

Nonthermal leptogenesis in the $B-L$ symmetric SM with inverse seesaw mechanism

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We investigate the viability of nonthermal leptogenesis in the gauged $U(1)_{B-L}$ extension of the Standard Model (BLSM) with an inverse seesaw (ISS) mechanism for neutrino mass generation. In this framework, right-handed neutrinos typically have $\mathcal{O}(1)$ Yukawa couplings, which induce strong washout effects and render conventional thermal leptogenesis ineffective. We demonstrate that a successful baryogenesis scenario can nevertheless be realized through nonthermal leptogenesis, where right-handed neutrinos are produced from the decay of the heavy $B-L$ Higgs boson χ . We explicitly analyze the interplay between the dilution factor T_R/M_χ and the washout parameter characteristic of the ISS, highlighting the tension between suppressing washout effects and maintaining sufficient reheating. We show that a viable lepton asymmetry can be generated provided the scalar mass spectrum is appropriately tuned—a condition that requires a fine-tuning of order one part in 10^5 against one-loop radiative corrections—allowing for a reduced reheating temperature while keeping washout under control. The resulting lepton asymmetry is efficiently converted into the observed baryon asymmetry of the Universe via sphaleron processes. Our results establish that the inverse-seesaw $B-L$ model remains a predictive and robust framework for nonthermal leptogenesis and baryogenesis.

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I. INTRODUCTION

The origin of the baryon asymmetry of the Universe (BAU) remains one of the most fundamental open problems in particle physics and cosmology. Observations of the cosmic microwave background and primordial nucleosynthesis indicate a small but nonzero baryon-to-photon ratio [1,2],

$$\eta_B^{\text{obs}} \simeq 6 \times 10^{-10}, \quad (1)$$

which cannot be explained within the Standard Model (SM). Any successful mechanism for baryogenesis must satisfy the Sakharov conditions [3]: baryon number violation, C and CP violation, and a departure from thermal equilibrium. Among the various proposals, leptogenesis provides an attractive and economical framework in which a

lepton asymmetry is first generated and later converted into a baryon asymmetry through electroweak sphaleron processes.

In its conventional realization, thermal leptogenesis relies on the out-of-equilibrium decay of heavy Majorana neutrinos that are thermally produced in the early Universe. Successful baryogenesis in this framework typically requires heavy-neutrino masses well above the TeV scale, $M_N \gtrsim 10^9$ GeV [4], together with sufficiently small washout effects [5–7]. While theoretically appealing, such high scales are difficult to probe experimentally and are in tension with low-scale realizations of neutrino mass generation. This has motivated extensive efforts to construct leptogenesis scenarios that operate at or near the TeV scale, where a closer connection to neutrino physics and collider phenomenology can be established [8–13].

Nonthermal leptogenesis provides a well-studied alternative to the thermal paradigm [14]. In this class of scenarios, the heavy right-handed neutrinos are produced out of equilibrium from the decay of a heavier field, most notably the inflaton [15]. In inflaton-induced nonthermal leptogenesis, the inflaton couples directly to right-handed neutrinos and decays into them after inflation, generating a

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lepton asymmetry through their subsequent decays. This mechanism naturally avoids strong thermal washout effects and allows successful baryogenesis even for reheating temperatures well below the heavy-neutrino mass, $T_R \ll M_N$. Such scenarios have been widely investigated in a variety of contexts, including minimal seesaw models, supersymmetric frameworks, and resonant leptogenesis, and have been shown to reproduce the observed BAU under well-defined assumptions on the inflaton sector, its couplings, and the reheating dynamics [15–18].

While inflaton-decay scenarios provide a compelling proof of principle for nonthermal leptogenesis, their assumptions about the inflationary sector and its direct couplings to the neutrino sector are practically not constrained from a particle-physics perspective. So, the exploration of alternative nonthermal mechanisms in which the source of right-handed neutrinos arises from well-defined and testable degrees of freedom associated with extensions of the SM gauge structure is still an interesting questions.

A particularly well-motivated framework in this regard is the gauged $U(1)_{B-L}$ extension of the Standard Model. In this theory, the existence of three right-handed neutrinos is required by anomaly cancellation, and light neutrino masses can be generated naturally via the inverse seesaw (ISS) mechanism [19,20]. The ISS allows the heavy neutrinos to reside at the TeV scale while accommodating realistic light neutrino masses with moderately large Yukawa couplings. As a result, the $B-L$ inverse seesaw model is highly predictive and experimentally testable [21–24]. However, these same features lead to a severe cosmological challenge: the large Yukawa couplings induce extremely strong washout effects, rendering conventional thermal leptogenesis catastrophically ineffective, even when resonant CP enhancement is taken into account.

This observation strongly suggests that the origin of the lepton asymmetry in the $B-L$ inverse seesaw framework must be nonthermal. In this work, we show that a baryogenesis scenario can be realized through nonthermal leptogenesis driven by the decay of the scalar field responsible for $B-L$ symmetry breaking. In this setup, the heavy (pseudo-Dirac) neutrinos are produced out of equilibrium from the decay of the $B-L$ Higgs boson, while the reheating temperature remains below the heavy-neutrino mass, inducing an exponential suppression of inverse decays and washout processes.

We show that the interplay of three key ingredients allows the observed baryon asymmetry to be reproduced: (i) nonthermal production of heavy neutrinos from $B-L$ Higgs decay, (ii) a naturally suppressed reheating temperature achieved when the scalar mass lies close to the kinematic threshold for neutrino production, and (iii) resonant enhancement of the CP asymmetry enabled by the small lepton-number-violating parameter intrinsic to the inverse seesaw.

The paper is organized as follows. In Sec. II, we review the $U(1)_{B-L}$ inverse seesaw model and show the failure of thermal leptogenesis due to extreme washout. Section III presents the nonthermal leptogenesis mechanism arising from $B-L$ Higgs decay and discusses the parametric conditions required for suppressing washout while maintaining a sufficient heavy-neutrino yield. Sections IV and V analyze the washout regimes and resonant CP violation in detail. Our numerical results are presented in Sec. VI, and we summarize our conclusions in Sec. VII.

II. BLSM-ISS AND THE FAILURE OF THERMAL LEPTOGENESIS

The gauged $U(1)_{B-L}$ extension of the SM provides a minimal and well-motivated framework in which three right-handed neutrinos (RHN) arise naturally from anomaly cancellation. This setup allows also for a low-scale realization of the inverse seesaw (ISS) mechanism, where light neutrino masses are generated without requiring ultraheavy Majorana states. The leptonic sector relevant for neutrino mass generation is described by the Lagrangian [20]

$$\begin{aligned} \mathcal{L}_{B-L} \supset & i\bar{\ell}_L \gamma^\mu D_\mu \ell_L + i\bar{e}_R \gamma^\mu D_\mu e_R + i\bar{N}_R \gamma^\mu D_\mu N_R \\ & + i\bar{S}_{1,2} \gamma^\mu D_\mu S_{1,2} + (D_\mu \phi)^\dagger D^\mu \phi + (D_\mu \chi)^\dagger D^\mu \chi \\ & - V(\phi, \chi) - (\lambda_\nu \bar{\ell}_L \tilde{\phi} N_R + \lambda_N \bar{N}_R^c \chi S_2 + \text{H.c.}) \\ & - \frac{\alpha}{M^3} \bar{S}_1^c \chi^{\dagger 4} S_1 - \frac{\beta}{M^3} \bar{S}_2^c \chi^4 S_2, \end{aligned} \quad (2)$$

where ϕ is the SM Higgs doublet and χ is the scalar field responsible for $B-L$ symmetry breaking. The parameters α and β are free coupling constants. Furthermore, in order to forbid a potentially large mass term $\mu S_1 S_2$ in the above Lagrangian, we assume that the particles, N_R , χ , and S_2 are even under matter parity, while S_1 is odd. Finally, $V(\phi, \chi)$ is the most general Higgs potential invariant under these symmetries, which can be found in Ref. [21]. Radiative $B-L$ breaking in supersymmetric realizations typically yields $v' = \langle \chi \rangle \sim \mathcal{O}(\text{TeV})$ [22].

After symmetry breaking, the inverse-seesaw neutrino mass matrix in the (N_R, S_2) basis reads

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & M_N \\ M_N^T & \mu_S \end{pmatrix}, \quad M_N = \frac{\lambda_N v'}{\sqrt{2}}, \quad (3)$$

where $\mu_S = \beta v'^4 / M^3 \sim 10^{-9}$ GeV. Consequently, one obtains quasi-Dirac heavy neutrino states with masses

$$m_{\nu_{H,H'}} \simeq \pm M_N + \frac{\mu_S}{2}, \quad \nu_{H,H'} \simeq \frac{1}{\sqrt{2}} (\mp N_R + S_2). \quad (4)$$

The light neutrino mass matrix generated via the inverse seesaw takes the form

$$m_\nu \simeq \frac{(y_\nu v_{\text{SM}})^2}{M_N^2} \mu_S, \quad (5)$$

where $v_{\text{SM}} \simeq 174$ GeV. Reproducing the atmospheric neutrino mass scale $\sqrt{\Delta m_{\text{atm}}^2} \simeq 0.05$ eV implies

$$y_\nu \simeq \frac{M_N}{v_{\text{SM}}} \sqrt{\frac{0.05 \text{ eV}}{\mu_S}}, \quad (6)$$

showing that for TeV-scale M_N and small lepton-number-violating parameter μ_S , the Dirac Yukawa couplings are naturally of $\mathcal{O}(10^{-1} - 1)$.

A. CP asymmetry and resonant enhancement

The CP asymmetry generated in the decay of the lightest heavy neutrino ν_{H_1} arises from the interference between tree-level and one-loop diagrams [5,25,26]

$$\epsilon_{CP} \simeq \frac{1}{4\pi} \frac{\sum_{k=2,3} \text{Im}[(\lambda_\nu \lambda_\nu^\dagger)_{1k}] F(x_k)}{(\lambda_\nu \lambda_\nu^\dagger)_{11}}, \quad x_k = \frac{m_{\nu_{H_k}}^2}{m_{\nu_{H_1}}^2}, \quad (7)$$

with

$$F(x) = \sqrt{x} \left[1 + \frac{1}{1-x} - (1+x) \ln \frac{1+x}{x} \right]. \quad (8)$$

In pseudo-Dirac systems such as the inverse seesaw, the interference term driving the CP asymmetry can be highly suppressed if the flavor structure is aligned. However, by introducing complex phases in the orthogonal matrix of the Casas-Ibarra parametrization, these catastrophic cancellations are avoided. We provide a concrete flavor realization and the explicit calculation of our benchmark asymmetry ($\epsilon_{CP} \simeq 0.48$) in the Appendix, ensuring full compatibility with low-energy neutrino oscillation data.

In the ISS, the heavy neutrinos form pseudo-Dirac pairs with a small mass splitting governed by μ_S ,

$$\Delta M \simeq \mu_S. \quad (9)$$

Resonant enhancement of CP violation occurs when

$$\Delta M \sim \frac{\Gamma}{2}, \quad (10)$$

in which case the CP asymmetry can reach its theoretical maximum, $\epsilon_{CP} \lesssim 0.5$ [9,26–30]. However, due to the pseudo-Dirac structure, the asymmetries generated by the two components partially cancel, making the net enhancement highly model-dependent.

B. Strong washout and breakdown of thermal leptogenesis

The efficiency of thermal leptogenesis is governed by the decay parameter

$$K = \frac{\Gamma_{D_1}}{H} \Big|_{T=M_N}, \quad Y_B \simeq -1.4 \times 10^{-3} \kappa(K) \epsilon, \quad (11)$$

where $\kappa(K)$ encodes washout effects [31]. The Hubble rate in standard cosmology is

$$H \simeq 1.66 \sqrt{g_*} \frac{T^2}{M_{\text{Pl}}}, \quad g_* \simeq \mathcal{O}(10^2). \quad (12)$$

For TeV-scale heavy neutrinos with $y_\nu = \mathcal{O}(1)$, one finds

$$H \sim 10^{-12} \text{ GeV}, \quad K \sim 10^{12} - 10^{14}, \quad (13)$$

placing the system deep in the strong-washout regime. In this limit:

- (i) Inverse decays ($LH \rightarrow N$) remain fully active in the thermal bath.
- (ii) The efficiency factor is exponentially suppressed, $\kappa(K) \sim 10^{-15}$.
- (iii) Even with maximal resonant CP asymmetry, $\epsilon_{CP} \sim 0.5$, the resulting baryon asymmetry is $\eta_B \sim 10^{-16}$, far below the observed value $\eta_B^{\text{obs}} \simeq 6 \times 10^{-10}$.

We therefore conclude that *thermal leptogenesis is catastrophically ineffective in the $B - L$ inverse seesaw model*. This failure is driven by the unavoidable combination of large Yukawa couplings and extreme washout, independently of possible resonant CP enhancement.

This observation motivates nonthermal leptogenesis scenarios, in which right-handed neutrinos are produced out of equilibrium [32–35], through the decay of the heavy $B - L$ Higgs scalar.

III. NONTHERMAL LEPTOGENESIS FROM $B - L$ HIGGS DECAY

As shown in the previous section, *thermal* leptogenesis is catastrophically ineffective in the inverse-seesaw (ISS) $B - L$ framework: the Yukawa couplings required to reproduce m_ν at the TeV scale are typically $y_\nu = \mathcal{O}(10^{-1} - 1)$, implying an enormous washout parameter $K \sim 10^{12} - 10^{14}$ and hence an exponentially suppressed efficiency. This motivates *nonthermal leptogenesis*, where the heavy (pseudo-Dirac) neutrinos are produced out of equilibrium through the decay of the $B - L$ breaking scalar χ , while the reheating temperature satisfies $T_R < M_N$ so that inverse decays are Boltzmann suppressed.

A. Basic idea and parametric requirements

Nonthermal leptogenesis avoids (or strongly reduces) the thermal problems by engineering the following conditions:

- (i) *Low reheating*, $T_R < M_N$: the Universe reheats to a temperature below the heavy-neutrino mass, so that $LH \rightarrow N$ inverse decays and related washout

processes are kinematically disfavored and Boltzmann suppressed.

- (ii) *Controlled washout*, $K_{\text{eff}} \ll K$: the effective washout is reduced relative to the thermal value by a factor as $\exp(-M_N/T_R)$, protecting the generated asymmetry.
- (iii) *Successful yield and asymmetry*: combining an enhanced CP asymmetry (e.g., resonant $\epsilon_{CP} \sim 0.1\text{--}0.3$) with a protected efficiency κ and a sizable nonthermal RHN yield can reproduce the observed baryon asymmetry $\eta_B^{\text{obs}} \simeq 6 \times 10^{-10}$.

B. Reheating temperature and nonthermal RHN yield

We assume that after inflation the energy density is dominated by coherent oscillations of the scalar field χ (matter-dominated era) and that when $H \simeq \Gamma_\chi$, where Γ_χ is the total decay width of χ , reheating is taking place. The reheating temperature is then [36–38]

$$T_R \simeq \left(\frac{90}{8\pi^3 g_*(T_R)} \right)^{1/4} \sqrt{\Gamma_\chi M_{\text{Pl}}}, \quad (14)$$

with $g_*(T_R) \sim \mathcal{O}(10^2)$. If a fraction of χ decays produces heavy neutrinos, $\chi \rightarrow NN$ (here N collectively denotes the heavy ISS states), the RHN yield at reheating is approximately

$$Y_N^{\text{non-th}} \equiv \frac{n_N}{s} \simeq \frac{3}{2} \text{Br}(\chi \rightarrow NN) \frac{T_R}{M_\chi}, \quad (15)$$

where s is the entropy density and $\text{Br}(\chi \rightarrow NN)$ is the branching ratio into heavy neutrinos. Equation (15) makes explicit the *dilution* factor T_R/M_χ discussed previously: lowering T_R helps washout, but it also reduces the injected RHN abundance.

C. Threshold kinematics and the reheating temperature

In our model, the heavy pseudo-Dirac neutrinos are produced out of equilibrium via the decays of the massive $B-L$ scalar field, $\chi \rightarrow NN$. To ensure that the generated asymmetry is not subsequently erased by inverse decays, the universe must never reach temperatures comparable to the heavy neutrino mass ($T_{\text{max}} \sim T_R \ll M_N$).

This low-temperature regime is achieved by adopting a near-threshold kinematic configuration. We assume the scalar mass M_χ is just slightly above the kinematic threshold for producing two heavy neutrinos

$$M_\chi = 2M_N + \delta, \quad \text{with } \delta \ll M_N. \quad (16)$$

To see how this suppresses the energy injection into the thermal bath, we must evaluate the tree-level decay width of the scalar into the heavy neutrinos. For a scalar decaying

into two fermions of mass M_N , the width is governed by the phase space velocity β

$$\Gamma(\chi \rightarrow NN) = \frac{Y_N^2}{8\pi} M_\chi \beta^3, \quad \text{where } \beta = \sqrt{1 - \frac{4M_N^2}{M_\chi^2}}. \quad (17)$$

By inserting the threshold condition $M_\chi = 2M_N + \delta$, we can expand the phase space parameter β in the limit $\delta \ll M_N$

$$\beta = \sqrt{1 - \frac{4M_N^2}{(2M_N + \delta)^2}} \simeq \sqrt{1 - \frac{4M_N^2}{4M_N^2(1 + \delta/M_N)}} \simeq \sqrt{\frac{\delta}{M_N}}. \quad (18)$$

Substituting this back into the decay width, and approximating $M_\chi \simeq 2M_N$, the partial width is suppressed by the small mass splitting

$$\Gamma(\chi \rightarrow NN) \simeq \frac{Y_N^2}{4\pi} M_N \left(\frac{\delta}{M_N} \right)^{3/2}. \quad (19)$$

To determine the total decay width of the scalar, Γ_χ , we must consider the full scalar sector. The physical $B-L$ scalar χ can, in principle, decay into Standard Model (SM) particles (e.g., W^+W^- , ZZ , hh , $f\bar{f}$) through its mixing with the SM Higgs doublet. Let this mixing angle be α . However, precision measurements of the 125 GeV Higgs boson properties at the LHC strictly require this mixing to be highly suppressed, typically $\sin \alpha \lesssim 0.1$. Because the decay widths of χ into SM states are proportional to $\sin^2 \alpha$, these channels are heavily penalized.

Consequently, the decay channel into the heavy neutrinos, governed by the unsuppressed $U(1)_{B-L}$ coupling Y_N , dominates. It is therefore a correct assumption to parametrize the branching ratio into heavy neutrinos as an $\mathcal{O}(1)$ fraction of the total χ decay rate.

Finally, the total decay width dictates the reheating temperature T_R of the universe. Approximating the onset of radiation domination as the moment when the Hubble parameter equals the scalar decay rate ($H \simeq \Gamma_\chi$), and using the standard radiation density relation $H = \sqrt{\pi^2 g_*/90} T^2/M_{\text{Pl}}$, the reheating temperature is given by

$$T_R = \left(\frac{90}{8\pi^3 g_*} \right)^{1/4} \sqrt{\Gamma_\chi M_{\text{Pl}}}. \quad (20)$$

Because Γ_χ is strongly suppressed by the factor $(\delta/M_N)^{3/2}$, the resulting reheating temperature is naturally driven far below the heavy neutrino mass scale, $T_R \ll M_N$, dynamically shutting off the thermal washout.

This mechanism keeps T_R strictly below M_N , which is precisely what ensures an exponential reduction of washout. It is also a strong signature of the nonthermal

leptogenesis scenario. In this scenario, the scalar mass spectrum is needed to be appropriately tuned allowing for a reduced reheating temperature while keeping washout under control and keeping the advantage of the ISS $B - L$ model.

We note that maintaining the condition $\delta \ll M_N$ requires parameter tuning against radiative corrections. The one-loop correction to the scalar mass from the heavy neutrinos is $\Delta M_\chi \simeq (Y_N^2/64\pi^2)M_N$. For our benchmark parameters yielding $T_R \approx 1$ TeV ($M_N = 5$ TeV, $Y_N \simeq 0.1$, $\delta = 10^{-6}$ GeV), the radiative shift is of order 10^{-2} GeV. So, a fine-tuning of ~ 1 part in 10^5 between the bare scalar mass and the one-loop corrections is required. This tuning is acceptable for exploring the viability of the mechanism.

D. Washout: From K to an effective K_{eff}

For reference, the decay width of a heavy neutrino into a lepton and Higgs, $N \rightarrow LH$, is

$$\Gamma_D \simeq \frac{1}{8\pi} (y_\nu^\dagger y_\nu) M_N, \quad (21)$$

and the Hubble rate in radiation domination is

$$H(T) = 1.66 \sqrt{g_*} \frac{T^2}{M_{\text{Pl}}}. \quad (22)$$

In thermal leptogenesis one evaluates $K = \Gamma_D/H$ at $T = M_N$, obtaining

$$K \equiv \frac{\Gamma_D}{H} \Big|_{T=M_N} \simeq \frac{(y_\nu^\dagger y_\nu) M_{\text{Pl}}}{8\pi \cdot 1.66 \sqrt{g_*} M_N}, \quad (23)$$

which becomes enormous for $y_\nu = \mathcal{O}(10^{-1} - 1)$ and $M_N \sim \text{TeV}$.

In the nonthermal case, however, the relevant washout operates around $T \simeq T_R < M_N$. The inverse decays are then Boltzmann suppressed, and a useful parametric estimate is [18]

$$K_{\text{eff}} \simeq K \exp\left(-\frac{M_N}{T_R}\right). \quad (24)$$

The reheating temperature inherits the strong phase-space suppression from Γ_χ ,

$$T_R(\delta) \propto \sqrt{\Gamma_\chi} \propto Y_N \beta^{3/2} \sqrt{M_\chi M_{\text{Pl}}} \sim Y_N \left(\frac{\delta}{M_N}\right)^{3/4} \sqrt{M_N M_{\text{Pl}}}, \quad (25)$$

E. Baryon asymmetry and the dilution-washout tension

In our scenario, the baryon asymmetry η_B can be written as

$$\eta_B \simeq d \epsilon_{CP} \kappa (K_{\text{eff}}) Y_N^{\text{non-th}}, \quad (26)$$

where:

- (i) d is the sphaleron conversion factor (numerically $d \simeq 0.35$), encoding the conversion of $B - L$ into B by electroweak sphalerons,
- (ii) ϵ_{CP} is the CP asymmetry per heavy-neutrino decay (in the ISS it can be resonantly enhanced when the pseudo-Dirac splitting is comparable to the decay width),
- (iii) $\kappa(K_{\text{eff}})$ is an efficiency factor controlled by washout, and is strongly correlated with K_{eff} , and
- (iv) $Y_N^{\text{non-th}}$ is the nonthermal RHN yield of Eq. (15).

Substituting Eq. (15) into Eq. (26) gives

$$\eta_B \simeq d \epsilon_{CP} \kappa \left(\frac{3}{2} \text{Br}(\chi \rightarrow NN) \frac{T_R}{M_\chi} \right). \quad (27)$$

Increasing T_R boosts production linearly via T_R/M_χ , but also increases washout through K_{eff} (roughly exponentially sensitive to T_R). Taking $M_\chi = 2M_N + \delta$ allows T_R to be small enough to suppress washout, while keeping $\text{Br}(\chi \rightarrow NN)$ sizable.

F. Competing decay channels and the Higgs-portal constraint

A successful nonthermal scenario also requires that χ decays mainly into heavy neutrinos. Since χ is a Higgs-like scalar in an extended sector, it can couple to the SM Higgs through a portal interaction, which induces $\chi \rightarrow hh$ decays after symmetry breaking. Parametrically one may write

$$\Gamma(\chi \rightarrow hh) \sim \frac{\lambda_{\chi H}^2 v'^2}{32\pi M_\chi}, \quad (28)$$

so that the branching ratio into neutrinos is

$$\text{Br}(\chi \rightarrow NN) = \frac{\Gamma(\chi \rightarrow NN)}{\Gamma(\chi \rightarrow NN) + \Gamma(\chi \rightarrow hh) + \dots}. \quad (29)$$

In the threshold regime, $\Gamma(\chi \rightarrow NN)$ is suppressed by β^3 , hence $\lambda_{\chi H}$ must be sufficiently small to prevent $\chi \rightarrow hh$ from dominating and driving $\text{Br}(\chi \rightarrow NN) \rightarrow 0$ [39,40]. This requirement is essential for maintaining a large $Y_N^{\text{non-th}}$ and, consequently, a viable η_B .

G. Benchmark illustration and figures

In our numerical analysis we fix, for illustration, $M_N = 5$ TeV and use the threshold relation $M_\chi \simeq 2M_N + \delta$ to obtain $T_R \sim 1$ TeV for correlated values of (Y_N, δ) . Figure 1 shows the parameter combinations that satisfy $T_R = 1$ TeV, while Fig. 2 displays the corresponding region where the correct baryon asymmetry is achieved once the resonant condition for CP enhancement is imposed.

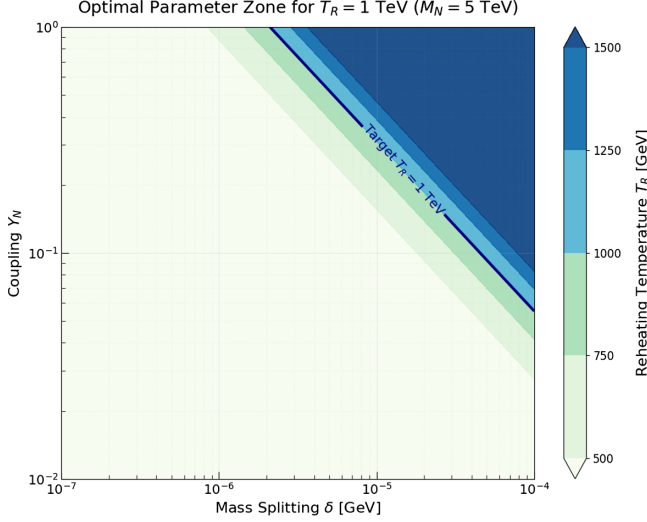


FIG. 1. Dynamical reheating temperature contours. Contour plot displaying the reheating temperature (T_R) as a function of the scalar mass splitting (δ) and the scalar-neutrino coupling (Y_N). The scanned parameter space spans $\delta \in [10^{-7}, 10^{-4}]$ GeV and $Y_N \in [10^{-2}, 1]$. The remaining parameters are fixed to the benchmark values: $M_N = 5$ TeV, $\text{Br}(\chi \rightarrow NN) = 0.5$, and $y_\nu = 10^{-2}$. The T_R values in this plot are computed dynamically by identifying the exact moment of radiation domination ($\rho_R = \rho_\chi$) via the full numerical integration of the coupled Boltzmann equations presented in Sec. IV. The solid navy line highlights the optimal target zone corresponding to $T_R = 1$ TeV, demonstrating that a near-threshold configuration ($\delta \sim 10^{-6}$ GeV) and a moderately large coupling ($Y_N \sim 0.1$) yield the low reheating temperatures required to avoid thermal washout.

In summary, nonthermal leptogenesis from χ decay can overcome the thermal catastrophe of the ISS $B-L$ model by enforcing $T_R < M_N$ (thereby suppressing washout) while retaining a sufficiently large RHN yield and branching ratio. The threshold (δ) regime provides an efficient and technically natural way to achieve this separation of scales, and it sets the stage for a successful realization of baryogenesis in the inverse-seesaw $B-L$ model.

While the parametric relation in Eq. (27) highlights the fundamental dilution-washout tension, it relies on standard thermal efficiency approximations. In a genuinely non-thermal setup, the heavy-neutrino abundance is not thermal, and the washout depends on the detailed time-evolution of the distributions. To establish the viability of this mechanism quantitatively, we must move beyond these estimates and treat the system dynamically.

IV. DYNAMICAL EVOLUTION AND BOLTZMANN EQUATIONS

In our previous parametric estimates, the survival of the baryon asymmetry relied on the assumption of an effective efficiency factor to model the washout

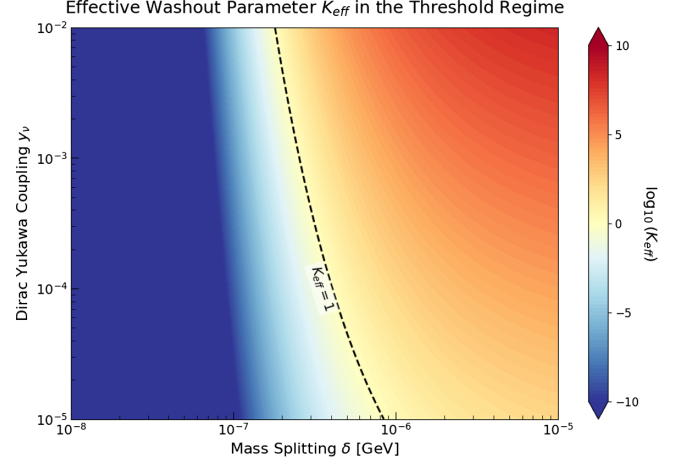


FIG. 2. Effective washout parameter K_{eff} in the threshold regime. Contour plot illustrating the logarithmic magnitude of the effective washout parameter, $\log_{10}(K_{\text{eff}})$, as a function of the scalar mass splitting δ and the Dirac Yukawa coupling y_ν . The parameter space is scanned over $\delta \in [10^{-8}, 10^{-5}]$ GeV and $y_\nu \in [10^{-5}, 10^{-2}]$. The remaining parameters are fixed to $M_N = 5$ TeV and the scalar-neutrino coupling $Y_N = 1.0$. To generate this, we utilize the analytical approximation $K_{\text{eff}} \simeq K \exp(-M_N/T_R)$, where K is the standard thermal washout factor evaluated at $T = M_N$, and T_R is the reheating temperature analytically derived from the threshold-suppressed scalar decay width ($\Gamma_\chi \propto \beta^3 \simeq (\delta/M_N)^{3/2}$). The dashed black contour delineates the critical boundary $K_{\text{eff}} = 1$. The region to the left (dark blue, $K_{\text{eff}} \ll 1$) represents the safe parameter space where the near-threshold kinematics enforce $T_R \ll M_N$, resulting in a severe exponential suppression of the inverse-decay washout processes.

suppression. To rigorously establish the viability of this nonthermal leptogenesis scenario and accurately track the energy transfer during the scalar decay, we now perform a full numerical integration of the coupled Boltzmann equations.

We track the proper time evolution of the comoving energy density of the decaying scalar field ρ_χ , the radiation bath ρ_R , the number density of the heavy pseudo-Dirac neutrinos n_N , and the $B-L$ asymmetry n_{B-L} . The system is governed by

$$\frac{d\rho_\chi}{dt} + 3H\rho_\chi = -\Gamma_\chi \rho_\chi \quad (30)$$

$$\frac{d\rho_R}{dt} + 4H\rho_R = (1 - \text{Br}_N)\Gamma_\chi \rho_\chi + M_N \Gamma_N (n_N - n_N^{\text{eq}}) \quad (31)$$

$$\frac{dn_N}{dt} + 3Hn_N = 2\text{Br}_N \frac{\Gamma_\chi \rho_\chi}{M_\chi} - \Gamma_N (n_N - n_N^{\text{eq}}) \quad (32)$$

$$\frac{dn_{B-L}}{dt} + 3Hn_{B-L} = \epsilon_{CP} \Gamma_N (n_N - n_N^{\text{eq}}) - \Gamma_W n_{B-L} \quad (33)$$

Γ_χ is the total decay width of the scalar field χ , Br_N is the branching ratio for the scalar decaying into a pair of heavy

neutrinos ($\chi \rightarrow NN$). Consequently, $(1 - \text{Br}_N)$ is the fraction of the scalar's energy that decays directly into the standard radiation bath. Γ_N is the total decay width of the heavy neutrinos. n_N^{eq} is the equilibrium number density of the heavy neutrinos at a given temperature T . For non-relativistic neutrinos ($M_N \gg T$), this follows a Maxwell-Boltzmann distribution proportional to $T^{3/2} e^{-M_N/T}$. ϵ_{CP} is the resonant CP asymmetry (see the Appendix for the explicit Yukawa texture realization).

To close the system, the Hubble expansion rate H is determined by the first Friedmann equation. Since the heavy neutrinos are strictly nonrelativistic in our regime ($T \ll M_N$), their energy density can be approximated by their rest mass times their number density ($\rho_N \simeq M_N n_N$), leading to

$$H = \sqrt{\frac{8\pi}{3M_{\text{Pl}}^2}(\rho_\chi + \rho_R + M_N n_N)} \quad (34)$$

The temperature of the thermal bath, $T(t)$, is dynamically extracted from the radiation energy density at each integration step, $\rho_R(t) = (\pi^2/30)g_* T(t)^4$.

The most critical aspect of the nonthermal evolution is the dynamical washout rate, Γ_W . Driven primarily by inverse decays ($LH \rightarrow N$), detailed balance in the thermal bath dictates that the macroscopic washout rate is proportional to the heavy neutrino decay width Γ_N , weighted by the thermodynamic penalty of extracting a heavy state from the bath [31]

$$\Gamma_W(t) = \frac{1}{2} \Gamma_N \frac{n_N^{\text{eq}}(T)}{n_\ell^{\text{eq}}(T)} \quad (35)$$

where $n_\ell^{\text{eq}}(T) = (3/4)(\zeta(3)/\pi^2)g_\ell T^3$ is the equilibrium density of the massless Standard Model leptons, and $n_N^{\text{eq}}(T)$ is the equilibrium density of the heavy neutrinos, which follows a Maxwell-Boltzmann distribution $n_N^{\text{eq}} \propto T^{3/2} e^{-M_N/T}$.

By continuously tracking $T(t)$ through ρ_R , our numerical integration captures the dynamical evolution of the washout rate throughout all phases of the evolution. The washout term Γ_W is evaluated at the instantaneous temperature $T(t)$ at every time step, with no approximation made at any stage. The detailed behavior of Γ_W/H across these phases is discussed in the following subsection.

A. Pre-reheating washout and the role of T_{max}

A subtle but important point concerns the washout during the pre-reheating phase, before radiation domination is established. Although the reheating temperature satisfies $T_R \ll M_N$, the bath temperature does not simply start at T_R . During the early matter-dominated phase, the radiation bath is built up gradually by the decay of χ , and the temperature rises from its initial value to a maximum T_{max} before falling back toward T_R as the scalar energy is depleted. For our

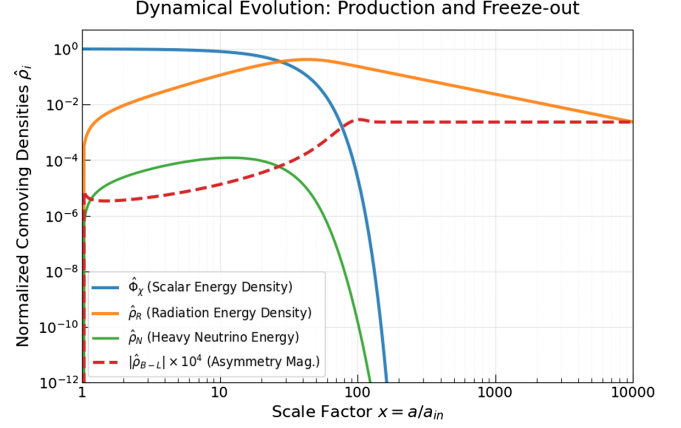


FIG. 3. Dynamical evolution of comoving densities normalized by the initial comoving scalar energy density, $\Phi_{\text{in}} = \rho_\chi(x=1)$, tracking energy transfer as a function of the dimensionless scale factor $x = a/a_{\text{in}}$. The solid blue line shows the normalized comoving energy density of the decaying scalar field, $\hat{\Phi}_\chi(x) = (\rho_\chi x^3)/\Phi_{\text{in}}$. The solid orange line indicates the normalized comoving radiation energy density, $\hat{\rho}_R(x) = (\rho_R x^3)/\Phi_{\text{in}}$, illustrating the reheating process. The solid green line represents the normalized energy density stored in the nonrelativistic heavy neutrinos, $\hat{\rho}_N(x) = (n_N \times M_N \times x^3)/\Phi_{\text{in}}$. Finally, the dashed red line depicts the absolute magnitude of the generated comoving $B - L$ energy density, $|\hat{\rho}_{B-L}(x)| = (|n_{B-L}| \times M_N \times x^3)/\Phi_{\text{in}}$, amplified by a factor of 10^4 for visibility on this logarithmic scale. a_{in} is the initial scale factor at the beginning of the decays. The parameters are fixed to the benchmark point: $M_N = 5$ TeV, $\delta = 10^{-6}$ GeV, and $y_\nu = 10^{-2}$. The $B - L$ asymmetry successfully freezes out and approaches a constant comoving value, confirming that the generated asymmetry is ultimately protected once $T \lesssim M_N$.

benchmark parameters ($M_N = 5$ TeV, $\delta = 10^{-6}$ GeV), we find numerically $T_{\text{max}} \approx 3.5 M_N$, reached very early at $x \approx 1.5$. This is a well-known feature of nonthermal reheating scenarios [41]: T_{max} can significantly exceed T_R even when $T_R \ll M_N$.

This raises the question of whether washout processes are active during the pre-reheating phase $T_{\text{max}} > T > T_R$. Our numerical integration already addresses this question rigorously, since the washout rate $\Gamma_W(t)$ is evaluated at the instantaneous dynamical temperature $T(t)$, extracted from $\rho_R(t)$ at every time step, with no approximation made regarding this phase.

Two features are worth noting. First, the bath temperature peaks early at $T_{\text{max}} \approx 3.5 M_N$ ($x \approx 1.5$), well above M_N , so inverse decays are kinematically accessible during the initial heating phase. Second, and more importantly, Γ_W/H itself does not peak at T_{max} but rather much later, at $x \approx 4 \times 10^3$ when $T \approx 0.2 M_N$, reaching $\Gamma_W/H \sim \mathcal{O}(10^7)$. This is because $\Gamma_W/H \propto (M_N^2/T^2) e^{-M_N/T}/H$ grows as H falls during matter domination, even while T is already decreasing, so the peak washout rate occurs deep in the cooling phase rather than at T_{max} .

Nevertheless, the generated $B - L$ asymmetry is not erased throughout this entire period, for the following reason. The evolution of n_{B-L} is governed by the competition between two terms in Eq. (33): the CP -violating source $\epsilon_{CP}\Gamma_N(n_N - n_N^{\text{eq}})$ and the washout term $\Gamma_W n_{B-L}$. As long as χ is actively decaying and injecting heavy neutrinos into the bath, $n_N \gg n_N^{\text{eq}}$, so the source term dominates over the washout term.

Only after the scalar has largely decayed and n_N approaches n_N^{eq} does the source switch off. At this point T has fallen sufficiently below M_N that Γ_W/H drops to zero exponentially, and the final $B - L$ asymmetry freezes in. This behavior is confirmed in Fig. 3: the comoving $B - L$ density reaches its asymptotic value and remains constant once $T \ll M_N$, demonstrating that the late-time exponential suppression of Γ_W ultimately protects the generated asymmetry.

V. RESONANT LEPTOGENESIS IN THE INVERSE SEESAW

In the inverse seesaw model, the heavy neutrinos form quasidegenerate pseudo-Dirac pairs. The CP asymmetry generated from the mixing of two nearly degenerate heavy neutrinos N_1 and N_2 is given by

$$\epsilon_{CP} \simeq \frac{\text{Im}[(h_{12})^2]}{|h_{11}|^2 + |h_{22}|^2} \frac{(M_1^2 - M_2^2)M_1\Gamma_2}{(M_1^2 - M_2^2)^2 + M_1^2\Gamma_2^2}, \quad (36)$$

where $h = Y_\nu^\dagger Y_\nu$ and Γ_2 is the decay width of N_2 .

Assuming maximal CP -violating phases, this expression simplifies to

$$\epsilon_{CP}^{\text{res}} \simeq \frac{\Delta M \Gamma}{\Delta M^2 + \Gamma^2/4}, \quad (37)$$

where $\Delta M = M_2 - M_1$ and Γ denotes the common decay width. The CP asymmetry reaches its theoretical maximum,

$$\epsilon_{CP}^{\text{max}} = \frac{1}{2}, \quad (38)$$

when the resonance condition

$$\Delta M \simeq \frac{\Gamma}{2} \quad (39)$$

is satisfied. In the Appendix, we shown how it is possible to get a texture for the Yukawa couplings that allowed to reach this maximal CP violation and still compatible with all light neutrinos data.

In the inverse seesaw, the mass splitting is controlled by the lepton-number violating parameter μ ,

$$\Delta M \simeq \mu, \quad (40)$$

while the decay width of a heavy neutrino is

$$\Gamma \simeq \frac{(y_\nu^\dagger y_\nu)}{8\pi} M_N. \quad (41)$$

The resonance condition therefore implies

$$\mu \simeq \frac{y_\nu^2 M_N}{16\pi}. \quad (42)$$

Since the Dirac Yukawa coupling is constrained by the light neutrino mass,

$$m_\nu \simeq \frac{(y_\nu v)^2}{M_N^2} \mu, \quad (43)$$

one finds

$$y_\nu^2 \simeq \frac{m_\nu M_N^2}{v^2 \mu}. \quad (44)$$

Substituting into the resonance condition yields the optimal mass splitting

$$\mu_{\text{res}} \simeq \sqrt{\frac{m_\nu M_N^3}{16\pi v^2}}. \quad (45)$$

VI. NUMERICAL RESULTS

We now summarize the numerical results for the baryon asymmetry η_B obtained via nonthermal resonant leptogenesis in the $U(1)_{B-L}$ inverse seesaw model, emphasizing the role of the threshold mechanism in suppressing washout at the TeV scale.

- (i) The reheating temperature T_R is generated by the decay of the $B - L$ scalar χ into heavy neutrinos. To avoid thermal washout, we impose the nonthermal condition $T_R < M_N$.
- (ii) Resonant enhancement of the CP asymmetry is achieved through the small lepton-number violating parameter μ
 - (a) For $M_N = 5$ TeV, the resonance condition is satisfied at $\mu \simeq 2$ MeV.
 - (b) Under these conditions, the CP asymmetry reaches $\epsilon_{CP} \simeq 0.48$.

The numerical results can be found in Figs. 4–6.

A. Phenomenological constraints on the benchmark scenario

Here, we evaluate our successful leptogenesis benchmark point ($M_N = 5$ TeV, $y_\nu = 10^{-2}$) against the relevant experimental constraints:

Heavy-neutrino mixing: The active-heavy mixing parameter governs the nonunitarity of the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix and direct collider

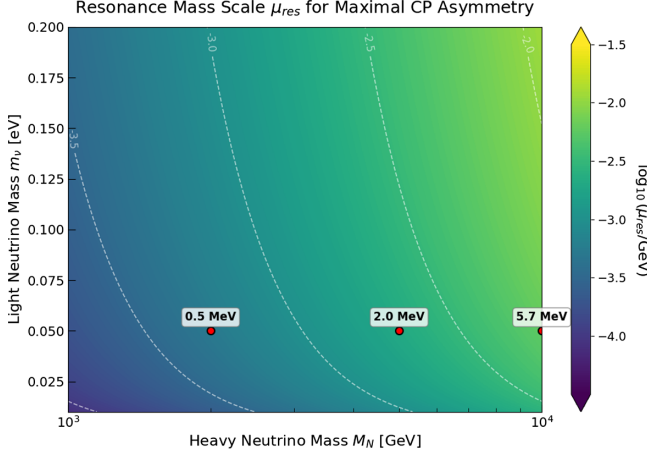


FIG. 4. Contour plot illustrating the lepton-number-violating mass scale, μ_{res} , required to satisfy the exact resonance condition ($\Delta M \simeq \Gamma_N/2$) for maximal CP asymmetry. The required scale is shown as a function of the heavy pseudo-Dirac neutrino mass M_N and the active light neutrino mass scale m_ν . The color map displays $\log_{10}(\mu_{\text{res}}/\text{GeV})$, computed using the analytical tree-level inverse seesaw matching condition $\mu_{\text{res}} = \sqrt{m_\nu M_N^3/(16\pi v^2)}$, where $v = 174 \text{ GeV}$ is the SM Higgs vacuum expectation value. The red markers indicate specific benchmark scenarios for an atmospheric neutrino mass scale of $m_\nu = 0.05 \text{ eV}$, demonstrating that for heavy neutrino masses in the multi-TeV range ($M_N = 2, 5, \text{ and } 10 \text{ TeV}$), the necessary lepton-number-violating scale falls in the $\mathcal{O}(\text{MeV})$ regime.

production. For our benchmark, this yields $\theta \simeq y_\nu v/(\sqrt{2}M_N) \simeq 1.2 \times 10^{-4}$. The corresponding nonunitarity measure $\theta^2 \sim 10^{-8}$ is several orders of magnitude

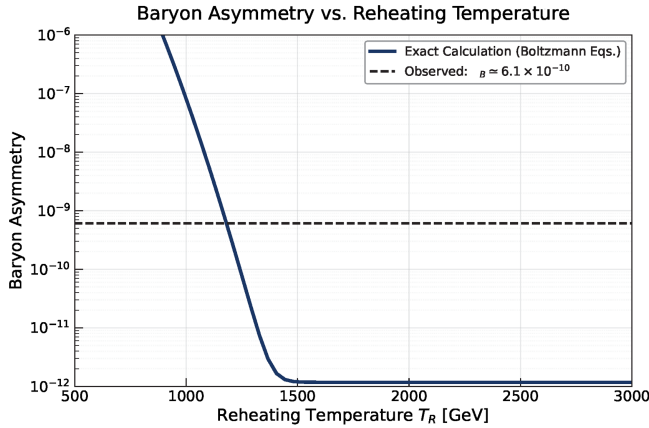


FIG. 5. Dependence of the baryon asymmetry on the reheating temperature. The final baryon-to-photon ratio η_B (solid blue curve) is shown as a function of the reheating temperature T_R . The heavy neutrino mass is fixed to $M_N = 5 \text{ TeV}$, with a Dirac Yukawa coupling of $y_\nu = 10^{-2}$ and a resonant CP asymmetry of $\epsilon_{CP} = 0.48$. The horizontal dashed line denotes the observed baryon asymmetry ($\eta_B \simeq 6.1 \times 10^{-10}$). As T_R approaches M_N , the exponential suppression of the equilibrium density $n_N^{\text{eq}}(T)$ weakens, activating the inverse-decay washout processes (Γ_W) and driving the asymmetry to zero.

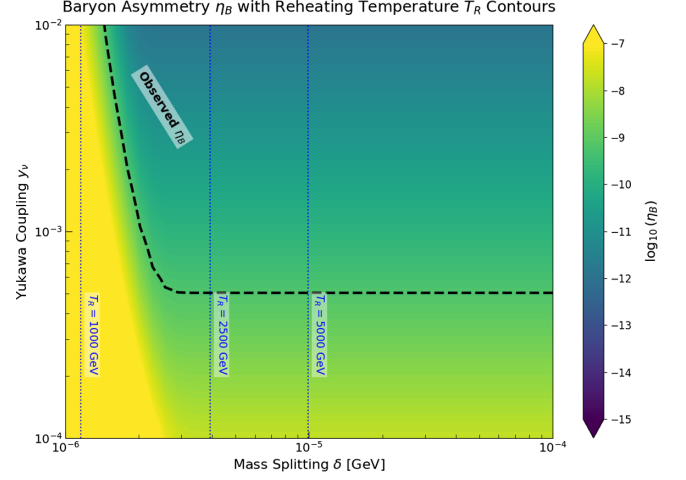


FIG. 6. Baryon asymmetry from full dynamical evolution. Contour plot displaying the logarithmic magnitude of the final baryon-to-photon ratio, $\log_{10}(\eta_B)$, as a function of the scalar mass splitting δ and the Dirac Yukawa coupling y_ν . The parameter space is scanned over $\delta \in [10^{-6}, 10^{-4}] \text{ GeV}$ and $y_\nu \in [10^{-4}, 10^{-2}]$. The remaining model parameters are fixed to our standard benchmark values: $M_N = 5 \text{ TeV}$, scalar-neutrino coupling $Y_N = 1.0$, branching fraction $\text{Br}(\chi \rightarrow NN) = 0.5$, and the resonant CP asymmetry $\epsilon_{CP} = 0.48$. Every point is computed by extracting the asymptotic $B - L$ yield from the exact numerical integration of the fully coupled Boltzmann equations (Sec. IV). The thick black dashed contour marks the parameter space that successfully reproduces the observed baryon asymmetry of the Universe ($\eta_B \simeq 6.1 \times 10^{-10}$). The vertical blue dotted lines indicate contours of constant reheating temperature ($T_R = 1000, 2500, \text{ and } 5000 \text{ GeV}$), visually confirming that the target asymmetry is achieved in the near-threshold regime ($\delta \sim \mathcal{O}(10^{-6}) \text{ GeV}$) where T_R is strictly bounded below the heavy neutrino mass scale, thereby dynamically evading thermal washout.

below current electroweak precision constraints, which require $\theta^2 \lesssim 10^{-3}$.

Lepton flavor violation (LFV): The most restrictive flavor constraint in the ISS is the radiative decay $\mu \rightarrow e\gamma$. The branching ratio depends on the used Casas-Ibarra matrix (as discussed in the Appendix). With an active-heavy mixing of $\theta \sim 10^{-4}$ and $M_N = 5 \text{ TeV}$, the induced LFV rates are below the current MEG limit of $\text{Br}(\mu \rightarrow e\gamma) < 4.2 \times 10^{-13}$, though they may fall within the reach of the future MEG II upgrade.

The Z' gauge boson: The $U(1)_{B-L}$ symmetry breaking generates a massive Z' boson. Current ATLAS and CMS dilepton resonance searches constrain $M_{Z'} \gtrsim 4.5 \text{ TeV}$ for a gauge coupling $g_{B-L} \sim 0.1$. The Z' mass is determined by the $B - L$ breaking scale v_{B-L} , which also sets the heavy neutrino mass. Choosing v_{B-L} such that $M_{Z'} \sim 6 \text{ TeV}$ evades LHC bounds without modifying the leptogenesis dynamics, as the Z' decoupled from the thermal bath earlier in the cosmic history.

Extended Higgs sector: The physical scalar χ mixes with the Standard Model Higgs boson h with a mixing angle α . To respect the LHC measurements of the 125 GeV Higgs signal strengths, $\sin \alpha \lesssim 0.1$ is needed. In our model, the near-threshold condition ($\delta \ll M_N$) assumes a suppressed coupling between the scalar sector and the heavy neutrinos, making a suppressed scalar mixing angle consistent and natural.

In conclusion, in our nonthermal leptogenesis scenario, the parameter space is among the allowed parameter spaces regions by current theoretical and experimental requirements.

VII. CONCLUSIONS

We have shown that the gauged $U(1)_{B-L}$ inverse seesaw model provides a phenomenologically viable model for baryogenesis through nonthermal resonant leptogenesis. Although the large Yukawa couplings inherent to the inverse seesaw constrain standard thermal leptogenesis due to strong washout processes, this difficulty is overcome by producing the heavy pseudo-Dirac neutrinos out of equilibrium via the decay of the $B - L$ Higgs scalar.

Two mechanisms are needed to get a successful baryogenesis in this scenario. First, setting the scalar mass close to the kinematic threshold, $M_\chi \simeq 2M_N + \delta$, lowers the maximum temperature of the radiation bath to $T_R \ll M_N$ (while $T_{\max}$ can transiently exceed M_N , as discussed in Sec. IVA). The CP -violating source dominates over washout while χ is actively decaying, and once the source switches off the bath temperature has fallen sufficiently below M_N that the washout rate becomes exponentially suppressed, protecting the frozen-in asymmetry. Second, the quasidegenerate structure of the heavy neutrinos enables resonant CP violation. By employing a modified Casas-Ibarra parametrization, we confirmed that the CP asymmetry can reach macroscopic values ($\epsilon_{CP} \sim 0.45$) without suffering from flavor cancellations when the resonance condition $\Delta M \simeq \Gamma_N/2$ is met.

Our exact numerical analysis demonstrates that for heavy neutrino masses $M_N \sim 5$ TeV, reheating temperatures of order $T_R \sim 1$ TeV, and a lepton-number-violating mass scale $\mu_S \sim \mathcal{O}(\text{MeV})$, the observed baryon asymmetry of the Universe is successfully reproduced. Maintaining the narrow near-threshold condition ($\delta \sim 10^{-6}$ GeV) requires a parameter tuning of roughly one part in 10^5 against one-loop radiative corrections. However, this degree of tuning is quite common for TeV-scale phenomenological models and is less severe than the electroweak hierarchy problem.

The scale of this scenario of nonthermal leptogenesis is the TeV scale. We have verified that the required parameter space evades current phenomenological constraints from lepton flavor violation (e.g., $\mu \rightarrow e\gamma$), LHC Z' resonance searches, and active-heavy neutrino mixing bounds. This model provides a consistent connection between neutrino mass generation and the cosmological origin of matter,

compatible with current data and open to future experimental scrutiny.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [1,2].

APPENDIX: CONCRETE YUKAWA TEXTURES AND RESONANT CP ASYMMETRY

In pseudo-Dirac systems characteristic of the Inverse Seesaw (ISS) mechanism, the resonant CP asymmetry can suffer from severe cancellations. In this appendix, we explicitly demonstrate that our phenomenological benchmark of an unsuppressed resonant CP asymmetry [$\epsilon_{CP} \sim \mathcal{O}(0.1-0.5)$] is possible without fine-tuning the flavor structure, provided that the heavy neutrino mixing matrix contains $\mathcal{O}(1)$ phases.

The active neutrino mass matrix in the ISS model is given by the exact tree-level formula

$$m_\nu \simeq \frac{v^2}{2} Y_\nu M_N^{-1} \mu_S (M_N^T)^{-1} Y_\nu^T$$

where $v = 174$ GeV is the Standard Model Higgs vacuum expectation value, M_N is the heavy Dirac mass matrix, and μ_S is the lepton-number-violating mass matrix. We adopt a modified Casas-Ibarra parametrization suitable for the ISS

$$Y_\nu = \frac{\sqrt{2}}{v} U_{\text{PMNS}} \sqrt{m_\nu^{\text{diag}}} R \sqrt{\mu_S^{-1}} M_N$$

Here, U_{PMNS} is the unitary neutrino mixing matrix, m_ν^{diag} contains the light active neutrino masses, and R is an arbitrary complex orthogonal matrix ($RR^T = \mathbb{I}$) that encodes the high-scale CP -violating phases.

The CP asymmetry for the decay of the heavy states ν_{H_i} is due to the interference of tree-level and one-loop self-energy diagrams. In the resonant regime ($\Delta M \simeq \Gamma_N/2$), the asymmetry is proportional to the interference term

$$\epsilon_{CP} \propto \frac{\text{Im}[(Y_\nu Y_\nu^\dagger)_{12}^2]}{(Y_\nu Y_\nu^\dagger)_{11}(Y_\nu Y_\nu^\dagger)_{22}}$$

To avoid the catastrophic pseudo-Dirac cancellations, the matrix R must be complex. We consider a minimal scenario with two heavy pseudo-Dirac pairs, defining R with a complex angle $\theta = a + ib$

$$R = \begin{pmatrix} \cos(a + ib) & \sin(a + ib) \\ -\sin(a + ib) & \cos(a + ib) \end{pmatrix}$$

For typical $\mathcal{O}(1)$ values of the parameters (e.g., $a \simeq \pi/4$ and $b \simeq 1.2$), the imaginary components of R prevent the elements of $Y_\nu Y_\nu^\dagger$ from becoming strictly real or aligned.

Inserting our benchmark values, $M_N = 5$ TeV and $\mu_S = 2$ MeV into the Casas-Ibarra parametrization yields elements of the Dirac Yukawa matrix of order $|(Y_\nu)_{ij}| \sim 10^{-2}$.

Evaluating the exact CP asymmetry with this texture produces a nonvanishing interference term $\text{Im}[(Y_\nu Y_\nu^\dagger)_{12}^2] \neq 0$, yielding a macroscopic resonant asymmetry of

$$\epsilon_{CP} \simeq 0.45$$

This explicit flavor realization confirms that the unsuppressed CP asymmetry utilized in our Boltzmann integration is consistent with the stringent low-energy neutrino oscillation physics.

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