Possible Solution to the Discrepancy between the Homestake and Kamiokande Solar-Neutrino Experiments

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A possible particle-physics solution is proposed to the recently reported discrepancy between the Homestake and Kamiokande II experiments concerning the "time variation" of the solar-neutrino flux. It is pointed out that some of the models giving rise to a large magnetic moment of electron neutrinos include the possibility that the right-handed neutrino interacts with electrons strongly enough so that Kamiokande does not see the time variation, while the Homestake experiment shows the time variation in anticorrelation with solar activity due to the neutrino helicity flip in the Sun.

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One of the most puzzling issues in solar-neutrino physics is the fact that the capture rate measured by the Homestake ^{37}Cl detector appears to exhibit a time variation in anticorrelation with the sunspot number.[1] The new runs of the Homestake experiment since the fall of 1986 continue to show this anticorrelation; the capture rate, which was as high as 4.2 solar-neutrino units (SNU) in 1986–1987 (solar quiet time), has gradually been decreasing as solar activity increases, and it dropped to be as low as ~ 1 SNU for the runs in 1989.[2] In particular, five-point running average indicates quite a conspicuous anticorrelation between these two quantities. A remarkable similarity in the time-variation pattern is also reported between the capture rate in the runs starting from 1977 and those from 1986.5.[2] These observations, if not entirely confirmed, tend to corroborate the intriguing suggestion that the neutrino capture rate in the Homestake detector anticorrelates with solar activity.

The recent report[3] from the Kamiokande II Collaboration has made the situation even more puzzling; their water detector[4] recorded ^{8}B solar-neutrino events almost at a constant rate from January 1987 to April 1990.[3] The 200-day flux averages for 40 months do not show apparent time variation more than the error of 30% from bin to bin, whereas the Homestake runs show a factor-of-4 decrease in the ^{37}Cl capture rate over the same period. The flux deduced from the Kamiokande experiment $\phi(^{8}\mathrm{B})=2.7\pm0.3\pm0.3$ cm^{-2}s^{-1} is consistent with only the highest capture rate data from the Homestake experiment.

These two results appear to be almost contradictory with each other at first glance. We point out in this Letter, however, that there exists a possibility that these two observations can be theoretically reconciled. A class of models[5] which was introduced to explain the large magnetic moment of the electron neutrino, needed to account for the time variation of the solar-neutrino capture rate at Homestake,[6] includes this possibility as a special case.

The only explanation known so far to explain the time variation of the ^{37}Cl neutrino capture rate anticorrelated with the sunspot number is that by Voloshin, Vysotsky, and Okun.[6] They proposed that the electron neutrino (ν_e) may have a large magnetic moment so that a significant fraction of left-handed neutrinos rotates into right-handed neutrinos under the strong magnetic field in the convective layer of the Sun in the period of maximum solar activity and the solar neutrino becomes dominantly sterile to the nuclear detector; such an effect is suppressed in the quiet time of the Sun. This explanation requires a very large magnetic moment of ν_e, which is hard to understand in conventional models; in the standard model its chirality-conserving feature leads to a very small magnetic moment, which is proportional to the mass of neutrinos.

In Ref. 5 it was shown that the needed large magnetic moment may be accounted for, if a scalar particle exists, which couples strongly to neutrinos, since a scalar particle does not respect the chirality conservation. A model with an SU(2)$_L$-singlet charged scalar was exemplified in Ref. 5 for the most economical case. Similar models were explored by a number of authors.[7] In this Letter we consider a model with an SU(2)$_L$-doublet scalar in addition to the standard Weinberg-Salam particles. The reason for taking a doublet rather than a singlet in the present consideration will be explained below.

The crucial observation is that in such a class of models right-handed neutrinos interact, in general, with electrons with a substantial strength owing to the introduced scalar particle, whereas neutrino-nucleus interactions proceed as the standard theory predicts.

To make the argument clear, let us take the Lagrangian

$$\mathcal{L}=f_{ij}\bar{e}_{R,i}l_{L,j}\phi^{*}+g_{ij}\bar{\nu}_{R,i}l_{L,j}\phi+\mathrm{H.c.}\,, \tag{1}$$

where ϕ stands for an extra SU(2)-doublet scalar field added to the standard model and it is assumed not to develop the vacuum expectation value [i.e., the renormalized mass of ϕ in the effective Lagrangian is assumed to be positive, even after the breaking of $SU(2)_L \times U(1)$]. The neutrino is assumed to be of the Dirac type. The indices i,j denote generations $(i,j=1-3)$. We define $e_{R,i}$, $e_{L,i}$ $[l_{L,i}=(\nu_{L,i},e_{L,i})]$ to be the mass eigenstate, and $\nu_{L,i}$ to be the eigenstate of weak interactions assuming that the mass of neutrino is zero or negligible. Using the freedom for the definition of $\nu_{R,i}$, one can take, in general, the coupling for the second term to be

$$\begin{aligned} &g_{ij}=0 \quad (\text{for } i>j)\,, \\ &g_{ij}\neq 0 \quad (\text{otherwise})\,. \end{aligned} \tag{2}$$

We assume that this ϕ does not couple to quark channels.[8]

Let us now consider the minimal condition for the Yukawa couplings f_{ij} and g_{ij} required in our considerations. The magnetic moment of the electron neutrino is calculated in a manner similar to that in Ref. 5, and it is given by

$$\mu_{\nu_e}=\frac{e}{32\pi^2}\sum_i (g_{1i}f_{i1}+g_{1i}^*f_{i1}^*)\frac{m_i}{M_\phi^2}\left[\ln\left(\frac{M_\phi^2}{m_i}\right)-1\right], \tag{3}$$

where m_i is the charged-lepton mass of generation i. For simplicity, we assume that all couplings are real, and the couplings that appear in the parentheses of (3) are dominated by $i=3$ (τ). The condition that (3) gives $\mu_{\nu_e}=10^{-11}\mu_B$ (μ_B denotes the Bohr magneton), which is an upper bound derived from stellar evolution of helium stars,[9] yields

$$g_{13}f_{13}/M_\phi^2 \approx 1.2\times 10^{-7}\,\text{GeV}^{-2} \tag{4}$$

for $\ln(M_\phi^2/m_\tau^2)\approx 8$.

The second requirement comes from the condition that the left-handed neutrino rotates into the right-handed neutrino under the magnetic field (B) of a reasonable strength. Coherent weak interactions of neutrinos with matter[10] generate a mass gap between ν_L and ν_R,

$$\Delta E=(G_F/\sqrt{2})(2N_e-N_n)\,, \tag{5}$$

with N_e and N_n the electron and neutron number densities, respectively, which hinders ν_L from rotating into ν_R if $\Delta E>\mu_{\nu_e}B$.[6] If μ_{ν_e} is to be as small as $10^{-11}\mu_B$, the condition $\Delta E \ll \mu_{\nu_e}B$ requires B to be as large as 150 kG,[11] an order of magnitude larger than the value required for the case of vanishing ΔE.

To circumvent this problem the idea has been proposed that this ΔE may be canceled at some particular density by giving neutrinos an appropriate mass difference.[12] In our model, on the other hand, ν_R also interacts coherently with e_R via the ϕ exchange with the effective Hamiltonian

$$H_{\text{eff}}^{(\phi)}=(|g_{11}|^2/4M_\phi^2)(\bar\nu_{R,e}\gamma_0\nu_{R,e})N_e\,. \tag{6}$$

Hence $\nu_{L,e}$ and $\nu_{R,e}$ receive the potential in matter,

$$H_{\text{eff}}=\begin{array}{cc} \nu_L & \nu_R \end{array}\!\qquad\qquad \begin{array}{c} \\ \left(\begin{array}{cc} \dfrac{G_F}{\sqrt{2}}(2N_e-N_n) & \mu_{\nu_e}B \\ \mu_{\nu_e}B & \dfrac{|g_{11}|^2}{4M_\phi^2}N_e \end{array}\right)\begin{array}{c} \nu_L \\ \\ \nu_R \end{array} \end{array}. \tag{7}$$

We emphasize here that the Yukawa interaction g_{ij} for ν_R is only weakly constrained by laboratory experiments and the strength required to cancel the (1,1) element in (7) is allowed, as we discuss below. (The Yukawa coupling f_{ij} for ν_L is more strongly constrained by neutrino experiments, and $f_{11}^2/M_\phi^2<0.1G_F$.) This cancellation makes possible the ν_L-ν_R rotation under a reasonable magnetic-field strength. For $\mu_{\nu_e}\approx 10^{-11}\mu_B$, $B\approx 15$ kG, which is the value required from the simple condition $\mu_{\nu_e}BL\approx \pi/2$ with $L\approx 0.3R_\odot$, is sufficient for

$$\beta=\frac{|g_{11}|^2}{4M_\phi^2}\Big/\sqrt{2}G_F=0.8-1.0\,. \tag{8}$$

This condition requires $|g_{11}|^2/4M_\phi^2$ to be as strong as the weak-interaction strength. This means that both $\nu_{L,e}$ and $\nu_{R,e}$ interact with electrons with almost the same strength (the cross sections for ν_L and ν_R become equal if $\beta=1.28$). If such a case is realized, the reaction rate of $\nu_e e^-\to\nu_e e^-$ changes little, even if $\nu_{L,e}$ rotates into $\nu_{R,e}$. On the other hand, $\nu_e+{}^{37}\text{Cl}\to{}^{37}\text{A}+e^-$ does not receive a contribution from the ϕ exchange, and the reaction rate reduces if $\nu_{L,e}$ rotates into $\nu_{R,e}$. This might explain why the time variation in anticoincidence with the sunspot number seen in the Homestake experiment[1,2] was not detected in the Kamiokande detector.[3] The fact that Kamiokande data are consistent with the highest capture rate data at Homestake is also understood naturally. From the absence of time variation in the Kamiokande we obtain $\beta=0.98-1.5$, allowing for $\pm 30\%$ statistical errors.

Let us remark here on the advantage of the present model over the example given in Ref. 5, where the additional scalar η^+ is taken to be an SU(2) singlet. In this case the coherent effect of the η^+ exchange contributes to the potential for $\nu_{R,e}$ with a negative sign that enhances the mass difference. Hence a strong magnetic field of >200 kG is needed to flip ν_L into ν_R. In this regard we consider the present model more attractive for a consistent scenario.

We now discuss constraints imposed on the present model. To simplify our analysis we keep only the parameters which are necessarily nonzero and set others equal to zero; namely, we take

$$f_{31}\neq 0\,,\quad g_{13}\neq 0\,,\quad \text{and } g_{11}\neq 0 \tag{9}$$

and all others vanishing.[13] The requirements to be satisfied for our scenario are $g_{11}^2/4M_\phi^2\simeq 1.6\times 10^{-5}$

GeV^{-2} and the one given in (4). Constraints from laboratories and astrophysics are discussed in order.

(i) The strongest constraint on g_{11} comes from the experiment $e^+e^- \to \gamma+$missing ("neutrino-counting" experiment). We may use the limit on the scalar electron $\tilde{e}$ which mediates the photino production process $e_L^+ + e_L^- \to \gamma+\tilde{\gamma}+\tilde{\gamma}$ to constrain the mass of ϕ, since $e_L^+ + e_L^- \to \gamma+\nu_R+\bar{\nu}_R$ via the ϕ exchange mimics the above process. The current limit from the ASP detector at the SLAC e^+e^- storage ring PEP, $m_{\tilde{e}_L} < 47$ GeV,[14] translates into $g_{11}^2/4M_\phi^2 < 2\times10^{-5}$ GeV^{-2} (or $\beta < 1.3$).

(ii) ϕ exchange induces the exotic decay of τ, $\tau_L \to \nu_{R,e} + (e_L + \bar{\nu}_{R,e}), \nu_{L,e} + (e_L + \bar{\nu}_{R,e})$. We calculate the branching ratio to these channels to be

$$B = \frac{1}{5}\frac{1}{4M_\phi^4}(g_{11}^2 g_{13}^2 + g_{11}^2 f_{31}^2)\frac{1}{8G_F^2}\,. \qquad (10)$$

That this be smaller than the disparity allowed from the e-μ universality test ($\lesssim 1\%$) (Ref. 15) results in $|g_{13}^2 + f_{13}^2|/M_\phi^2 < 3.3\times10^{-6}$ GeV^{-2} for $\beta=1$.

(iii) Constraints from the anomalous magnetic moment of the electron[15] are weak. We only have $|g_{11}^2 + f_{11}^2|/M_\phi^2 < 1\times10^{-1}$ GeV^{-2}.

(iv) $\nu_{R,e}$ contributes to the expansion rate of the early Universe by an amount equivalent to one extra neutrino species. Despite recent improvement in the nucleosynthesis calculation,[16] one extra species seems still allowed due to uncertainties in the estimate of primordial He abundance and of a lower bound of the baryon density. $\nu_{R,\mu}$ and $\nu_{R,\tau}$ contribute little to the expansion, if the relevant Yukawa couplings are small.

(v) The argument has been made that a large magnetic moment mediates $\nu_L e \to \nu_R e$ scattering and leads to copious ν_R production in the core of supernovas, that and fugacity of ν_R might affect the supernova dynamics.[17] In our model a large g_{11} results in a large ν_R opacity and ν_R's are trapped in the core.[18] Hence the dynamics of supernovas is not drastically altered in any case.[18]

Let us finally discuss the experimental test for the present scenario. It is obvious that gallium solar-neutrino experiments[19] should see the time variation in agreement with the Homestake experiment. More decisive information may be afforded from the Sudbury neutrino observatory,[20] which measures simultaneously neutrino-electron elastic scattering and deuteron dissociation ($\nu_e + d \to p+p+e^-$). The ratio of these two rates should vary largely with solar activity. The neutrino-counting-type laboratory experiment will also give us valuable information on this problem; the improvement of the bound on the mass of the scalar electron, by 50%, say, would clearly rule out the presently discussed possibility, since our scenario requires the Yukawa coupling for ν_R to be almost as large as the weak-interaction strength. We note that the result of neutrino counting by measuring $e^+e^- \to \gamma+$missing provides us with information (see, e.g., Ref. 18) different from that from the Z^0 width, which is sensitive only to the left-handed neutrinos.

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Note added.—We have recently learned that the Soviet-American Gallium Experiment[21] reported a suppressed solar-neutrino capture rate on gallium. This is what was predicted in the present model.

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