受入
92-11-149
高エ研図書室

CERN-TH.6626/92
August 1992

PHYSICS AT NEW ACCELERATORS: LOOKING BEYOND THE STANDARD MODEL

J. W. F. Valle[1] *
Theoretical Physics Division, CERN
CH-1211 Geneve 23, Switzerland

ABSTRACT

Neutrino masses can have important implications in nuclear and particle physics, astrophysics and cosmology. Apart from the effects related to solar neutrinos, neutrino oscillations, dark matter, beta and double beta decays, massive neutrinos can also produce signals in the new accelerators. Here I focus on μ and τ number violating processes, very promising for muon and tau factories, as well as on the signatures associated with spontaneously broken R parity supersymmetry and neutral heavy leptons. These include the possibility of high rates for single chargino and neutralino production at LEP, LHC/SSC, as well as new signatures involving invisibly decaying Higgs bosons.

Invited Talk at the XV International Conference on Neutrino Physics and Astrophysics, Neutrino 92, Granada, Spain, June, 1992. To appear in the Proceedings.

CERN-TH.6626/92
August 1992

[1] Permanent address: Instituto de Física Corpuscular - C.S.I.C.
Dept. de Fisica Teòrica, Universitat de València
46100 Burjassot, València, SPAIN

---

*Bitnet VALLE@EVALUN11 - Decnet 16444::VALLE

Physics at New Accelerators: Looking Beyond the Standard Model

J. W. F. Valle[a] *

[a]Instituto de Física Corpuscular - IFIC/C.S.I.C.
Departament de Física Teòrica, Universitat de València
46100 Burjassot, València, SPAIN
Bitnet VALLE@EVALUN11 - Decnet 16444::VALLE

Neutrino masses can have important implications in nuclear and particle physics, astrophysics and cosmology. Apart from the effects related to solar neutrinos, neutrino oscillations, dark matter, beta and double beta decays, massive neutrinos can also produce signals in the new accelerators. Here I focus on μ and τ number violating processes, very promising for muon and tau factories, as well as on the signatures associated with spontaneously broken R parity supersymmetry and neutral heavy leptons. These include the possibility of high rates for single chargino and neutralino production at LEP, LHC/SSC, as well as new signatures involving invisibly decaying Higgs bosons.

1. INTRODUCTION

Our present understanding of particle physics is based upon the standard $SU(2) \otimes U(1)$ model. Although extremely successful wherever it has been tested, the standard model leaves several open puzzles that motivate further extensions. Perhaps the most fundamental problem in particle physics today is that of understanding what lies behind the mechanism of mass generation. By far the simplest possibility is to rely on the Higgs mechanism which, as is well known, implies the existence of a fundamental scalar boson [1]. If this is the case it is widely believed that some stabilizing principle - e.g. supersymmetry (SUSY) - should be operative at the electroweak scale in order to explain the stability of this scale against quantum corrections associated with physics at superhigh energies. At present there is only circumstantial evidence, provided by the joining of the three gauge coupling strengths at a common $\sim 10^{16}\,GeV$ energy, provided SUSY sets in at $M_{SUSY} \sim 10^3\,GeV$ [2]. Despite the lack of direct experimental support, the above argument has been taken as a strong motivation to carry out SUSY searches at higher energies. Unfortunately there is no clue as to how SUSY is realized. The most popular *ansatz* - called the minimal supersymmetric standard model (MSSM) [3] - realizes SUSY in the presence of a discrete R parity (R_p) symmetry. In this case neutrinos are massless, as in the standard model. However, this choice has no firm theoretical basis and it is of great interest to investigate theories that avoid it [4]. Here we focus on the case of spontaneous R_p breaking in the $SU(2) \otimes U(1)$ theory. The viability of this possibility has been recently demonstrated [5]. The breaking of R-parity is driven by right-handed *isosinglet* sneutrino vacuum expectation values (VEVS) [6], so that the associated Goldstone boson (majoron) is mostly singlet and as a result the Z does not decay by majoron emission, in agreement with LEP observations [7].

There is a wealth of neutrino-mass-related phenomena covering a broad range of energies and of experimental situations. Apart from providing a natural explanation of the solar neutrino data through the MSW effect, and accounting for the hot dark matter component suggested by COBE data [8], neutrino masses may be seen in neutrino oscillations, in observable rates for neutrinoless double beta decays (with and without majoron emission), and (if they are large enough) also as distortions in beta decay spectra. A large number of related processes can also manifest themselves at muon and tau factories and at high energy e^+e^- collisions (e.g. LEP). There are also good prospects to observe some of these signatures at the upcoming hadron supercolliders LHC/SSC. In this talk I will give some examples of these processes ranging from μ and τ number

---

*Work supported by CICYT

violating decays, up the high energy processes associated with the single production of SUSY fermions or neutral heavy leptons (NHLS) at LEP or at a future hadron supercollider, LHC/SSC. I also note in this context the rather peculiar possibility that the Higgs boson may decay dominantly by two majoron emission, leading to missing energy events. New Higgs search strategies will be needed to cover this possibility. All of the above effects related to nonstandard neutrino properties may well be accessible to experiment.

2. PRELIMINARIES

Before describing the signatures of interest for the high energy accelerators it is useful to recall the main constraints on neutrino masses that follow from observation. This will be useful in constraining the processes to be discussed in the next sessions. The present neutrino mass laboratory bounds may be summarized as [9]

$$m_{\nu_e} \lesssim 9eV, \;\; m_{\nu_\mu} \lesssim 500keV, \;\; m_{\nu_\tau} \lesssim 31MeV \quad (1)$$

In addition there is a cosmological limit that follows from considerations related to the abundance of relic neutrinos [10]

$$\sum_i m_{\nu_i} \lesssim 50\, eV \quad (2)$$

However the limit in eq. (2) only holds if neutrinos are stable. Indeed, there are many ways [4] to make neutrinos sufficiently unstable in such a way as to avoid the limit in eq. (2) The models rely on the existence of fast neutrino decays involving majoron emission [11–14], *e.g.*,

$$\nu_\tau \to \nu_\mu + J \quad (3)$$

The resulting lifetime can be made sufficiently short so that neutrino mass values as large as eq. (1) are fully consistent with astrophysics and cosmology. Examples of seesaw type models where this is possible have been discussed in ref. [12, 13]. Another more recent example is provided by the spontaneously broken R parity model [14]. Curve C in Fig. 1 shows the minimum estimated ν_τ lifetime consistent with observation in this model, plotted versus the ν_τ mass. The ν_τ decay lifetime required in order to efficiently suppress the relic ν_τ contribution is shown as the solid straight line in Fig. 1. Clearly the decay lifetimes can be shorter than required. Moreover, since these decays are ***invisible***, they are consistent with all astrophysical observations [10]. An additional constraint, shown as the dashed line in Fig. 1 [15], may be derived by demanding that the universe has become matter-dominated by a redshift of 1000 at the latest, so that fluctuations have grown sufficiently by today. Again, one sees that shorter lifetimes are possible. However, this lifetime limit is less reliable than the one derived from the critical density. It is worth noting in this context, the hints, albeit controversial, from recent beta decay experiments in favour of a 17 keV neutrino [16] which would require a decay of the type in eq. (3). The importance that such an observation would have justifies the effort necessary to clear the controversy. In addition to the above, there are limits on neutrino mass and mixing that follow from the nonobservation of neutrino oscillations [9]. Further improvements on ν_μ - ν_τ oscillation limits are expected from new experiments such as CHORUS and NOMAD at CERN and the proposed P803 at Fermilab. These experiments are especially relevant in relation to

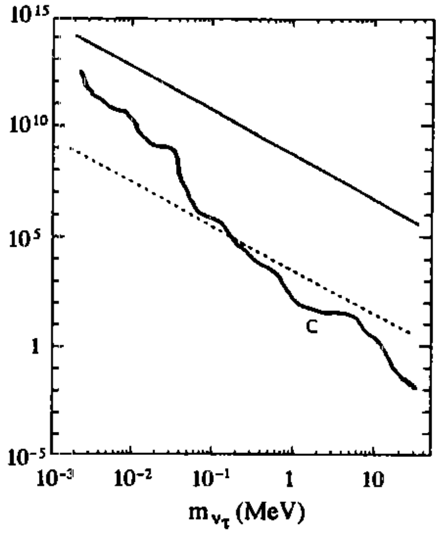

Figure 1. ν_τ lifetime versus observational limits

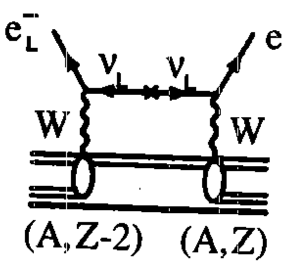

Figure 2. Mass mechanism for $\beta\beta_{0\nu}$ decay.

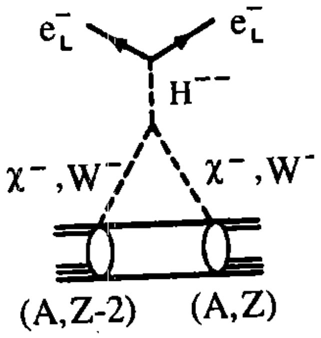

Figure 3. Scalar-induced $\beta\beta_{0\nu}$ decay.

the hints in favour of the existence of a hot dark matter component in the universe [8]. In this connection, I wish to emphasize the theoretical interest in searching for neutrino oscillations also in the ν_e to ν_τ channel, as evident in the model discussed in ref. [17].

A process of great interest in neutrino physics is neutrinoless double beta decay, $(A, Z-2) \to (A, Z) + 2\,e^-$, so far never observed. Its existence would signal the violation of total lepton number in nature, as expected in many gauge theories. The standard way to induce this decay is via neutrino exchange, as shown in Fig. 2. Although highly favoured by phase space over the usual 2ν mode, the neutrinoless process proceeds only if the virtual neutrino is a Majorana particle. The decay amplitude is proportional to

$$\langle m \rangle = \sum_\alpha K_{e\alpha}{}^2 m_\alpha \tag{4}$$

where α runs over the light neutrinos. The parameter $\langle m \rangle$ above may differ substantially from the neutrino mass inferred from beta decay since in eq. (4) there can be a destructive interference between contributions of different neutrino types. The simplest symmetry implying such cancellations is the lepton number symmetry inherently assigned to a Dirac neutrino mass term, which can be decomposed as two Majorana neutrinos degenerate in mass in such a way that $\langle m \rangle$ automatically vanishes [18].

The non-observation of $\beta\beta_{0\nu}$ leads to the limit

$$\langle m \rangle \lesssim 1 - 3\ eV \tag{5}$$

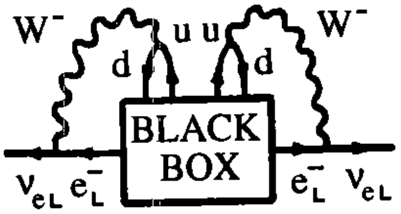

Figure 4. $\beta\beta_{0\nu}$ decay and Majorana neutrinos.

depending on nuclear matrix elements [19]. A better sensitivity is expected from the enriched germanium experiments.

In addition to the above "mass mechanism" in gauge theories there are other ways to engender the $\beta\beta_{0\nu}$ decay process, for example, through the exchange of scalars, as illustrated in Fig. 3. This raises an important question regarding the relationship that the $\beta\beta_{0\nu}$ decay process bears to the Majorana nature of the neutrinos [20]. A simple but essentially rigorous proof showing that, in a gauge theory, whatever the origin of neutrinoless double beta decay is, it requires neutrinos to be Majorana particles, is illustrated in Fig. 4. Any generic "black box" mechanism inducing neutrinoless double beta decay in gauge theories is bound to also produce a diagram generating a nonzero Majorana neutrino mass, so the relevant neutrino will, at some level, be a Majorana particle [20].

Gauge theories may lead to new varieties of neutrinoless double beta decay involving the *emission* of light scalars, such as the majoron [21] [†]

$$(A, Z-2) \rightarrow (A, Z) + 2\, e^- + J\,. \qquad (6)$$

Since such light scalars are very weakly coupled to matter, their emission would only be detected through their effect on the β spectrum.

The simplest model leading to sizeable majoron emission in $\beta\beta$ decays involving an isotriplet majoron [22] is no longer phenomenologically viable, since it leads to a new invisible decay mode for the neutral gauge boson by the emission of light scalars,

$$Z \rightarrow \rho + J, \qquad (7)$$

now ruled out by LEP measurements of the invisible Z width [7]. However it has been recently shown that a large majoron-neutrino coupling leading to observable majoron emission in neutrinoless double beta decay can easily be reconciled with the LEP results in models where the majoron is an isosinglet and lepton number is broken at a low scale [23]. This is specially interesting now in view of the puzzling features recently hinted in some double beta decay experiments, which might indicate the presence of very light scalars [24].

In addition to limits, observation also provides some positive hints for neutrino masses. Perhaps the most significant are the solar neutrino data. It is interesting to remark that the recent results of the GALLEX experiment on low energy pp neutrinos reported here do not really "eliminate" the solar neutrino puzzle, in view of the persisting deficit of high energy neutrinos seen in Kamiokande and Homestake. The astrophysical explanation of these data would require not only too large a drop in the temperature of the solar core, but also would predict wrongly the relative degree of suppression observed in these two experiments [26]. The most attractive way to interpret the existing data is via the MSW effect which requires neutrino masses $\sim 10^{-3}$ eV. The region of parameters allowed by present experiments is illustrated in Fig. 5 [25].

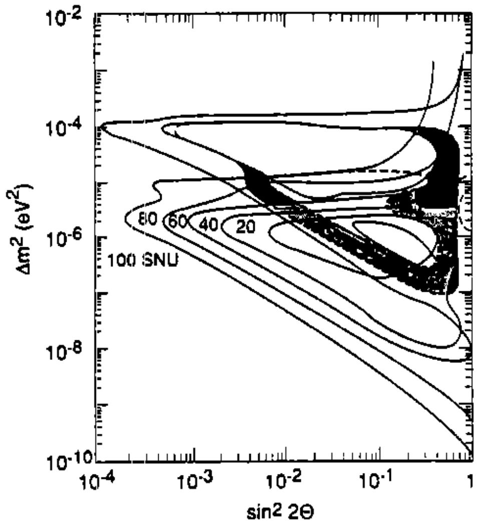

Figure 5. Allowed MSW oscillation parameters

3. SUSY WITHOUT R PARITY

In its most conventional *ansatz* SUSY is realized assuming the existence of a discrete R_p symmetry under which all standard model particles are even while their partners are odd. This leads to important selection rules *i.e.*, that SUSY particles must always be produced in pairs and the lightest must be stable. Taking this for granted, hadron colliders $Sp\bar{p}S$ and the Fermilab Tevatron constrain strongly interacting SUSY particle (squarks and gluino) masses to be in excess of 120 GeV or so. The limits on electrically charged weakly interacting SUSY particles are set by LEP and restrict the masses of sleptons and charginos (charged partner of the $SU(2) \otimes U(1)$ gauge and Higgs bosons) to be in excess of 45 GeV. Similarly, LEP limits on electrically neutral weakly interacting SUSY particles (sneutrinos and neutralinos) are somewhat weaker. In any case it is clear by now that present high energy facilities

---

[†] A related light scalar boson ρ should also be emitted.

have almost reached their limits in finding effects of SUSY particles. Here I stress that one can still perform meaningful searches for broken R parity SUSY physics at present facilities.

There are two ways to break R_p: explicitly [27] and spontaneously [6, 28, 29]. The second provides a more systematic way to include R parity violating effects. Moreover it automatically respects low energy baryon number conservation and evades restrictions based on cosmological baryogenesis arguments [30–32] inasmuch as the breaking of R parity sets in only as an electroweak scale phenomenon. Here also there are two cases to consider, depending on whether lepton number is part of the gauge symmetry or not. In the first case [28, 29] there is a Z' gauge boson which acquires mass via the Higgs mechanism at a scale related to that which characterizes R-parity violation. In the second case [6] we can consider just the simplest $SU(2) \otimes U(1)$ gauge structure. The breaking of R-parity is driven by

$$v_R = \langle \tilde{\nu^c}_\tau \rangle \lesssim 1\, Tev, \quad v_L = \langle \tilde{\nu}_{L\tau} \rangle \lesssim 100\, MeV \tag{8}$$

where $v_R \gg v_L$ is an *isosinglet* vacuum expectation value (VEV). As a result the associated majoron is mostly singlet and as a result the decay in eq. (7) is highly suppressed, in agreement with LEP observations [7]. A superpotential leading to these features can also be found in ref. [6]. Although it conserves lepton number symmetry, the presence of new $SU(2) \otimes U(1)$ singlet superfields such as ν^c can drive the spontaneous violation of R parity and L at or below the electroweak scale [5, 6]. The hierarchy between v_R and v_L is needed in order to adequately suppress the stellar energy loss that would result from majoron emission processes *e.g.*, the Compton-like reaction $\gamma + e \to e + J$. This hierarchy follows naturally, since $v_L = \langle \tilde{\nu}_{L\tau} \rangle$ is related to a new Yukawa coupling h_ν and vanishes as $h_\nu \to 0$ [5] [6].

In spontaneously broken R parity models L is violated, so neutrinos get masses, from mixing with the heavy neutral R-odd fermions. The basic feature of the model is that the heavier neutrino is the ν_τ and its mass scales as

$$m_{\nu_\tau} \propto v_R^2 \tag{9}$$

whereas the ν_μ mass scales only as

$$m_{\nu_\mu} \propto v_L^2 \tag{10}$$

and ν_e is massless in the tree approximation. Thus the ν_e -ν_μ mass difference is therefore very small on the scale of the ν_τ mass. Choosing $v_R = 1\, TeV$ and $v_L = 100\, MeV$ one gets a ν_μ - ν_τ mass ratio

$$\frac{m_{\nu_\mu}}{m_{\nu_\tau}} = \mathcal{O}\,(10^{-8}) \tag{11}$$

Thus for a ν_τ mass in the $10\, KeV - 1\, MeV$ range (consistent with cosmological limits, since the heavy ν_τ decays via majoron emission) one gets $m_{\nu_\mu} \sim 10^{-3\pm1}\, eV$, as needed for the MSW effect. On the other hand the ν_e -ν_μ mixing angle can easily be chosen to correspond to one of the small or large angle solutions left by GALLEX and shown in Fig. 5. In addition, in these SUSY models there are R_p violating couplings in the electroweak charged and neutral currents, as discussed in ref. [33]. Their magnitude is restricted by observational constraints from high energy colliders LEP/$Sp\bar{p}S$/Tevatron as well as from those related to flavour and/or total-lepton-number-violating processes, e.g. the non-observation of neutrino oscillations and neutrinoless double β decay, the failure to observe anomalous peaks on the energy distribution of the electrons and muons coming from decays such as $\pi, K \to e\nu$ and $\pi, K \to \mu\nu$, etc. These constraints have been systematically studied in ref. [33] and the result is that these couplings can be relatively strong, as discussed in ref. [33]. This leads to the possibility of large R_p violating effects, involving:

1. *single* chargino production in Z decays [14, 29],

   $$Z \to \chi^\pm \tau^\pm\,. \tag{12}$$

   where $\chi^\pm$ denotes a chargino. These decays could originate recognizable events at LEP.

2. *single* chargino (or neutralino) production in hadron colliders [33], associated to a τ or ν_τ

   $$pp \to \chi^+ \tau^- X, \quad pp \to \chi^- \nu_\tau X, etc \tag{13}$$

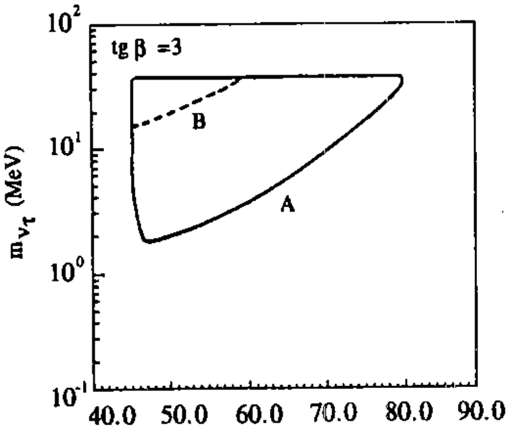

Figure 6. Attainable $Z \to \chi\tau$ branching ratios

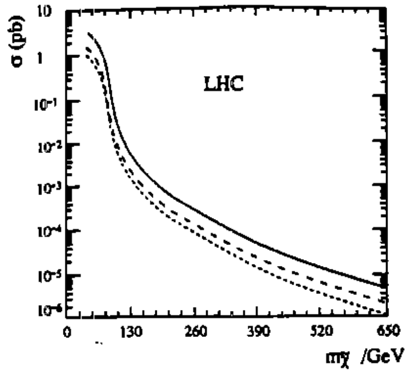

Figure 7. Single chargino production at LHC

3. muon and τ decays involving majoron emission [34], *e.g.*,

   $$\mu \to e + J \ , \tau \to \mu + J \ , \tau \to e + J \ . \quad (14)$$

   These could well be observed at muon or tau factories.

The attainable branching ratios for $Z \to \chi\tau$ have been evaluated in ref. [14, 29]. Fig. 6 shows that the corresponding branching ratio can well be larger than 10^{-5} and therefore within the sensitivity of the LEP1 experiments [29]. The figure corresponds to $\tan\beta = 3$ where $\tan\beta = \frac{v_u}{v_d}$ is the ratio of the two SUSY Higgs doublet VEVS, expected to lie somewhere between 1 and the top-bottom quark mass ratio. Clearly the $Z \to \chi\tau$ decay is a measure of the ν_τ mass. Similarly, neutralinos can also be singly produced from $Z \to \chi^0\nu_\tau$. Moreover, the lightest of them, χ^0, is unstable and is therefore not necessarily an origin of events with missing energy. For example in the process $Z \to \chi^0\chi^0$ one χ^0 can decay visibly as $\chi^0 \to \nu f \bar{f}$ while the other invisibly, as $\chi^0 \to \nu_\tau J$, leading to enhanced rates for *zen* events, which may reach values as large as $BR(zen) \gtrsim 10^{-4}$ [14, 29]. Although such events are expected in the MSSM, their origin is quite different in that case, since there the missing energy is carried by stable "photinos" while in the spontaneously broken R parity models it is always carried by neutrinos or majorons [29, 14].

The single production of SUSY fermions would also take place at hadron colliders, LHC/SSC. The maximum rates consistent with experimental limits have been worked out in ref. [33]. As an example, we show in Fig. 7 the allowed cross sections for single chargino production at LHC. The three curves correspond the channels $\chi\tau X$ (upper), $\chi^+\nu_\tau X$ (middle) and $\chi^-\nu_\tau X$ (lower), respectively, for the Drell-Yan mode of production [33]. Single neutralino production gives a specially interesting signature involving like-sign taus [33].

Now turning to the processes in eq. (14) we note that, since the majoron is a massless weakly interacting pseudoscalar Goldstone boson, its effect would be manifest only insofar as its emission would substantially affect the spectra of the decay-produced leptons, leading to bumps at half of the parent mass in its rest frame. The processes in eq. (14) could have sizeable branching ratios, within present experimental sensitivities [34]. For example the muon may decay by majoron emission with rates similar to the the present TRIUMF limit. Here I focus on the case of tau de-

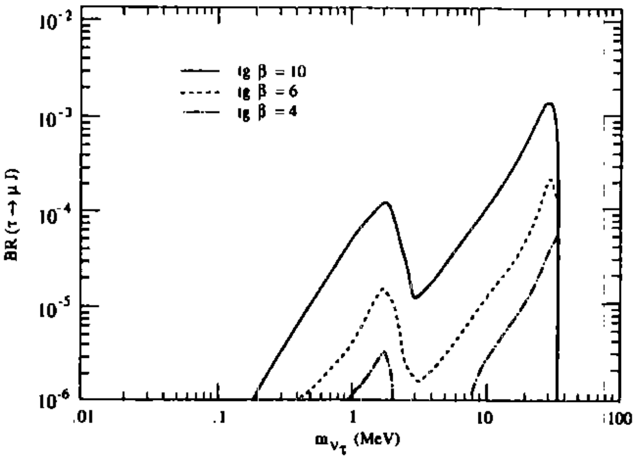

Figure 8. Majoron emitting τ decay branching ratios

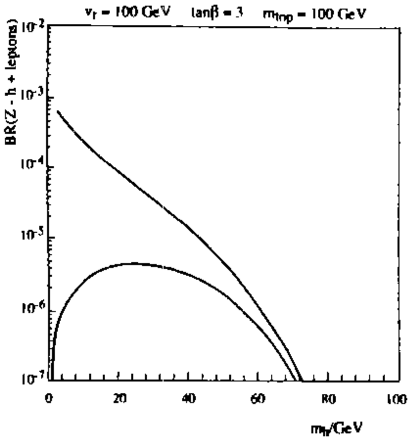

Figure 9. $BR(Z \to h + leptons)$ in broken R_p SUSY

cays, of obvious relevance for a tau-charm factory (TCF). In Fig. 8 we display the results for $\tau \to \mu + J$ [34] and compare with the present ARGUS limit [35] $BR(\tau \to \mu + J) \leq 5.8 \times 10^{-3}$. All points under the solid contour denote branching ratios that can be reached for $\tan\beta = 10$ without violating any of the relevant observational constraints. Those inside the dashed contour correspond to $\tan\beta = 6$ and those inside the dot-dashed contour correspond to $\tan\beta = 4$. For the case of the decay $\tau \to e + J$ the attainable branching ratios vary in the range $10^{-6} \leq BR(\tau \to e + J) \leq few \times 10^{-4}$, as β varies between $\tan\beta = 10$ and $\tan\beta = 40$ [34]. This is to be contrasted with the present ARGUS limit $BR(\tau \to e + J) \leq 3.2 \times 10^{-3}$. Upcoming facilities such as a TCF will be ideally suited to probe for the possible existence of these new processes.

4. INVISIBLY DECAYING HIGGS

One interesting aspect of the above SUSY model with spontaneous R_p violation is the novel possibility that the main Higgs decay channel is likely to be invisible [36]. At LEP1 the main Higgs production mechanism in this model is the conventional Bjorken process as in the standard model. However, the corresponding production rates can be weaker than in the standard model (SM), especially in the low mass region. This will substantially weaken the Higgs boson mass limits derived from LEP1 [37]. However one finds that the dominant Higgs decay channel is "invisible", over most of the mass range accessible at LEP1, leading to events with large missing energy carried by majorons. Fig. 9 shows the branching ratios for the production of the lightest CP-even Higgs (h) plus leptons from Z decay. The upper line is the standard model prediction and the lower one is the maximum branching ratio possible for $\tan\beta = 3$ and $m_{top} = 100 GeV$. All points underneath this curve are possible in the model, and correspond to the unknown parameters being randomly varied over reasonable ranges, consistent with observational constraints and with the minimization of the scalar potential of the theory. For larger masses, accessible at LEP2, LHC/SSC, there is no necessary suppression in the production rate and the favored Higgs decays can still be invisible. This very peculiar possibility should be taken into account in the planning of Higgs boson search strategies at future colliders, and could be particularly promising for intermediate mass Higgs search at hadron colliders [38].

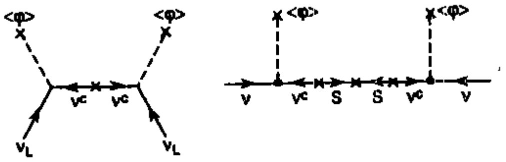

Figure 10. How to generate small neutrino masses.

5. NEUTRAL HEAVY LEPTONS

It is not possible in the standard $SU(2) \otimes U(1)$ electroweak model to write any gauge invariant interactions that can lead to nonzero neutrino mass [4]. Thus all lepton flavours as well as total lepton number, L are exactly conserved. Since we have strong bounds on neutrino masses, eq. (1), it is reasonable to expect that, if nonzero, they arise in a different way from those of the charged fermions of the standard model. If neutrinos are Majorana particles the relative smallness of their masses may be understood through seesaw-like mechanisms, as illustrated in Fig. 10. The required right handed (RH) neutrinos arise in many theories beyond the standard model in order to realize a larger symmetry e.g. left-right, or grand unified symmetries [4]. Alternatively Majorana neutrino masses could arise only as the result of radiative corrections and therefore be naturally smaller than the charged fermion masses.

Here I will focus on the case of seesaw-inspired models at the electroweak scale. For simplicity, I assume the simplest $SU(2) \otimes U(1)$ gauge and Higgs structure and add to the standard model a number of isosinglet NHLS. These may have $SU(2) \otimes U(1)$ invariant mass terms breaking L symmetry. The precise nature of this breaking is important in determining many aspects of lepton physics. In the simplest seesaw model one adds one isosinglet right handed neutrino ν_i^c sequentially, giving it a bare, $SU(2) \otimes U(1)$ invariant, but L violating, Majorana mass term. The physical neutrino masses are determined from the 6×6 neutrino mass matrix (in the basis (ν, ν^c)) [39]

$$\begin{pmatrix} 0 & D \\ D^T & M_R \end{pmatrix} \tag{15}$$

where the matrix D_{ij} is the Dirac mass term for the three RH neutrinos, and M_R is an isosinglet Majorana mass matrix. The light neutrino mass matrix that results after diagonalizing out the heavy fields is $M_L{}^{seesaw} = -DM_R^{-1}D^T$.

Alternatively, one can have an L-conserving Model, where we include 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. This restriction leads to the following form for the neutral mass matrix (in the basis ν, ν^c, S)[‡]

$$\begin{pmatrix} 0 & D & 0 \\ D^T & 0 & M \\ 0 & M^T & 0 \end{pmatrix} \tag{16}$$

where the Dirac mass term coupling the $SU(2) \otimes U(1)$ doublets with the singlets ν_i^c is described by the matrix D proportional to the standard Higgs VEV $\langle \phi \rangle$ responsible for electroweak breaking and fermion masses, while the other Dirac mass matrix M couples together two electroweak singlets ν^c and S. Although these M-coefficients can be large compared with the characteristic scale for the elements D_{ij} they need not be very large. The three light neutrinos are massless *Weyl* neutrinos while the other 6 neutral 2-component leptons combine to form 3 heavy *Dirac* fermions. *Despite the fact that physical neutrinos are strictly massless* individual leptonic flavours are violated in this model [42–44]. Similarly for CP symmetry [45, 46].

A variant of the previous model may be obtained by introducing total lepton number viola-

---

[‡] A similar mass matrix for one generation has been suggested in several theoretical frameworks such as in superstring inspired models [40, 41]. In this case the zero entries naturally arise due to the lack of Higgs fields that could provide the usual Majorana mass terms needed in the seesaw mechanism.

tion through a non-zero Majorana mass μ_{ij} for the S_i in eq. (16) [13] as follows

$$\begin{pmatrix} 0 & D & 0 \\ D^T & 0 & M \\ 0 & M^T & \mu \end{pmatrix} \tag{17}$$

This leads to small neutrino masses determined from $M_L = DM^{-1}\mu M^{T^{-1}}D^T$. In this "$\mu$-model" 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 [18]. 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. In addition, if μ is generated from a spontaneous violation of lepton number at a very low scale $\mu \sim \langle\sigma\rangle \lesssim 50$ keV, then we also expect to have substantial majoron emission effects in 0ν double beta decay [23].

The pattern of electroweak currents in the above NHL models is much more complex than the familiar quark mixing. The form of the charged and neutral currents for the above models was given in ref. [45–48]. As an example, I mention that if the physical mass-eigenstate neutrinos are Majorana type, CP is violated even in a 2 generation world [49]. The most important point is that these neutral heavy leptons also couple, via mixing, in the electroweak currents, similarly to charginos and neutralinos discussed above. This allows for many novel possibilities. For example, CP violation (just with two generations of leptons) may occur even when the physical neutrinos are strictly massless [45, 46]! Here I will focus on the possibility of enhanced flavour violation.

The prime example of flavour violating decays is $\mu \to e + \gamma$. A typical contribution to this decay comes W exchange. The branching ratio for this process in the above models is

$$BR(\mu \to e + \gamma) = \frac{3\alpha}{32\pi}|K_{\mu\alpha}F_\alpha K^\dagger_{\alpha e}|^2 \tag{18}$$

where F is a form factor depending on the mass of the neutrinos and the K-factors denote the appropriate mixing coefficients that describe the charged current. In the simplest case where we consider Dirac neutrino masses (no NHLS) there is an almost exact cancellation [§] amongst the various massive neutrino contributing to eq. (18) so that F scales as m_ν^2 and consequently the $BR(\mu \to e + \gamma)$ as m_ν^4, too small therefore due to eq. (2) which applies in this case.

The presence of NHLS destroys the above cancellation so that one can have $F \sim 1$. However, even in this case the $BR(\mu \to e + \gamma)$ scales as m_ν^2 due to the fact that the NHL admixture in the charged current is expected to scale as m_ν. Thus, although substantially enhanced with respect to the simplest Dirac neutrino case above, we are still far out of experimental reach.

Here I want to stress that one can have models where *both* F and the mixing coefficients are unrelated to m_ν. This is the case in the model of eq. (16) above. In this case $BR(\mu \to e + \gamma)$ can be large since we avoid altogether any limits on neutrino masses. Similarly, for the model described by eq. (17) with small μ. The relevant limiting factor is the validity of universality in the charged weak current [4]. Consequently all effects due to the NHL admixture in the electroweak currents are enhanced. These include flavour violation [42, 44], CP violation [45, 46], NHL production at LEP [47, 48] or at hadron supercolliders LHC/SSC, etc.

As examples of lepton flavour violating decays we consider $\mu \to e + \gamma$. The maximum estimated branching ratio which is consistent with universality limits is shown as the dashed line in Fig. 11 [44]. We have used in our estimates the numbers in ref. [50]. We see that for M bigger than 10 GeV this branching ratio will exceed the present experimental limit, and therefore restrict the allowed branching ratio for $BR(\mu \to 3e)$, as shown by the solid line. In contrast, barring the presence of additional Higgs bosons, the simplest seesaw model expectations for these processes are well below detectability.

The τ lepton would also exhibit similar flavour violating effects, e.g.

$$\tau \to \pi^0 l_i \tag{19}$$

$$\tau \to l_i \gamma\,, \tag{20}$$

[§] This "GIM-mechanism" follows from the unitarity of the leptonic mixing matrix.

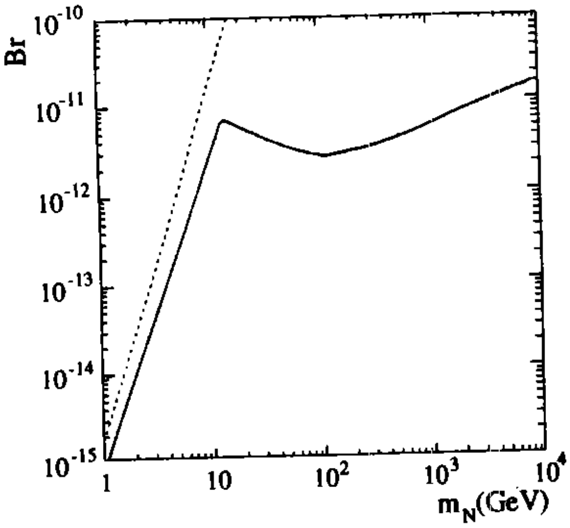

Figure 11. $BR(\mu \to 3e)$ and $BR(\mu \to e + \gamma)$ consistent with weak universality.

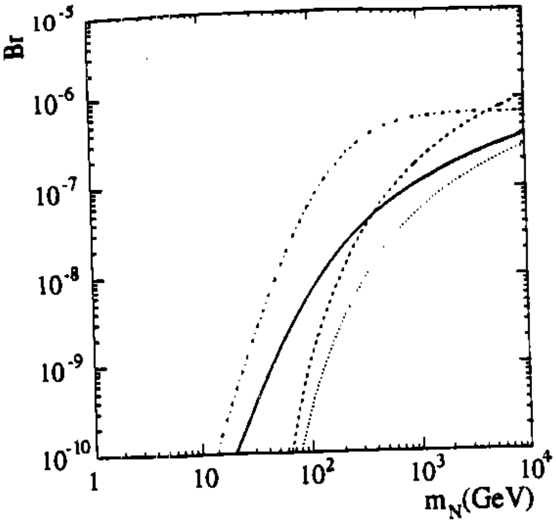

Figure 12. τ decay branching ratios consistent with weak universality.

as well as $\tau \to l_i l_j^+ l_j^-$ and $\tau \to \eta l_i$, where $l_i = \mu, e$. Present weak universality limits allow the corresponding decay branching ratios to be as large as 10^{-6} [44] as shown in Fig. 12. The solid line corresponds to the attainable branching for the decay $l \to l_i l_j^+ l_j^-$, the dashed line corresponds to $\tau \to \pi^0 l_i$, the dotted line to $\tau \to \eta l_i$ and the dash-dotted line to $\tau \to l_i \gamma$. All possible final-state leptons have been summed over in each case. The most favorable of all these channels are $\tau \to e\pi^0$ and $\tau \to e\gamma$, the second being the dominant channel for lower NHL mass values in the range $100\,GeV - 10\,TeV$. The above decays are not out of reach of a tau factory of the type now under discussion [51]. These examples illustrate that the physics of lepton flavour violation is *complementary to the physics of neutrino mass per se*, since these processes may occur even if neutrinos are strictly massless.

Flavour and/or CP violating effects at LEP energies have also been considered [43, 46]. Although too small for realistic detection, there have been dedicated searches for flavour violation at the Z peak, and limits have been obtained [52].

If the NHLS are lighter that the Z they should be singly-produced [47, 48] in Z decays

$$Z \to N_\tau + \nu_\tau \,. \tag{21}$$

The subsequent N_τ decay would give rise to large missing energy events, called zen-events. These have also been searched at LEP and very good limits have been placed [53].

6. PROBING MSW PARAMETERS AT ACCELERATORS?

As seen above, there can be large effects related e.g. with the tau lepton even in cases where neutrino masses are zero or very small. This opens an interesting possibility, i.e. that one may be able to probe MSW parameters and, say, distinguish between the large and small mixing solutions left out by the recent GALLEX results by means of high energy accelerator experiments. This would provide an independent check upon the solar neutrino oscillation parameters as determined in solar neutrino experiments.

A model that illustrates this idea is SUSY with spontaneously broken R_p, discussed above. In this model there is a large hierarchy between the squared ν_e -ν_μ mass difference δm^2 and $m_{\nu_\tau}^2$. Fixing some of the free parameters as in ref. [54]

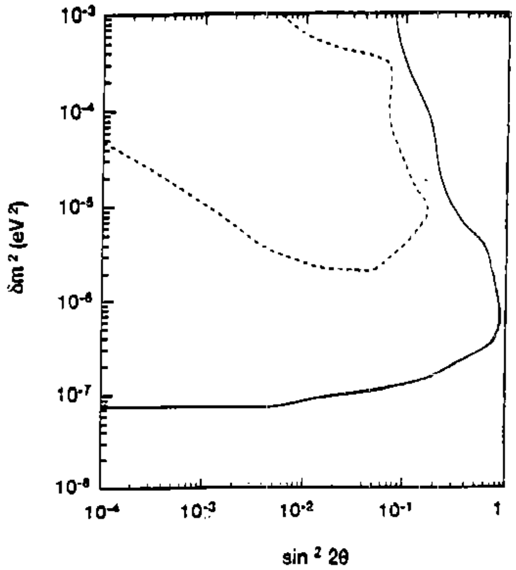

Figure 13. Distinguishing between large and small mixing MSW solutions by related tau effects

e.g. $\tan\beta = 10$, $v_R = 1\,TeV$ and $v_L = 100\,MeV$, and randomly varying others (such as the SUSY parameters, and the new Yukawa couplings $h_{\nu_{ij}}$ that determine the ν-masses) over a reasonable range one finds that the values of the oscillation parameters cover the range interesting for the MSW effect, as shown in Fig. 13. In addition, one notes that only in the small mixing solution one can have sizeable branching ratios for *both* $Z \to \chi + \tau$ and $\tau \to \mu + J$, in excess of 10^{-5} [54]. In Fig. 13 the contours correspond to $BR(Z \to \tilde{\chi} + \tau) = 10^{-6}$ (solid line) and $BR = 10^{-5}$ (dashed line).

7. CONCLUSION

The lepton sector can exhibit substantial departures from standard model expectations both in SUSY models with spontaneously broken R_p as well as in a wide class of NHL models. In addition to effects directly associated to nonzero neutrino masses, such as neutrino oscillations in vacuo, $\beta\beta_{0\nu}$ decays, etc we have also potentially detectable phenomena associated to the existence of NHLS or SUSY fermions and their admixture in the electroweak currents. There are good prospects for searching for them at accelerators, e.g. searching directly for the NHLS or SUSY fermions at the new colliders, LEP, LHC/SSC. Moreover, there are effects more related to flavour physics in the lepton sector. From this point of view it is of interest to pursue further searches for processes such as $\mu \to e + \gamma$, $\mu \to 3e$ and the related lepton flavour violating τ decays. These effects are *complementary* to direct neutrino mass searches in the sense that better neutrino mass limits do not necessarily constrain these possibilities. However, improved tests of weak universality could limit the space for the new physics described above, especially for the case of the τ lepton.

REFERENCES

1 P. W. Higgs, *Phys. Lett.* **12**, 132 (1964).
2 U Amaldi, W de Boer, H. Fursteneau, *Phys. Lett.* **B260**, 417 (1991).
3 H. Haber and G. Kane, *Phys. Rev.* **117**, 75 (1985); H. P. Nilles, *Phys. Rep.* **110**, 1 (1984).
4 J. W. F. Valle, *Prog. Part. Nucl. Phys.* **26**, 91 (1991) and references therein.
5 J. C. Romao, C. A. Santos, and J. W. F. Valle, *Phys. Lett.* **B288**, 311 (1992).
6 A. Masiero and J. W. F. Valle, *Phys. Lett.* **B251**, 273 (1990).
7 J. Steinberger, in *Electroweak Physics Beyond the Standard Model*, ed. J. W. F. Valle and J. Velasco (World Scientific, Singapore, 1992).
8 G. F. Smoot et al., Berkeley preprint (1992). E. L. Wright et al., Berkeley preprint (1992).
9 Particle Data Group, *Phys. Rev.* **D45**, S1 (1992).
10 E. Kolb and M. Turner (Addison Wesley, California, 1990) and references therein.
11 J. W. F. Valle, *Phys. Lett.* **B131**, 87 (1983).
12 G. Gelmini and J. W. F. Valle, *Phys. Lett.* **B142**, 181 (1984). A. Joshipura, S.Rindani, PRL-TH/92-10.
13 M. C. Gonzalez-Garcia and J. W. F. Valle, *Phys. Lett.* **B216**, 360 (1989).

14 P. Nogueira, J. C. Romao, and J. W. F. Valle, *Phys. Lett.* **B251**, 142 (1990).

15 G. Steigman and M. Turner, *Nucl. Phys.* **B253**, 375 (1985).

16 F.Boehm and A.Hime, *these proceedings.*

17 J. Peltoniemi, D. Tommasini, J. W. F. Valle, CERN-TH.6624/92.

18 J. W. F. Valle, *Phys. Rev.* **D27**, 1672 (1983).

19 W. Haxton, *these proceedings.*

20 J. Schechter and J. W. F. Valle, *Phys. Rev.* **D25**, 2951 (1982).

21 H. Georgi, S. Glashow, and S. Nussinov, *Nucl. Phys.* **B193**, 297 (1981).

22 G. Gelmini and M. Roncadelli, *Phys. Lett.* **B99**, 411 (1981).

23 Z. Berezhiani, A. Smirnov, and J. W. F. Valle, FTUV/92-20 (1992), Phys. Lett. B, in press.

24 M. Moe, *these proceedings.*

25 T. Kirsten, *these proceedings*

26 A. Y. Smirnov, *private communication.*

27 L. Hall and M. Suzuki, *Nucl. Phys.* **B231**, 419 (1984); R. Barbieri, etal *Phys. Lett* **B238**, 86 (1990).

28 J. W. F. Valle, *Phys. Lett.* **B196**, 157 (1987).

29 M. C. Gonzalez-Garcia and J. W. F. Valle, *Nucl. Phys.* **B355**, 330 (1991).

30 V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, *Phys. Lett.* **155B**, 36 (1985).

31 B. A. Campbell, S. Davison, J. Ellis, and K. Olive, *Phys. Lett.* **B256**, 457 (1991).

32 W. Fischler, G. Giudice, R. Leigh, and S. Paban, *Phys. Lett.* **B258**, 45 (1991).

33 M. C. Gonzalez-Garcia, J. C. Romao and J. W. F. Valle, FTUV/91-42 (1991), *Nucl. Phys.* **B**, (1992), in presss.

34 J. C. Romao, N. Rius, and J. W. F. Valle, *Nucl. Phys.* **B363**, 369 (1991).

35 ARGUS coll., *Phys. Lett.* **B246**, 278 (1990).

36 J. C. Romao, F. Campos, and J. W. F. Valle, CERN-TH.6586/92, *Phys. Lett.* **B**, (1992), in presss.

37 ALEPH Coll., *Phys. Lett.* **B236**, 233 (1990); **B237**, 291 (1990); **B241**, 141 (1990); **B245**, 289 (1990); **B246**, 306 (1990); preprints CERN-PPE/91-149; DELPHI Coll., *Nucl. Phys.* **B342**, 1 (1990); *Z. Physics* . **C51**, 25 (1991); and preprint CERN-PPE/91-132; L3 Coll., preprints CERN-PPE/92-40 and references therein; OPAL Coll., *Phys. Lett.* **B236**, 224 (1990); **B251**, 211 (1990); **B253**, 511 (1991); **B268**, 122 (1991)

38 F. Campos, L. Diaz Cruz, J. C. Romao, and J. W. F. Valle, in preparation.

39 M Gell-Mann, P Ramond, R. Slansky, in *Supergravity*, ed. D. Freedman et al. (North Holland, Singapore, 1979); T. Yanagida, in *KEK lectures*, ed. O. Sawada et al. (1979); R. Mohapatra and G. Senjanovic, *Phys. Rev. Lett.* **44**, 912 (1980)

40 J. W. F. Valle, *Nucl. Phys. Proc. Suppl.* **11**, 118 (1989), and references therein.

41 R. Mohapatra and J. W. F. Valle, *Phys. Rev.* **D34**, 1642 (1986).

42 J. Bernabeu, A. Santamaria, J. Vidal, A. Mendez, and J. W. F. Valle, *Phys. Lett.* **B187**, 303 (1987).

43 M. Dittmar and J. W. F. Valle, yellow report CERN-91/02 (1991).

44 M. C. Gonzalez-Garcia and J. W. F. Valle, *Mod. Phys. Lett.* **A7**, 477-488 (1992).

45 G. C. Branco, M. N. Rebelo, and J. W. F. Valle, *Phys. Lett.* **B225**, 385 (1989).

46 N. Rius and J. W. F. Valle, *Phys. Lett.* **246B**, 249 (1990).

47 M. Dittmar, M. C. Gonzalez-Garcia, A. Santamaria, and J. W. F. Valle, *Nucl. Phys.* **B332**, 1 (1990).

48 M. C. Gonzalez-Garcia, A. Santamaria, and J. W. F. Valle, *Nucl. Phys.* **B342**, 108 (1990).

49 J. Schechter and J. W. F. Valle, *Phys. Rev.* **D22**, 2227 (1980); *Phys. Rev.* **D23**, 1666 (1981)

50 M. Gronau, C. Leung, and J. Rosner, *Phys. Rev.* **D29**, 2539 (1984).

51 R. Alemany, J. J. Gomez-Cadenas, M. C. Gonzalez-Garcia, and J. W. F. Valle, in preparation.

52 OPAL coll., *Phys. Lett.* **B254**, 293 (1991).

53 OPAL coll., *Phys. Lett.* **B247**, 448 (1990).

54 J. C. Romao and J. W. F. Valle, *Phys. Lett.* **B272**, 436 (1991); *Nucl. Phys.* **B** (1992) , in press.