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NEUTRINOS BEYOND THE STANDARD MODEL*

J. W. F. VALLE

Stanford *Linear Accelerator Center*
Stanford University, Stanford, California 94309
and

Departament de Física Teòrica
Universitat de València
46100 Burjassot, València, Spain

ABSTRACT

I review some basic aspects of neutrino physics beyond the Standard Model such as neutrino mixing and neutrino non-orthogonality, universality and CP violation in the lepton sector, total lepton number and lepton flavour violation, ***etc.***. These may lead to neutrino decays and oscillations, exotic weak decay processes, neutrinoless double β decay, ***etc.***. Particle physics models are discussed where some of these processes can be sizable even in the absence of measurable neutrino masses. These may also substantially affect the propagation properties of solar and astrophysical neutrinos.

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1 Neutrino Masses and Non-Orthogonality

In the standard $SU(2)_L \otimes U(1)$ electroweak model there are no gauge invariant interactions that can lead to non-zero neutrino mass [1]. As a result all lepton flavours as well as total lepton number, L *, are exactly conserved. On the other hand, it is clear that neutrino masses are much smaller than those characteristic of the charged fermions of the standard model, if not zero. The clue to this puzzle may lie in the fact that neutrinos are the only electrically neutral fermions in the standard model, most likely expected to be Majorana particles. How can we then understand this relative smallness of neutrino masses?

An attractive way to understand the smallness of the neutrino masses is through the addition of right handed neutrinos. While this may seem *ad* hoc from the point of view of the standard theory, the situation changes in theories beyond the standard model where often one is forced to add right handed (RH) neutrinos in order to realize a larger gauge symmetry such as left-right, grand unified, or superstring-inspired symmetries. Alternatively Majorana neutrino masses could arise as radiative corrections and therefore be naturally smaller than the charged fermion masses. These mechanisms are illustrated in Fig. 1.

When the physical mass-eigenstate neutrinos are Majorana fermions total lepton number is broken and CP is violated even in a 2 generation world [2, 3], in contrast with the quark sector. CP violation at this level may also occur even when the physical neutrinos are massless, provided iso-singlets exist. Here I will focus on the possibility that deviations from the minimal massless-neutrino scenario arise from the existence of extra leptons. In this case the gauge currents contain additional couplings involving these neutral heavy leptons, as shown in Fig. 2. The expected pattern of weak interactions of massive neutrinos is then considerably richer than that characteristic of the quark sector.

New physics in the lepton sector may arise in neutrino mixing models both from

*Total lepton number is defined as $L = L_e + L_\mu + L_\tau$ In unified models the combination that often appears is B-L, baryon number minus lepton number. From our point of view here these are equivalent.

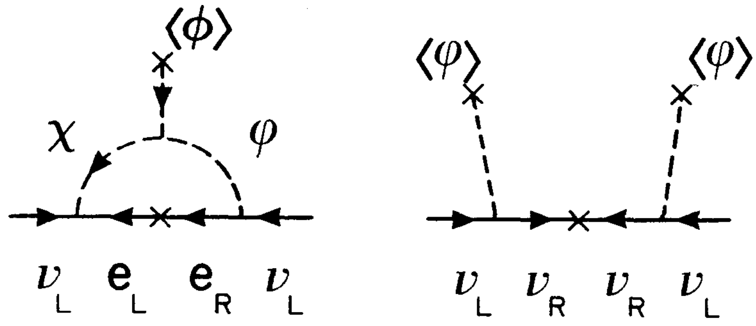

Fig. **1** : Mechanisms generating small Majorana neutrino masses

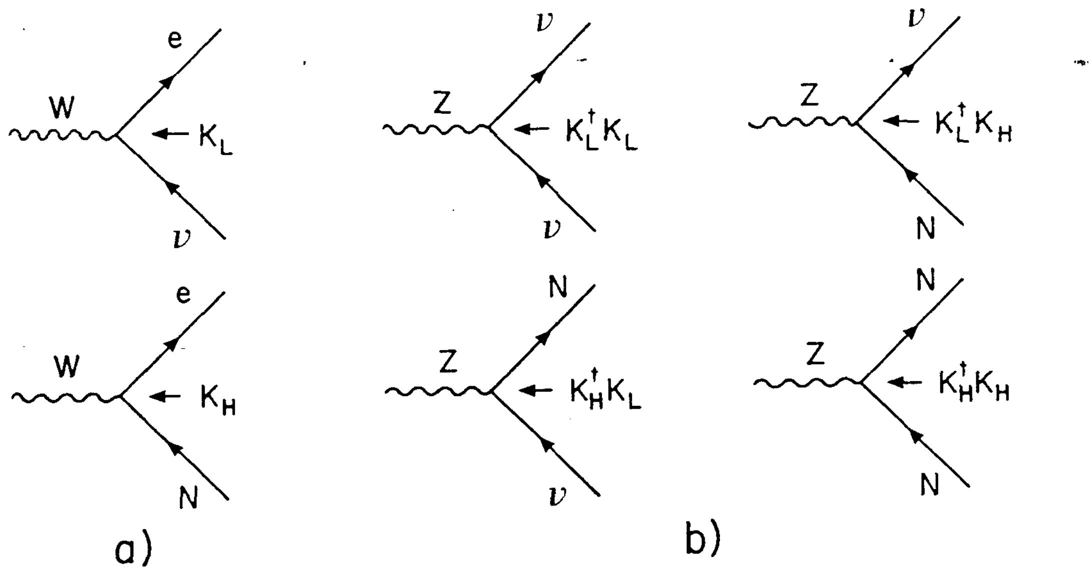

Fig. 2 : Couplings of the neutral fermions to the standard charged and neutral intermediate vector bosons.

neutrino mass and from the effective non-orthogonality of the flavour neutrinos produced in weak decays. Although these features are often present simultaneously they are a prioriindependent. For example, processes such as $\mu \rightarrow e+\gamma$, $\mu \rightarrow 3e$, μ-e conversion in nuclei, etc. can proceed at observable rates even when the physical neutrinos have no mass. It follows that experimental searches of "rare" processes should be carried out in addition. to direct searches for neutrino mass, since they probe complementary aspects of lepton physics.

2 Models

The models discussed here can be implemented in several gauge theoretic frameworks. For simplicity and generality I base the discussion on the simplest $SU(2)_L \otimes$ U(1) gauge structure. To the fermion sector of the standard electroweak model, consisting of the following set of (left-handed) fermions (repeated over generation index, i)

$$\begin{pmatrix} \nu_i \\ e_i \end{pmatrix}, \quad e^c_i$$
$$\begin{pmatrix} u_i \\ d_i \end{pmatrix}, \quad d^c_i, \quad u^c_i \tag{1}$$

and manifestly asymmetric between quarks and leptons, we now add a number of iso-singlet neutral leptons. There is no constraint on the number of isosinglets that can be added to eq. (1). The models also illustrate the important role played by B-L symmetry in lepton physics. To simplify the discussion I first consider the simplest possibility where B-L is broken explicitly or unbroken. The basic Yukawa terms leading to charged fermion masses are completely standard and these models differ only in the neutral fermion sector, as follows.

2.1 The Seesaw Model

In this model one adds one isosinglet right handed neutrino denoted ν_i^c † for each generation of isodoublet leptons, thus completing the table in eq. (1). The existence of isosinglet leptons brings in the possibility not only of Dirac mass terms for the neutrinos (analogous to those of the charged fermions) but also of gauge invariant, L violating, Majorana mass terms. The physical neutrino masses are determined from the corresponding neutrino mass matrix (in the basis (ν, ν^c)) [4]

$$\begin{pmatrix} 0 & D \\ D^T & M_R \end{pmatrix} \tag{2}$$

where the matrix D_{ij} is the Dirac mass term for the three RH neutrinos, and M_{Rij} is an isosinglet mass term, added as a bare mass. The resulting neutrino mass matrix after diagonalizing out the heavy fields is $M_L{}^{seesaw} = DM_R^{-1} D^T$ where M_R is the Majorana mass for the right handed neutrinos.

2.2 The L-Conserving Model

This model includes in addition to right handed neutrinos, an equal number of gauge singlet leptons S_i. The neutral fermion masses are restricted by imposing the exact conservation of total lepton number, thus ensuring the masslessness of neutrinos in the presence of the extra singlets. In such models total lepton-number symmetry is imposed while in the standard model it is automatic. This restriction leads to the following form for the neutral mass matrix (in the basis ν, ν^c, S)

$$\begin{pmatrix} 0 & D & 0 \\ D^T & 0 & M \\ 0 & M^T & 0 \end{pmatrix} \tag{3}$$

where the Dirac mass term coupling an $SU(2)_L$ doublet ν_L to an $SU(2)_L$ singlet ν^c is described by the matrix D while the other Dirac mass term described by the matrix M couples together two electroweak singlets ν^c and S. Hence one expects

† Right handed fermions are described in charge-conjugate notation.

the corresponding coefficients M_{ij} to be large compared with the characteristic scale for the elements D_{ij} which are proportional to the standard Higgs VEV responsible for electroweak breaking and charged fermion masses ‡. The three light neutrinos are massless *Weyl* neutrinos while the other 6 neutral 2-component leptons combine exactly into 3 heavy ***Dirac*** fermions. However individual leptonic flavours are violated in this model ***despite the fact that physical neutrinos are strictly massless*** [7, 8]. The corresponding form of the gauge currents is given in ref (7, 8, 9] and is briefly discussed below.

2.3 The p-Model

A variant of the previous model may be obtained by introducing total lepton number violation e.g. through a non-zero Majorana mass μ_{ij} for the S_i [8] in eq. (3)

$$\begin{pmatrix} 0 & D & 0 \\ D^T & 0 & M \\ 0 & M^T & \mu \end{pmatrix} \tag{4}$$

This leads to small neutrino masses determined from $M_L = DM^{-1}\mu M^{T^{-1}}D^T$ and to the possibility of neutrinoless double-beta decay. In this case the six heavy Weyl leptons split and no longer form three Dirac particles exactly. For sufficiently small μ they will form 3 quasi-Dirac heavy leptons [1 1]. Most of the results also apply to this variant of the model. In particular leptonic flavour and CP violating effects need not be suppressed by the smallness of neutrino masses.

‡A mass matrix of the form given in eq. (3) has been suggested in several theoretical frameworks such as in superstring-inspired models [1, 5, 6]. In this case the zeroes of these entries can naturally arise e.g. due to the lack of Higgs fields that could provide the usual Majorana mass terms needed in the seesaw mechanism.

3 Weak Gauge Currents

The physical mass eigenstate leptons are identified by diagonalizing the mass matrices in each of the above models. This leads to the following general form for the charged current leptonic weak interaction

$$+\frac{ig}{\sqrt{2}}W_{\mu}^{-}\ \bar{e}_L\ \gamma_\mu\ [K_L\nu'_L\ +\ K_H N_L] + \textit{\textbf{H.C.}} \tag{5}$$

where ν'_L are the light neutrinos and N_L are the neutral heavy leptons. The mixing matrix K describing the charged current leptonic weak interaction is made up of two blocks [2]

$$K\ =\ (K_L, K_H) \tag{6}$$

where the second block describes the coupling of the heavy leptons to the standard gauge bosons, Fig. 2. If one takes as starting point the weak basis where the charged lepton mass matrix is diagonal, these submatrices may be written only in terms of the matrices that diagonalize the neutral leptons. The expected magnitude of the elements of the n **x** n submatrix K_H may be severely limited by observational restrictions on the neutrino mass. For example in the seesaw model one has

$$K_H^{seesaw} = O\left(\frac{m_\nu}{M_R}\right)^{1/2} \tag{7}$$

while in the p-model one has

$$K_H^{\mu-model} = O\left(\frac{m_\nu}{\mu}\right)^{1/2} \tag{8}$$

The smallness of the neutrino masses implies the smallness of the NHL couplings in eq. (7) and eq. (8). However, in the second case when μ 4 0, m_ν 4 0 so that $K_H \rightarrow const$, as in the massless neutrino model of section 2.2. In this case the only constraints that apply on K_H are those that follow from universality considerations (see below) [§]. In contrast the m_ν constraints that apply on eq. (7) are much more restrictive due to the large M_R value required in the seesaw mechanism.

[§]The detailed forms for K in each of the models discussed in sections 2.1 and 2.2 above have been determined in ref [1, 12] by diagonalizing out the heavy fermions with the method of ref [13].

The corresponding neutral current expressed in terms of mass-eigenstate neutrinos is determined by a projective hermitian matrix P given in terms of K as $P = K^{\dagger}K$ and takes the form [2]

$$P = \begin{pmatrix} K_L^{\dagger}K_L & K_L^{\dagger}K_H \\ K_H^{\dagger}K_L & K_H^{\dagger}K_H \end{pmatrix} \tag{9}$$

There is no Glashow-Iliopoulos-Maiani mechanism due to the admixture of fermions of different weak isospin in the currents. As a result, there are neutral current couplings connecting light to heavy neutrinos, and in general off-diagonal couplings involving only light (or heavy) neutrinos among themselves [2].

In the L-conserving model or the p-model these new couplings may be considerably enhanced relative to the seesaw model expectations. As a result many processes become potentially detectable.

In the absence of a detailed model for the mass matrices all one can say about the parameters involved in the rectangular matrix K is that they consist of n orthonormal row vectors, $i.e.$ $K_L K_L^{\dagger} + K_H K_H^{\dagger} = 1$ (n is the number of generations. This imposes restrictions on the products of these n row vectors that eliminate $n(n-1)+n$ parameters [n(n − 1) are orthogonality relations and n are normalizations]. Of these, there are n unphysical phases that can be eliminated by redefining the phases of the charged leptons through a diagonal matrix of phases w,. This leads to a total of $n(3n-1)$ real parameters describing the seesaw model weak leptonic currents! This counting plus an explicit parametrization was given in ref [2]. The charged current weak interaction of leptons is clearly far more complex than that of quarks because of the possible Majorana nature of the neutrinos [2, 3]. In addition there can be light-heavy mixing as well as CP violation [12].

Depending on the model, there may be additional parameters that can be eliminated by redefining the neutral leptons. For example, in the model of section 2.2, as a result of B-L conservation, it is possible to transform the heavy Dirac neutral leptons through a diagonal matrix of phases ω_N and the massless neutrinos through an arbitrary unitary matrix U_ν. Under these transformations the submatrices K_L, K_H change as follows, $K_L \rightarrow \omega_e^{\dagger} K_L U_\nu$ and $K_H \rightarrow \omega_e^{\dagger} K_H \omega_N$. The net number of physical parameters in K is then obtained as $2n \times 2n - n(n-1) - n - n^2 - (2n-1) = n^2 + (n-1)^2$ independent parameters. These split as follows: n^2 mixing angles and $(n-1)^2$ CP violating phases. This number of parameters still exceeds ¶ those needed to describe the charged current interactions of quarks despite the strict masslessness of the physical neutrinos. In this model one can not eliminate mixing ***even if the physical neutrinos are massless.*** Moreover, ***there is CP violation*** even if there are only 2 families of massless neutrinos involved.

Laboratory experiments constrain isosinglet NHL admixture K_H in the gauge currents eq. (6) and eq. (9). For example in low energy weak decay processes only the light neutrinos can be kinematically produced and from eq. (5) one sees that the coupling of a given light neutrino to the corresponding charged lepton is decreased by a certain factor. This would lead to universality violation in low energy weak decays [14] such as β decay and μ decay, $e-\mu$ universality in $\pi \rightarrow l_2$ decay, τ lifetime, ***etc.*** These constraints have been summarized in ref [14] as $(K_H K_H^\dagger)_{ee} \lesssim 4.3 \times 10^{-2}$, $(K_H K_H^\dagger)_{\mu\mu} \lesssim 0.8 \times 10^{-2}$ and $(K_H K_H^\dagger)_{\tau\tau} \lesssim 10 \times 10^{-2}$. If light enough the isosinglets would also be produced in charm and beauty decays; ***etc.*** and these have been looked for in beam dump experiments. A tighter constraint extending up to the 10 GeV range has recently been obtained at Fermilab using the wide band neutrino beam. So far no direct constraint exists above 20 GeV or so. To extend the limits on the possible existence of neutral heavy leptons to high mass values (or discover them!) high energy ***accelerator*** experiments are needed. LEP/SLC experiments can probe isosinglet NHL mass and coupling strength parameters far beyond the range accessible to other laboratory experiments [10].

When the presence of neutral heavy leptons also engenders non-zero neutrino masses, as in sections 2.1 and 2.3 there are additional, ***stronger*** limits on the attainable rates for new physics that follow from astrophysics and cosmology as in eq. (10).

¶Note however that it is considerably smaller than the corresponding number of parameters describing the weak interaction of the leptons in a model without total lepton number conservation, such as the seesaw model or the μ model.

4 Majorons and Neutrino Stability

Many aspects of lepton physics depend on whether B-L symmetry is broken and in this case, what is the corresponding mass scale and the nature of this symmetry breaking. In a. theory where B-L is a local gauge symmetry one expects it to be broken as there is no evidence for the existence of a light intermediate gauge boson coupled to it. This is the situation in left-right theories [1] and also in some superstring-inspired models [15]. Alternatively, when an ***ungauged*** B-L symmetry is violated in a spontaneous way it leads to a massless Goldstone boson - the Majoron - denoted J. It provides neutrinos with new interactions [16, 13, 8, 17] that can play a very important role in astrophysics and cosmology.

Majoron emission generates new stellar energy loss mechanisms: once produced in a stellar environment, in Compton-like processes e.g. γ + e $\rightarrow$ e + J the weakly coupled Majorons easily escape. Suppressing the resulting energy-loss requires a stringent limit on its coupling to electrons [18].

Light neutrinos (of mass less than $O(100\ KeV)$ or so, depending on the specific model) that only have the interactions prescribed in Fig. 2 are cosmologically stable. Their contribution to the present density of the universe implies [19]

$$\sum_i \frac{g_i}{2} m_{\nu_i} \leq 97\, \Omega_\nu h^2\, eV \tag{10}$$

where $\Omega_\nu h^2 \leq 1$ and the multiplicity factor g_i = 2 for the models discussed in section 2.

The possible existence of non-standard interactions of neutrinos due to their couplings to the Majoron brings in the possibility of fast invisible neutrino decays [20, 8].

$$\nu' \rightarrow \nu + J. \tag{11}$$

where J denotes the Majoron. These can be much faster [||] than the neutral-current-mediated neutrino decay $\nu' \rightarrow 3\nu$ [2] that follows from Fig. 2. Invisible

[||] Although this was shown in ref [13] not to hold in the minimal versions [16] of the Majoron model it has been demonstrated to hold in the extended models considered in references [20] and [8].

neutrino decay has interesting astrophysical and cosmological implications [21]. It avoids the astrophysical restrictions based on nucleosynthesis considerations, on the cosmic background radiation spectrum, *etc.* as long as it is the dominant form of decay. The decay lifetime due to eq. (11) can be made sufficiently short as to satisfy the cosmological constraint following from the critical density argument and allows neutrinos to have any mass that is consistent with experiment.

5 New Physics

There is a broad range of new phenomena that can take place in various models of neutrino mass and mixing [1].

- lepton-flavour-violating effects associated to non-zero neutrino masses such as neutrino oscillations. These may be affected by the presence of matter.
- total-lepton-number-violating processes such as $\beta\beta_{0\nu}$ decay and neutrino-anti-neutrino oscillations [22, 3].
- lepton-flavour-violating effects such as $\mu \rightarrow e + \gamma$, $\mu \rightarrow 3e$, μ-e conversion in nuclei, *etc.* following from NHL admixture in the weak currents.
- leptonic CP-violating observables such as a non-zero electric dipole moment for the electron, d_e [12, 23].
- invisible neutrino decays [20, 8].
- non-standard neutrino propagation properties affecting the solar and/or supernova neutrino fluxes [24] *etc.*

The smallness of neutrino masses suggests that, with the possible exception of neutrino oscillations and $\beta\beta_{0\nu}$ decay, most of the above processes are expected to be non-observable due to experimental bounds on the neutrino mass that follow from cosmology eq. (10). However the possibility of larger neutrino masses being cosmologically acceptable due to the decay in eq. (11) allows an enhancement in the rates of some of the new effects listed above. The expected rates depend on the constraints that apply on the neutral-heavy-lepton admixtures shown in Fig. 2 and these in turn depend on the neutrino mass. For example, in the seesaw model these couplings are still too small to lead to observable flavour violating effects because they are restricted by laboratory limits on neutrino masses [7, 8].

Alternatively, the relationship between the magnitude of the rates for "new physics" in the lepton sector and the magnitude of the neutrino mass is very much ***dependent*** upon the details of neutrino mixing. For example, as noted above, both flavour and CP violation can occur in the lepton sector even when the physical neutrinos are kept strictly *massless* (for example, as a result of an exact B-L symmetry). The new couplings in Fig. 2 can then be substantially enhanced, being restricted only by universality considerations. Correspondingly, many of the new effects can be enhanced without unnatural fine-tuning of the lepton parameters nor conflict with any laboratory, cosmological or astrophysical limits on neutrino masses and lifetimes [7, 8].

5.1 Neutrino Oscillations and Solar Neutrinos .

Neutrino oscillations are one of the most basic manifestations of non-zero neutrino masses. If the neutrino mass differences are sufficiently small, then a given neutrino flavour produced by the weak interaction is a coherent mixture of all mass-eigenstate neutrinos and can, in the course of time, develop a different flavour component. This is the phenomenon of neutrino oscillations. So far no oscillations have been seen and this places constraints on the mass diferences and mixing angles, shown in Fig. 3, ref [25] **.

Oscillations could play an important role in the propagation of solar and supernova neutrinos. If neutrinos travel through matter then oscillations can be affected due to charged current (CC) forward elastic scattering on electrons, which exists for ν_e but not for ν_μ (or ν_τ) [26]. The system of evolution equations has been

**For a brief discussion of the possibility of oscillations into sterile neutrinos see ref [1] and references therein.

discussed by several authors after Mikheyev and Smirnov noted [27] that when the CC effect is included there is the possibility of a ***resonant*** neutrino conversion whenever the condition

$$\delta m^2 \cos(2\theta) = 2\sqrt{2} G_F E N_e \tag{12}$$

is satisfied, where N_e is the electron number density, E is the neutrino energy, δm^2 is the squared neutrino mass difference and θ is the mixing angle in *vacuo*. As a result sizeable neutrino conversion in matter is possible even for small values of the vaccuum mixing angle θ. To discuss solar neutrino propagation the system of evolution equations describing oscillations in matter of ***varying*** density was solved both analytically (for given density profiles), and numerically [28]. When the density decreases slowly enough a ν_e born e.g. in the solar core can simply "follow" the slowly changing Hamiltonian, find the resonance region, and end up still in the same matter eigenstate, but this time as a ν_μ or v,. This requires for the sun,$\left(\frac{\delta m_\nu}{eV}\right)^2 \simeq 10^{-4}$, and corresponds to the upper ***horizontal*** lines in Fig. 3.a. The conversion of neutrinos due to this MSW effect has been analysed also in the *non-adiabatic* (NAD) approximation where it happens for$\left(\frac{\delta m_\nu}{eV}\right)^2 \sin^2 2\theta \simeq 3 \times 10^{-8}$, leading to the ***tilted*** lines in Fig. 3.a. There is also a "large mixing" solution indicated by the ***vertical*** lines on the right of Fig. 3.a. A big effect requires the solar core density to exceed the resonance value so only neutrinos whose $E/\delta m^2 \gtrsim 10^5$ (E in MeV and δm in eV) are converted. For AD one needs $\delta m \simeq 10^{-2} eV$ ***so high*** energy neutrinos are the ones converted, as can be seen from the iso-SNU contours shown in Fig. 3.a. However for the NAD solution low energy neutrinos can be strongly suppressed for ***larger*** mixings (*lower* masses) since in this case the threshold for conversion is low as can be seen again from Fig. 3. In this figure we also show the predictions of a prototype model suggested in ref [31, 30] where the solar neutrino transitions occur preferrably in the NAD regime. The oscillation parameters are predicted by relating the tiny neutrino mass to some new particles (such as the Majoron) whose couplings can cause effects measurable in the laboratory. The prediction is relatively ***stable*** in a wide class of models. The neutrino mass arises from gauge interactions due to the exchange of heavy supersymmetric fermions as a result of ***spontaneous violation of total lepton number.***

Fig. 3.a

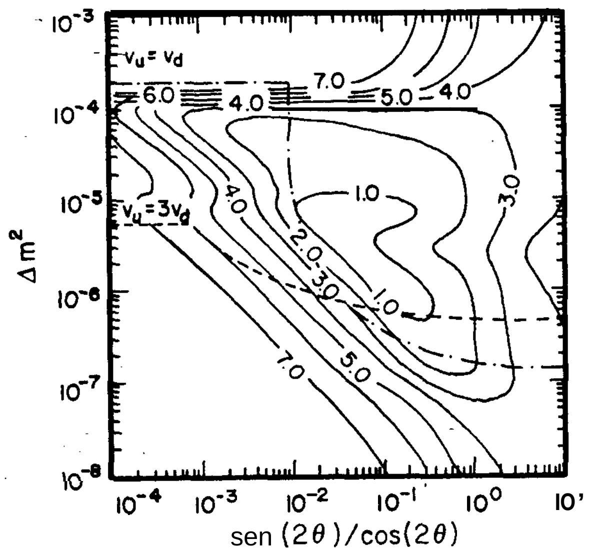

Fig. 3. b

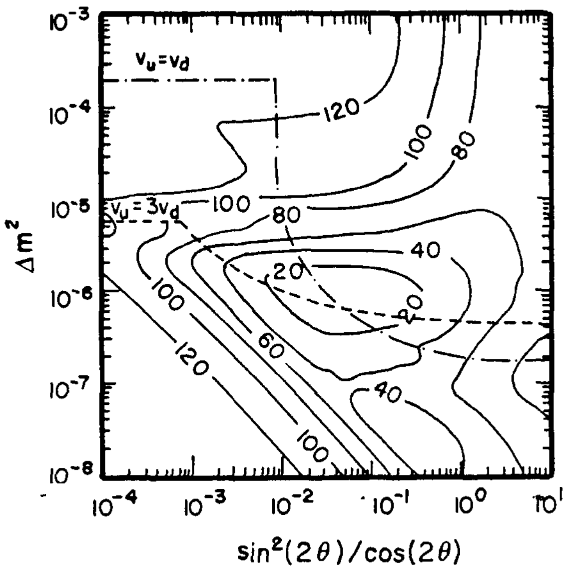

Fig. 3 Fig. 3.a shows the iso-SNU contours for the chlorine experiment in the recent 7.9 SNU solar model [29] (taken from Baltz and Weneser [28] who also included earth effects averaged over day/night, and seasonal variations). Shown in Fig. 3.b are gallium iso-SNU contours, for night. In the model of ref [31, 30] only the region left and below curve labelled $v_u = v_d$ or $v_u = 3v_d$ is allowed. Comparing 3.a and 3.b shows that even when the high energy neutrino count rate is high, the model favours a large suppression in the low energy neutrino flux.

To generate neutrino oscillations one introduces some mechanism of explicit flavour violation. Because in this model the Majoron carries $L = 1$ one needs a *double* breaking to generate the left handed Majorana neutrino mass. This fact, combined with the astrophysical limit [18], then implies a neutrino mass in the range where the MSW effect can substantially affect the solar neutrino flux for-very reasonable, yet *restricted* choices of the parameters. Out of the three neutrinos, to a very good approximation, only one acquires a mass, thus reducing to three the parameters describing neutrino oscillations: two mixing angles and one mass parameter, m_ν. One of the two angles specifies whether the oscillation channel is $\nu_e \rightarrow \nu_\mu$ or $\nu_e \rightarrow \nu_\tau$ and is not affected by matter, while the other is the angle shown in Fig. 3. The astrophysical constraint [18] combined with restrictions on the oscillation parameters that follow from conventional *laboratory* experiments then lead to the region of oscillation parameters shown in Fig. 3 ††. The *upper* bound on m_ν reflects a strong trend towards non-adiabaticity of high energy solar neutrino conversions, especially for the favoured case $v_u > v_d$. For $v_u = v_d$ the region of large m_ν and *small* θ is allowed. However, for $v_u = 3\ v_d$ this AD region disappears, the only allowed region being that of very small m_ν and relatively large θ. In this case we expect a depletion of low energy neutrinos up to a factor 7 or so *below* the standard solar model prediction, in sharp contrast to the AD case. In addition, in this region of parameters, regeneration of solar neutrinos at the earth could be important [28]. This large reduction in the expected pp and ^7Be neutrino fluxes is expected even in cases where the high-energy-neutrino-count rate in chlorine is fairly large, say $5SNU$ or so, thus stressing the importance of gallium experiments. There is however still a window for an adiabatic high energy neutrino conversion due to the uncertain determination of the astrophysical limit used above.

In the MSW effect the NC, being the same for all neutrinos, can be consistently ignored. However if gauge-singlet leptons e.g. right-handed neutrinos exist in nature, then as we saw in section 3, the resulting lepton mixing matrix is non-unitary, and the NC in the neutrino sector is non-trivial [2]. In this case the neutrinos pro-

†† Another test of this model arises from the fact that the Majoron belongs to an isodoublet and increases the width of the Z^0 by 80 MeV i.e. ***one half*** generation of neutrinos, which will be measurable at LEP.

duced in low-energy weak-decay processes can be *effectively* non-orthonormal and lead to a new type of ***resonant*** oscillation effect in matter that may happen even for massless neutrinos [24]. The massless oscillation effect is very different from the MSW effect. First it has an energy ***independent*** resonance condition

$$N_e^{net} = \delta h^2 N_n \tag{13}$$

where $\delta h << 1$ due to limits on weak universality violation, N_e^{net} is the net electron number density and N_n is the neutron density. This is to be contrasted with eq. (12) which holds in the MSW model. In matter, such as the sun, with a "normal" concentration of electrons relative to nucleons, the large CC contribution masks the NC effect making it impossible for the massless neutrino system to undergo resonant oscillation in the $\nu_e \rightarrow \nu_{\mu,\tau}$ channel of interest to the ***solar*** neutrino problem [24]. Oscillations such as $\nu_\mu \rightarrow \nu_\tau$ or $\bar{\nu}_\mu \rightarrow \bar{\nu}_\tau$ could be large, since there is no CC coherent scattering term for ν_μ or ν_τ on electrons. While the MSW effect can affect ***either*** neutrinos or anti-neutrinos, ***but not both,*** the massless resonant oscillation, if it occurs, will affect ***both*** $\nu's$ and $\bar{\nu}'s$. It could then substantially affect the propagation of $\nu - \bar{\nu}$ pairs emitted in the late phase of a supernova explosion, and the resulting neutrino signal.

5.2 Neutrinoless Double-Beta Decay

The existence of the neutrinoless double-beta decay $(A, Z - 2) \rightarrow (A, Z) + 2\ e^-$ would signal the violation of total lepton number in nature, as expected in many gauge theories. Moreover it would shed light on the nature of neutrinos themselves. This lepton number violating decay has usually been interpreted in terms of the ordinary (lepton number conserving) second order nuclear β decay process, $(A, Z - 2) \rightarrow (A, Z) + 2\ e^- + 2\ \nu$ by assuming that a ***virtual*** neutrino (or a combination of these) is exchanged between the decaying neutrons. The nuclei for which $\beta\beta$ decay is expected to be detectable are those where the transition to the intermediate nucleus via single β decay is kinematically forbidden or strongly suppressed, as in ^{48}Ca, ^{76}Ge, ^{82}Se, ^{100}Mo, ^{128}Te, ^{130}Te ***etc.*** Although the neutrinoless process is highly favoured by phase space, it proceeds only if the virtual neutrino is a Majorana particle. Since this requires non-zero Majorana neutrino mass this mechanism is sometimes called the "mass mechanism". It is the simplest mechanism that can engender $\beta\beta_{0\nu}$ in the standard model. The large phase space advantage of the $\beta\beta_{0\nu}$ process makes it a sensitive probe of the Majorana character of neutrinos. The relevant combination of parameters that governs the decay amplitude is

$$< m >= \sum_{\alpha=1}^{n} {K_{e\alpha}}^2 m_\alpha \tag{14}$$

where α runs over all the light neutrinos. The non-observation of $\beta\beta_{0\nu}$ e.g. in ^{76}Ge, ^{82}Se, ^{100}Mo, ^{128}Te, and ^{130}Te leads to the limit [32]

$$< m > \leq 1 - 5 \ eV \tag{15}$$

illustrating the uncertainties in the determination of nuclear matrix elements [33]. A better sensitivity is expected in the planned enriched germanium experiments.

The parameter < m > in eq. (14) may differ substantially from the neutrino mass inferred from tritium decay since in eq. (14) there can be a destructive interference between contributions of different neutrino types. To understand the simplest example where this can take place, consider the case where there is only one Dirac neutrino coupled to the charged current. A Dirac neutrino is equivalent to two Majorana neutrinos degenerate in mass in such a way that < m >= 0, as expected [11]. There are ***other*** ways to achieve the vanishing of < ***m*** >, for example whenever the mass matrix in the weak basis has a zero in the ee entry. Gauge models where either of these types of cancellations may take place have been considered [1].

If there are neutrinos heavier than about 10 MeV one needs to correct eq. (14) as discussed in ref [34]. In this case the relevant parameter that controls the decay amplitude is no longer universal, but depends on the specific nucleus considered.

In addition to the "mass mechanism" in gauge theories there are new ways to engender the $\beta\beta_{0\nu}$ decay process. One of these only involves the exchange of scalars and raises an important question of principle [22]: since it is possible to induce the $\beta\beta_{0\nu}$ process without virtual Majorana neutrinos being exchanged, perhaps we can do away altogether with the Majorana requirement. A simple but essentially rigorous proof showing that this is ***not*** so was given in ref [22]. There we showed that any generic "black box" mechanism inducing neutrinoless double-beta decay in gauge theories is bound to also produce a diagram generating a finite Majorana neutrino mass, so the relevant neutrino will, at some level, be a particle of Majorana type [22]. This result is completely general, and has been subsequently elaborated and made quantitative by several authors. If neutrinoless double-beta decay is induced by the exchange of vector currents one may set a lower limit on one of the neutrino masses $m_\nu \gtrsim 1\ eV \times (10^{24} yr/\tau_{Ge})^{1/2}$ [35]. This bound does not however apply in general. For example it fails if the neutrinoless double-beta decay is dominated by scalar exchange. A model of this type has recently been suggested in ref [36].

Many other mechanisms, involving right handed currents, Quasi-Dirac neutrinos and supersymmetric particles could also induce the $\beta\beta_{0\nu}$ [1]. In addition gauge theories may lead to other, genuinely new, possible varieties of neutrinoless double-beta decay where light scalars are ***emitted.*** For example, in theories that contain a Majbron $\beta\beta_{0\nu}$ decay may proceed with Majoron emission. The dominant Majoron emission mode depends on the lepton number carried by the Major-on. If the Majoron belongs to an isotriplet coupled to two lepton doublets, it carries $|\Delta L| = 2$ and the dominant scalar emission mode is

$$(A, Z-2) \rightarrow (A, Z) + 2\ e^- + J \qquad (16)$$

while in the case that the Majoron has $|\Delta L| = 1$ the favored scalar emission process involves double Majoron emission [37]

$$(A, Z-2) \rightarrow (A, Z) + 2\ e^- + 2\ J \qquad (17)$$

An interesting example of this situation is provided by the model discussed in ref [31, 30]. It is not yet clear whether this process is large enough to be observable in this model [38]

Since the Majoron is very weakly coupled to matter it will not be detected except through its indirect effect on the β spectrum, characteristically distinct for each one of the double-beta decay processes.

Another process quite similar in structure to neutrinoless double-beta decay is the neutrino-anti-neutrino-oscillation process suggested in ref [3]. It involves the propagation of real rather than the ***virtual*** neutrinos of the $\beta\beta_{0\nu}$ process. Oscillation probabilities are however too small to observe due to helicity suppression present in the standard model.

5.3 Flavour and CP Violation

Lepton flavour violation can be induced at one-loop due to neutral heavy lepton admixture in the weak currents and lead to flavour violating effects at low energies [7, 8]. These effects include processes such as $\mu \rightarrow e + \gamma$, $\mu \rightarrow 3e$, μ-e conversion in nuclei, ***etc.*** These processes can occur even when the physical neutrinos are massless [7]. In Fig. 4, taken from ref [8], we illustrate the branching ratios for the flavour violating process $\mu \rightarrow e + \gamma$ produced by non-zero neutrino masses in the seesaw model and in μ models discussed in sections 2.1 and 2.3. The non observation of $\mu \rightarrow e + \gamma$ places strong constraints on the parameters of the models discussed in sections 2.2 and 2.3. Note in contrast that the seesaw model expectation is below detectability.

In addition we have the possibility of flavour violating decays of the Z^0 [7] such as $Z^0 \rightarrow e + \overline{\tau}$, $Z^0 \rightarrow \mu + \overline{\tau}$ which might lead to observable signatures at high energies in Z^0 factories such as LEP [7].

Another type of flavour violating decays involves Majoron emission. This is possible, for example, when the Majoron is the supersymmetric partner of the neutrino as in the model described in ref [31, 30]. There are then various spectral distortions in weak decay processes due to Majoron emission in processes such as $\mu \rightarrow e + J$, $\mu \rightarrow e + J + J$, $\tau \rightarrow e + J + J$, ***etc.*** The accuracy of the present determination of the Michel parameter places stringent constraints on the branching ratios for the double Majoron emission processes, while single emission could be barely measurable in the case of μ decay.

As discussed above leptonic CP violating effects may also be present even when

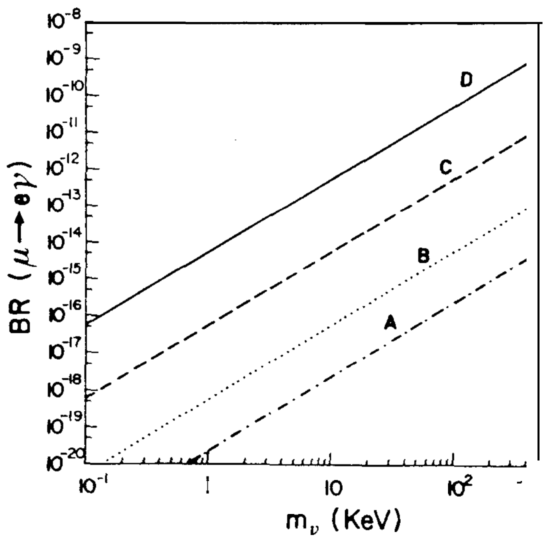

Fig. 4 : Large $\mu \to e+\gamma$ decay branching ratio estimated in the μ model assuming a typical NHL mass of 20 to $60 GeV$ and different μ values of $50 GeV$(A), $10 GeV$(B), $1 GeV$(C) and O.lGeV(D). For large μ values the seesaw-model prediction is recovered.

neutrinos are massless and thus need not be suppressed by the smallness of neutrino masses. This leads to CP asymmetries in weak processes such as the decays of the intermediate vector bosons. Another promising possibility is the existence of a sizable electric dipole moment for the electron [12].

5.4 Neutrino Decay

When a Majoron exists the invisible neutrino decay lifetime due to eq. (11) can be made sufficiently short so that large neutrino mass values (consistent with experiment) are fully consistent with astrophysics and cosmology. Models where this is possible have been discussed in ref [20, 8] and a typical lifetime versus mass relationship is illustrated in Fig. 5, taken from ref [8]. It is worth noting that a recent experiment [39] reports a finite neutrino mass that violates eq. (10). The importance that such an observation would have justifies the effort necessary to obtain an independent experimental confirmation of this result.

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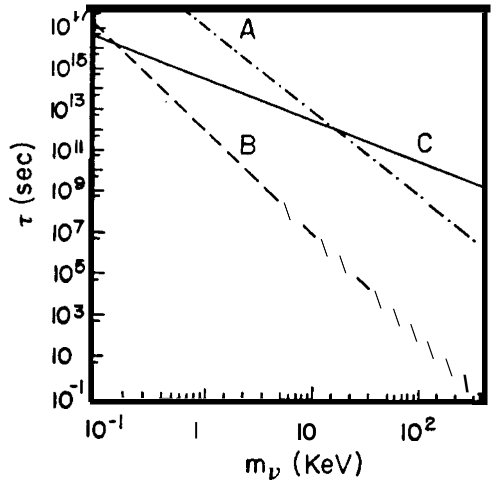

Fig. 5 : Neutrino decay lifetime estimated in two versions of the Majoron model where lepton number is broken spontaneously in the seesaw model (line A, M_R = ***50GeV)*** or in the μ model (line B, $M = 10GeV$). Line C denotes the cosmological constraint. *All* neutrino mass values consistent with experiment are cosmologically allowed in the μ model and neutrino masses larger than $O(20\ KeV)$ are allowed in the seesaw model.

the hospitality extended to me at Montreal.

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