Leptonic long-range forces and the MSW process in the Sun

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The observed duration of the neutrino pulse from the Supernova 1987A may be used to constrain hypothetical long-range forces coupling neutrinos and galactic matter. If these forces are mediated by massless vector bosons, this can lead to a drastic change in the usual MSW analysis for the Sun, depending on screening effects by the cosmic neutrino background, electrons, protons, or dark-matter particles accumulated inside the Sun.

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The detection of about a dozen (anti)neutrinos from the SN 1987A [1,2] in the Large Magellanic Cloud have been used by many authors to constrain particle properties from the observed duration of the neutrino pulse. The absence of any anomalous dispersion of the neutrino pulse has been used to scrutinize neutrino properties such as their mass [3], electric charge [4], and also some sort of "fifth force charge" [5] for neutrinos and galactic matter constituents.

The long-range interaction coupled to a novel neutrino charge was assumed to be mediated by a massless spin-one particle [5]. If electrons or protons in our Galaxy or dark-matter particles in its halo also carry such a charge, the novel long-range interaction may be responsible for bending of the neutrino trajectory, causing a larger path to be traveled by the less energetic neutrinos than the more energetic ones. From the fact that the energy-dependent time delay does not exceed the expected spread of the emission times, about ten seconds, one can derive, depending on the nature of the source particles and their distribution in our Galaxy, the constraint

$$|q_{e,p,x}q_\nu| \lesssim 3\times 10^{-40}\times \begin{cases} 1 & \text{for } e,p, \\ m_x/m_p & \text{for dark matter.}\end{cases} \tag{1}$$

Recently, it has been observed [6] that the CP-symmetric cosmic neutrino background constitutes a neutral plasma with regard to the new interaction, so that the screening of a source charge should be taken into account in derivation of the SN 1987A limit as given by Eq. (1). Since the screening effects operate on galactic scales, one finds that, in most cases, it is impossible to derive any SN 1987A bound [6].

It is clear that any additional interaction between neutrinos and galactic matter can also affect the important Mikheyev-Smirnov-Wolfenstein (MSW) mechanism [7,8] inside the Sun. Here we derive an extra MSW potential for the Sun corresponding to the new long-range interaction, assuming the different source particles (e,p,x). We show that for the coherent forward neutrino scattering mediated by massless vector bosons, the screening effects are crucial for the determination of the new MSW contribution. In the massless-soft regime, the resummation program developed by Braaten and Pisarski [9,10] must be applied, resulting in the use of the resummed propagator for massless vector bosons. Hence, mass singularities are shielded owing to the screening property of the resummed propagator. On the other hand, the MSW interaction in the standard theory is mediated by the heavy vector bosons $(W^\pm, Z^0)$, resulting in the use of bare propagators.

We perform our calculation using the covariant field-theoretical treatment, known as finite temperature and density (FTD) field theory (for convenience we use the real-time version). In addition, we choose the self-energy approach (instead of the amplitude approach) and consider neutrino self-energies presented in Figs. 1 and 2.

First we consider the case when the huge electron number in our Galaxy acts as a source of the new long-range force. Consequently, electrons in the Sun would induce the self-energy for neutrinos (see Fig. 1) whose contribution reads

$$-i\Sigma_5^r(k) = q_{\nu_r}q_e\gamma_\alpha iD^{\alpha\beta}(q=0)\int\frac{d^4p}{(2\pi)^4}\mathrm{tr}\{\gamma_\beta iS^e(p)\}\ . \tag{2}$$

Since solar neutrinos ν_e can oscillate into either active (ν_μ,ν_τ) or sterile neutrinos (ν_s), $r=e,\mu,\tau$ or s. The resummed propagator $D^{\alpha\beta}$, obtained by summing all the one-particle irreducible self-energy diagrams, can be written in a covariant gauge, exhibiting the transverse, longitudinal, and gauge-fixing part [11]

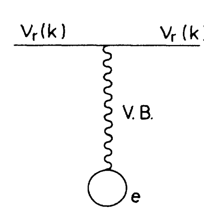

FIG. 1. The contribution to the neutrino self-energy due to nonvanishing neutrino and electron charges.

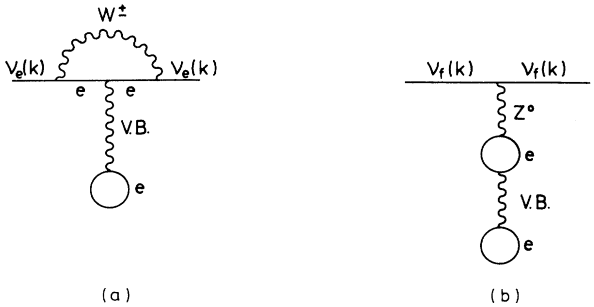

FIG. 2. The contribution to the neutrino self-energy due to its charge induced by background electrons: (a) charged-current contribution; (b) neutral current contribution.

$$iD^{\alpha\beta}(q_0,|\vec{q}|) = \frac{(-i)}{G(q_0,|\vec{q}|)-q^2}P_T^{\alpha\beta} + \frac{(-i)}{F(q_0,|\vec{q}|)-q^2}P_L^{\alpha\beta} + \frac{i\rho}{q^2}\frac{q_\alpha q_\beta}{q^2} , \quad (3)$$

with the explicit expressions

$$P_T^{00} = P_T^{0i} = 0; \quad P_T^{ij} = \delta_{ij} - \frac{q_i q_j}{|\vec{q}|^2} ,$$
$$P_L^{\alpha\beta} = \frac{q_\alpha q_\beta}{q^2} - g^{\alpha\beta} - P_T^{\alpha\beta} \quad (4)$$

and ρ is the gauge parameter. The real-time propagators are given by

$$S^e(p) = (\not{p}+m)\left\{\frac{1}{p^2-m^2+i\varepsilon} + 2\pi i\delta(p^2-m^2)\eta(p\cdot v)\right\} \quad (5)$$

with

$$\eta(x) = \frac{\theta(x)}{e^{\beta(x-\mu)}+1} + \frac{\theta(-x)}{e^{-\beta(x-\mu)}+1} , \quad (6)$$

where $\theta(x)$ is the unit step function and v^μ is the four-velocity of the center of mass of the medium. Since the momentum transfer is a spacelike four-vector it is correct to take the static limit in Eq. (2), which is defined by setting $q_0 = 0; |\vec{q}| \to 0$. The static limit of the polarization tensors G and F given by Eq. (3) can be obtained as [12,13]

$$G(q_0 = 0; |\vec{q}| \to 0) = 0 ,$$
$$F(q_0 = 0; |\vec{q}| \to 0) = m_{\rm el}^2 , \quad (7)$$

where $m_{\rm el}$ is the inverse Debye length corresponding to the well-known phenomenon of screening of static electric fields. Calculation of the FTD part of the trace factor in Eq. (2) shows that only the D^{00} component of the vector boson propagator survives. This is of importance here since, according to Eq. (7), spacelike transverse excitations are not screened (absence of screening for static magnetic fields).

Now, the contribution to the external potential for any neutrino type can be read off from Eq. (2):

$$V_5^r = (q_{\nu_r} q_e / m_{\rm el}^2) N_e , \quad (8)$$

where N_e is the number density of electrons in the Sun and for classical and nonrelativistic electron gases $m_{\rm el}$ is given by

$$m_{\rm el}^2 = \frac{q_e^2}{T} N_e . \quad (9)$$

Hence, V_5^r depends only on the ratio of charges but not on their absolute values, and is given by

$$V_5^r = (q_{\nu_r}/q_e)T . \quad (10)$$

If all charges q_{ν_r} are equal then the effective potential (the crucial quantity in the theory of neutrino oscillations in matter; see, e.g., [14]) vanishes and there is no new contribution to the MSW effect. The effective potential in the standard MSW theory can be calculated to be [15]

$$\delta V^e = \sqrt{2}G_F(N_e - N_n) \quad (\text{for } \nu_e - \nu_s \text{ oscillations}) ,$$
$$\delta V^e = \sqrt{2}G_F N_e \quad (\text{for } \nu_e - \nu_{\mu(\tau)} \text{ oscillations}) , \quad (11)$$

where N_n is the number density of neutrons. For the case of nonequal charges, one finds that the standard MSW pattern remains valid if

$$\left|\frac{q_\nu}{q_e}\right| \lesssim G_F(N_e/T) , \quad (12)$$

where q_ν may represent q_{ν_e} (if $q_{\nu_e} \gg q_s$) for the case of ν_e-ν_s oscillations or, for example, q_{ν_μ} (if $q_{\nu_\mu} \gg q_{\nu_e}$) for ν_e-ν_μ oscillations. Depending on the position inside the Sun, one gets a strong bound

$$\left|\frac{q_\nu}{q_e}\right| \lesssim (1-5)\times 10^{-15} . \quad (13)$$

If all charges are equal except that for ν_s, one finds that, even for masses $m_s \sim 100$ eV (an excellent warm dark-matter candidate; see Ref. [16]), the resonant temperature is still lower than the temperature of the solar surface.

Even for the case of equal neutrino charges [or if the bound given by Eq. (13) is respected] there is an extra contribution to V_5 as shown in Fig. 2. It is due to the fact that electrons in the Sun can induce a nonzero vertex function for neutrinos. The contribution depicted in Fig. 2(a) affects the ν_e only, whereas the contribution shown in Fig. 2(b) affects all flavors equally. Hence, the latter is relevant for ν_e-ν_s oscillations only, but at the same time it is suppressed by an additional factor of $-1 + 4\sin 2\theta_w$. Following the notation of Ref. [17] we find that the diagram shown in Fig. 2(a) gives rise to the self-energy (considering the FTD part only)

$$-i\Sigma_5^{e,\rm ind}(k) = q_e\Gamma_\alpha^{(W)}(k,k,v)iD^{\alpha\beta}(q=0) \times \int \frac{d^4p}{(2\pi)^4}{\rm tr}\{\gamma_\beta iS^e(p)\} . \quad (14)$$

In the rest frame of the medium, which is defined by setting $v^\mu = (1,\vec{0})$, the off-shell vertex function can be written as

$$\Gamma_\alpha^{(W)} = \tau_{\alpha\beta}^{(W)}\gamma^\beta L \equiv \tau_{\alpha\beta}^{(W)}\gamma^\beta \frac{1}{2}(1-\gamma_5) \ , \qquad (15)$$

with

$$\tau_{\alpha\beta}^{(W)} = \tau_T^{(W)} R_{\alpha\beta} + \tau_L^{(W)} Q_{\alpha\beta} + \tau_P^{(W)} P_{\alpha\beta} \ , \qquad (16)$$

such that

$$\begin{aligned} R_{\alpha\beta} &= \tilde{g}_{\alpha\beta} - Q_{\alpha\beta} \ , \\ Q_{\alpha\beta} &= \frac{\tilde{v}_\alpha \tilde{v}_\beta}{\tilde{v}^2} \ , \\ P_{\alpha\beta} &= \frac{i}{|\vec{q}|}\varepsilon_{\alpha\beta\gamma\delta} q^\gamma v^\delta \end{aligned} \qquad (17)$$

and

$$\begin{aligned} \tilde{g}_{\alpha\beta} &= g_{\alpha\beta} - \frac{q_\alpha q_\beta}{q^2} \ , \\ \tilde{v}_\alpha &= \tilde{g}_{\alpha\beta} v^\beta \ . \end{aligned} \qquad (18)$$

The form factors of the neutrino can be written as

$$\begin{aligned} \tau_T^{(W)} &= 2\sqrt{2} q_e G_F \left(A - \frac{B}{\tilde{v}^2}\right) \equiv \tau_T^{(W)}(q_0, |\vec{q}|) \ , \\ \tau_L^{(W)} &= 4\sqrt{2} q_e G_F \frac{B}{\tilde{v}^2} \equiv \tau_L^{(W)}(q_0, |\vec{q}|) \ , \\ \tau_P^{(W)} &= 4\sqrt{2} q_e |\vec{q}| C \equiv \tau_P^{(W)}(q_0, |\vec{q}|) \ , \end{aligned} \qquad (19)$$

where

$$\begin{aligned} A &= \int \frac{d^3p}{(2\pi)^3 2E}(f_- + f_+)\left\{\frac{2m^2 - 2p\cdot q}{q^2 + 2p\cdot q} + (q \to -q)\right\}, \\ B &= \int \frac{d^3p}{(2\pi)^3 2E}(f_- + f_+) \\ &\quad \times \left\{\frac{2(p\cdot v)^2 + 2(p\cdot v)(q\cdot v) - p\cdot q}{q^2 + 2p\cdot q} \right. \\ &\quad \left. + (q \to -q)\right\} \ , \\ C &= \int \frac{d^3p}{(2\pi)^3 2E}(f_- - f_+)\frac{p\cdot\tilde{v}}{\tilde{v}^2} \\ &\quad \times \left\{\frac{1}{q^2 + 2p\cdot q} + (q \to -q)\right\} \ . \end{aligned} \qquad (20)$$

Here $f_\mp$ are the electron and positron distributions: i.e.,

$$f_\mp(p) = \frac{1}{e^{\beta(p_0 \mp \mu)} + 1} \ . \qquad (21)$$

In the static limit we find that

$$\begin{aligned} \tau_{T,P}^{(W)} &\to 0 \ , \\ \tau_L^{(W)} &\to 4\sqrt{2} q_e G_F B(q_0 = 0; |\vec{q}| \to 0) \ , \end{aligned} \qquad (22)$$

and

$$\Gamma_\alpha^{(W)} \to g_{\alpha 0}\tau_L^{(W)}\gamma_0 L \ , \qquad (23)$$

where the form factor B is given by the integral [17]

$$B(q_0 = 0; |\vec{q}| \to 0) = \frac{1}{2}\int \frac{d^3p}{(2\pi)^3}\frac{d}{dE}(f_- + f_+) \ . \qquad (24)$$

One can be convinced that, in the static limit, the form factor B is proportional to $m_{\rm el}^2$. Indeed, for classical and nonrelativistic electron gases ($f_+ \simeq 0, f_- \simeq e^{-(E-\mu)/T}$),

$$(-4)q_e^2 B(q_0 = 0; |\vec{q}| \to 0) = \frac{q_e^2}{T}N_e = m_{\rm el}^2 \ . \qquad (25)$$

Finally, we find the result for $V_5^{e,\rm ind}$ as

$$V_5^{e,\rm ind} = -\sqrt{2}G_F N_e \ , \qquad (26)$$

which is independent of the electron charge q_e. It is interesting to note that for active-active oscillations in this case [compare with Eq. (11)] there is no MSW effect at all.

Now, we turn to the case where protons in our Galaxy act as a source of the new force. For the electron gas in the Sun (i.e., a plasma with a positive ion background) we have the electrical neutrality condition

$$N_{^1\rm H} + N_{^4\rm He} \simeq N_e \ , \qquad (27)$$

but everywhere in the Sun $N_{^1\rm H} > N_{^4\rm He}$. Ignoring the effect of nuclei heavier than ^{1}H, one gets the same constraint as given by Eq. (13) (with the replacement $q_e \to q_p$). In addition, only the diagram shown in Fig. 2(b) contributes in this case.

Finally, let us consider a possibility that a source of the new field are dark matter particles. Their density distribution in our Galaxy probably behaves as r^{-2} as unobserved mass increasing linearly with r must exist. The success of big bang nucleosynthesis in predicting the abundance for all of the light elements in the universe gives confidence that the ratio of the mass density in ordinary baryonic matter to the density required to just close the universe is $0.01 \leq \Omega_b \leq 0.08$ [18,19]. However, observations give much less visible matter ($\Omega_{\rm vis} \simeq 0.007$) which means that there could mathematically be much unseen baryonic matter in our Galaxy since observations indicate that $\Omega_{\rm halo} \simeq 0.1$. This is most likely in form of "Jupiters," but not gas, dust, or "snowballs." Hence, for the case of baryonic dark matter in our Galaxy, the previous bound given by Eq. (13) still applies. However, it seems more likely that some of whatever dark matter makes the global density up to the critical value also appears in the galactic halo whose local density is 0.20 GeV/cm$^3 \leq \rho^h \leq 0.43$ GeV/cm^3 [20]. Cold dark-matter particles in the galactic halo passing by the neighborhood of the Sun may be trapped when they cross the Sun, and can accumulate in the course of time. Besides, after trapping, repeated collisions with solar nuclei give the dark-matter particles a quasithermal distribution in the core of the Sun. In such a scenario the relic particles must have a mass m_x larger than 4 GeV, otherwise they will quickly evaporate [21] from the Sun.

The number density of dark-matter particles in the core of the Sun at present may be estimated to be [22]

$$N_x \simeq f_\sigma \frac{\rho_{0.3}}{\tilde{v}_{300}} \frac{7.2\times10^{-11}}{m_x} \frac{\alpha}{\alpha+Y_{\rm now}} (N_{^1{\rm H}} + N_{^4{\rm He}}) \ , \qquad (28)$$

where $Y_{\rm now}$ is the present dark-matter antiparticle abundance, $\rho_{0.3} = (\rho^h_x + \rho^h_{\bar{x}})/0.3\ {\rm GeV\,cm^{-3}}$, $\tilde{v}_{300}$ is the rms particle velocity in the halo (expressed in 300 $\rm km\,s^{-1}$) and m_x is expressed in GeV. f_σ is defined as

$$f_\sigma \simeq 0.89(\sigma_{\rm eff}/\sigma_C), \quad \sigma_{\rm eff} < \sigma_C \ ,$$
$$f_\sigma \simeq 1, \quad \sigma_{\rm eff} > \sigma_C \ , \qquad (29)$$

where $\sigma_{\rm eff}$ is the effective scattering cross section in the Sun and $\sigma_C = 4\times10^{-36}$ cm^2. The cosmic asymmetry parameter α is defined as

$$\alpha = \frac{(N_x - N_{\bar{x}})_{\rm relic}}{2s} \qquad (30)$$

and s is the entropy density. We have assumed the relic particle asymmetry ($N_x \gg N_{\bar{x}}$) both in order to have the MSW potential different from zero (it is proportional to $N_x - N_{\bar{x}}$) and to suppress appreciable dark-matter annihilation in the Sun [note the α dependence of Eq. (28)].

Now, the contribution to V^r_5 from the diagram shown in Fig. 1 (but now with dark-matter particles circulating inside the loop instead of electrons) reads

$$V^r_5 = (q_{\nu_r} q_x/m^2_{\rm el}) N_x \ , \qquad (31)$$

with

$$m^2_{\rm el} = \frac{1}{3}\sum_r q^2_{\nu_r} T^2_{\nu_r} \langle v^{-1}_{\nu_r}\rangle + q^2_x \frac{N_x}{T_x} \ , \qquad (32)$$

and $\langle v^{-1}_{\nu_r}\rangle = m_{\nu_r}/3.15 T_{\nu_r}$ if $m_{\nu_r} \gtrsim T_{\nu_r}$, otherwise $\langle v^{-1}_{\nu_r}\rangle \simeq 1$. Here we have included screening effects by the cosmic neutrino background [first term in Eq. (32)] with $T_\nu = 1.9$ K since the number density N_x is much less than the number density of ordinary solar matter constituents. Note that the first term in Eq. (32) can be dominated by a neutrino weighing about 10 eV. Namely, such a neutrino is by far the most plausible candidate for the hot dark part of the matter and then $T_\nu \simeq 50$ K [23]. Because the cosmic neutrino background is likely CP symmetric it enters Eq. (32) only through $m^2_{\rm el}$.

If the second term in Eq. (32) still dominates the first one, the bound given by Eq. (13) again applies (with $q_e \to q_x$) since $T_x \simeq T_{\rm core}$. In fact, only the upper value (which corresponds to the solar core) in Eq. (13) is relevant as the distribution of dark-matter particles inside the Sun is roughly barometric: i.e.,

$$N_x(r) = N_x(0)\exp\{-r^2/r^2_0\} \ , \qquad (33)$$

with a scale height r_0,

$$r_0 \simeq 0.04 R_{\rm Sun}\left(\frac{10 m_p}{m_x}\right)^{1/2} \ . \qquad (34)$$

Hence, in the main body of the Sun one expects the first term in Eq. (32) to dominate.

Finally, in the more general case where all these background particles possess a nonvanishing leptonic charge (but neglecting the existence of baryonic photons), the external neutrino potential can be written as a sum of contributions given by Eqs. (8) and (31), and $m^2_{\rm el}$ is a sum of contributions given by Eqs. (9) and (32). Since $N_e \gg N_x$ [see Eq. (28)], V^r_5 would be dominated by the electron contribution unless

$$\left|\frac{q_x}{q_e}\right| \gtrsim \frac{N_e}{N_x} \ . \qquad (35)$$

In conclusion, we have investigated the effect of novel long-range forces, constrained by delay limits inferred from the $\bar{\nu}_e$ burst from SN 1987A, on the MSW effect in the Sun. We have found that extra neutrino potentials in the Sun either depend on the ratio of the neutrino charge to the charge of a source particle or are charge independent. In order to retain the MSW effect as predicted by the standard theory, the value of the neutrino charge should be much less than the charge of a source particle.

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