

## Tensor interactions and $\tau$ decays

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We study the effects of charged tensor weak currents on the strangeness-changing decays of the  $\tau$  lepton. First, we use the available information on the  $K_{e3}^+$  form factors to obtain  $B(\tau^- \rightarrow K^- \pi^0 \nu_\tau) \sim 10^{-4}$  when the  $K\pi$  system is produced in an antisymmetric tensor configuration. Then we propose a mechanism for the direct production of the  $K_2^*(1430)$  in  $\tau$  decays. Using the current upper limit on this decay we set a bound on the symmetric tensor interactions.

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### I. INTRODUCTION

The  $\tau$  lepton is the only charged lepton massive enough to decay into hadrons. This property serves to test interesting standard model (SM) predictions in a clean way. In particular, one can also study several properties of the charged vector [ $\rho(770)$ ,  $K^*(892)$ ] or axial vector [ $a_1(1260)$ ] mesons produced in  $\tau$  decays because their production mechanism is free from strong interaction complications.

Experimentally, the production mechanism for tensor mesons is of hadronic origin. For example, the  $a_2(1320)$  is observed in  $\pi p$  collisions or in  $J/\psi$  decays while the  $K_2^*(1430)$  is produced in  $Kp$  experiments [1]. These tensor mesons cannot be produced by a leptonic mechanism because of the  $V$  or  $A$  character of the electromagnetic and weak interactions in the standard model.

The aim of this paper is to estimate the effects of tensor interactions in strangeness-changing decays of the  $\tau$  lepton. For the purposes of this paper, it is convenient to start by introducing some terminology. We will call an *antisymmetric* tensor interaction to the low energy effective Lagrangian which involves the product of antisymmetric fermionic currents of the form  $J_{\{\mu\nu\}} \sim \bar{\psi}\sigma_{\mu\nu}\psi'$ , while the *symmetric* tensor interaction will involve the product of the currents  $J_{\{\mu\nu\}} \sim \bar{\psi}\Sigma_{\mu\nu}\psi'$ , where  $\Sigma_{\mu\nu}$  is a symmetric tensor involving Dirac  $\gamma$  matrices.

Let us first argue that on-shell tensor particles ( $J^P = 2^+$ ) cannot be produced by the SM interactions. The  $V - A$  structure of the weak charged currents at tree level does not allow the production of tensor mesons in  $\tau$  decays [the hadronic matrix element  $\langle T|\bar{q}\gamma_\mu(1 - \gamma_5)u|0\rangle$ ,  $q = d, s$  vanish identically]. At the one-loop level, a tensor vertex of the form  $\sigma_{\mu\nu}q^\nu$  ( $q^\nu$  the four-momentum transfer) can be induced in the SM by the first order QCD corrections to the vertex  $\bar{q}q'W^\pm$  [2]. However, the orthogonality conditions on the polarization tensor  $\varepsilon_{\{\mu\nu\}}$  describing the tensor particle, forbid the production of this particle when it is on shell [the only antisymmetric tensor that can be written out of  $q^\mu$  and  $\varepsilon_{\{\mu\nu\}}$  is  $\langle T|\bar{q}\sigma_{\mu\nu}u|0\rangle = \varepsilon_{\mu\nu\alpha\lambda}q_\beta(\varepsilon^{\{\alpha\beta\}}q^\lambda - \varepsilon^{\{\lambda\beta\}}q^\alpha)$  which vanishes because  $q_\mu\varepsilon^{\{\mu\nu\}} = 0$ ]. As an alternative, in this paper we will consider the  $|\Delta S| = 1$  flavor partner of the energy-momentum tensor as a possible mechanism

for the production of the  $K_2^*(1430)$  meson in  $\tau$  decays. Note that symmetric tensor interactions cannot be generated from radiative corrections to the  $V - A$  vertices [2].

The search for tensor currents dates from the beginning of the weak interaction theory. Recently, the existence of tensor fermionic interactions has been raised in several contexts. For instance, the presence of tensor antisymmetric interactions has been suggested in order to explain the apparent problems observed in (a) the  $\pi \rightarrow e\nu\gamma$  decay rate [3] and (b) the measurement of a nonzero tensor term in the decay  $K^+ \rightarrow \pi^0 e^+ \nu_e$  [4] (see also Ref. [5]). One also observes the presence of tensor fermionic currents in the context of the effective Lagrangian formulation for the low energy weak interactions [6]. Since the tensor interactions in  $K_{e3}^+$  decays is closely related to this work, let us first discuss it in more detail.

When the  $V - A$  requirement for the weak interactions is relaxed, the decay amplitude for the  $K \rightarrow \pi l^+ \nu_l$  ( $K_{l3}$ ) decays can be written as (Ref. [1], pp. 1530–1531)

$$\begin{aligned} \mathcal{M} \propto & f_+(q^2)[(P_K + P_\pi)_\mu \bar{l}\gamma_\mu(1 + \gamma_5)\nu] \\ & + f_-(q^2)[m_l \bar{l}(1 + \gamma_5)\nu] + 2m_K f_S \bar{l}(1 + \gamma_5)\nu \\ & + \frac{2f_T}{m_K} (P_K)_\lambda (P_\pi)_\mu \bar{l}\sigma_{\lambda\mu}(1 + \gamma_5)\nu, \end{aligned} \quad (1)$$

where  $q^2 = (P_K - P_\pi)^2$ . The form factors  $f_+$ ,  $f_-$  are associated with the vector hadronic current of the SM and, in the SU(3) limit, they are normalized such that  $f_+(0) = (1, 1/\sqrt{2})$  and  $f_-(0) = (0, 0)$  for  $(K_{l3}^0, K_{l3}^+)$ , respectively [7].  $f_S$ ,  $f_T$  are *scalar* and *tensor* form factors; their nonzero values would indicate signals for physics beyond the standard model.

Observe that the last term in Eq. (1) is just a convenient parametrization introduced to analyze the experiment. This amplitude could be associated, for example, to an effective Lagrangian that couples two antisymmetric currents:

$$\mathcal{L} = -\sqrt{2}G_F \tilde{f}_t \bar{u}\sigma_{\alpha\lambda}d' \left( \frac{q^\alpha q^\beta}{q^2} \right) \bar{l}_R \sigma_{\beta\lambda} \nu_L$$

as proposed by Chizhov in Ref. [5] (see the Appendix). In this case, the last term of Eq. (1) would arise from

the following parametrization of the hadronic matrix element:

$$\langle \pi | \bar{u} \sigma_{\mu\nu} s | K \rangle \sim (P_K)_\mu (P_\pi)_\nu - (P_\pi)_\mu (P_K)_\nu.$$

In the case of  $K_{e3}$  decays, the observables are not sensitive to  $f_-$ , and this allows in principle studying the effects of  $f_S$  and  $f_T$ . Surprisingly, the experimental results reported in [1] indicate  $|f_T/f_+(0)| = 0.38 \pm 0.11$  or equivalently,

$$f_T \equiv f_T(0) = 0.27 \pm 0.08 \quad (2)$$

for the  $K_{e3}^+$  decay, which is more than three standard deviations above zero [8]. In passing, let us mention that some other discrepancies between theory and experiment are observed in  $K$  semileptonic decays, namely, isospin breaking in the ratio  $f_+(0, K_{e3}^+)/f_+(0, K_{e3}^0)$  and isospin breaking in the slopes of the scalar form factors of  $K_{\mu 3}^0$  and  $K_{\mu 3}^+$  [9].

In order to clarify this possible experimental evidence for scalar or tensor antisymmetric interactions, it would be interesting to have new measurements of the form factors  $f_S$  and  $f_T$  at the  $\phi$  factory, where one expects the production of about  $10^{10}$  pairs of  $K^+K^-$ /year [10]. In the following we first will provide an estimate for the SM contribution to the  $\tau^- \rightarrow K\pi\nu_\tau$ . In Sec. III we assume the existence of the tensor antisymmetric interactions and consider its effects in  $\tau$  decays. In Sec. IV we will use the current upper limit on  $\tau^- \rightarrow K_2^-\nu_\tau$  to set a bound on symmetric tensor interactions.

## II. SM CONTRIBUTION TO $\tau^- \rightarrow K\pi\nu_\tau$

The SM contribution to the  $\tau \rightarrow K\pi\nu_\tau$  decay is given by the following amplitude:

$$\mathcal{M}_{\text{SM}} = \frac{G_F V_{us}}{\sqrt{2}} \bar{\nu} \gamma^\mu (1 - \gamma_5) \tau \langle K\pi | \bar{u} \gamma_\mu s | 0 \rangle, \quad (3)$$

where  $G_F$  denotes the Fermi constant and  $V_{us}$  the relevant Kobayashi-Maskawa matrix element.

The hadronic matrix element above can be written as  $\langle K(k)\pi(k') | \bar{u} \gamma_\mu s | 0 \rangle$

$$\begin{aligned} &= f_+(q^2)(k - k')_\mu + f_-(q^2)q_\mu \\ &= f_+(q^2) \left[ (k - k')_\mu - \frac{\Delta^2}{q^2} q_\mu \right] + \frac{\Delta^2}{q^2} f_0(q^2)q_\mu, \end{aligned} \quad (4)$$

where  $q = k + k'$  is the momentum transfer to the hadronic system,  $\Delta^2 \equiv m_K^2 - m_\pi^2$ , and  $f_+$ ,  $f_0$  are form factors associated to the  $J^P = 1^-, 0^+$  configuration of  $K\pi$ . Unlike  $K_{e3}$  decays where  $q^2 \leq (m_K - m_\pi)^2$ , in the  $\tau$  decay under consideration  $(m_K + m_\pi)^2 \leq q^2 \leq m_\tau^2$ . This allows the possibility of producing the  $K\pi$  system in a resonant way: for example, the  $K^*(892)$ ,  $K_0^*(1430)$ , and the  $K_2^*(1430)$  in the  $J^P = 1^-, 0^+$ , and  $2^+$  channels, respectively.

The decay rate corresponding to Eqs. (3) and (4) is given by

$$\Gamma_{\text{SM}}(\tau \rightarrow K\pi\nu_\tau) = \frac{G_F^2 m_\tau^5}{768\pi^3} |V_{us}|^2 I_{\text{SM}}, \quad (5)$$

where

$$\begin{aligned} I_{\text{SM}} &= \frac{1}{m_\tau^2} \int \frac{dq^2}{q^6} (m_\tau^2 - q^2)^2 \\ &\times \left\{ |f_+|^2 (m_\tau^2 + 2q^2) \lambda^{3/2}(q^2, m_K^2, m_\pi^2) + 3|f_0|^2 \Delta^4 m_\tau^2 \lambda^{1/2}(q^2, m_K^2, m_\pi^2) \right\} \end{aligned} \quad (6)$$

and  $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ .

We can estimate the decay rates by assuming a simple Breit-Wigner<sup>1</sup> form for the form factors in Eq. (6): namely,

$$f_i(q^2) = \frac{f_i(0)m_*^2}{m_*^2 - q^2 - im_*\Gamma_*}, \quad i = +, 0,$$

where  $m_*$ ,  $\Gamma_*$  are the resonant parameters of the  $K^*(892)$  or  $K_0^*(1430)$  when  $i = +$  or  $0$ , respectively. The form factors at  $q^2 = 0$  are taken from Ref. [11] to be  $f_+(0) = f_0(0) = 0.961 \pm 0.008$  and  $(0.982 \pm 0.008)/\sqrt{2}$  for the  $\bar{K}^0\pi^-$  and  $K^-\pi^0$  cases, respectively.

Using  $V_{us}$  and the  $\tau$  lifetime given in [1] we obtain

$$B(\tau^- \rightarrow K^-\pi^0\nu_\tau) = (3.35 \pm 0.11) \times 10^{-3}, \quad (7)$$

$$B(\tau^- \rightarrow \bar{K}^0\pi^-\nu_\tau) = (6.18 \pm 0.21) \times 10^{-3}, \quad (8)$$

where the quoted errors arise from the uncertainties in  $m_\tau$ ,  $\tau_\tau$ ,  $V_{us}$ , and  $f_i(0)$ . Adding both results we obtain  $B(\tau^- \rightarrow (K\pi)^-\nu_\tau) = (9.53 \pm 0.25) \times 10^{-3}$ , which compares reasonably well with the experimental value  $B(\tau^- \rightarrow K^{*-}(892)\nu_\tau) = (1.33 \pm 0.09)\%$  [12]. The numerical discrepancy between both results, if real, should be attributed to the simple Breit-Wigner form used to parametrize the form factors. Finally, the  $q^2$  dependence of  $f_0$  is not important because it contributes only 3% to Eq. (6).

## III. ANTISYMMETRIC TENSOR INTERACTIONS

Let us now consider the antisymmetric tensor contribution to the decay  $\tau^- \rightarrow K^-\pi^0\nu_\tau$ . If we assume  $e - \tau$  universality for tensor interactions, we can write the fol-

<sup>1</sup>A Breit-Wigner form with an energy-dependent width can be chosen as well.

lowing decay amplitude for the tensor contribution to this decay:

$$\mathcal{M} = \frac{G_F V_{us}}{\sqrt{2}} \frac{2f_T(q^2)}{m_K} k_\lambda k'_\mu \bar{\nu} \sigma^{\lambda\mu} (1 + \gamma_5) \tau, \quad (9)$$

where  $f_T$  is the  $q^2$ -dependent tensor form factor.

The decay rate corresponding to Eq. (9) can be written in the form (one can easily check that the tensor amplitude does not interfere with  $f_+$ ,  $f_0$  in the decay rate)

$$\Gamma(\tau \rightarrow K\pi\nu_\tau) = \frac{G_F^2 |V_{us}|^2 |f_T(0)|^2}{768\pi^3 m_\tau^2 m_K^2} I_{AS}, \quad (10)$$

where the integral  $I_{AS}$  is given by

$$I_{AS} = \int dq^2 (m_\tau^2 - q^2)^2 (2m_\tau^2 + q^2) \lambda^{3/2}(q^2, m_K^2, m_\pi^2) \times \left| \frac{f_T(q^2)}{f_T(0)} \right|^2. \quad (11)$$

Notice that the  $q^2$  distribution, Eq. (11), contains a factor  $(2m_\tau^2 + q^2)$  instead of  $(m_\tau^2 + 2q^2)$  obtained for the  $1^-$  channel, Eq. (6). This could help to isolate the tensor contribution in  $\tau \rightarrow K\pi\nu_\tau$  and have an independent measurement of  $f_T$ .

If we set  $f_T$  to a constant (see the following paragraph for the possibility of a  $q^2$ -dependent form factor), given in Eq. (2), and compute the branching ratio corresponding to Eq. (10), we obtain

$$B(\tau \rightarrow K^- \pi^0 \nu_\tau) = (1.8 \pm 1.0) \times 10^{-4}, \quad (12)$$

which lies one order of magnitude below the SM contribution, Eq. (7). Thus, if present, it seems possible to achieve a measurement of the nonresonant production of  $K\pi$  in the antisymmetric tensor configuration in a high statistic experiment such as a  $\tau$ -charm factory.

The decay rate given in Eq. (10) would receive an enhancement if the  $K\pi$  system were produced in a resonant way. An exotic candidate for this resonance would be, for example, the strange hybrid meson of the  $q\bar{q}g$  ( $J^P = 1^+$ ) family (this exotic meson can be described by an antisymmetric tensor  $\varepsilon_{[\mu\nu]}$ ) [13]. However, this contribution is inhibited because the hybrid mesons decay preferentially to final states containing excited  $q\bar{q}$  mesons [14].

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$$\Gamma(\tau \rightarrow K_2^* \nu_\tau) = \frac{G_F^2 |V_{us}|^2}{16\pi M^3} g_t^2 g_{K_2^*}^2 m^2 (2M^2 + 3m^2) (M^2 - m^2)^4, \quad (17)$$

where  $M$  denotes the mass of the  $\tau$  lepton.

Finally, if we compare Eqs. (16), (17), and the current upper limit on the  $\tau^- \rightarrow K_2^{*-} \nu_\tau$  decay [ $\Gamma^{\text{exp}}(\tau^- \rightarrow K_2^*(1430)\nu_\tau) < 6.7 \times 10^{-12}$  MeV [1]] we obtain the bound

#### IV. SYMMETRIC TENSOR INTERACTIONS

We propose that the current  $\times$  current form of the symmetric tensor interaction Lagrangian for the strangeness-changing  $\tau$  decays is given by

$$\mathcal{L}^{\Delta S=1} = \frac{G_F}{\sqrt{2}} V_{us} g_t (\bar{\nu} \Sigma_{\mu\nu} l) (\bar{u} \Sigma^{\mu\nu} s), \quad (13)$$

where the symmetric tensor  $\Sigma_{\mu\nu} \equiv i(\gamma_\mu \overleftrightarrow{\partial}_\nu + \gamma_\nu \overleftrightarrow{\partial}_\mu)$  involves first-order derivatives. The dimension of the effective tensor coupling  $g_t$  is  $(\text{mass})^{-2}$ . We have chosen the above Lagrangian in order to provide a mechanism responsible for the  $\tau \rightarrow K_2^* \nu_\tau$  decay. Although this choice does not exclude the existence of other tensor structures, we have used this Lagrangian for simplicity. It should be noted that it does not arise from radiative corrections to vertices with  $V - A$  currents [2].

From the above Lagrangian we get the following amplitude for  $\tau(p) \rightarrow K_2^*(k) \nu_\tau(p')$ :

$$\mathcal{M} = \frac{G_F}{\sqrt{2}} V_{us} g_t \bar{\nu} (\gamma_\mu Q_\nu + \gamma_\nu Q_\mu) \bar{l} \langle K_2^* | \bar{u} \Sigma^{\mu\nu} s | 0 \rangle, \quad (14)$$

where  $Q = p + p'$ . The hadronic matrix element in the previous equation can be parametrized as

$$\langle K_2^* | \bar{u} \Sigma^{\mu\nu} s | 0 \rangle = g_{K_2^*} m^3 \varepsilon^{\{\mu\nu\}}, \quad (15)$$

where  $m$  denotes the  $K_2^*$  mass and  $\varepsilon^{\{\mu\nu\}}$  its (symmetric) polarization tensor. With the above definition  $g_{K_2^*}$  becomes a dimensionless coupling.

Let us address a comment on the evaluation of the hadronic matrix element. Although it is not a popular idea, it has been suggested in the literature [15] that the *tensor meson dominance* of the energy-momentum operator can be assumed in order to give a single parameter description of the  $\pi\pi$  and  $\gamma\gamma$  decays of the  $J^P = 2^+$  meson  $f_2(1270)$ . Since the  $K_2^*(1430)$  meson and the  $\bar{u}\Sigma_{\mu\nu}s$  operator are flavor partners of the  $f_2(1270)$  and the energy-momentum tensor, respectively, we can assume the nonet symmetry in order to relate Eq. (15) and the corresponding annihilation amplitude of the  $f_2$ . The use of nonet symmetry gives

$$g_{K_2^*} = \frac{\sqrt{3}}{4} \left( \frac{m_{f_2}}{m} \right)^3 g_{f_2}, \quad (16)$$

where  $g_{f_2} = 0.103 \pm 0.011$  has been estimated by Terazawa [15] by using the  $f_2 \rightarrow \pi^+ \pi^-$  decay rate.

The decay rate corresponding to Eqs. (14) and (15) is

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$$g_t < 2.6 \times 10^{-6} \text{ MeV}^{-2}. \quad (18)$$

Equation (13) will also give a contribution to  $\tau^- \rightarrow K^- \pi^0 \nu_\tau$ . In this case, the hadronic matrix element can

be parametrized as [15]

$$\langle K\pi|\bar{u}\Sigma_{\mu\nu}s|0\rangle = \frac{g_{K_2^*}g_{K_2^*K\pi}m^2}{m^2 - q^2 - im\Gamma}(k - k')_\mu(k - k')_\nu, \quad (19)$$

where  $k(k')$  is the momentum of the  $K^-(\pi^0)$ , and  $\Gamma$  is the total width of the  $K_2^*$ . The strong coupling constant  $g_{K_2^*K\pi}$  can be determined from  $\Gamma^{\text{exp}}(K_2^{*-} \rightarrow K^-\pi^0) = (16.3 \pm 0.6) \text{ MeV}$  [1] and the expression

$$\Gamma(K_2^* \rightarrow K\pi) = \frac{2g_{K_2^*K\pi}^2|\vec{k}|^5}{5\pi m^4}. \quad (20)$$

We can compute the decay rate for  $\tau^- \rightarrow K_2^{*0}\nu \rightarrow K^-\pi^0\nu$  using the matrix element given in Eq. (19) and the upper bound given in Eq. (18). We obtain

$$B(\tau^- \rightarrow (K^-\pi^0)_{\text{symm}}\nu_\tau) < 9.9 \times 10^{-4}. \quad (21)$$

This upper limit is at the same level as the antisymmetric tensor contribution given in Eq. (12).

## V. CONCLUSIONS

We summarize our results. We have studied the effects of tensor interactions in strangeness-changing  $\tau$  decays. Using the information on the antisymmetric tensor interactions measured in  $K_{e3}^+$  decays we get a branching fraction for  $\tau^- \rightarrow [K^-\pi^0]_{\text{antisym}}\nu_\tau$  which is one order of magnitude below the SM contribution. On the other hand, we have proposed a mechanism for the direct production of the  $K_2^*(1430)$  in  $\tau$  decays. Using the current upper limit on the  $\tau \rightarrow K_2^*\nu_\tau$  decay mode we are able to set a bound on the intensity of the symmetric tensor interactions. Using this upper bound we have estimated  $B(\tau^- \rightarrow [K^-\pi^0]_{\text{symm}}\nu_\tau) < 9.9 \times 10^{-4}$ .

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## APPENDIX

The model of Chizhov [5] was proposed in order to account simultaneously for the destructive interference observed in  $\pi^+ \rightarrow e^+\nu_e\gamma$  [3] and a tensor term in  $K_{e3}^+$  reported in [4].

In Ref. [5], the SM is extended by introducing two Higgs doublets and two doublets of antisymmetric tensor

fields,  $T_{\mu\nu} = (T_{\mu\nu}^+, T_{\mu\nu}^0)$  and  $U_{\mu\nu} = (U_{\mu\nu}^0, U_{\mu\nu}^-)$ , having opposite hypercharges [ $Y(T) = -Y(U) = +1$ ] in order to cancel the anomalies. In this brief summary we use only the interactions of  $\nu_e, e, u, d, s$  fermions with the tensor fields that are relevant for the semileptonic processes.

By assuming quark-lepton universality of the coupling constant  $t$ , the  $SU(2)_L \times U(1)$  invariant interaction gives rise to the following interaction Lagrangian of the charged tensor fields with fermions [5]:

$$\mathcal{L} = \frac{t}{2}\{(\bar{\nu}_L\sigma^{\mu\nu}e_R + \bar{u}_L\sigma^{\mu\nu}d_R)T_{\mu\nu}^+ + \bar{u}_R\sigma^{\mu\nu}d_LU_{\mu\nu}^+ + \text{H. c.}\}, \quad (A1)$$

where  $d, u$  are interaction eigenstates. The tensor field  $U_{\mu\nu}$  couples only to quarks, because only left-handed neutrinos are present.

After spontaneous symmetry breaking the charged tensor fields become mixed and the corresponding matrix of propagators, in the  $q^2 \ll m^2, M^2$  approximation, is given by [5]

$$\mathcal{P} = \begin{pmatrix} (T^+T^-)_0 & (T^+U^-)_0 \\ (U^+T^-)_0 & (U^+U^-)_0 \end{pmatrix} = \frac{2i}{m^2 - M^2} \begin{pmatrix} \Pi(q) & -\mathbf{I} \\ -\mathbf{I} & M^2\Pi(q)/m^2 \end{pmatrix}, \quad (A2)$$

where  $m, M$  are the mass parameters associated with the vacuum expectation values (VEV's) of the two Higgs doublets,  $(XY)_0$  denote the corresponding Green functions of  $X$  and  $Y$ , and

$$\Pi_{\mu\nu\alpha\beta}(q) = \mathbf{I}_{\mu\nu\alpha\beta} - \frac{q_\mu q_\alpha g_{\nu\beta} - q_\mu q_\beta g_{\nu\alpha} - q_\nu q_\alpha g_{\mu\beta} + q_\nu q_\beta g_{\mu\alpha}}{q^2} \quad (A3)$$

with  $\mathbf{I}_{\mu\nu\alpha\beta} = \frac{1}{2}(g_{\mu\alpha}g_{\nu\beta} - g_{\mu\beta}g_{\nu\alpha})$ .

After diagonalization of  $\mathcal{P}$  and of the quark mass matrix, Eq. (A1) gives rise to the following four-fermion effective Lagrangian [5]:

$$\mathcal{L} = -\sqrt{2}G_F\bar{f}_t\bar{u}\sigma_{\mu\lambda}d\left(\frac{q^\mu q_\nu}{q^2}\right)\bar{l}_R\sigma^{\nu\lambda}\nu_L, \quad (A4)$$

where  $G_F\bar{f}_t/\sqrt{2} = t^2/(M^2 - m^2)$  and  $d = V_{ud}d' + V_{us}s'$ , with  $d', s'$  the quark mass eigenstates. Observe that the hadronic current in Eq. (A4) does not include a pseudotensor term; this will give rise to tensor contributions in  $\pi^+ \rightarrow e^+\nu_e\gamma$  and  $K^+ \rightarrow \pi^0e^+\nu_e$ , as required by experiment, but would leave unchanged the  $\pi^+ \rightarrow e^+\nu_e$  decay rate.

- [1] Particle Data Group, L. Montanet *et al.*, Phys. Rev. D **50**, 1173 (1994).
- [2] A. Halprin, B. W. Lee, and P. Sorba, Phys. Rev. D **14**, 2343 (1976); M. B. Gavela *et al.*, *ibid.* **22**, 2906 (1980); J. L. Díaz Cruz and G. López Castro, Mod. Phys. Lett. A **6**, 1605 (1991).
- [3] V. N. Bolotov *et al.*, Phys. Lett. B **243**, 308 (1990); A. A. Poblaguev, *ibid.* **238**, 108 (1990).

- [4] S. A. Akimenko *et al.*, Phys. Lett. B **259**, 225 (1991).
- [5] Some unconventional explanations for the origin of these tensor couplings have been raised in, for example, V. M. Belyaev and I. I. Kogan, Phys. Lett. B **280**, 238 (1992); M. V. Chizhov, Mod. Phys. Lett. A **8**, 2753 (1993).
- [6] W. Buchmüller and D. Wyler, Nucl. Phys. B **268**, 621 (1986); R. Martínez, M. A. Pérez, and J. J. Toscano, Phys. Lett. B **340**, 91 (1994).

- [7] L. M. Chounet, J.-M. Gaillard, and M. K. Gaillard, Phys. Rep. C **4**, 90 (1972).
- [8] Note also that the scalar form factor is found to be  $|f_S/f_+(0)| = 0.084 \pm 0.023$  (see Ref. [1]). Those averages are dominated by a single high statistic experiment (Ref. [4]). However, the effects of tensor or scalar interactions have not been observed [1] for the  $K_L \rightarrow \pi^\pm e^\mp \nu_e$  case.
- [9] G. López Castro and J. Pestieau, Mod. Phys. Lett. A **4**, 2237 (1989); G. Ordaz Hernández and G. López Castro, *ibid.* **5**, 755 (1990).
- [10] *The Daφne Physics Handbook*, edited by L. Maiani, G. Pancheri, and N. Paver (Servizio Documentazione dei Lab. Nazionali di Frascati, Frascati, 1992).
- [11] H. Leutwyler and M. Roos, Z. Phys. C **25**, 91 (1984).
- [12] R. Alemany, in *Proceedings of the 27th International Conference on High Energy Physics*, Glasgow, Scotland, 1994, edited by P. J. Bussey and J. G. Knowles (IOP, London, in press).
- [13] P. Hasenfratz, Phys. Lett. **95B**, 299 (1980); F. de Viron and J. Weyers, Nucl. Phys. **B185**, 391 (1981); D. Horn and J. Mandula, Phys. Rev. D **17**, 898 (1982).
- [14] N. Isgur, R. Kokoski, and J. Paton, Phys. Rev. Lett. **54**, 869 (1985).
- [15] H. Terazawa, Phys. Lett. B **246**, 503 (1990); M. Suzuki, Phys. Rev. D **47**, 1043 (1993).