FCNC, CP violation, and impure Majorana neutrinos

Dan-Di Wu

HEP, Prairie View A&M University, Prairie View, TX 77446-0355, USA

The tree level diagonalization of a neutrino mass matrix with both Majorana and Dirac masses is discussed in a general context. Flavor changing neutral currents in such models are inevitable. Rephasing invariant quantities characterizing CP violation in FCNC Fermion-Higgs interactions are identified. At the one loop level, the mass eigenstates become an impure Majorana type. The possibility of a significant change in the mass spectrum for the left-handed neutrinos is explored, with an example of two species of neutrinos. Neutrino oscillations with impure Majorana neutrinos are also discussed.

1 E-mail: wu@hp75.pvamu.edu or danwu@physics.rice.edu
Data from more and more neutrino oscillation experiments hint at non-zero neutrino masses\[1\]. Furthermore, the available data seem to indicate that the mass squared differences of the neutrinos, if they exist, form a hierarchical spectrum. This is to be compared with the fact that the masses of each type of charged fermions have a hierarchical spectrum. If neutrino oscillations exist, this will be a very strong motivation for introducing neutrino masses in the theory. Another motivation for massive neutrinos is that they could be a candidate for the dark matter in the universe.

Neutrinos are very different from other fermions, because it is electrically neutral. Since neutrinos are neutral, it may have two different types of masses, the Dirac masses, which is similar to the masses of all charged fermions; and the Majorana masses, which require the violation of the fermion number conservation. Since fermion number conservation is not regarded a fundamental principle in most gauge interaction models, neutrinos are very likely to have Majorana masses, if they are massive.

In the minimal standard model (SM) neutrinos are massless. There is no way to write a mass term for the neutrinos in SM without breaking the gauge symmetry of this model. However, one can give neutrinos masses with a minimal extension of the SM. One can, for example, add a Higgs triplet $\Delta_L$ to SM and obtain a new term

$$\kappa^L_{ij} \psi^T_{Li} C \psi_{Li} \psi_{Lj} + h.c. \quad (1)$$

where $\psi_L = \left( \nu_L \ e_L \right)$ is the lepton doublet. Note, $\nu_L$ here as an eigenstate of interactions is a Weyl spinor. $\kappa^L_{ij}$ is a $3 \times 3$ matrix of the Majorana coupling constants. This matrix must be symmetric, because $\psi^T_{Li} C \psi_{Li} = \psi^T_{Lj} C \psi_{Lj}$. The anti-symmetric part of the coupling matrix, if it exists, does not contribute (The anti-symmetric part is cancelled out by itself.). If the vacuum expectation value of $\Delta_L$ is non-zero, then neutrinos will be massive, with a symmetric mass matrix $M_L$, and the lepton number conservation will be spontaneously broken. One needs both small coupling constants and small VEV of $\Delta_L$ to accommodate for the observed tiny neutrino masses in such a model.

One can also introduce, instead of $\Delta_L$, right-handed neutrinos $\nu_{Ri}$. One then has

$$g_{ij} \nu^c_{Ri} C \phi^T \psi_{Lj} + \kappa^R_{ij} \nu^c_{Ri} \nu^c_{Rj} + h.c. \quad (2)$$
where $\kappa_{ij}^R$ is a symmetric Majorana coupling matrix, and $\mu$ is a mass scale. Here again we meet with Majorana mass terms, the mass terms for the right-handed neutrinos. These mass terms do not vanish, unless fermion number conservation is imposed on the model. The Dirac-type Yukawa coupling constants $g_{ij}$ are an arbitrary matrix. A possible criticism to this model says that since the fields $\nu_{ri}$ do not have any gauge interactions of the SM, it does not look appealing to have them; unless there are some other interactions beyond the SM interactions which involve $\nu_{ri}$. After $\phi$ develops VEV, we obtain a $6 \times 6$ symmetric mass matrix, if there are three generations of fermions. The mass matrix of mass terms (2) is, in the form of a $2 \times 2$ block matrix

$$
\begin{pmatrix}
0 & M \\
M^T & M_R
\end{pmatrix}
$$

(3)

where $M_{ij} = \frac{1}{2} g_{ij} v$, and $M_{Rij} = \kappa_{ij}^R \mu$. When $\mu = 0$ this is the model with only Dirac masses.

Combining mass terms (1) and (2), a general mass matrix of the neutrinos has the form

$$
\begin{pmatrix}
M_L & M \\
M^T & M_R
\end{pmatrix}
$$

(4)

which is symmetric[2], with $M_{Lij} = \kappa_{ij}^L < \Delta_L >$.

Models like the left-right symmetric models, the $SO(10)$ grand unification models[3], and their supersymmetric extensions, necessarily require massive neutrinos. A characteristic property of these models is that they all need neutrinos to have Majorana masses, if neutrino masses of the order of the masses of the charged fermions are to be avoided. Indeed, in these models, neutrinos obtain Dirac masses which are “naturally” compatible with the masses of charged fermions. In order to produce the acceptable tiny masses for left-handed neutrinos, out of these much too large Dirac neutrino masses, a see-saw mechanism must be applied[2]. Consequently, one needs at least an original large right-handed Majorana neutrino mass matrix. The effective neutrino mass terms in these models can be represented by Eq (2), or by the combination of Eqs (1) and (2). The latter case with a non-zero VEV for $\Delta_L$ is not popular. In either case, it is not difficult to make the model satisfy the condition $\kappa_{ij}^L = \kappa_{ij}^R$ by suitably applying the L - R symmetry of the models.

This note will be divided into two parts. The first part is limited to the tree-level discussions. Flavor changing neutral currents (FCNC) interactions mediated by Higgs bosons
and neutrino related CP violation are discussed in a general context. First loop effects will be discussed in the second part. These effects include impure Majorana states as mass eigenstates and the asymmetry of decay products due to CP violation. The potential of a significant change in the mass splitting between Majorana neutrinos is explored.

First let us discuss the general diagonalization problem in order to see the FCNC neutrino interactions on a precise basis. A symmetric matrix $S$, like that in Eq (3) or (4), can be diagonalized by one unitary matrix in a symmetric way[4],

$$U^T S U = D$$

where $D$ is a diagonalized matrix with all elements real and no less than zero, and $U = U_1^T \times e^{i\phi}$ with $U_1$, a unitary matrix which satisfies

$$U_1 S S^\dagger U_1^\dagger = D^2;$$

and $\phi$ is a real diagonal matrix such as to make $D$ real and definitely non-negative.

If the see-saw mechanism is at work, block diagonalizing the complete neutrino mass matrix of the type of Eq (3) with $M_R >> M$, one ends up with a tiny Majorana mass matrix for the left-handed neutrinos, as expressed in the following formula[2]

$$M^{NL} = - M M_R^{-1} M^T$$

where $M_R$ is the large $3 \times 3$ right-handed Majorana mass matrix, and $M$ is the $3 \times 3$ Dirac mass matrix of the neutrinos. Here we assume the original left-handed Majorana mass matrix is either zero or negligible ($M_L = 0$). Otherwise there will be an additional $3 \times 3$ original left-handed Majorana mass matrix at the right-hand side of the equation. Note that Eq (6) is symmetric, which is consistent with its Majorana property. This formula is widely used in the literature, but sometimes wrongly recorded.

Second, one notices that the flavor changing neutral currents (FCNC) in the neutrino sector is inevitable and copious, when Dirac and Majorana masses are both present in the models, particularly in any see-saw models of neutrino masses. Indeed, expressing the unitary matrix which diagonalizes the full neutrino mass matrix in a $2 \times 2$ block form

$$U = \begin{pmatrix} U^N_L & U_{12} \\ U_{21} & U^N_R \end{pmatrix},$$


one finds that the following matrices are diagonalized very accurately after the full mass
diagonalization

\[ M_D^R = U_R^{NT} M_R U_R^N \]  
(8)

and

\[ M_D^L = U_L^{NT} M^{NL} U_L^N \]  
(9)

where \( M^{NL} \) is defined in Eq (6). However the new Majorana coupling constants between
the left-handed neutrinos and the left-handed triplet \( \Delta_L \)

\[ \kappa^N = U_L^{NT} \kappa_L U_L^N = U_L^{NT} U_R^{N*} \frac{M_D^R}{V_R} U_R^{N*} U_L^N \]  
(10)

is not diagonalized, where left-right symmetry is assumed in the second step. Neither is the
new Yukawa coupling constants among the left-handed neutrinos, right-handed neutrinos
and the doublet Higgs \( \phi^0 \)

\[ g^N = U_L^{NT} g U_R^N. \]  
(11)

Consequently, flavor changing neutral currents (FCNC) mediated by the neutral component
of the triplet \( \Delta^0 \) and by the neutral component of the doublet \( \phi^0 \) exist in general.

One may wonder how these FCNC interactions have affected the abundance of different
species of relic neutrinos from the big-bang universe. A quick examination tells that these
interactions are too weak to have such an effect. Actually, the average rate of the relevant
process (e.g. \( \nu_{L3} + \nu_{L1} \rightarrow \Delta_{L0} \rightarrow 2 \nu_{L1} \)) \( \langle \sigma v \rangle_{T_d} \) at the temperature \( T_d \sim 10 \) eV, the assumed
heaviest mass of left-handed neutrinos, is much smaller than the then Hubble constant \( H(T_d) \)
for all reasonable masses of \( \Delta_{L0} \).

As CP violation involved in massive neutrino models is concerned, one notices that CP
violation comes from several sources. There are CP violations mediated by \( W_L \) and \( W_R \),
which are characterized by their corresponding CKM matrices. Of special interest is CP
violation rooted in the non-diagonal neutrino interactions of (10) and (11). CP violation
in Higgs coupling constants is defined by the imaginary parts of the rephasing invariant
quartets\[5\], e.g.

\[ \Delta_{\alpha \beta} = \varepsilon_{ijk} \varepsilon_{\alpha \beta \gamma} \kappa^N_{ij} \kappa^N_{\alpha \beta} \kappa^N_{\gamma \gamma} \kappa^N_{k \beta} \]  
(no summation)  
(12)

\[ ^2 \text{Or } U_L^{NT} M^{NL} U_L^N + U_L^{NT} M^{NL} U_L^N \text{ if } M_L \text{ in (4) is tiny but still significant.} \]
It is found that it is impossible to make a non-trivial quartet by interfering two tree diagrams for the Majorana couplings in Eq (10). Therefore, tree level CP violation with FCNC Majorana couplings does not exist. Such CP violation exists for the Yukawa couplings in Eq (11), which is quite similar to the charged current gauge couplings, except for the lack of universality for the Yukawa couplings. Note, since $\kappa^N$ (or $g^N$) here is not unitary, therefore $\text{Im}\Delta_{\alpha}^{\kappa}$ (or $\text{Im}\Delta_{\alpha}^{g}$) for different processes are different. These quantities are small if there is a hierarchy in the matrix elements. Since all the phases in the coupling matrices are subject to change by rephasing the neutrino fields, there are only four independent useful phases in $g^N$ and three in $\kappa^N$, if these matrices are of dimension $3 \times 3$. Consequently, when one $\Delta_{\alpha}$ is purely imaginary (so-called having a maximal CP violation), the others may not. Such a concept of maximal CP violation is widely used in the estimation of baryon excess due to CP violation[6]. The number of pure imaginary $\Delta_{\alpha}$ is limited to three for $\kappa^N$ and four for $g^N$ respectively. One can also find a matrix with all its $\Delta_{\alpha}$ having significant phases, for instance, a matrix with the following phase distribution:

\[
\begin{pmatrix}
0 & \sqrt{i} & \sqrt{i} \\
\sqrt{i} & 0 & \sqrt{i} \\
\sqrt{i} & \sqrt{i} & 0
\end{pmatrix},
\]

where a zero element means the matrix element is real. This is to be compared with maximal CP violation in the CKM matrix, where none of the quartets can be made purely imaginary because of the unitarity constraint.

Other interesting new CP violation sources are in the charged Majorana couplings. For example, the $\Delta_L$ coupling of Eq (1)

\[
L_T^C \left( U_L^{T} \kappa_L^N U_L^T \right) \Delta_L^+ \nu_{Lj} + \text{transposed} + \text{h.c.}
\]  

(13)

There will be a similar term for the right-handed neutrinos, if the model is L-R symmetric, and there is only one pair of $\Delta_L - \Delta_R$,

\[
L_R^C \left( U_R^{T} \frac{M_R^D}{V_R} U_R^T \right) \Delta_R^+ \nu_{Rj} + \text{transposed} + \text{h.c.}
\]  

(14)

Of course each of these coupling constants have their own corresponding quartets to be discussed.
The above discussions are only valid at the tree level. When loop effects are introduced, there will be some more interesting physics, in particular, the physics somehow resembles that in the $K^0 - \bar{K}^0$ system. This piece of physics, especially that of right-handed neutrinos, has recently aroused some enthusiasm due to a brilliant paper by Flanz, Paschos, Sarkar, and Weiss (FPSW)[7]. Essentially, the Weyl fields $\nu_R$ and $\nu^c_R$, which are the eigenstates of interactions, are mutually CPT conjugated. This system would have been an exact copy of the $K^0 - \bar{K}^0$ system with asymmetric decay products (e.g. the $l^+ / l^-$ ratio is not 1), if mixing mass terms between $\nu_R$ and $\nu^c_R$ were not forbidden by Lorentz invariance. FPSW found a two species system, which may fulfill that kind of mixing. A further study of their system will be given here. An extension of their discussion to the left-handed neutrinos will also be attempted.

Consider first two neutrino states of the same chirality, the Hamiltonian at the tree level after mass diagonalization is ($\nu_R \equiv \nu$, and assuming CPT)

$$
(\nu^c_1, \nu^c_2, \nu_1, \nu_2) \hat{H}^0 \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu^c_1 \\ \nu^c_2 \end{pmatrix},
$$

(15)

where

$$
\hat{H}^{(0)} = \begin{pmatrix} 0 & 0 & M_{11} & 0 \\ 0 & 0 & 0 & M_{22} \\ M^*_{11} & 0 & 0 & 0 \\ 0 & M^*_{22} & 0 & 0 \end{pmatrix}.
$$

(16)

Note here only neutrinos with a specific chirality (e.g. $\nu_R$ and $\nu^c_R$, instead of $\nu^c_L$) are considered. This form is convenient for separating the absorptive part from the dispersive part of the Hamiltonian, which will become clear in a moment. At the tree level, the Weyl and the Majorana states are equivalent, so far as the mass (and the decay life time) eigenstates are concerned. This is not true when loop effects are considered. The loop corrections (see Fig. 1, which is copied from Ref[7].) to the zeros in $\hat{H}^{(0)}$ are convergent[8], if the theory is renormalizable and these corrections do not have counter terms. Including loop effects, the
total Hamiltonian is expressed as

$$\hat{H} = \hat{H}^{(0)} + \hat{H}^{(corr)} = \begin{pmatrix}
0 & 0 & \alpha_{11} & \alpha_{12} \\
0 & 0 & \alpha_{12} & \alpha_{22} \\
\tilde{\alpha}_{11} & \tilde{\alpha}_{12} & 0 & 0 \\
\tilde{\alpha}_{12} & \tilde{\alpha}_{22} & 0 & 0
\end{pmatrix}$$

(17)

There are no odd terms of the eigenvalue \(\lambda\) in the eigenequation for \(\hat{H}^{(0)} + \hat{H}^{(corr)}\). One therefore has

$$\lambda_2 = -\lambda_1, \quad \lambda_4 = -\lambda_3. \quad (\text{Im}\lambda_1 \leq 0, \quad \text{Im}\lambda_3 \leq 0)$$

(18)

The eigenvectors turn out to be almost Majorana states, or impure Majorana states as they are called. These eigenstates \(M_\beta = (x_\beta, y_\beta, z_\beta, w_\beta)\), with \(\beta = 1, 2, 3, 4\), are expressed as, up to normalization constants

$$x_\beta = (\alpha_{11}\alpha_{22} - \alpha_{12}^2)\tilde{\alpha}_{12} + \lambda_2^2\alpha_{12},$$
$$y_\beta = - (\alpha_{11}\alpha_{22} - \alpha_{12}^2)\tilde{\alpha}_{11} + \lambda_3^2\alpha_{22},$$
$$z_\beta = \lambda_3(\alpha_{12}\tilde{\alpha}_{12} + \alpha_{12}\tilde{\alpha}_{11}),$$
$$w_\beta = - \lambda_3(\alpha_{11}\tilde{\alpha}_{11} + \alpha_{12}\tilde{\alpha}_{12}) + \lambda_3^3.$$  

(19)

A phase redefinition of the eigenstates with \(\beta = 2, 4\), which are almost CP odd, will change the signs of their eigenvalues, so to let them have positive widths. In other words, \(M_1\) and \(M_2\) are two orthogonal mass eigenstates with exactly the same mass and width. In the following we will take a phase convention to make \(\tilde{\alpha}_{11} = \alpha_{11}, \tilde{\alpha}_{22} = \alpha_{22}\).

It is easy to discuss the solutions in two special cases.

**Case 1:** \(|\alpha_{22}| \gg |\alpha_{11}| \gg |\alpha_{12}|\). One finds, \(\lambda_1 \approx \alpha_{22}, \lambda_3 \approx \alpha_{11}\), and to the leading orders (\(\lambda_3\) must be calculated to the next leading orders in order to obtain the following answer.), the eigenvectors are: \((\delta = \alpha_{12}/\tilde{\alpha}_{22}, \tilde{\delta} = \tilde{\alpha}_{12}/\alpha_{22}, \gamma = \alpha_{12}/\tilde{\alpha}_{12})\)

$$M_1 \sim (\delta \quad 1 \quad \tilde{\delta} \quad 1),$$
$$M_2 \sim (-\delta \quad -1 \quad \tilde{\delta} \quad 1),$$
$$M_3 \sim (1 \quad -\gamma\delta \quad 1 \quad -\delta),$$
$$M_4 \sim (-1 \quad \gamma\tilde{\delta} \quad 1 \quad -\tilde{\delta}).$$

(20)
Note that the mass (decay) eigenstates are of impure Majorana type. Assuming charged leptons in \( \nu_i \) decays while an equal amount of anti-charged leptons in \( \nu_i^c \) decays are found, one then has the lepton-anti-lepton asymmetry in the decays of these impure Majorana particles:

\[
\delta_1 = \frac{\Gamma(M_1 \to l^- + x) - \Gamma(M_1 \to l^+ + x)}{\Gamma(M_1 \to l^- + x) + \Gamma(M_1 \to l^+ + x)} = \frac{|\tilde{\alpha}_{12}|^2 - |\tilde{\alpha}_{11}|^2}{2|\alpha_{22}|^2} = \frac{\text{Im} M_{12} \Gamma_{12}^*}{M_{22}^2 + \Gamma_{22}^2/4},
\]

\[
\delta_3 = \frac{\Gamma(M_3 \to l^- + x) - \Gamma(M_3 \to l^+ + x)}{\Gamma(M_3 \to l^- + x) + \Gamma(M_3 \to l^+ + x)} = \delta_1.
\]

The formula of \( \delta_1 \) can be further expressed by nontrivial \( \Delta_{i\alpha} \) as illustrated in Ref[4, 1986].

**Case 2:** \( |\alpha_{12}| \gg |\alpha_{22}| \gg |\alpha_{11}| \).

In this case, the \( \Gamma \) part of the Hamiltonian is negligible. The interesting new physics is twofold: First, the masses are enhanced from \( \alpha_{11} \) and \( \alpha_{22} \) to about \( \alpha_{12} \). Second, the splitting is enhanced from \( \alpha_{22} \) to \( \sqrt{\alpha_{22}\alpha_{12}} \). Since both masses are now of the order of \( \alpha_{12} \), this is a fascinating mechanism to obtain almost degenerate masses (in terms of their mass ratio being close to 1) and a large mixing. It seems that this mass spectrum is not favored by the present data, although the present data are still to be clarified. The mass eigenstates are now

\[
\begin{pmatrix}
M_1 \\
M_2 \\
M_3 \\
M_4
\end{pmatrix} \approx \frac{1}{2} \begin{pmatrix}
1 & 1 & 1 & 1 \\
1 & 1 & -1 & -1 \\
-1 & 1 & 1 & -1 \\
-1 & 1 & -1 & 1
\end{pmatrix} \begin{pmatrix}
\nu_1 \\
\nu_2 \\
\nu_1^c \\
\nu_2^c
\end{pmatrix},
\]

(22)

which are pure Majorana states, if the decay rates are neglected. Looking at the \( 4 \times 4 \) full mixing, one may wonder whether it necessary to work on a \( 4 \times 4 \) mixing matrix when discussing neutrino oscillations of two species of neutrinos. To discuss this, let the gauge couplings be, in the \( 4 \times 4 \) form,

\[
\left( \bar{L}_L W^+ \bar{L}_L W^+ \bar{L}_L W^- \bar{L}_L W^- \right) V \begin{pmatrix}
M_1 \\
M_2 \\
M_3 \\
M_4
\end{pmatrix}
\]

(23)

where the \( 4 \times 4 \) mixing matrix \( V \) is

\[
V = \frac{1}{2} \begin{pmatrix}
\rho_+ & \rho_+ & -\rho_- & -\rho_- \\
\rho_- & \rho_- & \rho_+ & \rho_+ \\
\rho_+ & -\rho_+ & \rho_- & -\rho_- \\
\rho_- & -\rho_+ & \rho_- & \rho_+
\end{pmatrix}, \quad (\rho_\pm = \cos \theta \pm \sin \theta)
\]

(24)
where $\theta$ is the tree level Cabbibo angle between the two species. Suppose at the production point a neutrino is produced together with $l^+_1$, then its wave function at a later time $t$ will be

$$\psi(t, \nu_1) = \frac{1}{2}[(\rho_+ M_1 + \rho_+ M_2)e^{-im_1 t} - (\rho_- M_3 + \rho_- M_4)e^{-im_3 t}]$$

(25)

One then has, at the detector,

$$|\langle \psi(t, \nu_1)|l_{1,2}\rangle|^2 = \frac{1}{2} \{1 \pm \cos 2\theta \cos[(m_1 - m_3)t]\}.$$  

(26)

A special situation is when $\Delta m \Delta t >> 2\pi$, where $\Delta t$ is the uncertainty of the time measurement, the oscillation part is wiped out and the two species seem to be 45$^0$ mixed.

Finally, let us calculate the loop diagrams in Fig. 1 in order to estimate the size of the effects. The boson and fermion in the loops of Fig. 1 can be $(\Delta^-/\Delta^0, l^+ / \nu)$, or $(\phi^{+}/\phi^0, l^- / \nu)$, and the first combination has a potential to contribute a significant effect. It has been realized that the Higgs-fermion couplings may be large, since the discovery of the top quark. As an example, the first combination will be considered here. For right-handed neutrinos, in the basis where the right-handed neutrino masses are already diagonalized at the tree level,

$$\alpha_{12} = \frac{m_{R1} m_{R2}}{V^2_R} \sum_i M^R_i I_R$$

(27)

where $M^R_i = (U_{1i}^0 U_{2i}^{l*})_{R}$, $I_R = \frac{m^2_{R}}{16\pi^{2} \mu_{R}^{2}}(-1 - i\pi \frac{m^2_{R}}{\mu_{R}^{2}})$ and $\mu_{R}$ is the mass of $\Delta_{R}$. The calculation is scale ($P^2$) sensitive because the outside propagator is $\frac{P + m_{R}}{P^2 - m_{R}^{2}}$. The uncertainty in the momentum flow ($P^2$) will disappear in special physical situations. $P^2$ is chosen to be $P^2 = \frac{1}{2}(m^2_{R1} + m^2_{R2})$ in (27) and the assumption of $P^2 >> \mu_{R}^{2} + m^2_{l}$ is made.

The diagrams in Figure 1 for left-handed neutrinos do not enjoy GIM suppression, and therefore are divergent. A more careful calculation is needed in order to obtain loop corrections. A guess is that because of the smallness of the left-handed neutrino masses, which are produced by the see-saw mechanism, the corrections can be large to satisfy the condition for Case 2, perhaps for one pair of the left-handed neutrinos.

Some of these discussions may apply to models with both Dirac and Majorana masses for charged particles[9]. In these models, opposite charged leptons coexist.
The FCNC and Majorana interactions among neutrinos may play a role in the neutrino scattering in the early universe when the temperature is very high. These interactions do not respect lepton number conservation, e.g. one may have

\[ \nu_R + \nu_R \rightarrow \Delta^0_R \rightarrow \bar{\nu}_R + \bar{\nu}_R. \]

This process is possible because \( \Delta_R \), which is a component of the right-handed triplet, develops VEV. \( \Delta_R \) is part of the 126-plet in the \( SO(10) \) models (in some models, 126 is a composite field). These interactions provide a vehicle for lepton and anti-lepton numbers to reach an equilibrium at extremely high temperatures, even if there is a large lepton number excess at the beginning. On the other hand, the existence of CP violation in the neutrino sector plus lepton number nonconservation processes may contribute to a baryon number excess immediately after the decoupling of some heavy particles[6].

In conclusion, massive neutrino models are likely to have flavor changing neutral currents in the Higgs mediated neutrino interactions. In the models that use the see-saw mechanism to explain the smallness of the neutrino masses, FCNC are inevitable.

The author acknowledges R. Arnowitt for conversations and discussions at an early stage of this work and Z.Z. Xing for communications and comments. He thanks X.M. Zhang and H.Q. Zheng and D. Lichtenberg for encouragement. This work began during a summer visit to Texas A&M University. The National Science Foundation partially supported this work by a NSF HRD grant.
Figure 1: A one loop contribution to the mass matrix
References

1. B. Pontecovo, Zh. Eksp. Teor. Fiz. 33, (1957) 549; F. Nezrick and F. Reines, Phys. Rev. 142 (1966) 852. For a recent review see, e.g. M. Narayan, M.V.N. Murthy, G. Rajaskaran, S. Uma Sankar, Phys. Rev. D 53 (1996) 2809; see also, G.L. Gogli, E. Lisi, D. Montanino, and G. Scioscia, [hep-ph/9607251]; LSND Collaboration, Las Alamos Bulletin Board, [hep-ph/9605001].

2. M. Gell-Mann, P. Ramond, and R. Slansky, in Supergravity, ed. D. Freedman and P. van Nuenhuizen, (North Holland, 1979); T. Yanagida, KEK Proceedings (1979); R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44 (1980) 912.

3. H. Georgi, Proc. AIP, Ed. C.E. Carlson, Meeting at William & Mary College, 1974; H. Fritzsch and P. Minkowsky, Ann. Phys. (NY) 93 (1975) 193; M. Gell-Mann, P. Ramond, and R. Slansky, Rev. Mod. Phys. 50 (1978) 721; S. Rajpoot, and P. Sithikong, Phys. Rev. D22 (1980) 2244; F. Wilczek and A. Zee, Phys. Rev. D25 (1982) 553; J.A. Harvey, D. B. Reiss, and P. Remond, Nucl. Phys. B199 (1982) 223; R.N. Mahapatra, Unification and Supersymmetry, Springer-Verlag, 1986; D.D. Wu and Y.L. Wu, PVAMU-HEP-12-95 and [hep-ph/9603418]. Mod. Phys. Lett. A11 (1996) 2703.

4. D.D. Wu, Nucl. Phys. B199 (1982) 533.

5. D.D. Wu, Phys. Lett. B90B (1980)451. D.D. Wu, Phys. Rev. D33 (1986) 860R; D.D. Wu and Y.L. Wu, Chinese Phys. Lett. 4 (1987) 441.

6. See, e.g. T. Yanagida, and M. Yoshimura, Phys. Rev. D23 (1981) 2048.

7. See e.g. M. Flanz, E. A. Paschos, U. Sarkar. and J. Weiss, Los Alamos Bulletin Board, [hep-ph/9607310] and its citations, Phys. Lett. B, forthcoming.

8. S. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. D2 (1960)1285; A. de Rujula, and S. Glashow, Phys. Rev. Lett. 45 (1980) 942.

9. D.D. Wu and T.Z. Li, Nucl. Phys. B245 (1984) 532.