No-Go Theorem for Nonstandard Explanations of the $\tau \to K_S \pi \nu_\tau$ *CP* Asymmetry

Vincenzo Cirigliano,[1] Andreas Crivellin,[2] and Martin Hoferichter[3]
[1] *Theoretical Division, Los Alamos National Laboratory, Los Alamos, New Mexico 87545, USA*
[2] *Paul Scherrer Institut, PSI, CH-5232 Villigen, Switzerland*
[3] *Institute for Nuclear Theory, University of Washington, Seattle, Washington 98195-1550, USA*

(Received 22 December 2017; published 6 April 2018)

The *CP* asymmetry in $\tau \to K_S \pi \nu_\tau$, as measured by the *BABAR* collaboration, differs from the standard model prediction by 2.8σ. Most nonstandard interactions do not allow for the required strong phase needed to produce a nonvanishing *CP* asymmetry, leaving only new tensor interactions as a possible mechanism. We demonstrate that, contrary to previous assumptions in the literature, the crucial interference between vector and tensor phases is suppressed by at least 2 orders of magnitude due to Watson's final-state-interaction theorem. Furthermore, we find that the strength of the relevant *CP*-violating tensor interaction is strongly constrained by bounds from the neutron electric dipole moment and D–$\bar{D}$ mixing. These observations together imply that it is extremely difficult to explain the current $\tau \to K_S \pi \nu_\tau$ measurement in terms of physics beyond the standard model originating in the ultraviolet.

DOI: 10.1103/PhysRevLett.120.141803

Introduction.—The presence of the baryon asymmetry in the Universe is one clear indication that there has to be physics beyond the standard model (SM) of particle physics [1], since *CP* violation within the SM, originating solely from the CKM matrix [2], is far too small to explain the observed asymmetry [3,4]. This need for additional *CP* violation renders *CP*-violating observables particularly interesting probes of beyond-the-SM (BSM) physics, with potentially profound implications for the SM and the physics of the early Universe.

CP violation was first observed in the neutral kaon system [5]. K^0 and $\bar{K}^0$ mix into the mass eigenstates K_S and K_L, which decay predominantly into 2π and 3π, respectively. However, K^0–$\bar{K}^0$ oscillations induce the *CP*-violating decays $K_L \to \pi\pi$ at a level of $\mathcal{O}(10^{-3})$. In addition to this indirect mechanism, the SM also permits direct *CP* violation, suppressed by another 3 orders of magnitude compared to indirect *CP* violation [6,7].

The focus of this article is the *CP* asymmetry in the decay width Γ of $\tau \to K_S \pi \nu_\tau$,

$$A_{CP}^{\tau} = \frac{\Gamma(\tau^+ \to \pi^+ K_S \bar{\nu}_\tau) - \Gamma(\tau^- \to \pi^- K_S \nu_\tau)}{\Gamma(\tau^+ \to \pi^+ K_S \bar{\nu}_\tau) + \Gamma(\tau^- \to \pi^- K_S \nu_\tau)}. \tag{1}$$

In the SM the dominant contribution again arises indirectly from K^0–$\bar{K}^0$ mixing [8], and the same statement holds for the analogous decays of D mesons [9],

$$A_{CP}^{D} = \frac{\Gamma(D^+ \to \pi^+ K_S) - \Gamma(D^- \to \pi^- K_S)}{\Gamma(D^+ \to \pi^+ K_S) + \Gamma(D^- \to \pi^- K_S)} = -4.1(9) \times 10^{-3}, \tag{2}$$

where the experimental number refers to the average of [10–13]; see [14]. In fact, the amplitude governing the indirect *CP* violation can be extracted very accurately from semileptonic kaon decays [15],

$$A_L = \frac{\Gamma(K_L \to \pi^- \ell^+ \nu_\ell) - \Gamma(K_L \to \pi^+ \ell^- \bar{\nu}_\ell)}{\Gamma(K_L \to \pi^- \ell^+ \nu_\ell) + \Gamma(K_L \to \pi^+ \ell^- \bar{\nu}_\ell)} = 3.32(6) \times 10^{-3}, \tag{3}$$

where $\ell = e, \mu$, and, neglecting small corrections from direct *CP* violation,

$$A_{CP}^{\tau,\mathrm{SM}} = -A_{CP}^{D,\mathrm{SM}} = A_L. \tag{4}$$

In each case, the signs follow from analyzing the quark content: for $K_L \to \pi^- \ell^+ \nu_\ell$ and $\tau^+ \to \pi^+ K_S \bar{\nu}_\tau$ the neutral kaon is produced as $K^0 = \bar{s}d$, while $D^+ \to \pi^+ K_S$ requires $\bar{K}^0 = \bar{d}s$ (in this case, however, there are corrections from the Cabibbo-suppressed decay mode [16]). Indeed, for the D-meson decay the corresponding prediction $A_{CP}^{D,\mathrm{SM}} = -3.32(6) \times 10^{-3}$ agrees well with the experimental result (2). In contrast, while earlier searches had not found evidence for *CP* violation [17,18], the latest result by the *BABAR* collaboration [19],

$$A_{CP}^{\tau,\mathrm{exp}} = -3.6(2.3)(1.1) \times 10^{-3}, \tag{5}$$

revealed a striking disagreement with the SM prediction. As pointed out in [20], since the intermediate K_S is reconstructed in terms of a final-state $\pi^+\pi^-$ pair with invariant mass around M_K and a decay time consistent with the K_S lifetime, the prediction (4) might be altered by the exact experimental conditions. However, the corresponding shift to $A_{CP}^{\tau,\mathrm{SM}} = 3.6(1)\times 10^{-3}$ even slightly increases the discrepancy to 2.8σ [19].

Optimistically, this tension could be considered a hint for BSM physics and it is natural to ask in a first step whether it is possible to account for the difference with nonstandard interactions. In general, for producing a nonvanishing CP asymmetry one needs the interference of two amplitudes,

$$\mathcal{A}_j = |\mathcal{A}_j| e^{i\delta_j^s} e^{i\delta_j^w}, \qquad j\in\{1,2\}, \tag{6}$$

with relative strong and weak phases $\delta^s = \delta_1^s - \delta_2^s$ and $\delta^w = \delta_1^w - \delta_2^w$. Both phases have to be nonvanishing, i.e.,

$$\begin{aligned} A_{CP} &\propto |\mathcal{A}_1 + \mathcal{A}_2|^2 - |\bar{\mathcal{A}}_1 + \bar{\mathcal{A}}_2|^2 \\ &= -4|\mathcal{A}_1||\mathcal{A}_2| \sin\delta^s \sin\delta^w. \end{aligned} \tag{7}$$

Here the $\bar{\mathcal{A}}_j$ denote the amplitudes with opposite weak phase. As argued in [21–23], due to the lack of a strong phase, this excludes an explanation using scalar operators, but new tensor interactions were found to be admissible. In this paper we demonstrate that this conclusion relies on erroneous assumptions for the πK tensor form factor, provide the corrected expression of the CP asymmetry in terms of the tensor Wilson coefficient, and study the consequences for a possible BSM explanation of (5).

Kinematics and conventions.—We define momenta according to

$$\tau(p_\tau) \to K_S(p_K) + \pi(p_\pi) + \nu_\tau(p_\nu), \tag{8}$$

with invariant mass $s = (p_K + p_\pi)^2$ of the πK system. In the following, we are only interested in the singly differential decay rate $d\Gamma/ds$, which is obtained after integrating over the remaining angular dependence of the three-body phase space.

For effective operators and form factors we largely follow the conventions of [24]. Because of parity conservation in the $K\to\pi$ matrix elements it suffices to consider the effective Lagrangian

$$\begin{aligned} \mathcal{L}_{su}^{\Delta S=1} = -\frac{G_F}{\sqrt{2}} V_{us} [& c_V (\bar{s}\gamma^\mu u)(\bar{\nu}\gamma_\mu \ell) + c_A (\bar{s}\gamma^\mu u)(\bar{\nu}\gamma_\mu\gamma_5 \ell) \\ & + c_S (\bar{s}u)(\bar{\nu}\ell) + i c_P (\bar{s}u)(\bar{\nu}\gamma_5\ell) \\ & + c_T (\bar{s}\sigma^{\mu\nu}u)(\bar{\nu}\sigma_{\mu\nu}(1+\gamma_5)\ell)] + \mathrm{H.c.}, \end{aligned} \tag{9}$$

where the Fermi constant G_F and the CKM element V_{us} have been factored out, and we have ignored all operators that vanish in the absence of right-handed neutrinos. The Wilson coefficients are defined at the weak scale in such a way that in the SM $c_V(M_W) = -c_A(M_W) = 1$ and all others equal to 0. The hadronic matrix elements are parametrized via form factors,

$$\begin{aligned} \langle \bar{K}^0(p_K)\pi^-(p_\pi)|\bar{s}\gamma^\mu u|0\rangle &= (p_K - p_\pi)^\mu f_+(s) \\ &\quad + (p_K + p_\pi)^\mu f_-(s), \\ \langle \bar{K}^0(p_K)\pi^-(p_\pi)|\bar{s}u|0\rangle &= \frac{M_K^2 - M_\pi^2}{m_s - m_u} f_0(s), \\ \langle \bar{K}^0(p_K)\pi^-(p_\pi)|\bar{s}\sigma^{\mu\nu}u|0\rangle &= i\frac{p_K^\mu p_\pi^\nu - p_K^\nu p_\pi^\mu}{M_K} B_T(s), \end{aligned} \tag{10}$$

where

$$f_-(s) = \frac{M_K^2 - M_\pi^2}{s}[f_0(s) - f_+(s)]. \tag{11}$$

Taking everything together, we obtain for the differential decay width for $\tau^- \to K_S\pi^-\nu_\tau$ (see also [21,25,26])

$$\begin{aligned} \frac{d\Gamma}{ds} = G_F^2 |V_{us}|^2 S_{\mathrm{EW}} & \frac{\lambda_{\pi K}^{1/2}(s)(m_\tau^2 - s)^2 (M_K^2 - M_\pi^2)^2}{1024\pi^3 m_\tau s^3} \\ \times & \left[\xi(s)\left(|V(s)|^2 + |A(s)|^2 + \frac{4(m_\tau^2 - s)^2}{9 s m_\tau^2}|T(s)|^2 \right) \right. \\ & \left. + |S(s)|^2 + |P(s)|^2 \right], \end{aligned} \tag{12}$$

where $\lambda_{\pi K}(s) = \lambda(s, M_\pi^2, M_K^2)$, $\lambda(a,b,c) = a^2 + b^2 + c^2 - 2(ab + ac + bc)$,

$$\xi(s) = \frac{(m_\tau^2 + 2s)\lambda_{\pi K}(s)}{3m_\tau^2 (M_K^2 - M_\pi^2)^2}, \tag{13}$$

$S_{\mathrm{EW}} = 1.0194$ [27–29] encodes the electroweak running down to m_τ, and

$$\begin{aligned} V(s) &= f_+(s) c_V - T(s), \\ A(s) &= f_+(s) c_A + T(s), \\ S(s) &= f_0(s)\left(c_V + \frac{s}{m_\tau(m_s - m_u)} c_S \right), \\ P(s) &= f_0(s)\left(c_A - i\frac{s}{m_\tau(m_s - m_u)} c_P \right), \\ T(s) &= \frac{3s}{m_\tau^2 + 2s}\frac{m_\tau}{M_K} c_T B_T(s). \end{aligned} \tag{14}$$

In the SM case $c_V = -c_A = 1$ (and $c_S = c_P = c_T = 0$) this reduces to

$$\begin{aligned} \left.\frac{d\Gamma}{ds}\right|_{\mathrm{SM}} = G_F^2 |V_{us}|^2 S_{\mathrm{EW}} & \frac{\lambda_{\pi K}^{1/2}(s)(m_\tau^2 - s)^2 (M_K^2 - M_\pi^2)^2}{512\pi^3 m_\tau s^3} \\ \times & [\xi(s)|f_+(s)|^2 + |f_0(s)|^2]. \end{aligned} \tag{15}$$

The general decomposition of the decay width (12) already shows why the scalar-vector (pseudoscalar–axial-vector) interference encoded in $S(s)$ $(P(s))$ cannot produce a CP asymmetry: the hadronic form factor $f_0(s)$ factorizes, so that the relative strong phase vanishes. This leaves the interference with the tensor operator in $V(s)$ and $A(s)$ as the only possible source for a strong phase. (The interference of vector and scalar operator could still contribute to the CP asymmetry due to long-distance QED corrections [30]. We estimate $|A_{CP}^{\tau,\mathrm{BSM}}| \lesssim 10^{-4}|\mathrm{Im}\, c_S|$, which is strongly suppressed due to the kinematic factor $\xi(s)$, see (15), the suppression of $f_0(s)$ compared to $f_+(s)$, and the QED factor $\mathcal{O}(\alpha/\pi)$. The branching ratio for $\tau \to K_S \pi \nu_\tau$ itself already excludes $|\mathrm{Im}\, c_S| \gtrsim 1$, so that even without further input the scalar contribution to A_{CP}^τ is of little phenomenological relevance.)

Upon neglecting direct CP violation in the SM, the total CP asymmetry can be written as [21]

$$A_{CP}^\tau = \frac{A_{CP}^{\tau,\mathrm{BSM}} + A_{CP}^{\tau,\mathrm{SM}}}{1 + A_{CP}^{\tau,\mathrm{BSM}} A_{CP}^{\tau,\mathrm{SM}}}, \tag{16}$$

where

$$A_{CP}^{\tau,\mathrm{BSM}} = \frac{\sin\delta_T^w |c_T|}{\Gamma_\tau \mathrm{BR}(\tau \to K_S \pi \nu_\tau)} \times \int_{s_{\pi K}}^{m_\tau^2} ds' \kappa(s') |f_+(s')||B_T(s')| \sin[\delta_+(s') - \delta_T(s')], \tag{17}$$

$s_{\pi K} = (M_\pi + M_K)^2$, and

$$\kappa(s) = G_F^2 |V_{us}|^2 S_{\mathrm{EW}} \frac{\lambda_{\pi K}^{3/2}(s)(m_\tau^2 - s)^2}{256\pi^3 m_\tau^2 M_K s^2}. \tag{18}$$

Moreover, δ_T^w denotes the phase of c_T relative to $c_V = -c_A = 1$, and $\delta_+(s)$, $\delta_T(s)$ are the phases of $f_+(s)$ and $B_T(s)$.

Hadronic form factors.—In [21] it was assumed that $B_T(s)$ is constant in such a way that only the phase of $f_+(s)$ remains and produces a sizable CP asymmetry via (17). That this assumption is incorrect can be argued in several ways. In the context of a vector-meson-dominance picture, the form factor $f_+(s)$ is dominated by the isospin-$I = 1/2$, spin-1 resonances $K^*(892)$ and the $K^*(1410)$, Breit-Wigner (BW) approximations of which are indeed used to parametrize the experimental decay width for the $\tau \to K_S \pi \nu_\tau$ process [31]. However, spin-1 resonances can be described equivalently by vector or antisymmetric tensor fields [32,33], so that the same resonances that contribute to $f_+(s)$ appear in $B_T(s)$ as well, most notably the $K^*(892)$.

Beyond the model approach, this conclusion can be derived by analyzing the unitarity relation for the form factors. For the vector current, such constraints from dispersion relations are frequently used to derive a parametrization of the form factor with good analytic properties [30,34–37]. In particular, πK intermediate states generate an imaginary part according to

$$\mathrm{Im}\, f_+(s) = \frac{\lambda_{\pi K}^{1/2}(s)}{s} f_+(s)[f_1^{1/2}(s)]^* \theta(s - s_{\pi K}), \tag{19}$$

where the πK partial waves $f_l^I(s)$ (with angular momentum l) obey the elastic unitarity relation

$$\mathrm{Im}\, f_l^I(s) = \frac{\lambda_{\pi K}^{1/2}(s)}{s} |f_l^I(s)|^2 \theta(s - s_{\pi K}), \tag{20}$$

and can thus be parametrized in terms of the πK phase shifts $\delta_l^I(s)$,

$$f_l^I(s) = \frac{s}{\lambda_{\pi K}^{1/2}(s)} e^{i\delta_l^I(s)} \sin\delta_l^I(s). \tag{21}$$

The unitarity relation for the form factor (19) then implies that, in the elastic region, the phase of $f_+(s)$ has to coincide with $\delta_1^{1/2}(s)$, a manifestation of Watson's final-state theorem [38]. In this way, $\delta_1^{1/2}(s)$ can be considered a model-independent implementation of the $K^*(892)$, which indeed decays almost exclusively to the $K\pi$ channel.

In fact, the tensor form factor obeys the exact same unitarity relation,

$$\mathrm{Im}\, B_T(s) = \frac{\lambda_{\pi K}^{1/2}(s)}{s} B_T(s)[f_1^{1/2}(s)]^* \theta(s - s_{\pi K}), \tag{22}$$

which simply reflects the fact that the $K^*(892)$ is equally well described by a vector or an antisymmetric tensor field. The relation (22) can be derived explicitly from the πK loop integral using standard Cutkosky rules. To actually construct a parametrization for $B_T(s)$ this relation is not sufficient because it does not determine the normalization. However, it shows that as long as πK states dominate the unitarity relation, the phases of $f_+(s)$ and $B_T(s)$ are identical, so that the corresponding CP asymmetry vanishes. This statement is exact as long as inelastic states, most notably $\pi\pi K$, are negligible. The next resonance, $K^*(1410)$, decays predominantly via $K^*(1410) \to K^*(892)\pi \to K\pi\pi$, and this indeed requires inelastic contributions that result in a nonvanishing CP asymmetry. Given the dominance of the $K^*(892)$ resonance, the cancellation in the elastic region strongly suppresses the amount of CP asymmetry that a tensor operator can produce.

Empirically, information on the form factor $f_+(s)$ can be derived from the $\tau \to \pi K_S \nu_\tau$ spectrum [31], which is strongly dominated by the $K^*(892)$ resonance. For the modulus of the form factors in (17) we can therefore ignore any inelastic corrections and use the elastic solution of the unitarity relation,

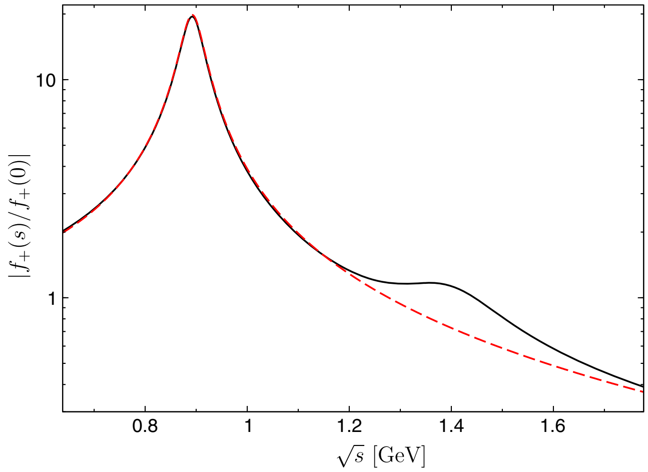

FIG. 1. $|f_+(s)/f_+(0)|$ from [31] (black solid line) in comparison to the Omnès factor (24) (red dashed line).

$$f_+(s) = f_+(0)\Omega(s), \qquad B_T(s) = B_T(0)\Omega(s), \tag{23}$$

in terms of the Omnès factor [39]

$$\Omega(s) = \exp\left(\frac{s}{\pi}\int_{s_{\pi K}}^{\infty}\frac{\delta(s')}{s'(s'-s)}\right). \tag{24}$$

The phase shift $\delta(s)$ can be identified with $\delta_1^{1/2}(s)$, and be approximated by a BW phase with parameters as determined in [31]. The resulting modulus is virtually indistinguishable from the experimental fit below the $K^*(1410)$ resonance, see Fig. 1, and the relative size of the two resonance peaks serves as an indication for the size of the inelastic effects.

The phase $\delta_+(s)$ cannot be directly taken from experiment, which is only sensitive to the modulus, and its extraction requires the use of a fit function that preserves the analytic structure of the form factor. This is not the case for the fit function used in [31] (a superposition of BW functions with complex coefficients), see Fig. 2, and indeed the corresponding phase cannot be physical because it does not vanish at threshold and violates Watson's theorem long before the $K^*(1410)$ can possibly have an effect. Still, the deviation between the phase found to be compatible with the spectrum [31] (blue dot-dashed line in Fig. 2), and the elastic phase (red dashed line in Fig. 2) provides a useful indication of the size of inelastic effects. As a simple estimate of the inelastic contribution $\delta_+^{\text{inel}}(s)$ to $\delta_+(s)$ we add the BW phase for $K^*(1410) \to K^*(892)\pi$ with a coefficient that allows for a similar phase motion in the vicinity of the $K^*(1410)$, to arrive at the band shown in Fig. 2 (consistent with more refined estimates along the lines of [30,34–37]).

Assuming that inelastic contributions in $\delta_T(s)$ are of similar size (but potentially opposite in sign), we take $\delta_+(s) - \delta_T(s) \sim 2\delta_+^{\text{inel}}(s)$. With this at hand, using

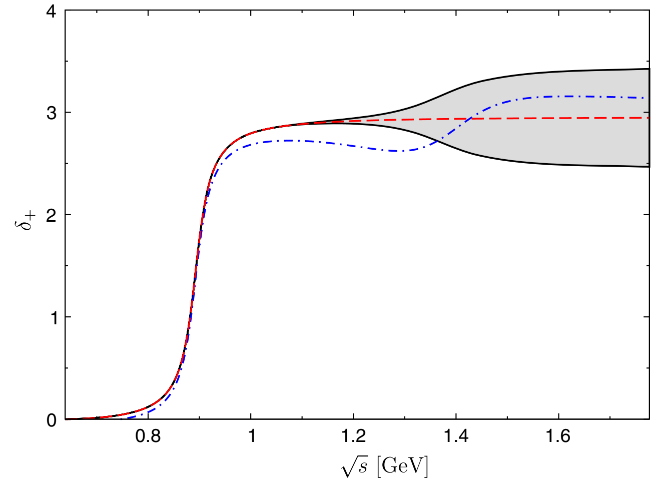

FIG. 2. δ_+ from a BW approximation for the $K^*(892)$ (red dashed line) in comparison to the phase from the experimental fit [31] (blue dot-dashed line). The band represents our estimate of inelastic effects; see the main text for details.

$\text{BR}(\tau \to K_S\pi\nu_\tau) = 4.04(13) \times 10^{-3}$ [31], $B_T(0)/f_+(0) = 0.676(27)$ from lattice QCD [40] (see [41] for an earlier calculation), and $f_+(0)|V_{us}| = 0.2165(4)$ as well as particle masses and couplings from [15] (see also [42]), we estimate for the CP asymmetry (17)

$$|A_{CP}^{\tau,\text{BSM}}| \lesssim 0.03|\text{Im}\, c_T|, \tag{25}$$

about 2 orders of magnitude less than for the maximum hadronic phase assumed in [21].

Limits on $\text{Im}\, c_T$.—To further appraise (25) we now turn to phenomenological constraints on $\text{Im}\, c_T$. By exploiting $SU(2)$ invariance of the weak interactions very strong limits follow from the electric dipole moment (EDM) of the neutron and D–$\bar{D}$ mixing, as we demonstrate in the following.

At a high scale $\Lambda \gg v$, where $v = 246$ GeV is the vacuum expectation value of the Higgs field, the tensor operator contributing to $\tau \to K_S\pi\nu_\tau$ arises from the following $SU(3) \times SU(2) \times U(1)$ gauge-invariant Lagrangian,

$$\mathcal{L}_T = C_{abcd}\bar{L}^i_{La}\sigma_{\mu\nu}e_{Rb}\epsilon^{ij}\bar{q}^j_{Lc}\sigma^{\mu\nu}u_{Rd} + \text{H.c.}, \tag{26}$$

where L_L and q_L denote the lepton and quark $SU(2)_L$ doublets, e_R and u_R are the charged lepton and up-quark $SU(2)_L$ singlets, i, j are $SU(2)_L$ indices, and a, b, c, d are generation indices. (In the notation of [43,44] this is the operator $Q^{(3)}_{\ell equ}$.) The tensor operator in (9) is generated from

$$\mathcal{L}_T = C_{3321}[(\bar{\nu}_\tau\sigma_{\mu\nu}R\tau)(\bar{s}\sigma^{\mu\nu}Ru) - V_{us}(\bar{\tau}\sigma_{\mu\nu}R\tau)(\bar{u}\sigma^{\mu\nu}Ru)] + \text{H.c.}, \tag{27}$$

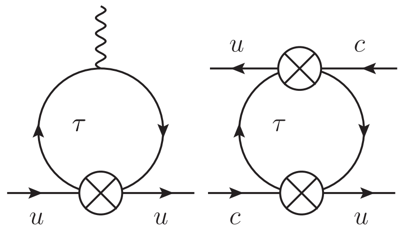

FIG. 3. Diagrammatic representation of the electromagnetic dipole operator contributing to the neutron EDM produced by inserting the $(\bar{\tau}\sigma_{\mu\nu}R\tau)(\bar{u}\sigma^{\mu\nu}Ru)$ operator (left), and the contribution to D–$\bar{D}$ mixing originating from the double insertion of the operator $(\bar{\tau}\sigma_{\mu\nu}R\tau)(\bar{c}\sigma^{\mu\nu}Ru)$ (right, the second permutation is omitted).

where $R = (1+\gamma_5)/2$, in the second line terms involving the charm and top quark have been neglected, and the Wilson coefficient C_{3321} is related to c_T by

$$C_{3321} = -\sqrt{2}G_F V_{us} c_T = -V_{us}\frac{c_T}{v^2}. \qquad (28)$$

In this way, $SU(2)$ symmetry relates the tensor operator relevant for $\tau \to K_S\pi\nu_\tau$ to a neutral current operator involving the τ and the up quark only. The renormalization group (RG) evolution [45] of this operator then produces an up-quark EDM $d_u(\mu)$,

$$\mathcal{L}_D = -\frac{i}{2}d_u(\mu)\bar{u}\sigma^{\mu\nu}\gamma_5 u F_{\mu\nu}, \qquad (29)$$

via the diagram shown in Fig. 3. Solving the RG following [46–48] we find

$$d_u(\mu) = \frac{em_\tau}{v^2}\frac{V_{us}^2}{\pi^2}\,\mathrm{Im}\,c_T(\mu)\log\frac{\Lambda}{\mu}$$
$$\simeq 3.0 \times \mathrm{Im}\,c_T(\mu)\log\frac{\Lambda}{\mu} \times 10^{-21}e\text{ cm}. \qquad (30)$$

Using the 90% C.L. bound $d_n = g_T^u(\mu)d_u(\mu) < 2.9 \times 10^{-26}e$ cm [49,50] and the recent lattice result [51] $g_T^u(\mu = 2\text{ GeV}) = -0.233(28)$ we obtain ($\mu_\tau = 2$ GeV)

$$|\mathrm{Im}\,c_T(\mu_\tau)| \leq \frac{4.4\times 10^{-5}}{\log\frac{\Lambda}{\mu_\tau}} \lesssim 10^{-5}, \qquad (31)$$

where the last inequality holds for $\Lambda \gtrsim 100$ GeV. This bound is based on the assumption that there are no other contributions to the neutron EDM canceling the effect of c_T. However, for values of $\mathrm{Im}\,c_T(\mu_\tau) \sim 0.1$ required to explain the tau CP asymmetry, the c_T contribution alone would predict a neutron EDM 4 orders of magnitude larger than the current bound, requiring an extraordinary cancellation at the level of one part in 10^4.

Such a cancellation could, in principle, occur with operators related to the flavor structure C_{3311} in (26), since the neutron EDM is sensitive to the combination $V_{ud}\mathrm{Im}\,c_T^{11} + V_{us}\mathrm{Im}\,c_T^{21}$, where $c_T^{21} = c_T$ and c_T^{11} is defined analogously to (28). However, yet another combination

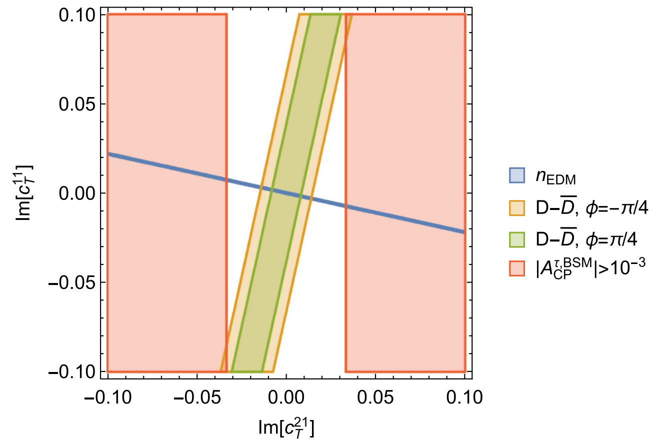

FIG. 4. Allowed regions in the $\mathrm{Im}\,c_T^{21}$–$\mathrm{Im}\,c_T^{11}$ plane from the neutron EDM and D–$\bar{D}$ mixing (for $\phi = \pm\pi/4$ and $\Lambda = 1$ TeV), compared to the favored region from the $\tau \to K_S\pi\nu_\tau$ CP asymmetry. The exclusion regions for $\phi = \pm\pi/4$ differ due to the asymmetric form of the fit result in [53].

appears in D–$\bar{D}$ mixing, which is very sensitive to the imaginary part of the Wilson coefficients (as for example defined in [52]),

$$C_2' = \frac{1}{2}C_3' = 4G_F^2\frac{m_\tau^2}{\pi^2}\log\frac{\Lambda}{\mu_\tau}V_{us}^2(V_{cd}c_T^{11} + V_{cs}c_T^{21})^2, \qquad (32)$$

where we have neglected the effect of external momenta, i.e., the mass of the charm quark. Using the global fit of [53] and assuming the phase of $V_{cd}c_T^{11} + V_{cs}c_T^{21}$ to be equal to $\phi = \pm\pi/4$ (in general, the constraint is diluted by $\sqrt{|\tan\phi|}$ and therefore disappears for $\phi = \pm\pi/2$), this leads to the situation depicted in Fig. 4. Since (32) requires the insertion of two effective operators, the leading contribution here is of dimension 8, while in an ultraviolet complete model there is in general already a dimension-6 contribution, making the bounds from D–$\bar{D}$ mixing even stronger than the one shown in Fig. 4. To evade all bounds, one would therefore not only have to cancel the c_T contribution to the neutron EDM at the level of 10^{-4}, but also tune the combination $V_{cd}c_T^{11} + V_{cs}c_T^{21}$ close to purely imaginary to evade the constraint from D–$\bar{D}$ mixing.

Conclusions.—In this Letter we examined nonstandard contributions to the CP asymmetry in $\tau \to K_S\pi\nu_\tau$. We find that at the dimension-6 level only the tensor operator can lead to direct CP violation, with negligible QED corrections from the scalar operator. However, the effect of the tensor operator is much smaller than previously estimated as a consequence of Watson's final-state-interaction theorem. Therefore, a very large imaginary part of the Wilson coefficient of the tensor operator would be required in order to account for the current tension between theory and experiment. In fact, we find in a model-independent analysis that this is in general in conflict with the bounds from the neutron EDM and D–$\bar{D}$ mixing, making a BSM explanation (realized above the electroweak breaking scale) highly improbable.

Nonetheless, a confirmation of the current *BABAR* measurement by Belle and/or Belle II would have intriguing consequences. In the absence of fine-tuning, it would point towards the existence of light BSM physics (realized below the electroweak breaking scale) so that our model-independent bounds could be evaded. We hope that the present analysis provides additional motivation to pursue such a measurement.

This research is supported by the U.S. Department of Energy, Office of Science, Office of Nuclear Physics, under Awards No. DE-AC52-06NA25396 and No. DE-FG02-00ER41132. A. C. is supported by an Ambizione grant of the Swiss National Science Foundation (Grant No. PZ00P2_154834).

[1] A. D. Sakharov, Pis'ma Zh. Eksp. Teor. Fiz. **5**, 32 (1967) [JETP Lett. **5**, 24 (1967)]; A. D. Sakharov, Usp. Fiz. Nauk **161**, 61 (1991); Sov. Phys. Usp. **34**, 392 (1991).
[2] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1973).
[3] A. G. Cohen, D. B. Kaplan, and A. E. Nelson, Annu. Rev. Nucl. Part. Sci. **43**, 27 (1993).
[4] A. Riotto and M. Trodden, Annu. Rev. Nucl. Part. Sci. **49**, 35 (1999).
[5] J. H. Christenson, J. W. Cronin, V. L. Fitch, and R. Turlay, Phys. Rev. Lett. **13**, 138 (1964).
[6] A. Alavi-Harati *et al.* (KTeV Collaboration), Phys. Rev. D **67**, 012005 (2003); **70**, 079904(E) (2004).
[7] J. R. Batley *et al.* (NA48 Collaboration), Phys. Lett. B **544**, 97 (2002).
[8] I. I. Bigi and A. I. Sanda, Phys. Lett. B **625**, 47 (2005).
[9] H. J. Lipkin and Z. z. Xing, Phys. Lett. B **450**, 405 (1999).
[10] J. M. Link *et al.* (FOCUS Collaboration), Phys. Rev. Lett. **88**, 041602 (2002); **88**, 159903(E) (2002).
[11] P. del Amo Sanchez *et al.* (*BABAR* Collaboration), Phys. Rev. D **83**, 071103 (2011).
[12] B. R. Ko *et al.* (Belle Collaboration), Phys. Rev. Lett. **109**, 021601 (2012); **109**, 119903(E) (2012).
[13] G. Bonvicini *et al.* (CLEO Collaboration), Phys. Rev. D **89**, 072002 (2014); **91**, 019903(E) (2015).
[14] Y. Amhis *et al.*, Eur. Phys. J. C **77**, 895 (2017).
[15] C. Patrignani *et al.* (Particle Data Group), Chin. Phys. C **40**, 100001 (2016).
[16] D. Wang, F. S. Yu, and H. N. Li, Phys. Rev. Lett. **119**, 181802 (2017).
[17] G. Bonvicini *et al.* (CLEO Collaboration), Phys. Rev. Lett. **88**, 111803 (2002).
[18] M. Bischofberger *et al.* (Belle Collaboration), Phys. Rev. Lett. **107**, 131801 (2011).
[19] J. P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **85**, 031102 (2012); **85**, 099904(E) (2012).
[20] Y. Grossman and Y. Nir, J. High Energy Phys. 04 (2012) 002.
[21] H. Z. Devi, L. Dhargyal, and N. Sinha, Phys. Rev. D **90**, 013016 (2014).
[22] L. Dhargyal, arXiv:1605.00629.
[23] L. Dhargyal, arXiv:1610.06293.
[24] M. Antonelli *et al.* (the FlaviaNet Kaon Working Group), arXiv:0801.1817.
[25] J. H. Kühn and E. Mirkes, Z. Phys. C **56**, 661 (1992); **67**, 364(E) (1995).
[26] M. Finkemeier and E. Mirkes, Z. Phys. C **72**, 619 (1996).
[27] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **56**, 22 (1986).
[28] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **61**, 1815 (1988).
[29] E. Braaten and C. S. Li, Phys. Rev. D **42**, 3888 (1990).
[30] M. Antonelli, V. Cirigliano, A. Lusiani, and E. Passemar, J. High Energy Phys. 10 (2013) 070.
[31] D. Epifanov *et al.* (Belle Collaboration), Phys. Lett. B **654**, 65 (2007).
[32] G. Ecker, J. Gasser, A. Pich, and E. de Rafael, Nucl. Phys. **B321**, 311 (1989).
[33] G. Ecker, J. Gasser, H. Leutwyler, A. Pich, and E. de Rafael, Phys. Lett. B **223**, 425 (1989).
[34] B. Moussallam, Eur. Phys. J. C **53**, 401 (2008).
[35] D. R. Boito, R. Escribano, and M. Jamin, Eur. Phys. J. C **59**, 821 (2009).
[36] D. R. Boito, R. Escribano, and M. Jamin, J. High Energy Phys. 09 (2010) 031.
[37] V. Bernard, D. R. Boito, and E. Passemar, Nucl. Phys. B, Proc. Suppl. **218**, 140 (2011).
[38] K. M. Watson, Phys. Rev. **95**, 228 (1954).
[39] R. Omnès, Nuovo Cimento **8**, 316 (1958).
[40] I. Baum, V. Lubicz, G. Martinelli, L. Orifici, and S. Simula, Phys. Rev. D **84**, 074503 (2011).
[41] D. Becirevic, V. Lubicz, G. Martinelli, and F. Mescia (SPQcdR Collaboration), Phys. Lett. B **501**, 98 (2001).
[42] M. Antonelli *et al.*, Phys. Rep. **494**, 197 (2010).
[43] W. Buchmüller and D. Wyler, Nucl. Phys. **B268**, 621 (1986).
[44] B. Grzadkowski, M. Iskrzyński, M. Misiak, and J. Rosiek, J. High Energy Phys. 10 (2010) 085.
[45] E. E. Jenkins, A. V. Manohar, and M. Trott, J. High Energy Phys. 01 (2014) 035.
[46] S. Bellucci, M. Lusignoli, and L. Maiani, Nucl. Phys. **B189**, 329 (1981).
[47] G. Buchalla, A. J. Buras, and M. K. Harlander, Nucl. Phys. **B337**, 313 (1990).
[48] V. Cirigliano, S. Davidson, and Y. Kuno, Phys. Lett. B **771**, 242 (2017).
[49] C. A. Baker *et al.*, Phys. Rev. Lett. **97**, 131801 (2006).
[50] J.M. Pendlebury *et al.*, Phys. Rev. D **92**, 092003 (2015).
[51] T. Bhattacharya, V. Cirigliano, R. Gupta, H. W. Lin, and B. Yoon, Phys. Rev. Lett. **115**, 212002 (2015).
[52] M. Bona *et al.* (UTfit Collaboration), J. High Energy Phys. 03 (2008) 049.
[53] A. J. Bevan *et al.* (UTfit Collaboration), J. High Energy Phys. 03 (2014) 123.