for the Φ multiplets. In order to match the present day value of $Y_{\Delta B}^0 = (8.58 \pm 0.17) \times 10^{-11}$ [70], we would thus require $Y_{\Delta L}^0 = (1.92 \pm 0.04) \times 10^{-10}$.

Similar to the case of the Type-III seesaw, these fermion triplets have gauge interactions that are in equilibrium for $T \gg M_\Sigma$, but get doubly Boltzmann suppressed at $T < M_\Sigma$. Unless the Yukawa interactions are strong enough [$Y_{H,\Phi} \gtrsim \mathcal{O}(0.01)$], these gauge annihilations typically dominate and deplete the Σ abundance. Subsequently, the decays of the remnant triplet population are not enough to generate the lepton asymmetry required to match the observed BAU.

As the Yukawa couplings in the BNT model can be much larger compared to the canonical seesaws, one can easily be in the Yukawa-dominated regime to reproduce the BAU for $M_\Sigma \lesssim 10^{10}$ GeV—below the DI-equivalent bound for Type-III seesaw leptogenesis. The consequence of having large Yukawas is analogous to Type-III seesaw leptogenesis with a second Higgs doublet with $v_{H_2} \ll v_\Phi$, where the lower bound on the Majorana triplet was found to be $\sim \mathcal{O}(10^7)$ GeV for a hierarchical spectrum of triplet masses [80].

A. CP Asymmetry

For a hierarchical spectrum, $M_{\Sigma_2} \gg M_{\Sigma_1}$, we consider only the decays of the lightest fermion triplet Σ_1 , as the asymmetry generated in Σ_2 decays is washed out by Σ_1 interactions. Therefore, we drop the subscript to simplify the notation: $\Sigma_1 \equiv \Sigma$. The CP asymmetry is generated via the interference of the tree-level and loop-level diagrams shown in Fig. 3. Notice that the scalar in the external state is different from the one in the internal loop due to the vectorlike nature of Σ .

The CP asymmetry for the two different decay channels is defined as

$$\begin{aligned} \varepsilon_H &= \frac{\sum_a [\Gamma(\Sigma \rightarrow \bar{L}H) - \Gamma(\bar{\Sigma} \rightarrow LH^*)]}{2\Gamma_\Sigma}, \\ \varepsilon_\Phi &= \frac{\sum_a [\Gamma(\Sigma \rightarrow L\Phi) - \Gamma(\bar{\Sigma} \rightarrow \bar{L}\Phi^*)]}{2\Gamma_\Sigma}, \end{aligned} \quad (29)$$

where $\Gamma_\Sigma = \Gamma_{\bar{\Sigma}}$ is the total decay width of $\Sigma, \bar{\Sigma}$ (hence the factor of 2 in the denominator). An explicit calculation leads to the following result:

$$\begin{aligned} \varepsilon_H &= \frac{1}{2\pi} \frac{\text{Im}[\mathbf{Y}]f_s(z_\Sigma, z_\Phi) + 2\text{Im}[\mathbf{Y}']f_v(z_\Sigma, z_\Phi)}{3(Y_H^\dagger Y_H)_{11} + 4(Y_\Phi^\dagger Y_\Phi)_{11}}, \\ \varepsilon_\Phi &= \frac{1}{2\pi} \frac{\text{Im}[\mathbf{Y}^\dagger]f_s(z_\Sigma, z_H) + 2\text{Im}[\mathbf{Y}']f_v(z_\Sigma, z_H)}{3(Y_H^\dagger Y_H)_{11} + 4(Y_\Phi^\dagger Y_\Phi)_{11}}, \end{aligned} \quad (30)$$

where $z_\Sigma = M_{\Sigma_2}^2/M_{\Sigma_1}^2$, $z_\Phi = M_\Phi^2/M_{\Sigma_1}^2$, $z_H = M_H^2/M_{\Sigma_1}^2$, and

$$\begin{aligned} \mathbf{Y} &= (Y_H^\dagger Y_H)_{12} (Y_\Phi^\dagger Y_\Phi)_{21}, \\ \mathbf{Y}' &= (Y_H^\dagger Y_\Phi^*)_{12} (Y_H^T Y_\Phi)_{21}. \end{aligned} \quad (31)$$

The loop functions $f_{s,v}$ corresponding to the contribution from the self-energy and vertex correction, respectively, are given by

$$\begin{aligned} f_v(x, y) &= \sqrt{x} \left[1 - y - (1 + x - y) \ln \left(1 + \frac{1}{x} - \frac{y}{x} \right) \right], \\ f_s(x, y) &= \frac{\sqrt{x}}{1-x} (1-y)^2. \end{aligned} \quad (32)$$

It is straightforward to see that $\varepsilon_\Phi = -\varepsilon_H$ in the limit $z_\Phi, z_H \rightarrow 0$, a consequence of CPT and unitarity [81,82]. The total CP asymmetry generated in Σ decays thus becomes

$$\varepsilon_\Sigma = \varepsilon_\Phi - \varepsilon_H = 2\varepsilon_\Phi. \quad (33)$$

The form of the products $\mathbf{Y}$ and $\mathbf{Y}'$ suggest that in the presence of a large hierarchy among the Yukawas, i.e., $Y_H \gg Y_\Phi$ or $Y_H \ll Y_\Phi$, the CP asymmetry gets suppressed.⁶ Instead, it becomes maximal when $Y_\Phi \simeq Y_H$, implying that $\text{Br}_\Phi \simeq 0.6$ and $\text{Br}_H \simeq 0.4$, see Appendix A.⁷

Now we discuss whether there exists a bound on the maximal CP asymmetry in the model, analogous to the DI bound for the Type-I seesaw [57]. Unlike the case of Type-I or Type-III seesaw, the CP asymmetry and neutrino masses involve different contractions of Yukawa matrices and a standard DI bound can not be defined.⁸ As the factors of λ_5 and loop suppression f' that suppress the neutrino masses allowing for $\mathcal{O}(1)$ Yukawas do not appear in the expressions for ε_Σ , the CP asymmetry can be as large as $\mathcal{O}(1)$. However, such large Yukawas would also enhance $\Delta L = 2$ scatterings such as $L\Phi \leftrightarrow \bar{L}H$ mediated by Σ_i that can wash out the produced lepton asymmetry. These scattering processes scale as

$$\Gamma_{\Delta L=2} \simeq \frac{|Y_H|^2 |Y_\Phi|^2}{32\pi M_\Sigma^2} T^3. \quad (34)$$

Demanding that these scatterings are out-of-equilibrium, $\Gamma_{\Delta L=2}/H < 1$ at $T \sim M_{\Sigma_1}$ gives us

$$|y| \lesssim 0.1 \left(\frac{M_{\Sigma_1}}{10^{12} \text{ GeV}} \right)^{1/4}, \quad (35)$$

where $|Y_\Phi| \sim |Y_H| \sim |y|$, and $H = 1.66\sqrt{g_*}T^2/M_{\text{Pl}}$ is the Hubble rate with $M_{\text{Pl}} = 1.22 \times 10^{19}$ GeV being the Planck scale.

⁶This is the case when Φ does not carry L and the lepton number is violated in either of the Yukawa couplings.

⁷If $Y_{\Phi ai} = Y_{H ai}$, the imaginary part of $\mathbf{Y}$ and $\mathbf{Y}'$ is zero, and the CP asymmetry vanishes.

⁸Similarly, we cannot relate the Γ_Σ to m_ν for the analysis as opposed to the case of Type-II and Type-III seesaw, see Ref. [60] for more details.

B. Boltzmann equations

The cosmological evolution of the asymmetries discussed above are determined by solving the following set of coupled Boltzmann equations (BEs)

$$sHx \frac{dY_\Sigma}{dx} = -2\gamma_\Sigma \left(\frac{Y_\Sigma}{Y_\Sigma^{\text{eq}}} - 1 \right) - 4\gamma_A \left(\frac{Y_\Sigma^2}{Y_\Sigma^{\text{eq}2}} - 1 \right), \quad (36)$$

$$sHx \frac{dY_{\Delta\Sigma}}{dx} = 2\gamma_\Sigma (\text{Br}_\Phi \mathcal{X}_\Phi - \text{Br}_H \mathcal{X}_H), \quad (37)$$

$$sHx \frac{dY_{\Delta\Phi}}{dx} = \gamma_\Sigma \left[\varepsilon_\Sigma \left(\frac{Y_\Sigma}{Y_\Sigma^{\text{eq}}} - 1 \right) - 2\text{Br}_\Phi \mathcal{X}_\Phi \right] - \gamma_{3H} \left(\frac{Y_{\Delta\Phi}}{Y_\Phi^{\text{eq}}} - \frac{3Y_{\Delta H}}{Y_H^{\text{eq}}} \right) - 2\gamma_{\text{sub}} (\mathcal{X}_\Phi + \mathcal{X}_H), \quad (38)$$

$$sHx \frac{dY_{\Delta H}}{dx} = -\gamma_\Sigma \left[\varepsilon_\Sigma \left(\frac{Y_\Sigma}{Y_\Sigma^{\text{eq}}} - 1 \right) - 2\text{Br}_H \mathcal{X}_H \right] + 3\gamma_{3H} \left(\frac{Y_{\Delta\Phi}}{Y_\Phi^{\text{eq}}} - \frac{3Y_{\Delta H}}{Y_H^{\text{eq}}} \right) + 2\gamma_{\text{sub}} (\mathcal{X}_\Phi + \mathcal{X}_H), \quad (39)$$

$$sHx \frac{dY_{\Delta L}}{dx} = 2\gamma_\Sigma \left[\varepsilon_\Sigma \left(\frac{Y_\Sigma}{Y_\Sigma^{\text{eq}}} - 1 \right) - \text{Br}_\Phi \mathcal{X}_\Phi - \text{Br}_H \mathcal{X}_H \right] - 4\gamma_{\text{sub}} (\mathcal{X}_\Phi + \mathcal{X}_H), \quad (40)$$

where $x \equiv M_{\Sigma_1}/T$ and

$$\mathcal{X}_\Phi = \frac{Y_{\Delta L}}{Y_L^{\text{eq}}} + \frac{Y_{\Delta\Phi}}{Y_\Phi^{\text{eq}}} - \frac{Y_{\Delta\Sigma}}{Y_\Sigma^{\text{eq}}}, \quad (41)$$

$$\mathcal{X}_H = \frac{Y_{\Delta\Sigma}}{Y_\Sigma^{\text{eq}}} + \frac{Y_{\Delta L}}{Y_L^{\text{eq}}} - \frac{Y_{\Delta H}}{Y_H^{\text{eq}}}. \quad (42)$$

The equilibrium yield is given by

$$Y_i^{\text{eq}} = \frac{45g_i}{4\pi^4 g_*} x_i^2 K_2(x_i), \quad (43)$$

where $x_i \equiv M_i/T$ and $K_n(x)$ is the modified Bessel function of the second kind of order n .

The first BE (36) takes into account the evolution of $\Sigma \equiv \Sigma_1 + \bar{\Sigma}_1$ which includes its decay, inverse decay and scattering processes. The rest of the BEs (37)–(40) track the evolution of asymmetries in different species. The decay and scattering rate densities are given by

$$\gamma_{a \rightarrow ij} = n_a^{\text{eq}} \frac{K_1(m_a/T)}{K_2(m_a/T)} \Gamma_{a \rightarrow ij},$$

$$\gamma_{ab \rightarrow ij} = \frac{T}{64\pi^4} \int_{s_{\min}}^\infty ds \sqrt{s} \hat{\sigma}(s) K_1\left(\frac{\sqrt{s}}{T}\right). \quad (44)$$

Here $\hat{\sigma}$ is the reduced cross section and $s_{\min} = (m_a + m_b)^2$. Analogous to the case of Type-II seesaw [68], not all of

these equations are independent as there exists a sum rule from hypercharge conservation, which leads to the following relation among the asymmetries in various species

$$2\Delta\Sigma + 3\Delta\Phi + \Delta H - \Delta L = 0. \quad (45)$$

Below, we discuss the effects of various terms included in the BEs above.

Gauge interactions— Due to their nonzero hypercharge, the fermion triplets in this model have stronger gauge interactions than that of Type-III seesaw. The gauge interactions include $\Sigma\bar{\Sigma} \leftrightarrow f\bar{f}, HH^*, GG'$ where $G(f)$ denotes gauge bosons (fermions). The total reduced cross section for these annihilations that tend to suppress the Σ population is given by [83]

$$\hat{\sigma}_A = \frac{3(2g^2 + g'^2)^2}{2\pi} \left[2\beta(\beta^2 - 2) - (\beta^4 - 3) \ln\left(\frac{1+\beta}{1-\beta}\right) \right], \quad (46)$$

with $\beta \equiv \sqrt{1 - 4/y}$ and $y \equiv s/M_\Sigma^2$. The reaction density can then be written in terms of $\hat{\sigma}_A$ as

$$\gamma_A = \frac{M_\Sigma^4}{64\pi^4} \int_{y_{\min}}^\infty dy \sqrt{y} \frac{K_1(x\sqrt{y}) \hat{\sigma}_A}{x}. \quad (47)$$

As mentioned above, these gauge scatterings/annihilations are initially not in thermal equilibrium for $T \gg M_\Sigma$, and are doubly Boltzmann suppressed at late times, $T \ll M_\Sigma$. The strength of these interactions is maximal at $T \sim M_\Sigma$. This behavior can be seen from the blue curves in Fig. 4. For the triplet with hypercharge 1, it can be checked that for $M_\Sigma < 10^{15}$ GeV, the triplets always thermalize irrespective of their initial abundance, hence, they play an important role in the evolution of Σ abundance. However, they do not appear in any of the asymmetry equations as they always involve pair annihilations.

Yukawa interactions. The effect of Yukawa interactions that leads to the decays of Σ is incorporated in ε_Σ , and γ_Σ which includes the total decay width of Σ . These interactions compete against the gauge interactions; if $\gamma_A/(n_\Sigma H) > 1$ and $\gamma_\Sigma < \gamma_A$, the gauge interactions dominate and the triplets annihilate before decaying, thus significantly suppressing the efficiency of leptogenesis. At $T < M_\Sigma$, the decays have a single Boltzmann suppression in contrast with the double suppression of the gauge interactions. Hence, if $\gamma_\Sigma \geq 4\gamma_A$, the decays dominate and the triplets will decay before annihilating, making leptogenesis more efficient. We find that $\gamma_\Sigma \geq 4\gamma_A$ around $x \equiv x_A \sim 10$ for $\mathcal{O}(0.1)$ Yukawa couplings.

Since the Yukawa couplings in the model are not suppressed by neutrino mass, and can take larger values than in canonical seesaws, it is easy to satisfy the condition $\gamma_\Sigma \geq 4\gamma_A$ for $|Y_H|, |Y_\Phi| \gtrsim 0.1$. In this case, the γ_A term can be neglected in the BEs, and the suppressed efficiency of leptogenesis is mainly due to the washout from inverse

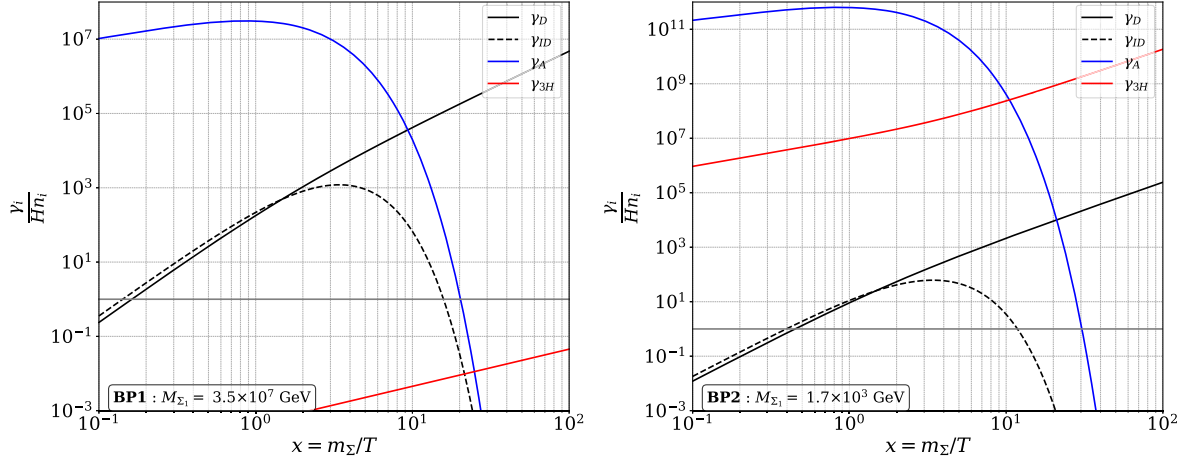


FIG. 4. Thermalization rates for gauge, Yukawa and scalar interactions relevant for the Boltzmann equations for two benchmark points indicated in Table III.

decays that scale as $|y|^2$, which become suppressed at $T < M_{\Sigma}$. The inverse decays are taken into account by the terms in the square brackets not multiplied by ε_{Σ} in Eqs. (38)–(40). The efficiency is enhanced if these inverse decays also decouple around x_A . The decay and inverse decay rates are shown by black solid and dashed curves, respectively, in Fig. 4 for our choice of benchmark points.

Next, we discuss the $\Delta L = 2$ scatterings involving the Yukawa interactions. The vectorlike nature of Σ allows only $\Delta L = 2$ scattering $L\Phi \leftrightarrow \bar{L}H$ and its conjugate process. The on-shell rate of these processes has already been accounted for via the inverse decays and decays in the BEs; therefore, only the on-shell or real intermediate state (RIS)-subtracted part of these rates needs to be tracked. We define $\gamma_{\text{sub}} = \gamma_s + \gamma_t$ corresponding to the s - and t -channel processes. The complete expressions for these can be found in Appendix A. These rates are not only suppressed by $|y|^4$, but also fall off at $x \gg 1$, therefore their effects are subdominant compared to the decays/inverse decays for $|y| \lesssim \mathcal{O}(0.1)$, and can be neglected for $M_{\Sigma} \gtrsim \mathcal{O}(\text{TeV})$, as discussed above.

Scalar interactions. The λ_5 term in Eq. (9) required for neutrino masses leads to the three-body decays $\Phi \rightarrow HHH$ as well as $2 \leftrightarrow 2$ scatterings such as $\Phi H^* \leftrightarrow HH$. The decay rate for $\Phi \rightarrow HHH$ in the limit $M_{\Phi} \gg M_H$ can be written as

$$\Gamma_{\Phi \rightarrow 3H} \simeq \frac{3\lambda_5^2}{256\pi^3} M_{\Phi}. \quad (48)$$

The reduced scattering cross section for the contact interaction in the limit $s \gg M_{\Phi}^2, M_H^2$ is given by

$$\hat{\sigma}_{\Phi H^* \rightarrow HH} \simeq \frac{9\lambda_5^2}{16\pi}. \quad (49)$$

Hence, we define

$$\begin{aligned} \gamma_{3H} &\equiv \gamma_{\Phi \rightarrow 3H} + \gamma_{\Phi H^* \rightarrow HH} \\ &= \frac{3\lambda_5^2 M_{\Phi}^4}{128\pi^5 x_{\Phi}^2} [x_{\Phi} K_1(x_{\Phi}) + 3K_2(x_{\Phi})]. \end{aligned} \quad (50)$$

The γ_{3H} term appears only in the BEs for $\Delta\Phi$ and ΔH , where the factor of 3 in the BE for ΔH appears as these processes change H number by three units. In Fig. 5, we show the region where these interactions are in equilibrium in the λ_5 and x_{Φ} plane for different values of M_{Φ} . The thermalization rate for these interactions is shown by red solid lines in Fig. 4.

V. RESULTS

The asymptotic lepton asymmetry generated in Σ decays can be determined by numerically solving the coupled BEs (36)–(40). Let us first discuss the implications for different asymmetry equations. If we ignore all the contributions from scalar interactions and the washout from inverse decays, the set can be simplified by ignoring the scalar asymmetries to only solve for Y_{Σ} , $Y_{\Delta\Sigma}$ and $Y_{\Delta L}$. In this case, it is easy to see that $Y_{\Delta L} = Y_{\Delta\Phi} - Y_{\Delta H}$, and $Y_{\Delta H} = -Y_{\Delta\Phi}$, as all these asymmetries have the same source.

Next, let us include the γ_{3H} term which ties the scalar asymmetries together. If these interactions are in equilibrium at $T \sim M_{\Sigma}$, then the asymmetry in Φ will be quickly washed out. Conversely, if they remain out of equilibrium at all times, they can be neglected. In Fig. 5, we see that these interactions are always in equilibrium at $T \sim M_{\Phi}$ for $\lambda_5 > 10^{-6}$; therefore, in order to track the lepton asymmetry correctly, we must carefully account for these interactions and track the scalar asymmetries.⁹

⁹Since the γ_{3H} interactions come into equilibrium already when the sphalerons are active, we include their effect to determine the modified sphaleron conversion factor in the model; see Appendix B.

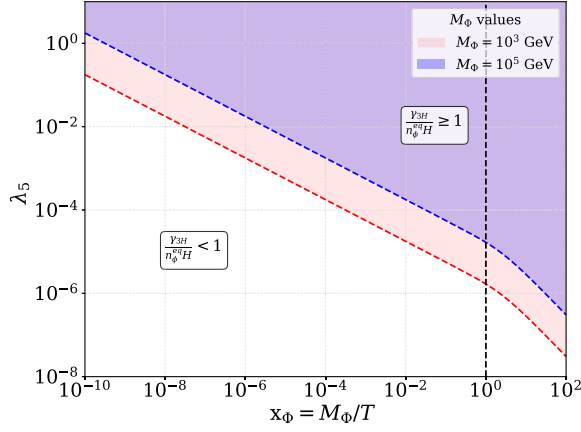


FIG. 5. Region in the (x_Φ, λ_5) plane where the scalar quartic interactions are in equilibrium (shaded regions) for different values of M_Φ . The vertical line is for $T = M_\Phi$.

Similarly, including the inverse decay contribution encoded in $\mathcal{X}_{\Phi,H}$ also relates the scalar asymmetries to the lepton asymmetry.

We find that including the γ_{3H} processes can change $Y_{\Delta L}$ by $\mathcal{O}(1)$ factors for the range of parameters we work with. This is due to the fact these interactions suppress the Φ population and thus modify the washout from inverse decays involving Φ . In other words, Φ produced in the decays of Σ can quickly decay to the SM Higgs or scatter before it can participate in the inverse decay contributing to washout, thus slightly enhancing the final lepton asymmetry.

In order to analyze the parameter space where leptogenesis is viable, we numerically solve the BEs for a random scan of the following set of free parameters:

TABLE II. Summary of the model parameters and their corresponding variation ranges explored in the numerical scans.

Parameters	Range
M_{Σ_1}	$[10^3, 10^{15}]$ GeV
r	$[2, 10]$
M_Φ	$[10^3, 10^6]$ GeV
λ_5	$[10^{-7}, 1]$
θ_{ij}	$[10^{-5}, 10^5] \times \exp(i[0, 2\pi])$

$$\{M_\Sigma, M_\Phi, \lambda_5, \theta_{ij}, r \equiv M_{\Sigma_2}/M_{\Sigma_1}\},$$

as the Yukawa couplings are already fixed by the GCI parametrization. The ranges over which the parameters of interest are scanned are given in Table II. We select those points that lead to $Y_{\Delta L}^0 = (1.92 \pm 0.04) \times 10^{-10}$. The parameter space where leptogenesis can work in the model is shown in Fig. 6. The scale of M_Σ above which leptogenesis becomes viable depends on the hierarchy among the triplet masses, we show this dependence in the right plot of Fig. 6. Note that for $r \simeq 1$, we enter the quasidegenerate regime which requires a different approach, as we discuss in the next Section.

For $r = 2$, we find that leptogenesis is viable for

$$M_{\Sigma_1} \gtrsim 3.5 \times 10^7 \text{ GeV}, \quad (51)$$

while for larger values of r , the lower bound goes up slightly, as shown in Fig. 6 right panel. A representative set of values of other parameters for this fixed value of M_Σ are shown in Table III, and we label the set as BP1. The corresponding Yukawa matrix structure comes out to be

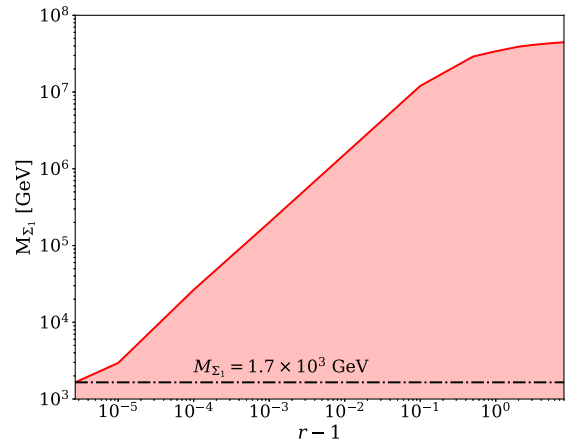
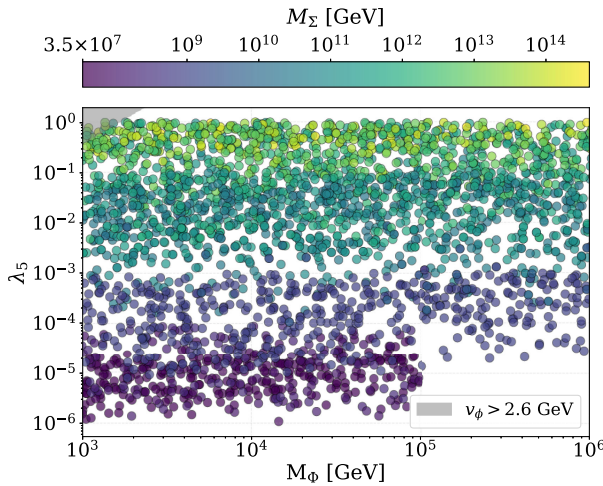


FIG. 6. *Left*: Viable parameter space for hierarchical leptogenesis in the BNT model in the (M_Φ, λ_5) plane, where each point corresponds to a different value of M_Σ as shown in the color bar. The gray region in the top-right corner is excluded from the constraints on the ρ parameter. *Right*: The dependence of the lower bound on M_Σ on the hierarchy among the triplet masses. Leptogenesis is not viable in the shaded region.

TABLE III. Set of benchmark points associated with the lowest possible scale for M_{Σ_1} above which one can successfully reproduce the baryon asymmetry via leptogenesis in the BNT model. BP1 and BP2 correspond to the hierarchical and resonant cases, respectively. The remaining angles θ_{23} and θ_{24} in the GCI parametrization are set to zero here.

	M_{Σ_1} (GeV)	r	M_Φ (GeV)	λ_5	v_Φ (GeV)	θ_{12}	θ_{13}	θ_{14}	θ_{34}
BP1	3.5×10^7	2	1000	10^{-5}	7.4×10^{-5}	$-0.003 - 0.005i$	0.001	$0.241 - 0.179i$	$-1.56 + 0.018i$
BP2	1.7×10^3	$\simeq 1$	1000	0.01	0.074	$-0.003 - 0.005i$	0.001	$0.241 - 0.179i$	$-1.56 + 0.018i$

$$Y_\Phi = \begin{pmatrix} -0.008 - 0.015i & -0.98 - 2.2i \\ 0.026 + 0.080i & 5.38 - 1.86i \\ 0.041 + 0.067i & 4.76 + 1.77i \end{pmatrix} \times 10^{-2},$$

$$Y_H = \begin{pmatrix} 0.002 + 0.020i & -0.98 - 1.2i \\ 0.052 - 0.062i & 5.55 - 1.86i \\ 0.036 - 0.063i & 4.91 + 1.77i \end{pmatrix} \times 10^{-2}. \quad (52)$$

First of all, we see that $Y_\Phi \sim Y_H$, consistent with the discussion above for maximizing the CP asymmetry. Second, the structure exhibits a strong hierarchy among the Yukawa couplings of Σ_1 and Σ_2 , $|Y_{i1}| \ll |Y_{i2}|$, implying that the inverse decays involving Σ_1 are further suppressed; this pattern is also seen in other low-scale leptogenesis scenarios [84]. Moreover, the similar values of the Yukawa couplings along with a small λ_5 for this BP suggest that giving Φ a lepton number, $L = -2$ is more natural in the hierarchical case.

We also show the evolution of abundances and asymmetries for BP1 in the left panel of Fig. 7 in order to understand the dynamics. It can be seen that the lepton and the scalar asymmetries (green, blue and violet curves) get dominantly produced at $x > 1$, once the decays of Σ go out of equilibrium and its abundance gets suppressed (black curve). The solid orange line shows the evolution of the asymmetry in Σ , with the lepton asymmetry freezing around the peak of the orange curve, after which the Σ

asymmetry decays away. On the other hand, the scalar interactions will deplete the asymmetry in Φ once they come into equilibrium, as seen from the downward gradient of the blue curve. Notice that $Y_{\Delta L}$ and $Y_{\Delta\Phi}$ curves are following each other closely, therefore, in order to satisfy Eq. (45), $Y_{\Delta H} = Y_H - Y_{H^*}$ comes out to be negative. Indeed, our numerical solution agrees with this, and we thus plot $|Y_{\Delta H}|$ in violet. Hence, for $x > 1$, the downward slope of the violet curve corresponds to the increase in the H population coming from $\Phi \rightarrow 3H$ decays.

VI. LOW-SCALE LEPTOGENESIS

The results from the hierarchical scenario show that leptogenesis is viable a few orders above the TeV scale, thus undermining the original motivation of the BNT model to realize a realistic neutrino mass generation mechanism along with interesting TeV-scale phenomenology. To alleviate this tension, we now consider the possibility of lowering the leptogenesis scale via resonant enhancement of the CP asymmetry by considering a quasidegenerate spectrum of Σ , similar to the resonant leptogenesis in canonical seesaw models [63,85]. For the resonant enhancement, we require

$$\Delta M_\Sigma \simeq \frac{\Gamma_{\Sigma_2}}{2} \ll M_{\Sigma_{1,2}}. \quad (53)$$

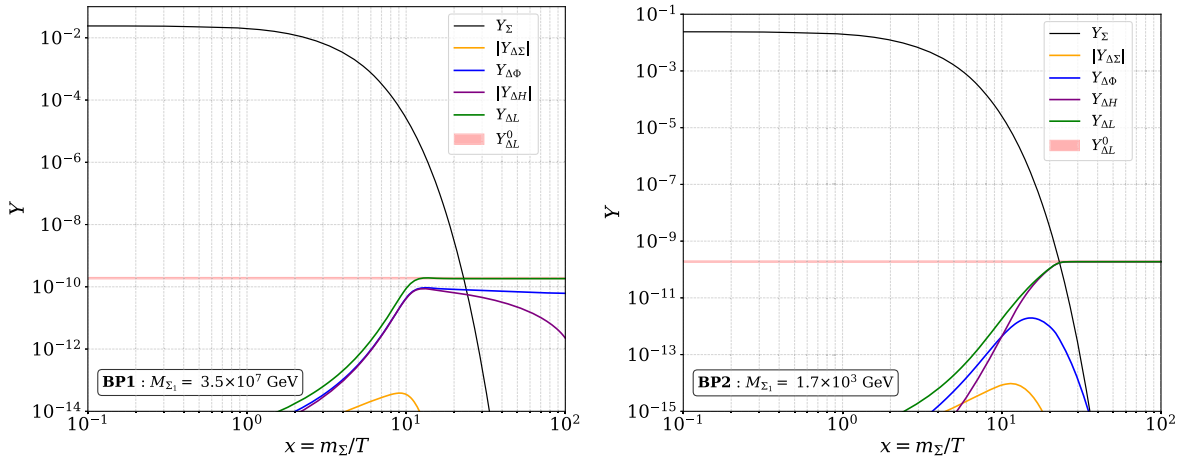


FIG. 7. Evolution of abundances and asymmetries in the hierarchical (left) and quasi-degenerate regime (right) for our two BP choices shown in Table III. The solid red line corresponds to the required asymmetry today $Y_{\Delta L}^0 = (1.92 \pm 0.04) \times 10^{-10}$.

It can be seen that in the limit $M_{\Sigma_1} \simeq M_{\Sigma_2}$, the self energy loop function in Eq. (32) diverges. Hence, the loop function must be regulated by including the Σ_2 decay width, so that we have [86]

$$f_s^{\text{reg}}(z_\Sigma, z_\Gamma) = \frac{(1 - z_\Sigma)z_\Gamma}{(1 - z_\Sigma)^2 + z_\Gamma^2}, \quad (54)$$

where $z_\Sigma = M_{\Sigma_2}^2/M_{\Sigma_1}^2$ and $z_\Gamma = \Gamma_{\Sigma_2}/M_{\Sigma_1}$. The CP asymmetry in the resonant scenario can be written as

$$\begin{aligned} \varepsilon_H &= \frac{1}{2048\pi^3} \frac{\text{Im}[\mathbf{Y}] f_s^{\text{reg}}(z_\Sigma, z_\Gamma) M_{\Sigma_1} M_{\Sigma_2}}{\Gamma_{\Sigma_1} \Gamma_{\Sigma_2}}, \\ \varepsilon_\Phi &= \frac{1}{2048\pi^3} \frac{\text{Im}[\mathbf{Y}^\dagger] f_s^{\text{reg}}(z_\Sigma, z_\Gamma) M_{\Sigma_1} M_{\Sigma_2}}{\Gamma_{\Sigma_1} \Gamma_{\Sigma_2}}, \end{aligned} \quad (55)$$

and can be as large as $\mathcal{O}(1)$. This enhancement thus allows us to have leptogenesis at the TeV scale. We find that for $r \simeq 1$, $M_{\Sigma_1} \geq 1.7$ TeV (see Fig. 6 right panel), similar to that for Type-II and Type-III seesaw leptogenesis [69], thus restoring the original motivation for the BNT model as providing a new neutrino mass generation mechanism at the TeV scale, while simultaneously explaining the BAU.

The representative values of other parameters corresponding to this value of M_Σ are labeled as BP2 in Table III. The corresponding Yukawa matrices are given by

$$\begin{aligned} Y_\Phi &= \begin{pmatrix} -0.013 - 0.024i & -1.13 - 2.52i \\ 0.041 + 0.13i & 6.22 - 2.16i \\ 0.064 + 0.11i & 5.51 + 2.04i \end{pmatrix} \times 10^{-5}, \\ Y_H &= \begin{pmatrix} 0.003 + 0.032i & -1.13 - 1.38i \\ 0.082 - 0.098i & 6.42 - 2.15i \\ 0.058 - 0.099i & 5.68 + 2.04i \end{pmatrix} \times 10^{-5}. \end{aligned} \quad (56)$$

Once again, the Yukawa matrices are of similar order $Y_H \sim Y_\Phi$ and $|Y_{i1}| \ll |Y_{i2}|$ like the hierarchical case above. However, notice that they are much smaller $\sim \mathcal{O}(10^{-5})$ compared to BP1. This decrease in Yukawas is instead compensated by a larger value of $\lambda_5 \simeq \mathcal{O}(10^{-2})$. Therefore, in this scenario, it is natural to shift LNV directly into the Yukawa sector instead of the scalar sector.

The evolution of the abundances and asymmetries corresponding to BP2 are shown in the right panel of Fig. 7, and can be contrasted with the hierarchical case shown in the left panel. The consequence of large λ_5 is immediately seen by the rise and decline of the asymmetry in Φ , while the lepton asymmetry freezes in. For BP2, as γ_{3H} is always in equilibrium during the evolution, $Y_{\Delta H}$ is positive in contrast to BP1. In addition, for $x > 1$, as $Y_{\Delta\Phi}$ quickly goes to zero, the populations of $Y_{\Delta H}$ and $Y_{\Delta L}$ become exactly the same, which once again validates Eq. (45).

Moreover, models with low-scale seesaws are in general subjected to constraints from lepton flavor violation, and in

this case also Higgs lepton flavor violating couplings, see Ref. [87]. This is attributed to the presence of a dimension-6 operator that is obtained after integrating the heavy fermions out, yielding a nonunitary PMNS matrix. Such constraints typically place limits on the Yukawa couplings of the form [77]

$$Y_{\beta i}^* Y_{\alpha i} \left(\frac{\text{TeV}}{M_\Sigma} \right)^2 \leq \begin{cases} 0.0004 & \text{for } \alpha\beta = \mu e \\ 0.24 & \text{for } \alpha\beta = \tau e \\ 0.28 & \text{for } \alpha\beta = \tau\mu, \end{cases} \quad (57)$$

which is easily satisfied by our choice of BP2. Further, the collider phenomenology for the TeV-scale BNT model has already been discussed in Ref. [45–50] which applies to our resonant case as well. We show the values of v_Φ for our BPs in Table III. In the hierarchical case, as λ_5 is suppressed, we get $v_\Phi = 74$ keV, implying that the dominant decay channel is to dileptons, whereas for the resonant case, we get $v_\Phi = 74$ MeV and the dominant decays will be to a pair of gauge bosons.

VII. CONCLUSIONS

In this work, we have studied leptogenesis in the Babu-Nandi-Tavartkiladze (BNT) model with fermion triplets and a scalar quadruplet. In this model, neutrino masses are generated at the tree level via a dimension-7 effective operator, as well as at the one-loop level via the dimension-5 Weinberg operator. Although this model was proposed many years ago and its collider phenomenology has been extensively studied, to the best of our knowledge, this is the first detailed study of leptogenesis in the BNT model and, more generally, in a framework where both fermionic and scalar multiplets larger than $SU(2)_L$ triplets participate in neutrino mass generation and leptogenesis.

The presence of vectorlike fermions and two independent Yukawa couplings makes this scenario distinct from canonical seesaw-based leptogenesis. Whereas in canonical Type-I leptogenesis the heavy Majorana mass provides the source of lepton-number violation, the original BNT realization contains vectorlike fermions with Dirac masses, while lepton number is broken by the scalar interaction. Moreover, the small induced vacuum expectation value of the scalar quadruplet allows significantly larger Yukawa couplings while reproducing the observed neutrino masses. This enhances the CP asymmetry and lowers the scale required for successful leptogenesis compared with canonical seesaw scenarios.

We also consider the alternative realization discussed in Ref. [46], in which lepton-number violation originates in the Yukawa sector. We show that both realizations can successfully account for the observed baryon asymmetry, with the Yukawa-induced realization being particularly well suited for resonant leptogenesis.

We find that for a hierarchical spectrum of triplet masses $M_{\Sigma_2} \geq 2M_{\Sigma_1}$, leptogenesis is viable for $M_{\Sigma_1} \gtrsim \mathcal{O}(10^7)$ GeV which, although lower than the DI bound of 10^9 GeV in canonical seesaw, is still beyond the experimental reach, thus undermining the original TeV-scale testability of the model. However, a quasidegenerate spectrum of fermion triplets allows us to resonantly enhance the asymmetry production and lowers the viable scale down to $\mathcal{O}(\text{TeV})$. While we do not consider flavor effects [88] in our analysis, which can indeed be important for the scales mentioned above, we expect the bound in the hierarchical case to be further lowered by about an order of magnitude, analogous to the case of Type-III leptogenesis with an additional Higgs doublet getting a small VEV [80]. On the other hand, the bound in the resonant case is extremely close to the absolute lower bound on leptogenesis with triplet fermions and scalars evaluated after considering Sommerfeld corrections [69], and is therefore not expected to be lowered further once the flavor effects are included.

The results of this work can be analogously applied to other neutrino mass models with higher-dimensional operators involving multiplets beyond the triplets. The viable leptogenesis scale could be similar or even smaller by a few orders of magnitude as they would be free from the constraints on the SM Higgs, unlike the BNT model. Future searches for multicharge fermions and scalars at the colliders as well as probes of lepton flavor violation at low scale can shed some light on the parameter space of these higher-dimensional operator-based neutrino mass models.

ACKNOWLEDGMENTS

We thank Kaladi Babu, Sudip Jana and Shaikh Saad for useful discussions. The work of B. D. was supported in part by the U.S. Department of Energy under Grant No. DE-SC0017987 and by a Humboldt Fellowship from the Alexander von Humboldt Foundation. S. G. acknowledges the J. C. Bose Fellowship (No. JCB/2020/000011) of the Anusandhan National Research Foundation, and the Department of Space, Government of India. The work of D. P. is funded by the Indian Space Research Organisation (ISRO) under Grant No. STC/PHY/2024427Q. D. V. is supported by the Canada First Research Excellence Fund through the Arthur B. McDonald Canadian Astroparticle Physics Research Institute, and a Subatomic Physics Discovery Grant (individual) from the Natural Sciences and Engineering Research Council of Canada. B. D. thanks the organizers of DTP-TALENT 2026 at ECT*, Trento for local hospitality during the final stages of this work.

DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

APPENDIX A: DECAY AND SCATTERING RATES

The total decay width of Σ can be written as

$$\Gamma_{\Sigma_i} = \Gamma(\Sigma_i \rightarrow \bar{L}_\alpha H) + \Gamma(\Sigma_i \rightarrow L_\alpha \Phi) = \frac{3}{32\pi} M_{\Sigma_i} \left[(Y_H^\dagger Y_H)_{ii} + \frac{4}{3} (Y_\Phi^\dagger Y_\Phi)_{ii} \right], \quad (\text{A1})$$

and $\Gamma_{\Sigma_i}^- = \Gamma_{\Sigma_i}$. The branching ratios are given by

$$\text{Br}_\Phi^i = \frac{4(Y_\Phi^\dagger Y_\Phi)_{ii}}{3(Y_H^\dagger Y_H)_{ii} + 4(Y_\Phi^\dagger Y_\Phi)_{ii}}, \quad \text{Br}_H^i = \frac{3(Y_H^\dagger Y_H)_{ii}}{3(Y_H^\dagger Y_H)_{ii} + 4(Y_\Phi^\dagger Y_\Phi)_{ii}}. \quad (\text{A2})$$

If the Yukawas are similar, the branching ratios become $\text{Br}_\Phi \simeq 0.6$ and $\text{Br}_H \simeq 0.4$.

The thermally averaged RIS-subtracted rate for $\Delta L = 2$ scattering process in the s channel is

$$\gamma_s(\bar{L}_\beta H \rightarrow L_\alpha \Phi) = \frac{3M_{\Sigma_1}^4}{1024\pi^5 x} |Y_{\Phi_{\beta 1}}|^2 |Y_{H_{\beta 1}}|^2 \int dy y^{3/2} \times \left[\frac{(y-1)^2 - a_\Gamma}{[(y-1)^2 + a_\Gamma]^2} \right] K_1(\sqrt{y}x), \quad (\text{A3})$$

where we have defined $y = s/M_\Sigma^2$ and $a_\Gamma = \Gamma^2/M_\Sigma^2$. On the other hand, for the t channel we have

$$\gamma_t(\bar{L}_\beta H \rightarrow L_\alpha \Phi) = \frac{3M_{\Sigma_1}^4}{512\pi^5 x} |Y_{\Phi_{\beta 1}}|^2 |Y_{H_{\alpha 1}}|^2 \int dy y^{-1/2} \times (y - \ln(1+y)) K_1(\sqrt{y}x). \quad (\text{A4})$$

APPENDIX B: SPHALERON REDISTRIBUTION COEFFICIENT

The asymmetry created in the standard model leptons can be redistributed among the other fermions and scalars via the Yukawa interactions that are in equilibrium. These processes also known as spectator processes can lead to the conservation laws for the individual number densities, n_x which for ultrarelativistic bosons and fermions is given by

$$n_x - n_{\bar{x}} = \frac{\mu_x g_x}{3} T^2 \begin{cases} \frac{1}{2} & \text{for fermions} \\ 1 & \text{for bosons} \end{cases}, \quad (\text{B1})$$

where μ_x and g_x stand for chemical potential and internal degrees of freedom of the species x , whereas for a massive particle the number density is expressed as

$$n_x - n_{\bar{x}} = \frac{\mu_x g_x}{\pi^2} T^2 F_\pm \left(\frac{m_x}{T} \right), \quad (\text{B2})$$

with

$$F_{\pm} = \int_x^{\infty} dy \frac{y \sqrt{y^2 - z^2} e^y}{(1 \pm e^y)^2}. \quad (\text{B3})$$

For our model the appropriate conservation laws read as,

(i) Hypercharge neutrality of the plasma,

$$3(\mu_q + 2\mu_u - \mu_d - \mu_{\ell} - \mu_e) + 2\mu_H + 12\mu_{\Phi} = 0. \quad (\text{B4})$$

(ii) $SU(2)_L$ sphalerons:

$$3\mu_q + \mu_{\ell} = 0 \text{ from } \mathcal{O}_{\text{sph}} = \Pi_{i=1}^3 \ell_i q_i q_i q_i. \quad (\text{B5})$$

(iii) $SU(3)_c$ sphalerons:

$$2\mu_q - \mu_u - \mu_d = 0. \quad (\text{B6})$$

(iv) Yukawa interactions:

$$\begin{aligned} \mu_{\ell} - \mu_H - \mu_e &= 0 \text{ from } \bar{L} H e_R, \\ \mu_q - \mu_H - \mu_d &= 0 \text{ from } \bar{Q}_L H d_R, \\ \mu_q + \mu_H - \mu_u &= 0 \text{ from } \bar{Q}_L \tilde{H} u_R. \end{aligned} \quad (\text{B7})$$

(v) Scalar interaction:

$$\mu_{\Phi} - 3\mu_H = 0 \text{ from } H^3 \Phi^*. \quad (\text{B8})$$

Here q, ℓ stand for the $SU(2)$ quark and lepton doublets, while u, d, e stand for the right-handed up, down and electron singlets in the SM. Solving the above system of equations and defining

$$\begin{aligned} \mu_B &= 3 \times (2\mu_q + \mu_u + \mu_d), \\ \mu_L &= 3 \times (2\mu_{\ell} + \mu_e) - 8\mu_{\Phi}, \end{aligned} \quad (\text{B9})$$

we find that

$$\mu_B = \frac{100}{217} \mu_{B-L}, \quad \text{or,} \quad \frac{\mu_B}{\mu_{B-L}} = c_{\text{sph}} \simeq 0.461. \quad (\text{B10})$$

This is the sphaleron conversion factor used in our calculation of the final baryon asymmetry.

-
- [1] S. Weinberg, Baryon and lepton nonconserving processes, *Phys. Rev. Lett.* **43**, 1566 (1979).
 - [2] P. Minkowski, $\mu \rightarrow e\gamma$ at a rate of one out of 10^9 muon decays?, *Phys. Lett.* **67B**, 421 (1977).
 - [3] R. N. Mohapatra and G. Senjanovic, Neutrino mass and spontaneous parity nonconservation, *Phys. Rev. Lett.* **44**, 912 (1980).
 - [4] T. Yanagida, Horizontal symmetry, and masses of neutrinos, *Prog. Theor. Phys.* **64**, 1103 (1980).
 - [5] M. Gell-Mann, P. Ramond, and R. Slansky, Complex spinors and unified theories, *Conf. Proc. C* **790927**, 315 (1979).
 - [6] W. Konetschny and W. Kummer, Nonconservation of total lepton number with scalar bosons, *Phys. Lett.* **70B**, 433 (1977).
 - [7] M. Magg and C. Wetterich, Neutrino mass problem and gauge hierarchy, *Phys. Lett.* **94B**, 61 (1980).
 - [8] J. Schechter and J. W. F. Valle, Neutrino masses in $SU(2) \times U(1)$ theories, *Phys. Rev. D* **22**, 2227 (1980).
 - [9] T. P. Cheng and L.-F. Li, Neutrino masses, mixings and oscillations in $SU(2) \times U(1)$ models of electroweak interactions, *Phys. Rev. D* **22**, 2860 (1980).
 - [10] G. Lazarides, Q. Shafi, and C. Wetterich, Proton lifetime and fermion masses in an $SO(10)$ model, *Nucl. Phys.* **B181**, 287 (1981).
 - [11] R. N. Mohapatra and G. Senjanovic, Neutrino masses and mixings in gauge models with spontaneous parity violation, *Phys. Rev. D* **23**, 165 (1981).
 - [12] R. Foot, H. Lew, X. G. He, and G. C. Joshi, Seesaw neutrino masses induced by a triplet of leptons, *Z. Phys. C* **44**, 441 (1989).
 - [13] G. 't Hooft, Naturalness, chiral symmetry, and spontaneous chiral symmetry breaking, *NATO Sci. Ser. B* **59**, 135 (1980).
 - [14] J. Kersten and A. Y. Smirnov, Right-handed neutrinos at CERN LHC and the mechanism of neutrino mass generation, *Phys. Rev. D* **76**, 073005 (2007).
 - [15] A. Ibarra, E. Molinaro, and S. T. Petcov, TeV scale see-saw mechanisms of neutrino mass generation, the Majorana nature of the heavy singlet neutrinos and $(\beta\beta)_{0\nu}$ -decay, *J. High Energy Phys.* **09** (2010) 108.
 - [16] R. N. Mohapatra and J. W. F. Valle, Neutrino mass and baryon number nonconservation in superstring models, *Phys. Rev. D* **34**, 1642 (1986).
 - [17] E. K. Akhmedov, M. Lindner, E. Schnapka, and J. W. F. Valle, Left-right symmetry breaking in NJL approach, *Phys. Lett. B* **368**, 270 (1996).
 - [18] E. K. Akhmedov, M. Lindner, E. Schnapka, and J. W. F. Valle, Dynamical left-right symmetry breaking, *Phys. Rev. D* **53**, 2752 (1996).
 - [19] M. Malinsky, J. C. Romao, and J. W. F. Valle, Novel supersymmetric $SO(10)$ seesaw mechanism, *Phys. Rev. Lett.* **95**, 161801 (2005).
 - [20] A. Zee, Quantum numbers of Majorana neutrino masses, *Nucl. Phys.* **B264**, 99 (1986).

- [21] K. S. Babu, Model of 'calculable' Majorana neutrino masses, *Phys. Lett. B* **203**, 132 (1988).
- [22] A. Pilaftsis, Radiatively induced neutrino masses, and large Higgs neutrino couplings in the standard model with Majorana fields, *Z. Phys. C* **55**, 275 (1992).
- [23] Z.-j. Tao, Radiative seesaw mechanism at weak scale, *Phys. Rev. D* **54**, 5693 (1996).
- [24] L. M. Krauss, S. Nasri, and M. Trodden, A model for neutrino masses and dark matter, *Phys. Rev. D* **67**, 085002 (2003).
- [25] E. Ma, Verifiable radiative seesaw mechanism of neutrino mass, and dark matter, *Phys. Rev. D* **73**, 077301 (2006).
- [26] P. S. B. Dev and A. Pilaftsis, Minimal radiative neutrino mass mechanism for inverse seesaw models, *Phys. Rev. D* **86**, 113001 (2012).
- [27] Y. Cai, J. Herrero-García, M. A. Schmidt, A. Vicente, and R. R. Volkas, From the trees to the forest: A review of radiative neutrino mass models, *Front. Phys.* **5**, 63 (2017).
- [28] C. Klein, M. Lindner, and S. Ohmer, Minimal radiative neutrino masses, *J. High Energy Phys.* **03** (2019) 018.
- [29] K. S. Babu, P. S. B. Dev, S. Jana, and A. Thapa, Non-standard interactions in radiative neutrino mass models, *J. High Energy Phys.* **03** (2020) 006.
- [30] F. F. Deppisch, P. S. B. Dev, and A. Pilaftsis, Neutrinos and collider physics, *New J. Phys.* **17**, 075019 (2015).
- [31] Y. Cai, T. Han, T. Li, and R. Ruiz, Lepton number violation: Seesaw models and their collider tests, *Front. Phys.* **6**, 40 (2018).
- [32] K. S. Babu and C. N. Leung, Classification of effective neutrino mass operators, *Nucl. Phys. B* **619**, 667 (2001).
- [33] A. de Gouvea and J. Jenkins, A survey of lepton number violation via effective operators, *Phys. Rev. D* **77**, 013008 (2008).
- [34] K. S. Babu, S. Nandi, and Z. Tavartkiladze, New mechanism for neutrino mass generation and triply charged higgs bosons at the LHC, *Phys. Rev. D* **80**, 071702 (2009).
- [35] F. Bonnet, D. Hernandez, T. Ota, and W. Winter, Neutrino masses from higher than $d = 5$ effective operators, *J. High Energy Phys.* **10** (2009) 076.
- [36] I. Picek and B. Radovic, Novel TeV-scale seesaw mechanism with Dirac mediators, *Phys. Lett. B* **687**, 338 (2010).
- [37] Y. Liao, Cascade seesaw for tiny neutrino mass, *J. High Energy Phys.* **06** (2011) 098.
- [38] K. Kumericki, I. Picek, and B. Radovic, Exotic seesaw-motivated heavy leptons at the LHC, *Phys. Rev. D* **84**, 093002 (2011).
- [39] K. Kumericki, I. Picek, and B. Radovic, TeV-scale seesaw with quintuplet fermions, *Phys. Rev. D* **86**, 013006 (2012).
- [40] K. L. McDonald, Minimal tree-level seesaws with a heavy intermediate fermion, *J. High Energy Phys.* **07** (2013) 020.
- [41] W. Wang and Z.-L. Han, Naturally small Dirac neutrino mass with intermediate $SU(2)_L$ multiplet fields, *J. High Energy Phys.* **04** (2017) 166.
- [42] R. Cepedello, M. Hirsch, and J. C. Helo, Loop neutrino masses from $d = 7$ operator, *J. High Energy Phys.* **07** (2017) 079.
- [43] G. Anamiati, O. Castillo-Felisola, R. M. Fonseca, J. C. Helo, and M. Hirsch, High-dimensional neutrino masses, *J. High Energy Phys.* **12** (2018) 066.
- [44] I. Doršner and S. Saad, Towards minimal $SU(5)$, *Phys. Rev. D* **101**, 015009 (2020).
- [45] A. Giarnetti, J. Herrero-Garcia, S. Marciano, D. Meloni, and D. Vatsyayan, Neutrino masses from new Weinberg-like operators: Phenomenology of TeV scalar multiplets, *J. High Energy Phys.* **05** (2024) 055.
- [46] G. Bambhaniya, J. Chakraborty, S. Goswami, and P. Konar, Generation of neutrino mass from new physics at TeV scale and multilepton signatures at the LHC, *Phys. Rev. D* **88**, 075006 (2013).
- [47] K. Ghosh, S. Jana, and S. Nandi, Neutrino mass generation at TeV scale and new physics signatures from charged Higgs at the LHC for photon initiated processes, *J. High Energy Phys.* **03** (2018) 180.
- [48] T. Ghosh, S. Jana, and S. Nandi, Neutrino mass from Higgs quadruplet and multicharged Higgs searches at the LHC, *Phys. Rev. D* **97**, 115037 (2018).
- [49] J. Pan, J.-H. Chen, X.-G. He, G. Li, and J.-Y. Su, Triply charged Higgs bosons at a 100 TeV pp collider, *Eur. Phys. J. C* **81**, 43 (2021).
- [50] A. Chakraborty, S. Chowdhury, N. Kumar, and V. Sahdev, Search for quadruplet scalars using boosted decision trees at the LHC, *Phys. Rev. D* **113**, 095017 (2026).
- [51] C. Arbeláez, R. Cepedello, R. M. Fonseca, and M. Hirsch, $(g - 2)$ anomalies and neutrino mass, *Phys. Rev. D* **102**, 075005 (2020).
- [52] S. Ashanujjaman and S. P. Maharathy, Vacuum structure of the Babu-Nandi-Tavartkiladze model of neutrino mass generation, *Phys. Rev. D* **113**, 075040 (2026).
- [53] M. Fukugita and T. Yanagida, Baryogenesis without grand unification, *Phys. Lett. B* **174**, 45 (1986).
- [54] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, On the anomalous electroweak baryon number nonconservation in the early universe, *Phys. Lett.* **155B**, 36 (1985).
- [55] K. Dick, M. Lindner, M. Ratz, and D. Wright, Leptogenesis with Dirac neutrinos, *Phys. Rev. Lett.* **84**, 4039 (2000).
- [56] H. Murayama and A. Pierce, Realistic Dirac leptogenesis, *Phys. Rev. Lett.* **89**, 271601 (2002).
- [57] S. Davidson and A. Ibarra, A lower bound on the right-handed neutrino mass from leptogenesis, *Phys. Lett. B* **535**, 25 (2002).
- [58] W. Buchmuller, P. Di Bari, and M. Plumacher, A bound on neutrino masses from baryogenesis, *Phys. Lett. B* **547**, 128 (2002).
- [59] W. Buchmuller, P. Di Bari, and M. Plumacher, Leptogenesis for pedestrians, *Ann. Phys. (Amsterdam)* **315**, 305 (2005).
- [60] T. Hambye, Leptogenesis: Beyond the minimal type I seesaw scenario, *New J. Phys.* **14**, 125014 (2012).
- [61] K. Moffat, S. Pascoli, S. T. Petcov, H. Schulz, and J. Turner, Three-flavored nonresonant leptogenesis at intermediate scales, *Phys. Rev. D* **98**, 015036 (2018).
- [62] E. K. Akhmedov, V. A. Rubakov, and A. Y. Smirnov, Baryogenesis via neutrino oscillations, *Phys. Rev. Lett.* **81**, 1359 (1998).
- [63] A. Pilaftsis and T. E. J. Underwood, Resonant leptogenesis, *Nucl. Phys. B* **692**, 303 (2004).
- [64] E. J. Chun *et al.*, Probing leptogenesis, *Int. J. Mod. Phys. A* **33**, 1842005 (2018).

- [65] T. Hambye, Y. Lin, A. Notari, M. Papucci, and A. Strumia, Constraints on neutrino masses from leptogenesis models, *Nucl. Phys.* **B695**, 169 (2004).
- [66] E. Ma and U. Sarkar, Neutrino masses and leptogenesis with heavy Higgs triplets, *Phys. Rev. Lett.* **80**, 5716 (1998).
- [67] T. Hambye and G. Senjanovic, Consequences of triplet seesaw for leptogenesis, *Phys. Lett. B* **582**, 73 (2004).
- [68] T. Hambye, M. Raidal, and A. Strumia, Efficiency and maximal CP-asymmetry of scalar triplet leptogenesis, *Phys. Lett. B* **632**, 667 (2006).
- [69] A. Strumia, Sommerfeld corrections to type-II, and III leptogenesis, *Nucl. Phys.* **B809**, 308 (2009).
- [70] Particle Data Group, Review of particle physics, *Phys. Rev. D* **110**, 030001 (2024).
- [71] CMS Collaboration, A search for doubly-charged Higgs boson production in three and four lepton final states at $\sqrt{s} = 13$ TeV, Technical Report No. CMS-PAS-HIG-16-036, CERN, 2017, <https://cds.cern.ch/record/2242956>.
- [72] ATLAS Collaboration, Search for doubly charged Higgs boson production in multi-lepton final states using 139 fb^{-1} of proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Eur. Phys. J. C* **83**, 605 (2023).
- [73] CMS Collaboration, Observation of electroweak production of same-sign W boson pairs in the two jet and two same-sign lepton final state in proton-proton collisions at $\sqrt{s} = 13$ TeV, *Phys. Rev. Lett.* **120**, 081801 (2018).
- [74] ATLAS Collaboration, Search for doubly and singly charged Higgs bosons decaying into vector bosons in multi-lepton final states with the ATLAS detector using proton-proton collisions at $\sqrt{s} = 13$ TeV, *J. High Energy Phys.* **06** (2021) 146.
- [75] ATLAS Collaboration, Search for type-III seesaw heavy leptons in leptonic final states in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Eur. Phys. J. C* **82**, 988 (2022).
- [76] CMS Collaboration, Inclusive nonresonant multilepton probes of new phenomena at $\sqrt{s} = 13$ TeV, *Phys. Rev. D* **105**, 112007 (2022).
- [77] A. Giarnetti, J. Herrero-García, S. Marciano, D. Meloni, and D. Vatsyayan, Neutrino masses from new seesaw models: Low-scale variants and phenomenological implications, *Eur. Phys. J. C* **84**, 803 (2024).
- [78] J. Herrero-García, S. Marciano, J. Racker, and D. Vatsyayan, Generalized Casas-Ibarra parametrization for Majorana neutrino masses, *Phys. Rev. D* **113**, 035035 (2026).
- [79] J. A. Casas and A. Ibarra, Oscillating neutrinos and $\mu \rightarrow e, \gamma$, *Nucl. Phys.* **B618**, 171 (2001).
- [80] D. Vatsyayan and S. Goswami, Lowering the scale of fermion triplet leptogenesis with two Higgs doublets, *Phys. Rev. D* **107**, 035014 (2023).
- [81] E. W. Kolb and S. Wolfram, Baryon number generation in the early universe, *Nucl. Phys.* **B172**, 224 (1980).
- [82] M. Berbig, S. M. A. S. H. E. D.: Standard model axion seesaw Higgs inflation extended for Dirac neutrinos, *J. Cosmol. Astropart. Phys.* **11** (2022) 042.
- [83] M. Cirelli, A. Strumia, and M. Tamburini, Cosmology and astrophysics of minimal dark matter, *Nucl. Phys.* **B787**, 152 (2007).
- [84] J. Racker, Low-scale leptogenesis in the scotogenic model: Spectator processes and benchmark points, *Phys. Rev. D* **111**, L081301 (2025).
- [85] P. S. B. Dev, M. Garny, J. Klaric, P. Millington, and D. Teresi, Resonant enhancement in leptogenesis, *Int. J. Mod. Phys. A* **33**, 1842003 (2018).
- [86] P. S. B. Dev, P. Millington, A. Pilaftsis, and D. Teresi, Flavour covariant transport equations: An application to resonant leptogenesis, *Nucl. Phys.* **B886**, 569 (2014).
- [87] J. Herrero-Garcia, N. Rius, and A. Santamaria, Higgs lepton flavour violation: UV completions and connection to neutrino masses, *J. High Energy Phys.* **11** (2016) 084.
- [88] P. S. B. Dev, P. Di Bari, B. Garbrecht, S. Lavignac, P. Millington, and D. Teresi, Flavor effects in leptogenesis, *Int. J. Mod. Phys. A* **33**, 1842001 (2018).