Probing nonstandard neutrino physics by two identical detectors with different baselines

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The Kamioka-Korea two-detector system is a powerful experimental setup for resolving neutrino parameter degeneracies and probing *CP* violation in neutrino oscillation. In this paper, we study sensitivities of the same setup to several nonstandard neutrino physics such as quantum decoherence, tiny violation of Lorentz symmetry, and nonstandard interactions of neutrinos with matter. We show that it can achieve significant improvement on the current bounds on nonstandard neutrino physics. In most cases, the Kamioka-Korea two-detector setup is more sensitive than the one-detector setup, either in Kamioka or in Korea, except for the cases when Lorentz symmetry is broken in a CPT-violating manner and the nonstandard neutrino interactions with matter is present.

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I. INTRODUCTION

A variety of the neutrino experiments, atmospheric [1], solar [2], reactor [3], and accelerator [4], have been successful in identifying the oscillation of neutrinos induced by their masses as a dominant mechanism for neutrino disappearances. After passing through the discovery era, the neutrino physics will enter the epoch of precision study, as the Cabibbo-Kobayashi-Maskawa phenomenology and a detailed study of *CP* violation have blossomed in the quark sector. The Maki-Nakagawa-Sakata [5] lepton flavor mixing matrix elements, including the *CP* phase(s), will be measured with high precision, and the properties of neutrinos such as the absolute values of their masses and their interactions with matter, etc., will be studied with greater accuracies. During the course of precision studies, it will become natural to investigate nonstandard physics which could be possessed by neutrinos, including flavor changing neutral/charged current interactions [6–8], the effects of quantum decoherence [9–11] that may possibly arise due to quantum gravity at short distance scale [9], and violation of Lorentz and CPT invariance [12–14], to name a few.

It is well known in the history of physics experiments based on interference effects played very crucial roles in advancing our understanding of the physical laws. The well-known examples include the famous two-slit experiment by Young, the Michelson-Morley experiment which demonstrated that there is no ether, the Davidson-Germer experiment on electron diffraction, and $K^0 - \bar{K}^0$ oscillation, etc. Likewise, neutrino oscillation experiments may probe another important structure of fundamental physics by observing tiny effects due to CPT violation or quantum decoherence, which may be rooted in quantum gravity. Along with the neutral meson systems ($K^0 - \bar{K}^0$ and $B^0 - \bar{B}^0$), neutrino oscillations could provide competing and/or complementary informations on those exotic effects. See Ref. [15] for a recent review on this subject.

In a previous work [16,17] we introduced and explored in detail the physics potential of the Kamioka-Korea two-detector setting which receive an intense neutrino beam from J-PARC. We demonstrated that the setting is powerful enough to resolve all the eight-fold parameter degeneracy [18–20], if θ_{13} is in reach of the next generation accelerator [21,22] and the reactor experiments [23,24]. The degeneracy includes the parameters θ_{13} and δ, and it is doubled by the ambiguities which arise due to the unknown sign of Δm^2_{31} and octant of θ_{23}. The detector in Korea plays a decisive role to lift the last two degeneracies. For related works on the Kamioka-Korea two-detector complex, see, for example, [25–28].

The Kamioka-Korea identical two-detector setting is a unique apparatus for studying nonstandard physics (NSP). As will be elaborated in Sec. II, the deviation from the

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expectation by the standard mass-induced neutrino oscillation can be probed with high sensitivity by comparing yields at the intermediate (Kamioka) and the far (Korea) detectors. In this paper, we aim at exploring the physics potential of the Kamioka-Korea setting in a systematic way. By this we mean that we examine several NSP effects in a single framework by concentrating on a $\nu_\mu - \nu_\tau$ subsystem in the standard three-flavor mixing scheme. We focus on the ν_μ disappearance measurement in our analysis. While we will work within the limited framework, it will allow us to treat the problem in a coherent fashion.

In analyzing nonstandard physics in this paper, we demonstrate the powerfulness of the Kamioka-Korea identical two-detector setting. For this purpose, we systematically compare the results obtained with the following three settings:

(i) Kamioka-Korea setting: Two identical detectors one at Kamioka and the other in Korea each 0.27 Mton.
(ii) Kamioka-only setting: A single 0.54 Mton detector at Kamioka.
(iii) Korea-only setting: A single 0.54 Mton detector somewhere in Korea.

Among the cases we have examined, the Kamioka-Korea setting always gives the best sensitivities, apart from two exceptions of violation of Lorentz invariance in a CPT-violating manner and the nonstandard neutrino interactions with matter. Whereas the next best case is sometimes Kamioka-only or Korea-only settings depending upon the problem.

This paper is organized as follows. In Sec. II, we illustrate how we can probe nonstandard physics with the Kamioka-Korea two-detector setting, with a quantum decoherence as an example of nonstandard neutrino physics. In Sec. III, we describe the statistical method which is used in our analyses in the following sections. In Sec. IV, we discuss quantum decoherence. In Sec. V, we discuss possible violation of Lorentz invariance. In Sec. VI, we discuss nonstandard neutrino matter interactions, and the results of the study are summarized in Sec. VII.

II. PROBING NONSTANDARD PHYSICS WITH THE KAMIOKA-KOREA TWO-DETECTOR SETTING

In this section, we describe how we proceed in the following sections. For the purpose of illustration, we consider quantum decoherence (QD) as nonstandard neutrino physics. In this case, the ν_μ survival probability (and the $\bar{\nu}_\mu$ survival probability assuming CPT invariance in the presence of QD) is given by [10,11],

$$P(\nu_\mu \to \nu_\mu) = P(\bar{\nu}_\mu \to \bar{\nu}_\mu) = 1 - \frac{1}{2}\sin^2 2\theta\left[1 - e^{-\gamma(E)L}\cos\left(\frac{\Delta m^2 L}{2E}\right)\right], \quad (1)$$

where θ is the mixing angle, δm^2 is the mass squared difference, E is the neutrino energy, and L is the baseline, and we consider the case $\gamma(E) = \gamma/E$ (see Sec. IV where we consider also the cases of $\gamma(E)$ which has other energy dependences). Then, one can calculate the number of ν_μ and $\bar{\nu}_\mu$ events observed at two detectors placed at Kamioka (with a baseline of 295 km) and Korea (with an assumed baseline of 1050 km), using the above survival probability and the neutrino beam profiles.

For definiteness, let us consider the number of observed neutrino events both at Kamioka and Korea, for each energy bin (with 50 MeV width) from $E_\nu = 0.2$ GeV up to $E_\nu = 1.2$ GeV. In Fig. 1, we show the ν_μ event spectra at detectors located at Kamioka and Korea for the pure oscillation $\gamma = 0$ (left column) and the oscillation plus QD with two different QD parameters, $\gamma = 1 \times 10^{-4}$ GeV/km (middle column) and $\gamma = 2 \times 10^{-4}$ GeV/km (right column).[1] One finds that the event energy spectra are modulated for nonvanishing γ. Most importantly, the spectral changes are different between detectors in Kamioka and Korea due to the different L/E values at the two locations. We observe that the decoherence effect is visible for the range of E/L larger than $\sim$0.6 GeV/1000 km. For E/L smaller than this value, the oscillation frequency (as a function of neutrino energy) is large and what one can observe is the averaged effect due to the finite energy resolution of the experiment. In this case, we do not expect to see clear differences between the pure oscillation and the oscillation-decoherence coexisting cases.

Assuming the actual data at detectors in Kamioka and in Korea are given (or well described) by the pure oscillation with $\sin^2 2\theta = 1$ and $\Delta m^2 = 2.5 \times 10^{-2}$ eV^2, we could claim that $\gamma = 1 \times 10^{-4}$ GeV/km (shown in the middle column), for example, would be inconsistent with the data. One can make such a claim in a more proper and quantitative manner using the χ^2 analysis, as described in details in the following sections.

III. ANALYSIS METHOD

In studying nonstandard physics through neutrino oscillations we restrict ourselves into the $\nu_\mu - \nu_\tau$ subsystem, rather than dealing with the full three generation problem. The reason for truncation is partly technical and partly physics motivated; it is to simplify analysis in a manner not to spoil the most important features of the problem. First of all, $\nu_\mu \to \nu_\mu$ and $\bar{\nu}_\mu \to \bar{\nu}_\mu$ probabilities do not depend much on the yet unknown neutrino mixing parameters, namely, θ_{13}, CP phase, and the sign of Δm^2_{31}. Therefore, as long as we restrict our analysis to the particular parameters of new physics which have a strong impact in this channel, it is likely that we obtain approxi-

[1] In order to convert this γ in units of GeV/km to γ defined in Eq. (5), one has to multiply 0.197×10^{-18}.

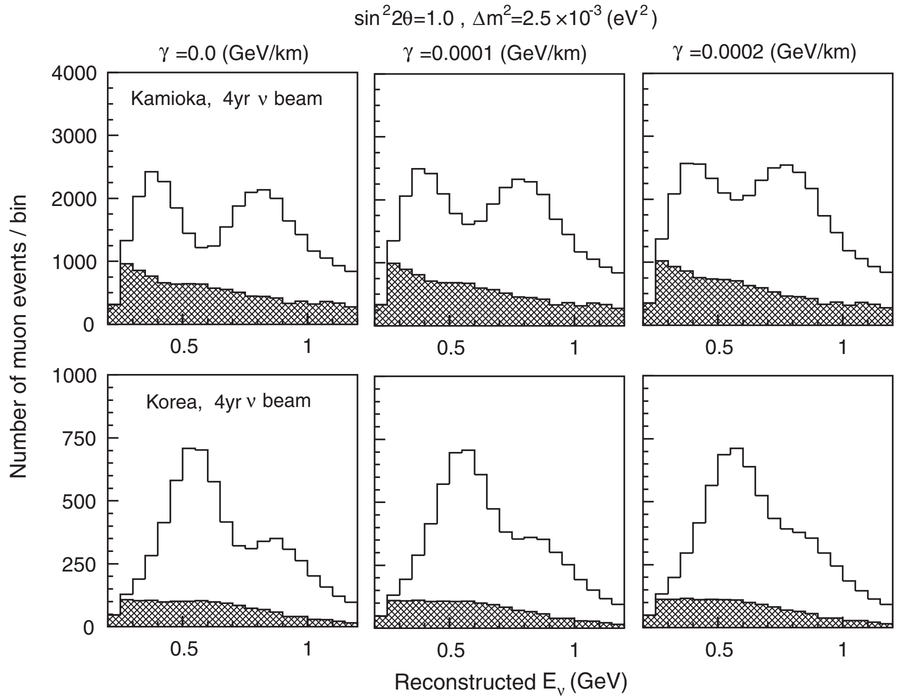

FIG. 1. Event spectra of neutrinos at Kamioka (top panel) and Korea (bottom panel) for $\gamma = 0$ (left column), 1×10^{-4} GeV/km (middle column), and $\gamma = 2 \times 10^{-4}$ GeV/km (right column). The hatched areas denote the contributions from non-quasielastic events.

mately correct bounds despite the truncation. Ignoring the appearance channel is also expected not to produce important change in the results because the statistics is much less than that in the disappearance channel. Second, in most cases the earth matter effect is a subleading effect in the $\nu_\mu - \nu_\tau$ subsystem so that by restricting to the subsystem we need not to worry about complications due to the matter effect. Moreover, many of the foregoing analyses were carried out under the truncated framework. By working under the same approximation the comparison between ours and the existing results becomes much simpler. Thus, the Δm^2 which will appear in neutrino oscillation probabilities in the following sections is meant to be $\Delta m^2 \equiv \Delta m^2_{32} = m^2_3 - m^2_2$. θ will be a mixing angle in the unitary matrix diagonalizing the Hamiltonian in the $\nu_\mu - \nu_\tau$ subsystem, which is essentially equal to θ_{23}.

A. Method of statistical analysis

In order to estimate the sensitivity of the experiment with the two-detector system at 295 km (Kamioka) and 1050 km (Korea), we carry out a χ^2 analysis. In the present analysis, we only included ν_μ and $\bar{\nu}_\mu$ disappearance channels. In short, the definition of the statistical procedure is similar to the one used in [17] excluding the electron events. The assumption on the experimental setting is also identical to that of the best performance setting identified in Ref. [16]. Namely, 0.27 Mton fiducial masses for the intermediate site (Kamioka, 295 km) and the far site (Korea, 1050 km). For the reference, we also consider 0.54 Mton detector for Kamioka or Korea only. The neutrino beam is assumed to be a 2.5° off-axis one produced by the upgraded J-PARC 4 MW proton beam. It is assumed that the experiment will continue for 8 years with 4 years of neutrino and 4 years of antineutrino runs.

We use various numbers and distributions available from references related to T2K [29], in which many of the numbers are updated after the original proposal [21]. Here, we summarize the main assumptions and the methods used in the χ^2 analysis. We use the reconstructed neutrino energy for single-Cherenkov-ring muon events. The resolution in the reconstructed neutrino energy is 80 MeV for quasielastic events. We take $\Delta m^2 = 2.5 \times 10^{-3}$ eV^2 and $\sin^2 2\theta_{23} = 1.0$ for our reference value. However, whenever we expect that there is a correlation between the expected sensitivity to new physics and the oscillation parameters, we scan Δm^2_{31} between 2.0 and 3.0×10^{-3} eV^2 and $\sin^2 2\theta_{23}$ between 0.9 and 1.0. The shape of the energy spectrum for the antineutrino beam is assumed to be identical to that of the neutrino beam. The event rate for the antineutrino beam in the absence of neutrino oscillations is smaller by a factor of 3.4 due mostly to the lower neutrino interaction cross sections and partly to the slightly lower flux. In addition, the contamination of the wrong sign muon events is higher in the antineutrino beam.

We stress that in the present setting the detectors placed in Kamioka and in Korea are not only identical but also receive neutrino beams with essentially the same energy distribution (due to the same off-axis angle of 2.5°) in the absence of oscillations. However, it was realized recently that, due to a noncircular shape of the decay pipe of the J-PARC neutrino beam line, the flux energy spectra viewed at detectors in Kamioka and in Korea are expected to be slightly different even at the same off-axis angle, especially in the high-energy tail of the spectrum [30]. The possible difference between fluxes in the intermediate and the far detectors is taken into account as a systematic error in the analysis in [17]. We follow the same prescription in the present analysis.

We compute neutrino oscillation probabilities by numerically integrating a neutrino evolution equation under the constant density approximation. The average density is assumed to be 2.3 and 2.8 $\mathrm{g/cm^3}$ for the matter along the beam line between the production target and Kamioka and between the target and Korea, respectively [16]. We assume that the number of electrons with respect to that of nucleons to be 0.5 to convert the matter density to the electron number density.

The statistical significance of the measurement considered in this paper was estimated by using the following definition of χ^2:

$$\chi^2 = \sum_{k=1}^{4}\left(\sum_{i=1}^{20}\frac{(N(\mu)_i^{\mathrm{obs}} - N(\mu)_i^{\mathrm{exp}})^2}{\sigma_i^2}\right) + \sum_{j=1}^{4}\left(\frac{\epsilon_j}{\tilde{\sigma}_j}\right)^2, \quad (2)$$

where

$$N(\mu)_i^{\mathrm{exp}} = N_i^{\mathrm{non-QE}}\cdot\left(1 + \sum_{j=1,3,4} f(\mu)_j^i\cdot\epsilon_j\right) + N_i^{\mathrm{QE}}\cdot\left(1 + \sum_{j=1,2,4} f(\mu)_j^i\cdot\epsilon_j\right). \quad (3)$$

In Eq. (2), $N(\mu)_i^{\mathrm{obs}}$ is the number of single-ring muon events to be observed for the given (oscillation) parameter set, and $N(\mu)_i^{\mathrm{exp}}$ is the expected number of events for the assumed parameters in the χ^2 analysis. $k = 1$, 2, 3 and 4 correspond to the four combinations of the detectors in Kamioka and in Korea with the neutrino and antineutrino beams, respectively. The index i represents the reconstructed neutrino energy bin for muons. The energy range for the muon events covers from 200 to 1200 MeV. Each energy bin has 50 MeV width. σ_i denotes the statistical uncertainties in the expected data. The second term in the χ^2 definition collects the contributions from variables which parameterize the systematic uncertainties in the expected number of signal and background events.

$N_i^{\mathrm{non-QE}}$ are the number of non-quasielastic muon events for the ith bin whereas N_i^{QE} are the number of quasielastic muon events. We treat the non-quasielastic and quasielastic muon events separately, since the neutrino energy cannot be properly reconstructed for non-quasielastic events. To reproduce the Monte Carlo simulation results by the T2K Collaboration, we assumed that the reconstructed neutrino energies for non-quasielastic events are shifted to lower energies by 300 ± 160 MeV independent of the neutrino energy, where 160 MeV is the 1σ width of the reconstructed neutrino energy distribution relative to the true neutrino energy. Whereas for quasielastic events the reconstructed neutrino energies agree with the true neutrino energy with 80 MeV energy resolution. See the right panel of Fig. 3.4 of Ref. [31] where the energy resolution is shown for quasielastic events (solid histogram) and non-quasielastic (hatched histogram). Hence both N_i^{QE} and $N_i^{\mathrm{non-QE}}$ depend on neutrino (oscillation) parameters but in a different way, namely, the former (latter) being affected in a direct (indirect) manner and hence dependence is strong (weak) as we can see from the solid histogram ($N_i^{\mathrm{QE}} + N_i^{\mathrm{non-QE}}$) and the hatched region ($N_i^{\mathrm{non-QE}}$) in Fig. 1. The key to high sensitivity to NSP is the different oscillation parameter dependence between the Kamioka and the Korean detectors due to different baselines. The uncertainties in $N_i^{\mathrm{non-QE}}$ and N_i^{QE} are represented by four parameters ϵ_j ($j = 1$ to 4).

During the fit, the values of $N(\mu)_i^{\mathrm{exp}}$ are recalculated for each choice of the (oscillation) parameters which are varied freely to minimize χ^2, and so are the systematic error

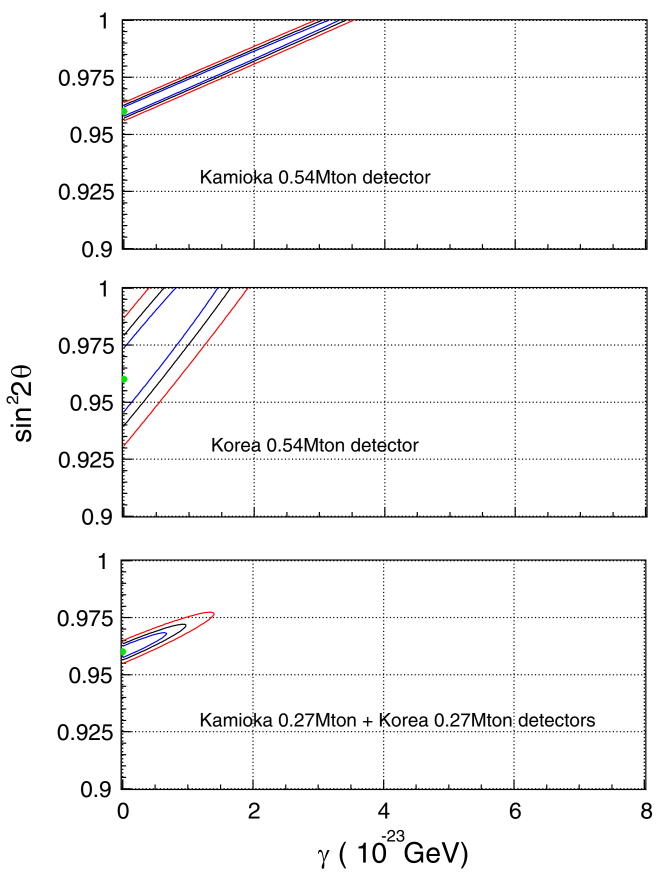

FIG. 2 (color online). The correlations between γ and $\sin^2 2\theta$ for the three experimental setups we consider: Kamioka only, Korea only, and Kamioka-Korea. Inner (blue), middle (black), and outer (red) curves represent the contours for 68%, 90% and 99% C.L., respectively, for 2 degrees of freedom. Input values are $\gamma = 0$, $\Delta m^2 = 2.5\times10^{-3}\ \mathrm{eV}^2$, and $\sin^2 2\theta = 0.96$.

parameters ϵ_j. The parameter $f(\mu)^i_j$ represents the fractional change in the predicted event rate in the ith bin due to a variation of the parameter ϵ_j. We assume that the experiment is equipped with a near detector which measures the unoscillated muon neutrino spectrum. The uncertainties in the absolute normalization of events are assumed to be 5% ($\tilde{\sigma}_1 = 0.05$). The functional form of $f(\mu)^i_2 = (E_\nu(\text{rec}) - 800\ \text{MeV})/800\ \text{MeV}$ is used to define the uncertainty in the spectrum shape for quasielastic muon events ($\tilde{\sigma}_2 = 0.05$) [17].

The uncertainty in the separation of quasielastic and non-quasielastic interactions in the muon events is assumed to be 20% ($\tilde{\sigma}_3 = 0.20$). In addition, for the number of events in Korea, the possible flux difference between Kamioka and Korea is taken into account in $f(\mu)^i_4$. The predicted flux difference [30] is simply assumed to be the 1σ uncertainty in the flux difference ($\tilde{\sigma}_4$).

Finally, in this work, the sensitivity at 90% (99%) confidence level (C.L.) is defined by

$$\Delta\chi^2 \equiv \chi^2_{\min}(\text{osc.} + \text{ nonstandard physics}) - \chi^2_{\min}(\text{osc.}) \geq 2.71(6.63),$$

corresponding to the 1 degree of freedom, except for the results shown in Figs. 2 and 6. Similarly, the criterion for the 2 degrees of freedom is $\Delta\chi^2 \geq 4.61(9.21)$ which is used to obtain allowed regions shown in Figs. 2 and 6.

IV. QUANTUM DECOHERENCE

The study of quantum decoherence is based on a hypothesis that somehow there may be a loss of coherence due to environmental effect or quantum gravity and spacetime foam, etc. [9]. We do not discuss the origin of decoherence in this paper, but concentrate on how this effect can be probed by the Kamioka-Korea setting. The same statement also applies to other nonstandard physics considered in this paper. For previous analyses of decoherence in neutrino experiments, see e.g., [10,32–35].

As discussed in Sec. I we consider the $\nu_\mu - \nu_\tau$ two-flavor system. Since the matter effect is a subleading effect in this channel we employ vacuum oscillation approximation in this section. The ν_μ and $\bar{\nu}_\mu$ survival probability we consider is already shown in Eq. (1) containing the parameter $\gamma(E) > 0$ which controls the strength of the decoherence effect. Notice that the conventional two-flavor oscillation formula is reproduced in the limit $\gamma(E) \to 0$. Notice that we assume CPT invariance in this section so that $\gamma(E)$ for $\bar{\nu}_\mu$ is the same as for ν_μ and as the survival probabilities. Since the total probability is still conserved in the presence of QD, the relation $P(\nu_\mu \to \nu_\tau) = 1 - P(\nu_\mu \to \nu_\mu)$ holds.

Unfortunately, nothing is known about the energy dependence of $\gamma(E)$. Therefore, we examine, following [10], several typical cases of energy dependence of $\gamma(E)$:

$$\gamma(E) = \gamma\left(\frac{E}{\text{GeV}}\right)^n (\text{with } n = 0, 2, -1). \quad (4)$$

In this convention, the overall constant γ has a dimension of energy or $(\text{length})^{-1}$, irrespective of the values of the exponent n. We will use γ in GeV units in this section. In the following three subsections, we analyze three different energy dependences, $n = 0, -1, 2$ one by one.

A. Case of $\frac{1}{E}$ dependence of $\gamma(E)$

First, we examine the case with $\frac{1}{E}$ dependence of $\gamma(E)$. In Fig. 2, we show the correlations between γ and $\sin^2 2\theta$ at three experimental setups for the case where the input values are $\gamma = 0$, $\Delta m^2 = 2.5 \times 10^{-3}\ \text{eV}^2$, and $\sin^2 2\theta = 0.96$. We immediately find that there are strong correlations between $\sin^2 2\theta$ and γ for the Kamioka-only and the Korea-only setups. We also note that the slope of the correlation for the Kamioka-only setup is different from that for the Korea-only setup. Therefore, the Kamioka-Korea setup can give a stronger bound than each single-detector setup. This advantage can be seen in Fig. 3, where we present the sensitivity regions of γ as a function of $\sin^2 2\theta$ (left panel) and Δm^2 (right panel), which is the allowed range of γ whose input value is zero (or the maximum allowed value of γ which is consistent with the case of vanishing γ). Note that the sensitivity to γ in the Kamioka-Korea setting is better than the Korea-only and the Kamioka-only settings by a factor of about 3 and 6, respectively.

B. Case of energy independent $\gamma(E)$

Next, we examine the case of energy independent γ, $\gamma(E) = \gamma = \text{constant}$. Repeating the same procedure as before, we find that the sensitivity to γ in the Kamioka-Korea setting is better than the Korea-only and the Kamioka-only settings by a factor of about 3 and 8, respectively, as summarized in Table I.

C. Case of E^2 dependence of $\gamma(E)$

Finally, we examine the case with the E^2 dependence of γ. The qualitative features of the sensitivities are similar to those of the previous two cases. The result is that the sensitivity to γ in the Kamioka-Korea setting is better than the Korea-only and the Kamioka-only settings by a factor greater than 3 and 5, respectively.

D. Comparison between sensitivities of Kamioka-Korea setting and the existing bound on γ

In Table I, we list, for the purpose of comparison, the upper bounds on γ at 90% C.L. obtained by analyzing the atmospheric neutrino data in [10].[2] We summarize in the

[2] We do not quote the bounds on γ obtained from solar and KamLAND neutrinos, since they are derived from $\nu_e \to \nu_e$ [36]. In this case, the neutrino energy is quite low, so that the constraint for $n = -1$ becomes quite strong.

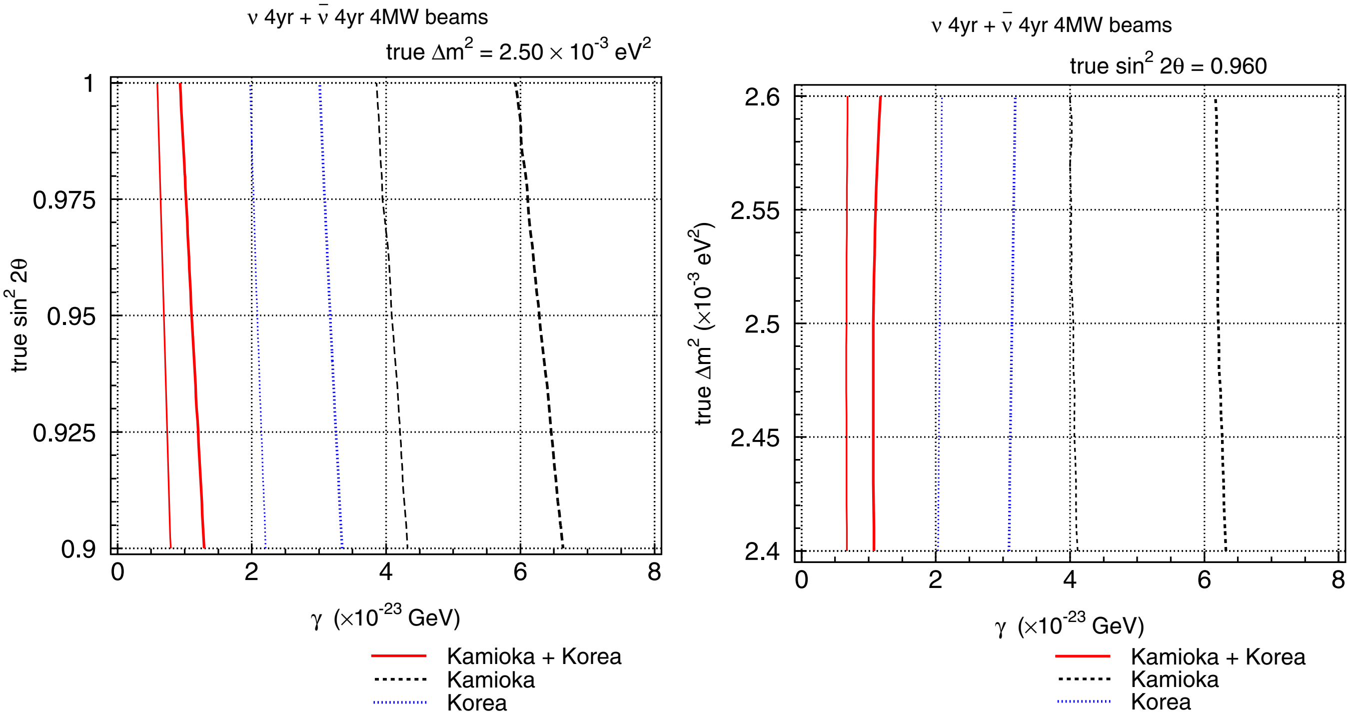

FIG. 3 (color online). The sensitivity to γ as a function of the true (input) value of $\sin^2 2\theta \equiv \sin^2 2\theta_{23}$ (left panel) and $\Delta m^2 \equiv \Delta m^2_{32}$ (right panel) for the case of $1/E$ dependence of $\gamma(E)$ obtained by fitting the input data corresponding to the vanishing $\gamma(E)$ with the nonzero $\gamma(E)$. The red solid lines are for Kamioka-Korea setting with each 0.27 Mton detector, while the dashed black (dotted blue) lines are for Kamioka (Korea) only setting with 0.54 Mton detector. The thick and the thin lines are for 99% and 90% C.L., respectively. Four years of neutrino plus 4 years of antineutrino running are assumed. In obtaining the results shown in the left and right panels, the input value of Δm^2_{32} is taken as $+2.5 \times 10^{-3}$ eV^2 (with positive sign indicating the normal mass hierarchy) and that of $\sin^2 2\theta_{23}$ as 0.96, respectively.

table the bounds on γ at 90% C.L. achievable by the Kamioka-only, the Korea-only, and the Kamioka-Korea settings. We use the same ansatz as in [10] for parameterizing the energy dependence of $\gamma(E)$, $n = 0, -1$ and 2.

In the case of $\frac{1}{E}$ dependence of $\gamma(E)$, all three settings can improve the current bound almost by 2 orders of magnitude. Note that the best case (Kamioka-Korea) is a factor of 6 better than the Kamioka-only case. This case demonstrates clearly that the two-detector setup is more powerful than the Kamioka-only setup. We notice that in the case of energy independence of $\gamma(E)$, the Kamioka-Korea two-detector setting can improve the current bound by a factor of $\sim$3.

In the case of E^2 dependence of $\gamma(E)$ the situation is completely reversed; the bound imposed by the atmospheric neutrino data surpasses those of our three settings by almost $\sim$4 orders of magnitude. Because the spectrum of atmospheric neutrinos spans a wide range of energy which extends to 100–1000 GeV, it gives much tighter constraints on the decoherence parameter for quadratic energy dependence of $\gamma(E)$. In a sense, the current Super-Kamiokande experiment is already a powerful neutrino spectroscope with a very wide energy range, and could be sensitive to nonstandard neutrino physics that may affect higher energy neutrinos such as QD with $\gamma(E) \sim E^2$ or Lorentz symmetry violation (see Sec. V for more details).

V. VIOLATION OF LORENTZ INVARIANCE

In the presence of Lorentz symmetry violation by a tiny amount, neutrinos can have both velocity mixings and

TABLE I. Presented are the upper bounds on decoherence parameters γ defined in (4) for three possible values of n. The current bounds are based on [10] and are at 90% C.L.. The sensitivities obtained by this study are also at 90% C.L. and correspond to the true values of the parameters $\Delta m^2 = 2.5 \times 10^{-3}$ eV^2 and $\sin^2 2\theta_{23} = 0.96$.

Ansats for $\gamma(E)$	Current bound (GeV)	Kamioka only (GeV)	Korea only (GeV)	Kamioka-Korea (GeV)
$\gamma(E) = \gamma(\text{const})$	$<3.5 \times 10^{-23}$	$<8.7 \times 10^{-23}$	$<3.2 \times 10^{-23}$	$<1.1 \times 10^{-23}$
$\gamma(E) = \gamma/(E/\text{GeV})$	$<2.0 \times 10^{-21}$	$<4.0 \times 10^{-23}$	$<2.0 \times 10^{-23}$	$<0.7 \times 10^{-23}$
$\gamma(E) = \gamma(E/\text{GeV})^2$	$<0.9 \times 10^{-27}$	$<9.2 \times 10^{-23}$	$<6.0 \times 10^{-23}$	$<1.7 \times 10^{-23}$

mass mixings, both are CPT conserving [12]. Also there could be CPT-violating interactions in general [12–14]. Then, the energy of neutrinos with definite momentum in an ultrarelativistic regime can be written as

$$\frac{mm^\dagger}{2p} = cp + \frac{m^2}{2p} + b, \tag{5}$$

where m^2, c, and b are 3×3 Hermitian matrices, and the three terms represent, in order, the effects of velocity mixing, mass mixing, and CPT violation [13]. The energies of neutrinos are eigenvalues of (5), and the eigenvectors give the "mass eigenstates." Notice that while c is dimensionless quantity, b has dimension of energy. For brevity, we will use the GeV unit for b.

Within the framework just defined above it was shown by Coleman and Glashow [13] that the ν_μ disappearance probability in vacuum can be written as

$$P(\nu_\mu \to \nu_\mu) = 1 - \sin^2 2\Theta \sin^2(\Delta L/4), \tag{6}$$

where the "mixing angle" Θ and the phase factor Δ depend upon seven parameters apart from energy E and Δm^2:

$$\begin{aligned} \Delta \sin 2\Theta &= \Delta m^2 \sin 2\theta_m / E + 2\delta b e^{i\eta} \sin 2\theta_b \\ &\quad + 2\delta c e^{i\eta'} E \sin 2\theta_c, \\ \Delta \cos 2\Theta &= \Delta m^2 \cos 2\theta_m / E + 2\delta b \cos 2\theta_b \\ &\quad + 2\delta c E \cos 2\theta_c. \end{aligned} \tag{7}$$

As easily guessed, what is relevant in neutrino oscillation is the difference in mass squared and b and c between two mass eigenstates, $\delta b \equiv b_2 - b_1$ and $\delta c \equiv c_2 - c_1$, where $c_{i=1,2}$ and $b_{i=1,2}$ are the eigenvalues of the matrix c and b. The angles θ_m, θ_b, and θ_c appear in the unitary matrices which diagonalize the matrices m^2, b, and c, respectively. There are also two phases η and η' that cannot be rotated away by field redefinition. We work in the convention in which $\cos 2\theta_m$ and $\cos 2\theta_b$ are positive, and $\Delta m^2 \equiv m_2^2 - m_1^2$, $\delta b \equiv b_2 - b_1$, and $\delta c \equiv c_2 - c_1$ can have either signs.

The survival probability for the antineutrino is obtained by the following substitution:

$$\delta c \to \delta c, \qquad \delta b \to -\delta b. \tag{8}$$

The difference in the sign changes signify the CPT conserving vs CPT-violating nature of c and b terms.

The two-flavor oscillation given by (6) and (7) is too complicated for full analysis. Therefore, we make some simplifications in our analysis. We restrict ourselves into the case $\theta_m = \theta_b = \theta_c \equiv \theta$ and $\eta = \eta' = 0$, for which one recovers the case treated in [37]:

$$P(\nu_\mu \to \nu_\mu) = 1 - \sin^2 2\theta \sin^2\left[L\left(\frac{\Delta m^2}{4E} + \frac{\delta b}{2} + \frac{\delta c E}{2}\right)\right], \tag{9}$$

which still depends on four parameters, θ, Δm^2, δb, and δc. One has a similar expression for $\bar{\nu}_\mu$ with $\delta b \to -\delta b$. As pointed out in [38], the analysis for violation of Lorentz invariance with the δc term is equivalent to testing the equivalence principle [39]. The oscillation probability in (9) looks like the one for conventional neutrino oscillations due to Δm^2, with small corrections due to the Lorentz symmetry violating δb and δc terms. In this sense, it may be the most interesting case to examine as a typical example with the Lorentz symmetry violation. Note that the sign of δb and δc can have different effects on the survival probabilities, so that the bounds on δb and δc could depend on their signs, although we will find that the difference is rather small.

For ease of analysis and simplicity of presentation, we further restrict our analysis to the case of either $\delta b = 0$ and $\delta c \neq 0$, or $\delta b \neq 0$ and $\delta c = 0$. Notice that the former is CPT conserving while the latter is CPT violating.

A. Case with $\delta b = 0$ and $\delta c \neq 0$ (CPT conserving)

We first examine violation of Lorentz invariance in the case of $\delta b = 0$ and $\delta c \neq 0$. In Fig. 4 we present the region of allowed values of δc as a function of $\sin^2 2\theta$ (left panel) and Δm^2 (right panel). The sensitivities to δc achieved by the Kamioka-Korea setting is better than those of the Korea-only and the Kamioka-only settings, but only slightly unlike the case of quantum decoherence. The sensitivity is weakly correlated to θ, and the best sensitivity is achieved at the maximal θ. There is almost no correlation to Δm^2.

B. Case with $\delta c = 0$ and $\delta b \neq 0$ (CPT violating)

The allowed regions with violation of Lorentz invariance in the case of $\delta c = 0$ and $\delta b \neq 0$ presented in Fig. 5 have several unique features. First of all, unlike the system with decoherence, the sensitivity is greatest in the Kamioka-only setting, though the one by the Kamioka-Korea setting is only slightly less by about 15%–20%. Whereas the sensitivity by the Korea-only setting is much worse, more than a factor of 2 compared to the Kamioka-only setting. The reason for this lies in the structure of the ν_μ and $\bar{\nu}_\mu$ survival probabilities. In this scenario, the effect of the nonvanishing δb appears as the difference in the oscillation frequency between neutrinos and antineutrinos, if the energy dependence is neglected. In this case, the measurement at different baselines is not very important. Then the Kamioka-only setup turns out to be slightly better than the Kamioka-Korea setup. This case is also unique by having the worst sensitivity at the largest value of Δm^2 (right panel). Also, the correlation of sensitivity to $\sin^2 2\theta$ (left panel) is strongest among the cases examined in this paper, with maximal sensitivity at maximal θ.

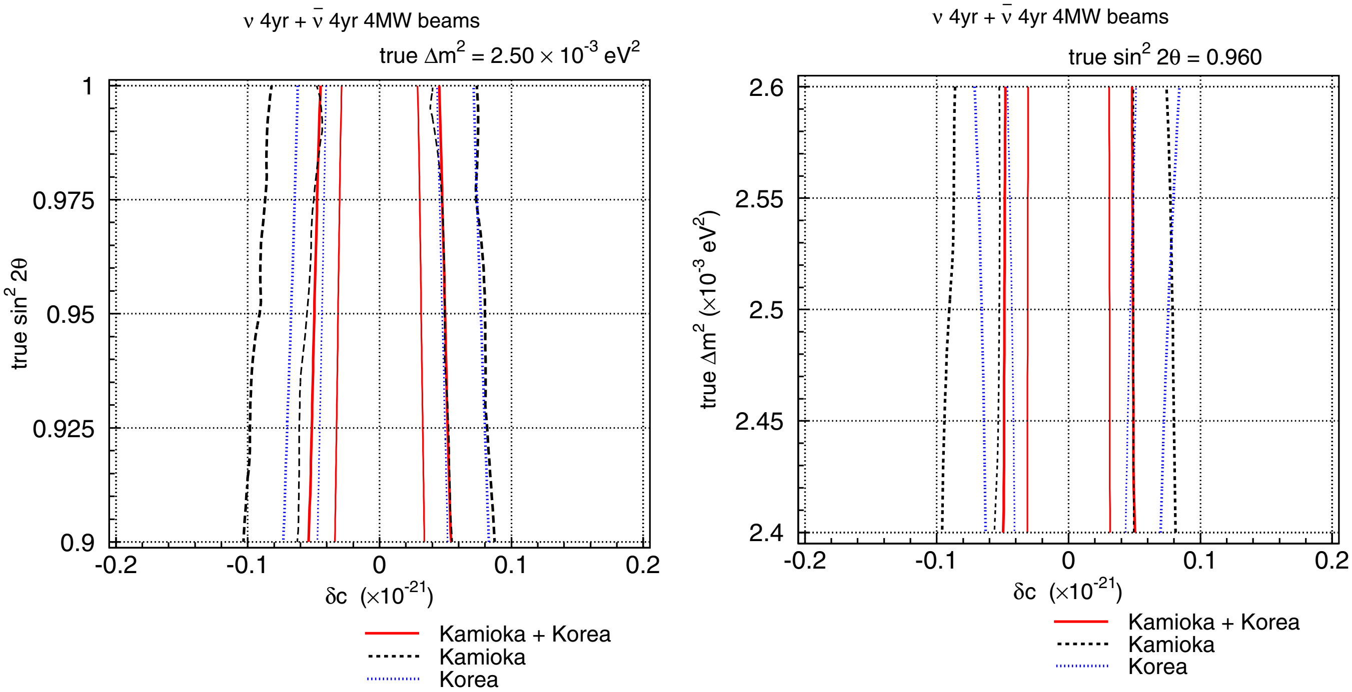

FIG. 4 (color online). The sensitivities to δc as a function of the true (input) value of $\sin^2 2\theta \equiv \sin^2 2\theta_{23}$ (left panel) and $\Delta m^2 \equiv \Delta m^2_{32}$ (right panel) obtained by fitting the input data corresponding to the vanishing δc with the nonzero δc. The red solid lines are for Kamioka-Korea setting with each 0.27 Mton detector, while the dashed black (dotted blue) lines are for Kamioka (Korea) only setting with 0.54 Mton detector. The thick and the thin lines are for 99% and 90% C.L., respectively. Four years of neutrino plus 4 years of antineutrino running are assumed. The other input values of the parameters are identical to those in Fig. 3.

C. Comparison between sensitivities of Kamioka-Korea setting and the existing bounds

We summarize the results of the previous subsections in Table II, along with the presently obtained bounds on δc and δb. We quote the current bounds on δc's from Refs. [40,42] which was obtained by the atmospheric neutrino data,

$$|\delta c_{\mu\tau}| \lesssim 3 \times 10^{-26}. \tag{10}$$

We note that the current bound on $\delta c_{\mu\tau}$ obtained by the atmospheric neutrino data is quite strong. The reason why the atmospheric neutrino data give a much stronger limit is that the relevant energy is much higher (typically $\sim$100 GeV) than the one we are considering ($\sim$ 1 GeV) and the baseline is larger, as large as the Earth's diameter.

For the bound on δb, Barger *et al.* [41] argue that

$$|\delta b_{\mu\tau}| < 3 \times 10^{-20}\ \text{GeV} \tag{11}$$

from the analysis of the atmospheric neutrino data. Let us compare the sensitivity on δb within our two-detector setup with the sensitivity at a neutrino factory. Barger *et al.* [41] considered a neutrino factory with 10^{19} stored muons with 20 GeV energy, and 10 kton detector, and concluded that it can probe $|\delta b| < 3 \times 10^{-23}$ GeV. The Kamioka-Korea two-detector setup and Korea-only setup have 5 and 6 times better sensitivities compared with the neutrino factory with the assumed configuration. Of course, the sensitivity achievable by a neutrino factory could be improved with a larger number of stored muons and a larger detector. A more meaningful comparison would be possible, only when one has configurations for both experiments which are optimized for the purposes of each experiment. But, we can still conclude that the two-detector setup is a powerful probe to the Lorentz symmetry violation, and could be competitive to neutrino factories.

VI. NONSTANDARD NEUTRINO INTERACTIONS WITH MATTER

It was suggested that neutrinos might have nonstandard neutral current interactions with matter [6–8,43], $\nu_\alpha + f \rightarrow \nu_\beta + f$ $(\alpha, \beta = e, \mu, \tau)$, with f being the up quarks, the down quarks, and electrons. This effect may be described by a low energy effective Hamiltonian for new nonstandard interactions (NSI) of neutrinos:

$$H_{\rm NSI} = 2\sqrt{2}G_F(\bar{\nu}_\alpha\gamma_\rho\nu_\beta)(\varepsilon^{ff'L}_{\alpha\beta}\bar{f}_L\gamma^\rho f'_L + \varepsilon^{ff'R}_{\alpha\beta}\bar{f}_R\gamma^\rho f'_R) + \text{H.c.}, \tag{12}$$

where G_F is the Fermi constant, $\varepsilon^f_{\alpha\beta} \equiv \varepsilon^{fL}_{\alpha\beta} + \varepsilon^{fR}_{\alpha\beta}$, and $\varepsilon^{fP}_{\alpha\beta} \equiv \varepsilon^{ffP}_{\alpha\beta}$ ($P = L, R$ indicates chirality). It is known that the presence of such NSI can affect production and/or detection processes of neutrinos as well as propagation of

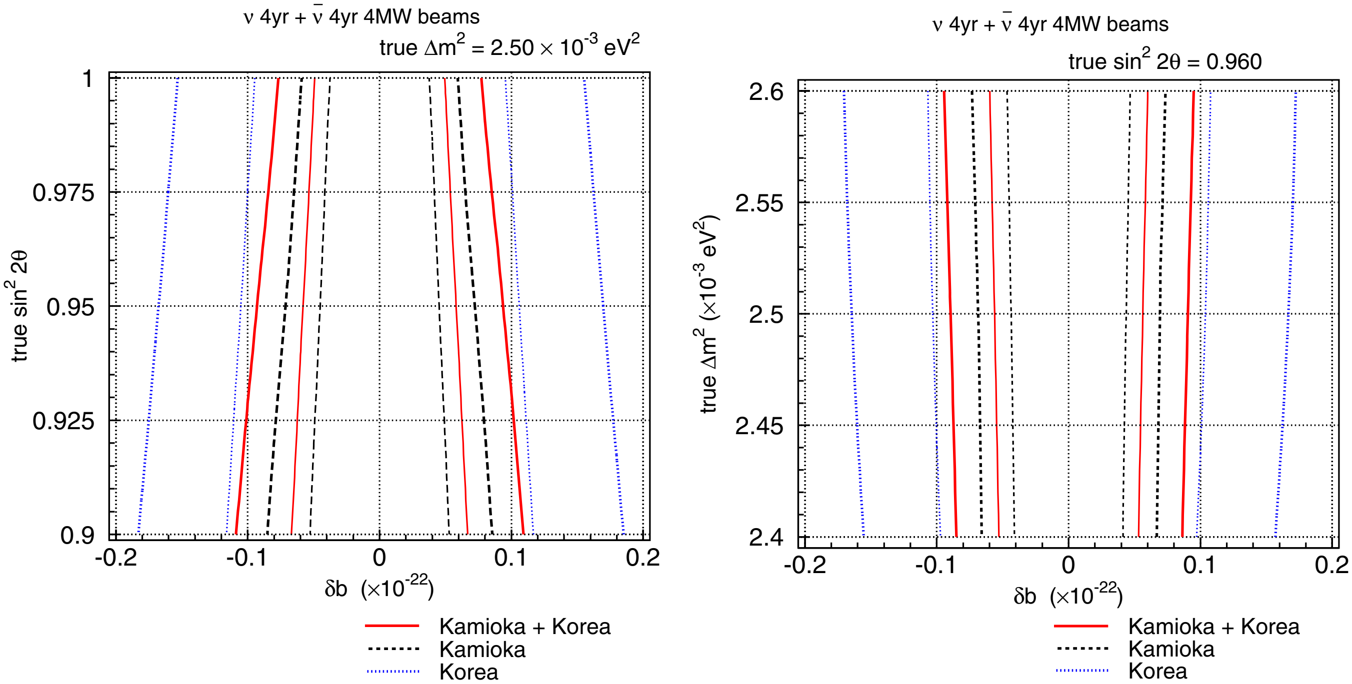

FIG. 5 (color online). The same as in Fig. 4 but for the case of nonvanishing δb.

neutrinos in matter. In this work, for simplicity, we consider the impact of NSI only for propagation. For previous analyses of the impact of NSI in long baseline experiments, see e.g., [44].

By using $\varepsilon_{\alpha\beta}$ defined as $\varepsilon_{\alpha\beta} \equiv \sum_{f=u,d,e} \varepsilon^f_{\alpha\beta} N_f/N_e$, where N_f denotes the number density of relevant fermions, the effects of NSI may be summarized by a term with dimensionless parameters $\varepsilon_{\alpha\beta}$ in the effective Hamiltonian

$$H_{\rm eff} = \sqrt{2} G_F N_e \begin{pmatrix} \varepsilon_{ee} & \varepsilon_{e\mu} & \varepsilon_{e\tau} \\ \varepsilon^*_{e\mu} & \varepsilon_{\mu\mu} & \varepsilon_{\mu\tau} \\ \varepsilon^*_{e\tau} & \varepsilon^*_{\mu\tau} & \varepsilon_{\tau\tau} \end{pmatrix}, \tag{13}$$

which is to be added to the standard matter term $\sqrt{2} G_F N_e {\rm diag.}(1, 0, 0)$ [6] in the evolution equation of neutrinos.

In this work we truncate the system so that we are confined into the $\mu - \tau$ sector of the neutrino evolution. Then, the time evolution of the neutrinos in flavor basis can be written as

$$i \frac{d}{dt} \begin{pmatrix} \nu_\mu \\ \nu_\tau \end{pmatrix} = \left[U \begin{pmatrix} 0 & 0 \\ 0 & \frac{\Delta m^2_{32}}{2E} \end{pmatrix} U^\dagger + a \begin{pmatrix} 0 & \varepsilon_{\mu\tau} \\ \varepsilon^*_{\mu\tau} & \varepsilon_{\tau\tau} - \varepsilon_{\mu\mu} \end{pmatrix} \right] \times \begin{pmatrix} \nu_\mu \\ \nu_\tau \end{pmatrix}, \tag{14}$$

where U is the flavor mixing matrix and $a \equiv \sqrt{2} G_F N_e$. The 2-2 element of the NSI term in the Hamiltonian is of the form $\varepsilon_{\tau\tau} - \varepsilon_{\mu\mu}$ because the oscillation probability depends upon ε's only through this combination. The evolution equation for the antineutrinos are given by changing the signs of a and replacing U and $\varepsilon_{\mu\tau}$ by U^* and $\varepsilon^*_{\mu\tau}$, respectively.

In fact, one can show that as long as we can ignore the effect of NSI of the order equal to or higher than ε^2, the truncation to the 2×2 subsystem is a good approximation. In the full three-flavor framework the ν_μ disappearance oscillation probability can be computed to leading order of NSI as [45]

TABLE II. Presented are the upper bounds on the velocity mixing parameter δc and the CPT-violating parameter δb (in GeV) for the case where $\theta_m = \theta_b = \theta_c \equiv \theta$ and $\eta = \eta' = 0$. The current bounds are based on [40,41] and are at 90% C.L.. The sensitivities obtained in this study are also at 90% C.L. and correspond to the true values of the parameters $\Delta m^2 = 2.5 \times 10^{-3}$ eV^2 and $\sin^2 2\theta_{23} = 0.96$.

LV parameters	Curent bound	Kamioka only	Korea only	Kamioka-Korea
$\vert\delta c\vert$	$\lesssim 3 \times 10^{-26}$	$\lesssim 5 \times 10^{-23}$	$\lesssim 4 \times 10^{-23}$	$\lesssim 3 \times 10^{-23}$
$\vert\delta b\vert$(GeV)	$<3.0 \times 10^{-20}$	$\lesssim 0.5 \times 10^{-23}$	$\lesssim 1 \times 10^{-23}$	$\lesssim 0.6 \times 10^{-23}$

$$P(\nu_\mu \to \nu_\mu) = 1 - \sin^2 2\theta_{23}\sin^2\Delta_{32} - |\varepsilon_{\mu\tau}|\cos\phi_{\mu\tau} \times \sin 2\theta_{23}(aL)\left[\sin^2 2\theta_{23}\sin 2\Delta_{32} + \cos^2 2\theta_{23}\frac{2}{\Delta_{32}}\sin^2\Delta_{32}\right] - \frac{1}{2}(\varepsilon_{\tau\tau} - \varepsilon_{\mu\mu})\sin^2 2\theta_{23}\cos 2\theta_{23}(aL) \times \left[\sin 2\Delta_{32} - \frac{2}{\Delta_{32}}\sin^2\Delta_{32}\right] + \mathcal{O}\left(\frac{\Delta m^2_{21}}{\Delta m^2_{31}}\right) + \mathcal{O}(s_{13}) + \mathcal{O}(\varepsilon^2), \quad (15)$$

where $\Delta_{32} \equiv \frac{\Delta m^2_{32}L}{4E}$ and $\phi_{\mu\tau}$ is the phase of $\varepsilon_{\mu\tau}$. The result in (15) indicates that the truncation is legitimate if ε's are sufficiently small. Also note that the dependence on $\varepsilon_{\tau\tau} - \varepsilon_{\mu\mu}$ goes away for $\sin^2\theta = 0.5$, so that the muon disappearance becomes insensitive to this combination of ε's. We will find that the sensitivity to $\varepsilon_{\tau\tau} - \varepsilon_{\mu\mu}$ strongly depends on $\sin^2\theta$ (θ being maximal or not) for this reason. In the following, we set $\epsilon_{\mu\mu} = 0$ so that the sensitivity contours presented for $\epsilon_{\tau\tau}$ actually means those for $\epsilon_{\tau\tau} - \epsilon_{\mu\mu}$. Moreover, for simplicity, we assume that $\epsilon_{\mu\tau}$ is real by ignoring its phase.

Since we work within the truncated 2 by 2 subsystem, we quote here only the existing bounds of NSI parameters which are obtained under the same approximation. By analyzing the super-Kamiokande and MACRO atmospheric neutrino data the authors of [46] obtained

$$|\varepsilon_{\mu\tau}| \lesssim 0.09, \qquad |\varepsilon_{\tau\tau}| \lesssim 0.15, \quad (16)$$

at 99% C.L. for 2 degrees of freedom.[3] See [48,49] for the current status of constraints on other NSI elements $\varepsilon_{\alpha\beta}$.

In Fig. 6, presented are the allowed regions in $\varepsilon_{\mu\tau} - \varepsilon_{\tau\tau}$ space for 4 years neutrino and 4 years antineutrino running of the Kamioka-only (upper panels), the Korea-only (middle panels), and the Kamioka-Korea (bottom panels) settings. The input values $\varepsilon_{\mu\tau}$ and $\varepsilon_{\tau\tau}$ are taken to be vanishing.

As in the CPT-Lorentz violating case studied in Sec. V B and unlike the system with decoherence, the Korea-only setting gives much worse sensitivity compared to the other two settings. Again the Kamioka-only setting has a slightly better sensitivity than the Kamioka-Korea setting. However, we notice that the Kamioka-only setting has multiple $\varepsilon_{\tau\tau}$ solutions for $\sin^2\theta_{23} = 0.45$. The fake solutions are nearly eliminated in the Kamioka-Korea setting.

The sensitivities of three experimental setups at 2σ C.L. can be read off from Fig. 6. The approximate 2σ C.L. sensitivities of the Kamioka-Korea setup for $\sin^2\theta = 0.45(\sin^2\theta = 0.5)$ are

[3] A less severe bound on $|\varepsilon_{\tau\tau}|$ is derived in [47] by analyzing the same data but with $\varepsilon_{e\tau}$ and without $\varepsilon_{\mu\tau}$

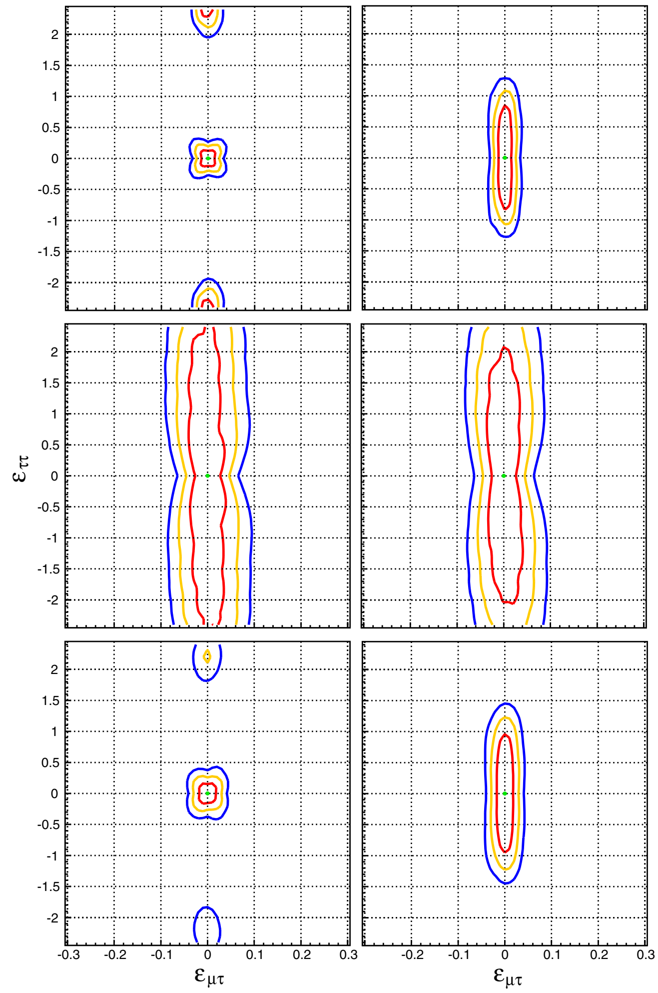

FIG. 6 (color online). The allowed regions in $\varepsilon_{\mu\tau} - \varepsilon_{\tau\tau}$ space for 4 years neutrino and 4 years antineutrino running. The upper, middle, and bottom three panels are for the Kamioka-only setting, the Korea-only setting, and the Kamioka-Korea setting, respectively. The left and the right panels are for cases with $\sin^2\theta \equiv \sin^2\theta_{23} = 0.45$ and 0.5, respectively. The inner (red), middle (yellow), and outer (blue) curves indicate the allowed regions at 1σ, 2σ, and 3σ C.L., respectively, for 2 degrees of freedom. The input value of Δm^2_{32} is taken as 2.5×10^{-3} eV^2.

$$|\epsilon_{\mu\tau}| \lesssim 0.03(0.03), \qquad |\epsilon_{\tau\tau}| \lesssim 0.3(1.2). \quad (17)$$

Here, we neglected a barely allowed region near $|\epsilon_{\tau\tau}| = 2.3$, which is already excluded by the current data. Note that the sensitivity on $\varepsilon_{\tau\tau}$ becomes weak for maximal mixing ($\sin^2\theta = 0.5$), for the above mentioned reason [see Eq. (15) and the subsequent discussions]. The Kamioka-Korea setup can improve the current bound on $|\epsilon_{\mu\tau}|$ by factors of $\sim$5 whereas the bound on $|\epsilon_{\tau\tau}|$ we obtained is comparable to (worse than) the current bound for $\sin^2\theta = 0.45(\sin^2\theta = 0.5)$. A similar statement applies to the case for the Kamioka-only setup.

There are a large number of references which studied the effects of NSI and the sensitivity reach to NSI parameters by the ongoing and the various future projects. We quote here only the most recent ones which focused on sensitivities by superbeam and reactor experiments [45] and neutrino factory [50]. The earlier references can be traced back through the bibliographies of these papers.

By combining the future superbeam and the reactor experiments, T2K [21] and Double-Chooz [24], the authors of [45] obtained the "discovery reach" of $|\epsilon_{\mu\tau}|$ to be $\sim$0.25 when it is assumed to be real (no CP phase) while essentially no discovery potential for $\epsilon_{\tau\tau}$ is expected (see Fig. 11 of [45]).[4] The same authors also consider the case of the NOνA experiment [22] combined with some future upgraded reactor experiment with larger detector as considered, e.g., in [51,52], and obtained the discovery reach of $\epsilon_{\mu\tau}$ to be about 0.05 which is comparable to what we obtained (see Fig. 12 of [45]).

While essentially no sensitivity of $\epsilon_{\tau\tau}$ is expected by superbeam, a future neutrino factory with the so called golden channel $\nu_e \to \nu_\mu$ and $\bar{\nu}_e \to \bar{\nu}_\mu$ could reach the sensitivity to $\epsilon_{\tau\tau}$ at the level of $\sim$0.1–0.2 [50]. Despite that the sensitivity to $\epsilon_{\mu\tau}$ by neutrino factory was not derived in [50], from Fig. 1 of this reference, one can naively expect that the sensitivity to $\epsilon_{\mu\tau}$ is similar to that of ϵ_{ee} which is $\sim$0.1 or so. We conclude that the sensitivity we obtained for $\epsilon_{\mu\tau}$ is quite reasonable.

VII. CONCLUDING REMARKS

The Kamioka-Korea two-detector system was shown to be a powerful setup for lifting neutrino parameter degeneracies and probing CP violation in neutrino oscillation in a robust manner. In this paper, we study sensitivities of this setup to nonstandard neutrino physics such as quantum decoherence, tiny violation of Lorentz symmetry, and nonstandard interactions of neutrinos with matter. In most cases, a two-detector setup is more sensitive than a single detector at Kamioka or Korea, except for the Lorentz violation with $\delta b \neq 0$, and the nonstandard neutrino interactions with matter. The sensitivities of three experimental setups at 90% C.L. are summarized in Tables I and II for quantum decoherence and Lorentz symmetry violation with/without CPT symmetry, respectively. We can say that future long baseline experiments with two-detector setups can improve the sensitivities on nonstandard neutrino physics in many cases. We believe that it is a useful addition to the physics capabilities of the Kamioka-Korea two-detector setting that are already demonstrated, namely, resolution of the mass hierarchy and resolving CP and the octant degeneracies.

[4]Note that the discovery reach defined in Ref. [45] is expected to be similar to our "sensitivity" but they are not exactly the same. Both are defined as a range of the ε parameter which gives acceptable fit to its assumed value at a given confidence level. The only difference is that while they use nonvanishing input values of NSI parameters, we use vanishing ones.

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