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Evading the Dirac-neutrino-mass constraint from SN 1987A

K. S. Babu
Bartol Research Institute, University of Delaware, Newark, Delaware 19716

Rabindra N. Mohapatra and I. Z. Rothstein
Department of Physics, University of Maryland, College Park, Maryland 20742
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If the right-handed singlet components (ν_R) of Dirac neutrinos participate in an interaction with quarks via the exchange of scalar leptoquarks, supernova 1987A constraints on their masses can be evaded without running into gross conflict with nucleosynthesis results on the number of allowed neutrino species. Owing to an approximate SU(3)-flavor symmetry of the effective ν_R-quark four-fermion interaction, the ν_R's freeze out sufficiently early from the hadronic matter ($\pi^\pm$, $K^\pm$, nucleons, etc.) during the evoution of the Universe, contributing effectively as 0.5 of a neutrino species to nucleosynthesis, and yet are trapped inside the supernova core via their interaction with the nucleons.

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A massive neutrino can either be a Dirac particle or a Majorana particle. In the former case, one of the two-component spinors (ν_R) needed to generate a mass term in the Lagrangian will be inert with respect to weak interaction [1]. It has recently been shown in several papers [2–5] that the observed supernova 1987A neutrino luminosity provides a very stringent constraint on the Dirac mass of a neutrino ($m_\nu^D \leq 28$ keV, Ref. [2], and $m_\nu^D \leq 1$–10 keV, Ref. [4]). The basic argument goes as follows: In the supernova core, neutrino spin flip can lead to production of the sterile partner ν_R via processes such as bremsstrahlung, ν_e-e scattering, etc. Since the ν_R interaction with the medium is negligible, the energetic ν_R's from the core escape, thereby providing a new mechanism for supernova energy loss. Observed $\bar{\nu}_e$ luminosities by IMB and Kamiokande II experiments, which are consistent with the standard model, therefore provide an upper limit on the spin-flip cross section and hence an upper limit on the Dirac mass.

Interest in the upper limits on the Dirac mass m_ν^D has been heightened by the recent reports of a 17-keV neutrino present as a 10% admixture with ν_e in the study of the β-decay spectra of ^{3}H, ^{35}S, ^{14}C, and ^{71}Ge. In most theoretical models, the 17-keV neutrino is interpreted as a Dirac [6] or a pseudo Dirac [7] neutrino. This would be in conflict with SN 1987A observations [8,9], if the limit is indeed as stringent as in Ref. [4]. It is therefore of interest to see if there are ways to evade the supernova mass constraint.

Supernova cooling arguments have also been used to deduce upper limits on the Dirac-type magnetic moment of the neutrino [10]. The inferred upper limit, viz., $\mu_\nu \leq (10^{-12}$–$10^{-13})\mu_B$, would exclude neutrino spin precession via a Dirac-type magnetic moment as the explanation for the apparent anticorrelation of the neutrino flux with the sunspots reported in the chlorine experiment. Although Majorana-type transition magnetic moments can still account for the anticorrelation, it is of interest to see if the supernova bound on the Dirac magnetic moment can be evaded.

In order to evade the mass (or the magnetic moment) constraint, we must find a mechanism to trap the right-handed sterile neutrino. This implies the existence of new interactions of the ν_R with electrons or nucleons with a strength of order $(10^{-1}$–$1)G_F$. Such interactions can potentially keep the ν_R's in equilibrium during the nucleosynthesis era of the Universe leading to $N_\nu = 4$ for the effective number of neutrino species. This would be in conflict with recent detailed analysis of the abundance of light elements which quotes [11] $N_\nu \leq 3.4$.

In this paper, we suggest leptoquark-mediated interactions of ν_R as a way to evade the supernova constraint on the Dirac mass (or magnetic moment) of the neutrino. We show that if the vector piece of the resulting effective ν_R-quark four-fermion interaction is dominantly a singlet of flavor SU(3), then the ν_R's freeze out from the hadronic matter in the Universe (i.e., π, K, nucleons, etc.) sufficiently early contributing only as half an extra neutrino species ($\Delta N_\nu = 0.5$), which is not too far off the present upper limit.

It is important to note that the effective four-fermion interaction of ν_R must be mediated by a leptoquark Higgs boson. If it were mediated, for instance, by a scalar boson without baryon number [12], ν_R would interact with electrons with a strength close to the Fermi strength. This would keep it in equilibrium at the time of nucleosynthesis leading to $N_\nu = 4$ [13]. The SU(3)-singlet nature of the hadronic current coupling to ν_R is also crucial, since otherwise ν_R interactions with the pions, which are still abundant after muon decoupling, and with the kaons, which are abundant after the quark-hadron phase transition, would lead to $N_\nu = 4$.

Consider the following leptoquark-mediated interaction of ν_R with the right-handed u, d, and s quarks:

$$L_{\nu_R} = f_u \nu_R^T C^{-1} u_R \Delta_u + f_d \nu_R^T C^{-1} d_R \Delta_d + f_s \nu_R^T C^{-1} s_R \Delta_s + \text{H.c.} \quad (1)$$

Here C is the charge-conjugation matrix. If we assume further that $|f_u| = |f_d| = |f_s|$ and $M_{\Delta_u} \simeq M_{\Delta_d} \simeq M_{\Delta_s}$, the

ν_R-quark four-fermion interaction can be written as

$$H_{\rm eff}=\frac{f_u^2}{2M_{\Delta_u}^2}(\bar u_R\gamma_\mu u_R+\bar d_R\gamma_\mu d_R+\bar s_R\gamma_\mu s_R)\bar\nu_R\gamma^\mu\nu_R+\frac{f_u^2}{4M_{\Delta_u}^2}\left[\frac{M_{\Delta_d}^2-M_{\Delta_u}^2}{M_{\Delta_u}^2}\right](\bar u_R\gamma_\mu u_R-\bar d_R\gamma_\mu d_R)\bar\nu_R\gamma^\mu\nu_R$$
$$-\frac{f_u^2}{12M_{\Delta_u}^2}\left[\frac{M_{\Delta_u}^2+M_{\Delta_d}^2-2M_{\Delta_s}^2}{M_{\Delta_u}^2}\right](\bar u_R\gamma_\mu u_R+\bar d_R\gamma_\mu d_R-2\bar s_R\gamma_\mu s_R)\bar\nu_R\gamma^\mu\nu_R\,. \tag{2}$$

Let us now define $f_u^2/8M_{\Delta_u}^2\equiv\epsilon G_F/\sqrt{2}$. The ν_R's emitted by the weak spin-flip process in the core of the supernova get trapped by this interaction. To see the minimum required value of ϵ for trapping, we note that the cross section for the ν_R-nucleon scattering mediated by the leptoquarks is given by

$$\sigma_{\nu_R N}\simeq\frac{9}{\pi}(g_V^2+3g_A^2)G_F^2\epsilon^2E^2\,. \tag{3}$$

Recent European Muon Collaboration (EMC) data suggest that the isosinglet axial-vector coupling g_A is $\simeq 0.1$. Using this value, the cross section is found to be $10^{-39}(E/100\ \text{MeV})^2\ \epsilon^2\ \text{cm}^2$. For an average neutrino energy $E\simeq 100$ MeV, this leads to a ν_R mean free path $l_{\nu_R}\simeq 1.5/\epsilon^2$ cm. A rough estimate of ϵ required to trap the ν_R's inside the supernova core can be obtained by demanding that l_{ν_R} be less than a tenth of the supernova core radius R. For $R\sim 10$ km, this implies $\epsilon\gtrsim 4\times10^{-3}$. A more careful consideration using neutrino sphere [14] shows $\epsilon\gtrsim 3\times10^{-2}$ is needed. This number is obtained by assuming that the supernova luminosity in ν_R is less than 20 times the total $\bar\nu_{eL}$ luminosity. If thermal energy transport by the trapped ν_R's is too efficient it will result in a shortening of the observed neutrino signal [15]. The duration of the observed neutrino pulse would then set a lower limit on ϵ. Considering various uncertainties, we shall allow for $\epsilon\sim 0.1$.

We must now consider the effect of the sterile neutrino on nucleosynthesis. Prior to the QCD phase transition ($T\approx 200$ MeV) the ν_R's will be in equilibrium through interaction (2) with the free quarks. However, the ν_R's will not be in equilibrium during nucleosynthesis since there are no free quarks around, and the baryon-number density is so small ($n_B/n_\gamma\approx 10^{-10}$) [16]. Though not in equilibrium at the nucleosynthesis epoch, the sterile neutrino will contribute as a full species unless the photonic temperature is enhanced due to the decoupling of other particle species during the period between ν_R decoupling and nucleosynthesis. To determine the decoupling temperature, we must consider the interaction of ν_R with the hadronic matter (π, K, nucleons, etc.).

The sterile neutrino interacts with the mesons (π, K, $\eta,\ldots$) only via the SU(3)-nonsinglet current. If we choose the parameters such that

$$(M_{\Delta_u}^2-M_{\Delta_d}^2)/M_{\Delta_u}^2\lesssim 10^{-2}\,,$$
$$(M_{\Delta_u}^2-M_{\Delta_s}^2)/M_{\Delta_u}^2\lesssim 10^{-2}\,, \tag{4}$$

then the SU(3)-nonsinglet parts of the effective interaction will have strengths less than $\epsilon G_F\times10^{-2}$. For $\epsilon\simeq 0.1$ (which is sufficient for ν_R trapping inside the supernova core), $\nu_R\pi$ and $\nu_R K$ interactions go out of equilibrium above $T=150$ MeV. Both the scattering rate $\nu_R\phi\to\nu_R\phi$ and the production rate $\phi\phi\to\phi\phi\nu_R\bar\nu_R$ (where ϕ stands for a typical light scalar or vector meson such as π, $K,\ldots$) become slow compared to the Hubble expansion rate. One meson-initiated production process that is nonzero even for an SU(3)-singlet interaction is the process $\pi+\pi\to\pi+\nu+\bar\nu$. We have estimated this contribution in the omega-dominance approximation. Because of the fact that the $\omega\to e^+e^-$ rate is very small, we find that this process also goes out of equilibrium above 100 MeV. Contributions from η, η', etc., all vanish by G-parity conservation.

Let us now turn to the interaction of the ν_R's with the nucleons. Since the nucleons have a baryon number, the $\nu_R N\to\nu_R N$ four-fermion vertex will have a strength $\sim\epsilon G_F$. For $\epsilon\gtrsim 0.1$, the scattering rate will not become slower than the Hubble expansion rate until $T\simeq 60$ MeV or so. However, since scatterings do not change the number of ν_R's, this process remaining in thermal equilibrium until low temperatures does not effect the calculation of ΔN_ν at $T\simeq 1$ MeV. There are, however, production processes such as $NN\to NN\nu_R\bar\nu_R$, $\pi N\to N\nu_R\bar\nu_R$, etc., which change the number of ν_R's and affect ΔN_ν. It is, therefore, important that these processes decouple at reasonably high temperatures. We have estimated the collision rates for the above two processes. The nucleon-nucleon collision rate for $T\lesssim m_N$ is proportional to $e^{-2m_N/T}$. This process freezes out at $T=150$ MeV even if $\epsilon\simeq 1$.

Since the pions are relativistic for T between 100 and 150 MeV, the $\pi N\to N\nu_R\bar\nu_R$ process is important. There is, however, one important dynamical suppression of T/m_n in the rate coming from the fact that if the nucleon three-momenta are set to zero, then the vector-current contribution to the scattering cross section vanishes. The axial-vector contribution is down in amplitude by a factor of 20–30 relative to the vector part. Taking this suppression into account and the nucleon density factor of $e^{-m_N/T}$ in the rate, we estimate the decoupling temperature for the production reaction using the equation $H\sim n_\pi n_N\langle\sigma v\rangle/n_\nu$ and find that decoupling at $T=150$ MeV requires $\epsilon\lesssim 2\times10^{-2}$ and at $T=100$ MeV requires $\epsilon\lesssim 10^{-1}$. These estimates are obtained as follows. The vector-current contribution to the matrix element squared for the process $\pi N\to N\nu_R\bar\nu_R$ is

$$|M|^2=\frac{9G_F^2\epsilon^2f^2}{2m_\pi^2q_0^2}[2q\cdot k_1q\cdot k_2-m_\pi^2k_1\cdot k_2](\mathbf{p}_1-\mathbf{p}_2)^2\,, \tag{5}$$

where $\pi(q)+N(p_1)\to N(p_2)+\nu_R(k_1)+\bar\nu_R(k_2)$. Here $f\simeq 1.1$ is the pion-nucleon coupling. From this, we estimate the thermally averaged ν_R production rate to be

$$\langle\sigma v\rangle \approx c\frac{2187}{8\pi^3}\frac{G_F^2\epsilon^2 f^2 T^5}{m_\pi^2 m_N}. \tag{6}$$

Here we have collected all factors of π and 3 (coming from $E \simeq 3T$). The factor c is expected to be of order unity and arises from the phase-space integral, which we have not evaluated exactly. The estimates quoted above correspond to choosing $c=1$, which we believe is reliable to within a factor of 2 or so. We have assumed here that the pions and the nucleons are in equilibrium and that the pions are relativistic.

We shall therefore assume that after $T \sim 100$ MeV the total number of ν_R remains constant. But the ratio n_{ν_R}/n_γ decreases each time a heavy particle species disappears. At $T \sim 100$ MeV, there are many hadronic species (such as π, K, η, η', ρ, ω, K^*, $N, \ldots$) which are in thermal equilibrium via strong interaction, whereas at $T=1$ MeV they, as well as the muons, have all disappeared dumping their entropy into photons. This increases the photon temperature and therefore n_γ but n_{ν_R} remains unaffected. n_{ν_R}/n_γ decreases substantially from its value of $\frac{7}{8}$ (or $\Delta N_\nu = 1$) at $T \sim 100$ MeV. To calculate this decrease, we use the entropy conservation principle [17]. In Fig. 1 we plot the effective ΔN_ν as a function of the ν_R decoupling temperature, where we have included contributions from the muon as well as from all hadronic matter with masses below a GeV. In the range $T=100$–150 MeV, we find that $\Delta N_\nu \simeq 0.5$. This value is slightly higher than the upper limit on ΔN_ν quoted by Walker *et al.* [11]. In fact, $\Delta N_\nu = 0.5$ leads to a primordial helium abundance of 24.3% [11], which is to be compared with the observed value of 23% $\pm$ 1% at the $2-\sigma$ level.

Let us briefly discuss an extension of the standard model where such interactions of the leptoquarks arise with the required properties. Consider the gauge group $SU(2)_L \times SU(2)_R \times U(1)$. We shall impose an S_3 permutation symmetry acting in the generation space. The quarks and leptons transform under $SU(2)_L \times SU(2)_R \times U(1) \times S_3$ as follows:

$$\begin{pmatrix} u_L & c_L \\ d_L & s_L \end{pmatrix}:(2,1,\tfrac{1}{3})(2),\quad \begin{pmatrix} u_R & c_R \\ d_R & s_R \end{pmatrix}:(1,2,\tfrac{1}{3})(2),$$

$$\begin{pmatrix} \nu_{eL} & \nu_{\mu L} \\ e_L & \mu_L \end{pmatrix}:(2,1,-1)(2),\quad \begin{pmatrix} N_{eR} & N_{\mu R} \\ e_R & \mu_R \end{pmatrix}:(1,2,-1)(2),$$

$$\nu_R:(1,1,0)(1). \tag{7}$$

The numbers in the second parentheses represent transformation under the permutation symmetry S_3. We have exhibited only two families of fermions, the third family may be assumed to be a singlet of S_3.

Apart from the inclusion of a singlet neutrino ν_R to the fermionic spectrum (which should be identified with the sterile component of the 17-keV neutrino), the model is identical to the familiar left-right-symmetric model. We shall not assume manifest left-right symmetry; this would allow us to choose the $SU(2)_R$-breaking scale to be relatively low. A Higgs multiplet $\chi_R(1,3,2)$, an S_3 singlet, breaks the gauge symmetry down to $SU(2)_L \times U(1)_Y$. In addition, it gives large Majorana masses to N_{eR}, $N_{\mu R}$ fields. The electroweak symmetry breaking is achieved by

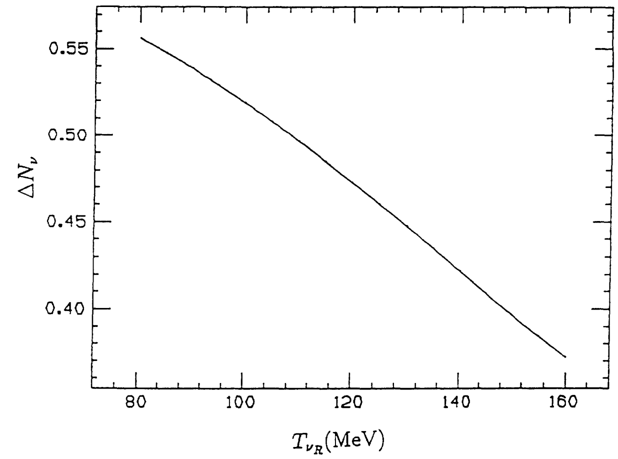

FIG. 1. The effective value of ΔN_ν at $T \approx 1$ MeV as a function of the ν_R decoupling temperature.

a $\Phi(2,2,0)(2)$ Higgs multiplet, which also breaks the S_3 symmetry. Its Yukawa couplings generate nonzero fermion masses and mixing angles. $SU(2)_L$-doublet fields $\varphi(2,1,1)$ are needed to generate the ν_R Dirac mass. The leptoquark (color-antitriplet) fields which couple ν_R to the right-handed quarks are denoted by

$$\Delta(1,2,-\tfrac{1}{3})(2) = \begin{pmatrix} \Delta_d & \Delta_s \\ \Delta_u & \Delta_c \end{pmatrix}. \tag{8}$$

Because of the presence of $SU(2)_R$ and S_3 symmetries, we have $|f_u| = |f_c| = |f_d| = |f_s|$ for the leptoquark Yukawa couplings. In addition, $M_{\Delta_u} = M_{\Delta_d} = M_{\Delta_c} = M_{\Delta_s}$ prior to symmetry breaking. After $SU(2)_R$ symmetry breaks, a mass splitting $(M_{\Delta_u}^2 - M_{\Delta_d}^2) = \lambda v_R^2$ results, where λ is the coefficient of the quartic scalar coupling $(\Delta^\dagger \chi_R^\dagger \chi_R \Delta)$ and v_R is the vacuum expectation value (VEV) of χ_R^0. Similarly, $(M_{\Delta_d}^2 - M_{\Delta_s}^2) = \lambda' v_L^2$ will be induced via the S_3 breaking VEV's of Φ field. If we choose the $SU(2)_R$ scale around 500 GeV or so, $M_{\Delta_u} \simeq M_{\Delta_d} \simeq M_{\Delta_s} \simeq M_{\Delta_c} \approx 500$ GeV. Then for $\lambda \simeq 10^{-2}$, the SU(3)-nonsinglet part of the hadronic currents will be sufficiently suppressed.

Without detailed elaboration, we wish to mention that the model does not lead to large K^0-$\bar{K}^0$ mixing for $M_{W_R} \simeq 500$ GeV as in the usual left-right-symmetric models, since we can choose the right-handed quark mixings to be much smaller than their left-handed counterparts. Furthermore, there are no electron-leptoquark interactions allowed in the model. The process $K^+ \to \pi^+ \nu\bar{\nu}$ can proceed in this model via leptoquark exchange. The decay amplitude is proportional to the mixing angle θ_R in the d_R-s_R sector, which we shall choose to be small. There is also a contribution from the Δ_d-Δ_s mixing, but this amplitude is proportional to $\epsilon G_F \lambda'(v_L^2/v_R^2)$. For $\epsilon \simeq 10^{-1}$, $\lambda' \simeq 10^{-2}$, and $v_L/v_R \simeq \frac{1}{4}$, the branching ratio is $\lesssim 10^{-8}$, consistent with present limits.

In conclusion, we wish to mention that the effective Hamiltonian of Eq. (2) can also be obtained in extensions of the standard model with an extra U(1) gauge symmetry under which ν_R has a nonzero charge and (u_R, d_R, s_R) have equal charges. The SU(3)-flavor-violating piece of Eq. (2) will be zero at the tree level in this case.

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