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A new three flavor oscillation solution of the solar neutrino deficit in R-parity violating supersymmetry

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Abstract

We present a solution of the solar neutrino deficit using three flavors of neutrinos within the R-parity non-conserving supersymmetric model. In vacuum, mass and mixing is restricted to the ν_μ-ν_τ sector only, which we choose in consistency with the requirements of the atmospheric neutrino anomaly. The ν_e is massless and unmixed. The flavor changing and flavor diagonal neutral currents present in the model and an energy-dependent resonance-induced ν_e-ν_μ mixing in the sun result in the new solution to the solar neutrino problem. The best fit to the solar neutrino rates and spectrum (1258-day SK data) requires a mass square difference $\sim 10^{-5}$ eV^2 in vacuum between the two lightest neutrinos. This solution cannot accommodate a significant day-night effect for solar neutrinos.

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Neutrino oscillation is the most popular solution of the solar neutrino problem [1, 2] and the atmospheric neutrino anomaly [3]. Oscillation in vacuum or in matter, through the MSW resonance mechanism, posits that neutrinos have non-vanishing, non-degenerate masses and that the basis defined by these eigenstates does not coincide with the flavor basis.

Supersymmetry (SUSY) with R-parity non-conservation is an extension of the Standard Model (SM) which is consistent with all particle physics experiments and is phenomenologically rich [4]. It carries within it new interactions between leptons and quarks which violate baryon (B) and lepton (L) number. In this work we show that the flavor changing neutral currents (FCNC) and flavor diagonal neutral currents (FDNC) due to the L-violating interactions induce mixing amongst neutrinos in matter, the key feature in this alternative solution of the solar neutrino discrepancy, even though, in vacuum, the ν_e state is massless and does not mix with the other neutrinos. We also indicate how in this model the parameters can be chosen to address the atmospheric neutrino anomaly, solving the solar neutrino problem at the same time.

Origins of neutrino oscillation other than mass-mixing, notable among them being non-standard interactions of neutrinos with matter, like FCNC, were examined by Wolfenstein [5].

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It is noteworthy that FCNC and FDNC interactions can drive neutrino oscillations even for massless neutrinos without any vacuum mixing through an energy-independent resonance effect [6]. However, these solutions require large L-violating couplings near their present experimental upper bounds [7, 8]. This has been examined earlier in connection with the solar [6, 7, 9] and atmospheric neutrino data [8, 10] in the *two* flavor oscillation framework. The new explanation of the solar neutrino deficit that we propose, in contrast, relies on an interplay between the $\not{R}$-interactions for *three* flavors of neutrinos with matter and their masses, keeping *vacuum mixing restricted to the $\nu_\mu - \nu_\tau$ sector only.*

Imposing baryon number conservation, we focus on the following L-violating terms in the superpotential:

$$W = \frac{1}{2}\lambda_{ijk}\, L_i\, L_j\, E^c_k + \lambda'_{ijk}\, L_i\, Q_j\, D^c_k, \tag{1}$$

assuming that bilinear terms have been rotated away with appropriate redefinition of superfields. Here, i,j, and k are generation indices, L and Q are chiral superfields containing left-handed lepton and quark doublets and E and D are chiral superfields containing right-handed charged-lepton and d-quark singlets. There are nine λ (antisymmetric in (ij)) and twenty-seven λ' couplings, only a few of which will be relevant for this analysis.

The interaction of neutrinos with the electrons and d-quarks in matter induces transitions (i) $\nu_i + e \to \nu_j + e$, and (ii) $\nu_i + d \to \nu_j + d$. (i) can proceed *via* W and Z exchange for $i = j$, as well as *via* λ couplings for all i, j, while process (ii) is possible through λ' couplings and squark exchange. Here we concentrate only on the λ-induced contributions.

The time evolution of the neutrino flavor eigenstates ($\nu_i,\ i = e, \mu, \tau$) is governed by

$$\begin{aligned} H &= H^0 + h^{matter} \\ &= \begin{pmatrix} E & 0 & 0 \\ 0 & E + S_+ - T_1 & T_2 \\ 0 & T_2 & E + S_+ + T_1 \end{pmatrix} + \begin{pmatrix} R_{11} + A_1 - A_2 & 0 & R_{13} \\ 0 & -A_2 & 0 \\ R_{13} & 0 & R_{33} - A_2 \end{pmatrix}, \end{aligned} \tag{2}$$

where $S_\pm = (m_3^2 \pm m_2^2)/4E$, $T_1 = S_- \cos 2\theta_{23v}$, $T_2 = S_- \sin 2\theta_{23v}$, $A_1 = \sqrt{2} G_F n_e$, $A_2 = G_F n_N/\sqrt{2}$, and $R_{ij} = \lambda_{ik1}\lambda_{jk1} n_e / 4\tilde{m}^2$. E is the neutrino energy and θ_{23v} the vacuum mixing angle in the $\nu_\mu - \nu_\tau$ sector. n_N and n_e are the neutron and electron number densities in matter and $\tilde{m}$ is the slepton mass. A_1 and A_2 in h^{matter} arise from SM charged and neutral current interactions, respectively. In vacuum, $h^{matter} = 0$ and H contains mixing only in the $\nu_\mu - \nu_\tau$ sector. In h^{matter}, we choose[1] $k = 2$ in the matter-induced contributions R_{ij}. For anti-neutrinos, the time evolution is determined by a similar total hamiltonian $\bar{H} = H^0 - h^{matter}$.

To obtain the mass eigenstates, first we rotate by $U' = U_{23} U_{13}$ (where U_{ij} is the standard rotation matrix) and write the effective mass squared matrix, $\frac{\widetilde{M}^2}{2E} = H - E - A_1 - A_2$, in the new basis as

$$\frac{\widetilde{M}^2}{2E} \approx \begin{pmatrix} R_{11} c_{13}^2 - 2 R_{13} c_{23} s_{13} c_{13} + \Lambda_+ s_{13}^2 & -R_{13} s_{23} c_{13} & 0 \\ -R_{13} s_{23} c_{13} & \Lambda_- & -R_{13} s_{23} s_{13} \\ 0 & -R_{13} s_{23} s_{13} & R_{11} s_{13}^2 + 2 R_{13} c_{23} s_{13} c_{13} + \Lambda_+ c_{13}^2 \end{pmatrix}, \tag{3}$$

[1] In view of the antisymmetry of λ_{ijk} in (i, j), in order to generate the mixing of the ν_e with the other neutrinos we have to choose $k = 2$ or 3. For the latter choice, mixings due to $\not{R}$ interactions are very small; for example, $\lambda_{131}\lambda_{231}$ is highly constrained from $\mu \to 3e$ decay [4].

where

$$\Lambda_\pm = \left[S_+ - A_1 + \frac{R_{33}}{2}\right] \pm \left[S_- \cos 2(\theta_{23\nu} - \theta_{23}) + \frac{R_{33}}{2}\cos 2\theta_{23}\right] \tag{4}$$

and $c_{ij} \equiv \cos\theta_{ij}$ and $s_{ij} \equiv \sin\theta_{ij}$. Furthermore,

$$\tan 2\theta_{23} = 2T_2/(2T_1 + R_{33}); \quad \tan 2\theta_{13} = 2R_{13}c_{23}/D_1; \quad D_1 = \Lambda_+ - R_{11}. \tag{5}$$

Note that $\theta_{23} \approx \theta_{23v}$ while[2] $\theta_{13} \approx 0$ except near a possible resonance, when $D_1 = 0$. We show below that this resonance condition cannot be achieved in the sun. Consequently, to a good approximation, the third state in this basis decouples in eq. (3). The upper left 2×2 block is readily diagonalised, resulting in three effective masses $\tilde{m}_i$ as:

$$\tilde{m}_1^2/(2E) = c_{12}^2\left(R_{11}c_{13}^2 - R_{13}c_{23}\sin 2\theta_{13} + \Lambda_+ s_{13}^2\right) + R_{13}s_{23}c_{13}\sin 2\theta_{12} + \Lambda_- s_{12}^2$$

$$\tilde{m}_2^2/(2E) = s_{12}^2\left(R_{11}c_{13}^2 - R_{13}c_{23}\sin 2\theta_{13} + \Lambda_+ s_{13}^2\right) - R_{13}s_{23}c_{13}\sin 2\theta_{12} + \Lambda_- c_{12}^2$$

$$\tilde{m}_3^2/(2E) = R_{11}s_{13}^2 + R_{13}c_{23}\sin 2\theta_{13} + \Lambda_+ c_{13}^2, \tag{6}$$

where

$$\tan 2\theta_{12} = \frac{-2R_{13}s_{23}c_{13}}{D_2}; \quad D_2 = \Lambda_- - R_{11}c_{13}^2 + R_{13}c_{23}\sin 2\theta_{13} - \Lambda_+ s_{13}^2. \tag{7}$$

A resonant enhancement of θ_{12} occurs when $D_2 = 0$.

The neutrino flavor eigenstates $\nu_\alpha = \nu_{e,\mu,\tau}$ are related to the mass eigenstates $\nu_i = \nu_{1,2,3}$ by

$$\nu_\alpha = \sum_i U_{\alpha i}\,\nu_i, \tag{8}$$

where $U_{\alpha i}$ are elements of the unitary mixing matrix

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13} & c_{12}c_{23} - s_{12}s_{23}s_{13} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13} & -c_{12}s_{23} - s_{12}c_{23}s_{13} & c_{23}c_{13} \end{pmatrix}. \tag{9}$$

We have chosen real L violating couplings and as such there is no CP violating phase in the above mixing matrix. Further, in order to satisfy $0 \leq \theta_{12} \leq \pi/2$ in eq. (9) for convenience, we take $\lambda_{121}\lambda_{321} < 0$.

As noted above, level crossings and resonance behavior, which are energy dependent due to neutrino masses, can occur in two situations, namely, (a) when $D_1 = 0$, and (b) when $D_2 = 0$. Of these, only the latter can be satisfied inside the sun, as we now discuss. The sub-GeV and multi-GeV zenith angle dependence of atmospheric neutrinos as well as the energy dependence of the up-down asymmetry require $\Delta m_{32} \approx m_3^2 \approx 10^{-3}$ eV^2 with maximal vacuum mixing in the $\nu_\mu - \nu_\tau$ sector [3]. The presence of L violating interactions does not alter this significantly (see later). On the other hand, n_e at the core of the sun is about 1.13×10^{12} eV^3. Thus even for E as high as 20 MeV, it is not possible to satisfy the (a) resonance condition and hence we consider only the (b) resonance in the subsequent

[2]This follows as $\Lambda_+ - R_{11} \sim m_3^2/(2E)$ is very large with respect to R_{13} in the sun.

discussion of the solar neutrino data. At resonance, $\theta_{12} = \pi/4$, while the other mixing angles are $\theta_{13} \sim 0$ and $\theta_{23} = \theta_{23v}$. Recall that away from resonance, $\theta_{12} \sim 0$ and for vacuum propagation only $\theta_{23} = \theta_{23v}$ is non-zero in eq. (9). At first glance, one might think that if U_{13} in vacuum is very small then solar neutrinos will be almost unaffected by the mass of ν_τ and analysis with three neutrino flavors may not be essential. However, unlike in the SM, where only ν_e interactions with matter are relevant for neutrino oscillation, in the $\not{R}$ supersymmetric Model, FCNC and FDNC interactions of all three flavors of neutrinos turn out to be important. In fact, one can see from eqs. (2) and (7) that R_{13}, arising from FCNC interactions, appears in $\tan 2\theta_{12}$ and plays a pivotal role.

We now turn to the oscillation of solar neutrinos due to their interaction with matter inside the sun. As already discussed, ν_e in the sun can experience only one of the two resonances. s_{13} in eq. (9) is very small as noted earlier and we use the survival probability of ν_e valid for a two flavor analysis:

$$P_{\nu_e \to \nu_e} = \frac{1}{2} + \left(\frac{1}{2} - P_{jump}\right) \cos 2\theta_{12}(x_1) \cos 2\theta_{12}(x_0), \tag{10}$$

where x_0 is the production point inside the sun and x_1 the detection point at earth[3]. The jump probability is $P_{jump} \approx \exp[-\pi\gamma_{res}F/2]$, γ_{res} being the adiabaticity parameter. $F = 1$ for the exponential density profile since the vacuum mixing angle is zero and

$$\gamma_{res} = \frac{\tilde{m}_2^2 - \tilde{m}_1^2}{4E\dot{\theta}_{12}} \simeq \frac{m_2^2}{E}\left(\frac{p}{\kappa}\right)^2 \left(\frac{n_e}{\dot{n}_e}\right)_{res}, \tag{11}$$

where $\kappa = (2\lambda_{121}^2 - \lambda_{321}^2)/(8\tilde{m}_\mu^2) + \sqrt{2}G_F \sim \sqrt{2}G_F$ [4] and $p = |\frac{\lambda_{121}\lambda_{321}}{4m_{\tilde{\mu}}^2}|$.

In order to obtain the best-fit values of $\Delta m_{12} \approx m_2^2$ and p, we have performed a χ^2 analysis using the Standard Solar Model (SSM) [11] and the solar neutrino rates from the Homestake (Cl), Gallex, Sage, and Kamiokande (K) experiments [2]. We have also used the latest SK rates and spectrum data for 1258 days [12]. Taking into account the production point distributions of neutrinos from the different reactions (*e.g.*, *pp*, *pep*, ^7Be, ^8B *etc.*)[5], we have calculated the averaged survival probabilities using eq. (10). Here, we include a parameter X_B to take into account a possible deviation of the overall normalization of the ^8B flux from its SSM value. We set $\theta_{23v} = \pi/4$. The best-fit values of the parameters are presented in Table 1 along with χ^2_{min}, the goodness of fit (gof), and the calculated rates using these values of the parameters. Case (1) is a fit to the total rates. Note that the best-fit parameters result in an unusually good fit to the Cl rate and the Ga prediction is right near the average of the Sage and Gallex data. We have found that the fit improves even more if the K-rate is excluded. In case (2) we have fitted the SK spectrum while (3) is a fit to the total rates and the SK spectrum[6]. In Fig. 1 is shown the calculated spectrum for SK for the best-fit parameters along with the experimental data. Also shown is the charged current spectrum expected at SNO for one sample case, the best-fit values in case (3)[7].

[3] Notice that $\cos 2\theta_{12}(x_1) = 1$, corresponding to $\theta_{12} = 0$ in vacuum.

[4] λ_{121} and λ_{321} are tightly constrained [4]. Besides significant cancellation between these terms is possible if they are of same order.

[5] We have dropped a small contribution from the *hep* process.

[6] We have checked that the fit (3) is essentially unchanged if the SK rate is excluded from the fit.

[7] For these best-fit values, the prediction for the neutral and charged current rates at SNO, normalized to the SSM, are 0.56 and 0.44, respectively.

Case	Best-fit Values					Corresponding Rates				
	p (10^{-24} eV^{-2})	m_2^2 (10^{-5} eV^2)	X_B	dof	χ^2 (gof)	Cl (0.33± 0.029)	$Gallex$ (0.52± 0.06)	$Sage$ (0.60± 0.06)	K (0.54± 0.07)	SK (0.451± 0.016)
1	0.595	1.063	0.845	2	1.71 (42.5)	0.326	0.561	0.561	0.478	0.455
2	0.009	0.01	0.446	16	18.76 (28.1)	0.582	0.947	0.947	0.446	0.446
3	0.360	0.980	0.560	21	25.53 (22.5)	0.364	0.558	0.558	0.456	0.447

Table 1: The best-fit values of the parameters, $p = |\frac{\lambda_{121}\lambda_{321}}{4m^2_{\tilde{\mu}}}|$, m_2^2, and X_B from fits to (1) all rates, (2) the SK spectrum, and (3) rates and SK spectrum. The rates for the different experiments obtained using these best-fit parameters are also shown.

The best-fit values of $\not{R}$ couplings in Table 1 are consistent with the existing constraints. For example, in case (3), choosing $m_{\tilde{\mu}} \sim 100$ GeV, we get $\lambda_{121}\lambda_{321} \approx 0.0144$. λ_{121} is constrained from $\mu \to e\bar{\nu}_e\nu_\mu$ decay (with selectron exchange tree level diagram apart from the SM W exchange diagram). The bound on λ_{321} is from $R \equiv \Gamma(\tau \to e\nu\bar{\nu})/\Gamma(\tau \to \mu\nu\bar{\nu})$ which gets a contribution from a selectron exchange diagram. For $m_{\tilde{e}} \sim 200$ GeV or more, the requirements are easily satisfied [4].

Turning now to atmospheric neutrinos [3], for a simple-minded analysis we can consider the earth to be a slab of a single density. n_e in earth lies in the range $(3-6)N_A$ cm^{-3}. So the resonance condition, $D_2 = 0$, cannot be met for atmospheric neutrinos having energy near the GeV range. In order to explain the observed zenith angle dependence, we must choose $\Delta m_{32} \sim 10^{-3}$ eV^2. This precludes the occurrence of the other resonance, $D_1 = 0$. Since neither resonance condition can be satisfied, there will be almost no effect on atmospheric neutrino oscillation due to the L violating interactions as the associated couplings are very small. So one can consider the mixing matrix in eq. (9) valid for vacuum for which only θ_{23v} is non-zero. Thus the solution to the atmospheric neutrino anomaly is just the standard two neutrino mass-mixing one.

The neutrino masses and mixing pattern in vacuum required in this solution can naturally arise in many models. For example, the trilinear couplings in eq. (1) contribute to the neutrino mass matrix at the one-loop level through slepton or squark exchange diagrams [13]. In particular, from the λ' couplings one obtains:

$$m_{ij}^{loop} = \frac{3\, m_b^2\, (A_b + \mu \tan\beta\,)}{8\pi^2\, \tilde{m}_b^2}\, \lambda'_{i33}\, \lambda'_{j33}. \tag{12}$$

where A_b and μ are soft SUSY-breaking parameters, $\tilde{m}_b$ is the b-squark mass and $\tan\beta$ is the ratio of two Higgs vacuum expectation values. The last two generation indices in λ' have been chosen as 3 for which the loop contributions are enhanced *via* the $b-$quark mass. We remark that m_{ij} is very small when $i = 1$ and/or $j = 1$ because of the more stringent constraint [4] on λ'_{133}. Notice that this mass matrix can correspond to almost maximal mixing for ν_μ and ν_τ if $\lambda'_{233} \approx \lambda'_{333}$, with two neutrino masses very small and one neutrino having significantly higher mass $m_3 \approx 2\, m_{33}^{loop}$, which can be suitably chosen by taking appropriate values of the different parameters in (12). It should be borne in mind that m_2 depends on the difference of λ'_{233} and λ'_{333} and can be several orders less than the mass of the heavier neutrino while there will be almost maximal mixing. The remaining neutrino mass is $m_1 \approx 0$. Thus masses and vacuum mixings can be as required in the model under consideration.

This neutrino mixing pattern also satisfies the bound ${U_{13}}^2 \leq 0.04$ in vacuum from the CHOOZ reactor experiment [14]. In fact, in vacuum ${U_{13}}^2 = 0$.

A comment about the earth regeneration effect for solar neutrinos is pertinent. The ν_e is unmixed with the other neutrinos in vacuum. As n_e in earth is about two orders less than that near the core of the sun, no resonance condition will be satisfied[8]. Hence, there will not be an earth effect for solar neutrinos. In comparison with the small angle MSW fits [15], the somewhat larger best-fit Δm_{12} and the zero value of θ_{12} in vacuum here, result in a smaller day-night effect.

Though our discussion has been within the framework of R-parity violating SUSY, there are other models [16] where FCNC and FDNC interactions are present. Our results can be adapted to these scenarios in a straight-forward manner.

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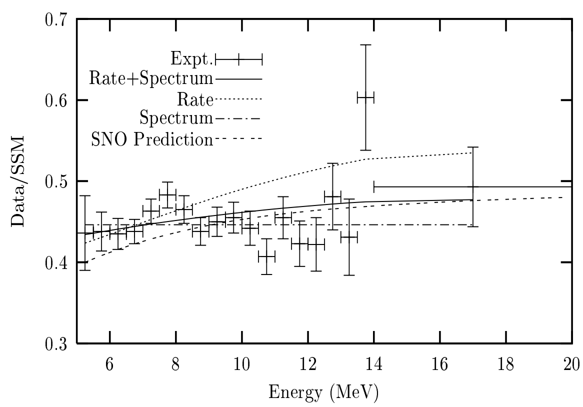

Figure 1: The calculated SK solar neutrino spectrum for the best-fit parameters Δm_{12}, p, and X_B from (1) rates, (2) SK spectrum, and (3) rates and spectrum. The SK 1258-day data [12] and the predicted SNO charged current spectrum for the best-fit (3) are also shown.