Can the observed CP asymmetry in $\tau \to K\pi\nu_\tau$ be due to nonstandard tensor interactions?

H. Zeen Devi,* L. Dhargyal,† and Nita Sinha‡

The Institute of Mathematical Sciences, Chennai 600113, India

(Received 18 September 2013; published 18 July 2014)

An intriguing opposite sign of the CP-violating asymmetry from that expected within the Standard Model was recently measured in the tau decay modes $\tau^\pm \to K_s\pi^\pm \overset{(-)}{\nu}_\tau$ by the *BABAR* collaboration. If this result is confirmed with higher precision, the observed decay rate asymmetry A_{CP} can only arise from some nonstandard interactions occurring possibly in both the hadronic as well as in the leptonic sectors. We illustrate that, while a simple charged scalar interaction cannot yield this rate asymmetry, it will be possible to generate this in the presence of a tensor interaction. Parametrizing the strength and weak phase of this nonstandard interaction contribution, the observed branching ratio and the decay-rate CP asymmetry for the particular mode $\tau^\pm \to K_s\pi^\pm \overset{(-)}{\nu}_\tau$ are used to determine the CP-violating weak phase and the coupling of a tensorial interaction that can give a consistent sign and magnitude of the asymmetry.

DOI: 10.1103/PhysRevD.90.013016 PACS numbers: 13.35.Dx, 14.40.Df, 14.60.Fg

I. INTRODUCTION

While the Kobayashi Maskawa ansatz for CP violation within the Standard Model (SM) [1] in the quark sector has been clearly verified by the plethora of data from the B factories, this is unable to account for the observed baryon asymmetry of the Universe. Hence, one needs to look for other sources of CP violation, including searches in the leptonic sector. Apart from the CP phases that may arise in the neutrino mixing matrix, the decays of the charged tau leptons may allow us to explore nonstandard CP-violating interactions.

In fact, CP violation through decays of tau leptons was studied in a series of papers by Tsai [2], around the time of the preliminary design for a tau-charm factory. Various experimental groups have been involved in exploring CP violation in tau decays in the last decade or more. In 2002, the CLEO collaboration [3], and more recently the Belle Collaboration [4], studied the angular distribution of the decay products in $\tau^\pm \to K_s^0\pi^\pm \overset{(-)}{\nu}_\tau$ in search of CP violation; however, neither study revealed any CP asymmetry. The *BABAR* collaboration [5] for the first time reported a measurement of the rate asymmetry A_{CP} in this decay mode to be

$$A_{CP}^{\text{Exp.}} = (-0.36 \pm 0.23 \pm 0.11)\%. \tag{1}$$

On the theoretical side, for $\tau^\pm \to K_s^0\pi^\pm \overset{(-)}{\nu}_\tau \to [\pi^+\pi^-]_K\pi^\pm \overset{(-)}{\nu}_\tau$, Bigi and Sanda [6] predicted the CP asymmetry to be

$$A_{CP}^{\text{SM}} = (+0.33 \pm 0.01)\%, \tag{2}$$

where the CP violation arises from the *known* $K^0 - \bar{K}^0$ mixing. Recently, Grossman and Nir [7], comparing the rate asymmetries for decays to neutral kaons of the taus with that of D mesons, pointed out that since $\tau^+(\tau^-)$ decays initially to a $K^0(\bar{K}^0)$ whereas $D^+(D^-)$ decays initially to $\bar{K}^0(K^0)$, the time-integrated decay-rate CP asymmetry (arising from oscillations of the neutral kaons) of τ decays must have a sign opposite to that of D decays. Further, they emphasized that the decay asymmetry is affected by the reconstruction efficiency as a function of the $K_s \to \pi^+\pi^-$ decay time. Using the parametrization of Ref. [7], *BABAR* has obtained a multiplicative correction factor for the decay-rate asymmetry and predicts the SM decay-rate asymmetry to be

$$A_{CP}^{\text{SM}} = (+0.36 \pm 0.01)\%. \tag{3}$$

As reported in Refs. [8–11], the CP asymmetry in $D^\pm \to \pi K_s$ decay has been measured to be $(-0.54 \pm 0.14)\%$.

The observation of a CP asymmetry in tau decays to K_s having the same sign as that in D decays, and moreover of the same magnitude but opposite in sign to the corrected SM expectation, implies that this asymmetry cannot be accounted for by the CP violation in $K^0 - \bar{K}^0$ mixing. Apart from this mixing contribution to CP violation, within the SM, since there is only a single amplitude with the W boson mediating the decay process, the observed CP asymmetry cannot be explained. Hence, this may be a signal of physics beyond the Standard Model, if confirmed to a higher statistical significance.

It should be pointed out that the *BABAR* collaboration has accounted for the modification of the decay-rate asymmetry due to different nuclear interaction cross sections of the K^0 and $\bar{K}^0$ mesons with the detector material

*zeen@imsc.res.in
†dhargyal@imsc.res.in
‡nita@imsc.res.in

through a correction, calculated on an event by event basis. Also, *BABAR* [5] claims that, using a control sample in data and Monte Carlo simulations, they have verified that no significant decay-rate asymmetry is induced by their detector or the selection criteria. Further, the Standard Model asymmetry is identical for decays with any number of π^0 mesons and hence can be searched for, in all the modes $\tau^\pm \to \pi^\pm K_s^0(\geq 0\pi^0)\overset{(-)}{\nu}_\tau$. We note that all the experiments CLEO, Belle, and *BABAR* assume that the *CP* asymmetry is conserved at the tau production vertex.

In the presence of an additional new physics (NP) amplitude with a complex coupling, along with the strong phases from the $K-\pi$ scattering, which can be large particularly in this resonance-dominated region, the observed *CP* asymmetry may be attainable. This nonstandard interaction (NSI) could possibly affect both the hadronic and leptonic currents. The Feynman diagrams for the SM decay mode of tau via W exchange and via exchange of some exotic particle X, in the presence of NP, are shown in Fig. 1.

Naively, one would expect a charged scalar boson exchange to provide the required additional diagram, where, with a complex weak coupling and the difference in the strong phases of the scalar and vector hadronic form factors, a *CP* violating asymmetry could arise. However, such an asymmetry appears only in the difference of the $\tau^\pm \to K_S\pi^\pm\nu_\tau(\bar{\nu}_\tau)$ decay angular distributions but vanishes in the integrated difference of the decay rates for τ^+ and τ^-, measured by *BABAR*. In fact, CLEO had used their nonobservation of an asymmetry in the distribution, to set a bound [3] on the imaginary part of the complex coupling of the scalar boson, such as a charged Higgs. Similar limits were set by Belle [4] on the *CP*-violation parameter modifying the scalar form factor, since their differential asymmetry was compatible with zero. The new physics amplitude that can account for a decay rate asymmetry therefore has to appear from a different kind of interaction, and we investigate whether a tensor interaction can produce the asymmetry reported by *BABAR* and, in fact, use the measured branching ratio and asymmetry to constrain the parameters of this kind of new interaction.

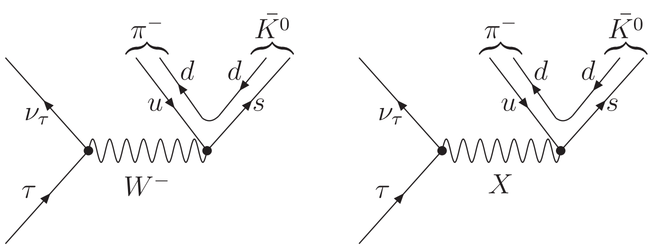


FIG. 1. The Feynman diagram for the Standard Model decay mode via W exchange is shown in the left panel, and that for the New physics exchange diagram via some exotic particle X is shown in the right panel.

In Sec. II, we evaluate the decay-rate asymmetry in the presence of any generic NP amplitude. The decay rate for the SM case is calculated in Sec. III, while in Sec. IV, the total rate is evaluated in the presence of the new additional tensor interaction term. The observables, the branching ratio, and the decay-rate asymmetry are used to estimate the parameters of this new interaction in Sec. V, and in Sec. VI, we conclude.

II. BRANCHING RATIO AND RATE ASYMMETRY IN $\tau \to K\pi\nu_\tau$ IN THE PRESENCE OF A GENERIC NEW PHYSICS AMPLITUDE

The amplitude for the decay of $\tau^+ \to K_s\pi^+\bar{\nu}_\tau$ in the presence of NSI can be written as

$$\mathcal{A} = \mathcal{A}^{\rm SM} + \mathcal{A}^{\rm NSI} e^{i\phi} e^{i\delta}, \tag{4}$$

where $A^{\rm SM}$ and $A^{\rm NSI}$ are the magnitudes of the SM and NP amplitudes, respectively, while ϕ is the weak phase of the NP contribution (since V_{us} is real, there is no weak phase in the SM contribution, except that coming from the neutral K meson mixing, which is accounted for, separately in the theoretical SM expectation). The relative strong phase of the $(K\pi)$ system, δ, between the NP amplitude with respect to the SM contribution is a function of the $K\pi$ invariant mass squared. The amplitude for the antiprocess $\tau^- \to K_s\pi^-\nu_\tau$ has the opposite weak phase but the same strong phase.

In the presence of a NSI for which the interference with SM is nonvanishing even after the angular integrations,[1] the general expression for the differential decay rate of $\tau \to K\pi\nu$ may be written as

$$d\Gamma \propto |\mathcal{A}|^2 \times dQ^2, \tag{5}$$

$$\propto [|\mathcal{A}^{\rm SM}|^2 + |\mathcal{A}^{\rm NSI}|^2 + 2|\mathcal{A}^{\rm SM}||\mathcal{A}^{\rm NSI}|\cos(\phi + \delta(Q^2))]dQ^2, \tag{6}$$

where Q is the sum of the hadron momenta. The differential decay rate for the antiprocess is

$$d\bar{\Gamma} \propto |\bar{\mathcal{A}}|^2 \times dQ^2, \tag{7}$$

$$\propto [|\mathcal{A}^{\rm SM}|^2 + |\mathcal{A}^{\rm NSI}|^2 + 2|\mathcal{A}^{\rm SM}||\mathcal{A}^{\rm NSI}|\cos(-\phi + \delta(Q^2))]dQ^2. \tag{8}$$

The branching ratio for $\tau \to K_s\pi\nu$ is the ratio of the average of the width of $\tau^+ \to K_s\pi^+\bar{\nu}$ and $\tau^- \to K_s\pi^-\nu$ to the total width of τ ($\Gamma_{\rm total}$). Hence,

$$BR(\tau \to K_s\pi\nu) = \frac{\Gamma + \bar{\Gamma}}{2\Gamma_{\rm total}}. \tag{9}$$

Now,

[1]Details about the integration variables are specified in the next section.

$$\frac{d\Gamma}{dQ^2}+\frac{d\bar{\Gamma}}{dQ^2}\propto 2[|\mathcal{A}^{\rm SM}|^2+|\mathcal{A}^{\rm NSI}|^2+2|\mathcal{A}^{\rm SM}||\mathcal{A}^{\rm NSI}|\cos(\phi)\cos(\delta(Q^2))], \tag{10}$$

$$\propto 2\mathcal{A}^{\rm SM}|^2[1+(r(Q^2))^2+2r(Q^2)\cos(\phi)\cos(\delta(Q^2))], \tag{11}$$

where $r(Q^2)$ is the ratio of the amplitude of the NSI contribution to the SM contribution. Therefore,

$$BR(\tau\to K_s\pi\nu)=\frac{\int dQ^2 2|\mathcal{A}^{\rm SM}|^2[1+(r(Q^2))^2+2r(Q^2)\cos(\phi)\cos(\delta(Q^2))]}{2\Gamma_\tau}. \tag{12}$$

Similarly, the difference is

$$\frac{d\Gamma}{dQ^2}-\frac{d\bar{\Gamma}}{dQ^2}\propto -4|A^{\rm SM}||A^{\rm NSI}|\sin(\phi)\sin(\delta(Q^2)), \tag{13}$$

$$\propto -4r(Q^2)\sin(\phi)\sin(\delta(Q^2))\frac{d\Gamma^{\rm SM}}{dQ^2}. \tag{14}$$

Hence, the integrated rate asymmetry is

$$A^\tau_{CP}=\frac{\Gamma-\bar{\Gamma}}{\Gamma+\bar{\Gamma}}=-\frac{\int 4r(Q^2)\sin(\phi)\sin(\delta(Q^2))\frac{d\Gamma^{\rm SM}}{dQ^2}dQ^2}{\Gamma+\bar{\Gamma}}. \tag{15}$$

As pointed out in Refs. [12] and [13], this *CP* asymmetry being linear in the new physics amplitude has a higher sensitivity to it than effects like lepton-flavor violation, electric dipole moments, etc., which depend quadratically on the NP amplitude, rendering this observation to play an important role in uncovering physics beyond the SM. Using Eqs. (12) and (15) and the Particle Data Group [14] value of the branching ratio and the rate asymmetry measured by *BABAR*, the weak phase and the magnitude of the new physics contribution can be estimated.

III. DECAY RATE OF $\tau\to K\pi\nu$ IN THE STANDARD MODEL

The tau leptonic and the hadronic decay amplitudes can be factorized into a purely leptonic part including the tau and the neutrino and a hadronic part, where the hadronic system is created from the QCD vacuum via the charged weak current.

Hence, the differential decay rate of the process $\tau(p_\tau)\to K(p_K)+\pi(p_\pi)+\nu(p_{\nu_\tau})$ [15] may be written as

$$d\Gamma(\tau\to K\pi\nu)=\frac{1}{2m_\tau}\frac{G_F^2}{2}\sin^2\theta_c\mathcal{L}_{\mu\nu}\mathcal{H}^{\mu\nu}dPS^{(3)}, \tag{16}$$

where $\mathcal{L}_{\mu\nu}$ is the leptonic term,

$$\mathcal{L}_{\mu\nu}=[\bar{\nu}_\tau\gamma_\mu(1-\gamma_5)\tau][\bar{\nu}_\tau\gamma_\nu(1-\gamma_5)\tau]^\dagger, \tag{17}$$

and the hadronic term,

$$\mathcal{H}^{\mu\nu}=\mathcal{J}^\mu(\mathcal{J}^\nu)^\dagger, \tag{18}$$

is given in terms of the hadronic vector current,

$$\mathcal{J}^\mu=\langle K(p_K)\pi(p_\pi)|V^\mu(0)|0\rangle. \tag{19}$$

The hadronic matrix element for the transition from vacuum to two pseudoscalar mesons state will have scalar and vector components of the weak charged current, corresponding to J^P values of 0^+ and 1^-, respectively. The hadronic vector current in Eq. (19) is parametrized in terms of the scalar and the vector form factors as

$$\mathcal{J}^\mu=F_V^{K\pi}(Q^2)\left(g^{\mu\nu}-\frac{Q^\mu Q^\nu}{Q^2}\right)(p_k-p_\pi)_\nu+F_s^{K\pi}Q^\mu, \tag{20}$$

where $Q^\mu=p_k^\mu+p_\pi^\mu$. The decay rate involves only the mod squared of the vector and scalar form factors but no scalar-vector interference term. In the hadronic rest frame where $\vec{p}_k+\vec{p}_\pi=0$, it takes the form

$$d\Gamma_{\rm SM}=\frac{1}{2m_\tau}\times\frac{G_F^2\sin^2\theta_c}{2}V_{us}S_{EW}(m_\tau^2-Q^2)\left\{\left(\frac{2Q^2+m_\tau^2}{3Q^2}\right)4[P(Q^2)]^2|F_V|^2+\frac{m_\tau^2}{Q^2}\frac{(m_K^2-m_\pi^2)^2}{Q^2}|F_S|^2\right\}dPS^{(3)}, \tag{21}$$

where $S_{EW}=1.02$ [16] is the electroweak correction factor and $P(Q^2)\equiv|\vec{p}_k|$ is the momentum of the kaon in the $K\pi$ rest frame, which is a function of the $K\pi$ invariant mass squared Q^2 and may be expressed as

$$P(Q^2) = \frac{1}{2\sqrt{Q^2}}\sqrt{[Q^2-(m_k+m_\pi)^2][Q^2-(m_k-m_\pi)^2]}. \tag{22}$$

After integrating out the neutrino momentum, the phase space in the $K\pi$ rest frame is

$$dPS^{(3)} = \frac{1}{(4\pi)^3}\frac{(m_\tau^2-Q^2)}{m_\tau^2}|\vec{p}_k|\frac{dQ^2}{\sqrt{Q^2}}\frac{d\cos\beta}{2}, \tag{23}$$

where β is the direction of the kaon with respect to the tau direction, denoted by $\hat{n}_\tau$, viewed from the hadronic rest frame, i.e., $\cos\beta = \hat{p}_K\cdot\hat{n}_\tau$, where $\hat{p}_K = \frac{\vec{p}_K}{|\vec{p}_K|}$. Hence, the differential decay rate takes the form

$$\frac{d\Gamma_{\rm SM}}{d\sqrt{Q^2}} = \frac{G_F^2\sin^2\theta_c m_\tau^3}{3\times 2^5\times\pi^3 Q^2}S_{EW}\left(1-\frac{Q^2}{m_\tau^2}\right)^2\left(1+\frac{2Q^2}{m_\tau^2}\right) \times P(Q^2)\left\{P(Q^2)^2|F_V|^2+\frac{3(m_k^2-m_\pi^2)^2}{4Q^2(1+\frac{2Q^2}{m_\tau^2})}|F_S|^2\right\}. \tag{24}$$

The hadronic current is dominated by many resonances, the vector ones [$K^*(892)$, $K^*(1410)$, and $K^*(1680)$] and the scalar ones [$K_0^*(800)$ and $K_0^*(1430)$]. The form factors can be parametrized in terms of Briet–Wigner forms with energy-dependent widths. Hence, the vector form factor may be written as [17]

$$F_V = \frac{1}{1+\beta+\chi}[BW_{K^*(892)}(Q^2)+\beta BW_{K^*(1410)}(Q^2)+\chi BW_{K^*(1680)}(Q^2)], \tag{25}$$

where β and χ are the complex coefficients for the relative contributions of $K^*(1410)$ and $K^*(1680)$ resonances, respectively, with respect to the dominant $K^*(892)$ contribution and $BW_R(s)$ is a relativistic Breit–Wigner function corresponding to R being $K^*(892)$, $K^*(1410)$, or $K^*(1680)$ for the vector case.

For each of the resonances, the Breit–Wigner function has the form

$$BW_R(Q^2) = \frac{M_R^2}{Q^2-M_R^2+i\sqrt{Q^2}\Gamma_R(Q^2)}, \tag{26}$$

where

$$\Gamma_R(Q^2) = \Gamma_{0R}\frac{M_R^2}{Q^2}\left(\frac{P(Q^2)}{P(M_R^2)}\right)^{(2l+1)}. \tag{27}$$

Here, $\Gamma_R(s)$ is the s-dependent total width of the resonance, and $\Gamma_{0R}(s)$ is the resonance width at its peak. The orbital angular momentum $l=1$, if the $K\pi$ system is in a p-wave or in a vector state and $l=0$, for the s-wave or scalar state. The scalar form factor F_S has $K_0^*(800)$ and $K_0^*(1430)$ contributions and has the similar form

$$F_S = \varkappa\frac{Q^2}{M^2_{K_0^*(800)}}BW_{K_0^*(800)}(Q^2) + \gamma\frac{Q^2}{M^2_{K_0^*(1430)}}BW_{K_0^*(1430)}(Q^2), \tag{28}$$

where, $\varkappa$ and γ are the complex constants that describe the relative contributions of the $K_0^*(800)$ and $K_0^*(1430)$ resonances, respectively. The Belle collaboration had performed fits to the $K_s\pi^-$ invariant mass spectrum (Q^2) distribution and had listed the values of the complex constants β, $\varkappa$, χ, and γ in Ref. [17]. Their fitted results (as well as those of *BABAR* reported in Ref. [18]) demonstrated that a $K^*_{(892)}$ alone is not enough to describe the $K_s\pi$ mass spectrum, but rather the distribution shows the clear evidence for a scalar contribution in the low invariant mass and another component at large Q^2. The fits were best explained with either $K^*(892)+K^*(1410)+K^*(800)$ or $K^*(892)+K^*(1430)+K^*(800)$. Therefore, we have used these two possibilities for our study and have excluded $K^*(1680)$ in our study as its inclusion worsens the fit quality in the Belle analysis. We would also like to point out that the Belle fit results are also consistent with a theoretical description using chiral perturbation theory, described in Ref. [19], and that of Ref. [20], which is based on analycity and $K-\pi$ scattering results.

IV. TOTAL RATE IN THE PRESENCE OF A NEW TENSOR INTERACTION

We now propose an additional tensor contribution to the amplitude. We explore if the interaction of the SM with the new tensorial interaction can account for a nonvanishing *CP* asymmetry in the decay mode $\tau\to K\pi\nu_\tau$. This tensor interaction could arise in various NP models; however, our approach is to study just the effect of this new structure and keep the analysis as model independent as possible. Interference of the SM with this tensor amplitude must give a nonvanishing *CP* asymmetry for this particular decay mode; moreover, the sign and magnitude the *CP* asymmetry must be consistent with the observed result.

The effective Hamiltonian due to this new operator is written as

$$\mathcal{H}_{\rm eff}^{\rm NSI} = G'\sin\theta_c(\bar{s}\sigma_{\mu\nu}u)(\bar{\nu}_\tau\sigma^{\mu\nu}(1+\gamma_5)\tau), \tag{29}$$

where G' is a complex coupling,

$$G' \equiv \frac{R_T G_F}{\sqrt{2}}. \tag{30}$$

The new physics amplitude is then given by

$$\mathcal{A}_T = G'[\langle K\pi|\bar{s}\sigma_{\mu\nu}u|0\rangle][\bar{u}(p_\nu)\sigma^{\mu\nu}(1+\gamma_5)u(p_\tau)], \tag{31}$$

where $\sigma^{\mu\nu} = \frac{i}{2}(\gamma^\mu\gamma^\nu - \gamma^\nu\gamma^\mu)$ and the Hadronic current is given by

$$\langle K(p_k)\pi(p_\pi)|\bar{s}\sigma_{\mu\nu}u|0\rangle = i\frac{F_T}{(m_K+m_\pi)}[p_k^\mu p_\pi^\nu - p_k^\nu p_\pi^\mu], \tag{32}$$

with F_T being the form factor due to the tensor interaction. In $\mathcal{A}_T$, we include only left-handed neutrinos, as the similar term with a right-handed neutrino, in the interference of the SM and NP contributions at the decay-rate level, will be suppressed by the neutrino mass. We note here that the A_T is a special case of the generalized A^{NSI} mentioned in Sec. III. More specifically now, the effective amplitude A is now given by

$$\begin{aligned}\mathcal{A} &= A_{\mathrm{SM}} + |A_T|e^{i\phi}e^{i\delta}\\ |\mathcal{A}|^2 &= |\mathcal{A}_{\mathrm{SM}}|^2 + |\mathcal{A}_T|^2 + 2\mathrm{Re}(\mathcal{A}_{\mathrm{SM}}\mathcal{A}_T^\dagger),\end{aligned} \tag{33}$$

where

$$\begin{aligned}Re(\mathcal{A}_{\mathrm{SM}}\mathcal{A}_T^\dagger) \propto 16m_\tau R_T[&(p_\nu\cdot p_k)\cdot(p_k\cdot p_\pi) - p_k^2(p_\nu\cdot p_\pi)\\ &- p_\pi^2(p_\nu\cdot p_k) + (p_\nu\cdot p_\pi)\\ &\cdot(p_k\cdot p_\pi)]|F_V||F_T|\cos(\delta_V-\delta_T)\end{aligned} \tag{34}$$

and

$$\begin{aligned}|\mathcal{A}_T|^2 \propto 32|R_TF_T|^2[&2(p_\nu\cdot p_\pi)(p_\tau\cdot p_k)(p_k\cdot p_\pi)\\ &+ 2(p_\nu\cdot p_k)(p_\tau\cdot p_\pi)(p_k\cdot p_\pi)\\ &- 2m_{pi}^2(p_\tau\cdot p_k)(p_\nu\cdot p_k) - 2m_k^2(p_\tau\cdot p_\pi)(p_\nu\cdot p_\pi)\\ &+ m_k^2m_\pi^2(p_\tau\cdot p_\nu) - (p_\tau\cdot p_\nu)(p_k\cdot p_\pi)].\end{aligned} \tag{35}$$

The observables used in this study are obtained after the integration over the angular variables of the different contributions in Eq. (33). In the interference term of the new tensor contribution with the SM contribution [Eq. (34)], the term that arises from the symmetric (Q^μ) part of the standard current vanishes after this angular integration, and hence only the interference of the tensor with the vector form factor appears. The full differential decay rate may hence be written as

$$\begin{aligned}d\Gamma &\equiv \frac{1}{2m_\tau}[|\mathcal{A}_{\mathrm{SM}}|^2 + |\mathcal{A}_T|^2 + 2\mathrm{Re}(\mathcal{A}_{\mathrm{SM}}\mathcal{A}_T^\dagger)]dQ^2\\ &= d\Gamma_1 + d\Gamma_2 + d\Gamma_3,\end{aligned} \tag{36}$$

where

$$d\Gamma_1 = G_F^2\sin^2\theta_c S_{EW}\frac{m_\tau^3}{64\pi^3}\left(\frac{m_\tau^2-Q^2}{m_\tau^2}\right)^2\frac{P(Q^2)}{(Q^2)^{3/2}}\times\left\{P(Q^2)^2\left(\frac{2Q^2+m_\tau^2}{3m_\tau^2}\right)|F_V|^2 + \frac{1}{4}\frac{(m_K^2-m_\pi^2)^2}{Q^2}|F_S|^2\right\}dQ^2, \tag{37}$$

$$d\Gamma_2 = G_F^2\sin^2\theta_c S_{EW}\frac{m_\tau^3}{64\pi^3}\left(\frac{m_\tau^2-Q^2}{m_\tau^2}\right)^2\frac{P(Q^2)}{(Q^2)^{3/2}}\times\left\{P(Q^2)^2Q^2\left(\frac{Q^2+2m_\tau^2}{3m_\tau^2}\right)R_T^2|F_T|^2\right\}dQ^2 \tag{38}$$

and

$$d\Gamma_3 = G_F^2\sin^2\theta_c S_{EW}\frac{m_\tau^3}{64\pi^3}\left(\frac{m_\tau^2-Q^2}{m_\tau^2}\right)^2\frac{P(Q^2)}{(Q^2)^{3/2}}\times\left\{2P(Q^2)^2R_T|F_V||F_T|\frac{Q^2}{m_\tau}\cos(\delta_T(Q^2)-\delta_V(Q^2)+\phi)\right\}dQ^2. \tag{39}$$

For the conjugate tau decay mode, only the interference term in Eq. (39) will differ, having the opposite weak phase ϕ.

As mentioned above, the interference of the scalar contribution of the SM and the antisymmetric tensor contribution vanishes. This is similar to the vanishing of the scalar and the vector interference contribution in the SM itself. Note that the scalar term is even under parity, while both vector and tensor are odd under parity, resulting in only the vector-tensor interference being even under parity and hence surviving after the full (parity-even) phase space integration. In other words, once the angular integration is performed, terms that are odd in $\cos\beta$ vanish; however, the parity-even interference of the vector and tensor terms contributes to the decay rate even after this integration.

We wish to point out that after completion of this work we became aware of some earlier papers in which tensor interactions had been introduced in semileptonic tau decays, namely, Refs. [21], [22], and [23]. We notice several differences in our approach and these earlier papers. First, Ref. [21], claims that the tensor amplitude does not interfere with the SM amplitude; however, as shown in our calculations in this section, this is not true. In fact, it is exactly this interference that can possibly account for the *CP*-violating rate asymmetry. Reference [22] generates a very tiny *CP* asymmetry ($\approx 10^{-12}$) by second-order weak interactions, and Ref. [23] has given a numerical estimate of such an asymmetry, but in the context of a supersymmetry model, unlike our analytical formulas for a generic tensor interaction. Moreover, since this paper preceded the *BABAR* rate asymmetry measurement, they have not used the observables to constrain the NP parameters, which we attempt in the following section.

V. ESTIMATION OF THE NEW PHYSICS PARAMETER FROM OBSERVABLES

To estimate the parameters of the new tensor interaction term, we need to numerically compute the total decay rates for $\tau^+ \to K_s\pi^+\bar{\nu}$ and $\tau^- \to K_s\pi^-\nu$, using Eqs. (36)–(39). We can write the vector and the scalar form factors in terms of the magnitude and the strong phases as

$$F_v = |F_v|e^{i\delta_v(Q^2)}, \qquad \text{and} \qquad F_s = |F_s|e^{i\delta_s(Q^2)}. \tag{40}$$

Having expressed the form factors in terms of the combinations of Briet–Wigner forms of the various resonances, given in Eqs. (25) and (28), the strong phases of the scalar and the vector form factors can be simply extracted from these complex forms. We have used this vector form factor strong phase (Q^2 dependent) for our results below. As mentioned in Sec. III, Belle proposed two fit models for the Q^2 distribution of their data, which had comparable χ^2 values, both having the vector $K^*(892)$ resonance where the data peak as well as the scalar $K^*(800)$. In the region around 1.4 GeV, since the data lay much higher than the fitted curve, the inclusion of either $K^*(1410)$ or $K^*(1430)$ resulted in a significant improvement in the goodness of fit. Hence, in our analysis, we consider both the possibilities, $K^*(892) + K^*(1410) + K^*(800)$ and $K^*(892) + K^*(1430) + K^*(800)$, considering these two cases one at a time and naming them as case I and case II, respectively.

Note that a relative orbital angular momentum $l = 2$ of the $K\pi$ system would get contribution from a symmetric 2^+ state; however, since our NP amplitude consists of an antisymmetric tensor contribution, such a resonant[2] contribution to the $K\pi$ is not feasible. We take the tensor form factor to be real, and hence δ appearing in Eq. (15) in the presence of the new tensor interaction will be $\delta = \delta_v$, since $\delta_T = 0$. A similar tensor form factor had been introduced in the analysis of K_{e3} and $K_{\mu 3}$ data, and in fact the Particle Data Group [14] gives the constraint on the ratio of $|f_T/f_V|$ for these decays. However, with no such existing analysis from experiments nor any theoretical lattice estimates of the tensor hadronic form factor in the Q^2 range relevant for the tau decay being considered here, we assume the tensor form factor to be a constant for simplicity and determine the product of this constant and its coupling strength from the tau decay observables.

In the presence of this tensorial NSI, the experimentally measured *CP* asymmetry A_{CP}^{Exp} will be a result of the combination of the *CP* asymmetry arising from the K^0-$\bar{K}^0$ mixing and direct *CP* asymmetry appearing in the tau decay. If the *CP* asymmetry is a consequence only of the mixing contribution, then this asymmetry depending on the integrated decay times may be expressed as

$$A_{CP}^{K} = \frac{\int_{t_1}^{t_2} dt[\Gamma(K^0(t) \to \pi\pi) - \Gamma(\bar{K}^0(t) \to \pi\pi)]}{\int_{t_1}^{t_2} dt[\Gamma(K^0(t) \to \pi\pi) + \Gamma(\bar{K}^0(t)(t) \to \pi\pi)]}. \tag{41}$$

The K_s produced in the tau decay is observed through the $(\pi^+\pi^-)$ final state with $m_{\pi\pi} = m_K$. The τ^+ decays to K^0 at the time t_1 of tau decay (τ^- decays to $\bar{K}^0$ at t_1), and the time difference between the tau decay and the Kaon decay is of the order of the K_s lifetime (τ_S). Hence, in presence of both direct and indirect *CP* violation, we may express the observed decay-rate asymmetries as

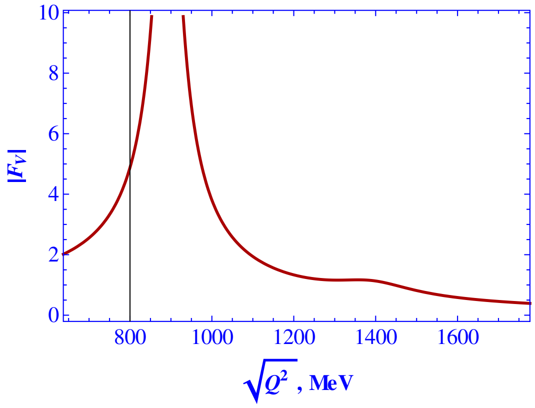
$$A_{CP}^{\text{Exp}} = \frac{\Gamma(\tau^+ \to K_s\pi^+\nu) \int_{t_1}^{t_2} dt\Gamma(K^0(t) \to \pi\pi) - \Gamma(\tau^- \to K_s\pi^-\nu) \int_{t_1}^{t_2} dt\Gamma(\bar{K}^0(t) \to \pi\pi)}{\Gamma(\tau^+ \to K_s\pi^+\nu) \int_{t_1}^{t_2} dt\Gamma(K^0(t) \to \pi\pi) + \Gamma(\tau^- \to K_s\pi^-\nu) \int_{t_1}^{t_2} dt\Gamma(\bar{K}^0(t) \to \pi\pi)}. \tag{42}$$

Interestingly, as shown below, this asymmetry can be factored into A_{CP}^K and A_{CP}^τ defined in Eqs. (15) and (41), where A_{CP}^τ is the *CP* violation due to the tensorial interaction. Defining

$$\Gamma^{\tau^\pm} \equiv \Gamma(\tau^\pm \to K_s\pi^\pm \overset{(-)}{\nu}_\tau); \qquad \Gamma^{\overset{(-)}{K^0}}(t) \equiv \Gamma(\overset{(-)}{K^0}(t) \to \pi\pi), \tag{43}$$

the difference of the decay rates in the numerator of A_{CP}^{Exp} in Eq. (42) can be written as

[2] A 2^- state cannot decay to $K\pi$ [for example, $K_2(1770)$ does not decay to $K\pi$ but to $K\pi\pi$], as expected by parity conservation.



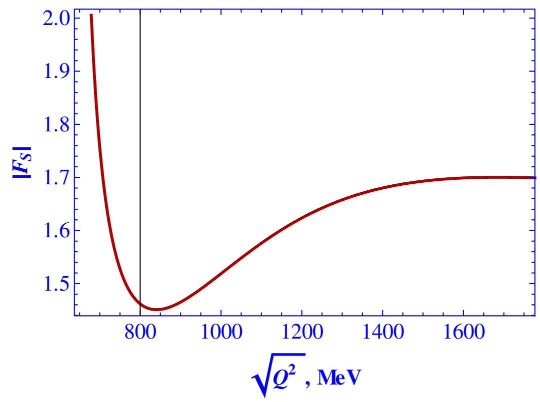


FIG. 2 (color online). The figures on the left and the right show the $\sqrt{Q^2}$ dependence of $|F_V|$ and $|F_S|$, respectively, where the parameters appearing in the combination of the Briet–Wigner forms for the resonances, $K^*(892)$, $K^*(1410)$, and $K^*(800)$, were used from the Belle fits for case I.

$$\begin{aligned}
&\Gamma(\tau^+ \to K_s\pi^+\nu)\int_{t_1}^{t_2} dt\Gamma(K^0(t)\to\pi\pi) - \Gamma(\tau^- \to K_s\pi^-\nu)\int_{t_1}^{t_2} dt\Gamma(\bar{K}^0(t)\to\pi\pi)\\
&= \Gamma^{\tau^+}\int_{t_1}^{t_2} dt\Gamma^{K^0}(t) - \Gamma^{\tau^-}\int_{t_1}^{t_2} dt\Gamma^{\bar{K}^0}(t)\\
&= 2\left\{\frac{\Gamma^{\tau^+}+\Gamma^{\tau^-}}{2}\frac{\int_{t_1}^{t_2} dt[\Gamma^{K^0}(t)-\Gamma^{\bar{K}^0}(t)]}{2} + \frac{\Gamma^{\tau^+}-\Gamma^{\tau^-}}{2}\frac{\int_{t_1}^{t_2} dt[\Gamma^{K^0}(t)+\Gamma^{\bar{K}^0}(t))]}{2}\right]\\
&= \frac{1}{2}\{\Gamma^{\tau^+}+\Gamma^{\tau^-}\}\int_{t_1}^{t_2} dt(\Gamma^{K^0}(t)+\Gamma^{\bar{K}^0}(t))[A_{CP}^K + A_{CP}^\tau]\\
&= \Gamma_{av}^{\tau^\pm}\int_{t_1}^{t_2} dt\{\Gamma^{K^0}(t)+\Gamma^{\bar{K}^0}(t)\}[A_{CP}^K + A_{CP}^\tau].
\end{aligned}$$

Similarly, the sum is

$$\begin{aligned}
&\Gamma(\tau^+ \to K_s\pi^+\nu)\int_{t_1}^{t_2} dt\Gamma(K^0(t)\to\pi\pi) + \Gamma(\tau^- \to K_s\pi^-\nu)\int_{t_1}^{t_2} dt\Gamma(\bar{K}^0(t)\to\pi\pi)\\
&= \Gamma^{\tau^+}\int_{t_1}^{t_2} dt\Gamma^{K^0}(t) + \Gamma^{\tau^-}\int_{t_1}^{t_2} dt\Gamma^{\bar{K}^0}(t)\\
&= \Gamma_{av}^{\tau^\pm}(1 + A_{CP}^K A_{CP}^\tau)\int_{t_1}^{t_2} dt\{\Gamma^{K^0}(t)+\Gamma^{\bar{K}^0}(t)\}.
\end{aligned}$$

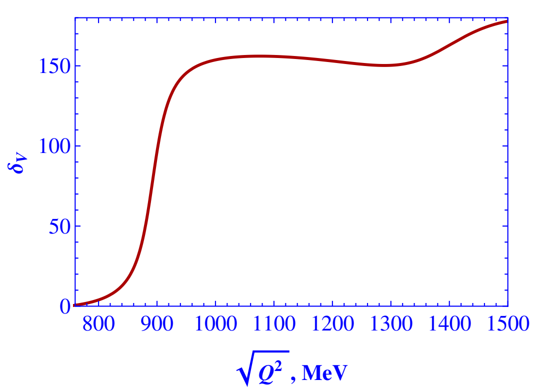


FIG. 3 (color online). The figure shows the $\sqrt{Q^2}$ dependence of δ_v as extracted from the complex form with contributions from the combination of the two vector resonances, $K^*(892)$ and $K^*(1410)$, with parameters from Belle fits for case I.

Therefore, the observed asymmetry and the branching ratio for this decay mode can be written as

$$A_{CP}^{\mathrm{Exp}} = \frac{A_{CP}^K + A_{CP}^\tau}{1 + A_{CP}^\tau A_{CP}^K} \tag{44}$$

and

$$BR = \frac{\Gamma^{\tau^+}+\bar{\Gamma}^{\tau^-}}{2\Gamma_{\mathrm{total}}}[1 + A_{CP}^\tau A_{CP}^K]\int_{t_1}^{t_2} dt\{\Gamma^{K^0}(t)+\Gamma^{\bar{K}^0}(t)\}.$$

Using the time dependence of the widths of K^0 and $\bar{K}^0$ to $\pi\pi$ for $t_1 \le \tau_S$ and $\tau_S \le t_2 \le \tau_L$, we can show that

TABLE I. Table showing the two solutions [(i) and (ii)] for the NSI parameters: the product of the ratio of the NSI coupling strength to the SM value and the tensor form factor, $R_T|F_T|$ and the cosine of the weak phase $\cos\phi$, allowed by the observables: the branching ratio and the CP asymmetry. Columns 4, 5, and 6 show the ratio of the contribution of the tensor mod-squared term, the interference terms involving the cosine of strong and weak phases, and that involving the sine of the phases, respectively, with respect to the SM contribution. The SM part uses the vector and scalar form factors corresponding to case I described in the text.

Sl.No	$R_T\|F_T\|$	$\cos\phi$	$\|\frac{T}{\mathrm{SM}}\|^2$	$\frac{\mathrm{Int(SM}*T)}{\mathrm{SM}^2}$ (cos term)	$\frac{\mathrm{Int(SM}*T)}{\mathrm{SM}^2}$ (sin term)
(i)	−0.303	−0.97	0.01837	0.09721	0.00753
(ii)	−1.945	−0.99	0.75853	0.63750	−0.02809

$$\int_{t_1}^{t_2} dt(\Gamma^K(t)+\Gamma^{\bar{K}}(t)) = \frac{\Gamma(K_s\to\pi\pi)}{\Gamma_{K_s}}\frac{|p|^2+|q|^2}{4|p|^2|q|^2} = BR(K_s\to\pi\pi), \tag{45}$$

where p, q, and ϵ are the standard K mixing parameters and we ignore terms of order ϵ^2 in the evaluation of the time integrals of the time-dependent decay rates of K^0 and $\bar{K}^0$ to $\pi\pi$, as was done in Ref. [7]. Hence, the branching ratio for tau decay to $K\pi\nu$ may be written in the following form:

$$BR = \frac{\Gamma_{av}^{\tau^\pm}}{\Gamma_{\mathrm{total}}}[BR(K_s\to\pi\pi)](1+A_{CP}^{\tau}A_{CP}^{K}). \tag{46}$$

A. Case I

Here, we have used $K^*(892)+K^*(1410)+K^*(800)$. For this case, Figs. 2 and 3 show the dependence of the form factors and the strong phase δ_V, respectively, on $\sqrt{Q^2}$. Substituting these Q^2-dependent form factors and the strong phase as well as the values of the masses, decay widths of the various resonances, and the branching ratio of the decay mode under consideration from PDG ([14,24]), we compute the Q^2-integrated results, $\Gamma^{\tau^\pm}$ for both the decay modes $\tau^\pm\to K_s\pi^\pm\overset{(-)}{\nu}_\tau$ within the kinematic limits $(m_K+m_\pi)^2$, and m_τ^2. This results in the average effective decay width dependent on the unknown parameters of the new interaction: the product of coupling constant R_T (ratio with respect to the SM) and the tensor form factor F_T, which is assumed to be a constant in this work and the CP-violating weak phase ϕ. Numerically, the effective widths in MeV are given by

$$\Gamma^{\tau^\pm} = 8.336\times10^{-12}+1.668\times10^{-12}(R_T|F_T|)^2+2.757\times10^{-12}R_T|F_T|\cos\phi\mp R_T|F_T|\sin(\phi)8.52674\times10^{-13}. \tag{47}$$

In the above equation, the first number is the integrated width for the SM, the second is the width corresponding to the mod squared of the tensor contribution, and the last two terms are from the interference of the SM and tensor parts and hence dependent on the strong phase (Q^2 dependent, which has been integrated out). The last three terms have been computed in terms of the unknown parameters of NP, the weak phase ϕ, and the product of the tensor coupling and form factor. From Eq. (15), we compute the second observable, the direct CP asymmetry, again in terms of the NP parameters to be

$$A_{CP}^{\tau} = \frac{2(R_T|F_T|)\sin\phi\times8.527\times10^{-13}}{\Gamma^{\tau^+}+\Gamma^{\tau^-}}, \tag{48}$$

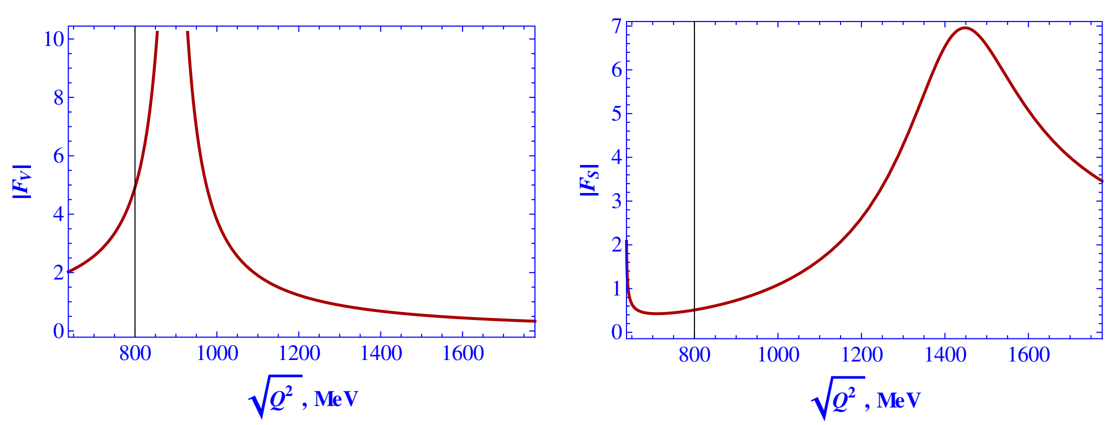


FIG. 4 (color online). The figure on the left shows the $\sqrt{Q^2}$ dependence of $|F_V|$, and that on the right shows the $\sqrt{Q^2}$ dependence of $|F_S|$, when these form factors include the Briet–Wigner contributions from the vector resonance $K^*(892)$ and the scalars, $K^*(1430)$ and $K^*(800)$, as used in the Belle fits for case II.

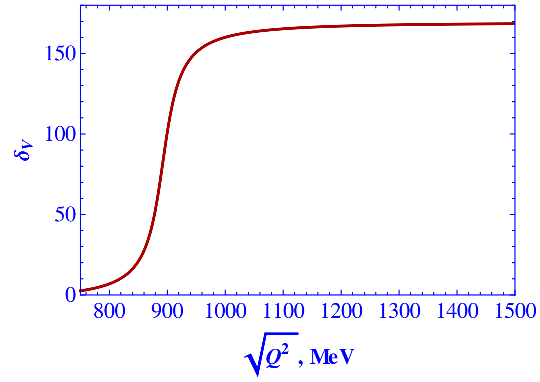


FIG. 5 (color online). The figure shows the $\sqrt{Q^2}$ dependence of δ_v corresponding to the vector form factor for the Belle fits of case II.

where the difference of the widths for τ^+ and τ^- is computed, using the difference of Γ_3 and $\bar{\Gamma}_3$ [after integrating the expressions in Eq. (39) corresponding to those for τ^+ and τ^-].

Using the average width, $(\Gamma^{\tau^+} + \Gamma^{\tau^-})/2$ evaluated from Eq. (47) and the *CP* asymmetry due to direct *CP* violation calculated in Eq. (48), in the expressions for the observed branching ratio and *CP* asymmetry derived earlier [Eqs. (46) and (44)], we can find the solutions for the unknowns $R_T|F_T|$ and $\cos\phi$. This results in two feasible solutions, displayed in Table I. Obviously, only the first solution is viable, as the NP contribution has to be much smaller than the SM contribution, since there is no glaring evidence of it, other than the unexpected direct *CP* violation seen. The smaller magnitude of the tensor mod squared and interference term relative to the SM contribution allows the Q^2 distribution of the SM alone to be reasonably consistent with the Belle data.

B. Case II

Here, the combination $K^*(892) + K^*(1430) + K^*(800)$ is used. Figures 4 and 5 show the dependence of the form factors and strong phases, on $\sqrt{Q^2}$ in this case. The complete decay rate is computed to be

$$\Gamma^{\tau^\pm} = 8.294\times 10^{-12} + 1.668\times 10^{-12}(R_T|F_T|)^2 + 2.633 \times 10^{-12} R_T|F_T|\cos\phi \mp 5.418\times 10^{-13} R_T|F_T|\sin\phi. \tag{49}$$

From Eq. (15), we get

$$A^{\tau}_{CP} = \frac{2R_T|F_T|\sin\phi\times 5.418\times 10^{-13}}{\Gamma^{\tau^+}+\Gamma^{\tau^-}}. \tag{50}$$

TABLE II. Table showing the allowed values of the NSI parameters, $R_T|F_T|$ and $\cos\phi$, as well as the ratio of the contribution of the tensor mod-squared term with respect to the SM contribution, as well as that of the interference contributions, corresponding to the SM form factors for case II.

Sl.No	$R_T\lvert F_T\rvert$	$\cos\phi$	$\lvert\frac{T}{SM}\rvert^2$	$\frac{\mathrm{Int(SM}*T)}{\mathrm{SM}^2}$ (cos term)	$\frac{\mathrm{Int(SM}*T)}{\mathrm{SM}^2}$ (sin term)
(i)	−0.213	−0.816	0.0091	0.05518	0.03909
(ii)	−3.333	0.999	2.2341	1.05703	0.04731

Similar to the first case, using the above two equations, we get two feasible solutions for $R_T|F_T|$ and $\cos\phi$ shown in Table II below, where again only the first solution is meaningful.

In the future, once the hadronic form factor for the tensor contribution is estimated theoretically, hopefully from lattice calculations or a fresh analysis of the larger data sample that may be available[3] is performed by the experimental groups, including the fits with a new tensorial contribution to the amplitude, then, with some handle on the tensor form factor (including its possible Q^2 dependence), the coupling strength as well as the weak phase of NP can be pinned down further. Note that the Q^2 dependence of the mod squared of the tensor amplitude is quite different from that of the other terms, which will enable its extraction from data.

VI. CONCLUSIONS

CP violation in the quark sector, observed through decays and mixing of K and B mesons, is consistent with its parametrization within the Standard Model. However, it fails to explain the large baryon asymmetry of the Universe and necessitates searches for *CP* violation beyond the Standard Model. Leptonic decays along with semileptonic decays of hadrons may offer a clean environment for searches of *CP*-violating new physics beyond the Standard Model.

The recent observation of a *CP*-rate asymmetry A_{CP} by *BABAR* [5] in the tau decay mode $\tau^\pm \to K_s\pi^\pm \overset{(-)}{\nu}_\tau$ seems to hint at some new physics, with the observed decay rate asymmetry being approximately 2.8 standard deviations away from the Standard Model predictions of an asymmetry that arises from $K^0 - \bar{K}^0$ mixing. The presence of various resonances in the vicinity of the decay hadrons invariant mass facilitates the availability of strong phases, while complex couplings in a new physics amplitude could provide the weak phase, enabling the possibility of a direct *CP* asymmetry. A charged scalar contribution can provide a

[3]The Belle collaboration, for example, has about three times larger statistics and has plans to repeat the $\tau \to K_S\pi\nu$ analysis. [25]

CP-violating asymmetry in the angular distribution but fails to produce an integrated rate asymmetry. However, this is achievable with a generic nonstandard tensorial interaction. We calculated the effective decay rate in the presence of the additional tensor interaction and in fact used the observed branching ratio and *CP* asymmetry to obtain the parameters of the new physics, the weak phase ϕ, and the product of tensorial coupling and the form factor.

ACKNOWLEDGMENTS

N. S. thanks D. Epifanov for discussions regarding the resonances used in the Q^2 distribution fits to data by the Belle collaboration. She thanks Swagato Banerjee and Randall Sobie for communication regarding the K_S lab angle constraints used by the *BABAR* collaboration in the asymmetry analysis. She is also grateful to N. G. Deshpande and Sandip Pakvasa for their comments.

[1] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1973).
[2] Y. S. Tsai, Phys. Rev. D **51**, 3172 (1995); Y.-s. Tsai, Phys. Lett. B **378**, 272 (1996); Y. S. Tsai, Nucl. Phys. B, Proc. Suppl. **55C**, 293 (1997).
[3] G. Bonvicini *et al.* (CLEO Collaboration), Phys. Rev. Lett. **88**, 111803 (2002).
[4] M. Bischofberger *et al.* (Belle Collaboration), Phys. Rev. Lett. **107**, 131801 (2011).
[5] J. P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **85**, 031102 (2012)**85**, 099904 (2012)].
[6] I. I. Bigi and A. I. Sanda, Phys. Lett. B **625**, 47 (2005).
[7] Y. Grossman and Y. Nir, J. High Energy Phys. 04 (2012) 002.
[8] P. del Amo Sanchez *et al.* (*BABAR* Collaboration), Phys. Rev. D **83**, 071103 (2011).
[9] B. R. Ko *et al.* (Belle Collaboration), Phys. Rev. Lett. **104**, 181602 (2010).
[10] H. Mendez *et al.* (CLEO Collaboration), Phys. Rev. D **81**, 052013 (2010).
[11] J. M. Link *et al.* (FOCUS Collaboration), Phys. Rev. Lett. **88**, 041602 (2002)**88**, 159903 (2002).
[12] I. I. Bigi, arXiv:1204.5817; arXiv:1210.2968.
[13] A. Pich, I. Boyko, D. Dedovich, and I. I. Bigi, Int. J. Mod. Phys. A **24** Suppl. 1, 715 (2009).
[14] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012)
[15] M. Finkemeier and E. Mirkes, Z. Phys. C **72**, 619 (1996); J. H. Kuhn and E. Mirkes, Z. Phys. C **56**, 661 (1992); **67**, 364 (1995).
[16] Y. Amhis *et al.* (Heavy Flavor Averaging Group Collaboration), arXiv:1207.1158.
[17] D. Epifanov *et al.* (Belle Collaboration), Phys. Lett. B **654**, 65 (2007).
[18] S. Paramesvaran (*BABAR* Collaboration), arXiv:0910.2884.
[19] M. Jamin, A. Pich, and J. Portoles, Phys. Lett. B **664**, 78 (2008).
[20] B. Moussallam, Eur. Phys. J. C **53**, 401 (2008).
[21] J. J. Godina Nava and G. Lopez Castro, Phys. Rev. D **52**, 2850 (1995).
[22] D. Delepine, G. Lopez Castro, and L.-T. Lopez Lozano, Phys. Rev. D **72**, 033009 (2005).
[23] D. Delepine, G. Faisl, S. Khalil, and G. L. Castro, Phys. Rev. D **74**, 056004 (2006).
[24] K. Nakamura *et al.* (Particle Data Group), J. Phys. G **37**, 075021 (2010).
[25] D. Epifanov (private communication).