Baryon Number Violation via Majorana Neutrinos in the Early Universe, at the LHC, and Deep Underground

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We propose and investigate a novel, minimal, and experimentally testable framework for baryogenesis, dubbed \textit{dexiogenesis}, using baryon number violating effective interactions of right-handed Majorana neutrinos responsible for the seesaw mechanism. The distinct LHC signature of our framework is same-sign top quark final states, possibly originating from displaced vertices. The region of parameters relevant for LHC phenomenology can also yield concomitant signals in nucleon decay experiments. We provide a simple ultraviolet origin for our effective operators, by adding a color-triplet scalar, which could ultimately arise from a grand unified theory.

Gauge singlet right-handed neutrinos (RHNs) must be included in the Standard Model (SM) to allow an explanation of non-zero neutrino masses, as demanded by a large and well-established body of neutrino oscillation data. Without these states, the SM cannot be considered a renormalizable theory of Nature. If neutrinos are Dirac particles, the associated RHNs only couple to the SM via Yukawa interactions of negligible strength, \( y_N \lesssim 10^{-12} \), and are not expected to be directly detectable in the foreseeable future. On the other hand, if the observed neutrinos are Majorana states, they most naturally get their mass from a seesaw mechanism [1]. In that case, RHNs may have \( O(1) \) couplings to SM neutrinos, but they would then have to be exceedingly heavy, \( \gtrsim 10^{14} \text{ GeV} \), far beyond the reach of terrestrial experiments. However, it may very well be that RHNs are much lighter, and they can be directly probed in lepton number violating processes in collider experiments [2,3], and perhaps other searches, such as those for proton decay [6].

In this letter, we assume that the RHNs associated with the seesaw mechanism are near the weak scale \( \lesssim 1 \text{ TeV} \). Higher dimensional operators involving the RHNs are generically present. In particular, we will further assume that the RHNs have baryon number violating interactions, mediated by dimension 6 operators involving right-handed quarks and suppressed by a scale \( \sim 1 \sim 10 \text{ TeV} \). We will show that these assumptions allow for direct generation of a baryon number asymmetry through RHN decays in the early Universe, which we dub \textit{dexiogenesis} (dexios: Greek for the right hand). This is in contrast to canonical leptogenesis [7] where the lepton asymmetry needs to be further processed into baryon number through electroweak sphalerons [8]. Our direct baryogenesis mechanism is most constrained by nucleon decay bounds. However, we show that for viable parameters one could have distinct collider signatures. This scenario, with dim-6 operators, is an effective field theory and can be embedded in a simple renormalizable model, and possibly a grand unified theory (GUT), where additional signals are expected to arise at the Large Hadron Collider (LHC) and future colliders.

Our point of departure is the SM augmented by two Majorana RHNs \( N_a, a = 1, 2 \), the minimum required for a realistic seesaw mechanism based on current neutrino data. We add the following terms to the Lagrangian

\[ \mathcal{L}_N = M_a \bar{N}_a^c N_a + y_{N}^{ij} H \bar{N}_a L_i + \text{h.c.}, \]  

(1)

where \( M_a \) is the mass of \( N_a \), \( i = 1, 2, 3 \) enumerates SM generations, and \( y_{N}^{ij} \) is a Yukawa matrix; \( H \) and \( L_i \) are the Higgs and lepton doublets of the SM, respectively.

Light neutrino masses \( m_{\nu} \gtrsim 0.1 \text{ eV} \), implied by the oscillation data, can be generated from the renormalizable interactions in Eq. (1), via the seesaw mechanism: \( \langle m_{\nu} \rangle_{ij} \sim y_{N}^{ai} y_{N}^{aj} \langle H \rangle M_a \). Nonetheless, the SM, henceforth defined to include the Lagrangian in Eq. (1), is widely expected to be an effective theory that is further enriched with new interactions at higher scales. This expectation is strongly motivated by the need for a dark matter candidate and also a baryogenesis mechanism to generate the observed baryon asymmetry of the Universe.

Baryogenesis requires a source of baryon number violation [13]. Hence, it may be necessary to extend the SM by effective operators that violate baryon number [14]. Such operators are suppressed by scales associated with new physics, often considered to be very high, \( \gtrsim 10^{15} \text{ GeV} \), as implied by nucleon decay constraints. However, it is compelling to look for scenarios where new physics arises at lower scales. For one thing, it is reasonable to assume that the physics underlying the Higgs potential is not very far from the weak scale. Such physics would then have the added benefit of being potentially testable.

Motivated by the above considerations, we assume the following baryon-number violating operators involving the RHNs, in addition to those made up of only the SM fields,

\[ \mathcal{L}_{\text{BV}} = \frac{\lambda_{ij}^{kh} N_a u_i d_j d_k}{\Lambda^2} [N_a u_i d_j d_k]_R + \frac{\lambda_{N}^{dim}}{\Lambda^2} [N_a d_i]_R [Q_l Q_m]_L + \text{h.c.}, \]

(2)
where $i,j,k$ are family numbers of right-handed quark ($u,d$) mass eigenstates and $l,m$ enumerate left-handed quark ($Q$) generations. Here, $\lambda^{ijk}_a$ and $\kappa^{ilm}_a$ are generally complex constants determined by the ultraviolet (UV) theory. These operators could arise from grand unified theories, as shown in a concrete example in the end. They are the lowest dimensional operators that allow RHNs to couple to baryon number directly [15].

In order to have successful dexiogenesis, the coefficients of the relevant dim-6 operators cannot be too small. To avoid excessive low energy baryon number violation for $\Lambda \gtrsim 1$ TeV, those operators would mainly involve right-handed third generation quarks, which help avoid severe constraints from nucleon decay data, discussed in more detail below. To see this, note that quark mass diagonalization can induce operators that involve light quarks, in the presence of left-handed fields. For this reason, we will not further consider tree-level dim-6 operators $NdQQ$ in Eq. (2). These operators can still be generated from $Ntb b$ via radiative corrections, but would not lead to severe constraints.

In light of the above discussion, throughout this work, we will focus on operators involving right-handed third generation top and bottom quarks (originating from the first term in Eq. (2), with explicit spinor contractions)

$$\mathcal{L}_{\text{UV}}^{3\text{R}} = \lambda_{ab} \left[ N \bar{P} R b \right] \frac{[\hat{v} P R b]}{\Lambda^2},$$

where $P_R = (1 + \gamma_5)/2$ is the right-handed projector. The dominance of the third generation could be expected from a connection to UV flavor dynamics. Operators with other combinations of chiralities and flavors can in principle be present, but they must be more suppressed [19].

Remarkably, the addition of the operators in Eq. (3) provides all the necessary ingredients encoded in Sakharov’s conditions [13] for baryogenesis: (i) these interactions are manifestly baryon number violating, (ii) their complex coefficients provide a source of CP violation, and (iii) if the Universe has a low reheat temperature $T_{\text{RH}} \ll M_a$, then the $N_a$ will decay out of equilibrium. This mechanism, dexiogenesis, allows $T_{\text{RH}} \ll 100$ GeV, since the baryon asymmetry is directly generated and hence electroweak sphalerons do not need to be active. Let $N_1$ be the lighter of the two RHNs in our setup. Then, the interference of the tree and the 2-loop diagrams in Fig. 1 will lead to a baryon asymmetry

$$\varepsilon = \frac{\Gamma(N_1 \to t b b) - \Gamma(N_1 \to \bar{t} \bar{b} b)}{2\Gamma_{N_1}},$$

where the width of $N_1$ is given by

$$\Gamma_{N_1} = \frac{|\lambda|^2 M_1^5}{1024 \pi^3 a^4} F \left( \frac{m_2^2}{M_1^2} \right),$$

with $F(x) = 1 - 8x - 12x^2 \log x + 8x^3 - x^4$.

In the presence of the higher dimensional operator Eq. (3) with a TeV scale cutoff, $N_1$ decays induced by neutrino Yukawa interactions [Eq. (1)] are subdominant, for values of $M_1$ near the weak scale. Given a realistic seesaw mechanism for the SM active neutrino masses, in general we have $y^2 \lesssim 10^{-6}$ in the absence of fine tuning [20]. The induced $N_1 \to W l$ decay rate is then estimated to be $\Gamma_{N_1 \to W l} \lesssim 10^{-12}$ GeV $(y^2/10^{-6})^2 [M_1/(200 \text{ GeV})]$. We find that, for $M_1$ of a few hundred GeV and $\Lambda/\sqrt{\Lambda_1} \lesssim 25$ TeV, that rate is smaller than the baryonic decay rate.

The baryon asymmetry can be conveniently obtained using the unitarity cut method [16]

$$\varepsilon = \frac{\text{Im}(\lambda^2 \lambda_2^2)}{3072 \pi^3 |\lambda|^2} \left( M_2 \right)^4 \frac{M_1 M_2}{(M_2^2 - M_1^2)}.$$

The baryon number to entropy ratio $\eta \equiv n_B/s \sim 10^{-10}$ is roughly given by $\eta \sim \varepsilon/g^* \sqrt{\Lambda}$ where $g^*$ is the number of relativistic degrees of freedom at the end of baryogenesis. We then see that for $M_1 \sim M_2$ and $\lambda_3 \sim 1$, we need $M_1/\Lambda \gtrsim 0.1$. Hence, for $M_1 \lesssim 1$ TeV, relevant for collider phenomenology, the cutoff scale must be sufficiently low, $\Lambda \lesssim 10$ TeV. Let us then examine the experimental constraints on $\Lambda$.

It turns out that nucleon decay limits provide the most stringent lower bound on $\Lambda$ in the above model where RHNs violate both lepton and baryon numbers. While the operators in Eq. (3) do not contain light quarks, quantum loop corrections can induce nucleon decay via these baryon number violating interactions. Fig. 2 provides a sample two-loop diagram for proton decay, with a rate given by

$$\Gamma(p \to \pi^+ \nu) = \frac{(1 + g_A)^2 \alpha^2 m_p |\xi|^2}{32 \pi f_\pi^2} \left( \frac{\Lambda}{\Lambda_1} \right)^{2 \alpha} \left( \frac{\sqrt{\Lambda}}{\Lambda_1} \right)^{4 \alpha}.$$

FIG. 1: Tree and two-loop diagrams for dexiogenesis.

FIG. 2: One of the leading diagrams that yield proton decay.
where $g_A = 1.27$ is the nucleon axial charge, $f_π = 131$ MeV is pion decay constant, and lattice calculations [21] yield the form factor $α ≈ −0.01125$ GeV$^3$. The quantity

$$\xi ≈ \frac{λ_{qcd}G_F^2m_π^2V_{ud}^2V_{us}^*V_{ub}^*}{(16\pi^2)^2Λ^2} λ_αθ_a$$  (8)

is the Wilson coefficient from estimating the two-loop diagram in Fig. 2. The angle $θ_a$ is the mixing between $N_a$ and the SM active neutrinos. The hadronic mass scale $Λ_{qcd} ≈ 200$ MeV must be introduced under a symmetry argument. The operator we started with is $[1]^{-c}P_{Rb}[\bar{f}P_{Rb}]$, and after the $W$-loop drawing as in Fig. 2, the operator for proton decay turns out to be $[N^c_{id}P_{Rd}][\bar{u}P_{Ld}]$ (which is the radiatively generated $NdQQ$ operator mentioned earlier). The fact that one of the down quark is still right-handed implies an external (constituent) quark mass insertion $≈ Λ_{qcd}$.

The resulting proton decay life time is

$$τ(p → π^+ν) ≈ 2.5 \times 10^{32} \text{yr} \left(\frac{Λ/\sqrt{N_a}}{1.5 \text{ TeV}}\right)^4 \left(\frac{θ_a}{10^{-6}}\right)^{-2}.$$  (9)

The current experimental lower limit on the $p → π^+ν$ decay channel is $1.6 \times 10^{31} \text{ yr}$ [22]. Hence, the above lifetime [9] is not far from the current limit and, in the region of parameters considered in our work, can be within the reach of future nucleon decay experiments [23, 24].

Here, we also address a potential bound from requiring that the asymmetry in Eq. [9] is not washed out after baryogenesis. The low reheat temperature assumption mentioned before can ensure that processes mediated by the operators in Eq. [9], such as $bb → N\bar{t}$, are effectively turned off. However, loop processes similar to those depicted in Fig. 2 can lead to baryon number violation mediated by lighter states. Let us assume, for illustrative purposes, that the reheat temperature is $T_{RH} ≈ 1$ GeV, well above the temperature at the onset of Big Bang Nucleosynthesis. We then have to make sure that the analog of neutron-anti-neutron oscillation mediated by $(c\bar{s}s)/(\bar{c}s\bar{s})/Λ^5$ is not active. A straightforward comparison with Eq. [8] yields

$$Λ^5 ≈ \left(\frac{θ_a}{ξ} λ_{qcd} \frac{V_{td}^2}{m_s} \frac{V_{ub}^*}{V_{ts}^*V_{cb}^*}\right)^2 M_1,$$  (10)

where $m_s ≈ 100$ MeV is the strange quark mass. For $M_1 = 200$ GeV, we find $Λ' ≈ 2 \times 10^8$ GeV. The rate of the $(c\bar{s}s)/(\bar{c}s\bar{s})$ process is estimated to be $Γ_{ΔB=2} \sim T_{RH}^8/Λ^{10} ≈ 10^{-93}$ GeV, for $T_{RH} ≈ 1$ GeV, which is completely negligible compared to the Hubble rate at this temperature, $H ≈ T_{RH}^2/M_{\text{Planck}} ≈ 10^{-19}$ GeV. We also note that bounds from neutron-anti-neutron oscillation would not constrain our model, since that process involves up and down quarks, for which the corresponding suppression scale is larger than the above $Λ'$ scale, and much higher than the current limit [25].

An immediate consequence of Eq. [3] is the possible production of same-sign top quarks at the LHC and future hadron colliders, due to the Majorana nature of RHNs, as shown in the left panel of Fig. 3. In this process, the RHN $N_a$ and a top quark are first produced, and then $N_a$ decays into another top quark and two bottom quarks. Because it is a Majorana particle, an on-shell $N_a$ is equally likely to decay into $t\bar{b}b$ or $\bar{t}b\bar{b}$ final states. The violation of baryon number is manifested in terms of the violation of top quark number (by two units). The sign of the top quark can be inferred from its leptonic decays. For a RHN with a few hundred GeV mass and the effective cutoff scale $Λ/\sqrt{N_a}$ of a few TeV, we find that the cross section for this process can be as large as $≈ 0.3 \text{ fb}$ in the LHC Run-II at 13 TeV. The main background for this signal is from $t\bar{b}b\bar{b}$ final states with the lepton charge from a top quark decay misidentified, which is suppressed by the small misidentification rate [26]. In Table [1] we list the leading order cross sections of our signal for several sample mass values of RHNs. These points have not been excluded by the existing LHC data. For example, with $m_N = 200$ GeV and $Λ/\sqrt{N_a} = 1.5$ TeV, the cross section at 8 TeV is $0.07 \text{ fb}$, which implies only 1–2 events given the existing integrated luminosity $≈ 27 \text{ fb}^{-1}$, and it is further suppressed by the top quark leptonic branching ratios.

| \( \sqrt{s} = 13 \text{ TeV} \) | 0.34 fb | 0.16 fb | 8 \times 10^{-2} \text{ fb} | 5 \times 10^{-2} \text{ fb} |
|-----------------------------------|----------------|---------|----------------|----------------|
| \( m_N \)                        | 200 GeV        | 500 GeV | 800 GeV        | 1 TeV          |

TABLE I: Same-sign top quark production cross section, at the 13 TeV LHC, via a Majorana RHN and the contact operators in Eq. [3]. The cutoff scale is fixed to be $Λ/\sqrt{N_a} = 1.5$ TeV.

Following the same logic as introducing RHNs to make the SM renormalizable, we now discuss a UV completion that generates the effective operator Eq. [3]. Given a TeV scale cutoff, it is possible to directly probe the heavy particles in such a model in LHC Run-II and future hadron colliders. The model is an extension of the SM that contains a color-triplet scalar, $T$, with quantum numbers $(1, 1/3)$. The corresponding Lagrangian is

$$\mathcal{L}_{UV} = f_a T \bar{N}^c_{id}P_{Rd} + f' T^c \bar{P}_{Rb} + M_{T}^2 |T|^2.$$  (11)

In fact, this is the simplest model that yields the flavor and color structures of the effective operators in Eq. [3]. After integrating out the color-triplet scalar $T$, corresponding to a cutoff $Λ^2 ≡ f_a f'/M_T^2$. The TeV scale cutoff as discussed above can be naturally obtained for $M_T ≈ \text{ TeV}$ and $f_a f' ≈ \mathcal{O}(1)$. We note that in the above UV model, it is possible to have baryogenesis through the decays of the $T$ particle [11]. We will not further explore such a possibility in this work.
The introduction of the scalar $T$ could offer richer phenomenology at colliders. If light enough, $T, T^*$ can be pair produced at hadron colliders. Each triplet will first decay into $N + b$, which is then followed by subsequent decay $N \to tbb$ via a virtual $T$. The above chain of processes are represented by the diagram in the right panel of Fig. [3]. These together result in same-sign top quark final states with many $b$-jets. In Table II, we give the leading-order QCD cross section for the $T, T^*$ pair production at the $13$ and $100$ TeV proton-proton collider, calculated with MadGraph [29].

Moreover, an additional distinct signal could be displaced vertices from the decay of RHNs, if we take a somewhat larger cutoff scale $\Lambda/\sqrt{\Lambda_1}$, where $\Lambda_1$ is of the order of $10^{12}$ GeV and $\Lambda_2$ is the strong scale, $\Lambda_2 \sim 10^{16}$ GeV. The recent discovery of the Higgs boson with the SM Yukawa interactions provides a good example of this. If the Higgs boson is as heavy as $125$ GeV, one could introduce a new $\bar{Q}_3 T_D^\pm H^{\pm} \sim (12)$ term. Without $\bar{Q}_3 T_D^\pm H^{\pm}$, $\Lambda_2/\sqrt{\Lambda_1}$ is much too large. If $\Lambda_1 < M_H$, $\Lambda_2/\sqrt{\Lambda_1}$ is of the order of $10^{12}$ GeV and $\Lambda_2/\sqrt{\Lambda_1}$ is of the order of $10^{12}$ GeV. Moreover, the Higgs boson can decay into a virtual $T$ and a virtual $T^*$, and the above chain of processes are represented by the diagram in Fig. [3]. These together result in same-sign top quark final states with many $b$-jets. In Table II, we give the leading-order QCD cross section for the $T, T^*$ pair production at the $13$ and $100$ TeV proton-proton collider, calculated with MadGraph [29].

TABLE II: Pair production cross sections of $T, T^*$ via strong interaction at the 13 and 100 TeV proton-proton colliders.

| $\sigma(pp \to T T^*)$ | $m_T$ |
|------------------------|-------|
|                        | 1.5 TeV | 2 TeV | 5 TeV | 10 TeV |
| $\sqrt{s} = 13$ TeV    | 0.16 fb | 0.01 fb | -- | -- |
| $\sqrt{s} = 100$ TeV   | 384 fb | 92 fb | 0.54 fb | $4 \times 10^{-3}$ fb |

Acknowledgments. We would like to thank B. Dev, P. Meade, R. Mohapatra, and G. Senjanovic for discussions. We also thank B. Dev and R. Mohapatra for informing us of their forthcoming paper on related topics [33]. The work of H.D. is supported in part by the United States Department of Energy under Grant Contracts DE-SC0012704. The work of Y.Z. is supported by the Gordon and Betty Moore Foundation through Grant #776 to the Caltech Moore Center for Theoretical Cosmology and Physics, and by the DOE Grant DE-FG02-92ER40701, and also by a DOE Early Career Award under Grant No. DE-SC0010255. Y.Z. thanks the BNL theory group for hospitality at the final stage of this work.
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