MSW Oscillations - LMA and Subdominant Effects

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Abstract. New physics near the TeV scale could modify neutrino-matter interactions or
generate a relatively large neutrino magnetic (transition) moment. Both types of effects have
been discussed since the 1970’s as alternatives to mass-induced neutrino flavor oscillations.
Nowadays, the availability of high-statistics data makes it possible to turn the idea around and
ask: How well do the simple mass-induced oscillations describe solar neutrinos? At what level
are the above-mentioned nonstandard effects excluded? Can we use solar neutrinos to constrain
physics beyond the Standard Model? These notes review the sensitivity of the present-day
solar neutrino experiments to the nonstandard neutrino interactions and transition moment
and outline progress that may be expected in the near future. Based on a talk given at the
Neutrino 2006 conference [1].

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1. Standard LMA solution: basic features
The most basic experimental fact about the neutrinos from the Sun is that the electron neutrino
survival probability, \( P_{\text{std}}^{\nu_e} \equiv P(\nu_e \rightarrow \nu_e) \), is measured to vary as a function of the neutrino
energy. At the high end of the spectrum (\( E_\nu \gtrsim 6 - 7 \) MeV) the SNO [2] and Super-Kamiokande [3] experiments have established that \( P_{\text{std}}^{\nu_e} \) is about \(~ 34 \pm 3\%\). The gallium experiments [4],
however, which are sensitive to both high- and low-energy neutrinos, see a higher survival
probability: the measured rate is \(~ 74 \pm 7\) SNU, whereas the standard solar model prediction [5] (before oscillations) is \(~ 131^{+12}_{-10}\).

This simple fact has highly non-trivial implications. Indeed, this behavior is not “generic”
for mass-induced oscillations, even when they combine with the MSW [6, 7] matter effect. \textit{A priori}, one might have expected solar neutrinos to be in one of these regimes:

- matter dominates at the production point, on the way out of the Sun neutrino flavor evolves
  adiabatically \( \rightarrow \) constant suppression (regime 1);
- matter dominates at the production point, on the way out of the Sun neutrino flavor evolves
  non-adiabatically \( \rightarrow \) vacuum oscillations
  \begin{itemize}
    \item vacuum oscillation length \( \ll 1 \) a.u. (astronomical unit) \( \rightarrow \) oscillations average out \( \rightarrow \)
    constant observed suppression (regime 2);
    \item vacuum oscillation length \( \gg 1 \) a.u. (astronomical unit) \( \rightarrow \) no time to oscillate \( \rightarrow \) no
    suppression (regime 3);
  \end{itemize}
- vacuum oscillations dominate everywhere, matter effects negligible even in the center of the
  Sun \( \rightarrow \) oscillations average out \( \rightarrow \) constant observed suppression (regime 4).
The observed energy-dependent $P_{\text{std}}$ then implies that solar neutrinos are in one of the several “special regimes”: the transition between regimes 1 and 4 (Large Mixing Angle – LMA – solution); the transition between regimes 2 and 3 (vacuum/quasi-vacuum oscillation solution); the transition between regimes 1 and 2 (Small Mixing Angle – SMA – solution); the regime where the density in the Earth is close to resonant, so that the flavor regeneration in the Earth is large (the LOW solution). These solutions were known for many years, in particular all four were allowed as recently as 2000, see, e.g., [9].

We now of course know that only the LMA solution survives. Let us consider the survival probability $P_{\text{std}}$ under the (a posteriori justified) assumption that the oscillations take place between just two eigenstates. One easily obtains that during the day time

$$P_{\text{std}}, 2\nu = \cos^2 \theta_{\odot} \cos^2 \theta + \sin^2 \theta_{\odot} \sin^2 \theta.$$

The probability of finding the neutrino in eigenstate 1(2) is $\cos^2 \theta_{\odot} \cos^2 \theta$, where $\theta_{\odot}$ is the mixing angle at the production point; in turn, the probability of detecting the neutrino already in eigenstate 1(2) as $\nu_e$ is $\cos^2 \theta \sin^2 \theta$. The key physical ideas here are that the evolution is adiabatic (no level jumping) and incoherent (interferences between 1 and 2 disappear upon integration over energies for $\Delta m^2 \gtrsim 10^{-9} - 10^{-8}$ eV$^2$ [10,8] and over the production region).

The angle $\theta_{\odot}$ is determined from the oscillation Hamiltonian $H_{\text{tot}} = H_{\text{vac}} + H_{\text{mat}}$, where

$$H_{\text{vac}} = \begin{pmatrix} -\Delta \cos 2\theta & \Delta \sin 2\theta \\ \Delta \sin 2\theta & -\Delta \cos 2\theta \end{pmatrix}, \quad H_{\text{mat}} = \begin{pmatrix} \sqrt{2} G_F n_e & 0 \\ 0 & 0 \end{pmatrix}.$$  \hspace{1cm} \hspace{1cm} \hspace{1cm} \hspace{1cm} \hspace{1cm} (2)

Here $\Delta \equiv \Delta m^2/(4E_\nu)$ and $\Delta m^2$ is the mass splitting between the first and second neutrino mass states: $\Delta m^2 \equiv m_2^2 - m_1^2$. The two limiting values are $\theta_{\odot} = \theta$ ($H_{\text{tot}}$ is dominated by $H_{\text{vac}}$) and $\theta_{\odot} = \pi/2$ ($H_{\text{tot}}$ is dominated by $H_{\text{mat}}$). The probability $P_{\text{std}}$ then varies from $\cos^4 \theta + \sin^4 \theta$ ($= 1 - (1/2) \sin^2 2\theta$, averaged vacuum oscillations) to $\sin^2 \theta$.

![Figure 1. The $\nu_e$ survival probability and day/night asymmetry for the LMA solution.](image_url)
large day/night variations of $P_{ee}$. The situation is illustrated in Fig. 1. Evidently, Nature chose to “tune” the mass splitting involved in solar neutrino oscillations to the density in the solar core! Remarkably, a completely independent reactor antineutrino measurement by KamLAND [11] showed that $\Delta m^2$ is indeed in this range.

The preceding discussion assumed that mass eigenstate 3 is not involved in the evolution of solar neutrinos. The correction due its presence is trivially computed if we notice that the splitting between this state and eigenstates 1 and 2 is significantly larger than the matter potential even in the center of the Sun ($\Delta m^2_{\text{atm}}/2E \gg \sqrt{2}G_F N_e(0)$), so that the $\nu_e$ content of that state is always given by $\sin^2 \theta_{13}$. Repeating the arguments that led to Eq. (1), one gets

$$P_{ee}^{\text{std, 3}\nu} = \sin^4 \theta_{13} + \cos^4 \theta_{13} P_{ee}^{\text{std, 2}\nu}.$$  

(3)

Given the bound $\sin^2 \theta_{13} \lesssim 0.02$ from CHOOZ [12], the first term is negligibly small. The effect of the third state then is to multiply the two-neutrino survival probability by $\cos^4 \theta_{13}$. The resulting correction is at most 4%; this correction is basically the probability that the original electron neutrino “disappears” into state 3. See, e.g., [13, 14] for recent data analyses and further references.

2. Searching for nonstandard neutrino interactions

The impact of nonstandard neutrino interactions with matter on solar neutrino oscillations was discussed already in the classical paper by L. Wolfenstein [6] and subsequently elaborated on by many authors ([15, 16, 17] and many others). The idea is that novel interactions due to a heavy charged lepton could modify the neutrino forward scattering amplitude and hence the oscillation Hamiltonian in matter. Regardless of their origin, at low energies relevant to neutrino oscillations, nonstandard interactions (NSI) are described by the effective Lagrangian

$$L^{\text{NSI}} = -2\sqrt{2}G_F (\bar{\nu}_\alpha \gamma_{\mu} \nu_\beta)(\epsilon^{fL}_{\alpha\beta} \bar{f}_L \gamma^\mu f_L + \epsilon^{fR}_{\alpha\beta} \bar{f}_R \gamma^\mu f_R) + \text{h.c.}$$  

(4)

Here $\epsilon^{fL}_{\alpha\beta}$ ($\epsilon^{fR}_{\alpha\beta}$) denotes the strength of the NSI between the neutrinos $\nu$ of flavors $\alpha$ and $\beta$ and the left-handed (right-handed) components of the fermions $f$.

Neutrino scattering tests, like those of NuTeV [18] and CHARM [19], constrain mainly the NSI couplings of the muon neutrino, e.g., $|\epsilon_{\mu\mu}| \lesssim 10^{-3}$, $|\epsilon_{\mu\mu}| \lesssim 10^{-3} - 10^{-2}$. In contrast, direct limits on $\epsilon_{ee}$, $\epsilon_{\tau\tau}$, and $\epsilon_{\tau\tau}$ are remarkably loose, e.g., $|\epsilon_{\tau\tau}| < 3$, $-0.4 < \epsilon_{ee} < 0.7$, $|\epsilon_{e\tau}| < 0.5$, $|\epsilon_{e\tau}| < 0.5$ [20]. Stronger constraints exist on the corresponding interactions involving the charged leptons. Those, however, are model-dependent and do not apply if the NSI come from the underlying operators containing the Higgs fields [21]. Here we only consider direct experimental bounds.

Even with the addition of NSI the splitting $\Delta m^2_{\text{atm}}/2E$ remains much greater than the matter potential anywhere along the neutrino trajectory. This means the solar neutrino analysis can still be reduced to two neutrino states, following the arguments of Sect. 1. Neglecting small corrections of order $\sin \theta_{13}$ or higher, the corresponding matter contribution to the two-neutrino oscillation Hamiltonian can be written as

$$H^{\text{NSI}}_{\text{mat}} = \frac{G_F n_e}{\sqrt{2}} \left( \begin{array}{cc} 1 + \epsilon_{11} & \epsilon_{12}^* \\ \epsilon_{12} & -1 - \epsilon_{11} \end{array} \right), \quad \text{where} \quad \epsilon_{11} = \epsilon_{ee} - \epsilon_{\tau\tau} \sin^2 \theta_{23}, \quad \epsilon_{12} = -2\epsilon_{e\tau} \sin \theta_{23}. \quad (5)$$

The epsilons are the sums of the contributions from the matter constituents: $\epsilon_{\alpha\beta} \equiv \sum f = u, d, e \epsilon^f_{\alpha\beta} n_f/n_e$. In turn, $\epsilon_{\alpha\beta}^f \equiv \epsilon^{fL}_{\alpha\beta} + \epsilon^{fR}_{\alpha\beta}$. Observe that only the vector component of the NSI enters the propagation effect; in contrast, the NC detection process at SNO depends on the axial coupling. The propagation and detection effects of the NSI are thus sensitive to different parameters, and the corresponding searches could be complementary.
Eq. (5) shows that the flavor changing NSI effect in solar neutrino oscillations comes from \( \epsilon_{\mu\tau} \), while the flavor preserving NSI effect comes from both \( \epsilon_{ee} \) and \( \epsilon_{\tau\tau} \). A useful parameterization is

\[
H_{\text{mat}}^{\text{NSI}} = \begin{pmatrix}
A \cos 2\alpha & A e^{-2i\phi} \sin 2\alpha \\
A e^{2i\phi} \sin 2\alpha & -A \cos 2\alpha
\end{pmatrix},
\]

where \( \tan 2\alpha = |\epsilon_{12}|/(1 + \epsilon_{11}) \), \( 2\phi = \text{Arg}(\epsilon_{12}) \), and

\[
A = G_F n_e \sqrt{[1 + \epsilon_{11}]^2 + |\epsilon_{12}|^2}/2.
\]

The effect of \( \alpha \) is to change the mixing angle in the medium of high density from \( \pi/2 \) to \( \pi/2 - \alpha \). The angle \( \phi \) (related to the phase of \( \epsilon_{\mu\tau} \)) is a source of CP violation. Solar neutrino experiments, just like terrestrial beam experiments [22, 23], are sensitive to its effects [24], while atmospheric neutrinos are not [25, 26].

\[\text{Figure 2.}\] The electron neutrino survival probability (left) and day/night asymmetry (right) for \( \Delta m^2 = 7 \times 10^{-5} \text{ eV}^2 \), \( \tan^2 \theta = 0.4 \) and several representative values of the NSI parameters:

1. \( \epsilon_{11}^e = \epsilon_{12}^d = \epsilon_{12}^u = 0 \); (2) \( \epsilon_{11}^e = \epsilon_{11}^d = -0.008, \epsilon_{12}^d = -0.06 \); (3) \( \epsilon_{11}^e = \epsilon_{11}^d = -0.044, \epsilon_{12}^d = 0.14 \); (4) \( \epsilon_{11}^e = \epsilon_{11}^d = -0.044, \epsilon_{12}^d = -0.14 \). Reproduced from [24].

The main effects of NSI on \( P_{ee} \) are as follows [24]: (i) the low-energy limit stays the same (vacuum oscillations); (ii) the high-energy limit changes, according to Eq. (1), \( P_{ee} \to \sin^2 \alpha \cos^2 \theta + \cos^2 \alpha \sin^2 \theta \); (iii) at intermediate energies, the transition from vacuum to matter dominated regime can shift in energy, with changing \( A \), and can become more or less abrupt, with changing \( \alpha \) and \( \phi \). The nonadiabatic regime occurs when \( \theta \to \alpha \), rather than \( \theta \to 0 \). Also, very importantly, the day/night effect can change with all three parameters. In particular, it becomes small either as \( A \to 0 \) [27] [28] or as \( \alpha \to \theta \) [24]. Thus, the LMA-0 region that is normally excluded by the non-observation of day/night asymmetry may become allowed [24] [27] [28].

Fig. 2 illustrates the impact of the NSI on \( P_{ee} \) and the day/night asymmetry. Curve 3 gives an example of parameters that can already be excluded by the current data. Curve 4 illustrates the suppression of the Earth effect described above. For technical details, including approximate analytical expressions for \( P_{ee} \) and day/night asymmetry, see [24].

The solar neutrino analysis of NSI cannot be done in isolation: the same NSI can also be probed with atmospheric neutrinos. Indeed, on general grounds, one expects the atmospheric neutrinos – particularly the high energy ones for which nonstandard matter effects can dominate over the vacuum oscillation effects – to be a very sensitive probe of NSI. Early two-neutrino (\( \nu_{\mu}, \nu_{\tau} \)) numerical studies [23] yielded \( \epsilon_{\mu\tau} \lesssim 0.08 - 0.12 \) and \( \epsilon_{\tau\tau} \lesssim 0.2 \). Clearly, these are very strong bounds; if they were to extend to \( \epsilon_{ee} \), the large NSI effects on solar neutrinos discussed above would be excluded. It turns out, however, that this is not the case: when the analysis is

\[1\] Notice the difference in normalization: our \( \epsilon \)'s are normalized per electron, while [29] gives \( \epsilon \)'s per \( d \) quark.
Figure 3. Left panel: A 2-D section ($\epsilon_{ee} = -0.15$) of the allowed region of the NSI parameters (shaded). We assumed $\Delta m^2_{23} = 0$ and $\theta_{13} = 0$, and marginalized over $\theta$ and $\Delta m^2$. The dashed contours indicate our analytical predictions. See text for details. Right panel: The effect of the NSI on the allowed region and best-fit values of the oscillation parameters. From [25].

properly extended to three flavors, one finds that very large values of both $\epsilon_{e\tau}$ and $\epsilon_{\tau\tau}$ are still allowed by the data [25]. This is illustrated in Fig. 3 (left panel), which shows that NSI with strengths comparable to the Standard Model interactions can be compatible with all atmospheric data. It must be noted that the compatibility is achieved as a result of adjusting the vacuum oscillation parameters: large NSI imply a smaller mixing angle and larger $\Delta m^2_{atm}$, as can be see in the right panel of Fig. 3.

The addition of the K2K data helps constrain the allowed NSI region somewhat [26]. While the addition of the first data from MINOS brings no further improvement [30], the future high-statistics MINOS dataset will be a very valuable probe of this parameter space [30].

3. Searching for neutrino transition moments
The idea that solar neutrinos could be affected by the neutrino spin precession (NSP) in the solar magnetic fields is even older [31] than the NSI idea. Remarkably, this idea – much improved with time [32, 33, 34, 35, 36, 37] – remained viable for the next three decades. While, by the late 1990’s, the lack of time variations in the Super-Kamiokande data gave strong evidence against large NSP in the solar convective zone, NSP in the radiative zone continued to give a good fit to all solar data [38].

Even after the confirmation of the LMA oscillation solution by KamLAND [11], the possibility of the NSP happening at a subdominant level remains of great interest, as a probe of the neutrino electromagnetic properties and, at the same time, of the magnetic fields in the solar interior. NSP coupled with flavor oscillations could lead to conversion $\nu_e \rightarrow \bar{\nu}_e$ in the Sun, on which recently, KamLAND [39] reported an upper bound. It is very important to understand what this bound implies for the neutrino magnetic (transition) moment and how it compares with other available bounds on the neutrino transition moment.

The laboratory bounds on the neutrino magnetic (transition) moment come from measuring the cross sections of $\nu e$ or $\bar{\nu}e$ scattering in nearly forward direction. The recent bound for the interaction involving the Majorana electron antineutrino is $2\mu_{e\beta} < 0.9 \times 10^{-10}\mu_B$ at the 90% confidence level [43], where $\mu_B \equiv e/(2m_e)$ is the Bohr magneton ($m_e$ is the electron mass, $e$ is its charge). Stronger bounds, $\mu \lesssim 3 \times 10^{-12}\mu_B$, exist from astrophysical considerations, particularly from the study of red giant populations in globular clusters [44]. Larger values of the transition moment would provide an additional cooling mechanism – via plasmon decay to $\nu \bar{\nu}$ – for the
red giant core and change the core mass at helium flash beyond what is observationally allowed.

An important theoretical consideration is that the magnetic moment operators will radiatively generate the neutrino mass. Requiring that the corresponding contribution to the neutrino masses be not much greater than their observed values, one obtains a model-independent “naturalness” bound on $\mu_{ab}$. For Dirac neutrinos, the bound obtained in this way turns out to be very stringent, $\mu_{ab} \lesssim 10^{-14} \mu_B$ [45]. However, very importantly, for Majorana neutrinos the bound is much weaker, only $\mu_{ab} \lesssim 10^{-10} \mu_B$ [46, 47], owing to the different flavor symmetry properties of the mass and the transition moment operators [48]. In fact, explicit models exploiting these very symmetry properties were discussed many years ago (e.g., [49, 50]).

It turns out that for the measured LMA oscillation parameters NSP in the radiative zone cannot produce the $\bar{\nu}_e$ flux above the KamLAND bound. This is illustrated in Fig. 4. This is a remarkable example that knowing neutrino oscillation parameters precisely can be very valuable: the answer would qualitatively change if the mixing angle were $20^\circ$ instead of $30^\circ$.

For NSP in the convective zone, the analysis is very different, though in the end the conclusion is similar: one should not have expected the flux of $\nu_e$ in excess of the published KamLAND bound. Put another way, the bound on the neutrino transition moment from the KamLAND bound is comparable to the direct laboratory and “naturalness” bounds, but still weaker than that from analysis of the red giant cooling. An updated analysis of a larger KamLAND dataset is needed. The reader is referred to [42] for details and further references.

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Figure 4. The regions of the oscillation parameter space where one may expect the electron antineutrino flux above the KamLAND bound [39] (three different shadings correspond to three different normalizations of the magnetic field, up to the upper bound [40]). An optimistic value of the transition moment, $\mu = 1 \times 10^{-11} \mu_B$, was taken. For comparison, the region allowed by the combined analysis of the KamLAND and solar neutrino data [41] is also shown. From [42].
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