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Can solar neutrino oscillation parameters be probed at LEP?

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Supersymmetry with spontaneously broken R parity naturally leads to the explanation of the observed deficit in the solar neutrino flux via matter-enhanced ν_e–ν_μ oscillations, while m_{ν_τ} is large enough to lead to related signatures associated with the τ lepton. These include single chargino production in Z decays with branching ratio $\mathrm{BR}(Z\to\tilde{\chi}\tau)$ as large as $\sim 6\times10^{-5}$, accessible at LEP1. The ν_τ is naturally much heavier than ν_μ and decays to it via majoron emission, with a lifetime short enough to obey cosmological limits.

1. Introduction

The possibility that the explanation of the solar neutrino deficit through the MSW effect could be independently (although indirectly) checked in accelerator experiments has attracted some attention recently [1,2]. Unfortunately, the suggested models are no longer phenomenologically viable, since they lead to a new invisible decay mode for the neutral gauge boson involving the emission of light scalars,

$$Z\to\rho+J\,, \qquad (1)$$

now ruled out by LEP measurements of the *invisible* Z width [3].

Here we point out that this idea can be *naturally* implemented in a slight variant of the original model [2], in a way completely consistent with observation *and* with the added bonus of having a richer set of high energy processes accessible to experiment. This should enable indirect tests of the validity of the MSW explanation of the solar neutrino deficit and constrain the relevant oscillation parameters. Here we only summarize the relevant results, postponing all detailed discussions to a separate publication. The model has been suggested in ref. [4] and is based on the idea of spontaneous R parity violation [5]. These theories have now acquired an extra interest also insofar as they are not restricted by considerations related to the erasure of the GUT-produced baryon asymmetry of the universe due to non-perturbative electroweak effects [6,7]. This is so as long as the relevant R_p breaking scale is lower than a few TeV. In this picture the universe has always been in the R_p-symmetric phase until slightly after the electroweak phase transition takes place. Proton stability is automatic and the phenomenological implications can be systematically derived. They provide an alternative to the minimal SUSY standard model (MSSM) that leads to new processes accessible to experiment. The signatures of supersymmetry in these spontaneously broken R parity models are different from those of the MSSM in several ways. First, these $SU(2)\otimes U(1)$ theories lead to the existence of a physical massless Nambu–Goldstone boson – a majoron – denoted J.

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The existence of the majoron enables the heavy ν_τ to decay to it via majoron emission, with a lifetime short enough to obey cosmological limits [8–10].

In addition, SUSY particles can be singly produced, for example in Z decays [11]. Here we suggest the possibility of using these *R* parity breaking effects to indirectly test the MSW oscillation parameters at LEP1. This possibility is rather remarkable, taking into account the smallness of the neutrino mass values needed for the MSW effect. However, spontaneously broken *R* parity models imply a natural mass hierarchy in the values of m_{ν_e}, m_{ν_μ} and m_{ν_τ} that naturally combines an ultralight ν_e–ν_μ sector with a fairly heavy Majorana ν_τ. If strongly mixed in the charged current weak interaction, ν_τ exchange would give too large a neutrinoless double β decay rate. More remarkably, even high energy accelerator experiments, such as LEP, would be sensitive to new effects related to the heavy ν_τ.

2. Spontaneous *R* parity violation

The minimal way to break *R* parity spontaneously is through a non-zero VEV for the scalar neutrino [2,5]

$$v_\mathrm{L} = \langle \tilde{\nu}_{\mathrm{L}\tau} \rangle . \tag{2}$$

In such a model there are majoron emission processes, such as the Compton-like reaction $\gamma + e \to e + J$, that lead to new mechanisms of stellar energy loss. To avoid this one has to constrain the magnitude of the left-handed sneutrino VEV [12]. We estimate the resulting limit to be

$$v_\mathrm{L} \leqslant \mathrm{O}(30\ \mathrm{keV}) . \tag{3}$$

However this implies that the decay in eq. (1) is unsuppressed, leading to a conflict with LEP results. A way out of this which avoids in a simple way the LEP constraint on $\Gamma_\mathrm{Z}^\mathrm{inv}$ was suggested in ref. [4]. In the new model the majoron is an isosinglet, and does not couple to the Z. The scale characterizing *R* parity breaking is now large, i.e.,

$$v_\mathrm{R} = \mathrm{O}(1\ \mathrm{TeV}) , \tag{4}$$

while the analogue of eq. (3) is estimated to be [4]

$$v_\mathrm{L} \leqslant \mathrm{O}(100\ \mathrm{MeV}) \tag{5}$$

and can be obtained without the need of unnatural fine-tuning of the parameters in the Higgs potential [4,13]. The superpotential terms are given as [#1]

$$\begin{aligned} &h_\mathrm{u} \mathrm{u}^\mathrm{c} \mathrm{Q} \mathrm{H}_\mathrm{u} + h_\mathrm{d} \mathrm{d}^\mathrm{c} \mathrm{Q} \mathrm{H}_\mathrm{d} + h_\mathrm{e} \mathrm{e}^\mathrm{c} \ell \mathrm{H}_\mathrm{d} + (h_0 \mathrm{H}_\mathrm{u} \mathrm{H}_\mathrm{d} - \epsilon^2) \Phi \\ &\quad + h_\nu \nu^\mathrm{c} \ell \mathrm{H}_\mathrm{u} + h \Phi \nu^\mathrm{c} \mathrm{S} + \hat{\mu} \mathrm{H}_\mathrm{u} \mathrm{H}_\mathrm{d} + M \nu^\mathrm{c} \mathrm{S} \\ &\quad + M_\Phi \Phi\Phi . \end{aligned} \tag{6}$$

The superfields $(\Phi, \nu_i^\mathrm{c}, \mathrm{S}_i)$ are singlets under $\mathrm{SU}(2) \otimes \mathrm{U}(1)$ and carry a conserved lepton number assigned as $(0, -1, 1)$ respectively. All couplings h_u, h_d, h_e, h_ν, h as well as the mass M are described by arbitrary matrices in generation space which explicitly break flavour conservation. These additional singlets, e.g. S, may arise in several extensions of the standard model [15], and may lead to interesting phenomenological signatures [16,17].

While the superpotential of eq. (6) conserves *total* lepton number as well as *R* parity, the presence of the new singlets can drive the spontaneous violation of *R* parity and electroweak symmetries [4,13]. This leads to the existence of a majoron given by the imaginary part of

$$\begin{aligned} &(v_\mathrm{L}^2 / V v^2)(v_\mathrm{u} \mathrm{H}_\mathrm{u} - v_\mathrm{d} \mathrm{H}_\mathrm{d}) \\ &\quad + (v_\mathrm{L}/V) \tilde{\nu}_\tau - (v_\mathrm{R}/V) \tilde{\nu}_\tau^\mathrm{c} + (v_\mathrm{S}/V) \tilde{\mathrm{S}}_\tau . \end{aligned} \tag{7}$$

The isosinglet VEVs

$$v_\mathrm{R} = \langle \tilde{\nu}_\tau^\mathrm{c} \rangle , \tag{8}$$

$$v_\mathrm{S} = \langle \tilde{\mathrm{S}}_\tau \rangle , \tag{9}$$

with $V = \sqrt{v_\mathrm{R}^2 + v_\mathrm{S}^2}$ set the scale of *R* parity or lepton number breaking, while the isodoublet VEVs

$$v_\mathrm{u} = \langle \mathrm{H}_\mathrm{u} \rangle , \tag{10}$$

$$v_\mathrm{d} = \langle \mathrm{H}_\mathrm{d} \rangle \tag{11}$$

drive electroweak breaking and the fermion masses. The combination $v^2 = v_\mathrm{u}^2 + v_\mathrm{d}^2$ is fixed by the W mass, while the ratio of isodoublet VEVs determines the parameter

$$\tan\beta = v_\mathrm{u}/v_\mathrm{d} . \tag{12}$$

[#1] Note that we have added some new terms that were not included in refs. [4,11,14]. These are allowed by our symmetries and give more flexibility in obeying all of the experimental constraints and/or justifying our approximate treatment of the neutral fermion sector.

A necessary ingredient for the consistency of this model is the presence of a small seed of R parity breaking in the SU(2) doublet sector. Using the results of ref. [12] we estimate that

$$v_L^2/v_R m_W \lesssim 10^{-7} \tag{13}$$

in order to adequately suppress the resulting stellar energy loss to acceptable levels. This may be easily satisfied for $v_R = O(1\ \text{TeV})$ provided $v_L \leqslant O(100\ \text{MeV})$. This constraint on v_L is more than three orders of magnitude less stringent than the one that holds in the model of ref. [2]. Moreover, the smallness of the ratio v_L/v_R may be achieved in a natural way, from the Higgs potential [4]. This follows because $v_L = \langle \tilde{\nu}_{L\tau} \rangle$ is related to the Yukawa coupling h_ν and vanishes as $h_\nu \to 0$. Last, but not least, the present model has only the canonical neutrino counting, since the majoron contribution is avoided.

3. Neutrino masses and decays

In spontaneously broken R parity models the L-number is violated, so that neutrinos acquire masses at the tree level, from mixing with the heavy neutral R-odd fermions [1]. The heavier neutrino is the ν_τ. Its mass m_{ν_τ} may be given in the approximation where we neglect v_L as

$$m_{\nu_\tau} \simeq \frac{\sum_i h_{\nu i\tau}^2 M_0 v_R^2 v_d^2}{\mu(2v_u v_d M_0 - \mu M_1 M_2)}, \tag{14}$$

where we have set $M_0 = g_1^2 M_2 + g_2^2 M_1$. In our numerical study, however, we have used the expression for m_{ν_τ} corrected for the general case where $v_L \neq 0$.

In order to determine the masses of the two lowest mass neutrinos relevant for the description of the propagation of solar neutrinos in our model – ν_e and ν_μ – we have to consider in detail the full neutralino mass matrix. One can show that, in complete generality, there is a massless neutrino in our model. To a good approximation, it is given by

$$\nu_1 \simeq \cos\theta\, \nu_e - \sin\theta\, \nu_\mu , \tag{15}$$

where

$$\tan\theta = h_{\nu 13}/h_{\nu 23} . \tag{16}$$

In this case the other light neutrino is

$$\nu_2 \simeq \alpha \sin\theta\, \nu_e + \alpha \cos\theta\, \nu_\mu + N_{23}\, \nu_\tau , \tag{17}$$

where

$$\alpha \simeq -(\mu v_L/v_d + h_{\nu 33} v_R)[(h_{\nu 13}^2 + h_{\nu 23}^2) v_R^2 + (\mu v_L/v_d + h_{\nu 33} v_R)^2]^{-1/2} \tag{18}$$

and

$$N_{23} \simeq \sqrt{h_{\nu 13}^2 + h_{\nu 23}^2}\, v_R [(h_{\nu 13}^2 + h_{\nu 23}^2) v_R^2 + (\mu v_L/v_d + h_{\nu 33} v_R)^2]^{-1/2} \tag{19}$$

and its mass can be given, in a reasonable approximation, by

$$m_2 \simeq (h_{\nu 13}^2 + h_{\nu 23}^2) v_u^2 v_L^2 [h_{33}(\mu + h_0 \langle \Phi \rangle) + h_0 M_{33}]^2 \times [M_\Phi (h_{33} \langle \Phi \rangle + M_{33})^2 h_{\nu 33}^2 v_d^2]^{-1} , \tag{20}$$

showing that it vanishes as $v_L \to 0$. For small values of θ the massless neutrino is mostly ν_e so that ν_2 corresponds to ν_μ. In the following we will always refer to ν_2 as ν_μ and to ν_1 as ν_e although this assignment is not valid for large values for the ν_e–ν_μ mixing angle θ. It follows from eq. (20) and eq. (14) that the ν_μ mass is extremely small on the scale of the ν_τ mass. For our choices of parameters, $v_R = 1$ TeV and $v_L = 100$ MeV, this implies a ν_μ–ν_τ mass ratio that can be as small as

$$m_{\nu_\mu}/m_{\nu_\tau} = O(10^{-8}) . \tag{21}$$

Thus for a ν_τ mass in the 10 keV–1 MeV range, natural in this model, we expect a ν_μ mass in the range 10^{-4}–10^{-2} eV, *just the one relevant for the MSW effect*. Similarly, the ν_e–ν_μ mixing angle given by eq. (16) may easily lie in the desired range where the MSW effect can effectively reduce the solar neutrino flux.

The expected size of the ν_τ mass in this model seems to be in conflict with the cosmological limit [18]

$$\sum_i m_{\nu_i} \lesssim 100\ \text{eV} \tag{22}$$

on the abundance of relic neutrinos. Fortunately, our model also has the solution to this apparent conflict. It relies on the existence of a new ν_τ decay mode involving majoron emission [#2]

$$\nu_\tau \to \nu_\mu + J \tag{23}$$

that naturally realizes the early proposal considered in refs. [8–10]. Unlike the situation in the originally proposed majoron model of ref. [8], the neutralino mass matrix is such that the majoron couplings to neutrinos is not simultaneously diagonal as the mass matrix [9], opening the possibility of a fast decay rate [10]. The attainable ν_τ lifetimes are given in fig. 1 as a function of the ν_τ mass. They should be compared to the cosmological limit on the ν_τ decay lifetime required in order to efficiently suppress the relic ν_τ contribution. This is shown as the solid line in fig. 1. Clearly the decay lifetimes can be much shorter than required by cosmology for a wide range of values of the parameters. Moreover, since these decays are *invisible*, they are consistent with all astrophysical observations [1]. If, however, the universe is to have become matter-dominated by a redshift of 1000 at the latest (so that fluctuations have grown by at least a factor of 1000 by today), the ν_τ lifetime has to be much shorter [19], as indicated by the dashed line in fig. 1. Again, lifetimes below the dashed line are allowed in the present model. However, this lifetime limit is much less reliable than the one derived from the critical density, since there is not yet an established theory for the formation of structure in the universe. We have therefore not imposed it in what follows.

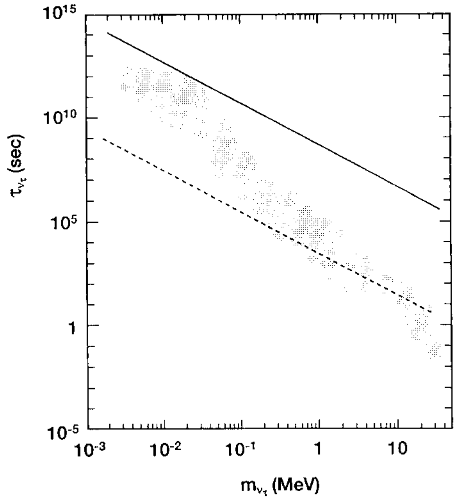

Fig. 1. The shaded area represents the attainable values of the ν_τ lifetime versus its mass, when the model parameters are varied as explained in the text. The curve delimiting the lower part of the shaded contour gives an estimate of the minimum lifetime consistent with observational constraints. For comparison we also show the cosmological critical density limit (solid line) and the *naive* limit that one may derive from galaxy formation (dashed line).

In short, in this model m_{ν_τ} can be larger than eq. (22), because of the existence of an efficient majoron decay channel, eq. (23). The large m_{ν_τ} values acceptable in this model enable the rare Z decays to be correspondingly enhanced. In addition, the large hierarchy between m_{ν_μ} and m_{ν_τ} naturally combines these measurable effects with the MSW oscillations, required to explain the solar neutrino data. This suggests the possibility of having an independent check upon the solar neutrino oscillation parameters as determined by solar neutrino experiments with searches performed in high energy laboratory experiments, such as LEP. We now turn to this question.

4. MSW effect

In ref. [20] we give a detailed study of the phenomenological implications of spontaneously broken R parity for the propagation of solar neutrinos. The channel in which solar neutrino oscillations are expected to take place is ν_e to ν_μ, since, as discussed, their mass difference is $\sim v_L^2$, much smaller than the heavy ν_τ mass, which scales as $\sim v_R^2$. The ν_τ is therefore decoupled from the solar neutrino oscillations.

For definiteness we have fixed in our analysis characteristic values for the following parameters:

$$\tan\beta = 10\,,$$

$$v_R = 1\ \mathrm{TeV}\,,$$

$$v_L = 100\ \mathrm{MeV}\,,$$

$$m_\Phi = 10\ \mathrm{TeV}\,,$$

$$M_{33} = 1\ \mathrm{TeV}\,, \qquad (24)$$

and assumed, for simplicity, that $v_R = v_S$.

The allowed region of oscillation parameters has been determined as a function of the relevant SUSY parameters μ and M_2. These have been randomly varied in the range [#3]

$$30 \leqslant M_2/\mathrm{GeV} \leqslant 250\,, \qquad (25)$$

[#2] One can also show that the decay $\nu_\tau \to \nu_e + J$ is strictly forbidden and that the mass of ν_e is zero.

[#3] It is always possible, if CP is conserved, to choose $M_2 > 0$, while μ may have either sign.

$$-250 \leqslant \mu/\text{GeV} \leqslant 250\,. \qquad (25\text{ cont'd})$$

The ν_τ and ν_μ masses are controlled by the parameters $h_{\nu ij}$ as can be seen from eq. (14) and eq. (20). These were varied randomly in the interesting range given by

$$10^{-10} \leqslant h_{\nu 13}, h_{\nu 23} \leqslant 10^{-1}\,,$$
$$10^{-4} \leqslant h_{\nu 33} \leqslant 10^{-1}\,. \qquad (26)$$

We have performed a careful sampling of the points in our parameter space that are allowed by all observational constraints, in order to evaluate the attainable magnitudes of m_{ν_μ} and the corresponding mixing angle θ. The constraints include all of the collider constraints relevant in any SUSY extension of the standard model as well as neutrino physics constraints. The allowed values of the oscillation parameters will determine the ν_e–ν_μ conversion probabilities and thus the solar neutrino predictions. We find that the regions of neutrino mass and mixing parameters characterizing ν_e–ν_μ oscillations presently allowed in our model cover the region where the resonant effects can play an important role in the oscillations of solar neutrinos. On the other hand, the constraints that follow from present solar neutrino observations, both from the Homestake experiment and from Kamiokande, are indicated by the shaded region shown in fig. 2, adapted from the paper by Barger, Phillips and Whisnant in ref. [21]. We also note that the preliminary results of the SAGE Collaboration point towards a very drastic reduction of the low energy pp neutrino flux, possible in the non-adiabatic branch of the shaded region. As fig. 2 shows, in the spontaneously broken R parity model the solar neutrino oscillation parameters cover the region of interest where the MSW oscillations between ν_e and ν_μ provide a successful explanation of the reduced solar neutrino flux. In addition, since we have a model where these MSW oscillations naturally coexist with a very massive ν_τ, we also expect to have other exotic effects that may be measurable in the laboratory, and *independently* check the underlying physics of the neutrino mass generation mechanism. We now turn to these signatures.

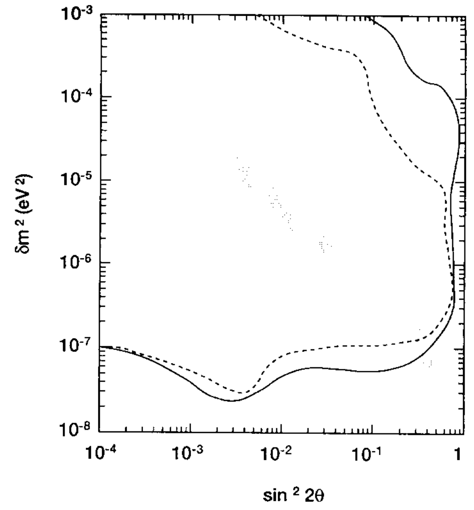

Fig. 2. Contour plots for the BR($Z \to \tilde{\chi} + \tau$). The solid line corresponds to BR $\geqslant 10^{-6}$, while the dashed line corresponds to BR $\geqslant 10^{-5}$. Also shown is the region allowed by present solar neutrino data. The figure shows that the decay $Z \to \tilde{\chi}\tau$ is accessible to the LEP experiments.

5. LEP1 signatures

In the spontaneously broken R parity model there are observable effects in Z decays that could well be observable at LEP1 [11]. Here we investigate whether their observability is consistent with the understanding of present solar neutrino observations in terms of matter-enhanced neutrino oscillations. Usually it is not easy to achieve models with these properties in view of the smallness of neutrino masses needed in the MSW effect. In the present model, however, the ν_τ mass is expected to be large, thus opening the way to this interesting complementarity between laboratory and astrophysical effects. The latter also involve stellar cooling mechanisms, since these are affected by majoron emission processes [12].

We now show that, indeed, in the spontaneously broken R parity model, ν_e–ν_μ MSW oscillations may be related with other processes that could be observable in the laboratory. We consider the example of LEP. If R parity does not hold exactly then SUSY particles may be singly-produced. In ref. [11] we made a study of the *single* chargino decay mode

$$Z \to \tilde{\chi}\tau \qquad (27)$$

and showed that it may be observable at LEP1 [4]. We have demonstrated that this decay is accessible to experiment even when the parameters of the MSW oscillations lie in the region indicated by present solar neutrino experiments. To do this we have performed a systematic analysis of the attainable values of this branching ratio once observational constraints are taken into account. These include those from laboratory, cosmology and astrophysics, including the requirement that the solar neutrino oscillation parameters lie in the range allowed by present solar neutrino observations. Our results are summarized in fig. 2. They show that the corresponding $Z \to \tilde{\chi}\tau$ branching ratio could possibly be measured at LEP1, thus enabling one to have an independent handle on the neutrino oscillation parameters!

6. Discussion

We conclude that the MSW oscillations responsible for the explanation of the reduced solar neutrino flux may well be *accompanied by effects that can be observable in the laboratory* and therefore these can be used as an additional tool to restrict the underlying physics. For example, the non-observation of the new effects would place independent restrictions upon solar neutrino oscillation parameters. We showed that the peculiar nature of the neutrino mass spectrum of the spontaneously broken R parity model, which combines a very small ν_e–ν_μ mass difference with a heavy Majorana ν_τ mass, naturally allows for the interesting possibility that the ν_e–ν_μ oscillations that can successfully explain present solar neutrino observations are accompanied by effects that can be probed not only in neutrinoless double β decay experiments, but also in high energy particle physics accelerator experiments, such as LEP. We identified a relevant process, i.e., the R parity violating Z decay mode in eq. (27). This decay may go with a branching ratio as large as $\mathrm{BR}(Z \to \tilde{\chi}\tau) \sim 6 \times 10^{-5}$, and can be studied at LEP1. In addition, since the lightest SUSY fermion is unstable, the origin and the rates for zen events may also be different than in the minimal supersymmetric standard model. Therefore these effects might be used as an independent tool to probe the physics underlying the explanation of the solar neutrino deficit. We showed that this idea does not contradict any laboratory, cosmological or astrophysical observation.

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