PRAMANA — journal of physics

Vol. 67, No. 5
November 2006
pp. 783–791

Observable N–$\bar{N}$ oscillation, high-scale see-saw and origin of matter

R N MOHAPATRA
Department of Physics and Center for String and Particle Theory, University of Maryland, College Park, MD 20742, USA
E-mail: rmohapat@physics.umd.edu

Abstract. See-saw mechanism has been a dominant paradigm in the discussion of neutrino masses. We discuss how this idea can be tested via a baryon number violating process such as N–$\bar{N}$ oscillation. Since the expected see-saw scale is high and the N–$\bar{N}$ amplitude goes like M_{R}^{-5}, one might think that this process is not observable in realistic see-saw models for neutrino masses. In this talk I show that in supersymmetric models, the above conclusion is circumvented leading to an enhanced and observable rate for N–$\bar{N}$ oscillation. I also discuss a new mechanism for baryogenesis in generic models for neutron–anti-neutron oscillation .

Keywords. Neutron–anti-neutron oscillation; baryogenesis.

PACS No. 11.30.Fs

1. Introduction

There are various reasons to suspect that baryon number is not a good symmetry of nature. They are: (i) nonperturbative effects of the Standard Model lead to $\Delta B \neq 0$, while keeping $\Delta(B-L) = 0$ [1]; (ii) an understanding of the origin of matter in the Universe requires $\Delta B \neq 0$ interactions [2] and (iii) many theories beyond the Standard Model lead to interactions that violate baryon number [3,4].

If indeed such interactions are there, an important question is: can we observe them in experiments? Two interesting baryon nonconserving processes of experimental interest are: (a) proton decay, e.g. $p \to e^+ + \pi^0, \bar{\nu} + K^0$ etc. [4,5] and (b) $N \leftrightarrow \bar{N}$ oscillation [6–8]. These two classes of processes probe two different selection rules for baryon nonconservation: $\Delta(B-L) = 0$ for proton decay and $\Delta(B-L) = 2$ for $N \leftrightarrow \bar{N}$ oscillation. They are signatures of two totally different directions for unification beyond the Standard Model. For example, observation of proton decay will point strongly towards a grand desert till about the scale of 10^{16} GeV whereas $N \leftrightarrow \bar{N}$ oscillation will require new physics at an intermediate scale at or above the TeV scale but much below the GUT scale. Further experimental search for both these processes can therefore provide key insight into the nature of unification beyond the Standard Model with or without supersymmetry.

While proton decay goes very naturally with the idea of eventual grand unification of forces and matter, recent discoveries of neutrino oscillations have made $N \leftrightarrow \bar{N}$ oscillation to be quite plausible theoretically if small neutrino masses are to be understood as a consequence of the see-saw mechanism [9]. This can be seen as follows: see-saw mechanism implies Majorana neutrinos implying the existence of $\Delta(B-L)=2$ interactions. In the domain of baryons, it implies the existence of $N \leftrightarrow \bar{N}$ oscillation as noted many years ago [8]. In fact an explicit model for $N \leftrightarrow \bar{N}$ oscillation was constructed in [8] by implementing the see-saw mechanism within the framework of the Pati–Salam [3] $SU(2)_{\rm L} \times SU(2)_{\rm R} \times SU(4)_{\rm c}$ model, where quarks and leptons are unified. It was shown that this process is mediated by the exchange of diquark Higgs bosons giving an amplitude (see figure 1 of [8]) ($G_{N\leftrightarrow\bar{N}}$) which scales like M_{qq}^{-5}. In the nonsupersymmetric version without fine-tuning, one expects $M_{qq} \propto v_{BL}$ leading to $G_{N\leftrightarrow\bar{N}} \simeq v_{BL}^{-5}$. So only if $M_{qq} \sim v_{BL} \sim$ 10–100 TeV, the $\tau_{N\leftrightarrow\bar{N}}$ is in the range of 10^6–10^8 s and is accessible to experiments. On the other hand, in generic see-saw models for neutrinos, one expects $v_{BL} \sim 10^{11}$–10^{14} GeV depending on the range of the third-generation Dirac mass for the neutrino of 1–100 GeV. An important question therefore is whether in realistic see-saw models, $N \leftrightarrow \bar{N}$ oscillation is at all observable. Another objection to the above nonsupersymmetric model for $N \leftrightarrow \bar{N}$ that was raised in the eighties was that such interactions will erase any baryon asymmetry created at high scales. It is therefore important to overcome this objection.

Several years ago, a high-scale see-saw model with observable N–$\bar{N}$ oscillation was presented using R-parity violating interactions [10]. Such models in general lead to difficulties in understanding the origin of matter and also do not have a naturally stable supersymmetric dark matter.

In this talk, I first discuss a recent paper [11] where it is shown that in a class of supersymmetric $SU(2)_{\rm L} \times SU(2)_{\rm R} \times SU(4)_{\rm c}$ models (called SUSY G_{224}), an interesting combination of circumstances improves the v_{BL} dependence of $G_{\Delta B=2}$ to $v_{BL}^{-2} v_{wk}^3$ instead of v_{BL}^{-5} making $N \leftrightarrow \bar{N}$ oscillation observable. This does not require the existence of R-parity violation and in fact in these models R-parity is naturally conserved giving rise to a stable dark matter. I then discuss a new mechanism for post-sphaleron baryogenesis where an observable N–$\bar{N}$ oscillation is necessary to generate the required baryon-to-photon ratio of the Universe.

The basic ingredients of such a theory was presented in [12] where it was shown that in the minimal supersymmetric $SU(2)_{\rm L} \times SU(2)_{\rm R} \times SU(4)_{\rm c}$ model, there exist accidental symmetries that imply that some of the M_{qq}'s which mediate N–$\bar{N}$ oscillation are in the TeV range even though $v_{BL} \simeq 10^{11}$–10^{12} GeV. The present work [11] points out that there exist a new class of Feynman diagrams which enhance the $N \leftrightarrow \bar{N}$ oscillation amplitude in generic supersymmetric models of this type making it observable. We then discuss the new mechanism for baryogenesis in models where the baryon number violation is mediated by a higher-dimensional operator such as in the case of N–$\bar{N}$ oscillation [13].

At present, the best lower bound on $\tau_{N\leftrightarrow\bar{N}}$ comes from ILL reactor experiment [14] and is 10^8 s. There are also comparable bounds from nucleon decay search experiments [15]. There are proposals to improve the precision of this search by at least two orders of magnitude [16]. We feel that the results of this paper [11,13] should give new impetus to a search for neutron–anti-neutron oscillation.

2. $SU(2)_{\rm L} \times SU(2)_{\rm R} \times SU(4)_{\rm c}$ model with light diquarks

The quarks and leptons in this model are unified and transform as ψ : $(\mathbf{2},\mathbf{1},\mathbf{4})\oplus\psi^c$: $(\mathbf{1},\mathbf{2},\bar{\mathbf{4}})$ representations of $SU(2)_{\rm L} \times SU(2)_{\rm R} \times SU(4)_{\rm c}$. For the Higgs sector, we choose, ϕ_1 : $(\mathbf{2},\mathbf{2},\mathbf{1})$ and ϕ_{15} : $(\mathbf{2},\mathbf{2},\mathbf{15})$ to give mass to the fermions. The Δ^c : $(\mathbf{1},\mathbf{3},\mathbf{10})\oplus\bar{\Delta}^c$: $(\mathbf{1},\mathbf{3},\overline{\mathbf{10}})$ to break the $B-L$ symmetry. The diquarks mentioned above which lead to $\Delta(B-L)=2$ processes are contained in the Δ^c : $(\mathbf{1},\mathbf{3},\mathbf{10})$ multiplet. We also add a $B-L$ neutral triplet Ω : $(\mathbf{1},\mathbf{3},\mathbf{1})$ which helps to reduce the number of light diquark states. The superpotential of this model is given by

$$W = W_Y + W_{H1} + W_{H2} + W_{H3}, \tag{1}$$

where

$$W_{H1} = \lambda_1 S(\Delta^c\bar{\Delta}^c - M^2) \ + \ \lambda_A\frac{(\Delta^c\bar{\Delta}^c)^2}{M_{\rm Pl}}, \tag{2}$$

$$W_{H2} = \lambda_B\frac{(\Delta^c\Delta^c)(\bar{\Delta}^c\bar{\Delta}^c)}{M_{\rm Pl}} + \lambda_C\Delta^c\bar{\Delta}^c\Omega + \mu_i{\rm Tr}\,(\phi_i\phi_i), \tag{3}$$

$$W_{H3} = \lambda_D\frac{{\rm Tr}\,(\phi_1\Delta^c\bar{\Delta}^c\phi_{15})}{M_{\rm Pl}}, \tag{4}$$

$$W_Y = h_1\psi\phi_1\psi^c + h_{15}\psi\phi_{15}\psi^c + f\psi^c\Delta^c\psi^c. \tag{5}$$

Note that since we do not have parity symmetry in the model, the Yukawa couplings h_1 and h_{15} are not symmetric matrices. When $\lambda_B = 0$, this superpotential has an accidental global symmetry much larger than the gauge group [12]; as a result, vacuum breaking of the $B-L$ symmetry leads to the existence of light diquark states that mediate $N \leftrightarrow \bar{N}$ oscillation and enhance the amplitude. In fact it was shown that for $\langle\Delta^c\rangle \sim \langle\bar{\Delta}^c\rangle \neq 0$, and $\langle\Omega\rangle \neq 0$ and all VEVs in the range of 10^{11}–10^{12} GeV, the light states are those with quantum numbers: $\Delta_{u^cu^c}$. The symmetry argument behind is that [12] for $\lambda_B = 0$, the above superpotential is invariant under $U(10,c)\times SU(2,c)$ symmetry which breaks down to $U(9,c)\times U(1)$ when $\langle\Delta^c_{\nu^c\nu^c}\rangle = v_{BL} \neq 0$. This results in 21 complex massless states; on the other hand these VEVS also break the gauge symmetry down from $SU(2)_{\rm R}\times SU(4)_{\rm c}$ to $SU(3)_{\rm c}\times U(1)_Y$. This allows nine of the above states to pick up masses of order gv_{BL} leaving 12 massless complex states which are the six $\Delta^c_{u^cu^c}$ plus six $\bar{\Delta}^c_{u^cu^c}$ states. Once $\lambda_B \neq 0$ and is of the order 10^{-2}–10^{-3}, they pick up mass (call $M_{u^cu^c}$) of the order of the electroweak scale.

3. $N \leftrightarrow \bar{N}$ oscillation – A new diagram

To discuss $N \leftrightarrow \bar{N}$ oscillation, we introduce a new term in the superpotential of the form [8]:

$$W_{\Delta B=2} \ = \ \frac{1}{M_*}\epsilon^{\mu'\nu'\lambda'\sigma'}\epsilon^{\mu\nu\lambda\sigma}\Delta^c_{\mu\mu'}\Delta^c_{\nu\nu'}\Delta^c_{\lambda\lambda'}\Delta^c_{\sigma\sigma'}\,, \tag{6}$$

where the μ,ν etc. stand for $SU(4)_{\rm c}$ indices (we have suppressed the $SU(2)_{\rm R}$ indices). *A priori* M_* could be of the order $M_{\rm Pl}$. However, the terms in eq. (2)

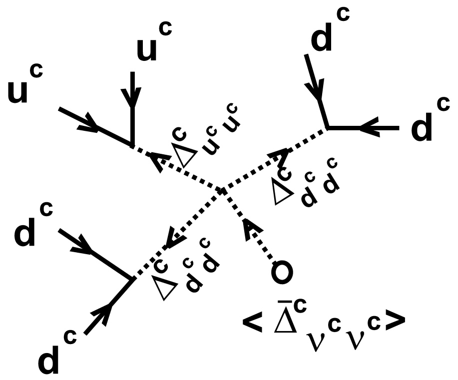

Figure 1. The Feynman diagram responsible for N–$\bar{N}$ oscillation as discussed in ref. [8].

are different from those in eq. (4). So they could arise from a different high-scale theory. The mass M_* is therefore a free parameter that we choose to be much less than $M_{\rm Pl}$. This term does not affect the masses of the Higgs fields. When $\Delta^c_{\nu^c\nu^c}$ acquires a VEV, $\Delta B = 2$ interactions are induced from this superpotential, and $N \leftrightarrow \bar{N}$ oscillations are generated by two diagrams given in figures 1 and 2. The first diagram (figure 1) in which only diquark Higgs fields are involved was already discussed in [8] and goes like $G_{N\leftrightarrow\bar{N}} \simeq \frac{f_{11}^3 v_{BL} M_\Delta}{M^2_{u^cu^c} M^4_{d^cd^c} M_*}$. Taking $M_{u^cu^c} \sim 350$ GeV, $M_{d^cd^c} \sim \lambda' v_{BL}$ and $M_\Delta \sim v_{BL}$ as in the argument [12], we see that this diagram scales like $v_{BL}^{-3} v_{wk}^{-2}$.

In ref. [11] a new diagram (figure 2) was pointed out which owes its origin to supersymmetry. We get for its contribution to $G_{\Delta B=2}$:

$$G_{N\leftrightarrow\bar{N}} \simeq \frac{g_3^2}{16\pi^2} \frac{f_{11}^3 v_{BL}}{M^2_{u^cu^c} M^2_{d^cd^c} M_{\rm SUSY} M_*}. \tag{7}$$

Using the same arguments as above, we find that this diagram scales like $v_{BL}^{-2} v_{wk}^{-3}$ which is therefore a significant enhancement over figure 1.

In order to estimate the rate for $N \leftrightarrow \bar{N}$ oscillation, we need not only the different mass values for which we now have an order of magnitude, but we also need the Yukawa coupling f_{11}. Now f_{11} is a small number since its value is associated with the lightest right-handed neutrino mass. However, in the calculation we need its value on the basis where quark masses are diagonal. We note that the N–$\bar{N}$ diagrams involve only the right-handed quarks, the rotation matrix need not be the CKM matrix. The right-handed rotations need to be large, e.g. in order to involve f_{33} (which is $O(1)$), we need $(V_{\rm R}^{(u,d)})_{31}$ to be large, where $V_{\rm L}^{(u,d)\dagger} Y_{u,d} V_{\rm R}^{(u,d)} = Y_{u,d}^{\rm diag.}$. The left-handed rotation matrices $V_{\rm L}^{(u,d)}$ contribute to the CKM matrix, but right-handed rotation matrices $V_{\rm R}^{(u,d)}$ are unphysical in the Standard Model. In this model, however, we get to see its contribution since we have a left–right gauge symmetry.

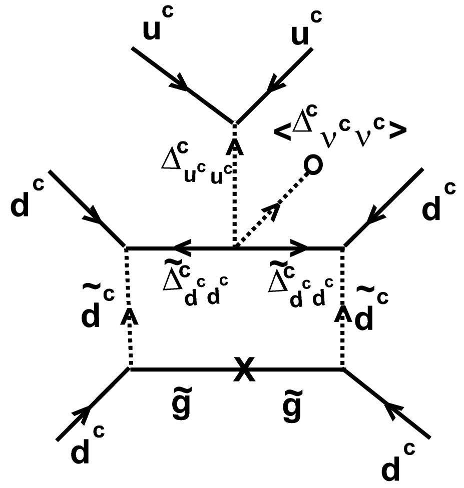

Figure 2. The new Feynman diagram for N–$\bar{N}$ oscillation.

Let us now estimate the time of oscillation. When we start with a f-diagonal basis (call the diagonal matrix $\hat{f}$), the Majorana coupling f_{11} in the diagonal basis of up- and down-type quark matrices can be written as $(V_{\mathrm{R}}^T \hat{f} V_{\mathrm{R}})_{11} \sim (V_{31}^R)^2 \hat{f}_{33}$. Now $\hat{f}_{33}$ is $O(1)$ and V_{31}^R can be $\sim$0.6, so f_{11} is about 0.4 in the diagonal basis of the quark matrices. We use M_{SUSY}, $M_{u^c u^c} \sim 350$ GeV and $v_{BL} \sim 10^{12}$ GeV. The mass of $\tilde{\Delta}_{d^c d^c}$, i.e. $M_{d^c d^c}$ is 10^9 GeV which is obtained from the VEV of Ω : $(\mathbf{1}, \mathbf{3}, \mathbf{1})$. We choose $M_* \sim 10^{13}$ GeV. Putting all the above numbers together, we get

$$G_{N \leftrightarrow \bar{N}} \simeq 1 \times 10^{-30}\ \mathrm{GeV}^{-5}. \tag{8}$$

Along with the hadronic matrix element [17], the N–$\bar{N}$ oscillation time is found to be about 2.5×10^{10} s which is within the reach of possible next generation measurements. If we choose, $M_* \simeq M_{\mathrm{Pl}}$, we will get for $\tau_{N-\bar{N}} \sim 10^{15}$ s unless we choose $M_{d^c d^c}$ to be lower (say 10^7 GeV). This is a considerable enhancement over the nonsupersymmetric model of [8] with see-saw scale of 10^{12} GeV.

We also note that as noted in [8] the model is invariant under the hidden discrete symmetry under which a field $X \to \mathrm{e}^{i\pi B_X} X$, where B_X is the baryon number of the field X. As a result, proton is absolutely stable in the model. Furthermore, since R-parity is an automatic symmetry of MSSM, this model has a naturally stable dark matter.

4. Baryogenesis and N–$\bar{N}$ oscillation

In the early 1980s when the idea of neutron–anti-neutron oscillation was first proposed in the context of unified gauge theories, it was thought that the high dimensionality of the $\Delta B \neq 0$ operator would pose major difficulty in understanding the origin of matter. The main reason for this is that the higher-dimensional operators remain in thermal equilibrium until late in the evolution of the Universe since the thermal decoupling temperature T_* for such interactions goes roughly like $v_{BL}(v_{BL}/M_{\rm Pl})^{1/9}$ which is in the range of temperatures where $B+L$ violating sphaleron transitions are in full thermal equilibrium. They will therefore erase any baryon asymmetry generated in the very early moments of the Universe (say close to the GUT time of 10^{-30} s or so) in the prevalent baryogenesis models. In models with observable N–$\bar{N}$ oscillation therefore, one has to search for new mechanisms for generating baryons below the weak scale. In this section, we discuss such a possibility [13] discussed in a recent unpublished work of K S Babu and S Nasri.

As an illustration of the way the new mechanism operates, let us assume that there is a scalar field that couples to the $\Delta B = 2$ operator, i.e. $L_{\rm I} = Su^c d^c d^c u^c d^c d^c/M^6$, where the scalar boson has mass of order of the weak scale. This leads to baryon number violation if $\langle S\rangle \neq 0$ and observable N–$\bar{N}$ transition if M is in the few hundred to 1000 GeV range. The direct decay of S in these models can lead to an adequate mechanism for baryogenesis.

To discuss how this comes about, let us first note that the high dimension of $L_{\rm I}$ allows the scalar $\Delta B \neq 0$ decay to go out of equilibrium at weak scale temperatures. This clearly satisfies the out-of-equilibrium condition given by Sakharov conditions for the origin of matter [2].

To outline the rest of the details of this mechanism [13], we consider an effective sub-TeV scale model that gives rise to the higher-dimensional operator for $N \leftrightarrow \bar{N}$ oscillation. It consists of the following color sextet, $SU(2)_{\rm L}$ singlet scalar bosons (X, Y, Z) with hypercharge $-\frac{4}{3}, +\frac{8}{3}, +\frac{2}{3}$ respectively that couple to quarks. We add to it a scalar field with mass in the 100 GeV range. One can now write down the following standard model invariant interaction Lagrangian:

$$\begin{aligned}\mathcal{L}_{\rm I} = {} & h_{ij} X d^c_i d^c_j + f_{ij} Y u^c_i u^c_j \\ & + g_{ij} Z(u^c_i d^c_j + u^c_j d^c_j) + \lambda_1 S X^2 Y + \lambda_2 S X Z^2. \end{aligned} \tag{9}$$

The scalar field S is assumed to have $B = 2$. To see the constraints on the parameters of the theory, we note that the present limits on $\tau_{N-\bar{N}} \geq 10^8$ s implies that the strength $G_{N-\bar{N}}$ of the $\Delta B = 2$ transition is $\leq 10^{-28}$ GeV^{-5}. From figure 3, we conclude that

$$G_{N-\bar{N}} \simeq \frac{\lambda_1 M_1 h^2_{11} f_{11}}{M^2_Y M^4_X} + \frac{\lambda_2 M_1 h_{11} g^2_{11}}{M^2_X M^2_Z} \leq 10^{-28}\ \text{GeV}^{-5}. \tag{10}$$

For $\lambda_{1,2} \sim h \sim f \sim g \sim 10^{-3}$, we have $M_1 \sim M_{X,Y,Z} \simeq 1$ TeV. In our discussion, we will stay close to this range of parameters and see how one can understand the baryon asymmetry of the Universe. The singlet field will play a key role in the generation of baryon asymmetry. We assume that $\langle S\rangle \sim M_X$ but $M_{S_{\rm r}} \sim 100$ GeV, where $S_{\rm r}$ is the real part of the S field after its VEV is subtracted. It can then decay into final states with $B = \pm 2$.

On the way to calculating the baryon asymmetry, let us first discuss the out of equilibrium condition. As the temperature of the Universe falls below the masses of the X, Y, Z particles, the annihilation processes $X\bar{X} \to d^c \bar{d}^c$ (and analogous processes for Y and Z) remain in equilibrium. As a result, the number density of X, Y, Z particles gets depleted and only the S particle survives along with the usual Standard Model particles. One of the primary generic decay modes of S is $S \to u^c d^c d^c u^c d^c d^c$. There could be other decay modes which depend on the details of the model. Those can be made negligible by a choice of parameters which do not enter our discussion of N–$\bar{N}$ and baryogenesis. For $T \geq M_S$, its decay rate is given by $\Gamma_S \sim (T^{13}/16\pi^9 M_X^{12})$ where we have set the masses of X, Y and Z particles to be of the same order for simplicity. This decay goes out of equilibrium around $T_* \simeq M_X \left(\frac{160\pi^9 M_X}{M_{\rm Pl}}\right)^{1/11}$. Here we have assumed that the coupling of the X, Y, Z particles to second- and third-generation quarks are of order 0.1–1. This gives $T_* \sim 0.1$–$0.2 M_X$ or in the sub-TeV range. Below this temperature the decay rate of S falls very rapidly as the temperature cools. However, as soon as $T \leq M_S$, the decay rate becomes a constant whereas the expansion rate of the Universe is slowing down. So at a temperature T_d, S will start to decay at

$$T_d \simeq \left(\frac{M_{\rm Pl} M_S^{13}}{(2\pi)^9 M_X^{12}}\right)^{1/2}. \tag{11}$$

The corresponding epoch must be above that of Big Bang nucleosynthesis. This puts a constraint on the parameters of the model. For instance, for $M_S \sim 200$ GeV and $M_X \sim 3$ TeV, we get $T_d \sim$ GeV.

It is well-known that baryon asymmetry can arise only via the interference of a tree diagram with a one-loop diagram. The tree diagram is clearly the one where $S \to 6q$. There are however two classes of loop diagrams that can contribute: one where the loop involves the same fields X, Y and Z. A second one involves W-exchange, which involves only Standard Model physics at this scale (figure 3). We find that the second contribution can actually dominate. It also has the advantage that it involves less number of arbitrary parameters. The baryon asymmetry is defined as follows

$$\epsilon_B \simeq \frac{n_S}{n_\gamma} \frac{\Gamma(S \to 6q) - \Gamma(S \to 6\bar{q})}{\Gamma(S)}. \tag{12}$$

We find that

$$\epsilon_B \simeq \begin{cases} 2\alpha_2 \, {\rm Im}(V_{tb} V_{cb}^* h_{33} h_{23}^*) \frac{m_c m_b^2 m_t}{M_W^2 M_S^2}; & M_S < m_t \\ 2\alpha_2 \frac{m_s m_b m_t^2}{M_W^2 M_S^2} {\rm Im}[(h_{33} h_{32}^*)(V_{tb}^* V_{cb})]; & M_S > m_t \end{cases}. \tag{13}$$

Note that the trace in the above equation has an imaginary part and therefore leads to nonzero asymmetry. The magnitude of the asymmetry depends on T_d/M_S as well as the detailed profile of the various coupling matrices h, g, f and we can easily get the desired value of the baryon asymmetry by appropriately choosing them. We have checked that there is no conflict between the desired magnitude of baryon asymmetry and the present lower bound on the N–$\bar{N}$ transition time of 10^8 s.

An interesting point worth noting is that as the masses of the X, Y, Z particles get larger, the amount of baryon asymmetry goes down for given M_S as does the strength of the $\Delta B = 2$ transition giving interesting correlation between the N–$\bar{N}$ process and baryon asymmetry.

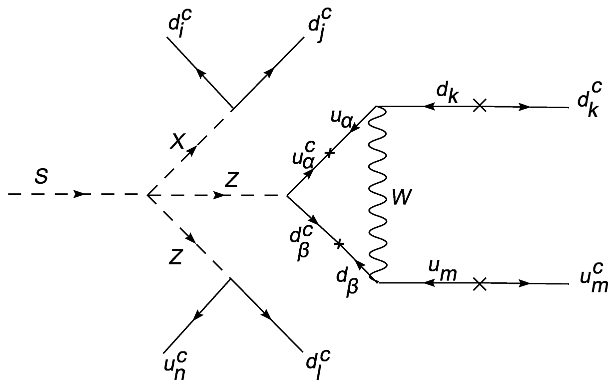

Figure 3. One-loop diagram for S decays.

Using $m_c(m_c) = 1.27$ GeV, $m_b(m_b) = 4.25$ GeV, $m_t = 174$ GeV, $V_{cb} \simeq 0.04$, $M_S = 200$ GeV and $|h_{33}| \simeq |h_{23}|$ $o1$ we find $\epsilon_B \sim 10^{-8}$.

There is a further dilution of the baryon asymmetry arising from the fact that $T_d \ll M_S$ since the decay of S also releases entropy into the Universe. In this case the baryon asymmetry reads

$$\eta_B \simeq \epsilon_B \frac{T_d}{M_S}. \tag{14}$$

In order that this dilution effect is not excessive, there must be a lower limit on the ratio T_d/M_S. From our estimate above we require that $T_d/M_S \geq 0.01$. Since the decay rate of the S boson depends inversely as a high power of $M_{X,Y}$, higher X, Y bosons would imply that $\Gamma_S \sim H$ is satisfied at a lower temperature and hence give a lower T_d/M_S. In figure 3 we plotted $M_{X,Y}$ vs. M_S using $T_d \geq M_S/100$, and the constraint $G_{N\bar{N}} \leq 10^{-28}$ GeV^{-5}. The coupling $\bar{\lambda}^4 \equiv \lambda_1 h_{11}^2 f_{11} \sim \lambda_2 h g_{11}^2$. This in turn implies that $\tau_{N-\bar{N}}$ must have an upper limit. For instance, for choice of the coupling parameters $\lambda \sim f \sim h \sim g \sim 10^{-3}$, and $M_S \simeq 200$ GeV we find $\tau_{N-\bar{N}} \leq 10^{10}$ s.

5. Conclusion

In conclusion, we have presented a realistic quark-lepton unified model where despite the high see-saw (v_{BL}) scale (in the range of $\sim 10^{12}$ GeV), the N–$\bar{N}$ oscillation time can be around 10^{10} s due to the presence of a new supersymmetric graph and accidental symmetries of the Higgs potential (also connected to supersymmetry). This oscillation time is within the reach of possible future experiments. We have also found a new way to generate the baryon asymmetry of the Universe for the case when N–$\bar{N}$ oscillation is observable. These results should provide a motivation to conduct a new round of search for N–$\bar{N}$ oscillation.

Acknowledgements

I would like to thank K S Babu, B Dutta, Y Mimura and S Nasri for collaborations that led to the results reported in this talk. I would like to thank Y Kamyshkov for many discussions on the experimental prospects for N–$\bar{N}$ oscillation. This work is supported by the National Science Foundation Grant No. Phy-0354401.

References

[1] G 't Hooft, *Phys. Rev. Lett.* **37**, 8 (1976)
[2] A D Sakharov, *JETP Lett.* **5**, 24 (1967)
[3] J C Pati and A Salam, *Phys. Rev.* **D10**, 275 (1974)
[4] H Georgi and S L Glashow, *Phys. Rev. Lett.* **32**, 438 (1974)
[5] S Dimopoulos, S Raby and F Wilczek, *Phys. Lett.* **B112**, 133 (1982)
[6] V A Kuzmin, *JETP Lett.* **12**, 228 (1970)
[7] S L Glashow, *Cargese Lectures* (1979)
[8] R N Mohapatra and R E Marshak, *Phys. Rev. Lett.* **44**, 1316 (1980)
[9] P Minkowski, *Phys. Lett.* **B67**, 421 (1977)
M Gell-Mann, P Ramond and R Slansky, *Supergravity* edited by P van Nieuwenhuizen *et al* (North Holland, Amsterdam, 1979) p. 315
T Yanagida, in *Proceedings of the Workshop on the Unified Theory and the Baryon Number in the Universe* edited by O Sawada and A Sugamoto (KEK, Tsukuba, Japan, 1979) p. 95
S L Glashow, The future of elementary particle physics, in *Proceedings of the 1979 Cargèse Summer Institute on Quarks and Leptons* edited by M Lévy *et al* (Plenum Press, New York, 1980) p. 687
R N Mohapatra and G Senjanović, *Phys. Rev. Lett.* **44**, 912 (1980)
[10] For a high scale seesaw model where N–$\bar{N}$ oscillation is observable, see K S Babu and R N Mohapatra, *Phys. Lett.* **B518**, 269 (2001). In this model however, the lightest SUSY particle, the neutralino is unstable and cannot be a dark matter candidate. Also leptogenesis cannot explain the origin of matter, since fast $\Delta B = 2$ interactions erase any baryons generated via this mechanism
[11] B Dutta, Y Mimura and R N Mohapatra, *Phys. Rev. Lett.* **96**, 061801 (2006)
[12] Z Chacko and R N Mohapatra, *Phys. Rev.* **D59**, 055004 (1999); hep-ph/9802388
[13] K S Babu, R N Mohapatra and S Nasri, to appear
See also talk by K S Babu, at this Symposium
[14] M Baldo-Ceolin *et al*, *Z. Phys.* **C63**, 409 (1994)
[15] KAMIOKANDE Collaboration: M Takita *et al*, *Phys. Rev.* **D34**, 902 (1986)
J Chung *et al*, *Phys. Rev.* **D66**, 032004 (2002), hep-ex/0205093
[16] Y A Kamyshkov, hep-ex/0211006
[17] S Rao and R Shrock, *Phys. Lett.* **B116**, 238 (1982)
J Pasupathy, *Phys. Lett.* **B114**, 172 (1982)
Riazuddin, *Phys. Rev.* **D25**, 885 (1982)
S P Misra and U Sarkar, *Phys. Rev.* **D28**, 249 (1983)