B-L Violating Proton Decay Modes and New Baryogenesis Scenario in $SO(10)$

K. S. Babu[1,*] and R. N. Mohapatra[2,†]

[1]*Department of Physics, Oklahoma State University, Stillwater, Oklahoma 74078, USA*
[2]*Maryland Center for Fundamental Physics, Department of Physics, University of Maryland, College Park, Maryland 20742, USA*
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We show that grand unified theories based on $SO(10)$ generate quite naturally baryon number violating dimension seven operators that violate B-L, and lead to novel nucleon decay modes such as $n \to e^- K^+$, $e^-\pi^+$ and $p \to \nu\pi^+$. We find that in two-step breaking schemes of nonsupersymmetric $SO(10)$, the partial lifetimes for these modes can be within reach of experiments. The interactions responsible for these decay modes also provide a new way to understand the origin of matter in the Universe via the decays of grand unified theory (GUT) scale scalar bosons of $SO(10)$. Their (B-L)–violating nature guarantees that the GUT scale induced baryon asymmetry is not washed out by the electroweak sphaleron interactions. In minimal $SO(10)$ models this asymmetry is closely tied to the masses of quarks, leptons and the neutrinos.

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Baryon number violation is a very sensitive probe of physics beyond the standard model (SM). Interactions that violate the baryon number (B) are not present in the renormalizable part of the SM Lagrangian but can arise through effective higher dimensional operators. The leading B-violating operators [1] have dimension six and are hence suppressed by two powers of an inverse mass scale. These operators arise naturally when SM is embedded in grand unified theories (GUTs) such as $SU(5)$ and $SO(10)$. They lead to nucleon decay modes such as $p \to e^+\pi^0$ and $p \to \bar{\nu}K^+$, which conserve baryon number minus lepton number, B-L symmetry. Experimental searches to date have primarily focused on these modes with the latest limits on proton lifetime constraining the masses of the heavy mediators to be larger than about 10^{15} GeV. This is in accord with the scale inferred from the unification of gauge couplings.

Going beyond the $d = 6$ B-violating operators, the next-to-leading ones have $d = 7$ and obey the selection rule $\Delta(B\text{-}L) = -2$ for nucleon decay [2]. These operators lead to novel nucleon decay modes such as $n \to e^- K^+$, $e^-\pi^+$, and $p \to \nu\pi^+$, which have received less attention. In this Letter, we show that these $d = 7$ operators arise naturally in unified theories based on $SO(10)$, upon the spontaneous breaking of B-L, which is part of the gauge symmetry. In particular, we find that in nonsupersymmetric $SO(10)$ models with an intermediate scale so that gauge couplings unify, the partial lifetime to these decay modes can be within reach of ongoing and proposed experiments. Furthermore, we show that these new modes provide a novel way to understand the origin of matter in the Universe. This mechanism relies on the fact that owing to their B-L breaking nature, a GUT scale-induced baryon asymmetry would not be affected by the electroweak sphalerons [3] and would survive down to low temperatures. The observed baryon number of the Universe then would carry the direct imprint of GUT scale physics. This is unlike the B-L–preserving baryon asymmetry induced in the decays of GUT mass particles such as in $SU(5)$, which is, however, washed out by the sphaleron interactions, leaving no trace of GUT physics. We show that in minimal $SO(10)$ models [4], which have been highly successful in predicting large neutrino oscillation angles, including a relatively large value of $\sin^2 2\theta_{13} \simeq (0.085\text{–}0.095)$, consistent with recent results [5], the baryon asymmetry of the right magnitude is generated by the new B-L–violating mechanism. The results of this Letter should provide motivations to search for B-L–violating semileptonic decay modes of the nucleon in the ongoing and next round of experiments. Their observation would furnish evidence against the simple one-step breaking of GUT symmetry, and could also resolve the mystery behind the origin of matter in the Universe.

We start by writing down the $d = 7$ B–violating effective operators in the SM [2] in the standard notation for fermion fields:

$$\begin{aligned}
\tilde{O}_1 &= (d^c u^c)^*(d^c L_i)^* H_j^* \epsilon_{ij}, \\
\tilde{O}_2 &= (d^c d^c)^*(u^c L_i)^* H_j^* \epsilon_{ij}, \\
\tilde{O}_3 &= (Q_i Q_j)(d^c L_k)^* H_l^* \epsilon_{ij}\epsilon_{kl}, \\
\tilde{O}_4 &= (Q_i Q_j)(d^c L_k)^* H_l^* (\vec{\tau}\epsilon)_{ij} \cdot (\vec{\tau}\epsilon)_{kl}, \\
\tilde{O}_5 &= (Q_i e^c)(d^c d^c)^* H_i^*, \\
\tilde{O}_6 &= (d^c d^c)^*(d^c L_i)^* H_i, \\
\tilde{O}_7 &= (d^c D_\mu d^c)^*(\bar{L}_i \gamma^\mu Q_i), \\
\tilde{O}_8 &= (d^c D_\mu L_i)^*(\bar{d}^c \gamma^\mu Q_i), \\
\tilde{O}_9 &= (d^c D_\mu d^c)^*(\bar{d}^c \gamma^\mu e^c).
\end{aligned} \tag{1}$$

Here D_μ stands for the covariant derivative, and $H(1, 2, 1/2)$ is the Higgs doublet. These operators obey $(B\text{-}L) = +2$ selection rule and mediate nucleon decays of the type $n \to e^- K^+$, $e^- \pi^+$, and $p \to \nu \pi^+$. We first show how these operators arise naturally in $SO(10)$ theories [6] when B-L symmetry contained in it is broken. This breaking may occur at the GUT scale as in models with supersymmetry, or at an intermediate scale M_I below the GUT scale, as in nonsupersymmetric $SO(10)$, which requires such a scale to be compatible with gauge coupling unification. To see the origin of Eq. (1) in $SO(10)$ via scalar boson exchange, we write down the Yukawa couplings in the most general setup. Noting that the fermion bilinears contain $16 \times 16 = 10_s + 120_a + 126_s$, the Yukawa couplings are [7]

$$\mathcal{L}(16_i 16_j 10_H) = h_{ij}\left[(u_i^c Q_j + \nu_i^c L_j)h - (d_i^c Q_j + e_i^c L_j)\bar{h} + \left(\frac{\epsilon}{2} Q_i Q_j + u_i^c e_j^c - d_i^c \nu_j^c\right)\omega + (\epsilon u_i^c d_j^c + Q_i L_j)\omega^c\right], \quad (2)$$

$$\begin{aligned}\mathcal{L}(16_i 16_j \overline{126}_H) = f_{ij}[&(u_i^c Q_j - 3\nu_i^c L_j)h - (d_i^c Q_j - 3e_i^c L_j)\bar{h} + \sqrt{3}i\left(\frac{\epsilon}{2} Q_i Q_j - u_i^c e_j^c + \nu_i^c d_j^c\right)\omega_1 \\ &+ \sqrt{3}i(Q_i L_j - \epsilon u_i^c d_j^c)\omega_1^c + \sqrt{6}(d_i^c \nu_j^c + u_i^c e_j^c)\omega_2 + 2\sqrt{3}i d_i^c L_j \rho - 2\sqrt{3}i \nu_i^c Q_j \bar{\rho} \\ &+ 2\sqrt{3} u_i^c \nu_j^c \eta - 2\sqrt{3}i u_i^c L_j \chi + 2\sqrt{3}i e_i^c Q_j \bar{\chi} - 2\sqrt{3} d_i^c e_j^c \delta + \sqrt{6}i Q_i L_j \bar{\Phi} + \ldots],\end{aligned} \quad (3)$$

$$\begin{aligned}\mathcal{L}(16_i 16_j 120_H) = g_{ij}[&(d_i Q^j + e_i^c L_j)\bar{h}_1 - (u_i^c Q_j + \nu_i^c L_j)h_1 - \sqrt{2} Q_i L_j \omega_1^c - \sqrt{2}(u_i^c e_j^c - d_i^c \nu_j^c)\omega_1 \\ &- \frac{i}{\sqrt{3}}(d_i^c Q_j - 3e_i^c L_j)\bar{h}_2 + \frac{i}{\sqrt{3}}(u_i^c Q_j - 3\nu_i^c L_j)h_2 - 2e_i^c Q_j \bar{\chi} + 2\nu_i^c Q_j \bar{\rho} - 2d_i^c L_j \rho + 2u_i^c L_j \chi \\ &- i\epsilon d_i^c d_j^c \bar{\eta} + 2i u_i^c \nu_j^c \eta + \sqrt{2}i\epsilon d_i^c u_j^c \omega_2^c + \sqrt{2}i(d_i^c \nu_j^c - e_i^c u_j^c)\omega_2 - \frac{\epsilon}{\sqrt{2}} Q_i Q_j \Phi - \sqrt{2} Q_i L_j \bar{\Phi} \\ &- 2i d_i^c e_j^c \delta + i\epsilon u_i^c u_j^c \bar{\delta} + \ldots],\end{aligned} \quad (4)$$

with h and f being symmetric and g being antisymmetric in flavor indices i, j. The $SU(3)_C \times SU(2)_L \times U(1)_Y$ quantum numbers of the various submultiplets in Eqs. (2)–(4) are $h(1,2,+1/2)$, $\bar{h}(1,2,-1/2)$, $\omega(3,1,-1/3)$, $\omega^c(\bar{3},1,1/3)$, $\rho(3,2,1/6)$, $\bar{\rho}(\bar{3},2,-1/6)$, $\eta(3,1,2/3)$, $\bar{\eta}(\bar{3},1,-2/3)$, $\Phi(3,3,-1/3)$, $\bar{\Phi}(\bar{3},3,1/3)$, $\chi(3,2,7/6)$, $\bar{\chi}(\bar{3},2,-7/6)$, $\delta(3,1,-4/3)$, and $\bar{\delta}(\bar{3},1,4/3)$.

The B-L generator of $SO(10)$ is broken by the vacuum expectation value (VEV) of the SM singlet field Δ^c in $\overline{126}_H$, which has B-$L = -2$. This VEV supplies large Majorana masses for the right-handed neutrinos through the coupling $f_{ij}\sqrt{6}\nu_i^c \nu_j^c \Delta^c$. It also generates trilinear scalar couplings of the type $\rho^* \omega H$, $\eta^* \rho H$, $\rho^* \Phi H$, and $\chi^* \eta H$, thereby inducing the $d = 7$ baryon number-violating operators of Eq. (1), via the Yukawa couplings of Eqs. (2)–(4). The flavor-symmetric Yukawa couplings of Eqs. (2) and (3) generate the operators $\tilde{\mathcal{O}}_3$ and $\tilde{\mathcal{O}}_1$ through the diagrams shown in Fig. 1.

The trilinear couplings in Fig. 1 have different sources in $SO(10)$. The quartic coupling $(126)^4$, which is invariant, contains the term $(2,2,15) \times (2,2,15) \times (1,1,6) \times (1,3,\overline{10})$ under $SU(2)_L \times SU(2)_R \times SU(4)_C$ subgroup. The $\rho^*(\bar{3},2,-1/6) \subset (2,2,15)$, $H(1,2,1/2) \subset (2,2,15)$, while $\omega(3,1,-1/3) \subset (1,1,6)$, and this coupling would contain the term $\rho^* \omega H \bar{\Delta}^c$. The three nontrivial invariants of the type $(126)^2 \times (126^*)^2$ also contain this trilinear term. Similarly, the three quartic couplings $(120)^2 \times (126)^2$ would generate the trilinear terms $\rho^* \eta H^*$, $\chi^* \eta H$, and $\Phi^* \rho H^*$ vertices, and along with Eq. (4) would induce the remaining nonderivative $d = 7$ operators of Eq. (1). For a detailed discussion, see Ref. [6].

Such trilinear couplings are also present when the 126_H is replaced by a 16_H albeit in a slightly different way. The 16_H contains an SM singlet filed with B-$L = +1$, which acquires a GUT scale VEV and a $\bar{h}(1,2,-1/2)$ field with B-$L = -1$. The trilinear scalar couplings of the types $16_H 16_H 10_H$ and $\overline{16}_H \overline{16}_H 10_H$ would mix the B-$L = 0$ Higgs doublet $h(1,2,1/2)$ from the 10_H and the $h(1,2,1/2)$ Higgs doublet from the $\overline{16}_H$, which has B-$L = +1$. The light SM Higgs doublet then would have no definite $B - L$ quantum number. The (1, 2, 4) component of $\overline{16}_H$ contains the field $\rho^*(\bar{3},2,-1/6)$, and the (2, 1, $\bar{4}$) of $\overline{16}_H$ contains $\omega(3,1,-1/3)$, and thus the coupling $\rho^* \omega H$ is generated via the $\overline{16}_H \overline{16}_H 10_H$ coupling.

The $d = 7$ operators of Eq. (1) can also arise by integrating out the vector gauge bosons $V_Q(3,2,1/6)$ and $V_{u^c}(\bar{3},1,-2/3)$ of $SO(10)$, which lie outside of $SU(5)$ [6].

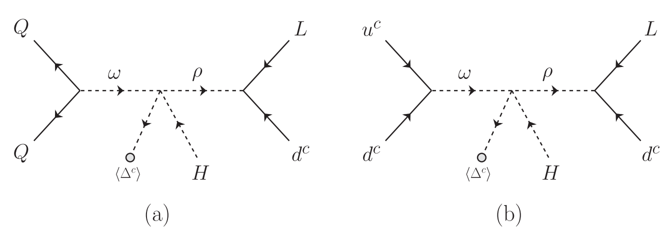

FIG. 1. Effective baryon number violating $d = 7$ operators induced by the symmetric Yukawa couplings of 10_H and $\overline{126}_H$ of $SO(10)$.

The covariant derivative for the 126_H would contain the term $V_Q V_{u^c} H(\Delta^c)^\dagger$, which generates the $d = 7$ operators. When 16_H is used instead of the 126_H, the covariant derivative would contain a similar term, but now with B-$L = +1$ and -1 for H and $\Delta^c \subset 16_H$, respectively.

Partial lifetime for B-L violating nucleon decay.—The diagrams of Fig. 1 lead to the following estimate for $n \to e^- \pi^+$ lifetime: ($B + L$–preserving nucleon decay has been studied in the context of R-parity breaking supersymmetry in Ref. [8].)

$$\Gamma(n \to e^- \pi^+)^{\text{Fig.1}} \approx \frac{|Y^*_{QQ\omega} Y_{Ld^c\rho}|^2}{64\pi}(1 + D + F)^2 \frac{\beta_H^2 m_p}{f_\pi^2}\left(\frac{\lambda v v_R}{M_\rho^2}\right)^2 \frac{1}{M_\omega^4}. \tag{5}$$

Here, we have defined the Yukawa couplings of the ω and ρ fields appearing in Fig. 1 to be $Y^*_{QQ\omega}$ and $Y_{Ld^c\rho}$. The factors D and F are chiral Lagrangian factors $D \simeq 0.8$ and $F \simeq 0.47$; $\beta_H \simeq 0.012\ \text{GeV}^3$ is the nucleon decay matrix element $v_R \equiv \langle \Delta^c \rangle$, and $v \equiv \langle H^0 \rangle = 174$ GeV. We have defined the trilinear coupling of Fig. 1 to have a coefficient λv_R. The mass of $\omega(3, 1, -1/3)$ is constrained to be relatively large, as it mediates $d = 6$ nucleon decay. For $Y \approx 10^{-3}$, $M_\omega > 10^{11}$ GeV must be met from the $d = 6$ decays. As an illustration, choosing $Y_{QQ\omega} = Y_{Ld^c\rho} = 10^{-3}$, $M_\omega = 10^{11}$ GeV, $M_\rho = 10^8$ GeV, and $\lambda v_R = 10^{11}$ GeV in Eq. (5), we find $\tau_n \approx 3 \times 10^{33}$ yr. Such a spectrum is motivated by the intermediate symmetry $SU(2)_L \times SU(2)_R \times SU(4)_C$, which is found to be realized at $M_I \approx 10^{11}$ GeV from gauge-coupling unification [9]. As a second example, take $M_\rho = 10^6$ GeV, $M_\omega = 10^{16}$ GeV, $\lambda v_R = 10^{16}$ GeV, and $Y_{QQ\omega} = Y_{Ld^c\rho} = 3 \times 10^{-3}$. This choice of spectrum leads to $\tau_n \approx 4 \times 10^{33}$ yr. This spectrum can arise as follows. Suppose the $\rho(3, 2, 1/6)$ particle, along with a pair of (1, 3, 0) scalar particles (contained in the 45_H, 54_H, or 210_H needed for symmetry breaking) survive down to $M_I = 10^6$ GeV. The SM gauge couplings are found to unify at a scale $M_X \approx 10^{15}$ GeV in this case, as shown in Fig. 2. This scenario would predict observable rates for both the B-L–conserving and B-L–violating nucleon decay modes. Analogous results are obtained from the exchange of $\eta(3, 1, 2/3) - \rho(3, 2, 1/6)$ scalar bosons from 120_H. Since these particles do not induce $d = 6$ baryon number violation, they can both have mass of order M_I, which would enhance the nucleon decay rate.

While the $\Delta(B\text{-}L) = -2$ nucleon lifetime is in the experimentally accessible range for reasonable choice of parameters as shown, it is quite sensitive to the precise values of the intermediate scalar masses. For example, an increase in M_ρ and M_ω by a factor of 3 will increase the lifetime by a factor of 10^4. Not finding these modes will not exclude this class of $SO(10)$ models, but a discovery of the $\Delta(B\text{-}L) = -2$ nucleon decay mode would lend strong

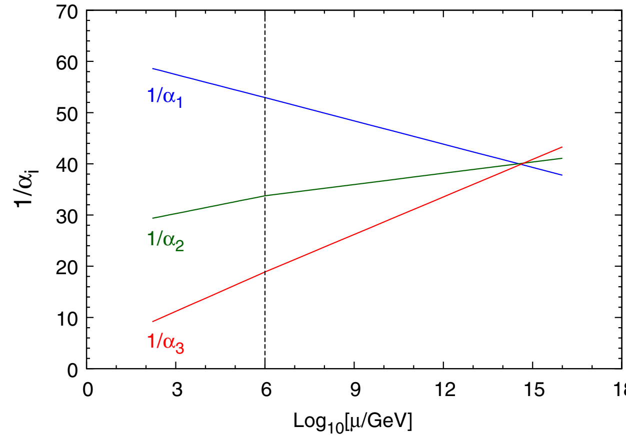

FIG. 2 (color online). Unification of the three SM gauge couplings obtained with a light $\rho(3, 2, 1/6)$ and two (1, 3, 0) scalar multiplets at $M_I = 10^6$ GeV.

support to a new mechanism of baryogenesis via B-L–violating decays of scalars, to which we now turn.

New baryogenesis scenario at the GUT epoch.—We now present a new baryogenesis scenario at the GUT epoch, using the B-L–violating decay of the scalar $\omega(3, 1, -1/3)$ with a GUT scale mass. The magnitude of the asymmetry is directly linked to the neutrino masses, since the Yukawa couplings that induce the asymmetry are the same couplings that are involved in neutrino mass generation. (B-L asymmetry in decays of specific heavy particles has recently been discussed in Ref. [10].) To be concrete, we shall work in the framework of nonsupersymmetric $SO(10)$, although our results would hold for SUSY $SO(10)$ as well, with some minor modifications. The couplings of Eqs. (2)–(4) imply that ω has two-body decays into fermions of the type $\omega \to \bar{Q}\,\bar{Q}$, $\bar{u}^c \bar{e}^c$, $\bar{\nu}^c \bar{d}^c$, $u^c d^c$, QL. These decays preserve B-L, as can be seen by assigning $(B\text{-}L)(\omega) = -2/3$. Now, ω also has a two-body scalar decay $\omega \to \rho H^*$, as shown in Fig. 3(a), which uses the

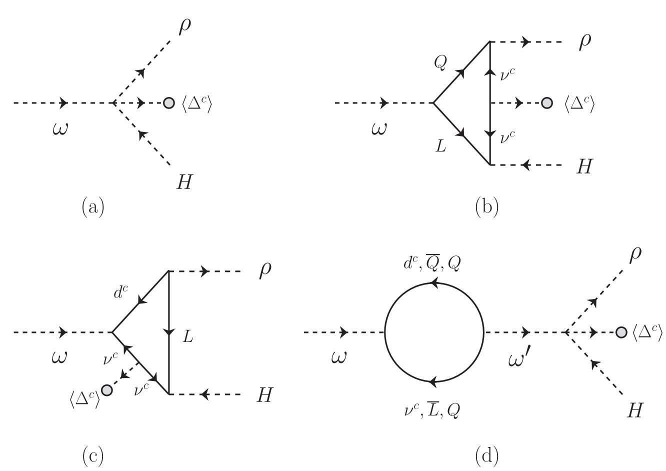

FIG. 3. Tree-level diagram and one-loop corrections responsible for generating B-L asymmetry in ω decay.

B-L breaking VEV of Δ^c. The scalar field ρ has two-body fermionic decays of the type $\rho \to \bar{L}\bar{d}^c$, $\nu^c Q$ (the latter if kinematically allowed), which define B-L charge of ρ to be $+4/3$. Thus, the decay $\omega \to \rho H^*$ would violate B-L by -2.

Let the branching ratio for $\omega \to \rho H^*$ be r, which produces a net B-L number of $4/3$, and that for $\omega^* \to \rho^* H$ be $\bar{r}$, with net B-$L = -4/3$. The branching ratio for the two-fermion decays $\omega \to ff$ is then $(1-r)$, which has B-$L = -2/3$, and that for $\omega^* \to \bar{f}\bar{f}$ is $(1-\bar{r})$, which has B-$L = 2/3$. Thus, in the decay of a $\omega + \omega^*$ pair, a net B-L number, defined as ϵ_{B-L}, is induced, with $\epsilon_{B\text{-}L} \equiv (B\text{-}L)_\omega + (B\text{-}L)_{\omega^*} = 2(r\text{-}\bar{r})$. The loop diagrams for $\omega \to \rho H^*$ are shown in Figs. 3(b) and 3(d), which involve the exchange of fermions. Since ω can also decay to two on-shell fermions, these loop diagrams have absorptive parts and also CP violation.

We evaluate Fig. 3 in a basis where the Majorana mass matrix M_{ν^c} of the ν^c fields is diagonal and real. The contributions of Figs. 3(b) and 3(d) to $\epsilon_{B\text{-}L}$ are found to be

$$\epsilon_{B\text{-}L} = -\frac{\mathrm{Br}}{\pi|\lambda v_R|^2}\Im[\lambda v_R \mathrm{Tr}\{Y^\dagger_{QL\omega^*}Y_{Q\nu^c\bar{\rho}}M_{\nu^c}F(M_\omega, M_\rho, M_{\nu^c_j})Y_{\nu^c LH} - Y_{d^c\nu^c\omega}Y^\dagger_{d^cL\rho}Y_{\nu^c LH}M_{\nu^c}F(M_\rho, M_\omega, M_{\nu^c_j}) - Y^\dagger_{d^c\nu^c\omega'}Y_{d^c\nu^c\omega}F'(M_\omega, M_{\omega'}, M_j)(\lambda' v_R)^*\}], \tag{6}$$

where the three terms are in order from Figs. 3(b)–3(d), respectively. Here, we have defined the trilinear scalar vertices of Figs. 3(a) and 3(d) to have coefficients of λv_R and $\lambda' v_R$ in the Lagrangian. $Y_{Q\nu^c\bar{\rho}}$ is the Yukawa-coupling matrix corresponding to the coupling $Q\nu^c\bar{\rho}$, etc.; Br stands for the branching ratio $\mathrm{Br}(\omega \to \rho H^*)$. A factor of 2 has been included here for the two $SU(2)_L$ final states in the decay. The functions F and F' are defined as

$$F(a,b,c) = \ln(1 + a^2/c^2) + \Theta(1 - c^2/b^2)(1 - c^2/b^2),$$
$$F'(a,b,c) = (1 - c^2/a^2)/(1 - b^2/a^2)\Theta(1 - c^2/a^2) \times (1 - c^2/a^2). \tag{7}$$

Here Θ stands for the step function, signaling additional ways of cutting the diagram when $M_j < M_\rho$ or $M_j < M_\omega$ in Fig. 3. Figure 3(d) arises because in any realistic $SO(10)$ model there are at least two ω fields. The heavier ω field is denoted as ω'. We have also assumed that $M_\omega - M'_\omega \gg \Gamma_\omega$, so that there is no resonant enhancement for the decay.

To estimate $\mathrm{Br} = \mathrm{Br}(\omega \to \rho H^*)$ appearing in Eq. (6), let us assume that ω is the field ω from 10_H with Yukawa couplings as given in Eq. (2). The partial widths for the decays $\Gamma_1(\omega \to \rho H^*)$ and $\Gamma_2(\rho \to ff)$ are then given by

$$\Gamma_1(\omega \to \rho H^*) = \frac{|\lambda v_R|^2}{8\pi M_\omega}\left(1 - \frac{M_\rho^2}{M_\omega^2}\right),$$
$$\Gamma_2(\omega \to ff) = \frac{\mathrm{Tr}(h^\dagger h)}{4\pi}M_\omega, \tag{8}$$

with $\mathrm{Br} = \Gamma_1/(\Gamma_1 + \Gamma_2)$. For $M_\omega = 10^{16}$ GeV, $h_{33} = 0.6$ (corresponding to the top quark Yukawa coupling at GUT scale) with other h_{ij} negligible, and $\lambda v_R = (10^{14}, 10^{15}, 10^{16})$ GeV, one gets $\mathrm{Br} = (1.4 \times 10^{-4}, 1.4 \times 10^{-2}, 0.58)$.

The B-L asymmetry $\epsilon_{B\text{-}L}$ of Eq. (6) will result in a baryon to entropy ratio Y_B given by

$$Y_B \equiv \frac{n_B\text{-}n_{\bar{B}}}{s} = \frac{\epsilon_{B\text{-}L}}{g_*}d, \tag{9}$$

where $g_* = 130$ is the total number of relativistic degrees of freedom at the epoch when these decays occur. The factor d in Eq. (9) is the dilution factor, which takes into account back reactions that would partially wash out the induced baryon asymmetry. Defining $K = \frac{\Gamma(\omega \to \rho H^*)}{2H}|_{T=M_\omega}$, where H is the Hubble expansion rate, $H = 1.66 g_*^{1/2}\frac{T^2}{M_{\mathrm{Pl}}}$, the dilution factors can be written as [11] $d \simeq 1 (K < 1)$ and $d \simeq \frac{0.3}{K(\ln K)^{0.6}} (K \gg 1)$. For $M_\omega = 10^{15}$ GeV and $\lambda v_R = (10^{14}, 10^{15}, 10^{16})$ GeV, we find $K = (0.12, 12.3, 1230)$ and the corresponding dilution factors to be $d = (1.0, 1.4 \times 10^{-2}, 7.5 \times 10^{-5})$, with $\mathrm{Br} = (1.3 \times 10^{-2}, 0.58, 1.0)$.

We now show how the GUT scale-induced baryon asymmetry in $\omega \to \rho H^*$ decay can consistently explain the observed value of $Y_B = (8.75 \pm 0.23) \times 10^{-11}$ in a class of minimal $SO(10)$ models. In these models, a single 10_H and a single $\overline{126}_H$ couple to fermions, as in Eqs. (2) and (3). It has been shown that these models lead to large mixing angles for solar and atmospheric neutrino oscillations. Furthermore, they predict $\sin^2 2\theta_{13} \approx (0.085$–$0.095)$, both in the non-SUSY and the SUSY versions [4], which is consistent with recent results from Daya Bay and other experiments [5]. To illustrate how a realistic choice of parameters can generate acceptable Y_B, we choose the ω field to be almost entirely in the 10_H. We also choose $\lambda' v_R$ that appears in Fig. 3(d) to be small, so that the leading contribution to $\epsilon_{B\text{-}L}$ is from Fig. 3(c). In this limit, we find $\epsilon_{B\text{-}L} \approx \frac{2\sqrt{3}}{\pi}\frac{|h_{33}f_3|^2}{|\lambda|}\{1 + \ln(1 + M_\rho^2/M_{\nu_3^c}^2)\}\sin\phi$. Here, we have kept only the third family Yukawa couplings and defined $\phi = \arg\{h_{33}^2 f_3^2 \lambda + \frac{\pi}{2}\}$. Choosing $h_{33} \simeq 0.6$ (the top quark Yukawa coupling at the GUT scale) and $\lambda = 0.25$, $v_R = 10^{16}$ GeV, $f_3 = 10^{-2}$ (so that $f_3 v_R = 10^{14}$ GeV, consistent with the light ν_τ mass arising via the seesaw mechanism), $\phi = 0.12$, we find $\epsilon_{B\text{-}L} = 1.6 \times 1.9 \times 10^{-5}$. If $M_\omega = 10^{15}$ GeV, then $\mathrm{Br} = 0.96$ and $K = 197$, so that the dilution factor is $d = 5.6 \times 10^{-4}$. This results in a net $Y_B = 8.2 \times 10^{-11}$, consistent with observations. We emphasize the intimate connection between ϵ_{B-L} and neutrino masses, since $Y^{\dagger}_{QL\omega^*}$, $Y_{Q\nu^c\bar{\rho}}$, etc., present in $\epsilon_{B\text{-}L}$ are the **126** couplings that determine the neutrino masses via the seesaw mechanism.

In conclusion, we have shown that all $d = 7$ baryon number-violating operators that lead to nucleon decay modes such as $n \to e^- K^+$, $e^- \pi^+$, and $p \to \nu \pi^+$ emerge naturally as effective low-energy operators in a wide class of $SO(10)$ models. In nonsupersymmetric $SO(10)$ models with an intermediate scale, we find the rates for these nucleon decay modes to be within reach of experiments. We have also shown that the existence of these B-L–violating interactions allows a new scenario for baryogenesis where a B-L asymmetry is generated in the decay of GUT mass particles that survives to low temperatures, unaffected by the sphaleron interactions. In minimal $SO(10)$ models that predict large neutrino mixing angles, including θ_{13}, this new mechanism can explain the observed baryon asymmetry of the Universe.

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*babu@okstate.edu
†rmohapat@umd.edu

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