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# Affleck-Dine Leptogenesis from Higgs Inflation

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We find that the triplet Higgs of the type-II seesaw mechanism can simultaneously generate the neutrino masses and observed baryon asymmetry while playing a role in inflation. We survey the allowed parameter space and determine that this is possible for triplet masses as low as a TeV, with a preference for a small vacuum expectation value for the triplet $v_\Delta < 10$ keV. This requires that the triplet Higgs must decay dominantly into the leptonic channel. Additionally, this model will be probed at the future 100 TeV collider, upcoming lepton flavor violation experiments such as Mu3e, and neutrinoless double beta decay experiments. Thus, this simple framework provides a unified solution to the three major unknowns of modern physics—inflation, the neutrino masses, and the observed baryon asymmetry—while simultaneously providing unique phenomenological predictions that will be probed terrestrially at upcoming experiments.

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*Introduction.*—Despite the great successes of the standard model (SM) at describing low-energy scales, there remain many open problems that demand the existence of new physics. These issues include the mechanism for the epoch of rapid expansion in the early Universe (inflation [1–5]), the origin of neutrino masses, and the source of the observed baryon asymmetry. Each of these mysteries is tied to early Universe physics, with any associated discoveries having significant implications for both particle physics and cosmology.

An exciting possibility to explore is whether each of these unknowns could be explained within a simple unified framework. There have been multiple attempts to do so in the past, but it is difficult to provide solutions to all three problems simultaneously with a single addition to the SM. For example, the SM plus three right-handed neutrinos can explain the neutrino masses via the seesaw mechanism [6–9] and generate the baryon asymmetry through leptogenesis [10] but not the inflationary sector. Inflationary baryogenesis has been widely investigated in the literature [11–27] but with few cases able to simultaneously explain each of the issues named above. Inflation with leptogenesis by the right-handed sneutrino has been considered in Refs. [28,29]. We present a model herein that represents a simple and well-motivated realization of this idea.

In this Letter, we study the possibility that these three problems can be solved through the simple extension of the SM by the triplet Higgs of the type-II seesaw mechanism. It has been known for a long time that with the addition of one triplet Higgs the baryon asymmetry cannot be generated through thermal leptogenesis but rather requires the introduction of a second triplet Higgs [30–34] or a right-handed neutrino [35–39]. In Ref. [40], a mechanism for nonthermal leptogenesis was proposed involving the Affleck-Dine mechanism but involved the addition of two triplet Higgses to the framework of supersymmetry. However, to explain the existence of neutrino masses, only one triplet Higgs is required. We propose a mechanism by which successful leptogenesis and neutrino mass generation can occur, with the addition of only a single triple Higgs. The triplet Higgs, in combination with the SM Higgs, will simultaneously give rise to a Starobinsky-like inflationary epoch [41–47]. This provides a unique connection between the high-energy dynamics of the early Universe and those at terrestrial colliders, which give novel phenomenological predictions that will be probed at future experiments.

*Baryogenesis from a complex inflaton.*—The fundamental feature of the Affleck-Dine mechanism is the generation of angular motion in the phase of a complex scalar field $\phi$ that is charged under a global U(1) symmetry [11]. Assuming it acquires a large initial field value in the early Universe, $\phi$ will begin to oscillate once the Hubble parameter becomes smaller than its mass $m$. If the scalar potential $V$ contains an explicit U(1) breaking term, a net U(1) charge asymmetry will be generated by this motion. The asymmetry number density associated with the U(1) charge is then given by

$$n_Q = 2Q\mathrm{Im}[\phi^\dagger \dot{\phi}] = Q\chi^2\dot{\theta}, \tag{1}$$

where $\phi = (1/\sqrt{2})\chi e^{i\theta}$. Therefore, in order to obtain a nonzero $n_Q$, we require nonzero vacuum value for $\chi$ and the motion of the complex phase $\theta$. This is easily realized if the $\phi$ field also plays the role of the inflaton, with an initial nonvanishing $\chi$ and $\theta$. For a general potential for the $\phi$ field, we can separate the U(1) conserving and nonconserving components:

$$U(\phi) \equiv U(\chi,\theta) = U_c(\chi) + U_b(\chi,\theta), \tag{2}$$

where $U_c(\chi)$ contains the U(1) conserving terms, which we assume dominate the potential during inflation, and $U_b(\chi,\theta)$ represents the U(1) breaking terms. If the kinetic term of $\phi$ is canonically normalized, then the Lagrangian can be written as

$$\mathcal{L} = -\frac{1}{2}g^{\mu\nu}\partial_\mu\chi\partial_\nu\chi - \frac{1}{2}f(\chi)g^{\mu\nu}\partial_\mu\theta\partial_\nu\theta - U(\chi,\theta), \tag{3}$$

where $f(\chi) = \chi^2$. Then the equations of motion for $\chi$ and $\theta$ are as follows:

$$\ddot{\chi} - \frac{1}{2}f'(\chi)\dot{\theta}^2 + 3H\dot{\chi} + U_{,\chi} = 0,$$
$$\ddot{\theta} + \frac{f'(\chi)}{f(\chi)}\dot{\theta}\dot{\chi} + 3H\dot{\theta} + \frac{1}{f(\chi)}U_{,\theta} = 0. \tag{4}$$

Assuming during inflation both $\chi$ and $\theta$ are slow rolling,

$$\dot{\chi} \simeq -\frac{M_p U_{,\chi}}{\sqrt{3U}} \quad \text{and} \quad \dot{\theta} \simeq -\frac{M_p U_{,\theta}}{f(\chi)\sqrt{3U}}. \tag{5}$$

From this, we may estimate the $Q$-number density at the end of inflation:

$$n_Q \approx -Q\chi_{\mathrm{end}}^2 \frac{M_p U_{,\theta}}{f(\chi_{\mathrm{end}})\sqrt{3U_{\mathrm{end}}}}. \tag{6}$$

Consequently, if the U(1) symmetry is composed of the global $\mathrm{U}(1)_B$ or $\mathrm{U}(1)_L$ symmetries, a baryon asymmetry can be generated prior to the electroweak phase transition. In the following, we will show that the $\phi$ field can be a mixed state of the SM and triplet Higgs, with a complex phase associated with the $\mathrm{U}(1)_L$ symmetry.

*Model framework.*—We now introduce the Lagrangian describing the SM Higgs doublet $H$ and the triplet Higgs $\Delta$. The scalars are parametrized by

$$H = \begin{pmatrix} h^+ \\ h \end{pmatrix}, \qquad \Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}, \tag{7}$$

where $h$ and $\Delta^0$ are the neutral components of $H$ and $\Delta$, respectively. The $\mathcal{L}_{\mathrm{Yukawa}}$ term contains not only the Yukawa interactions of the SM fermions, but also a new interaction between the left-handed leptons and the triplet Higgs $\Delta$:

$$\mathcal{L}_{\mathrm{Yukawa}} = \mathcal{L}_{\mathrm{Yukawa}}^{\mathrm{SM}} - \frac{1}{2}y_{ij}\bar{L}_i^c\Delta L_j + \mathrm{H.c.} \tag{8}$$

This interaction term will generate the neutrino mass matrix, once $\Delta^0$ obtains a nonzero vacuum expectation value (VEV). Through this interaction, we assign a lepton charge of $Q_L = -2$ to the triplet Higgs, thus opening the possibility for it to play a role in the origin of the baryon asymmetry.

The potential for the neutral Higgs' components is

$$\begin{aligned} V(h,\Delta^0) = {} & -m_H^2|h|^2 + m_\Delta^2|\Delta^0|^2 + \lambda_H|h|^4 + \lambda_\Delta|\Delta^0|^4 \\ & + \lambda_{H\Delta}|h|^2|\Delta^0|^2 + \bigg(\mu h^2\Delta^{0*} + \frac{\lambda_5}{M_p}|h|^2h^2\Delta^{0*} \\ & + \frac{\lambda_5'}{M_p}|\Delta^0|^2h^2\Delta^{0*} + \mathrm{H.c.}\bigg) + \cdots. \end{aligned} \tag{9}$$

Importantly for our model, the necessary $\mu$ coupling between the SM and triplet Higgs inherently violates the lepton number as defined through the triplet Higgs Yukawa interaction. Additionally, we have included dimension five lepton violating operators that are suppressed by $M_p$, since during inflation the field value is close to the Planck scale and, as such, they can dominate over the $\mu$ term. However, the higher-dimensional terms will play no role in low-energy physics.

The potential couplings are constrained by requiring the stability condition, and the nonvanishing $\Delta^0$ VEV can be approximated in the limit $m_\Delta \gg v_{\mathrm{EW}}$, $v_\Delta \equiv \langle\Delta^0\rangle \simeq -(\mu v_{\mathrm{EW}}^2/2m_\Delta^2)$, where the SM Higgs VEV is $v_{\mathrm{EW}} = 246$ GeV. The $\Delta^0$ VEV is bounded by $\mathcal{O}(1)$ GeV $> |\langle\Delta^0\rangle| \gtrsim 0.05$ eV, in order to generate the observed neutrino masses, while ensuring $y_\nu$ is perturbative up to $M_p$. The upper bound on the $\Delta^0$ VEV is derived from $T$-parameter constraints determined by precision measurements [48].

Although the above doublet-triplet Higgs model includes all the ingredients for generating the baryon asymmetry during inflation, the current data from cosmic microwave background (CMB) observations excludes their simple polynomial potential as the source of inflation [49]. One resolution to this problem is the addition of nonminimal couplings between the Higgs and the Ricci scalar. Then the full Lagrangian is

$$\begin{aligned} \frac{\mathcal{L}}{\sqrt{-g}} = {} & -\frac{1}{2}M_p^2R - F(H,\Delta)R - g^{\mu\nu}(D_\mu H)^\dagger(D_\nu H) \\ & - g^{\mu\nu}(D_\mu\Delta)^\dagger(D_\nu\Delta) - V(H,\Delta) + \mathcal{L}_{\mathrm{Yukawa}}. \end{aligned} \tag{10}$$

*Trajectory of inflation.*—The inflationary setting will be induced by both Higgs through their nonminimal couplings to gravity. These couplings act to flatten the scalar potential at large field values. This form of inflationary mechanism has been utilized in standard Higgs inflation and results in a Starobinsky-like inflationary epoch [41,50–61]. We consider the following nonminimal coupling:

$$F(H,\Delta) = \xi_H |h|^2 + \xi_\Delta |\Delta^0|^2 = \frac{1}{2}\xi_H \rho_H^2 + \frac{1}{2}\xi_\Delta \rho_\Delta^2, \tag{11}$$

where we have utilized the polar coordinate parametrization $h \equiv (1/\sqrt{2})\rho_H e^{i\eta}$, $\Delta^0 \equiv (1/\sqrt{2})\rho_\Delta e^{i\theta}$. An inflationary framework consisting of two nonminimally coupled scalars has been found to exhibit a unique inflationary trajectory [59]. In the large field limit, the ratio of the two scalars is fixed:

$$\frac{\rho_H}{\rho_\Delta} \equiv \tan\alpha = \sqrt{\frac{2\lambda_\Delta \xi_H - \lambda_{H\Delta}\xi_\Delta}{2\lambda_H \xi_\Delta - \lambda_{H\Delta}\xi_H}}. \tag{12}$$

To ensure the evolution of this trajectory, we require $2\lambda_\Delta \xi_H - \lambda_{H\Delta}\xi_\Delta > 0$ and $2\lambda_H \xi_\Delta - \lambda_{H\Delta}\xi_H > 0$. The derivation of this trajectory is given in Supplemental Material, Sec. I E [62]. The inflaton can then be defined as $\varphi$, through the relations

$$\rho_H = \varphi \sin\alpha, \qquad \rho_\Delta = \varphi\cos\alpha, \qquad \xi \equiv \xi_H \sin^2\alpha + \xi_\Delta \cos^2\alpha. \tag{13}$$

The Lagrangian becomes

$$\frac{\mathcal{L}}{\sqrt{-g}} = -\frac{1}{2}M_p^2 R - \frac{1}{2}\xi\varphi^2 R - \frac{1}{2}g^{\mu\nu}\partial_\mu\varphi\partial_\nu\varphi - \frac{1}{2}\varphi^2\cos^2\alpha g^{\mu\nu}\partial_\mu\theta\partial_\nu\theta - V(\varphi,\theta), \tag{14}$$

where

$$V(\varphi,\theta) = \frac{1}{2}m^2\varphi^2 + \frac{\lambda}{4}\varphi^4 + 2\varphi^3\left(\tilde{\mu} + \frac{\tilde{\lambda}_5}{M_p}\varphi^2\right)\cos\theta \tag{15}$$

and

$$\begin{aligned} m^2 &= m_\Delta^2 \cos^2\alpha - m_H^2 \sin^2\alpha, \\ \lambda &= \lambda_H \sin^4\alpha + \lambda_{H\Delta}\sin^2\alpha\cos^2\alpha + \lambda_\Delta \cos^4\alpha, \\ \tilde{\mu} &= -\frac{1}{2\sqrt{2}}\mu \sin^2\alpha\cos\alpha, \\ \tilde{\lambda}_5 &= -\frac{1}{4\sqrt{2}}(\lambda_5 \sin^4\alpha\cos\alpha + \lambda_5' \sin^2\alpha\cos^3\alpha). \end{aligned} \tag{16}$$

Since the generated lepton asymmetry is dependent upon the motion of $\theta$, we consider it to be a dynamical field. During inflation, $m \ll \varphi$, meaning that the quartic potential term dominates during the inflationary epoch.

We translate the Lagrangian in Eq. (14) from the Jordan frame to the Einstein frame, utilizing the transformations [63,64] $\tilde{g}_{\mu\nu} = \Omega^2 g_{\mu\nu}$ and $\Omega^2 = 1 + \xi\varphi^2/M_p^2$ and reparametrizing $\varphi$ in terms of the canonically normalized scalar $\chi$. We obtain the final Einstein frame Lagrangian

$$\frac{\mathcal{L}}{\sqrt{-g}} = -\frac{M_p^2}{2}R - \frac{1}{2}g^{\mu\nu}\partial_\mu\chi\partial_\nu\chi - \frac{1}{2}f(\chi)g^{\mu\nu}\partial_\mu\theta\partial_\nu\theta - U(\chi,\theta), \tag{17}$$

where

$$f(\chi) \equiv \frac{\varphi(\chi)^2\cos^2\alpha}{\Omega^2(\chi)} \quad \text{and} \quad U(\chi,\theta) \equiv \frac{V(\varphi(\chi),\theta)}{\Omega^4(\chi)}, \tag{18}$$

with the $\chi$ potential replicating the Starobinsky form in the large field limit,

$$U_{\inf}(\chi) = \frac{3}{4}m_S^2 M_p^2 \left(1 - e^{-\sqrt{2/3}(\chi/Mp)}\right)^2, \tag{19}$$

where $m_S = \sqrt{(\lambda M_p^2/3\xi^2)} \simeq 3\times 10^{13}$ GeV [46,65]. We will consider model parameters that ensure the inflationary trajectory is negligibly affected by the dynamics of $\theta$. Under this assumption, the resultant inflationary observables are consistent with the Starobinsky model and are in excellent agreement with current observational constraints [65]. See Supplemental Material, Sec. I [62], for details of the inflationary epoch and observational predictions.

*Lepton number density from the triplet Higgs.*—During inflation, we identify the inflaton field $\phi$ as a mixed state of the doublet Higgs and triplet Higgs. The corresponding potential becomes

$$U(\chi,\theta) = \frac{m^2\varphi^2(\chi) + \lambda\varphi^4(\chi)}{\Omega^4(\chi)} + \frac{2\tilde{\mu}\varphi^3(\chi) + 2\frac{\tilde{\lambda}_5}{M_p}\varphi^5(\chi)}{\Omega^4(\chi)}\cos\theta, \tag{20}$$

and the lepton number density is modified as

$$n_L = Q_L \varphi^2(\chi)\dot{\theta}\cos^2\alpha, \tag{21}$$

where $\alpha$ is the mixing angle between the doublet and triplet Higgs during inflation. During inflation, the $\chi$ field approaches $M_p$, and so we can ignore the subdominant $m$ and $\tilde{\mu}$ terms that are $\ll M_p$. Therefore,

$$\begin{aligned} n_{L\text{end}} &= Q_L \varphi_{\text{end}}^2 \dot{\theta}_{\text{end}} \cos^2\alpha \\ &\simeq -\mathcal{O}(1) Q_L \varphi_{\text{end}}^2 \frac{M_p U_{,\theta}}{f(\chi_{\text{end}})\sqrt{3U_{\text{end}}}}\cos^2\alpha \\ &\simeq -\mathcal{O}(1) Q_L \tilde{\lambda}_5 \varphi_{\text{end}}^3 \sin\theta_{\text{end}}/\sqrt{3\lambda}, \end{aligned} \tag{22}$$

where the $\mathcal{O}(1)$ factor accounts for the approximation of the slow roll relation at the end of inflation, used in the second step; see Supplemental Material, Sec. I B [62]. In the last step, we assume the quartic term dominates the inflationary potential and the $\tilde{\lambda}_5$ coupling dominates the breaking terms.

Numerically, we find that the $\mathcal{O}(1)$ factor is $\sim 3$, mainly originating from extending the slow roll approximation to the end of inflation, which we have defined as when the slow roll parameter is $\epsilon = 1$. After inflation, the lepton number density is just redshifted by $a^3$ with the inclusion of another $\mathcal{O}(1)$ factor $\Omega$; see Supplemental Material, Sec. I C [62].

*Baryon asymmetry parameter.*—After reheating, the generated nonzero $n_L$ will be present in the form of neutrinos, which will be redistributed by equilibrium electroweak sphalerons, with the ratio $n_B \simeq -(28/79)n_L$ [66–69]. To calculate the baryon to entropy ratio, we need the reheating temperature. The reheating process of Higgs inflation was first analyzed in Refs. [51,70], and it has been found that the parametric resonance production of $W/Z$ plays an important role for the preheating process. It has since been determined that the preheating process is more violent than previously expected [71–74] with unitarity being violated for $\xi > 350$. However, for such large $\xi$, the model can be UV completed in Higgs-$R^2$ inflation [56,75–78] for which the preheating process must be recalculated [79–81]. In our case, we choose $\xi = 300$ and, thus, the unitary problem is absent. A recent analyzes of the preheating in Higgs inflation [82] shows that, when $\xi > 100$, the reheating happens at an $e$-folding number of $\sim 3$ after the end of inflation, so we adopt $\Delta N = N_{\rm reh} - N_{\rm end} = 3$ for simplicity. The details of reheating may be different for our case due to the doublet and triplet Higgs mixing, and a comprehensive analysis is left for future work. For the typical parameters we consider, we find that reheating occurs at $t_{\rm reh} = 223/H_0$ and the corresponding Hubble parameter $H_{\rm reh} = 0.0047 H_0$. From $H^2_{\rm reh} \simeq (\pi^2/90) g_* (T^4_{\rm reh}/M_p^2)$ we obtain the reheating temperature $T_{\rm reh} \approx 2.2 \times 10^{14}$ GeV. Considering the entropy density just after reheating $s = (2\pi^2/45) g_* T^3_{\rm reh}$ [83], the baryon asymmetry parameter is then

$$\eta = \frac{n_B}{s}\bigg|_{\rm reh} = \eta_B^{\rm obs}\left(\frac{|n_{L\,\rm end}|/M_p^3}{1.3\times 10^{-16}}\right)\left(\frac{g_*}{112.75}\right)^{-1/4}, \qquad (23)$$

where $\eta_B^{\rm obs} \simeq 8.5 \times 10^{-11}$ is the observed baryon asymmetry parameter [84]. Equation (23) shows that, at the end of inflation, a lepton asymmetry of $1.3 \times 10^{-16} M_p^3$ is necessary to generate the observed baryon asymmetry, which corresponds to the example parameter sets $\tilde{\lambda}_5 = 7 \times 10^{-15}$ for $\theta_0 = 0.1$ and $\tilde{\lambda}_5 = 10^{-10}$ for $\theta_0 = 6.5 \times 10^{-6}$ from numerical calculations. Note that, in both of these cases, the typical parameters escape the isocurvature limits [85–87] placed by CMB observations [49]; see Supplemental Material, Sec. I D [62].

Comparing the quartic and dimension-5 terms, the cubic term becomes more relevant as $\varphi$ decreases. If the $\tilde{\mu}$ coupling becomes too large, the lepton asymmetry starts to rapidly oscillate and predictability breaks down. On the other hand, a small $\tilde{\mu}$ term helps to avoid washout of the lepton asymmetry. From the analysis in Supplemental Material, Sec. II [62], we require $|\tilde{\mu}| \lesssim 10^{-18} M_p$ for the initial $\theta_0 = 0.1$ to accommodate the observed baryon asymmetry.

For the typical parameters in our model, we assume $\lambda_H = 0.1$, and $\xi_H = \xi_\Delta = 300$ based on the argument above. For the other parameters, we set $\lambda_\Delta = 4.5 \times 10^{-5}$ to accommodate the inflation data, while there exist two options for $\lambda_{H\Delta}$. In the case of $\lambda_{H\Delta} > 0$, we require $2\lambda_\Delta \xi_H - \lambda_{H\Delta}\xi_\Delta > 0$ and $2\lambda_H \xi_\Delta - \lambda_{H\Delta}\xi_H > 0$ to ensure the mixing of $h$ and $\Delta^0$ during inflation. The typical parameter value we can consider is $\lambda_{H\Delta} = 10^{-5}$, giving the mixing angle $\alpha \simeq 0.02$. In the case of $\lambda_{H\Delta} < 0$, we need $|\lambda_{H\Delta}| < 2\sqrt{\lambda_H \lambda_\Delta}$ to avoid the potential becoming unbounded from below. We can then choose $\lambda_\Delta = 4.5 \times 10^{-5}$, $\lambda_{H\Delta} = -0.001$, giving $\alpha \simeq 0.07$.

Numerically, the observed baryon asymmetry is obtained for $\lambda_5 – \lambda_5' = 8.8 \times 10^{-11}$ $(7.9 \times 10^{-12})$ with $\theta_0 \sim 0.1$ and the typical parameter choices given above, in the case of $\lambda_{H\Delta} > 0$ $(\lambda_{H\Delta} < 0)$. In addition, we obtain an upper limit $|\mu| \lesssim 15(1.4)$ TeV. Note that all these parameters are defined at the renormalization scale near $M_p$. Given that the $\mathrm{U}(1)_L$ breaking term couplings are small, it is natural to consider that they originate from a spurion field which carries $\mathrm{U}(1)_L$ charge $+2$. The $\mathrm{U}(1)_L$ breaking terms are then generated by requiring that the spurion field obtains a VEV of the order of $\mathcal{O}(10^4)$ TeV.

*Parameter constraints.*—Since we expect a large reheating temperature, the triplet will be thermalized at the end of reheating, and we must consider possible washout processes. First, we require that the processes $LL \leftrightarrow HH$ are not effective:

$$\Gamma = n\langle\sigma v\rangle \approx y^2\mu^2/m_\Delta < H|_{T=m_\Delta}, \qquad (24)$$

where $H = \sqrt{(\pi^2 g_*/90)}(T^2/M_p)$.

The triplet Higgs generates the neutrino masses $m_\nu \simeq y(|\mu| v^2/2m_\Delta^2)$, where $m_\nu$ should be at least the order of $\sim 0.05$ eV. Combining this with the above relation, we obtain $m_\Delta < 10^{12}$ GeV for $m_\nu = 0.05$ eV.

The other processes that are necessary to consider are $LL \leftrightarrow \Delta$ and $HH \leftrightarrow \Delta$. They must not coexist; otherwise, the lepton number will be rapidly washed out. However, to maintain the lepton asymmetry, the process $LL \leftrightarrow \Delta$ must be efficient while $HH \leftrightarrow \Delta$ is out of equilibrium. This leads to the following requirement:

$$\Gamma_{ID}(HH \leftrightarrow \Delta)|_{T=m_\Delta} < H|_{T=m_\Delta}. \qquad (25)$$

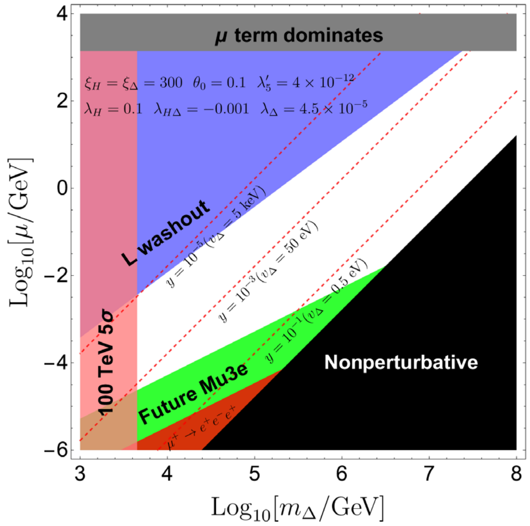


FIG. 1. The allowed region of parameter space is depicted (white), avoiding washout processes (blue), cubic term domination of $\mathrm{U}(1)_L$ breaking (gray), and nonperturbative neutrino Yukawa couplings (black). The red region denotes the current limits from lepton violating decays [88], with green indicating the future Mu3e experimental sensitivity [89]. The future 100 TeV collider constraints are depicted by the light red region [90].

Note that $\Gamma_{ID}(HH \leftrightarrow \Delta)|_{T=m_\Delta} \approx \Gamma_D(\Delta \to HH) \simeq (\mu^2/32\pi m_\Delta)$ and $v_\Delta \simeq -(\mu v_{\mathrm{EW}}^2/2m_\Delta^2)$. From Eq. (25), one can easily get

$$v_\Delta \lesssim 10^{-5}\ \mathrm{GeV}(m_\Delta/\mathrm{TeV})^{-1/2}; \tag{26}$$

hence, for $m_\Delta \gtrsim 1$ TeV, we generally require that $v_\Delta \lesssim 10$ keV to prevent the washout of the lepton asymmetry.

In Fig. 1, we depict the region of parameter space for which the generated lepton number density leads to successful baryogenesis. The black region is excluded by requiring perturbative neutrino Yukawa $y$ couplings up to the Planck scale $M_p$ ($y \lesssim 1$). However, a small Yukawa coupling $y < 0.1$ is preferred for the size of quartic coupling we consider, $\lambda_\Delta \simeq 4.5 \times 10^{-5}$, to avoid fine-tuning at the high-energy scale. The gray region is excluded by requiring that the $\mu$ term does not destroy the generated lepton asymmetry. The blue region describes the parameters that lead to washout of the lepton asymmetry. From Eq. (26), the blue region implies an upper limit on $v_\Delta$, namely, $v_\Delta \lesssim 10$ keV.

There is an additional limit from precision measurements on the vacuum expectation value $\langle \Delta^0 \rangle$; namely, it must be less than a few GeV. LHC searches apply lower bounds on the masses of the triplet Higgs components—for example, the current limit on the mass of the doubly charged Higgs is ~800 GeV [91]—and, thus, we depict only triplet masses $\geq 1$ TeV. This limit may be increased at the upgraded high-luminosity LHC or at future colliders [90]. Assuming the Yukawa couplings are of the same order of magnitude, the lepton flavor violating processes induced by the doubly charged Higgs already provide a limit on the parameter space [88,92]. The current limit will be improved by 2 orders of magnitude by the upcoming Mu3e experiment [89].

*Concluding remarks.*—We have shown that the introduction of the triplet Higgs of the type-II seesaw mechanism to the SM provides a simple framework in which inflation, neutrino masses, and the baryon asymmetry are all explained. We now summarize the unique combination of phenomenological predictions of this model. (1) Depending upon the vacuum value of the triplet and its Yukawa couplings, the triplet Higgs can decay mainly into gauge bosons or leptons. In our model, $v_\Delta$ can be accommodated only within the range 10 keV–eV, with the upper limit ensuring lepton asymmetry washout effects are negligible. Importantly, for this $v_\Delta$ range the triplet Higgs dominantly decays into leptons. If we observed the triplet Higgs in such a channel, it would provide a smoking gun for our model. (2) The associated doubly charged Higgs directly leads to lepton flavor violating processes such as $\mu \to e\gamma$ and $\mu \to eee$. The current experimental limits already provide constraints on the triplet Higgs properties (see Fig. 1), with future experiments such as Mu3e to improve upon the $\mu \to eee$ limits by 2 orders of magnitude. Thus, the allowed parameter space of our model will be tested in the near future. (3) In this model, the observed neutrino masses are of the Majorana type. This is in contrast to models which include right-handed neutrinos, where the observed neutrinos can have both Dirac- and Majorana-type mass terms. Thus, our model can be probed at near future neutrinoless double beta decay experiments. In addition, the baryon asymmetry generated in this model is independent of the leptonic $\mathcal{CP}$ phase, with $\mathcal{CP}$ spontaneously broken at early times of the Universe. There are currently conflicting measurements of the leptonic $\mathcal{CP}$ phase coming from the T2K and NOvA experiments, with T2K disfavoring a $\mathcal{CP}$ conserving angle [93], which is inconsistent with the NOvA result [94]. In this context, our model provides an interesting theoretical possibility for leptogenesis. (4) We assume that the inflationary period is induced by two scalar fields. Generally, such inflationary setups can generate nontrivial non-Gaussian features [87]. In addition, a sizable isocurvature signature could be produced if considerable washout effects are allowed at late times. These possibilities may be probed by future observations.

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