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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 and R-parity non-conserving supersymmetry. In this model, in vacuum, the ν_e is massless and unmixed, mass and mixing being restricted to the ν_μ–ν_τ sector only, which we choose in consistency with the requirements of the atmospheric neutrino anomaly. The flavor changing and flavor diagonal neutral currents present in the model and the three-flavor picture together produce an energy dependent resonance-induced ν_e–ν_μ mixing in the sun. This mixing plays a key role in the new solution to the solar neutrino problem. The best fit to the solar neutrino rates and spectrum (1258-day SK and 241-day SNO data) requires a mass square difference of $\sim 10^{-5}\,\mathrm{eV}^2$ in vacuum between the two lightest neutrinos. This solution cannot accommodate a significant day–night effect for solar neutrinos nor CP violation in terrestrial neutrino experiments.

Neutrino oscillation is the most popular solution of the solar neutrino problem [1–4] and the atmospheric neutrino anomaly [5]. The assumption of oscillations 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.

Origins of neutrino oscillation different from mass mixing are also possible, prominent among them being non-standard interactions of neutrinos with matter [6]. It is noteworthy that flavor changing neutral current (FCNC) and flavor diagonal neutral current (FDNC) interactions can drive neutrino oscillations even for massless neutrinos with no vacuum mixing [7]. This possibility has been examined earlier in connection with the solar [7–9] and atmospheric neutrino data [10,11] in the *two*-flavor oscillation framework. In contrast, in the new explanation of the solar neutrino deficit that we propose, only the ν_e is massless in vacuum. The mass and mixing in the ν_μ–ν_τ sector is determined by the oscillation solution to the atmospheric neutrino anomaly. An interplay between the R-parity non-conserving supersymmetry (SUSY) interactions for *three* flavors of neutrinos with matter and these masses is the novel ingredient in addressing the solar neutrino deficit, *vacuum mixing being restricted to the* ν_μ–ν_τ *sector only.*

Supersymmetry with R-parity non-conservation ($\not{R}$) carries within it new interactions between leptons and quarks which violate baryon (B) and lepton (L) number. In the model discussed in this paper, even though in vacuum the ν_e state is massless and does not mix with the other neutrinos, the L-violating FCNC and FDNC interactions induce mixing amongst neutrinos in solar matter. This is the key feature in this solution of the solar neutrino discrepancy. We find that the vacuum mass and mixing chosen to address the atmospheric neutrino anomaly are just right to explain the solar neutrino problem via this induced-mixing mechanism.

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 an appropriate redefinition of the 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 the transitions

(i) $\nu_i + e \to \nu_j + e$, and
(ii) $\nu_i + d \to \nu_j + d$.

(i) can proceed via the W for $i = j = e$ and the Z for $i = j$, as well as through λ couplings for all i, j, while process (ii) is possible through λ' couplings. Here we concentrate only on the λ-induced contributions.

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$$\frac{\widetilde{M}^2}{2E} \approx \begin{pmatrix} R_{11}c_{13}^2 - 2R_{13}c_{23}s_{13}c_{13} + \Lambda_+ s_{13}^2 & -R_{13}s_{23}c_{13} & 0 \\ -R_{13}s_{23}c_{13} & \Lambda_- & -R_{13}s_{23}s_{13} \\ 0 & -R_{13}s_{23}s_{13} & R_{11}s_{13}^2 + 2R_{13}c_{23}s_{13}c_{13} + \Lambda_+ c_{13}^2 \end{pmatrix} \quad (4)$$

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

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

where $S_\pm = (m_3^2 \pm m_2^2)/4E$, $T_1 = S_- \cos 2\theta_{23v}$, $T_2 = S_- \sin 2\theta_{23v}$, $A_1 = 2^{1/2} G_{\rm F} n_e$, $A_2 = G_{\rm F} n_N / 2^{1/2}$, E is the neutrino energy and θ_{23v} the vacuum mixing angle in the ν_μ–ν_τ sector. n_N and n_e are the neutron and electron number densities in matter. We have

$$R_{ij} = \lambda_{ik1}\lambda_{jk1} n_e / 4\tilde{m}^2. \quad (3)$$

$\tilde{m}$ is the exchanged slepton mass. The second index on the λs, k, identifies the exchanged slepton. Since the interaction of neutrinos with the electrons in the sun is of relevance here, the third index is always 1.

In vacuum, $h^{\text{matter}} = 0$ and H contains mixing only in the ν_μ–ν_τ sector. In h^{matter}, A_1 and A_2 result from SM charged and neutral current interactions, respectively, of which the latter is flavor independent. The R_{ij} arise from the $\not{R}$ interactions. In this model, non-zero mixing of the other neutrinos with the ν_e (i.e., $i = 1$ in (3)) is generated in the sun. In view of the antisymmetry of λ_{ik1} in (i, k), k can only be 2 or 3. For the latter choice, the existing constraints on the $\not{R}$ couplings result in mixings which are very small; for example, $\lambda_{131}\lambda_{231}$ is highly constrained from $\mu \to 3e$ decay [12]. $k = 2$ is therefore the preferred choice. This also entails $R_{12} = R_{21} = R_{23} = R_{32} = 0$. For anti-neutrinos, the time evolution is determined by a similar total hamiltonian $\bar{H} = H^0 - h^{\text{matter}}$.

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, $\widetilde{M}^2/(2E) = H - E - A_1 + A_2$, in the new basis as (see (4) on top of the page), where

$$\begin{aligned} \Lambda_\pm &= \left[S_+ - A_1 + \frac{R_{33}}{2} \right] \\ &\quad \pm \left[S_- \cos 2(\theta_{23\nu} - \theta_{23}) + \frac{R_{33}}{2} \cos 2\theta_{23} \right], \end{aligned} \quad (5)$$

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

$$\begin{aligned} &\tan 2\theta_{23} = 2T_2/(2T_1 + R_{33}), \quad \tan 2\theta_{13} = 2R_{13}c_{23}/D_1, \\ &D_1 = \Lambda_+ - R_{11}. \end{aligned} \quad (6)$$

Since $R_{33} \ll T_1$, $\theta_{23} \approx \theta_{23v}$ while $\theta_{13} \approx 0$ except near a possible resonance, when $D_1 = 0$. We show below that this resonance condition cannot be achieved, essentially because $\Lambda_+ \simeq m_3^2/(2E) - A_1$, which is much larger than the values which R_{11} can assume in the sun. Consequently, to a good approximation, the third state in this basis decouples in (4). The upper left 2×2 block is readily diagonalized, resulting in three effective masses $\tilde{m}_i$:

$$\begin{aligned} \tilde{m}_1^2/(2E) &= c_{12}^2 \left(R_{11}c_{13}^2 - R_{13}c_{23}\sin 2\theta_{13} + \Lambda_+ s_{13}^2 \right) \\ &\quad + R_{13}s_{23}c_{13}\sin 2\theta_{12} + \Lambda_- s_{12}^2, \\ \tilde{m}_2^2/(2E) &= s_{12}^2 \left(R_{11}c_{13}^2 - R_{13}c_{23}\sin 2\theta_{13} + \Lambda_+ s_{13}^2 \right) \\ &\quad - R_{13}s_{23}c_{13}\sin 2\theta_{12} + \Lambda_- c_{12}^2, \\ \tilde{m}_3^2/(2E) &= R_{11}s_{13}^2 + R_{13}c_{23}\sin 2\theta_{13} + \Lambda_+ c_{13}^2, \end{aligned} \quad (7)$$

where

$$\begin{aligned} \tan 2\theta_{12} &= \frac{-2R_{13}s_{23}c_{13}}{D_2}, \\ D_2 &= \Lambda_- - R_{11}c_{13}^2 + R_{13}c_{23}\sin 2\theta_{13} - \Lambda_+ s_{13}^2. \end{aligned} \quad (8)$$

A resonant enhancement of θ_{12} occurs when $D_2 = 0$. This resonance, occurring inside the sun, is crucial to explain the solar neutrino deficit, as we discuss now.

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

$$\nu_\alpha = \sum_i U_{\alpha i}\nu_i, \quad (9)$$

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

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13} & c_{12}c_{23} - s_{12}s_{23}s_{13} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13} & -c_{12}s_{23} - s_{12}c_{23}s_{13} & c_{23}c_{13} \end{pmatrix}. \quad (10)$$

We have chosen real L violating couplings, so there is no CP violating phase in the above mixing matrix. Further, in order to satisfy $0 \le \theta_{12} \le \pi/2$ in (10) 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 $\Delta m_{32} \approx m_3^2 \approx 10^{-3}\,\text{eV}^2$ with maximal vacuum mixing in the ν_μ–ν_τ sector [5]. The presence of L violating interactions does not alter this significantly (see below).

On the other hand, n_e at the core of the sun is about $1.13 \times 10^{12}\,\mathrm{eV}^3$. Thus even for E as high as 20 MeV, it is not possible to satisfy the resonance condition (a) and hence we consider only the resonance (b) in the subsequent discussion of the solar neutrino data. The condition for this latter resonance involves $\Lambda_- \simeq m_2^2/(2E) - A_1$ (see (8)). The preferred values of m_2 are obtained below from fits to 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 (10). 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 an 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 oscillations, 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 (2) and (8) 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, in the sun ν_e can experience only one of the two resonances. s_{13} in (10) is very small as noted earlier and we use the survival probability of ν_e valid for a two-flavor analysis:

$$P_{\nu_e\to\nu_e} = \frac{1}{2} + \left(\frac{1}{2} - P_{\mathrm{jump}}\right)\cos 2\theta_{12}(x_1)\cos 2\theta_{12}(x_0), \quad (11)$$

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

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

where $\kappa = (2\lambda_{121}^2 - \lambda_{321}^2)/(8\tilde{m}_\mu^2) + 2^{1/2}G_{\mathrm{F}} \sim 2^{1/2}G_{\mathrm{F}}$ [2] and $p = |\lambda_{121}\lambda_{321}/(4m_{\tilde{\mu}}^2)|$.

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) [13] and the solar neutrino rates from the Homestake (Cl), Gallex, Sage, GNO, and Kamiokande (K) experiments [2]. We have used the latest SK rates and spectrum data for 1258 days [3] and the SNO 241-day charged current spectral data [4]. The definition of χ^2 used for this analysis is

$$\chi^2 = \sum_{i,j=1,n}\left(F_i^{\mathrm{th}} - F_i^{\mathrm{exp}}\right)(\sigma_{ij}^{-2})\left(F_j^{\mathrm{th}} - F_j^{\mathrm{exp}}\right). \quad (13)$$

Here $F_i^\xi = T_i^\xi/T_i^{\mathrm{BP01}}$ where ξ is th (for the theoretical prediction of our model) or exp (for the experimental value). For the rates part, $n = 7$ and T_i is the total rate in the ith experiment. For the spectrum part, $n = 49$ corresponding to 38 SK day–night bins and 11 SNO CC bins; here T_i is the number of events in the ith energy bin. The error matrix σ_{ij} contains the experimental errors, the theoretical errors and their correlations. For evaluating the error matrix for the total rates and spectrum we use the procedure described in [14]. For the combined analysis of rates and spectrum we define the total χ^2 as

$$\chi^2_{\mathrm{total}} = \chi^2_{\mathrm{rate}} + \chi^2_{\mathrm{sp}}. \quad (14)$$

Taking into account the production point distributions of neutrinos from the different reactions (e.g., *pp*, *pep*, ^{7}Be, ^{8}B etc.)[3], we have calculated the averaged survival probabilities using (11). 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. In order to check the sensitivity of the result, 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. For every case, we present two sets of results:

(a) holding X_{B} fixed at unity, and
(b) letting X_{B} vary within its 3σ allowed range.

The experimental errors are taken at 1σ. Case (1) is a fit to the total rates. In case (2) we have fitted the SK and SNO measured spectra, while (3) is a fit to the total rates and the two spectra[4]. It can be seen from Table 1 that the best-fit parameters to the spectra alone – (2a) and (2b) – predict Cl and Ga rates inconsistent with the observations. This is because the best-fit parameters, especially m_2^2, obtained from the fit to the total rates – case 1 – differs significantly from those from the fits to the spectra – case 2. The rates and spectra fit – case 3 – is closer to the rates fit.

In Fig. 1 is shown the calculated spectrum for (a) SK electron scattering and (b) SNO charged current scattering for the best-fit parameters in Table 1 (for $X_{\mathrm{B}} = 1$) along with the experimental data.

Are the best-fit values of $\not R$ couplings in Table 1 consistent with the existing constraints? The constraints on these couplings depend on the mass of the selectron while the variable p involves $m_{\tilde{\mu}}$. For example, λ_{121} is constrained by $\mu \to e\bar{\nu}_e\nu_\mu$ 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. If $m_{\tilde{\mu}}$ and $m_{\tilde{e}}$ are degenerate then only case (2b) can be accommodated. The other cases require various extents of mass splitting between $\tilde{\mu}$ and $\tilde{e}$. For example, for case (3a), choosing $m_{\tilde{\mu}} \sim 100$ GeV, we get $\lambda_{121}\lambda_{321} \approx 0.0259$. For $m_{\tilde{e}} \sim 275$ GeV or more, the requirements are satisfied [12].

Turning now to atmospheric neutrinos [5], for a simple-minded analysis we can consider the earth to be a slab of a

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

[2] λ_{121} and λ_{321} are tightly constrained [12]. Besides, significant cancellation between these terms is possible if they are of same order

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

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

Table 1. The best-fit values of the parameters, $p = |\lambda_{121}\lambda_{321}/(4m_{\tilde{\mu}}^2)|$ and m_2^2 from fits to (1) all rates, (2) the SK and SNO spectra, and (3) rates and SK and SNO spectra. For every case, the ^{8}B flux normalization, X_{B}, has been (a) held fixed at the SSM value, and (b) varied within its 3σ allowed range. The rates obtained using these best-fit parameters are also shown

	Best-fit values					Corresponding rates			
Case	p (10^{-24} eV^{-2})	m_2^2 (10^{-5} eV^2)	X_{B}	χ^2/dof	gof	Cl (0.33± 0.029)	Ga (0.58± 0.04)	SK (0.451± 0.016)	SNO (0.347± 0.029)
1a	0.677	1.257	1.0	6.64/5	24.87	0.276	0.602	0.410	0.361
1b	0.729	1.348	1.33	4.99/4	28.88	0.291	0.618	0.462	0.376
2a	0.223	11.96	1.0	54.14/47	22.08	0.513	0.940	0.457	0.360
2b	0.10	19.24	0.52	44.98/46	51.49	0.554	0.946	0.446	0.407
3a	0.647	1.213	1.0	75.12/54	3.02	0.308	0.599	0.447	0.407
3b	0.335	1.026	0.541	65.63/53	11.42	0.365	0.589	0.442	0.441

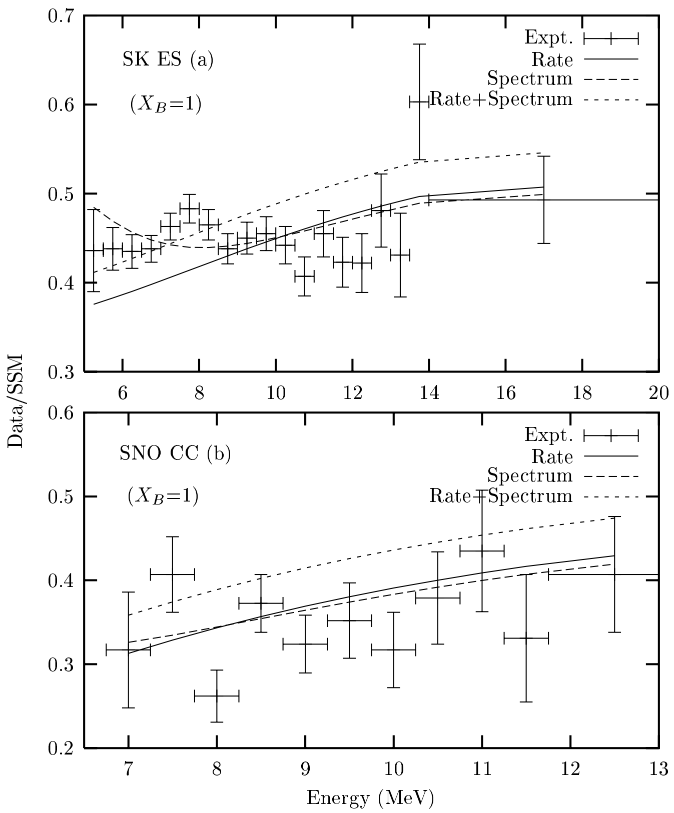

Fig. 1a,b. The calculated SK electron scattering spectrum **a** and the SNO charged current spectrum **b** for the best-fit parameters Δm_{12} and p, for cases (1a) rates, (2a) SK spectrum, and (3a) rates and spectrum. All these cases correspond to a ^{8}B flux normalization as in the SSM (i.e. $X_{\mathrm{B}} = 1$). The SK 1258-day data [3] and the 241-day SNO results [4] are also shown

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. Thus earth matter effects are insignificant and 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 the atmospheric neutrino oscillation due to the L violating interactions as the associated couplings are very small. Thus the mixing matrix in (10) valid for vacuum, for which only θ_{23v} is non-zero, is the appropriate one for this discussion and the solution to the atmospheric neutrino anomaly is 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 (1) contribute to the neutrino mass matrix at the one-loop level through slepton or squark exchange diagrams [15]. In particular, from the λ' couplings one obtains

$$m_{ij}^{\mathrm{loop}} = \frac{3m_b^2(A_b + \mu \tan\beta)}{8\pi^2 \tilde{m}_b^2} \lambda'_{i33}\lambda'_{j33}, \qquad (15)$$

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, even in the absence of any other symmetries, m_{ij} is very small when $i = 1$ and/or $j = 1$ because of the more stringent constraint [12] on λ'_{133}. If electron number symmetry is imposed on the λ' couplings then, in fact, the exact texture of H^0 (see (2)) is reproduced. 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, $m_3 \approx 2m_{33}^{\mathrm{loop}}$, which can be suitably chosen by taking appropriate values of the different parameters in (15). It should be borne in mind that m_2 depends on the difference of λ'_{233} and λ'_{333} and can easily be an order or two 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 [16]. 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[5]. Hence, there will not be an earth effect for solar neutrinos. In comparison with the small angle MSW fits [17], the somewhat larger best-fit Δm_{12} and the zero value of θ_{12} in vacuum here result in a smaller day–night effect.

How does the present analysis differ from previous work? The common practice so far has been to discuss FCNC effects in a *two*-flavor picture with massless neutrinos. This results in a resonance which is *energy independent*. The zenith angle dependence of the atmospheric neutrino data disfavors this scenario with an energy independent oscillation probability [18]. In this work, to address the atmospheric neutrino data, no new idea is invoked and massive neutrinos with appropriate mixing are chosen. We show for the first time that, in a three-generation picture, this permits a new solution to the solar neutrino problem in terms of $\not{R}$ interactions. The solar neutrino analyses within the framework of R-parity violating SUSY (or more generally in terms of FCNC and FDNC interactions) have also been made within a two-generation picture with massless neutrinos (see, e.g., [19]). In such a situation, the λ couplings result in a mass matrix where all entries are proportional to the product of $n_e(r)$ and E. Consequently, the mixing angle within the sun and the resonance are both energy independent and also independent of the position in the sun, r, and is determined by the relative strengths of the λ couplings and G_{F}. Such a scenario fails to satisfactorily meet the data. The λ' couplings also result in an energy independent oscillation probability but the resonance occurs at a fixed point in the sun determined by the electron and nucleon densities and the mixing angle is dependent on the production point. A solution to the solar neutrino problem is still possible within this latter framework for ν_e–ν_τ oscillations, ν_e–ν_μ being excluded by the existing bounds on the relevant $\not{R}$ couplings. This picture of solar neutrino oscillations undergoes a drastic modification, as we have shown, if only the ν_e is massless (and unmixed) in a three-flavor picture where the other neutrinos carry non-zero mass. The resonance is energy dependent here, ν_e–ν_μ oscillations are responsible, and the λ couplings are used.

Although some of the features of this model are reminiscent of the small angle MSW solution, there are testable differences. For example, in this model CP violation in the leptonic sector is absent and neutrinoless double beta decay is forbidden, whereas for the small angle MSW solution both possiblities exist.

Our discussion has been given within the framework of R-parity violating SUSY, but there are other models [20] where FCNC and FDNC interactions are present. Our results can be adapted to these scenarios in a straightforward manner.

In summary, the novelty of this model is that

(a) it is a three-generation analysis,
(b) the atmospheric anomaly is explained through mass mixing in the ν_μ–ν_τ sector,
(c) the non-zero ν_μ mass and the R-parity violating interactions add a distinctive energy dependence to the solar ν_e survival probability, even though the ν_e is massless and unmixed in vacuum.

The survival probability is such that the observed solar neutrino rates and spectral data can be explained.

[5] Except for neutrinos with $E > 10\,\mathrm{MeV}$ passing very near the center of the earth

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