Supersymmetric Majoron Signatures and Solar Neutrino Oscillations

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Spontaneous R-parity breaking in supergravity solves the solar neutrino problem through matter-enhanced neutrino oscillations. The model may be tested in collider experiments and through "dynamical" effects associated with the existence of a weakly interacting majoron. Apart from astrophysical effects, majoron emission can produce observable changes in μ and τ decay spectra for parameter values that substantially reduce the solar neutrino flux. A signature of the model is the possible observation of the decay $\mu \rightarrow e +$majoron.

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The solar neutrino problem may be resolved either with new physics to change the solar parameters[1] or with nonstandard weak interactions to modify neutrino propagation properties. These include neutrino decay,[2] neutrino magnetic moment,[3] and neutrino oscillations.[4-6] While it is possible to build consistent models for fast neutrino decay,[7] the required value of the neutrino magnetic moment needed to solve the solar neutrino problem is hard to reconcile with existing limits on neutrino mass. Neutrino decay, however, has also been rendered unlikely by the recent observation of neutrinos from the SN1987A supernova[8] although, for large mixings, the model may still survive.[9] This may leave oscillations as the most likely explanation of the solar neutrino problem in terms of nonstandard neutrino properties. Unfortunately, it will be difficult to test experimentally for the oscillation hypothesis. If oscillations occur *in vacuo,* then the relevant mass-squared differences are too small to be probed.[6] A lot of interest has been recently devoted to the possibility of *matter-enhanced* oscillations.[4] These are possible because of the effect of coherent neutrino scattering in the solar medium even when vacuum mixing angles are very small. This resonant enhancement is possible, however, only if the neutrino mass difference lies in the range[10]

$$10^{-7} \lesssim [\delta m/(1\ \mathrm{eV})]^2 \lesssim 10^{-4}. \quad (1)$$

Again the oscillation hypothesis indicates mass differences that are too small to be checked in existing experimental setups. Although effort will be pushed in this direction,[11] it will be a number of years before one will be able to start probing the range specified in Eq. (1). From this point of view it is interesting to ask the following questions: What are the physics options for having an experimentally testable solution to the solar-neutrino puzzle in terms of neutrino oscillations? What is the physical origin for the relevant neutrino mass scale?

The second question may be answered in a variety of ways, e.g., by introducing in the theory heavy singlet fermions such as right-handed neutrinos, at a mass scale *above* that of electroweak-symmetry breaking. This could, for example, be the grand unification scale in SO(10) models[12] or some intermediate scale such as is present, e.g., in some superstring-inspired E_6 models.[13,14] Alternatively, one may introduce a physical mass scale *below* the Fermi scale and have a see-saw mechanism in which the states relevant in the determination of the neutrino mass are at the Fermi scale.[15] A natural framework to implement this idea is supergravity, in which case the heavy states necessary in the see-saw mechanism are just the supersymmetric (SUSY) partners of gauge and Higgs fields. *In both cases we are still left with the first question.*

There are two generic types of models where the solar-neutrino oscillation hypothesis may be experimentally checked. First it can be checked "kinematically" in models where $m \gg \delta m$, where m represents a typical neutrino mass and δm denotes a neutrino mass difference. One way to model this hierarchical difference in the values of m and δm is to attribute it[5] to the presence of a large intermediate scale in superstring-inspired models.[13] The electron neutrino is just one Weyl component of a four-component quasi-Dirac neutrino[16] of mass $m \gg \delta m$. The other component may be another neutrino flavor or, as in Ref. 5, can be a *sterile* neutrino. If ν_e is a quasi-Dirac particle, it can be massive enough to be detectable in tritium decay experiments and also to be relevant for cosmology, without conflicting with any experimental limit. The solar-neutrino problem can then be solved by large vacuum oscillations from the *active* to the *sterile* component as ν_e propagates from sun to Earth.

The alternative possibility is $m \simeq \delta m$ so that "kinematical" neutrino mass effects will not be detectable but the theory contains some new "dynamical" degree of freedom whose "large" effects can be probed and used, in a sense, to "track" an otherwise undetectably small mass. The prototype of this situation is when $B-L$ is a spontaneously broken symmetry of the Lagrangian. If ungauged, as in the standard model, spontaneous $B-L$ breaking generates a Goldstone boson—a majoran —which we denote by J. (If the gauge group contains $B-L$ there will be instead an additional gauge boson, coupled to neutrinos.[17]) The majoron is a true dynamical degree of freedom and so it has interactions which are related in a well-defined way with the neutrino mass. Its emission generates new mechanisms of stellar energy loss. There are a variety of different majoron models; here we concentrate on a variant of the idea where the majoron is the SUSY partner of the neutrino.[15] In this case majoron emission will also produce small changes in the decay parameters of the muon and the τ lepton which could be seen in precision measurements. The present good agreement of the observations with the standard-model predictions leads to nontrivial constraints. Interestingly enough, these constraints may allow the possible observation of rare decays such as $\mu \to e+J$ which would provide an interesting signature for this scenario. The smallness of the neutrino mass, required by the resonance condition, Eq. (1), is dictated by constraints on the couplings of the majoron that follow from a variety of considerations thus giving a dynamical basis for the see-saw mechanism.

Spontaneously broken R-parity (SBRP) model.— The model is described by the minimal supergravity superpotential

$$h^u_{ij}u^c_iQ_jH_u+h^d_{ij}d^c_iQ_jH_d+h^e_{ij}e^c_il_jH_d+\mu H_uH_d, \qquad (2)$$

where the parameter μ is related with electroweak breaking driven by radiative corrections associated with the top quark. The first three terms give rise to masses for up- and down-type quarks and charged leptons, respectively, once the two Higgs fields H_u and H_d acquire their vacuum expectation values (VEV's) v_u and v_d. In general, supersymmetry and lepton flavor are broken explicitly in the scalar potential via soft scalar mass terms and possibly cubic scalar self-couplings, and also via $SU(3)\otimes SU(2)\otimes U(1)$-invariant gaugino mass terms M_i, $i=1,2,3$. Gaugino masses break the continuous R invariance of the theory down to a discrete symmetry, called R parity. R parity is even for all particles of the standard model (including the Higgs scalars) and odd for their SUSY partners. R parity too may be broken, either explicitly[18] or spontaneously by nonzero VEV's for the scalar neutrinos,[19]

$$v_i=\langle \tilde{l}^0_i\rangle, \quad i=e,\mu,\tau. \qquad (3)$$

In the minimal model, spontaneous R-parity breaking is very restrictive. Recent analysis[20] indicates that sleptons lighter than about 65 GeV and a top quark heavier than about 70 GeV are required. (This breaking may be far easier to achieve if one adds, e.g., terms that break total lepton number explicitly as well.)

Spontaneous breaking of $B-L$, an ungauged continuous symmetry, generates a Nambu-Goldstone boson—a majoron—given in Ref. 15, giving a neutrino mass

$$m=\mu M\textstyle\sum_i v_i^2/(2v_uv_dM-M_1M_2\mu), \qquad (4)$$

where $2M=g_1^2M_2+g_2^2M_1$ and g_i are gauge coupling constants. Note that since $B-L$ is broken by *one* unit via the scalar neutrino VEV, Eq. (3), it takes *two* such breakings to generate a (Majorana) mass for the (left-handed) neutrino; hence the square in the see-saw formula, Eq. (4). This contrasts with the nonsupersymmetric majoron model[21] in which a scalar Higgs triplet is introduced to generate neutrino masses directly and therefore *linear* in the lepton-number-breaking expectation value.

Another striking difference, which makes the SBRP model much more restrictive, is the fact that *one and only one* neutrino acquires mass, namely, the one which is related by supersymmetry to the majoron. Its mass is given by Eq. (4). (The full tree-level neutral-lepton mass matrix was given in Ref. 15 and radiative corrections are negligible for our purposes.) This simplifies considerably the structure of the charged-current weak interaction, reducing the parameters relevant for the description of the resonant neutrino oscillations in the sun to just *three* parameters: two mixing angles (the third angle is not a physical parameter, because of the mass degeneracy between the two massless neutrinos) and one neutrino mass parameter m. Moreover, CP is conserved in the charged current. The matrix K describing the charged-current weak interaction can be put in the canonical form[15,22]

$$K=\omega_{23}(\theta_{23})\omega_{13}(\theta_{13}), \qquad (5)$$

where ω_{ij} is a rotation by an angle θ_{ij} in the i-j plane.

For $v_\tau \gg v_\mu \gg v_e$ the mixing angles in Eq. (5) are small so that the massive state is mostly ν_τ and resonant solar-neutrino conversions occur from ν_e into ν_τ.[15] Because of such drastic simplification, there are interesting experimental signatures that make the SBRP model testable. In Ref. 15 we showed how the gaugino-Higgsino mass spectrum is restricted by Davis's experimental results and how the allowed range of parameters where resonant amplification of solar neutrino oscillations can occur may be directly explored by our searching for charged SUSY partners of gauge and Higgs particles, say, in electron-positron machines. Here we concentrate on low-energy tests of the model based on the dynamics of the majoron.

Constraints on majoron couplings.— The majoron couples to quarks through its Higgs-doublet admixture via the first two terms in Eq. (2). These couplings are always flavor diagonal. The situation is different for the charged leptons as a result of their mixing with charginos (charged SUSY partners of gauge and Higgs bosons), described by the mass matrix

$$\begin{array}{c|ccc} & e_j^+ & \tilde{H}_u^+ & \tilde{W}^+ \\ \hline e_i^- & h_{ij}v_d & 0 & g_2v_i \\ \tilde{H}_d^- & h_{ji}v_i & \mu & g_2v_d \\ \tilde{W}^- & 0 & g_2v_u & M_2 \end{array} . \qquad (6)$$

Since fermions of different weak isospin are mixed in

Eq. (6), it follows that in our model the majoron couplings to charged leptons are *not* flavor diagonal. *This is another important difference between the SBRP model and the triplet majoron model which, as we will see, may be tested experimentally.* The coupling of the majoron to physical charged leptons is described by the effective interaction Lagrangean,

$$i\frac{g_2^2}{\sqrt{2}}\frac{vm_j}{m_W^2}(\tfrac{1}{2}\delta_{ij}+fg_{ij})\bar{e}_{Li}e_{Rj}J+\text{H.c.}, \quad (7)$$

where the function f is given by

$$f=\frac{1+x^2}{(1-xy)^2}=m_W^2\left[\frac{\cos^2\phi}{\mu_1^2}+\frac{\sin^2\phi}{\mu_2^2}\right], \quad (8)$$

ϕ is a chargino mixing angle, and μ_i are their masses. x and y denote, respectively, the Higgsino mixing parameter μ and the supersymmetry-breaking SU(2) gaugino mass parameter M_2 in units of the W mass ($x=\mu/m_W$, $y=M_2/m_W$). For simplicity we took $v_u \approx v_d$. The coupling matrix g_{ij} is a projection matrix, given by

$$g_{ij}=\begin{bmatrix} s_{13}^2 & c_{13}s_{13}s_{23} & c_{13}s_{13}c_{23} \\ c_{13}s_{13}s_{23} & c_{13}^2s_{23}^2 & c_{13}^2c_{23}^2s_{23} \\ c_{13}s_{13}c_{23} & c_{13}^2c_{23}^2s_{23} & c_{13}^2c_{23}^2 \end{bmatrix}. \quad (9)$$

The first term in Eq. (7) comes from majoron admixture in the H_d Higgs scalar while the second comes from lepton-chargino mixing.

In this model stellar energy loss proceeds via *single* majoron emission in the Compton-type process $\gamma+e\to e+J$. From the cross section given by Georgi, Glashow, and Nussinov[23] and Ellis and Olive[23] and the coupling of Eq. (7) we have

$$v(\tfrac{1}{2}+fg_{ee}) \lesssim 30\ \text{keV}. \quad (10)$$

In addition, *double*-majoron emission occurs in $\gamma+e\to e+J+J$. This process has normal gauge strength couplings but is mediated by the heavy charginos. We have calculated the total cross section for the process in the approximation of small photon energies $E_\gamma \ll m_e$ and obtained the limit

$$f^2g_{ee}^2 \lesssim 10. \quad (11)$$

The same couplings also induce, mediated by the heavy charginos, double-majoron-emission processes $e_j\to e_i+J+J$, thus changing the spectrum of $e_j\to e_i+\nu+\nu$. Requiring the fractional change in the Michel parameter to be within the accuracy of present experimental determination gives, for the case of $\mu\to e+J+J$ decay, the limit[24]

$$f^2g_{e\mu}^2 \lesssim 10^{-2}, \quad (12)$$

while for the corresponding τ decays, the limits are

$$f^2g_{e\tau}^2 \lesssim 0.35 \quad (13)$$

and a similar bound on $fg_{\mu\tau}$. Because of the existence of another light scalar in the theory the limits could be up to a factor of 2 stronger. Thus, from this point of view, precision measurements of μ and τ decays are theoretically very interesting.

As a result of the couplings of Eq. (7) *single* majoron emission in μ and τ decays can also occur. The branching ratio B_{ij} for $e_j\to e_i+J$, relative to $e_j\to e_i+\nu+\nu$, is given by

$$B_{ij}=96\pi^2(v/m_j)^2f^2g_{ij}^2. \quad (14)$$

For the case of $\mu\to e+J$ decay the branching ratio may be as large as

$$B_{e\mu}\approx 3\times10^{-7}[v/(30\ \text{keV})]^2, \quad (15)$$

which is suggestively close to the present experimental limit from TRIUMF, $B_{e\mu}<2.6\times10^{-6}$.[25] For τ decays, however, the corresponding allowed branching ratios are far below what can be experimentally probed.

Discussion.—The supersymmetric majoron is markedly different from the triplet majoron[21] and also less likely to be excluded by accurate Z-width measurement. In this model the tiny neutrino mass needed for the Mikheyev-Smirnov-Wolfenstein effect[4,10] is accompanied by "large" dynamical effects associated with the existence of the majoron, thus providing a dynamical testing ground for the model. Majoron emission in μ and τ decays may be at the cutting edge of experimental test, for parameter values which solve the solar-neutrino problem by the enhanced oscillation effect. If mixing angles s_{23} and s_{13} are small, resonant solar neutrino conversions occur from ν_e into ν_τ.[15] A reduction in solar neutrino flux below 2.6 solar neutrino units implies severe restrictions on the supersymmetric spectrum[15] and potentially large effects in τ decays. As an example taking $v\approx 30$ keV, $s_{13}\approx s_{23}\approx 0.1$, and $f\approx 6$, corresponding to the lightest chargino mass values consistent with present limits from the DESY e^+e^- storage ring PETRA, we get a deviation in the τ-decay Michel parameter comparable to present limits, while the change for the μ decay is a factor of 3 or so below present experimental limits. For the same parameters the branching ratio for $\mu\to e+J$ is $\approx 3\times10^{-7}$. A larger value for s_{23} would increase both the deviation in the Michel parameter and the branching ratio for single majoron emission in μ decay. These effects are likely to be larger by up to a factor of 2 because of the existence of another light scalar in the theory. On the other hand, if we take a larger value for v the branching ratio $B_{e\mu}$ increases without changing the Michel parameters. Note that for these parameter values, the lightest *chargino* weighs less than half of the Z mass and thus might be seen at accelerator experiments.

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[1]G. West, talk at the Eleventh International Workshop on Weak Interactions, Santa Fe, NM, 1987 (unpublished); G. Gelmini, L. Hall, and M. Lin, Nucl. Phys. **B281**, 726 (1987).

[2]J. N. Bahcall, S. T. Petcov, S. Toshev, and J. W. F. Valle, Phys. Lett. B **181**, 369 (1986).

[3]L. Okun, M. Voloshin, and M. Vytotsky, Yad. Fiz. **44**, 677 (1986) [Sov. J. Nucl. Phys. **44**, 440 (1986)].

[4]P. Mikheyev and A. Smirnov, Nuovo Cimento Soc. Ital. Fis. **9C**, 17 (1986); H. Bethe, Phys. Rev. Lett. **56**, 1305 (1986); L. Wolfenstein, Phys. Rev. D **17**, 2369 (1978).

[5]R. Mohapatra and J. W. F. Valle, Phys. Lett. B **177**, 47 (1986).

[6]S. Glashow and L. Krauss, Phys. Lett. B **190**, 199 (1987).

[7]J. W. F. Valle, Phys. Lett. **B131**, 87 (1983); G. Gelmini and J. W. F. Valle, Phys. Lett. **142B**, 181 (1984); G. Gelmini, D. Schramm, and J. W. F. Valle, Phys. Lett. **146B**, 311 (1984); J. W. F. Valle, Phys. Lett. **159B**, 49 (1985); S. Nussinov, Phys. Lett. B **185**, 171 (1987).

[8]K. Hirata *et al.*, Phys. Rev. Lett. **58**, 1490 (1987); R. Bionta *et al.*, Phys. Rev. Lett. **58**, 1494 (1987).

[9]J. Frieman, H. Haber, and K. Freese, SLAC Report No. SLAC PUB 4261, 1987 (unpublished).

[10]Detailed analysis of the range of parameters required for the resonant amplification of neutrino oscillations to take place has been given by many authors; see, e.g., W. Haxton, Phys. Rev. Lett. **57**, 1271 (1986); S. P. Rosen and J. Gelb, Phys. Rev. D **34**, 969 (1986); V. Barger *et al.* Phys. Rev. D **34**, 980 (1986).

[11]F. Boehm, private communication.

[12]R. Johnson, S. Ranfone, and J. Schechter, Phys. Lett. B **179**, 355 (1986); P. Langacker *et al.*, Nucl. Phys. **B282**, 589 (1987).

[13]J. W. F. Valle, Phys. Lett. B **186**, 73 (1987); R. Mohapatra and J. W. F. Valle, Phys. Rev. D **34**, 1642 (1986); J. P. Derendinger *et al.*, Nucl. Phys. **B267**, 365 (1986); S. Nandi and U. Sarkar, Phys. Rev. Lett. **55**, 566 (1986); F. del Aguila *et al.*, Nucl. Phys. **B272**, 413 (1986).

[14]J. W. F. Valle, in *Proceedings of the International Symposium on Weak and Electromagnetic Interactions in Nuclei, Heidelberg, West Germany, 1986,* edited by H. Klapdor (Springer-Verlag, Berlin, 1987), p. 927.

[15]A. Santamaria and J. W. F. Valle, Phys. Lett. B **195**, 423 (1987).

[16]J. W. F. Valle, Phys. Rev. D **27**, 1672 (1983); J. W. F. Valle and M. Singer, Phys. Rev. D **28**, 540 (1983); S. Petcov, Phys. Lett. **110B**, 245 (1982); L. Wolfenstein, Nucl. Phys. **B186**, 147 (1981).

[17]J. W. F. Valle, Phys. Lett. B **196**, 157 (1987).

[18]L. Hall and M. Suzuki, Nucl. Phys. **B231**, 419 (1984); I. H. Lee, Nucl. Phys. **B246**, 120 (1984).

[19]R. Mohapatra and C. Aulakh, Phys. Lett. **119B**, 136 (1983); G. G. Ross and J. W. F. Valle, Phys. Lett. **151B**, 375 (1985); J. Ellis, G. Gelmini, C. Jarlskog, G. G. Ross, and J. W. F. Valle, Phys. Lett. **150B**, 142 (1985); S. Dawson, Nucl. Phys. **B261**, 297 (1985).

[20]A. Bouquet and P. Salati, Nucl. Phys. **B284**, 557 (1987); B. Gato, J. Leon, J. P. Mercader, and M. Quiros, Nucl. Phys. **B260**, 203 (1985).

[21]G. Gelmini and M. Roncadelli, Phys. Lett. **99B**, 411 (1981).

[22]J. Schechter and J. W. F. Valle, Phys. Rev. D **21**, 309 (1980), and **22**, 2227 (1980).

[23]H. Georgi, S. Glashow, and S. Nussinov, Nucl. Phys. **B193**, 297 (1981); J. Ellis and K. Olive, Nucl. Phys. **B223**, 252 (1983), and references therein.

[24]G. Kane and W. Rolnick, Nucl. Phys. **B232**, 21 (1984); W. Buchmuller and F. Scheck, Phys. Lett. **145B**, 421 (1984); A. Santamaria, J. Bernabeu, and A. Pich, Phys. Rev. D **36**, 1408 (1987).

[25]A. Jodidio *et al.*, Phys. Rev. D **34**, 1967 (1986).