The $\tau$ neutrino as a Majorana particle

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Abstract

A Majorana mass term for the $\tau$ neutrino would induce neutrino - antineutrino mixing and thereby a process which violates fermion number by two units. We study the possibility of distinguishing between a massive Majorana and a Dirac $\tau$ neutrino, by measuring fermion number violating processes in a deep inelastic scattering experiment $\nu p \rightarrow \tau X$. We show that, if the neutrino beam is obtained from the decay of high energetic pions, the probability of obtaining "wrong sign" $\tau$ leptons is suppressed by a factor $O(m_{\nu\tau}^2 \theta^2/m_\mu^2)$ instead of the naively expected suppression factor $\theta^2 m_{\nu\tau}^2/E_\nu^2$, where $E_\nu$ is the $\tau$ neutrino energy, $m_{\nu\tau}$ and $m_\mu$ are the $\tau$-neutrino and muon masses, respectively, and $\theta$ is the $\nu_\mu - \nu_\tau$ mixing angle. If $m_{\nu\tau}$ is of the order of 10 MeV and $\theta$ is of the order of $0.01 - 0.04$ (the present bounds are $m_{\nu\tau} < 35 MeV, \theta < 0.04$) the next round of experiments may be able to distinguish between Majorana and Dirac $\tau$-neutrinos.
Introduction. Neutrino masses and mixing angles may provide one of the simplest tests of physics beyond the Standard Model. Since neutrinos are chargeless, their mass can a priori be either of Dirac or Majorana type. Experimental measurements of neutrino masses are, however, very hard and, in fact, the only hints of nonzero neutrino masses come from astrophysical processes rather than from laboratory experiments. Nonzero neutrino masses provide a way of resolving the discrepancy between the predicted and the observed neutrino flux coming from the sun. Indeed, the so called MSW mechanism relies not only on nonzero masses but, in analogy to what happens in the quark sector of the electroweak theory, on nonzero mixing between the neutrinos of the first two families. If this mechanism were correct, very precise predictions for the masses and Cabbibo-like mixing angles of the first two generations could be obtained from the combination of the most recent experimental data.

A strong constraint on neutrino masses comes from a cosmological argument. If all neutrinos were stable, limits on their masses may be obtained by requiring the neutrino density to be below the critical density needed to close the universe, $\sum_i m_{\nu_i} < 100\text{eV}$. According to this argument heavy neutrinos with masses in the range $1\text{KeV} - 100\text{MeV}$ can only exist assuming a sufficiently fast decay rate. Thus, one has to assume that the same unknown mechanism which gives the neutrinos a mass induces suitable decay processes turning them unstable. Moreover, if neutrinos were Dirac particles, limits on their masses would be also obtained from Supernova events, $m_{\nu} < 30\text{KeV}$.

The bounds on the neutrino masses derived from laboratory experiments are, instead, much weaker than the astrophysical bounds. For the $\tau$ neutrino, for example, the present experimental limit is $m_{\nu_{\tau}} < 35\text{MeV}$. This bound comes from an analysis of final states $X$ in the $\tau$ decay, $\tau \rightarrow \nu_{\tau}X$. If the $\tau$ neutrino had a mass of the order of a few MeV, then supernova constraints would prevent it from being a Dirac particle. If, instead, it is a Majorana neutrino, its mass will not be restricted by supernova constraints, but it will have to decay sufficiently fast in order to fulfill the cosmological bound.

In addition, if our present understanding of nucleosynthesis in the early universe were correct, a more stringent limit on the $\tau$ neutrino mass would be obtained. As a matter
of fact, independently of whether the \( \tau \) neutrino is a Dirac or a Majorana particles, its mass would be bounded to be \( m_{\nu_\tau} < 1 \text{ MeV} \), if its lifetime were larger than \( \mathcal{O}(100 \text{ sec.}) \). On the contrary, if its lifetime were shorter or of the order of 10 sec., no bound on its mass would appear from these considerations [6]. It is easy to show that this requirement would not be fulfilled if the only physics beyond the standard model were the neutrino masses and mixing angles and hence a \( \tau \) neutrino mass in the MeV range would imply either a revision of our present understanding of nucleosynthesis, or exciting new physics in the TeV range.

In the following, we shall consider \( \tau \)-neutrinos with Majorana masses in the MeV range, produced in a deep inelastic experiment. Thus, we should also assume that the same mechanism which gives the \( \tau \)-neutrino a mass makes it unstable so that cosmological and astrophysical constraints do not apply. Furthermore, one should keep in mind that it is always important to crosscheck astrophysical predictions with laboratory experiments.

**Majorana masses and mixing.** As we just mentioned, for the extent of this work we shall consider the neutrinos to be massive Majorana particles. Therefore, neutrino-antineutrino mixing becomes possible. In fact, whereas a Dirac mass term

\[
\mathcal{L}_D = m_D \bar{\nu}_L \nu_R + h.c. \tag{1}
\]

flips the spin of the neutrinos, a Majorana mass term

\[
\mathcal{L}_M = m_M \nu^T_L C \nu_L + h.c. \tag{2}
\]

mixes neutrinos with antineutrinos, where \( C \) is the charge conjugation matrix and \( \nu_{L,R} \) are the left handed and right handed components of the neutrino field, respectively. This is possible due to the fact that a Majorana field is invariant under CP conjugation so that transitions between particle and antiparticle are allowed, with probability \( \simeq m_M^2 \). If one has more than two Majorana neutrinos, e.g. \( \nu_\mu \) and \( \nu_\tau \), then there is in general a mass matrix \( M \),

\[
\mathcal{L}_M = \begin{pmatrix} \nu_\mu^T & \nu_\tau^T \end{pmatrix} CM \begin{pmatrix} \nu_\mu \\ \nu_\tau \end{pmatrix} \tag{3}
\]
and mass eigenstates $\nu_1$, $\nu_2$ should be introduced

$$
\begin{pmatrix}
\nu_\mu \\
\nu_\tau
\end{pmatrix}
= \begin{pmatrix}
\cos \theta & -\exp(-i\delta) \sin \theta \\
\exp(i\delta) \sin \theta & \cos \theta
\end{pmatrix}
\begin{pmatrix}
\nu_1 \\
\nu_2
\end{pmatrix}.
\tag{4}
$$

Note that, in the above, contrary to the $2 \times 2$ Dirac type mixing matrix\cite{7}, in addition to the mixing angle $\theta$ a phase $\delta$ appears.

In the case of a free neutrino beam, one can show, based on simple dimensional arguments, that neutrino - antineutrino oscillations will be suppressed in comparison with the usual neutrino - neutrino oscillations by a factor $(m_{\nu_\tau}/E_\nu)^2$, where $E_\nu$ is the neutrino energy, and hence, for highly energetic neutrinos, these oscillations will be unobservably small \cite{8}. However, neutrino-antineutrino oscillations may be experimentally observed if the neutrinos proceed from a highly energetic pion beam. The physics behind this observation is extremely simple and is based on the helicity properties of the neutrinos as observed from the laboratory frame \cite{9}. Indeed, if one considers the frame in which the pion is at rest, the neutrinos will be emitted with an energy $E^* = (m_\pi^2 - m_\mu^2)/2m_\pi \simeq 30$ MeV, where $m_\pi$ and $m_\mu$ are the pion and muon masses, respectively, and with definite helicity. When observed from the laboratory frame, however, a fraction of them, namely those which were moving in the backward direction with respect to the beam in the rest frame, will appear to have opposite helicity and, hence, they behave like Majorana antineutrinos. In Ref.\cite{9}, this fraction was estimated to be of the order

$$
\epsilon \simeq \left(\frac{m_\nu}{E^*}\right)^2
\tag{5}
$$

If only the $\mu$ neutrino oscillations are studied, since due to experimental limits its mass can not be larger than a few tens of keV, then, as observed in Ref.\cite{9}, this fraction is unobservably small. If the $\mu$ neutrino has a small admixture with a massive $\tau$ neutrino, with mass in the MeV range, this process, although suppressed by mixing angles, can be significantly enhanced. In this letter, it is our intention to make a detailed analysis of deep inelastic scattering experiments, $\nu p \rightarrow \tau X$, in which the $\nu$ beam is obtained from high energetic pions. We shall show how, in the light of future neutrino oscillation experiments, such analysis may provide very useful information about the nature of massive neutrinos and their masses.
The Experiment. The present experimental search for $\nu_\mu - \nu_\tau$ mixing [10] is based on a very simple physical principle: If a high energy beam of charged pions is produced in the laboratory, the $\pi^+$ ($\pi^-$) will rapidly decay into a $\mu^+$ ($\mu^-$) and a $\nu_\mu$ ($\bar{\nu}_\mu$). If the neutrino beam subsequently collides with an isoscalar target, $\mu^-$ ($\mu^+$) particles will be produced, which may be easily detected experimentally. Moreover, considering that the electroweak eigenstate $\nu_\mu$ is an admixture of two massive particles $\nu_1$ and $\nu_2$, with $\nu_2$ being predominantly a $\nu_\tau$ particle, there will be a nonzero probability of producing $\tau^-$ ($\tau^+$) after the collision. The detection of $\tau$ particles in the final states would provide the first experimental clear signature of masses and mixing angles in the neutrino sector, and hence of physics beyond the standard model.

Experiments like this have been carried out at Fermilab and new experiments are designed at CERN and Fermilab [11]. The typical experimental setup is as follows: Assume the pions of energy $\mathcal{O}(100 \text{ GeV})$ move along the z axis in the direction of the target. After the pion decay pipe (of order $\leq 1 \text{ km}$), the neutrinos which are emitted mainly in the forward direction may oscillate along an oscillation pipe of length $\mathcal{O}(1 \text{ km})$ before they hit the target, which has a typical extension $\mathcal{O}(1-10 \text{ m}^2)$. For neutrino masses in the MeV range, the oscillation is too fast to be observable, so the exact length of the oscillation pipe is not relevant for our analysis. According to the experimental setup one may introduce cuts on the angle $\vartheta$ between the pion and neutrino beam, $\vartheta < \vartheta_{cut}$, so that the neutrino beam hits the target, and on the $\tau$ lepton energy, $E_\tau > E_{cut}$, so that $\tau$ lepton identification is guaranteed. Furthermore, as it is usual in deep inelastic scattering experiments, a cut $Q^2 > Q_{cut}^2$ with $Q_{cut} = \mathcal{O}(1 \text{ GeV})$, is introduced to ensure the validity of the quark parton model involved in the analysis.

The neutrino oscillation and mixing experiment under consideration leads in general to excluded areas in the $\theta - \Delta m$ plane, where $\Delta m = m_2 - m_1$ is the difference between the mass eigenvalues of the two neutrinos. For our case ($m_2 \simeq \mathcal{O}(1 \text{ MeV}), m_1 \simeq 0$) the oscillations are too fast, and consequently a limit $\theta \leq 0.04$ can be deduced from the non observation of $\tau$ leptons in the Fermilab data [12]. The proposed two new experiments are estimated to increase the experimental accuracy by at least a factor of ten. Therefore, it is conceivably that $\tau$ events will be observed, a small fraction of which will violate lepton
number conservation, so that the nature of the corresponding \( \tau \) neutrino mass term is revealed. The diagrams which describe the reactions are shown in Fig. 1. Also shown in Fig. 1 are the corresponding reactions with a \( \nu - \bar{\nu} \) mixing Majorana propagator which leads to \( \tau \) - leptons of opposite charge.

**Amplitudes and Cross Sections.** The total cross section consists on three different processes, the pion decay, the neutrino oscillations and finally the scattering process. Within a good approximation, for the case of \( \nu_\mu - \nu_\tau \) (\( \bar{\nu}_\mu - \bar{\nu}_\tau \)) oscillations the total cross section factorizes, that is to say, it is a product of the pion decay rate, the oscillation factor and the deep inelastic lepton scattering. In the case of neutrino - antineutrino oscillations, the process is a little more involved and the above factorization does not hold. Following Ref. [13], we compute the total amplitude as follows. We first consider the \( \pi^+ \) decay amplitude to be given by

\[
M_{i\mu} = \frac{G_F}{\sqrt{2}} f_\pi \bar{u}_i(k_i) [m_i(1 - \gamma_5) - m_{\mu}(1 + \gamma_5)] v_\mu(k_\mu)
\]

were \( G_F \) and \( f_\pi \) are the Fermi and pion decay constants, respectively, \( m_i \) is the mass of the neutrino mass eigenstates \( \nu_i \), \( k_i, \ k_\mu \), and \( k_\tau \) are the four momenta of the neutrino \( \nu_i \), the \( \mu \) and \( \tau \) leptons, respectively. The form of the \( \pi \mu \nu \) interaction vertex is essentially dictated by the pseudoscalar nature of the pion. Analogously, for the \( \pi^- \), its decay amplitude is given by

\[
M_{i\mu} = \frac{G_F}{\sqrt{2}} f_\pi \bar{u}_i(k_i) [m_i(1 + \gamma_5) - m_{\mu}(1 - \gamma_5)] v_\mu(k_\mu).
\]

The amplitude for the scattering process of the neutrino against an isoscalar target, with a \( \tau^- \) in the final state, reads,

\[
T_{\nu_i - N} = \frac{G_F}{\sqrt{2}} < J_\nu > \bar{u}(k_\tau)\gamma^\nu(1 - \gamma_5)u_i(k_i)
\]

where \( J_\nu \) is the hadronic charged current amplitude. For an antineutrino, the corresponding scattering process leads to

\[
T_{\bar{\nu}_i - N} = \frac{G_F}{\sqrt{2}} < J_\nu > \bar{v}(k_\tau)\gamma^\nu(1 - \gamma_5)v_i(k_i)
\]

The total process amplitude is a product of these two amplitudes multiplied by the factors associated to the mixing angles and the neutrino time evolution factors. Therefore,
the total amplitude for \( \pi^+ \to \mu^+ \tau^- X \) is given by

\[
T_{\pi^+ \to \tau^-} = \sum_i M_{\nu i} T_{\nu i} T_{\nu i}^* \exp(-iE_i t),
\]

(10)

where \( U_{ij} \) are the elements of the \( \nu_\mu - \nu_\tau \) mixing mass matrix,

\[
U = \begin{pmatrix}
\cos \theta & -\sin \theta \exp(-i\delta) \\
\sin \theta \exp(i\delta) & \cos \theta
\end{pmatrix}.
\]

(11)

For the lepton number violating case, instead, the total amplitude reads

\[
T_{\pi^+ \to \tau^+} = \sum_i M_{\nu i} T_{\nu i} T_{\bar{\nu} i} \exp(-iE_i t).
\]

(12)

The generalization for the \( \pi^- \) case is straightforward.

From the above expressions it is easy to get the total cross sections for the processes by computing the square of the probability amplitude. In order to do this the hadronic matrix elements of the charged currents should be given. These matrix elements may be parametrized by

\[
W_{\lambda \rho}^{(\nu, \bar{\nu})}(p, q) = \sum_X \frac{1}{2M} (2\pi)^3 \delta(p' - p - q) \sum_{S_{\text{pins}}} N(p)|J_\lambda^\pm|X(p') \langle X(p')|J_\rho^\pm|N(p) >
\]

\[
= \left( -g_{\lambda \rho} + \frac{q \lambda q \rho}{q^2} \right) W_1^{(\nu, \bar{\nu})}(\nu, Q^2) - \frac{i}{2M^2} \epsilon_{\lambda \rho \alpha \beta} p^\alpha q^\beta W_3^{(\nu, \bar{\nu})}(\nu, Q^2)
\]

\[
+ \left( p_\lambda - \frac{(pq) q_\lambda}{q^2} \right) \left( p_\rho - \frac{(pq) q_\rho}{q^2} \right) \frac{1}{M^2} W_2^{(\nu, \bar{\nu})}(\nu, Q^2),
\]

(13)

where \( q = k_\nu - k_\tau \), \( p \) is the nucleon momenta and \( \nu = pq/M \), with \( M \) the nucleon mass. We treat the W exchange in the Fermi limit, i.e. we assume that \( Q^2 = -q^2 \) is much smaller than \( m_\text{W}^2 \). This assumption is already implicit in eqs. 6 to 9. We have only considered the dominant structure functions, which are related to the scaling functions \( F_i(x) \) by

\[
F_1(x) = 2MW_1(\nu, Q^2) \quad F_2(x) = \nu W_2(\nu, Q^2) \quad F_3(x) = \nu W_3(\nu, Q^2),
\]

(14)

where \( x = Q^2/2M\nu \). In the parton model, the structure functions \( F_i(x) \) may be expressed as functions of the quark densities. In the following, in order to make quantitative estimates we have used the parametrization given in ref. [15].
One important consequence of the hadronic current structure is that the neutrino scattering process is enhanced in comparison to the antineutrino one. In fact, this can be observed, for example, in the approximation in which the charged lepton and neutrino masses are neglected, and ignoring small scaling violation terms appearing in the antineutrino cross section. The differential cross section for neutrino scattering against an isoscalar target is then given by
\[
\frac{\partial \sigma^{(\nu N)}}{\partial x \partial y} = \frac{G_F^2 M \nu}{\pi} \left( q(x) + \bar{q}(x)(1 - y)^2 \right),
\]
while for an antineutrino it reads,
\[
\frac{\partial \sigma^{(\bar{\nu} N)}}{\partial x \partial y} = \frac{G_F^2 M \nu}{\pi} \left( q(x)(1 - y)^2 + \bar{q}(x) \right),
\]
where \( q(x) = x[N_u(x) + N_d(x)] \), \( \bar{q}(x) = x[N_{\bar{u}}(x) + N_{\bar{d}}(x)] \), and \( N_{u(d)} \) and \( N_{\bar{u}(\bar{d})} \) are the up and down quark (antiquark) densities of the proton, respectively, and \( y = \nu/E_\nu \) is the fraction of neutrino energy transferred to the charged lepton. It is then obvious to observe, that since \( q(x) \gg \bar{q}(x) \), the total neutrino cross section is enhanced by approximately a factor three with respect to the antineutrino one.

Apart from the hadronic charged current matrix elements the neutrino oscillation process is important. For the lepton number conserving processes (l.c.) the relevant mass and mixing angle dependent factor is given by
\[
F_{osc}^{l.c.} = \sum_{ij} U_{i\mu} U^*_{i\tau} U_{j\mu} U^*_{j\tau} \exp \left[ i(E_j - E_i) t \right] 
\]
which can be approximately expressed as
\[
F_{osc}^{l.c.} = 2 \cos^2 \theta \sin^2 \theta \left[ 1 - \cos \left( \frac{\Delta m^2 t}{2E_\nu} \right) \right],
\]
where the second term between brackets gives the usual oscillation effect. For one of the masses in the MeV range, the oscillations in any of the planned experiments would be so fast that, once the average over the pion decay position is performed, the oscillation term gives no contribution to the total cross section.

For the lepton number violating process (l.v.), instead, the oscillation term depends quadratically on the neutrino masses. It is given by
\[
F_{osc}^{l.v.} = \sum_{ij} U_{i\mu} U^*_{i\tau} U_{j\mu} U^*_{j\tau} m_i m_j \exp \left[ i(E_i - E_j) t \right]
\]
\[ \sin^2 \theta \cos^2 \theta \left[ m_2^2 + m_1^2 - 2m_1 m_2 \cos \left( \frac{\Delta m^2}{2E_\nu} t + 2\delta \right) \right]. \quad (19) \]

Observe that, in this case the cross section has a nontrivial dependence on the CP violating parameter \( \delta \). This is a very relevant property of Majorana neutrinos, for which, in contrast to the Dirac neutrino cases, there is no need of mixing of the three generations to obtain CP violating effects. Unfortunately, it can be easily seen that, for neutrino mass values consistent with the present experimental constraints, the CP violating contribution is unobservably small and only the first term \( m_2^2 \) in the above expression gives a relevant contribution (\( m_2 \gg m_1 \)).

**Results.** The total process can now be straightforwardly computed. We shall not give the details here, but just refer to the relevant properties. The ratio of the lepton number conserving to the lepton number violating process rate, in the case of initial \( \pi^\pm \) state is given by

\[ \frac{\Gamma_{l.e.}}{\Gamma_{l.v.}} = \frac{m_2^2}{2} \frac{\int V_{\alpha\alpha'} W_{\alpha\alpha'} \overline{W}_{\alpha\alpha'} C_{\alpha\alpha'}}{\int C_{\alpha\alpha'} W_{\alpha\alpha'}} \quad (20) \]

where the first factor comes from the ratio of the oscillation factors, and the matrix elements \( V_{\alpha\alpha'} \) and \( C_{\alpha\alpha'} \) are given by

\[ V_{\alpha\alpha'} = 2 \left( m_\pi^2 + m_\mu^2 \right) Tr \left[ (1 \pm \gamma_5) \not{k}_\nu \gamma^\alpha \not{k}_\tau \gamma^{\alpha'} \right] \]

\[ + 2 m_\mu^2 Tr \left[ (1 \mp \gamma_5) \not{k}_\mu \gamma^\alpha \not{k}_\tau \gamma^{\alpha'} \right] \quad (21) \]

\[ C_{\alpha\alpha'} = 2 m_\mu^2 \left( m_\pi^2 - m_\mu^2 \right) Tr \left[ (1 \pm \gamma_5) \not{k}_\nu \gamma^\alpha \not{k}_\tau \gamma^{\alpha'} \right], \quad (22) \]

and an integration over phase space is understood. In order to get an idea of the order of magnitude of the above ratio, Eq.\( (20) \), we can use the approximation \( k_\nu \approx k_\mu \). In this case, it is easy to show that, approximately (\( m_2 \approx m_\nu \))

\[ \frac{\Gamma_{l.e.}}{\Gamma_{l.v.}} \approx \frac{m_\nu^2}{2m_\mu^2} \frac{\left( m_\pi^2 + 2m_\mu^2 \right) \sigma^{\nu(\bar{\nu})}}{\left( m_\pi^2 - m_\mu^2 \right) \sigma^{\bar{\nu}(\nu)}} \quad (23) \]

Observe that, as we anticipated, the lepton number violating process is suppressed with respect to the lepton number conserving one by a factor of order of the ratio of the massive neutrino mass to the muon mass squared. Indeed, if the above approximation were correct, and based on the approximate magnitude of the \( \nu \) and \( \bar{\nu} \) scattering cross
section we would get approximately a 6% effect in the case of having a $\pi^-$ in the initial state, and a 0.7% effect in the initial $\pi^+$ case, when choosing $m_{\nu_\tau} = 10$ MeV. The actual numbers for the above ratio, Eq. (20), as we shall show below, are somewhat larger.

In Figures 2, 3 and 4 we present the resulting ratios of cross section for different experimental cuts and initial pion energies. We observe that the percentage of lepton number violating events is roughly 1% for $\pi^+N$ and 8% for $\pi^-N$ scattering, rather independent of the value of $E_\pi$ (see Fig. 2) and $E_{\tau,\text{cut}}$ (see Fig. 3). From Fig. 2 it also follows that the ratio of the lepton number violating to the lepton number conserving process increases slightly if one relaxes the cut on the angle between the pion and neutrino beams, $\vartheta_{\text{cut}}$, while Fig. 4 shows that it depends somewhat stronger on the value of $Q^2_{\text{cut}}$.

Conclusions. In this letter we have worked out the following idea: If the neutrinos have masses, the $\tau$-neutrino is probably the one with the largest mass. A priori, this can be either of Dirac or of Majorana type. We have discussed an experiment which distinguishes between a Majorana and a Dirac mass term. The main problem is that, if obtained from $\tau$ decays, $\tau$ neutrinos are high energetic and a neutrino mass determination with high energy neutrinos is difficult since mass effects are in general suppressed by powers of $m^2_{\nu_\tau}/E^2_{\nu_{\tau}}$. Therefore, we chose the process $\pi N \rightarrow \tau X$, where instead of a suppression proportional to the inversed $\tau$-neutrino energy squared, we found a suppression proportional to the inverse of the muon mass squared. This is much less severe and allows to determine Majorana neutrino masses in the MeV range. The price we have to pay for this advantage is an assumption about $\nu_\mu - \nu_\tau$ mixing which is needed for the process to take place, so that a suppression factor $\theta^2$ appears throughout in the cross section. For neutrino masses of the order of a few MeV, the results look promising. The lepton number violation process rate $\pi^\pm N \rightarrow \mu^\pm\tau^\pm X$ is of the order of a few percent of the lepton number conserving $\pi^\pm N \rightarrow \mu^\pm\tau^\mp X'$ one, the exact percentage depending on the $\tau$-neutrino mass and the charge of the initial pion state. We have also commented on astrophysical and cosmological restrictions. The essential point is that cosmological arguments exclude stable neutrinos of mass $O(\text{MeV})$. This is due to the fact that they would contribute too much to the present mass energy density of the universe. We have
ignored this bound, together with the constraints on neutrino masses coming from our present understanding of nucleosynthesis, by assuming that the same mechanism which gives the $\nu_\tau$ a Majorana mass makes the $\tau$ neutrino unstable and with a lifetime such that these bounds do not apply.

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Figure Captions

Fig. 1
The pion nucleon scattering processes which are discussed in this letter. Lepton number conserving (Dirac) processes are on the left, lepton number violating (Majorana) processes on the right.

Fig. 2
The ratio of the Majorana (= lepton number violating) over the Dirac (= lepton number conserving) cross section as a function of the pion energy for fixed values of the angle cut, $\tau$ energy cut and the cut in $Q$. On the horizontal axis the ratios of the Dirac cross sections for the two processes $\pi^+N$ and $\pi^-N$ are given. The cross section for the lepton number violating process is denoted by $\sigma_M$, the cross section for the lepton number conserving process is denoted by $\sigma_D$. All results have a numerical error of about one percent.

Fig. 3
The same ratio as fig. 2, now as a function of the $\tau$ energy cut.

Fig. 4
The same ratio as fig. 2, now as a function of the cut in $Q$. 