Non-Standard Interaction Effects at Reactor Neutrino Experiments

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Abstract

We study non-standard interactions (NSIs) at reactor neutrino experiments, and in particular, the mimicking effects on θ_{13}. We present generic formulas for oscillation probabilities including NSIs from sources and detectors. Instructive mappings between the fundamental leptonic mixing parameters and the effective leptonic mixing parameters are established. In addition, NSI corrections to the mixing angles θ_{13} and θ_{12} are discussed in detailed. Finally, we show that, even for a vanishing θ_{13}, an oscillation phenomenon may still be observed in future short baseline reactor neutrino experiments, such as Double Chooz and Daya Bay, due to the existences of NSIs.

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I. INTRODUCTION

Neutrino oscillations have successfully turned into the most plausible description of neutrino flavor transitions. At the moment, the most important task in neutrino physics is to accurately determine the neutrino parameters, especially the neutrino mass hierarchy and the leptonic mixing angle θ_{13}. In this work, we will concentrate on the leptonic mixing parameters, and in particular, the parameter θ_{13}. Recently, so-called non-standard interactions (NSIs) have been extensively studied in the literature. Such NSIs could affect neutrino oscillations in a similar way as normal matter affects them. Therefore, if present, NSIs will affect the determination of the fundamental neutrino parameters.

In this work, we will mainly investigate measurements of the fundamental leptonic mixing angles θ_{13} and θ_{12} at reactor neutrino experiments. Since reactor neutrino experiments such as the future Double Chooz [1] and Daya Bay [2] experiments as well as the existing KamLAND experiment [3] have relatively short baseline lengths, normal matter effects are negligible. This also holds for NSI effects during propagation of neutrinos. Thus, we will only assume that the NSI effects are present at sources and detectors.

In Ref. [4], the basic formalism and different neutrino states for source and detector including NSIs (or "new physics") were first presented. Later, NSIs in sources and detectors have been discussed using amplitudes that describe the neutrino sources and detectors. Such works have been carried out for long baseline neutrino oscillation experiments in general [5] as well as for neutrino factories in particular [6]. Recently, a study on the impact of NSIs at reactor and accelerator neutrino experiments has been performed [7]. Especially, the authors of this work derive first-order series expansions for oscillation probabilities including NSIs from sources and detectors. Explicit upper bounds on parameters describing NSIs from sources and detectors exist. However, these bounds are only generic and given by $\varepsilon^s_{\alpha\beta} = \mathcal{O}(0.1)$ for NSIs at sources from universality in lepton decays and $\varepsilon^d_{\alpha\beta} = \mathcal{O}(0.2)$ for NSIs at detectors from universality in pion decays [8, 9].

This paper is organized as follows. In Sec. II, we will present general formulas for parameter mappings between the fundamental leptonic mixing parameters and the effective leptonic mixing parameters due to the effects of NSIs, and we will give expressions for oscillation probabilities. Then, in Sec. III, we will discuss reactor experiments and how these could be influenced by NSIs and what the outcome would be for the mixing angles θ_{13} and

θ_{12}. Finally, in Sec. IV, we will summarize our results and present our conclusions.

II. ANALYTIC FORMALISM

For a realistic neutrino oscillation experiment, in the presence of non-standard neutrino interactions, the neutrino states produced in the source and observed at the detector can be treated as superpositions of pure orthonormal flavor states:

$$|\nu_\alpha^s\rangle = \frac{1}{N_\alpha^s}\left(|\nu_\alpha\rangle + \sum_{\beta=e,\mu,\tau} \varepsilon_{\alpha\beta}^s |\nu_\beta\rangle\right) , \tag{1}$$

$$\langle\nu_\beta^d| = \frac{1}{N_\beta^d}\left(\langle\nu_\beta| + \sum_{\alpha=e,\mu,\tau} \varepsilon_{\alpha\beta}^d \langle\nu_\alpha|\right) , \tag{2}$$

where the superscripts 's' and 'd' denote the source and the detector, respectively, with the normalization factors being given by[1]

$$N_\alpha^s = \sqrt{\left[\left(\mathbb{1}+\varepsilon^s\right)\left(\mathbb{1}+\varepsilon^{s\dagger}\right)\right]_{\alpha\alpha}} , \tag{3}$$

$$N_\beta^d = \sqrt{\left[\left(\mathbb{1}+\varepsilon^{d\dagger}\right)\left(\mathbb{1}+\varepsilon^{d}\right)\right]_{\beta\beta}} . \tag{4}$$

Note that the states $|\nu_\alpha^s\rangle$ and $\langle\nu_\beta^d|$ are no longer orthonormal states because of NSIs. Since different physical processes take place at the source the and detector, the NSI parameter matrices ε^s and ε^d are arbitrary and non-unitary in general. In the minimal unitarity violation model (MUV) [10, 11, 12, 13, 14, 15], where the unitarity of the leptonic mixing matrix [16, 17] is slightly violated by possible new physics effects, the non-unitary effect can be regarded as one type of NSIs with the requirement $\varepsilon^s = \varepsilon^{d\dagger}$.[2]

Since in a terrestrial neutrino oscillation experiment, the Earth matter effects [18, 19] are more or less involved, the propagation of neutrino flavor states in matter is governed by the effective Hamiltonian

$$\begin{aligned}\hat{H} &= H_0 + H_m + H_{\rm NSI} \\ &= \frac{1}{2E} U \mathrm{diag}(m_1^2, m_2^2, m_3^2) U^\dagger + \mathrm{diag}(V_{\rm CC}, 0, 0) + V_{\rm CC}\varepsilon^m ,\end{aligned} \tag{5}$$

---

[1] Note that, in calculating the number of events, the normalization factors are canceled with the NSI factors in charged-current cross-sections. However, for running short baseline reactor neutrino experiments, the neutrino fluxes are directly measured by using a near detector, and not a Monte Carlo simulation. Hence, the normalization factors should be taken into account. (See also Ref. [10] for a detailed discussion.)

[2] In the MUV model, a neutral-current contribution cannot, in principle, be rewritten as a global phase in the oscillation amplitude, and thus, it affects the oscillation process.

where $V_{\rm CC} = \sqrt{2}G_F N_e$ arises from coherent forward scattering and N_e denotes the electron number density along the neutrino trajectory in the Earth. Different from ε^s and ε^d, ε^m is an exact Hermitian matrix describing NSIs in matter, and its current experimental bounds can be found in Refs. [20, 21, 22]. Here the superscript 'm' is used in order to distinguish NSI effects in the Earth matter from those in neutrino sources and detectors. The vacuum leptonic mixing matrix U is usually parametrized in the standard form by using three mixing angles and one CP violating phase [23]

$$
\begin{aligned}
U &= O_{23} U_\delta O_{13} U_\delta^\dagger O_{12} \\
&= \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}{\rm e}^{-{\rm i}\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}{\rm e}^{{\rm i}\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}{\rm e}^{{\rm i}\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}{\rm e}^{{\rm i}\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}{\rm e}^{{\rm i}\delta} & c_{23}c_{13} \end{pmatrix} ,
\end{aligned} \tag{6}
$$

where $U_\delta = {\rm diag}(1, 1, e^{i\delta})$, and O_{ij} is the orthogonal rotation matrix in the i, j plane with $c_{ij} \equiv \cos\theta_{ij}$ and $s_{ij} \equiv \sin\theta_{ij}$ (for $ij = 12$, 13 and 23). In analogy to the vacuum Hamiltonian H_0 in Eq. (5), the effective Hamiltonian in matter $\hat{H}$ can also be diagonalized through a unitary transformation

$$
\hat{H} = \frac{1}{2E} \hat{U} {\rm diag}\left(\hat{m}_1^2, \hat{m}_2^2, \hat{m}_3^2\right) \hat{U}^\dagger \ , \tag{7}
$$

where $\hat{m}_i^2$ $(i = 1, 2, 3)$ denote the effective mass squared eigenvalues of neutrinos and $\hat{U}$ is the effective leptonic mixing matrix in matter.

Now, we include all the NSI effects into the oscillation processes, and arrive at the amplitude for the process $\nu_\alpha^s \to \nu_\beta^d$

$$
\begin{aligned}
\mathcal{A}_{\alpha\beta}(L) &= \frac{1}{N_\alpha^s N_\beta^d} \langle \nu_\beta^d | {\rm e}^{-{\rm i}\hat{H}L} | \nu_\alpha^s \rangle = \frac{1}{N_\alpha^s N_\beta^d} (\mathbb{1} + \varepsilon^d)_{\rho\beta} A_{\gamma\rho} \left(\mathbb{1} + \varepsilon^s\right)_{\alpha\gamma} \\
&= \frac{1}{N_\alpha^s N_\beta^d} \left[(\mathbb{1} + \varepsilon^d)^T A^T \left(\mathbb{1} + \varepsilon^s\right)^T \right]_{\beta\alpha} = \frac{1}{N_\alpha^s N_\beta^d} \left[A + \varepsilon^s A + A\varepsilon^d + \varepsilon^s A \varepsilon^d \right]_{\alpha\beta} \ ,
\end{aligned} \tag{8}
$$

where L is the propagation distance and the explicit form of A is a coherent sum over the contributions of all the mass eigenstates ν_i

$$
A_{\alpha\beta} = \sum_i \hat{U}_{\alpha i}^* \hat{U}_{\beta i} {\rm e}^{-{\rm i}\frac{\hat{m}_i^2 L}{2E}} \ . \tag{9}
$$

Inserting Eq. (9) into Eq. (8), one can directly obtain

$$\begin{aligned}
\mathcal{A}_{\alpha\beta}(L) &= \frac{1}{N^s_\alpha N^d_\beta}\left[\sum_i \hat{U}^*_{\alpha i}\hat{U}_{\beta i}\mathrm{e}^{-\mathrm{i}\frac{\hat{m}^2_i L}{2E}} + \sum_{\gamma,i}\hat{U}^*_{\gamma i}\hat{U}_{\beta i}\varepsilon^s_{\alpha\gamma}\mathrm{e}^{-\mathrm{i}\frac{\hat{m}^2_i L}{2E}}\right. \\
&\quad \left. + \sum_{\gamma,i}\hat{U}^*_{\alpha i}\hat{U}_{\gamma i}\varepsilon^d_{\gamma\beta}\mathrm{e}^{-\mathrm{i}\frac{\hat{m}^2_i L}{2E}} + \sum_{\gamma,\rho,i}\varepsilon^s_{\alpha\gamma}\varepsilon^d_{\rho\beta}\hat{U}^*_{\gamma i}\hat{U}_{\rho i}\mathrm{e}^{-\mathrm{i}\frac{\hat{m}^2_i L}{2E}}\right] \\
&= \frac{1}{N^s_\alpha N^d_\beta}\sum_i \left[\hat{U}^*_{\alpha i}\hat{U}_{\beta i} + \sum_\gamma \varepsilon^s_{\alpha\gamma}\hat{U}^*_{\gamma i}\hat{U}_{\beta i}\right. \\
&\quad \left. + \sum_\gamma \varepsilon^d_{\gamma\beta}\hat{U}^*_{\alpha i}\hat{U}_{\gamma i} + \sum_{\gamma,\rho}\varepsilon^s_{\alpha\gamma}\varepsilon^d_{\rho\beta}\hat{U}^*_{\gamma i}\hat{U}_{\rho i}\right]\mathrm{e}^{-\mathrm{i}\frac{\hat{m}^2_i L}{2E}} \ . \qquad (10)
\end{aligned}$$

In order to compare Eq. (10) with the standard oscillation amplitude given in Eq. (9), we rewrite $\mathcal{A}_{\alpha\beta}(L)$ as

$$\mathcal{A}_{\alpha\beta}(L) = \sum_i \mathcal{J}^i_{\alpha\beta}\mathrm{e}^{-\mathrm{i}\frac{\hat{m}^2_i L}{2E}} \qquad (11)$$

with

$$\mathcal{J}^i_{\alpha\beta} = \frac{\hat{U}^*_{\alpha i}\hat{U}_{\beta i} + \sum_\gamma \varepsilon^s_{\alpha\gamma}\hat{U}^*_{\gamma i}\hat{U}_{\beta i} + \sum_\gamma \varepsilon^d_{\gamma\beta}\hat{U}^*_{\alpha i}\hat{U}_{\gamma i} + \sum_{\gamma,\rho}\varepsilon^s_{\alpha\gamma}\varepsilon^d_{\rho\beta}\hat{U}^*_{\gamma i}\hat{U}_{\rho i}}{N^s_\alpha N^d_\beta} \ . \qquad (12)$$

It can be clearly seen that only the αth row of ε^s and the βth column of ε^d are relevant to the transition amplitude. In the limit $\varepsilon \to 0$, Eq. (11) is reduced to the standard oscillation amplitude in matter.

With the definitions above, the oscillation probability is given by

$$\begin{aligned}
P(\nu^s_\alpha \to \nu^d_\beta) &= |\mathcal{A}_{\alpha\beta}(L)|^2 \\
&= \sum_{i,j}\mathcal{J}^i_{\alpha\beta}\mathcal{J}^{j*}_{\alpha\beta} - 4\sum_{i>j}\mathrm{Re}(\mathcal{J}^i_{\alpha\beta}\mathcal{J}^{j*}_{\alpha\beta})\sin^2\frac{\Delta\hat{m}^2_{ij}L}{4E} \\
&\quad +2\sum_{i>j}\mathrm{Im}(\mathcal{J}^i_{\alpha\beta}\mathcal{J}^{j*}_{\alpha\beta})\sin\frac{\Delta\hat{m}^2_{ij}L}{2E} \ . \qquad (13)
\end{aligned}$$

A salient feature of Eq. (13) is that, when $\alpha \neq \beta$, the first term in Eq. (13) is, in general, not vanishing, and therefore, a flavor transition would already happen at the source even before the oscillation process and is known as the zero-distance effect [24]. Although the effective mixing matrix in matter $\hat{U}$ is still unitary, the presences of NSIs in the source and detector prevent us from defining a unique CP invariant quantity like the standard Jarlskog invariant [25]. New CP non-conservation terms, which are proportional to the NSI parameters and have different dependences on L/E, will appear in the oscillation probability. Another peculiar feature in the survival probability is that, in the case of $\alpha = \beta$, CP violating terms in the last line of Eq. (13) should, in principle, not vanish. Note that Eq. (13) is also valid in the MUV model and could be very instructive for analyzing the CP violating effects in the MUV model in future long baseline experiments.

III. REACTOR NEUTRINO EXPERIMENTS

Reactor neutrino experiments with short or medium baselines are only sensitive to the survival probability $P(\bar{\nu}^s_e \to \bar{\nu}^d_e)$. The typical energy of antineutrinos produced in nuclear reactors is around a few MeV, which indicates that the Earth matter effects are extremely small and can safely be neglected. Hence, we take ($\hat{U} \simeq U$, $\hat{m}_i \simeq m_i$) or effectively set $V_{\rm CC} = 0$ in Eq. (5). As mentioned above, among all the NSI parameters, only $\varepsilon^s_{e\alpha}$ and $\varepsilon^d_{\alpha e}$ are relevant to our discussion. It has been pointed out that for realistic reactor neutrino experiments, the leading-order NSIs are of the $V \pm A$ type, and the relation $\varepsilon^s_{e\alpha} = \varepsilon^{d*}_{\alpha e}$ holds well [7]. Therefore, we assume $\varepsilon^s_{e\alpha} = \varepsilon^{d*}_{\alpha e} = |\varepsilon_{e\alpha}| {\rm e}^{{\rm i}\phi_{e\alpha}}$ in the current consideration and neglect the superscript 's' throughout the following parts of this work. It can be seen from Eq. (12) that the imaginary parts of the parameters $\mathcal{J}^i_{ee}$ disappear, and hence, the corresponding $\bar{\nu}^s_e \to \bar{\nu}^d_e$ oscillation is a CP conserved process.

Similar to the case without NSIs, one may define the effective mixing angles $\tilde{\theta}_{13}$ and $\tilde{\theta}_{12}$, in which all the NSI effects are included. For the smallest mixing angle $\tilde{\theta}_{13}$, we take $\alpha = \beta = e$ and $i = 3$ in Eq. (12) together with the standard parametrization defined by Eq. (6), and obtain the mapping between $\tilde{\theta}_{13}$ and θ_{13}

$$
\begin{aligned}
\tilde{s}^2_{13} \; = \; & s^2_{13} + 2 s_{13} c_{13} \left[s_{23} \cos(\delta - \phi_{e\mu}) |\varepsilon_{e\mu}| + c_{23} \cos(\delta - \phi_{e\tau}) |\varepsilon_{e\tau}| \right. \\
& \left. - s_{23} \cos(\delta - \phi_{ee} - \phi_{e\mu}) |\varepsilon_{ee}| |\varepsilon_{e\mu}| - c_{23} \cos(\delta - \phi_{ee} - \phi_{e\tau}) |\varepsilon_{ee}| |\varepsilon_{e\tau}| \right] \\
& + (s^2_{23} c^2_{13} - s^2_{13}) |\varepsilon_{e\mu}|^2 + (c^2_{23} c^2_{13} - s^2_{13}) |\varepsilon_{e\tau}|^2 \\
& + 2 s_{23} c_{23} c^2_{13} \cos(\phi_{e\mu} - \phi_{e\tau}) |\varepsilon_{e\mu}| |\varepsilon_{e\tau}| + \mathcal{O}(\varepsilon^3) \; , \qquad (14)
\end{aligned}
$$

where the third-order terms in ε are neglected. As for the effective mixing angle $\tilde{\theta}_{12}$, we take $\alpha = \beta = e$ and $i = 2$, and obtain

$$
\begin{aligned}
\tilde{s}^2_{12} \tilde{c}^2_{13} \; = \; & s^2_{12} c^2_{13} + 2 s_{12} c_{12} c_{13} \left[c_{23} \cos(\phi_{e\mu}) |\varepsilon_{e\mu}| - s_{23} \cos(\phi_{e\tau}) |\varepsilon_{e\tau}| \right] \\
& - 2 s^2_{12} s_{13} c_{13} \left[s_{23} \cos(\delta - \phi_{e\mu}) |\varepsilon_{e\mu}| + c_{23} \cos(\delta - \phi_{e\tau}) |\varepsilon_{e\tau}| \right] + \mathcal{O}(\varepsilon^2) \; . \qquad (15)
\end{aligned}
$$

Since the NSI parameters should not be comparable to the sizable mixing angle θ_{12}, only the first-order terms in ε are taken into account in Eq. (15).

With the help of the effective mixing angles $\tilde{\theta}_{13}$ and $\tilde{\theta}_{12}$, the survival probability reads

$$\begin{aligned} P(\bar{\nu}^s_e \to \bar{\nu}^d_e) \;=\; & 1 - \cos^4\tilde{\theta}_{13}\sin^2 2\tilde{\theta}_{12}\sin^2\frac{\Delta m^2_{21}L}{4E} - \cos^2\tilde{\theta}_{12}\sin^2 2\tilde{\theta}_{13}\sin^2\frac{\Delta m^2_{31}L}{4E} \\ & - \sin^2\tilde{\theta}_{12}\sin^2 2\tilde{\theta}_{13}\sin^2\frac{\Delta m^2_{32}L}{4E} \; . \end{aligned} \tag{16}$$

A. Short baseline reactor experiments and θ_{13}

The forthcoming two improved short baseline reactor neutrino experiments Double Chooz and Daya Bay are planned with the same goal of searching for the smallest leptonic mixing angle θ_{13}. Both of these two experiments make use of the same concept: one near detector is placed a few hundred meters from the core of the nuclear power plant in order to reduce systematic errors and one far detector is located at distance ($L \simeq 1-2$ km) close to the first maximum of the survival probability caused by the large mass squared difference Δm^2_{31}. The $\bar{\nu}_e \to \bar{\nu}_e$ channel is dominated by the atmospheric oscillation dip, which allows us to safely neglect the term containing Δm^2_{21} in Eq. (16), and we arrive at

$$P(\bar{\nu}^s_e \to \bar{\nu}^d_e) \simeq 1 - \sin^2 2\tilde{\theta}_{13}\sin^2\frac{\Delta m^2_{31}L}{4E} \; . \tag{17}$$

Since NSIs are only sub-leading order effects, higher-order terms proportional to εs^2_{13} can be ignored, and then the effective mixing angle $\tilde{\theta}_{13}$ in Eq. (14) approximates to

$$\begin{aligned} \tilde{s}^2_{13} \;=\; & s^2_{13} + 2s_{13}\left[s_{23}\cos(\delta-\phi_{e\mu})|\varepsilon_{e\mu}| + c_{23}\cos(\delta-\phi_{e\tau})|\varepsilon_{e\tau}|\right] \\ & + s^2_{23}|\varepsilon_{e\mu}|^2 + c^2_{23}|\varepsilon_{e\tau}|^2 + 2|\varepsilon_{e\mu}||\varepsilon_{e\tau}|s_{23}c_{23}\cos(\phi_{e\mu}-\phi_{e\tau}) + \mathcal{O}(\varepsilon^3, \varepsilon s^2_{13}) \; . \end{aligned} \tag{18}$$

Note that $\tilde{s}^2_{13}$ is invariant with respect to the exchange $\varepsilon_{e\mu} \leftrightarrow \varepsilon_{e\tau}$, and obviously, in the limit $\varepsilon \to 0$, $\tilde{\theta}_{13}$ equals θ_{13}. Equation (18) clearly shows how the mixing angle θ_{13} is modified by NSIs. Some comments are in order:

- The contributions coming from the NSI parameter ε_{ee} are always correlated with higher-order corrections, and hence cannot be well constrained in a reactor experiment. However, it induces an enhancement of the total neutrino flux at a near detector, which appears as an overall factor in the oscillation probability if we do not normalize neutrino states as in Eq. (1). Due to the flux uncertainty in reactor experiments, it is very hard for this enhancement to be observed [7].

- For a given set of NSI parameters, $\sin\tilde{\theta}_{13}$ is a quadratic function of $\sin\theta_{13}$. Thus, there exists a minimum of $\sin^2\tilde{\theta}_{13}$ at the position

$$s_{13}|_{\rm min} = -s_{23}\cos(\delta-\phi_{e\mu})|\varepsilon_{e\mu}| - c_{23}\cos(\delta-\phi_{e\tau})|\varepsilon_{e\tau}| \ , \tag{19}$$

and the minimum value of $\sin^2\tilde{\theta}_{13}$ is given by

$$\tilde{s}^2_{13}|_{\rm min} = \begin{cases} \begin{array}{l} 2\Big[s^2_{23}|\varepsilon_{e\mu}|^2\sin^2\frac{\delta-\phi_{e\mu}}{2} + c^2_{23}|\varepsilon_{e\tau}|^2\sin^2\frac{\delta-\phi_{e\tau}}{2} \\ +s_{23}c_{23}|\varepsilon_{e\mu}||\varepsilon_{e\tau}|\sin(\delta-\phi_{e\mu})\sin(\delta-\phi_{e\tau})\Big] \end{array} & \text{for } s_{13}|_{\rm min} > 0 \\ \\ \begin{array}{c} s^2_{23}|\varepsilon_{e\mu}|^2 + c^2_{23}|\varepsilon_{e\tau}|^2 \\ +2|\varepsilon_{e\mu}||\varepsilon_{e\tau}|s_{23}c_{23}\cos(\phi_{e\mu}-\phi_{e\tau}) \end{array} & \text{for } s_{13}|_{\rm min} \leq 0 \end{cases} . \tag{20}$$

Since the fundamental θ_{13} cannot be well distinguished from the effective $\tilde{\theta}_{13}$ measured in an experiment, the mimicking effects of NSIs play a very important role in the small θ_{13} region. Even if the true value of θ_{13} is too tiny to be detected, we may still hope to obtain an oscillation phenomenon in reactor experiments. On the other hand, compared to θ_{13}, $\tilde{\theta}_{13}$ may also be remarkably suppressed by NSIs, which makes the current experiments quite pessimistic. Note that mimicking (or "fake") values of θ_{13} due to so-called damping effects have been investigated in Ref. [26]. Such damping effects could arise from decoherence-like damping signatures (*e.g.* wave-packet decoherence related to production and detection processes). Thus, damping could fake values of θ_{13}, and therefore the value of θ_{13} would turn out to be smaller than one expects.

- We illustrate the mappings between $\tilde{\theta}_{13}$, ε, and θ_{13} in Fig. 1. In our numerical calculations, we use the exact analytical formulas and do not make any approximations. We also adopt the central values of other relevant parameters from the global fit given in Ref. [27]. Without loss of generality, we take $|\varepsilon_{e\mu}| = |\varepsilon_{e\tau}| = |\varepsilon|$ in our analysis[3], and allow all the CP violating phases to vary from 0 to 2π. For a given value of $\tilde{\theta}_{13}$, which is in fact the parameter measured in experiments, the true values of θ_{13} may be remarkably different, *i.e.*, there exists a degeneracy in θ_{13}. Therefore, one has to

[3] Since we do not make any constraint on the CP violating phases, the numerical results would almost be the same for the case $|\varepsilon_{e\mu}| \neq |\varepsilon_{e\tau}|$.

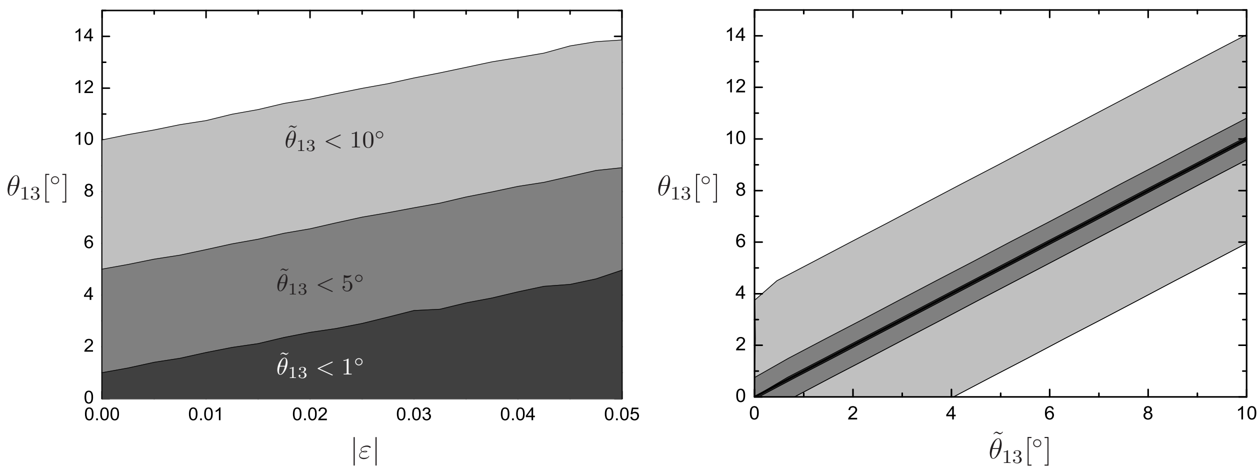

FIG. 1: Mappings between $\tilde{\theta}_{13}$, θ_{13}, and NSI parameters $\varepsilon_{\alpha\beta}$. In the left plot, we assume $0 < \tilde{\theta}_{13} < 10°$ as experimental constraints. The shaded areas correspond to different upper bounds on the effective mixing angles. For the right-hand plot, $|\varepsilon|$ is allowed to vary from 0 to 0.05, and the gray shadings represent $|\varepsilon| < 0.05$, $|\varepsilon| < 0.01$, and $|\varepsilon| < 0.001$, respectively, with darker regions for smaller $|\varepsilon|$. All CP violating phases are treated as free parameters and allowed to vary from 0 to 2π.

disentangle the parameter θ_{13} from the NSI parameters. Within the present upper bound $\tilde{\theta}_{13} < 10°$ [28], θ_{13} may approach 14° at large ε regions. In the case $\tilde{\theta}_{13} < 5°$, there is still a widely allowed range $1° < \theta_{13} < 8°$ with respect to a large ε. Even if $\tilde{\theta}_{13}$ is too small to be measured in a reactor experiment, *i.e.*, $\tilde{\theta}_{13} < 3°$, a discovery search of a non-vanishing θ_{13} may still be carried out at future neutrino factories, where the source of neutrinos is a muon storage ring with very clean muon decay and quite limited room for NSIs [29].

- In Fig. 2, we show the oscillation probabilities with respect to the NSI parameters. The upper plot in Fig. 2 indicates that mimicking oscillation effects, which are induced by sizable NSIs, can be observed in despite of a negligible θ_{13}. Once other type of neutrino oscillation experiments can help us to fix the true value of θ_{13}, the mimicking effects will provide us with the opportunity to search for NSIs in neutrino sources and detectors.

The oscillation process expressed in Eq. (17) is actually CP conserved. However, the CP violating phase δ enters the oscillation probability explicitly, and so does the leptonic mixing

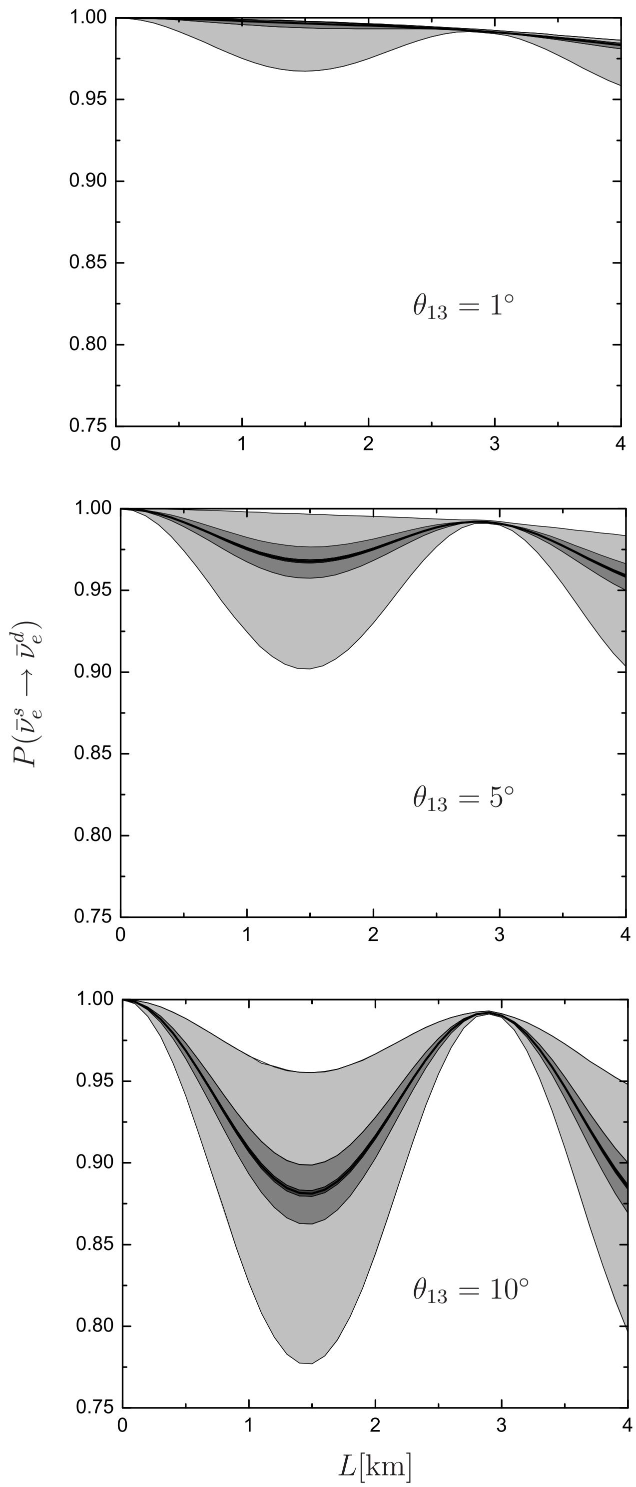

FIG. 2: NSI corrections to the oscillation probabilities $P(\bar{\nu}_e^s \to \bar{\nu}_e^d)$ in a short baseline experiment. The shadings correspond $\varepsilon < 0.05$, $\varepsilon < 0.01$, and $\varepsilon < 0.001$, respectively. The values of θ_{13} are labeled on the plots. For other mixing parameters, we use the central values given in Ref. [27]. Here we take the average energy of reactor neutrinos $E = 3$ MeV.

angle θ_{23}. It is then very helpful to extract information on leptonic CP violation and θ_{23} by analyzing the corresponding disappearance channel together with future long-baseline appearance experiments.

B. Medium baseline reactor experiments and θ_{12}

The current medium baseline reactor neutrino experiment KamLAND receives $\bar{\nu}_e$ from nuclear reactors located at an average distance $L \simeq 180$ km. In order to improve the accuracy of current measurements, the next generation experiments should take the baseline length of about 50 km, which is close to the first minimum of the survival probability related with the small mass squared difference Δm^2_{21}.

In neglecting contributions from θ_{13}, the corresponding oscillation probability reads

$$P(\bar{\nu}^s_e \to \bar{\nu}^d_e) \simeq 1 - \sin^2 2\tilde{\theta}_{12} \sin^2 \frac{\Delta m^2_{21} L}{4E} \ , \tag{21}$$

where the effective mixing angle $\tilde{\theta}_{12}$ approximates to

$$\tilde{s}^2_{12} = s^2_{12} + 2 s_{12} c_{12} \left[c_{23} \cos(\phi_{e\mu}) |\varepsilon_{e\mu}| - s_{23} \cos(\phi_{e\tau}) |\varepsilon_{e\tau}| \right] + \mathcal{O}(\varepsilon s_{13}, s^2_{13}) \ . \tag{22}$$

We now discuss how the NSIs affect the leptonic mixing angle θ_{12}:

- Similarly, there is no contribution coming from ε_{ee} at leading order. Thus, we can observe that reactor experiments are not sensitive to ε_{ee}.

- Compared to the tiny θ_{13}, the magnitude of θ_{12} is more sizable. Hence the NSI effects cannot mimic an effective $\tilde{\theta}_{12}$ with a vanishing θ_{12}. However, NSI effects may dramatically modify the observed mixing angle $\tilde{\theta}_{12}$. We plot θ_{12} as a function of $\tilde{\theta}_{12}$ and ε in Fig. 3. In the large ε regions, the true value of θ_{12} may be close to the bi-maximal mixing value 45° [30, 31, 32]. On the other hand, the lower bound $\theta_{12} > 26°$ deviates much from its tri-bimaximal mixing pattern [33, 34]. Figure 3 indicates that, in the presence of NSIs, even if $\tilde{\theta}_{12}$ can be well measured, there is still a large room of θ_{12} for various flavor symmetric models.

- The oscillation probabilities of medium baseline reactor experiments are illustrated in Fig. 4. We take the best-fit values of θ_{13} and θ_{23} in our numerical calculations [27]. Hence, the oscillation behavior around $L \simeq 0$ is mainly induced by Δm^2_{31}. It can be clearly seen that NSI corrections are more significant for a smaller θ_{12}.

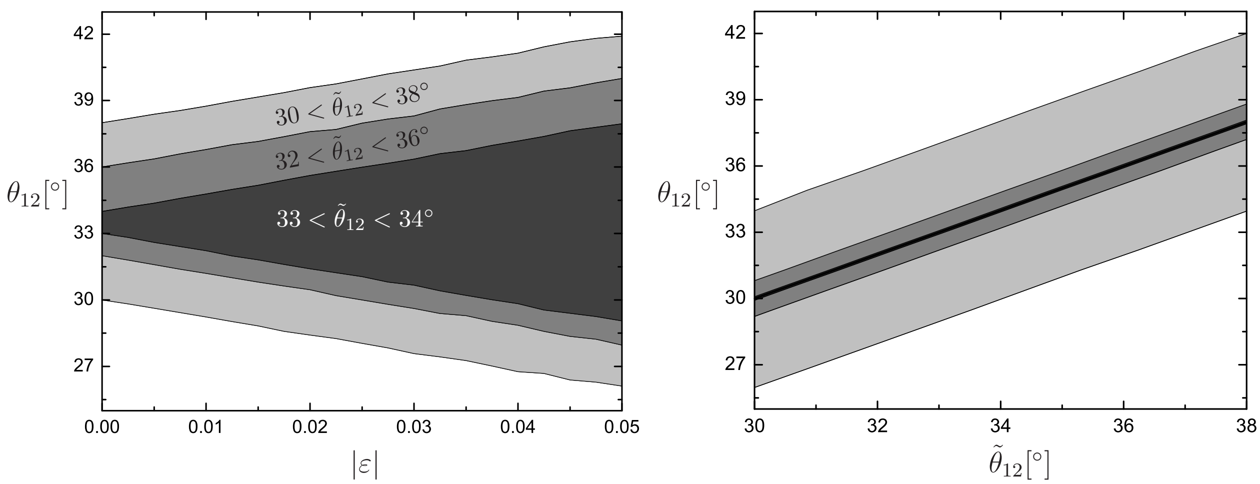

FIG. 3: Mappings between $\tilde{\theta}_{12}$, θ_{12}, and NSI parameters $\varepsilon_{\alpha\beta}$. The shaded areas in the left-hand plot correspond to different upper bounds on $\tilde{\theta}_{12}$. For the right-hand plot, $|\varepsilon|$ is allowed to vary from 0 to 0.05, and the gray shadings represent $|\varepsilon| < 0.05$, $|\varepsilon| < 0.01$, and $|\varepsilon| < 0.001$, respectively, with darker regions for smaller $|\varepsilon|$. As in Fig. 1, we allow all the CP violating phases to vary from 0 to 2π.

Unlike the mapping between $\tilde{\theta}_{13}$ and θ_{13}, only θ_{23} is entangled in Eq. (22). Hence, we may also acquire useful constraints on θ_{23} through precision measurements of θ_{12} and NSI parameters in future experiments.

C. Correlations between $\boldsymbol{\theta_{13}}$ and $\boldsymbol{\theta_{12}}$

A crucial question for future experiments is how to distinguish real mixing parameters from NSI effects. As discussed above, there is no hope to extract the fundamental mixing angles from a single reactor neutrino experiment. Since the NSI effects bring in intrinsic correlations between the effective mixing parameters, a combined analysis of both appearance and disappearance channels should be able to help us to determine NSI effects.

In Fig. 5, we show the confidence regions of θ_{13} and θ_{12}, constrained by the current global fit of neutrino oscillation data [27]. Because of the experimental uncertainties associated with $\tilde{\theta}_{13}$ and $\tilde{\theta}_{12}$, the allowed parameter spaces for θ_{13} and θ_{12} are quite wide. The true value of θ_{13} can achieve the range of the Cabibbo angle within 1σ confidence level, which shades some light on the quark-lepton complementary models [32, 35, 36]. We want to stress that

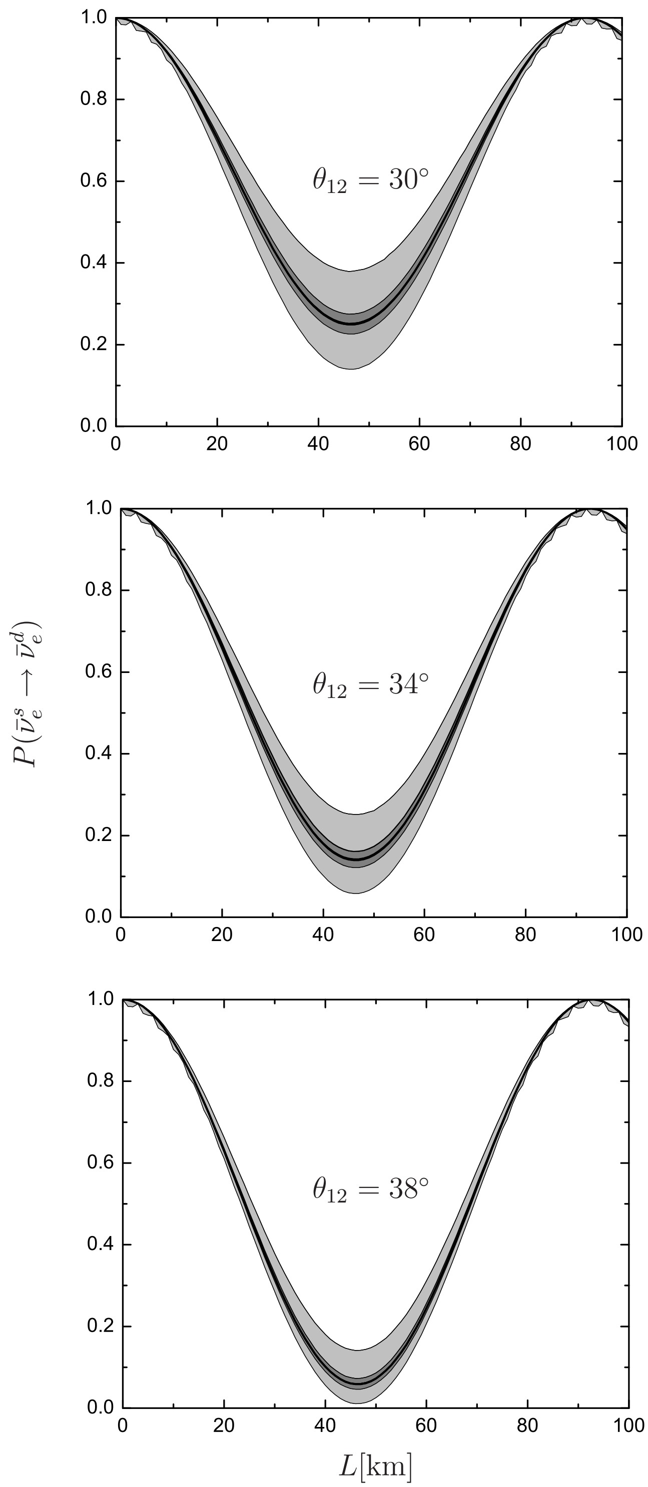

FIG. 4: NSI corrections to the oscillation probabilities $P(\bar{\nu}_e^s \to \bar{\nu}_e^d)$ in a medium baseline experiments. The shadings correspond $\varepsilon < 0.05$, $\varepsilon < 0.01$, and $\varepsilon < 0.001$, respectively. The values of $\tilde{\theta}_{12}$ are labeled on the plots. The other input parameters are the same as in Fig. 2.

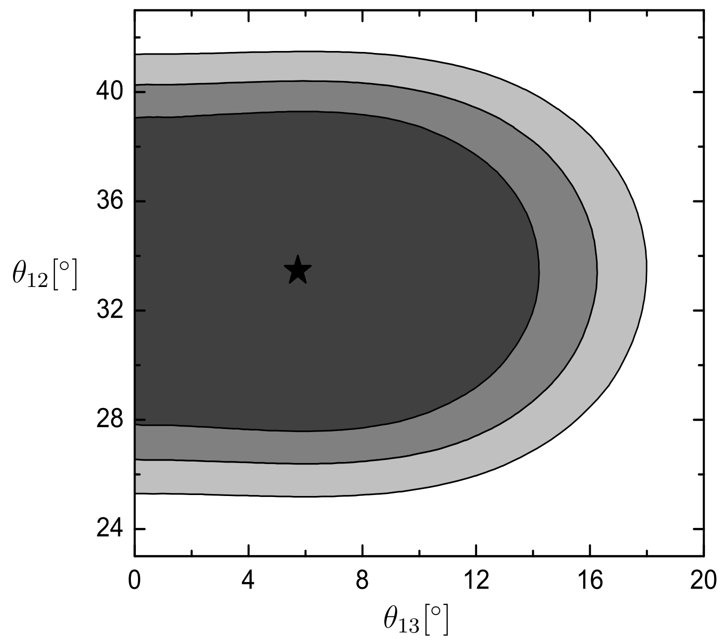

FIG. 5: The 1σ, 2σ, and 3σ confidence regions of fundamental neutrino mixing angles θ_{13} and θ_{12}, constrained by the current global fit of neutrino oscillation data. The value $|\varepsilon| = 0.05$ is assumed with all the CP violating phases being allowed to vary from 0 to 2π.

our computations depend on the input parameter ε and can just serve as a rough illustration.

D. A discussion on θ_{23} and δ

Finally, we briefly discuss the NSI corrections to θ_{23} and δ. One may define an effective mixing angle $\tilde{\theta}_{23}$ by using the analogous way that we performed above for θ_{13} and θ_{12}. However, since reactor neutrino experiments are only sensitive to the first row of the leptonic mixing matrix, the effective θ_{23} loses its meaning. In future long baseline β-beam experiments or neutrino factories, where different types of NSIs are involved in the production, propagation, and detection processes, one cannot simply employ the language of effective mixing parameters as in reactor neutrino experiments. However, the generic formulas given in Eq. (13) are still valid and very helpful for us in order to figure out NSI effects. A detailed and joint numerical analysis based on Eq. (13) should be very meaningful and will be elaborated elsewhere.

IV. SUMMARY

In this work, we have studied NSI effects in reactor neutrino experiments, and in particular, the mimicking effects on θ_{13}. We first presented the most general formulas of oscillation probabilities with all NSI effects at production, propagation, and detection processes being considered. Instead of directly discussing oscillation probabilities, we took use of a more straightforward method, which started from the effective amplitude and derived instructive mappings between fundamental mixing angles (θ_{13}, θ_{12}) and effective NSI corrected mixing angles ($\tilde{\theta}_{13}$, $\tilde{\theta}_{12}$) in reactor neutrino experiments. The analytical relations clearly show how these mixing angles are affected by NSIs. We have also illustrated the NSI effects at short and medium baseline reactor experiments. We found that the mixing angles measured in reactor neutrino experiments could be dramatically modified by NSIs at the neutrino source and detector. The mimicking effects induced by NSIs play a very important role in a short baseline experiment, especially in the case of a tiny θ_{13}. Even for a vanishing θ_{13}, the forthcoming Double Chooz and Daya Bay experiments could still perform a discovery search of an oscillation phenomenon, which should totally be attributed to NSI effects.

From the phenomenological point of view, two different and complementary oscillation experiments are needed in order to constrain corresponding NSIs. The measurement of NSI parameters should be one of the most interesting topics of experimental physics in future even before the discovery of leptonic CP violation.

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