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Observable neutron–antineutron oscillations in seesaw models of neutrino mass

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Abstract

We show that in a large class of supersymmetric models with spontaneously broken $B-L$ symmetry, neutron–antineutron oscillations occur at an observable level even though the scale of $B-L$ breaking is very high, $v_{B-L} \sim 2 \times 10^{16}$ GeV, as suggested by gauge coupling unification and neutrino masses. We illustrate this phenomenon in the context of a recently proposed class of seesaw models that solves the strong CP problem and the SUSY phase problem using parity symmetry. We obtain an *upper* limit on $N-\bar{N}$ oscillation time in these models, $\tau_{N-\bar{N}} \leq 10^9-10^{10}$ s. This suggests that a modest improvement in the current limit on $\tau_{N-\bar{N}}$ of 0.86×10^8 s will either lead to the discovery of $N-\bar{N}$ oscillations, or will considerably restrict the allowed parameter space of an interesting class of neutrino mass models.

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1. Introduction

It is widely believed that the most natural and appealing explanation of the recent neutrino oscillation results is provided by the seesaw mechanism [1] incorporated into extensions of the Standard Model that include a local $B-L$ symmetry. The simplest models with local $B-L$ symmetry are the left–right symmetric models [2] based on the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. These models have the additional virtue that they explain the origin of parity violation in weak interactions as a consequence of spontaneous symmetry breaking in very much the same way as one explains the strength of the weak interaction in the Standard Model. Stability of the

Higgs sector under radiative corrections calls for weak scale supersymmetry as in the Minimal Supersymmetric Standard Model (MSSM). It has recently been shown that if the MSSM is embedded into a left–right symmetric framework at a high scale $v_R \sim 10^{14}-10^{16}$ GeV, as suggested by neutrino oscillation data and by gauge coupling unification, it helps solve some important problems faced by the MSSM, viz., the SUSY CP problem [3], the strong CP problem [4] and the μ problem. Supersymmetric models with such a high scale embedding are, therefore, attractive candidates for physics beyond the Standard Model.

It was noted many years ago [5] that the electric charge formula of the left–right symmetric models, $Q = I_{3L} + I_{3R} + \frac{B-L}{2}$, allows one to conclude from pure group theoretic arguments that parity symmetry breaking implies a breakdown of $B-L$ symmetry as well with the constraint that $2\Delta I_{3R} = -\Delta(B-L)$. This simple relation is profoundly revealing. It says

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that the neutrinos must be Majorana particles since the lepton number breaking terms in the theory must obey $|\Delta L| = 2$ selection rule. This conclusion follows directly if Higgs triplets are used to break $SU(2)_R$ symmetry since $I_{3R} = 1$ for triplets, it also holds when Higgs doublets are used for this purpose, since gauge invariance requires the presence of two such doublets in the mass term for the neutrinos. Secondly, for purely hadronic baryon number violating processes, baryon number must change by at least two units, $|\Delta B| = 2$. This means that models based on left–right symmetric gauge structure can lead to the process where a neutron transforms itself into an antineutron ($N-\bar{N}$ oscillation [5–7]), while they may forbid the decay of the proton, which is a $\Delta B = 1$ process.

While the above group theory argument predicts the existence of $N-\bar{N}$ oscillation in left–right symmetric models, its strength will depend on the details of the model. Using simple-dimensional analysis it is easy to find that the lowest-dimensional operators that contributes to $N-\bar{N}$ oscillation are six quark operators, a typical one being $(u^c d^c u^c d^c d^c d^c)$. This operator has dimension 9 and, therefore, the coupling strength scales as $G_{\Delta B=2} \sim \frac{1}{M^5}$, where M is the scale of new physics. It is natural to identify M with the scale of $B-L$ (or parity) breaking. The current lower limit on $N-\bar{N}$ oscillation time, $\tau_{N-\bar{N}} \geq 0.86 \times 10^8$ s [8],¹ implies an upper limit $G_{\Delta B=2} \leq 3 \times 10^{-28}$ GeV⁻⁵. For $N-\bar{N}$ oscillations to be observable then, the scale M should be rather low, $M \leq 10^6$ GeV.

One class of models where $\Delta B = 2$ transition manifests itself through Higgs boson exchange has been discussed in Ref. [5]. There it was shown that if the $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model is embedded into the $SU(4)_C \times SU(2)_L \times SU(2)_R$ gauge group, then $N-\bar{N}$ oscillations can arise at an observable level if the $SU(4)_C$ breaking scale is in the 100 TeV range. In these models, $N-\bar{N}$ oscillation amplitude is intimately tied to an understanding of small neutrino masses via the seesaw mechanism as well as the breaking of quark–lepton degeneracy implied by $SU(4)_C$ symmetry. The same Higgs field

that breaks $SU(4)_C$ and generates heavy Majorana masses for the right-handed neutrinos also mediate $N-\bar{N}$ oscillations here. With the scale of $SU(4)_C$ breaking in the 100 TeV range, these models would appear to be incompatible with gauge coupling unification. Furthermore, such a low scale of parity breaking would not yield naturally neutrino masses in the range suggested by current experiments. If we raise the scale of parity/ $SU(4)_C$ breaking to values above 10^{12} GeV, so that small neutrino masses in the right range are generated naturally, then $N-\bar{N}$ transition amplitude becomes unobservably small in these models.²

Does the above arguments mean that $N-\bar{N}$ oscillations are beyond experimental reach based on current neutrino oscillation phenomenology? In this Letter we will show that this is not the case in a class of attractive seesaw models with local $B-L$ symmetry. We will see that in these models a new class of $\Delta B = 1$ operators is induced as a consequence of parity breaking. These operators lead to observable $N-\bar{N}$ oscillation despite the scale v_R of parity breaking being close to the conventional GUT scale of 2×10^{16} GeV. In fact, $G_{\Delta B=2}$ increases with v_R and, therefore, one has the inverse phenomenon that increasing v_R leads to stronger $N-\bar{N}$ oscillation amplitude. Interestingly, the scale v_R implied by neutrino masses is such that $N-\bar{N}$ oscillation should be accessible experimentally with a modest improvement in the current limit. We obtain an *upper* limit of $\tau_{N-\bar{N}} \leq 10^8-10^{10}$ s in this class of models. This prediction becomes sharper in a concrete model where flavor symmetries reduce considerably the uncertainties in the estimate of $\tau_{N-\bar{N}}$. We emphasize that our upper limit is derived in the context of conventional seesaw models of neutrino mass without using any special ingredients to enhance $N-\bar{N}$ oscillation amplitude. This should provide new impetus for an improved experimental search for $N-\bar{N}$ oscillations.

¹ This is the direct limit from free neutron oscillation searches. Indirect limits which involves some reasonable nuclear physics assumptions have been extracted from nucleon decay experiments which are slightly more stringent: $\tau_{N-\bar{N}} \geq 1.2 \times 10^8$ s [9].

² A counter example where a higher scale of parity violation can go hand in hand with observable $N-\bar{N}$ oscillation was noted in the context of a SUSY $SU(2)_L \times SU(2)_R \times SU(4)_C$ model in Ref. [10]. These models possess accidental symmetries that lead to light (~ 100 GeV) diquark Higgs bosons even though the scale of parity violation is high. As a result, the $N-\bar{N}$ oscillation operator can have observable strength. Unification of gauge couplings is, however, difficult to achieve in these models.

2. High scale seesaw model and $N-\bar{N}$ oscillation

The basic framework of our model involves the embedding of the MSSM into a minimal SUSY left–right gauge structure at a scale v_R close to the GUT scale. The electroweak gauge group of the model, as already mentioned, is $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ with the standard assignment of quarks and leptons—left-handed quarks and leptons (Q, L) transform as doublets of $SU(2)_L$, while the right-handed conjugate ones (Q^c, L^c) are doublets of $SU(2)_R$. The quarks Q transform under the gauge group as $(2, 1, 1/3)$ and Q^c as $(1, 2, -1/3)$, while the lepton fields L and L^c transform as $(2, 1, -1)$ and $(1, 2, +1)$, respectively. The Dirac masses of fermions arise through their Yukawa couplings to two Higgs bidoublet $\Phi_a(2, 2, 0)$, $a = 1 - 2$. The $SU(2)_R \times U(1)_{B-L}$ symmetry is broken down to $U(1)_Y$ in the supersymmetric limit by $B-L = \pm 1$ doublet scalar fields, the right-handed doublet denoted by $\chi^c(1, 2, -1)$ accompanied by its left-handed partner $\chi(2, 1, 1)$. Anomaly cancellation requires the presence of their charge conjugate fields as well, denoted as $\bar{\chi}^c(1, 2, 1)$ and $\bar{\chi}(2, 1, -1)$. The vacuum expectation values (VEVs) $\langle \chi^c \rangle = \langle \bar{\chi}^c \rangle = v_R$ break the left–right symmetry group down to the MSSM gauge symmetry. A singlet S is also used to facilitate symmetry breaking in the SUSY limit.

It has recently been shown that if there exists a Z_4 R symmetry, the minimal model just described will solve the strong CP problem and the SUSY phase problem based on parity symmetry [4]. Furthermore, the μ term will have a natural origin. One possible Z_4 assignment was given in Ref. [4]. Here we present a slight variant, which yields the same superpotential at the renormalizable level as in Ref. [4] and thus preserves all its success. Under this Z_4 , the superpotential W changes sign, as do $d^2\theta$ and $d^2\bar{\theta}$. The quark fields (Q, Q^c) are even, while (L, L^c) transform as $(i, -i)$. The fields ($\chi, \bar{\chi}, \chi^c, \bar{\chi}^c, \Phi_a, S$) are all odd under Z_4 .

The gauge invariant superpotential consistent with this Z_4 R symmetry at the renormalizable level is

$$\begin{aligned} W = & h_a Q \Phi_a Q^c + h'_a L \Phi_a L^c + \lambda_a \chi \Phi_a \chi^c \\ & + \lambda'_a \bar{\chi} \Phi_a \bar{\chi}^c \\ & + \kappa S (e^{i\xi} \chi^c \bar{\chi}^c + e^{-i\xi} \chi \bar{\chi} + a S^2 - M^2) \\ & + \mu_{ab} \text{Tr}(\Phi_a \Phi_b) S. \end{aligned} \quad (1)$$

This superpotential breaks the gauge symmetry to that of the Standard Model in the SUSY limit without leaving any unwanted Goldstone bosons and induces realistic quark masses and mixings.

The baryon number violating processes as well as neutrino masses arise in this model from higher-dimensional operators induced by Planck scale physics. They will be the main focus of the rest of the Letter. We shall pay special attention to the relation between the neutrino mass and the $N-\bar{N}$ oscillation time. The relevant dimension four operators in the superpotential which are scaled by M_{Pl}^{-1} and are allowed by the Z_4 symmetry are:

$$\begin{aligned} \mathcal{O}_1 &= f [(L^c \chi^c)^2 + (L \chi)^2], \\ \mathcal{O}_2 &= f' [Q^c Q^c Q^c \bar{\chi}^c + Q Q Q \bar{\chi}]. \end{aligned} \quad (2)$$

Operator \mathcal{O}_1 gives rise to Majorana masses for ν_R of order v_R^2/M_{Pl} . Combining this with the Dirac neutrino masses arising from Eq. (1), small neutrino masses will be generated by the seesaw mechanism. For $v_R \sim 10^{14}-10^{16}$ GeV, the magnitude of the light neutrino masses are in the right range to explain the atmospheric and the solar neutrino oscillation data. Operator \mathcal{O}_2 , which is also invariant under the Z_4 , leads to baryon number violation. While $\mathcal{O}_{1,2}$ could have their origin in quantum gravity, they may also be induced by integrating out vector states that have Z_4 -invariant masses of order the Planck scale.

Note that operators such as $L \Phi \chi^c, L L L^c \chi^c, Q L Q^c \chi^c$ are not allowed by the Z_4 symmetry. If they were present along with \mathcal{O}_2 , they would lead to rapid proton decay. Note also that the well known proton decay operator $Q Q Q L$ is not allowed by the Z_4 symmetry. In any case its presence would not have been a problem since it is scaled by the Planck mass and therefore can lead to a proton lifetime consistent with the present lower limit.

To see the connection between neutrino masses and the $N-\bar{N}$ oscillation time $\tau_{N-\bar{N}} \equiv (1/\delta m_{N-\bar{N}})$ qualitatively, first we note that the operator \mathcal{O}_1 leads to the Majorana mass for the right handed neutrino $M_R = f v_R^2/M_{\text{Pl}}$. The seesaw formula then leads to the relation $m_\nu = M_{\text{Pl}} m_{\nu,D}^2 / (f v_R^2)$. On the other hand, the operator \mathcal{O}_2 leads to a $\Delta B = 1$ operator with strength v_R/M_{Pl} . Leaving aside the details of the flavor structure of \mathcal{O}_2 and how actually $\delta m_{N-\bar{N}}$ arises, it is clear that we have a simple linear relation between

the neutrino masses and the $N-\bar{N}$ oscillation time:

$$m_\nu = C \frac{\tau_{N-\bar{N}}}{M_{\text{Pl}}}, \tag{3}$$

where C is a dimensional constant which depends only on the details of weak scale physics and does not involve the high scale v_R . We will evaluate C in the next section. This simple relation makes it clear that our present knowledge of the neutrino masses allows a direct prediction of the $N-\bar{N}$ oscillation time in the context of the supersymmetric left–right models broken by doublet Higgs fields.

3. From supersymmetric $\Delta B = 1$ operator to $N-\bar{N}$ oscillations

Let us now proceed to examine the expected $N-\bar{N}$ oscillation time resulting from the $\Delta B = 1$ operator \mathcal{O}_2 of Eq. (2). An important point to note here is that since \mathcal{O}_2 is a superpotential term with antisymmetric color contraction it must have antisymmetric flavor contraction as well. The flavor structure of this operator is then of the type $u^c d^c s^c$, $u^c d^c b^c$ or $u^c s^c b^c$ in terms of the superfields. We must then use flavor mixings to obtain the fermionic operator of the type $u^c d^c d^c$ and then the six quark $N-\bar{N}$ operator. The dominant contribution to this process comes from the Feynman diagram shown in Fig. 1 which proceeds through the exchange of a gluino and squarks [11] and involves two $\tilde{d}^c-\tilde{b}^c$ mixings. The strength of the $\Delta B = 2$ operator resulting from Fig. 1 can be estimated to be

$$G_{\Delta B=2} \simeq \frac{2g_3^2 [(\delta_{RR}^{13})]^2 f'^2}{M_{\tilde{g}} m_{\tilde{q}}^4}, \tag{4}$$

where $M_{\tilde{g}}$ is the gluino mass, $m_{\tilde{q}}$ is the squark mass and (δ_{RR}^{13}) is the $\tilde{d}^c-\tilde{b}^c$ mixing angle. The effective baryon number violating $u^c d^c b^c$ Yukawa coupling in the superpotential is parametrized here as $f'(v_R/M_{\text{Pl}})$ (see Eq. (2)).

Let us first discuss the origin of the flavor mixing that changes a $u^c d^c b^c$ operator to the required $u^c d^c d^c$ operator. The dominant source for this in the present context turns out to be the mixing of \tilde{b}^c with \tilde{d}^c . Such mixings occur in the left–right supersymmetric model since the right-handed quark mixings are physical above the scale v_R . The renormalization group

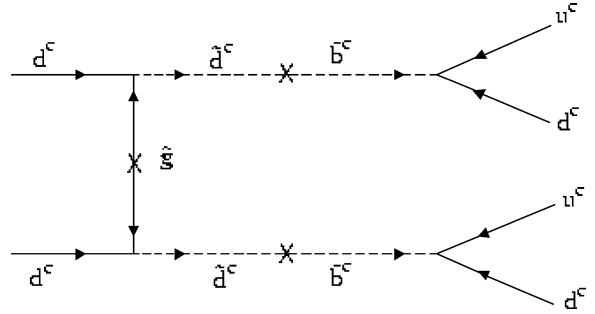


Fig. 1. Tree level gluino–squark diagram for $N-\bar{N}$ oscillations.

evolution of the soft SUSY breaking mass parameters between M_{Pl} and v_R will then induce mixings in the right-handed down squark sector proportional to the top-quark Yukawa coupling and the right-handed CKM mixings. This is analogous to the RGE evolution in the MSSM inducing squark mixing in the left-handed down squark sector proportional to the left-handed CKM angles and the top-quark Yukawa coupling. We estimate this right-handed $\tilde{d}^c-\tilde{b}^c$ mixing to be

$$\begin{aligned} (\delta_{RR}^{13}) &\simeq \frac{\lambda_t^2 (3m_0^2 + A_0^2)}{8\pi^2 (m_0^2 + 8M_{1/2}^2)} (V_{td}^* V_{tb}) \ln(M_{\text{Pl}}/v_R) \\ &\simeq 2 \times 10^{-4}. \end{aligned} \tag{5}$$

This estimate is obtained by integrating out the RGE between M_{Pl} and v_R assuming universality of masses at M_{Pl} [12]. Since above v_R , both t^c and b^c are part of the same $SU(2)_R$ multiplet, unlike in the MSSM, b^c Yukawa coupling is of order one. In this momentum range, the top-quark Yukawa coupling reduces the mass of \tilde{b}^c . In going to the physical basis of the quarks, this effect will induce the squark mixing quoted in Eq. (5). For the numerical estimate we took $m_0 = M_{1/2}$ and $A_0 = 0$ and $v_R \sim 2 \times 10^{16}$ GeV for illustration.

There is a second source of flavor violation that induces $\tilde{d}^c-\tilde{b}^c$ mixing in general SUSY models. That is the baryon number violating Yukawa couplings themselves. If we write in standard notation, the effective B -violating superpotential arising from Eq. (2) as $W \supset (\lambda''_{ijk}/2) u_i^c d_j^c d_k^c$, the RGE evolution from Planck scale to the weak scale will induce $\tilde{d}^c-\tilde{b}^c$ mixing proportional to λ''_{ijk} . For example, if we keep only the couplings involving the u^c quark, viz., $\lambda''_{123}, \lambda''_{113}$ and

λ''_{112} , we can estimate the induced (δ_{RR}^{13}) by integrating the relevant RGE [13] to be

$$(\delta_{RR}^{13}) \simeq \frac{\lambda''_{121}\lambda''_{123}}{4\pi^2} \frac{(3m_0^2 + A_0^2)}{(m_0^2 + 8M_{1/2}^2)} \ln(M_{\text{Pl}}/M_Z). \quad (6)$$

Recalling that $\lambda'' \sim v_R/M_{\text{Pl}}$, we see that while this source of flavor mixing may not be negligible, it would be typically smaller than the ones from the right-handed quark mixings of Eq. (5).

A third source of flavor violation relevant for $N-\bar{N}$ oscillations has been identified in Ref. [14] involving the exchange of the Wino. Such diagrams will have an electroweak loop suppression and a chirality suppression necessary to convert the left-handed squark to the right-handed one. We find that this contribution to $\delta m_{N-\bar{N}}$ has a suppression factor given approximately by $[(\alpha_2/4\pi)(m_b/m_{\tilde{q}})]^2 \sim 1 \times 10^{-9}$ (valid for small $\tan\beta$) which is about two orders of magnitude smaller in this class of models compared to the gluino exchange diagram of Fig. 1.

One has to calculate the hadronic matrix element of the six quark operator in order to obtain the $\tau_{N-\bar{N}}$. This has been discussed in several places in the literature [15]. The calculations of this “conversion” factor can be done using crude physical arguments, according to which one has to multiply the $G_{\Delta B=2}$ by $|\psi(0)|^4$ to obtain $\delta m_{N-\bar{N}}$ where ψ is the baryonic wave function for three quarks inside a nucleon. On dimensional grounds, one can deduce that $|\psi(0)|^4 \simeq \Lambda_{\text{QCD}}^6$, which implies that $\delta m_{N-\bar{N}} \sim 10^{-5} G_{\Delta B=2}$ GeV. More detailed bag model calculations have been carried out. Rao and Shrock in Ref. [15] quote this conversion factor to be $2.5 \times 10^{-5} G_{\Delta B=2}$. We shall use this number for our numerical illustrations.

Combining this matrix element with Eqs. (4), (5) we obtain

$$\tau_{N-\bar{N}} \simeq \frac{7 \times 10^8 \text{ s}}{f'^2} \left(\frac{2 \times 10^{14} \text{ GeV}}{v_R} \right)^2 \times \left(\frac{M_{\tilde{g}}}{500 \text{ GeV}} \right) \left(\frac{m_{\tilde{q}}}{500 \text{ GeV}} \right)^4. \quad (7)$$

We can rewrite Eq. (7) in a form that makes the connection with the neutrino mass more transparent. The mass of ν_τ can be expressed through the seesaw formula from Eq. (2) as $m_{\nu_\tau} = (m_{\nu_\tau^D})^2 M_{\text{Pl}}/(f v_R^2)$

where $m_{\nu_\tau^D}$ denotes the Dirac mass of ν_τ . Eliminating the high scale v_R from this, we have from Eq. (7),

$$\tau_{N-\bar{N}} \simeq 2.8 \times 10^4 \text{ s} \left(\frac{f}{f'^2} \right) \left(\frac{m_{\nu_\tau}}{0.06 \text{ eV}} \right) \left(\frac{m_t}{m_{\nu_\tau^D}} \right)^2 \times \left(\frac{M_{\tilde{g}}}{500 \text{ GeV}} \right) \left(\frac{m_{\tilde{q}}}{500 \text{ GeV}} \right)^4. \quad (8)$$

Since the value of m_{ν_τ} can be determined from the atmospheric neutrino data under certain assumptions, we conclude that within the seesaw framework, measurement of $N-\bar{N}$ oscillation will be a measure of the Dirac mass of the tau neutrino. This can then be used as a way to discriminate between models of neutrino masses.

To see the specific prediction for $N-\bar{N}$ oscillations within the context of the class of models under consideration, we need to know the ν_τ Dirac mass. We can estimate it from the following relations for the Dirac masses of the third generation quarks and leptons in the SUSY left–right model:

$$\begin{aligned} m_t &= (h_{t,1} \cos \alpha_u + h_{t,2} \sin \alpha_u) v_u, \\ m_b &= (h_{t,1} \cos \alpha_d + h_{t,2} \sin \alpha_d) v_d, \\ m_{\nu_\tau^D} &= (h_{\tau,1} \cos \alpha_u + h_{\tau,2} \sin \alpha_u) v_u, \\ m_\tau &= (h_{\tau,1} \cos \alpha_d + h_{\tau,2} \sin \alpha_d) v_d. \end{aligned} \quad (9)$$

Here $\alpha_{u,d}$ are the Higgs mixing parameters obtained from Eq. (1) (Eg: $\tan \alpha_u = \lambda_1/\lambda_2$) and $v_{u,d}$ are the VEVs of the MSSM doublets. From Eq. (9) it follows that in the limit of $h_{t,1} \gg h_{t,2}$ and $h_{\tau,1} \gg h_{\tau,2}$, we get $m_{\nu_\tau^D} \simeq m_\tau \frac{m_t}{m_b}$. Since at such high scales $m_b \simeq m_\tau$, this predicts $m_{\nu_\tau^D} \simeq m_t$. In fact, we find that unless the two terms in Eq. (9) for $m_{\nu_\tau^D}$ are precisely canceled, the Dirac mass of ν_τ will be approximately equal to m_t .

Using $m_{\nu_\tau^D} \simeq m_t$ and $f \sim 1$, $f' \sim 10^{-1}$, we get a value for the $N-\bar{N}$ oscillation time which is tantalizingly close to the present experimental lower limit [8]. For values of $m_{\nu_\tau^D}$ 10 times smaller than m_t , and taking the supersymmetric particle masses as large as 1 TeV, we see that $\tau_{N-\bar{N}}$ is less than 9×10^9 seconds,³ which is in the range accessible to a recently

³ We have not included the QCD evolution factor from the SUSY scale of few hundred GeV to the GeV scale. Based on the QCD factor of 1.33 for the three quark proton decay operator [16], we estimate that the corresponding factor for the $N-\bar{N}$ case should be about 2, which would reduce the estimate of $\tau_{N-\bar{N}}$ by a factor of 2.

proposed experiment [17]. It would thus appear that a search for neutron–antineutron oscillation will provide an enormously useful window into neutrino mass models and as such a powerful constraint on the nature of new physics beyond the Standard Model.

The prediction for $N-\bar{N}$ oscillations can be sharpened if we make use of flavor symmetries to determine the coefficients f and f' in Eq. (8). We illustrate this with a specific choice of flavor symmetry [18] taken to be $SU(2)_H \times U(1)_H$. The first two families of fermions form doublets of $SU(2)_H$ and have a $U(1)_H$ charge of +1 while the third family fermions are singlets under both groups. This flavor symmetry is broken by a pair of doublets $\phi(1) + \bar{\phi}(-1)$ and singlets $\chi(1) + \bar{\chi}(-1)$. Allowing for effective operators suppressed by a scale M larger than the VEVs of these fields provides a natural explanation of the fermion mass and mixing angle hierarchy. If we choose $\langle\phi\rangle = \langle\bar{\phi}\rangle = \epsilon_\phi M$ and $\langle\chi\rangle = \langle\bar{\chi}\rangle = \epsilon_\chi M$, a reasonable fit to all quark and lepton masses is obtained, including neutrinos, for $\epsilon_\phi \simeq 1/7$ and $\epsilon_\chi \simeq 1/20$ and all dimensionless couplings being order one [18]. In this model, we can estimate the couplings f and f' from the horizontal quantum numbers. They are $f \sim \epsilon_\phi^2$ and $f' \simeq \epsilon_\chi^2 \sin\theta_C$, so that⁴ $f/f'^2 \simeq 8 \times 10^4$. This estimate leads to $\tau_{N-\bar{N}} \simeq 2 \times 10^9$ s from Eq. (8). Allowing for uncertainties of order 1 in this estimate, we expect that $\tau_{N-\bar{N}}$ not to exceed about 10^{10} s.

Before we conclude a few comments are in order:

1. The model becomes unacceptable as soon as $SU(3)_C \times U(1)_{B-L}$ is embedded into a higher symmetry such as $SU(4)_C$ or $SO(10)$ group because in that case, the $\Delta B = 1$ operator described in Eq. (2) is accompanied by other R -parity violating operators coming from the same higher dimensional operator $\mathcal{O}_{1,2}$ due to the higher symmetry. Together, they would lead to unacceptable proton decay rate. Thus observation of $N-\bar{N}$ oscillation would be a signal of an explicit

$SU(3)_C \times U(1)_{B-L}$ symmetry all the way upto the Planck (or string) scale.

2. Baryogenesis has to proceed through a weak scale scenario since the $\Delta B = 1$ interactions in the model are in equilibrium down to the TeV scale and will wash out any primordial baryon or lepton asymmetry. We note that the baryon number violating interactions contained in \mathcal{O}_2 themselves can potentially be the source of weak scale baryogenesis [19].
3. The lightest neutralino in this model is unstable and will decay via $\chi^0 \rightarrow qq\bar{q}$ modes due to the presence of the effective $\Delta B = 1$ operator \mathcal{O}_2 . This prediction is directly testable at colliders. An alternative candidate for dark matter must be sought.

4. Conclusions

In conclusion, we have found that in a large class of seesaw models for neutrino masses, despite the high scale of seesaw dictated by the current neutrino oscillation data, neutron–antineutron oscillation is in the observable range. In fact, unless the Dirac masses of neutrinos are far below those deduced under simple and reasonable assumptions, we predict an upper bound on the neutron–antineutron oscillation time in the range of 10^9 – 10^{10} s. This is very close to the present experimental lower limit on $N-\bar{N}$ oscillations. In the most conservative theoretical scenario, the measurement of $N-\bar{N}$ oscillation time would be a measure of the Dirac mass for the tau neutrino, given the values of squark masses. This in itself would be an extremely interesting result, since it would discriminate among theoretical models of neutrino masses. This is apart from the fundamental importance that any observation of baryon number violation will carry. We, therefore, strongly urge a new experimental search for neutron–antineutron oscillation.

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⁴ Normally, the parameter f could have been of order one but in the horizontal model or Ref. [18], due to large $\nu_\mu-\nu_\tau$ mixing, it is the ν_μ flavor entry that dominates the atmospheric neutrino mass difference and hence the horizontal suppression factor ϵ_ϕ^2 . The ϵ_χ^2 factor is due the fact that the operator must be invariant under $U(1)_H$.

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