Resonant spin-flavor precession of solar and supernova neutrinos

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(Received 5 October 1987)

The combined effect of matter and magnetic fields on neutrino spin and flavor precession is examined. We find a potential new kind of resonant solar-neutrino conversion $\nu_{e_L} \to \nu_{\mu_R}$ or ν_{τ_R} (for Dirac neutrinos) or $\nu_e \to \bar{\nu}_\mu$ or $\bar{\nu}_\tau$ (for Majorana neutrinos). Such a resonance could help account for the lower than expected solar-neutrino ν_e flux and/or indications of an anticorrelation between fluctuations in the ν_e flux and sunspot activity. Consequences of spin-flavor precession for supernova neutrinos are also briefly discussed.

There has been a long-standing disagreement between the solar-neutrino ν_e flux monitored by Davis[1] and collaborators,

$$\text{average flux} = 2.1 \pm 0.3 \text{ solar-neutrino unit (SNU)} \quad (1 \text{ SNU} = 10^{-36} \text{ captures}/s \text{ atom}) , \tag{1}$$

via the reaction $\nu_e + {}^{37}\text{Cl} \to e^- + {}^{37}\text{Ar}$ and Bahcall's standard-solar-model prediction[2]

$$\text{predicted flux} = 7.9 \pm 2.5 \text{ SNU} \quad (3\sigma \text{ errors}) . \tag{2}$$

That discrepancy has come to be known as the solar-neutrino puzzle. Attempts to resolve it have given rise to many speculative ideas about unusual properties of neutrinos and/or the solar interior.

One rather recently proposed solution, the Mikheyev-Smirnov-Wolfenstein[3] (MSW) effect, is particularly elegant. It employs the changing solar density and the difference between ν_e and other neutrinos' interactions with matter to bring about an energy-level-crossing resonance. In that way, ν_e neutrinos propagating from the core of the Sun to its surface can encounter a resonance region (generally in the radiation zone which extends from $0.04R_\odot$ to $0.7R_\odot$, where $R_\odot \simeq 7\times 10^{10}$ cm) and be converted to ν_μ, ν_τ, or some as yet unknown flavor. Part of the MSW solution's appeal is that it naturally provides the required flux depletion [cf. Eqs. (1) and (2)] for a large range of neutrino mass differences $\Delta m_{21}^2 = m_2^2 - m_1^2 \simeq 10^{-7}-10^{-4}$ eV^2 and mixing angles, $\sin^2 2\theta \gtrsim 0.001$; so, it is not contrived.[4] In fact, the required parameters are very much in keeping with theoretical prejudices and expectations.

A more speculative solution to the solar-neutrino puzzle, originally advocated by Cisneros,[5] involves endowing neutrinos with magnetic moments such that spin precession in the strong interior solar magnetic fields can lead to $\nu_{e_L} \to \nu_{e_R}$. The sterility of ν_{e_R} would then lead to an effective depletion in the measured flux. That scenario has been recently revived and improved by Okun, Voloshin, and Vysotsky[6] (OVV). Their motivation came from the observation that there appears to be an anticorrelation between sunspot activity and variations in the detected solar ν_e flux.[7,8] During times of high sunspot activity (i.e., large-magnetic-field disturbances in the convection zone $> 0.7R_\odot$), the measured flux is smallest. It is, therefore, quite natural to correlate the flux variation with spin precession, which would of course be greatest when the magnetic fields are most intense.

The precession scenario has been studied in some detail by OVV (Ref. 6). They noted that either a magnetic- or electric-dipole moment could give $\nu_{e_L} \to \nu_{e_R}$ precession, while flavor transition moments between different species could result in the combined spin-flavor precession $\nu_{e_L} \to \nu_{\mu_R}$ or ν_{τ_R} (for four component Dirac neutrinos) and $\nu_e \to \bar{\nu}_\mu$ or $\bar{\nu}_\tau$ (Ref. 9) (for Majorana neutrinos). [Precession of Majorana neutrinos gives rise, for example, to $\nu_{e_L} \to (\nu_{\mu_L})^C$ with C being the charge-conjugation operator. Since $(\nu_{\mu_L})^C$ is right handed and generally called $\bar{\nu}_\mu$, we refer to it that way.] In all cases, however, they concluded that μB, where μ is a generic dipole or transition moment and B is the transverse solar magnetic field, must at least be of order $10^{-16}-10^{-15}$ eV to make such a scenario viable. Since they argued that B could be of order 10^3 G in the Sun's convection zone where precession was envisioned, they required[6]

$$|\mu| \simeq 0.3-1\times 10^{-10} e/2m_e . \tag{3}$$

A moment that large is consistent with direct experimental bounds,[10]

$$|\mu_{\nu_e}| \le 4\times 10^{-10} e/2m_e \tag{4}$$

(from $\overset{(-)}{\nu}_e e$ data[11]), but the upper range is slightly in conflict with astrophysical arguments which imply[12]

$$|\mu_{\nu_e}| < 8.5\times 10^{-11} e/2m_e \quad \text{(astrophysics bound)} . \tag{5}$$

In addition, it is generally difficult to generate such a large moment while keeping the neutrino mass very small. For example, the standard $SU(2)_L \times U(1)$ model with a singlet right-handed neutrino gives rise to[13]

$$|\mu_{\nu_e}| = \frac{3eG_\mu m_{\nu_e}}{8\sqrt{2}\pi^2} \simeq 3\times 10^{-19}(m_{\nu_e}/1\ \text{eV})e/2m_e \ , \quad (6)$$

which is much too small for a viable solar-neutrino precession scenario. Recently, however, a model has been proposed by Fukugita and Yanagida[14] in which an $SU(2)_L$ charged scalar singlet can induce at the one-loop level a neutrino magnetic moment in the range of Eq. (3) without conflicting with low-energy phenomenology. Therefore, in this paper we keep an open mind regarding the magnitude of μ.

It has also been noted[6] that neutrino interactions with matter can quench spin precession. In the case of $\nu_{e_L} \to \nu_{e_R}$ precession, the ν_{e_L} and ν_{e_R} interact differently with matter. The difference in their interactions effectively splits their degeneracy and suppresses precession. To illustrate that point, we consider the evolution equation connecting the chiral components ν_{e_L} and ν_{e_R}:

$$i\frac{d}{dt}\begin{bmatrix}\nu_{e_R}\\ \nu_{e_L}\end{bmatrix} = \begin{bmatrix}0 & \mu B\\ \mu B & a_{\nu_e}(t)\end{bmatrix}\begin{bmatrix}\nu_{e_R}\\ \nu_{e_L}\end{bmatrix} \ , \quad (7)$$

where B is the transverse magnetic field and $a_{\nu_e}(t)$ represents the "matter" potential experienced by ν_{e_L} as it propagates through the Sun. In the standard model (for an unpolarized medium),

$$a_{\nu_e}(t) = \frac{G_\mu}{\sqrt{2}}[(1+4\sin^2\theta_W)N_e + (1-4\sin^2\theta_W)N_p - N_n] \ , \quad (8)$$

where $G_\mu = 1.166\,36\times 10^{-5}$ GeV^{-2}, $\sin^2\theta_W \simeq 0.23$ and N_f is the fermion number density. [For other neutrino species $(1+4\sin^2\theta_W)N_e \to (-1+4\sin^2\theta_W)N_e$, while for antineutrinos, the sign of $a(t)$ is reversed.] For a neutral medium $N_e = N_p$ and one finds, from Eq. (8),

$$a_{\nu_e}(t) = \frac{G_\mu}{\sqrt{2}}(2N_e - N_n) \ . \quad (9)$$

The t (or spatial) dependence comes about because the densities in Eq. (9) vary as the neutrino propagates outward through the solar interior. In addition, B will also vary with t; but it is likely to exhibit a complicated dependence.

To obtain a feel for the matter effect, we can solve Eq. (7) for constant densities and constant B. In that case, the spin-precession probability of starting with ν_{e_L} at $t=0$ and finding ν_{e_R} at time t is given by

$$P(t)_{\nu_{e_L}\to\nu_{e_R}} = \frac{(2\mu B)^2}{a_{\nu_e}^2 + (2\mu B)^2} \times \sin^2\{[a_{\nu_e}^2 + (2\mu B)^2]^{1/2} t/2\} \ . \quad (10)$$

In a vacuum where $a_{\nu_e} = 0$, this expression reduces to a standard spin-precession formula with frequency μB; but if $a_{\nu_e}^2 \gg (2\mu B)^2$, precession is suppressed. Of course, to carry out a realistic calculation, one needs a density and magnetic-field profile for the solar interior. The densities of electrons and neutrons in the convection zone and upper radiation zone[15] are well approximated by

$$N_e \simeq 6N_n \simeq 2.4\times 10^{26}\exp(-r/0.09R_\odot)/\text{cm}^3 \ , \quad 0.2 < r/R_\odot \le 1 \ , \quad (11a)$$

while in the lower radiation zone the linear approximation

$$N_e \simeq 6\times 10^{25}\left[1 - \frac{10}{3}\frac{r}{R_\odot}\right] \ , \quad N_n \simeq 2\times 10^{25}\left[1 - \frac{21}{5}\frac{r}{R_\odot}\right]\Big/\text{cm}^3 \ , \quad 0.1 < r/R_\odot \le 0.2 \quad (11b)$$

works well. (The relative neutron density increases.) Unfortunately, little is known about magnetic fields in the core, radiation zone or convection zone, except that they may be quite large. We have examined the evolution in Eq. (7) for an average B of $\sim 10^3$ G in the convection zone (i.e., for a distance $\simeq 2\times 10^{10}$ cm) and find, for $\mu \simeq 10^{-10}e/2m_e$, one can obtain a ν_e flux depletion consistent with Eqs. (1) and (2). That finding is in keeping with the results of OVV (Ref. 6) and a more recent analysis by Barbieri and Fiorentini.[16] Of course, if that scenario is realized, the ν_e flux depletion would be strongly dependent on the magnitude of the magnetic field. Hence, one could expect a strong correlation between ν_e solar flux and sunspot activity, which is a measure of the convection-zone magnetic field.

We come now to the main focus of our work, the effect of matter on spin-flavor precession. To begin, we note that even if an electromagnetic transition moment between mass eigenstates ν_1 and ν_2 exists, one expects spin-flavor precession $\nu_{1L} \to \nu_{2R}$ in magnetic fields to be suppressed by the mass difference between ν_2 and ν_1, unless[6,9]

$$\mu_{21}B > \Delta m_{21}^2/2E_\nu \quad (12)$$

with μ_{21} the transition moment and E_ν the neutrino energy. For $\mu_{21} \simeq 10^{-10}e/2m_e$, $B \simeq 10^3$ G, and $E_\nu \simeq 10$ MeV, that condition requires $\Delta m_{21}^2 \lesssim 10^{-7}$ eV^2, which is below the MSW solutions but not prohibitively small. (Of course, the condition depends on energy.) Partly because of that mass difference suppression, neutrino spin-flavor precession seems not to have been thoroughly studied. However, here we will show that matter interactions of the distinct neutrino flavors can compensate for the mass difference. In fact, for a medium of changing density, such as the Sun, a resonance region can exist where neutrino spin-flavor precession may occur unimpeded. The physics of that resonance is quite similar to the MSW resonance[3] as we shall see.

To illustrate the resonant spin-flavor precession phenomenon, we first consider the case of two generations with four-component Dirac neutrinos. (Extensions to higher generations are straightforward but cumbersome.) For definiteness, we examine the ν_e-ν_μ system; but the results hold also for ν_e-ν_τ.

Using the chiral bases ν_{e_L}, ν_{μ_L}, ν_{e_R}, ν_{μ_R}, the evolution equation for neutrino propagation through matter and a transverse magnetic field B is

$$i\frac{d}{dt}\begin{bmatrix}\nu_{e_L}\\ \nu_{\mu_L}\\ \nu_{e_R}\\ \nu_{\mu_R}\end{bmatrix}=\begin{bmatrix}H_L & BM^\dagger\\ BM & H_R\end{bmatrix}\begin{bmatrix}\nu_{e_L}\\ \nu_{\mu_L}\\ \nu_{e_R}\\ \nu_{\mu_R}\end{bmatrix}, \tag{13}$$

where the 2×2 submatrices are

$$H_L=\begin{bmatrix}(\Delta m_{21}^2/2E_\nu)\sin^2\theta+a_{\nu_e} & (\Delta m_{21}^2/4E_\nu)\sin 2\theta\\ (\Delta m_{21}^2/4E_\nu)\sin 2\theta & (\Delta m_{21}^2/2E_\nu)\cos^2\theta+a_{\nu_\mu}\end{bmatrix},\quad H_R=\begin{bmatrix}0 & 0\\ 0 & \Delta m_{21}^2/2E_\nu\end{bmatrix},\quad M=\begin{bmatrix}\mu_{ee} & \mu_{e\mu}\\ \mu_{\mu e} & \mu_{\mu\mu}\end{bmatrix}, \tag{14}$$

and θ is the neutrino mixing angle

$$\begin{bmatrix}\nu_e\\ \nu_\mu\end{bmatrix}=\begin{bmatrix}\cos\theta & \sin\theta\\ -\sin\theta & \cos\theta\end{bmatrix}\begin{bmatrix}\nu_1\\ \nu_2\end{bmatrix}. \tag{15}$$

The different matter potentials are given by (for a neutral unpolarized medium)

$$a_{\nu_e}=\frac{G_\mu}{\sqrt{2}}(2N_e-N_n),\quad a_{\nu_\mu}=\frac{G_\mu}{\sqrt{2}}(-N_n), \tag{16}$$

and M represents the electromagnetic moments in the chiral-flavor bases. Note that because ν_{e_R} and ν_{μ_R} are sterile, i.e., do not interact electroweakly with matter, they can be considered vacuum mass eigenstates.

From Eq. (13) we can easily determine the energy-level crossings of the four neutrino chiral states by equating terms along the diagonal. Those crossings are illustrated in Fig. 1, where we have used the approximation $N_n\simeq\frac{1}{6}N_e$ [cf. Eq. (11a)] and assume $\Delta m_{21}^2>0$ is the more plausible scenario. There are three potential crossing resonances. The usual MSW $\nu_{e_L}\to\nu_{\mu_L}$ oscillation resonance occurs at density (for small θ)

$$N_e\simeq\frac{\Delta m_{21}^2}{2\sqrt{2}G_\mu E_\nu}\quad\text{(MSW resonance)}, \tag{17}$$

while the $\nu_{e_L}\to\nu_{\mu_R}$ spin-flavor resonance occurs at a somewhat higher density:

$$N_e\simeq\frac{12}{11}\frac{\Delta m_{21}^2}{2\sqrt{2}G_\mu E_\nu}\quad(\nu_{e_L}\to\nu_{\mu_R}\text{ resonance}). \tag{18}$$

The $\nu_{\mu_L}\to\nu_{e_R}$ spin-flavor resonance occurs at the much higher density:

$$N_e\simeq 12\frac{\Delta m_{21}^2}{2\sqrt{2}G_\mu E_\nu}\quad(\nu_{\mu_L}\to\nu_{e_R}\text{ resonance}). \tag{19}$$

Note that the MSW and $\nu_{e_L}\to\nu_{\mu_R}$ resonance regions are not so far from one another and may in some cases have to be studied together.

To illustrate the above scenario, we consider an adiabatic MSW solution $\Delta m_{21}^2\simeq 10^{-4}$ eV^2 and $\sin^2 2\theta$ small. In that case, the MSW resonance occurs for $E_\nu\simeq 10$ MeV at $r\simeq 0.10R_\odot$ while the $\nu_{e_L}\to\nu_{\mu_R}$ resonance is at $r\simeq 0.065R_\odot$. [These values are directly obtained from the bare density profiles N_e and N_n (Ref. 15), without approximation (11a) or (11b).] Both are in the inner radiation zone. The $\nu_{\mu_L}\to\nu_{e_R}$ resonance requires too large a N_n and therefore does not exist in the Sun. In this example, the MSW and $\nu_{e_L}\to\nu_{\mu_R}$ resonances are far enough spaced to be treated separately. The $\nu_{e_L}\to\nu_{\mu_R}$

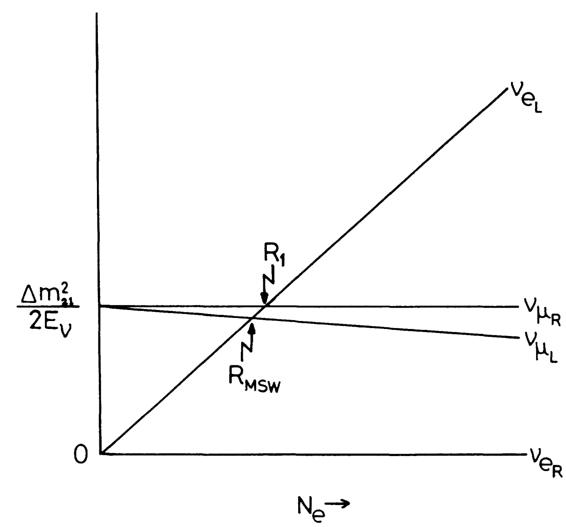

FIG. 1. The energy-level crossings of four neutrino chiral states in the two-generation model with Dirac neutrinos. The medium is assumed to be proton rich, say $N_e\simeq 6N_n$, for simplicity. The crossings denoted by $R_{\rm MSW}$ and R_1 correspond to the MSW resonance and the resonant $\nu_{e_L}\to\nu_{\mu_R}$ spin-flavor precession. R_2, the crossing of ν_{μ_L} and ν_{e_R}, does not appear for solar neutrinos because the required density is too high.

resonant conversion is therefore governed by a 2×2 submatrix

$$\begin{bmatrix} (\Delta m_{21}^2/2E_\nu)\sin^2\theta + a_{\nu_e} & \mu_{\mu e}^* B \\ \mu_{\mu e} B & \Delta m_{21}^2/2E_\nu \end{bmatrix} . \quad (20)$$

That problem is the same as the MSW resonant conversion one if we recall that a_{ν_e} gets contributions not only from the charged current but also from the neutral current in the present case [cf. Eq. (16)]. We can therefore utilize the analysis of Parke[17] to determine the average $\nu_{e_L} \to \nu_{\mu_R}$ transition probability

$$P(\nu_{e_L} \to \nu_{\mu_R}) = \tfrac{1}{2} - (\tfrac{1}{2} - P_{LZ})\cos 2\tilde{\theta}_N \cos 2\tilde{\theta} , \quad (21)$$

where

$$\tan 2\tilde{\theta} = 4E_\nu \,|\mu_{\mu e}|\, B/\Delta m_{21}^2 \cos^2\theta , \quad (22a)$$

$$\tan 2\tilde{\theta}_N = 4E_\nu \,|\mu_{\mu e}|\, B/[\Delta m_{21}^2 \cos^2\theta - 2E_\nu a_{\nu_e}(0)] , \quad (22b)$$

$$P_{LZ} = \exp\left[-\frac{4\pi E_\nu |\mu_{\mu e}|^2 B^2}{\Delta m_{21}^2 \cos^2\theta \,|d(\ln a_{\nu_e})/dr|} \right] , \quad (22c)$$

and $d(\ln a_{\nu_e})/dr$ in Eq. (22c) should be the value at the resonance point.

To be more quantitative, we have given the numerical solar-neutrino depletion results in Table I for several cases. For illustrative purpose, we have assumed $\Delta m_{21}^2 = 10^{-4}$ eV^2, $E_\nu = 10$ MeV and only $\mu_{\mu e} = 10^{-10} e/2m_e$ has been taken into account as the electromagnetic moment. Note that for relatively small B, as in the cases (a) and (b), the MSW resonance shown in the case (d) essentially provides the ν_{e_L} depletion. This is because the magnetic fields are too small for the adiabatic $\nu_{e_L} \to \nu_{\mu_R}$ transition to take place. For larger values of B, however, the $\nu_{e_L} \to \nu_{\mu_R}$ conversion only can account for the depletion, as we learn from the case (c). In the radiation zone, such large magnetic fields $\simeq 10^4 - 10^5$ G may in fact be feasible. It is also interesting to note that very large magnetic fields are entirely

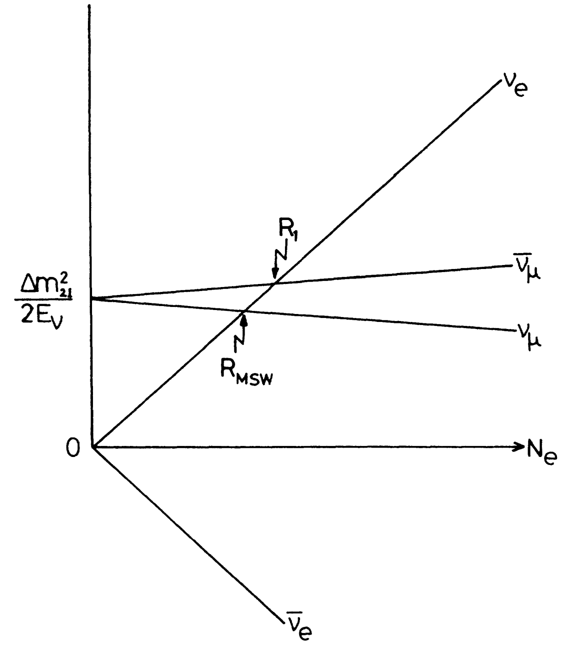

FIG. 2. The energy-level crossings in the two-generation Majorana neutrino model. The proton-rich medium has been assumed as in Fig. 1. For Majorana neutrinos the third crossing R_2 between ν_μ and $\bar{\nu}_e$ does not exist (for $N_e \geq N_n$). For $N_e < N_n$, the R_1 disappears while the R_2 crossing takes place.

reasonable in the supernova and the neutron stars. The value of $P_{\nu_{\mu_R}}$ given in (c) is in good agreement with the analytic formula Eq. (22). The final example, case (e), combines both types of resonances (c) and (d), to result in enough fractions of both $P_{\nu_{\mu_R}}$ and $P_{\nu_{\mu_L}}$.

In the above example, the $\nu_{e_L} \to \nu_{\mu_R}$ precession occurs deep in the Sun's radiation zone. Hence, it is unlikely to be correlated with sunspot activity. However, one could speculate that in a three-generation scenario $\nu_{e_L} \to \nu_{\tau_L}$ or ν_{τ_R} occurs in the radiation zone, while $\nu_{e_L} \to \nu_{\mu_R}$ occurs in the convection zone with the latter process correlated with sunspot activity. [The resonance in the convection

TABLE I. The probabilities $P_{\nu_{e_L}}$, $P_{\nu_{\mu_L}}$, and $P_{\nu_{\mu_R}}$ of a solar neutrino produced at the solar core as ν_{e_L} to be observed as ν_{e_L}, ν_{μ_L}, and ν_{μ_R}, respectively, at the solar surface in the two-generation model with Dirac neutrinos. For illustrative purpose, we have assumed $\Delta m_{21}^2 = 10^{-4}$ eV^2, $E_\nu = 10$ MeV, and only $\mu_{\mu e} = 10^{-10} e/2m_e$ has been switched on among electromagnetic moments. Several cases (a),(b), . . . , (e) with different sets of parameters B (G) and $\sin\theta$ are shown.

Case	$\sin\theta$	B	$P_{\nu_{e_L}}$	$P_{\nu_{\mu_L}}$	$P_{\nu_{\mu_R}}$
(a)	7.18×10^{-3}	10^3	0.33	0.67	0
(b)	7.18×10^{-3}	10^4	0.32	0.62	0.06
(c)	0	5.25×10^4	0.33	0	0.67
(d)	7.18×10^{-3}	0	0.33	0.67	0
(e)	7.18×10^{-3}	5.25×10^4	0.11	0.37	0.52

zone would necessitate, through Eq. (18), $\Delta m^2_{21} \lesssim 2 \times 10^{-7}$ eV^2 (for $E_\nu \simeq 10$ MeV and small θ). This condition happens to be similar to the one derived by OVV (Ref. 6) from Eq. (12). They are, however, based on very different physical processes. For instance, in the spin-flavor precession without resonance[6] the average probability of $\nu_{e_L} \to \nu_{\mu_R}$ does not become as large as $\frac{2}{3}$ even if Δm^2_{21} is small enough.] Careful monitoring of the ν_e flux in future experiments could shed light on such a scenario.

In the case of Majorana neutrinos, only transition moments which change lepton number by $|\Delta L| = 2$ can exist. They lead to $\nu_e \to \bar{\nu}_\mu$ or $\bar{\nu}_\tau$ spin-flavor precession and if mixing is large, $\nu_e \to \bar{\nu}_e$. Since antineutrinos can interact with matter, the evolution equation for ν_e, ν_μ, $\bar{\nu}_e$, and $\bar{\nu}_\mu$ is governed by the Hamiltonian 4×4 matrix

$$H_{\text{Maj}} = \begin{bmatrix} a_{\nu_e} & \frac{\Delta m^2_{21}}{4E_\nu}\sin 2\theta & 0 & \mu_\nu^* B \\ \frac{\Delta m^2_{21}}{4E_\nu}\sin 2\theta & \frac{\Delta m^2_{21}}{2E_\nu}\cos 2\theta + a_{\nu_\mu} & -\mu_\nu^* B & 0 \\ 0 & -\mu_\nu B & -a_{\nu_e} & \frac{\Delta m^2_{21}}{4E_\nu}\sin 2\theta \\ \mu_\nu B & 0 & \frac{\Delta m^2_{21}}{4E_\nu}\sin 2\theta & \frac{\Delta m^2_{21}}{2E_\nu}\cos 2\theta - a_{\nu_\mu} \end{bmatrix} . \quad (23)$$

Therefore, for small mixing θ, the $\nu_e \to \bar{\nu}_\mu$ level crossing occurs at somewhat higher densities [using $N_n \simeq (\frac{1}{6}) N_e$]

$$N_e \simeq \frac{3\sqrt{2}}{10} \frac{\Delta m^2_{21}}{G_\mu E_\nu} \quad (\nu_e \to \bar{\nu}_\mu \text{ resonance}) , \quad (24)$$

and there is no $\nu_\mu \to \bar{\nu}_e$ resonance (see Fig. 2). Most of what was said for $\nu_{e_L} \to \nu_{\mu_R}$ above carries over to $\nu_e \to \bar{\nu}_\mu$ precession. However, $\bar{\nu}_\mu$ or $\bar{\nu}_\tau$ are, in principle, detectable by measuring $\bar{\nu}$-e scattering cross section. Although that is very difficult, we are optimistic that it will one day be possible. That optimism is based on the tremendous progress in measuring $\nu_e - e$ scattering recently reported.[18] Also, if $\bar{\nu}_e$ (from mixing) is appreciable, one might try to detect $\bar{\nu}_e + p \to e^+ + n$, although it will be hard to disentangle from backgrounds.

In closing, we should emphasize that spin or spin-flavor precession of the solar neutrino is still a long shot. It requires large magnetic fields and a very large electromagnetic dipole or the transition moment. However, we feel that the existence of spin-flavor precession resonant regions makes such a scenario at least plausible. One should, therefore, keep an open mind and search for antineutrinos from the Sun as well as monitor the energy dependence of ν_e and its correlation with solar activity. Finally, we note that even if spin-flavor precession is too weak to affect solar neutrinos (say if μ is too small), it could influence other phenomena. In particular, spin-flavor precession of Majorana neutrinos could have remarkable consequences for supernova bursts[19] where very intense magnetic field ($\sim 10^{12}$ G) are quite likely. For example, it could lead to large asymmetries in the ν_μ and $\bar{\nu}_\mu$ or ν_τ and $\bar{\nu}_\tau$ fluxes, which are otherwise expected to be equal. An even more dramatic possibility stems from the fact that in the core region of a supernova $N_n > N_e$ is realized. In such a neutron-rich region $\nu_\mu \leftrightarrow \bar{\nu}_e$ or $\nu_\tau \leftrightarrow \bar{\nu}_e$ become the more plausible spin-flavor precessions (rather than $\nu_e \leftrightarrow \bar{\nu}_\mu, \bar{\nu}_\tau$). That could lead to a higher than expected $\bar{\nu}_e$ flux and/or effectively higher average energy $\bar{\nu}_e$. Also, potential sources of very high-energy neutrinos (such as Cygnus X-3) could lead to $\nu_\mu \to \bar{\nu}_\tau$ or $\bar{\nu}_\mu \to \nu_\tau$ (depending on the sign of the mass difference). We, of course, strongly advocate experimental searches for all such phenomena.

We are very grateful to A. Baltz, S. P. Rosen, J. Weneser, and M. Yoshimura for useful and informative discussions. One of us (C.S.L.) would like to thank K. Nomoto and T. Sakurai for very helpful and introductory arguments on the density profiles and the magnetic fields of the Sun and the supernova. He also would like to thank N. Ohnishi, and members of Theory Group at the Institute for Nuclear Studies, University of Tokyo, for their useful discussions. W.J.M. was supported by the Department of Energy under Contract No. DE-AC02-76-CH-0016.

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