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Supernova Bounds on Supersymmetric R -parity Violating Interactions

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Abstract

We re-examine resonant massless-neutrino conversions in a dense medium induced by flavour changing neutral current (FCNC) interactions. We show how the observed $\bar{\nu}_e$ energy spectra from SN1987a and the supernova r -process nucleosynthesis provide constraints on supersymmetric models with R parity violation, which are much more stringent than those obtained from the laboratory. We also suggest that resonant massless-neutrino conversions may play a positive role in supernova shock reheating. Finally, we examine the constraints on explicit R -parity-violating FCNCs in the presence of non-zero neutrino masses in the eV range, as indicated by present hot dark matter observations.

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1 Introduction

Under suitable circumstances neutrinos can oscillate in the presence of matter [1] or undergo resonant conversions [2] even when they are strictly massless. In some models even unmixed neutrinos can resonantly convert in matter [3, 4]. Massless-neutrino resonant conversions are distinct from the usual MSW conversions [1, 5] in that they are independent of neutrino energy and affect simultaneously neutrinos as well as anti-neutrinos. For this reason this mechanism is expected to play an important role in supernova physics [2, 6, 7]. Moreover, in some of these models with flavour changing neutrino neutral current (FCNC) interactions with matter constituents it has been suggested that, for a certain range of the corresponding parameters, they may account for the observed deficit of solar neutrinos [3, 4, 8].

The required ingredients can naturally emerge in the context of various models beyond the standard model [9]. In particular, in this paper we consider this type of interactions mediated by the scalar partners of quarks and leptons in supersymmetric extensions of the standard model with explicitly broken R parity [10, 11].

The presence of R parity breaking interactions induce resonant neutrino conversions of the type $\nu_e \leftrightarrow \nu_\alpha$ as well as $\bar{\nu}_e \leftrightarrow \bar{\nu}_\alpha$. Such conversions have important implications for the supernova r -process nucleosynthesis [12] as well as the observed $\bar{\nu}_e$ energy spectra from SN1987a [13, 14, 15].

In a recent work [7], we have investigated the constraints on massless neutrino resonant conversions that follow from supernova considerations. In the present paper we apply the same considerations in order to constrain models with explicit R parity violating supersymmetric interactions which can effectively induce resonant conversions even when neutrino masses are neglected. We also suggest that resonant massless-neutrino conversion may play a positive role in supernova shock reheating. In addition, we generalize this approach in order to include the possibility of non-zero neutrino masses. These are typically expected to arise in these models and could help to explain present observations. We derive the corresponding constraints on flavour changing neutral current couplings generated by explicit R parity violating interactions.

The paper is structured as follows. In Sect. 2 we briefly present the form of the FCNC and flavour diagonal neutral current (FDNC) interactions emerging from the R parity violating terms and the new effective neutrino evolution Hamiltonian in matter. In particular we consider two possible scenarios:

1. massless and unmixed neutrinos ($\delta m^2 = 0, \sin 2\theta = 0$) with FCNC as well as *non standard* FDNC interactions of neutrinos with matter;
2. massive neutrinos ($\delta m^2 \neq 0$) assuming negligible mixing in vacuum ($\sin 2\theta = 0$), but with FCNC interactions.

Sect. 3 is devoted to a discussion of resonant massless-neutrino conversions for supernova neutrino detection and r -process nucleosynthesis. We show how the observed $\bar{\nu}_e$ energy spectra from SN1987a and the supernova r -process nucleosynthesis place important restrictions on the parameters of R parity violating models. In sect. 4 we discuss the second scenario above and derive the corresponding restrictions. In sect. 5 we briefly suggest resonant massless-neutrino conversion as a way to power supernova shock reheating. Finally, we summarize our results and conclude in Sect. 6.

2 The MSW effect with FCNC interactions

R parity is a quantum number which is +1 for all standard particles and -1 for the superpartners. It is directly related to the baryon (B) and lepton (L) number as $R = (-1)^{3B+L+2S}$, where S is the particle spin. In the Minimal Supersymmetric Standard Model (MSSM) [16] the R discrete symmetry is imposed to enforce the L and B number conservation and no tree-level flavour changing interactions exist. However no fundamental principle precludes the possibility to violate these symmetries [10, 11]. Within the particle content of the MSSM R parity can be broken explicitly by renormalizable (and hence *a priori* unsuppressed) operators. The following extra L violating couplings in the superpotential are directly relevant for neutrino propagation through matter:

$$\lambda_{ijk} L_i L_j E_k^c \tag{1}$$

$$\lambda'_{ijk} L_i Q_j D_k^c \tag{2}$$

where L, Q, E^c and D^c are (chiral) superfields which contain the usual lepton and quark $SU(2)$ doublets and singlets, respectively, and i, j, k are generation indices. In the next we focus only on the second term eq. (2) because the first is much more constrained by experimental data. Note that the simultaneous presence of the $\lambda'' U^c U^c D^c$ and $\lambda' L Q D^c$ -type couplings is very strongly constrained ($\lambda', \lambda'' \leq 10^{-10}$) from non-observation of proton decay [17, 18]. However, the constraints on λ' (see below) are rather weak in the absence of the B violating $\lambda''_{ijk} U_i^c U_j^c D_k^c$ term. We will adopt this choice throughout this paper.

The couplings in eq. (2) at low energy ($< 100\text{GeV}$) give rise to the following four-fermion effective Lagrangian for neutrinos interactions with d -quark ¹:

$$L_{eff} = -2\sqrt{2}G_F \sum_{\alpha,\beta} \xi_{\alpha\beta} \bar{\nu}_{L\alpha} \gamma^\mu \nu_{L\beta} \bar{d}_R \gamma^\mu d_R \quad \alpha, \beta = e, \mu, \tau, \quad (3)$$

where the parameters $\xi_{\alpha\beta}$ represent the strength of the effective interactions normalized to the Fermi constant G_F . For our purpose we consider explicitly the following *non standard* FDNC couplings:

$$\xi_{ee} = \sum_j \frac{|\lambda'_{1j1}|^2}{4\sqrt{2}G_F m_{\tilde{q}_{jL}}^2}, \quad (4)$$

$$\xi_{\mu\mu} = \sum_j \frac{|\lambda'_{2j1}|^2}{4\sqrt{2}G_F m_{\tilde{q}_{jL}}^2}, \quad (5)$$

$$\xi_{\tau\tau} = \sum_j \frac{|\lambda'_{3j1}|^2}{4\sqrt{2}G_F m_{\tilde{q}_{jL}}^2}, \quad j = 1, 2, 3, \quad (6)$$

and the FCNC ones:

$$\xi_{e\mu} = \sum_j \frac{\lambda'_{1j1} \lambda'_{2j1}}{4\sqrt{2}G_F m_{\tilde{q}_{jL}}^2}, \quad (7)$$

$$\xi_{e\tau} = \sum_j \frac{\lambda'_{1j1} \lambda'_{3j1}}{4\sqrt{2}G_F m_{\tilde{q}_{jL}}^2}, \quad j = 1, 2, 3, \quad (8)$$

where $m_{\tilde{q}_{jL}}$ are the masses of the exchanged squarks and $j = 1, 2, 3$ denotes $\tilde{d}_L, \tilde{s}_L, \tilde{b}_L$, respectively. These effective neutral current interactions contribute to the neutrino scattering off d quarks in matter, providing new flavour conserving as well as flavour changing terms for the matter potentials of neutrinos.

The phenomenological implications of the R parity violating couplings have been extensively studied and constraints on the coupling constants λ' from low-energy processes (charged current universality, $e - \mu - \tau$ universality, $\nu_\mu - e$ scattering, atomic parity violation) has been obtained [17]. Recently, new bounds have been derived from LEP electroweak observables to constrain λ'_{i3k} (for all i, k) and from D - decays to constrain λ'_{12k} and λ'_{22k} as well as from τ decays to restrict λ'_{31k} (for all k) (see [19] and refs. therein). In summary, the most stringent bounds on the coupling constants entering our study are the following ² (at 1

¹ For simplicity we omit in the λ' -type Yukawa couplings the terms $\lambda'_{i1k} (\nu_{iL})^c d_{1L} (\tilde{d}_{kR})^*$ ($i, k = 1, 2, 3$). However, the coupling constants λ'_{i1k} are much more constrained than λ'_{ik1} [17].

²In ref. [20] stringent bounds, $\lambda'_{113} \lambda'_{131} \leq 1.1 \times 10^{-7}$, $\lambda'_{112} \lambda'_{121} \leq 3.2 \times 10^{-5}$, $\lambda'_{111}^2 \leq 6.4 \times 10^{-5}$, are obtained from the non-observation of $0\nu\beta\beta$ decay for squark masses of 100 GeV. However, these limits suffer from some theoretical uncertainties on nuclear matrix elements.

σ level):

$$\begin{aligned}
\lambda'_{12k} &\leq 0.29, & \lambda'_{13k} &\leq 0.26, \\
\lambda'_{22k} &\leq 0.18, & \lambda'_{23k} &\leq 0.44, \\
\lambda'_{33k} &\leq 0.26, \\
\lambda'_{i1k} &\leq 0.05, & (i = 1, 2, k = 1, 2, 3)
\end{aligned}
\tag{9}$$

normalized to a 100 GeV *reference* squark mass.

The most general Schroedinger neutrino evolution equation in matter takes the form

$$i \frac{d}{dr} \begin{pmatrix} \nu_e \\ \nu_x \end{pmatrix} = \begin{pmatrix} H_e & H_{ex} \\ H_{ex} & H_x \end{pmatrix} \begin{pmatrix} \nu_e \\ \nu_x \end{pmatrix}, \quad x = \mu(\tau)
\tag{10}$$

The entries of the Hamiltonian reads as

$$H_e = V_e - \frac{\delta m^2}{2E} \cos 2\theta, \quad H_x = V_x, \quad H_{ex} = V_{ex} + \frac{\delta m^2}{4E} \sin 2\theta
\tag{11}$$

where E is the neutrino energy, δm^2 is the mass squared difference, θ is the neutrino mixing angle in vacuum and V_e, V_x and V_{ex} are the effective matter potentials as given by

$$V_e = \frac{\sqrt{2}G_F\rho}{m_p} \left[\frac{3Y_e - 1}{2} + \xi_{ee}(2 - Y_e) \right],
\tag{12}$$

$$V_x = \frac{\sqrt{2}G_F\rho}{m_p} \left[\frac{Y_e - 1}{2} + \xi_{xx}(2 - Y_e) \right],
\tag{13}$$

$$V_{ex} = \frac{\sqrt{2}G_F\rho}{m_p} \xi_{ex}(2 - Y_e).
\tag{14}$$

Here m_p is the nucleon mass, ρ is the matter density, Y_e is the electron number per nucleon and charge neutrality is assumed ³. For the corresponding anti-neutrino states the sign of matter potentials is opposite.

Let us note that the matter potential induced by the *non standard* FDNC interactions plays the role of an extra effective mass, whereas those induced by the FCNC couplings play the role of a new *mixing* term. As a result, in principle even for strictly massless neutrinos ($\delta m^2 = 0$) and vanishing θ , these new matter potentials make the resonant neutrino conversion in medium possible [1, 3, 4].

Let us now turn to the application of the above to the neutrino conversions in a supernova. Let us discuss separately the cases of $\delta m^2 = 0$ and $\delta m^2 \neq 0$.

³Here the d quark number density N_d in the medium is understood to be expressed as $N_e + 2N_n$.

3 Massless neutrino resonant conversion in supernovae

We now turn to the application of the previous formalism to resonant neutrino conversion in supernovae.

By equating the diagonal terms in the Hamiltonian matrix of eq. (11) one can infer, for the case of massless neutrinos, that the resonance condition is given by

$$\xi' \equiv \xi_{xx} - \xi_{ee} = \frac{Y_e}{2 - Y_e} \quad (15)$$

which is clearly energy independent. Here we should note that a positive value of ξ' is necessary for the above equation to hold. It is important to note that the same resonance condition holds also for the anti-neutrino system $\bar{\nu}_e \leftrightarrow \bar{\nu}_x$. As a result, *both* neutrinos *and* anti-neutrinos can simultaneously undergo resonant conversions as discussed in ref. [2]. As a result, this can affect in an important way supernova neutrino emission.

The mixing angle θ_m and the neutrino oscillation length L_m in matter are given by

$$L_m = \frac{\pi \sin 2\theta_m}{V_{ex}}, \quad (16)$$

$$\tan 2\theta_m = \frac{2\xi_{ex}(2 - Y_e)}{Y_e - \xi'(2 - Y_e)}, \quad (17)$$

respectively.

In our subsequent discussion, we will employ the simple Landau-Zener approximation [21, 22] to estimate the conversion probability after the neutrinos cross the resonance. Under this approximation, the probability for $\nu_e \leftrightarrow \nu_x$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_x$ conversions is given by

$$\begin{aligned} P &= 1 - \exp\left(-\frac{\pi^2}{2} \frac{\delta r}{L_m^{\text{res}}}\right) \\ &\approx 1 - \exp\left[-5 \times 10^4 \times \left(\frac{\rho_{\text{res}}}{10^{12} \text{g/cm}^3}\right) \left(\frac{h}{\text{cm}}\right) \frac{\xi_{ex}^2}{\xi'}\right], \\ \delta r &= 4h \frac{\xi_{ex}}{\xi'}, \quad h \equiv \left|\frac{d \ln Y_e}{dr}\right|_{\text{res}}^{-1}, \end{aligned} \quad (18)$$

where L_m^{res} is the neutrino oscillation length at resonance and ρ_{res} is the corresponding matter density.

Let us briefly review the supernova process we are going to consider. A few seconds after the bounce, the electron number density Y_e is very low just above the neutrinosphere, $Y_e \sim 10^{-2}$, while at large radii it saturates to an asymptotic value ~ 0.4 (see Sect. 4.1 in ([7]).

This implies, from eq. (15), that the resonance condition requires lepton universality to be at least violated at the 1% level, $\xi' \gtrsim 10^{-2}$ which is not in contradiction with present bounds outlined in eq. (9). To keep the discussion simple and more conservative, we consider, for each flavour conversion ($\nu_e \rightarrow \nu_\mu$ or $\nu_e \rightarrow \nu_\tau$), only the contribution due to the exchange of one left-handed \tilde{q} at a time in the corresponding effective couplings $\xi_{ee}, \xi_{xx}, \xi_{ex}$.

After the bounce of the core, all neutrinos, emitted from the neutrinosphere, have approximately equal luminosities but rather different energy spectra. Correspondingly, the average neutrino energies satisfy the following hierarchy:

$$\langle E_{\nu_e} \rangle < \langle E_{\bar{\nu}_e} \rangle < \langle E_{\nu_{\tau(\mu)}} \rangle \approx \langle E_{\bar{\nu}_{\tau(\mu)}} \rangle. \quad (19)$$

Typically, the average supernova neutrino energies are:

$$\langle E_{\nu_e} \rangle \approx 11 \text{ MeV}, \quad \langle E_{\bar{\nu}_e} \rangle \approx 16 \text{ MeV}, \quad \langle E_{\nu_{\tau(\mu)}} \rangle \approx \langle E_{\bar{\nu}_{\tau(\mu)}} \rangle \approx 25 \text{ MeV}. \quad (20)$$

As a result, a considerable conversion $\bar{\nu}_e \leftrightarrow \bar{\nu}_{\mu,\tau}$ leads to a permutation of the neutrino energy spectra which would provide a high energy tail in the anti-neutrino energy spectrum from the supernova SN1987a [23, 24]. Comparison with the SN1987A observations leads to an upper bound for the transition probability P close to 0.35 [14]. Following the same reasoning, we will constrain the effective FCNC couplings that can arise in supersymmetric models with explicitly broken R-parity. Using the density and Y_e profiles from Wilson's supernova model (see Fig. 1 in ref. [7]), we plot in Fig. 1 two contours of the conversion probability in the $((|\lambda'_{ij1}|^2 - |\lambda'_{1j1}|^2), \lambda'_{1j1}\lambda'_{ij1})$ parameter space ($i = 2, 3; j = 1, 2, 3$). Here the *reference* squark mass has been chosen to be 100 GeV. Should the squark mass be different the plot should be appropriately re-scaled. The solid line is for a conversion probability of $P \approx 0.5$, and the dashed one is for $P \approx 0.35$. We see from the figure that, provided the violation of universality induced by the new diagonal interactions is sufficiently high that the resonant conversions take place, i.e. if $(|\lambda'_{ij1}|^2 - |\lambda'_{1j1}|^2) \gtrsim 10^{-2}$ one can rule out $\lambda'_{1j1}\lambda'_{ij1} \gtrsim 10^{-6} \div 10^{-4}$. Note, that this bound on $\lambda'_{1j1}\lambda'_{ij1}$ is about three orders of magnitude stronger than the present experimental one in eq. (9).

In addition, the region above the neutrinosphere is also supposed to be the site for the synthesis of heavy elements (with mass number $A > 70$) through r processes [25]. A necessary condition required for this to occur is $Y_e < 0.5$ in the nucleosynthesis region. The value of Y_e is controlled by the charged current reactions:

$$\nu_e + n \rightleftharpoons p + e^-, \quad (21)$$

$$\bar{\nu}_e + p \rightleftharpoons n + e^+. \quad (22)$$

Roughly speaking, the rates $\Gamma_{\nu_e N}$ of the above reactions are proportional to the products of the ν_e and $\bar{\nu}_e$ luminosities and average energies,

$$\Gamma_{\nu_e N} \approx \phi_\nu \langle \sigma_{\nu_e N} \rangle \propto \frac{L_\nu}{\langle E_\nu \rangle} \langle E_\nu^2 \rangle \propto L_\nu \langle E_\nu \rangle, \quad (23)$$

where ϕ_ν is the neutrino flux, $\sigma_{\nu_e N} \propto E_\nu^2$ is the neutrino absorption cross section, and $\langle \rangle$ denotes the averaging over the neutrino energy distribution. As a result, the relevant expression for Y_e turns out to be very simple:

$$Y_e \approx \frac{\Gamma_{\nu_e n}}{\Gamma_{\bar{\nu}_e p} + \Gamma_{\nu_e n}} \approx \frac{1}{1 + \langle E_{\bar{\nu}_e} \rangle / \langle E_{\nu_e} \rangle}. \quad (24)$$

Using the average energies in eq. (20), we obtain $Y_e \approx 0.41$, in good agreement with the numerical supernova models.

However, in the presence of neutrino conversions, average energies of $\bar{\nu}_e$ and/or ν_e can be affected and consequently the value of Y_e can deviate from the predicted one.

As a result, in the nucleosynthesis region Y_e should be replaced by

$$Y_e \approx \frac{1}{1 + \langle E_{\bar{\nu}_e} \rangle_{\text{eff}} / \langle E_{\nu_e} \rangle_{\text{eff}}}, \quad (25)$$

where

$$\begin{aligned} \langle E_{\bar{\nu}_e} \rangle_{\text{eff}} &\equiv \langle E_{\bar{\nu}_e} \rangle (1 - P) + \langle E_{\bar{\nu}_\tau} \rangle P, \\ \langle E_{\nu_e} \rangle_{\text{eff}} &\equiv \langle E_{\nu_e} \rangle (1 - P) + \langle E_{\nu_\tau} \rangle P. \end{aligned} \quad (26)$$

Due to the the *simultaneous* occurrence of resonant $\nu_e \leftrightarrow \nu_\tau$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_\tau$ conversions, there is a trend to equalize the average ν_e and $\bar{\nu}_e$ energies, and as a result, to increase Y_e with respect to the standard model case with no neutrino or anti-neutrino conversions.

For conversion probabilities of $P \approx 0.15$, 0.3 , and 0.8 , we obtain $Y_e \approx 0.43$, 0.45 , and 0.49 . In Fig. 2, we present the contour lines corresponding to these Y_e values. The dotted, dashed, and solid lines in this figure are for $Y_e \approx 0.43$, 0.45 , and 0.49 , respectively. If we take $Y_e < 0.45$ as a criterion for a successful r -process, then $\lambda'_{131} \lambda'_{ij1} \gtrsim 10^{-6} \div 10^{-4}$ is excluded for $(|\lambda'_{ij1}|^2 - |\lambda'_{131}|^2) \gtrsim 10^{-2}$. This excluded region is similar to the previous one obtained by considering the $\bar{\nu}_e$ energy spectra from SN1987a, because the limits on the conversion probability are about the same in both cases. However, we note that if the r -process indeed occurs in supernovae, then the resulting limits on the effective FCNC couplings are much less dependent on the predicted average neutrino energies than the previous one. This is because the r -process argument relies only on the ratio of the average neutrino energies [cf. Eq. (24)].

A remark is in order. The parameter space we have explored in this section is complementary to the one relevant for the solar neutrino problem [4, 8]. Indeed, in the solar case much larger values of the FDNC couplings ($|\lambda'_{331}|^2 - |\lambda'_{131}|^2 \sim 0.4 \div 0.6$) are necessary to satisfy the resonance condition in the inner solar core where $Y_e \sim 0.7$. Certainly, the $\bar{\nu}_e$ energy spectrum consideration could be used to exclude, at least partially, the resonant massless neutrino conversion as a solution to the solar neutrino problem ⁴, as suggested in [6]. In that case the value of the effective FDNC couplings should be much larger in order to allow the resonant neutrino conversion to take place, i.e. ($|\lambda'_{331}|^2 - |\lambda'_{131}|^2 \geq 0.5$). This would correspond to massless resonant neutrino conversion very far from the neutrinosphere, unlike the case studied in the present paper. On the other hand, no complementary information can be obtained from the r-process nucleosynthesis argument, since this requires neutrinos to undergo the resonance just above the neutrinosphere.

4 Massive Neutrino Conversion in Supernovae

In models with explicitly broken R parity neutrino masses are induced radiatively at the one-loop level due to the exchange of down-type quarks and squarks [9]. A simple estimate of the corresponding diagram shown in Fig. 3, leads to a typical neutrino mass parameter $\lambda'^2 m_d^2 / m_{SUSY}$. For reasonable choices of m_{SUSY} and λ' (see below) one can see that the resulting neutrino masses could lie in the eV range for which they could play an important role in neutrino propagation in the supernova environment. Moreover, such mass could account for the hot dark matter in the Universe. In this section we include the effect of non-zero δm^2 on our previous evolution Hamiltonian of eq. (11). Let us assume, for definiteness, that the vacuum mixing angle characterizing the two-neutrino system is negligible and, moreover, that one of the two neutrino species is much heavier than the other. In our description we will neglect the *non standard* FDNC contributions in the Hamiltonian matrix eq. (11), this way evading the constraints given in eq. (9). In contrast, the FCNCs generated by the R-parity breaking interactions provide the required mixing term in the evolution Hamiltonian, through the matter potential V_{ex} . In this case the resonant condition reduces to the familiar one for the MSW effect with vanishing mixing, i.e.

$$\frac{\delta m^2}{2E} = \frac{\sqrt{2} G_F \rho Y_e}{m_p} \quad (27)$$

⁴Note that such solution is already disfavoured, since it predicts an energy-independent neutrino suppression, contrary to what is indicated by present solar neutrino observations.

A simple numerical check shows that the relevant neutrino mass scale for which the corresponding resonant neutrino conversions will occur in the supernova environment includes neutrino mass range of few eV, which is precisely the one required in order that one of the two neutrino species, ν_e or ν_τ play a role as hot dark matter [26].

The neutrino wave length is still given by eq. (17) where the mixing angle is now given by:

$$\tan 2\theta_m = \frac{2\xi_{ex}\rho(2 - Y_e)}{\rho Y_e - \delta m^2 m_p / (2\sqrt{2}G_F E)}. \quad (28)$$

Therefore, the transition probability is given by

$$\begin{aligned} P &= 1 - \exp\left(-\frac{\pi^2}{2} \frac{\delta r}{L_m^{\text{res}}}\right) \\ &\approx 1 - \exp\left[-1.6 \times 10^{-2} \times \left(\frac{\delta m^2}{1\text{eV}^2}\right) \left(\frac{10\text{MeV}}{E}\right) \left(\frac{2 - Y_e}{Y_e}\right)_{\text{res}}^2 \left(\frac{h}{\text{cm}}\right) \xi_{ex}^2\right], \\ \delta r &= 4h\xi_{ex} \left(\frac{2 - Y_e}{Y_e}\right), \quad h \equiv \left|\frac{d \ln(\rho Y_e)}{dr}\right|_{\text{res}}^{-1}, \end{aligned} \quad (29)$$

This way we will constrain the $(\delta m^2, \lambda'_{1j1} \lambda'_{ij1})$ parameter space irrespective of any universality violation.

Let us note that for a given sign ⁵ of δm^2 only one kind of resonant conversion, either $\nu_e \leftrightarrow \nu_x$ (for $\delta m^2 > 0$), or $\bar{\nu}_e \leftrightarrow \bar{\nu}_x$ (for $\delta m^2 < 0$), can occur. Therefore to discuss $\bar{\nu}_e$ energy spectra distortion from SN1987a we have to assume $\delta m^2 < 0$. The upper bound on $\bar{\nu}_e$ mass from β decay experiment, $m_{\nu_e} < 4.35$ eV (95% C.L.) [27] cut off our relevant δm^2 range in Fig. 3. One sees from this figure that for $\delta m^2 \lesssim 1 \div 20\text{eV}^2$ the FCNC couplings are restricted to be $\lesssim 10^{-3}$ irrespective of any lepton non-universality. From this point of view the limits derived in this section are of more general validity than those of section 3. For this mass hierarchy the resonant neutrino conversion would not conflict with the nucleosynthesis process for any choice of parameters, and therefore no constraint can be obtained.

On the other hand, for $\delta m^2 > 0$ one expects that $\nu_e \leftrightarrow \nu_x$ transitions will occur and they can affect the nucleosynthesis process. In contrast, in this case the $\bar{\nu}_e$ spectra would be unaffected. In Fig. 4 we plot the iso-contours for different values of the electron abundance Y_e . One can see that in the interesting range $\delta m^2 \sim 1 \div 20\text{eV}^2$, favoured by the hot plus cold dark matter scenario [26], we can rule out the FCNC couplings $\lambda'_{1j1} \lambda'_{ij1}$ at the level of few 10^{-3} .

⁵Here we set $\delta m^2 > 0$ for $m_{\nu_x} > m_{\nu_e}$.

5 Resonant Massless Neutrino Conversion and Supernova Shock Re-heating

We would like to briefly address an interesting open problem related with the energetics of supernova explosion. It is now generally accepted that the prompt shock stalls at a radius ~ 100 kilometres, due to photo-dissociation, neutrino losses, and accretion [28]. The main aspect of a supernova explosion is the transfer of energy from the core to the mantle. The mantle is less bound than the core, whose binding energy can grow during the delay to explosion. The core is the protoneutron star that will evolve due to neutrino cooling and deleptonization over many seconds. Bethe & Wilson [29] showed how neutrino heating of the accreted material near the shock could lead to an explosion. It seems compelling that neutrinos mediate this energy transfer and are the agents of explosion [28].

If neutrinos have only standard model interactions the energy they carry seems insufficient to re-energise the shock material. It has been argued that the occurrence of $\nu_e \rightarrow \nu_{\mu,\tau}$ MSW neutrino conversions behind the shock would increase the energy deposited by ν 's. This is due to the fact that the average energy of $\nu_{\mu,\tau}$ is about twice larger than that of ν_e . The capture processes in eq. (21) and eq. (22) are mostly responsible for the energy deposit.

Our scenario is rather distinct from the MSW effect. Unlike in the MSW case, the simultaneous $\nu_e \rightarrow \nu_{\mu,\tau}$ and $\bar{\nu}_e \rightarrow \bar{\nu}_{\mu,\tau}$ conversions can power both reactions eq. (21) and eq. (22) and as a result the effect may be larger than for the standard MSW or resonant spin-flavour precession [30, 31].

We adopt the argument given by Fuller *et al.* in [30] for providing the total heating rate by ν_e and $\bar{\nu}_e$. Qualitatively, the heating rate $\dot{\epsilon}$ is just the product $\langle E \rangle \Gamma_{\nu N} Y_N$ (see eq. (23)), namely

$$\dot{\epsilon} \approx L_\nu \left(Y_n \langle E_{\nu_e} \rangle^2 + Y_p \langle E_{\bar{\nu}_e} \rangle^2 \right) \quad (30)$$

In the presence of complete resonant conversions $\nu_\mu \rightarrow \nu_e$ and $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$ the rate can be increased by the amount

$$\frac{\dot{\epsilon}'}{\dot{\epsilon}} \approx \frac{Y_n \langle E_{\nu_\tau} \rangle^2 + Y_p \langle E_{\bar{\nu}_\tau} \rangle^2}{Y_n \langle E_{\nu_e} \rangle^2 + Y_p \langle E_{\bar{\nu}_e} \rangle^2} = \left(\frac{\langle E_{\nu_\tau} \rangle}{\langle E_{\nu_e} \rangle} \right)^2 \sim 2, \quad (31)$$

where it is assumed $\langle E_{\nu_\tau} \rangle = \langle E_{\bar{\nu}_\tau} \rangle \sim 21$ MeV and $\langle E_{\nu_e} \rangle = \langle E_{\bar{\nu}_e} \rangle \sim 15$ as typical average energies for the earlier epoch after the bounce $t \gtrsim 0.1$ s. At this epoch, the Y_e value is somewhat larger than that characteristic of the later epoch discussed above $Y_e \sim 10^{-2}$.

However, the present experimental bounds on $\lambda'_{1j1}, \lambda'_{3j1}$ allow $\xi' \gtrsim 0.1$, needed in order to have resonant neutrino conversions (see eq. (15)) at $t \gtrsim 0.1$ s if $Y_e \sim 0.15$ at neutrino sphere.

We can notice that in the usual $\nu_e \leftrightarrow \nu_x$ MSW conversion ⁶ the gain in reheating rate with respect to that of the standard model is [30] $\dot{\epsilon}'/\dot{\epsilon} \approx 5/3$ whereas in the resonant spin-flavour precession scenario [31] $\dot{\epsilon}'/\dot{\epsilon} \approx 4/3$.

Clearly, for the massive neutrino case we can also expect analogous effects. Actually the scenario, depending on the sign of δm^2 looks like the usual MSW picture.

6 Conclusions

Supersymmetry with explicitly broken R parity breaking provides a variety of novel possibilities for neutrino propagation properties in the presence of matter, even when they are strictly massless. The supernova matter background seems to be one where most likely resonant conversions of massless neutrinos can play an important role.

We have re-examined the resonant massless-neutrino conversion in a supernova medium in the presence of flavour changing neutral current (FCNC) couplings present in explicit R parity violating supersymmetric models. We have shown how the observed $\bar{\nu}_e$ energy spectra from SN1987a and the supernova r -process nucleosynthesis argument may provide very stringent constraints on such new FCNC interactions. Typically they are much more stringent than previously obtained at the laboratory. From this point of view the SN1987a event provides a strong sensitivity in restricting neutrino properties in supersymmetric models with R parity violation. Our results here are summarized in Figs. 1 and 2.

We have also generalised the description of MSW massive-neutrino conversions in supernovae so as to account for the presence of explicit R -parity-violating FCNCs and determined the corresponding restrictions in the limit of vanishing vacuum mixing. Our results are summarized in Figs. 3 and 4. The relevant neutrino mass scale could play an important role in connection with hot dark matter. While these constraints we derive on R parity violating interactions are weaker than the ones obtained in the massless limit they are still stronger than those available from laboratory experiments. More importantly, they are of wider validity than those obtained in the massless limit.

⁶ Our estimates of the heating rates are somewhat qualitative but they are sufficient for our discussion.

Last but not least, our discussion of massless-neutrino conversions in supernovae should highlight the interest in improving the present laboratory limits on universality violation and flavour changing R-parity breaking interactions.

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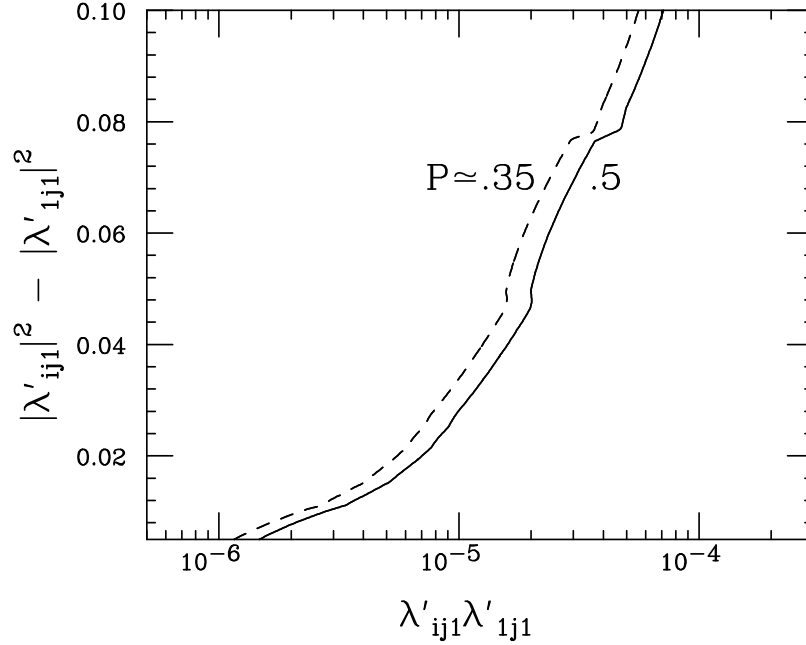


Figure 1: Constraints on the R -parity violating couplings from the observed SN1987a $\bar{\nu}_e$ energy spectra. Here, $i = 2, 3; j = 1, 2, 3$. The dashed (solid) lines correspond to an allowed conversion probability of $P = 0.35$ (0.5). The region to the right of these lines are excluded by the requirement $P < 0.35$ (0.5), as indicated by the SN1987a data.

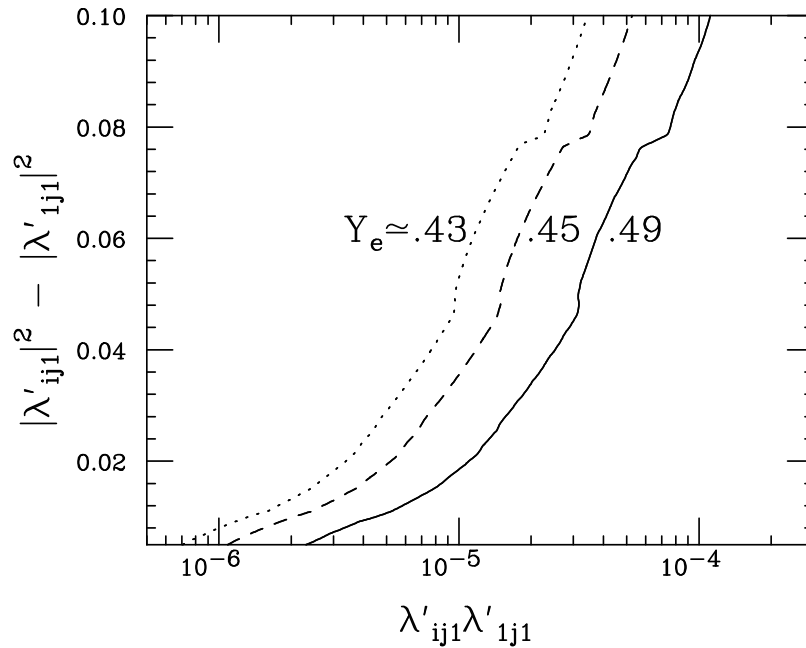


Figure 2: Constraints on the R -parity violating couplings from the supernova r -process nucleosynthesis. The region to the right of the dotted, dashed and solid lines are excluded for the required values of $Y_e < 0.43$, 0.45, and 0.49, respectively, in the r -process.

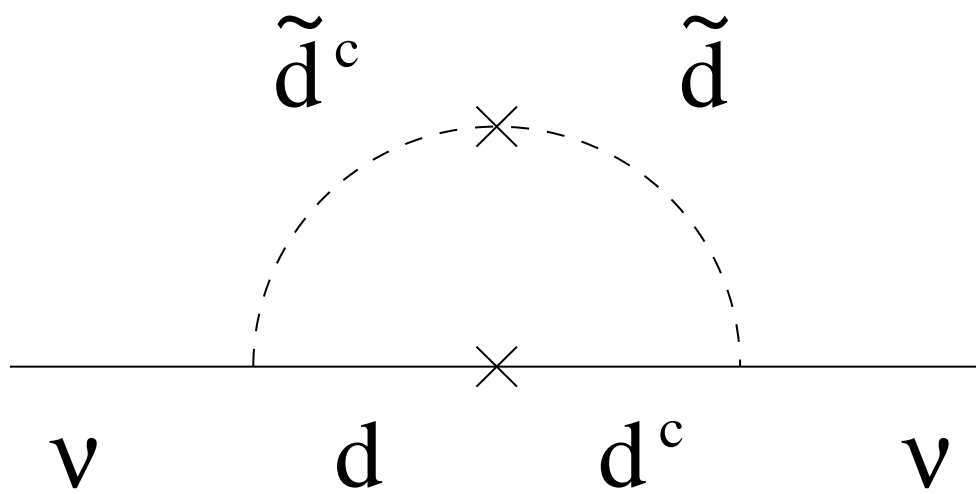


Figure 3: Typical diagram generating neutrino mass in a supersymmetric model with explicitly broken R-parity.

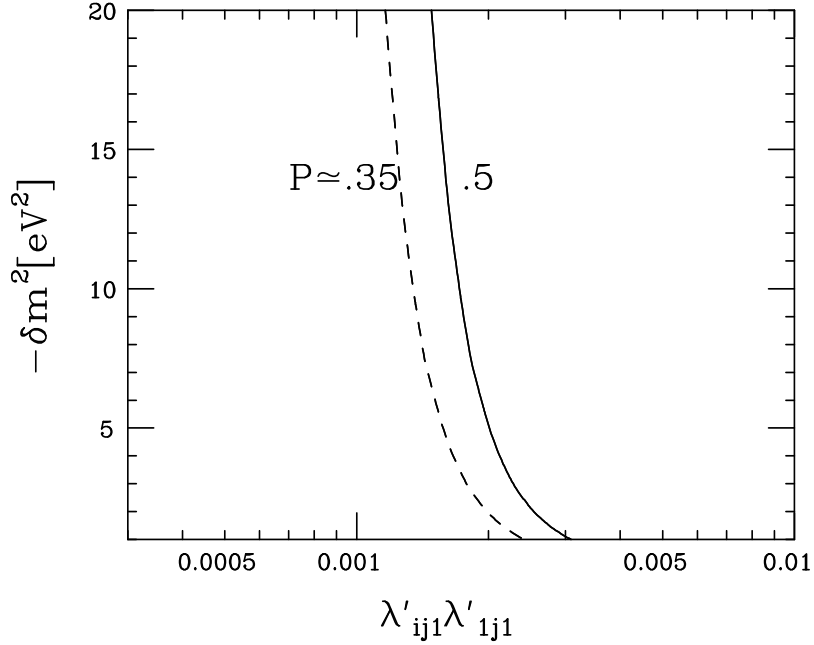


Figure 4: SN1987a $\bar{\nu}_e$ energy spectra constraints on the FCNC R -parity violating couplings for as a function of δm^2 and for negligible vacuum neutrino mixing. The region to the right of the dashed (solid) lines are excluded by the data for an allowed conversion probability of $P < 0.35$ (0.5) irrespective of any laboratory restriction on R -parity-violating interactions.

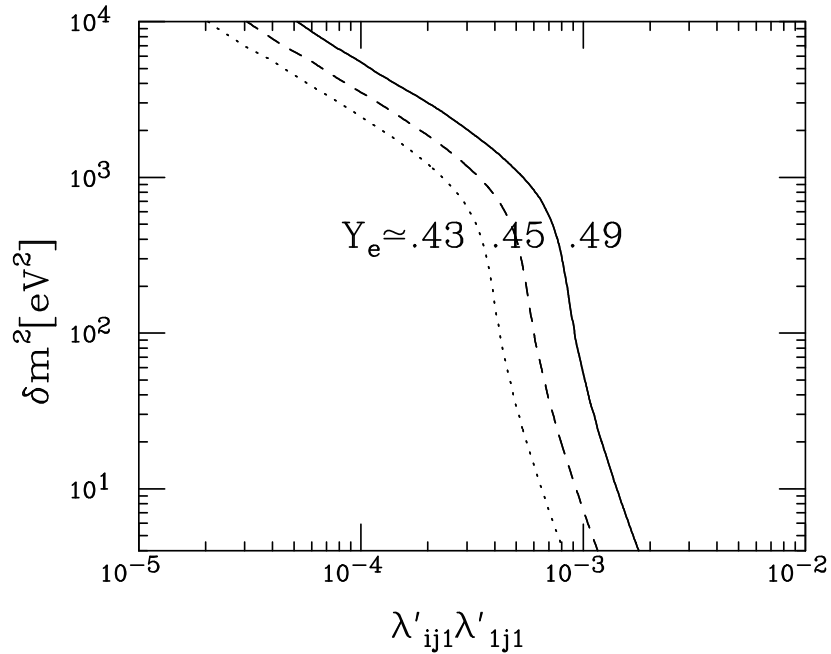


Figure 5: Constraints on the R -parity violating couplings from the supernova r -process nucleosynthesis. The region to the right of the dotted, dashed and solid lines are excluded for the required values of $Y_e < 0.43$, 0.45, and 0.49, respectively, in the r -process.