New Tensor Particles

from $\pi^- \rightarrow e^- \bar{\nu} \gamma$ and $K^+ \rightarrow \pi^0 e^+ \nu$ Decays

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Abstract

The inserting of antisymmetric tensor fields into the standard electroweak theory may explain the recent experiments on $\pi^- \rightarrow e^- \bar{\nu} \gamma$ and $K^+ \rightarrow \pi^0 e^+ \nu$ decays. New intermediate particles can induce the destructive interference in the pion decay and the nonzero tensor and scalar form factors in the kaon decay.
1 Introduction

Recently the interest in possible tensor interactions in weak decays increased. This is connected with the last experiments on $\pi^- \rightarrow e^- \bar{\nu} \gamma$ and $K^+ \rightarrow \pi^0 e^+ \nu$ decays (see also Ref. 3). The experimentally obtained form factors cannot be explained in the framework of the standard electroweak theory. This led to numerous discussions on the possible existence of tensor terms in the effective Fermi interaction. In Ref. 4 results of the experiment on $\pi^- \rightarrow e^- \bar{\nu} \gamma$ decay were explained by introducing an additional tensor interaction in the Fermi Lagrangian

$$L_T = \sqrt{2} G_F V_{ud} f_T \bar{u}_R \sigma_{\mu\nu} d_L \cdot \bar{e}_R \sigma^{\mu\nu} \nu_L,$$

where $V_{ud}$ is an element of the Kobayashi-Maskawa mixing matrix and $\sigma_{\mu\nu} = \frac{i}{2} [\gamma_{\mu}, \gamma_{\nu}]$. The constant $f_T$ is estimated in the framework of the relativistic quark model: $f_T = -(4.2 \pm 1.3) \times 10^{-2}$. The independent calculation by QCD techniques leads to $f_T = -(1.4 \pm 0.4) \times 10^{-2}$. It was demonstrated in Ref. 6 that using (1) an agreement between all data of the previous experiments on $\pi^- \rightarrow e^- \bar{\nu} \gamma$ can be obtained. However, in Ref. 7 a mechanism for generating a (pseudo)scalar interaction

$$L_S = \sqrt{2} G_F V_{ud} f_0 (\bar{u}_R d_L) (\bar{e}_R \nu_L)$$

from (1) was presented leading to $|f_T| < 10^{-4}$, when the restrictions on $f_0$ from $\pi_e^2$ decay are accounted for:

$$|f_0| < 2.8 \times 10^{-6}.$$

In this paper we demonstrate that the simultaneous description of the results from the meson decays experiments is possible if the standard electroweak theory is extended by an additional Higgs doublet and two doublets of antisymmetric tensor particles: $(T^{+}_{\mu\nu} T^{\alpha}_{\mu\nu})$ and $(U^{\alpha}_{\mu\nu} U^{\mu\nu}_{\nu})$. This allows to explain both the destructive interference in the amplitude of the $\pi^- \rightarrow e^- \bar{\nu} \gamma$ decay and the appearance of the tensor and the scalar form factors in the $K^+ \rightarrow \pi^0 e^+ \nu$ decay.

2 Antisymmetric Tensor Fields

The antisymmetric tensor fields have been discussed in literature for a long time. However, now they are used only in the supersymmetric theory of gravitation. Two types of second rank antisymmetric tensor fields exist. The first ones play the role of the gauge fields and are described by a gauge-invariant Lagrangian. The second ones are matter fields and their Lagrangian is a conformal invariant one

$$L_o = \frac{1}{4} \partial_\lambda T_{\mu\nu} \partial^\lambda T^{\mu\nu} - \partial_\mu T^{\mu\lambda} \partial^\nu T_{\nu\lambda}.$$
Henceforth we use these fields. We want to note here that this Lagrangian is ghost free despite the widely spread belief that it is not. The six independent components of the field \( T_{\mu\nu} \) describe two massless particles: the vector one \( A_i = T_{\alpha i} \) and the pseudovector one \( B_i = \frac{1}{2} e_{ijk} T_{jk} \) \((i, j, k = 1, 2, 3)\). The particles must not be made massive by simple writing a mass term \( M^2 T_{\mu\nu} T^{\mu\nu} \) by hand, because the Hamiltonian will not be positively defined anymore. Therefore, the mechanism of the spontaneous symmetry breaking must be used to make the fields massive. For this purpose it is necessary to introduce an interaction of the tensor field with a massless particle, that in its intermediate state has a pole \( 1/q^2 \). In contrast to the usual Higgs mechanism of symmetry breaking in the case of a tensor field it is necessary to use both a scalar field with a nonzero vacuum expectation value \( \sqrt{2} \langle \varphi \rangle = M/2 g \), and a massless vector particle \( V_{\mu} \), giving the pole \( 1/q^2 \). Then the renormalizable gauge-invariant interaction has the form:

\[
\mathcal{L}_{SB} = 2g\varphi T_{\mu\nu} F_{\mu\nu} - 2g^2 \varphi^2 T_{\mu\nu} T^{\mu\nu}, \tag{4}
\]

where \( F_{\mu\nu} = \partial_{\mu} V_{\nu} - \partial_{\nu} V_{\mu} \). The tensor field propagator, taking into account the diagrams in Figs. 1a and 1b, is:

\[
\langle T(T_{\mu\nu} T_{\alpha\beta}) \rangle_0 = \frac{2i\Pi_{\mu\nu\alpha\beta}}{q^2 - M^2}, \tag{5}
\]

where

\[
\Pi_{\mu\nu\alpha\beta}(q) = \frac{1}{2} (g_{\mu\alpha} g_{\nu\beta} - g_{\mu\beta} g_{\nu\alpha}) - \frac{q_{\mu} q_{\alpha} q_{\nu} q_{\beta} - q_{\mu} q_{\beta} q_{\nu} q_{\alpha} + q_{\nu} q_{\alpha} q_{\mu} q_{\beta} + q_{\nu} q_{\beta} q_{\mu} q_{\alpha}}{q^2}.
\]

It is natural that the vector fields acquire mass by an exchange of a tensor particle (Fig. 1c).
The Lagrangian (3) is invariant under the axial transformation $T_{\mu\nu}^\pm \to \exp(\pm i\theta)T_{\mu\nu}^\pm$, where $T_{\mu\nu}^\pm$ are complex combinations of the real field $T_{\mu\nu}$ and its dual field $\tilde{T}_{\mu\nu} = \frac{1}{2}e_{\mu\alpha\beta}T^{\alpha\beta}$: $T_{\mu\nu}^\pm = (T_{\mu\nu} \pm i\tilde{T}_{\mu\nu})/\sqrt{2}$. The gauge fields can be easily introduced by rewriting the Lagrangian (3) in terms of the charged fields $T_{\mu\nu}^\pm$ and replacing the derivatives $\partial^\mu T_{\mu\nu}^\pm$ by the covariant derivatives: $D^\mu T_{\mu\nu}^\pm = (\partial^\mu \mp iA^\mu)T_{\mu\nu}^\pm$. The Lagrangian has the form:

$$L = -D^\mu T_{\mu\lambda}^+ D^\nu T_{\nu\lambda}^-.$$

(6)

In case where, besides the tensor and the axial gauge fields, a scalar and a vector gauge fields with an interaction of the type (4) are included, the theory will then be plagued with chiral anomalies (Fig. 2). Usually, in order to compensate these anomalies, additional fields with opposite transformations $U_{\mu\nu}^\pm \to \exp(\mp i\theta)U_{\mu\nu}^\pm$ are introduced. This concerns the scalar fields as well (Fig. 2b). Therefore, we can conclude that the extended electroweak model with tensor fields can be constructed anomaly free if corresponding pairs of tensor fields, like $T_{\mu\nu}$ and $U_{\mu\nu}$, are introduced. The Higgs sector of (pseudo)scalar fields must be duplicated as well.

**Fig. 2.**

3 The Extended Electroweak Model

We suppose a local $SU(2) \times U(1)$ electroweak symmetry with gauge fields $A_\mu$ and $B_\mu$. The fermion sector consists of several generations (noted by the index $a$) two component spinors: the lepton doublets $L_a = (\nu L_e L)$ and the lepton singlets $e_{Ra}$; the quark doublets $Q_a = (u_L d_L^t)_{a}$ and the quark singlets $u_{Ra}^t, d_{Ra}^t$. We suppose that neutrino is massless. Therefore, the mixing of the generations concerns only quark sector. Let $u_{La} = [S_u]_{ab} u_{La}$, $u_{Ra}^t = [T_u]_{ab} u_{Ra}^t$, $d_{La} = [S_d]_{ab} d_{La}$ and $d_{Ra}^t = [T_d]_{ab} d_{Ra}^t$, where the non-primed fields are the eigenstates of the mass matrix. Using the Kobayashi-Maskawa mixing matrix $V_{ab} = [S_u^t S_d]_{ab}$ and introducing another two matrices $U_{ab} = [S_u^t T_d]_{ab}$
and $W_{ab} = [T^a_u S_d]_{ab}$, a mixing either for the up-type quarks or for the down-type quarks will result.

The tensor interaction does not conserve chirality, hence it couples the left doublets and the right singlets of the fermion fields. In order to have $SU(2) \times U(1)$ invariant Yukawa Lagrangian for the spinor and the tensor fields, the tensor fields must be doublet ones. The requirement for anomaly free extended model can be fulfilled by introducing two doublets of scalar Higgs fields $H_1 = (H^+_1 H^0_1)$ and $H_2 = (H^+_2 H^0_2)$, and two doublets of tensor fields $T_{\mu\nu} = (T^{\mu+}_{\nu} T^{\nu+}_{\mu})$ and $U_{\mu\nu} = (U^{\mu}_{\nu} U^{\nu}_{\mu})$ with opposite hypercharges: $Y(T) = Y(H_1) = +1$, $Y(U) = Y(H_2) = -1$. Their interactions with the gauge fields are introduced by the covariant derivative $D_{\mu} = \partial_{\mu} - igT \cdot A_{\mu} - ig'Y/2 B_{\mu}$, where $g$ and $g'$ are the coupling constants, and $T$ and $Y$ are the generators of the $SU(2)$ and $U(1)$ groups, correspondingly. A triple interaction is also allowed. It is of the type of the first term of the Lagrangian (4) and it appears among the scalar fields, the tensor fields and the gauge fields strength tensors $F_{\mu\nu} = \partial_{\mu} B_{\nu} - \partial_{\nu} B_{\mu}$, $G_{\mu\nu} = \partial_{\mu} A_{\nu} - \partial_{\nu} A_{\mu} + g A_{\mu} \times A_{\nu}$:

$$\mathcal{L}_3 = (f_1 H^1_1 T^{\mu\nu} + f_2 H^2_2 U^{\mu\nu}) \cdot F_{\mu\nu} + (g_1 H^1_1 \tau T^{\mu\nu} + g_2 H^2_2 \tau U^{\mu\nu}) \cdot G_{\mu\nu} + \text{h.c.} \quad (7)$$

The requirement for the photon to be massless after the symmetry breaking, when $\langle H^0_1 \rangle = v_1$ and $\langle H^0_2 \rangle = v_2$, sets the following relations $f_1/g_1 = -f_2/g_2 = \tan \theta_W$.

In the following we will discuss the tensor fields interaction to the fermions. The $SU(2) \times U(1)$ invariant Lagrangian with Yukawa interactions is

$$\mathcal{L}_Y = \frac{1}{2} \left[ t^a \left( \bar{L}_a \sigma^{\mu\nu} e_{Ra} \right) + t^a \left( \bar{Q}_a \sigma^{\mu\nu} d'_{Ra} \right) \right] \left( \begin{array}{c} T^{\mu+}_{\nu} \\ \tau^0 \end{array} \right) + \frac{v_q}{2} \left( \bar{Q}_a \sigma^{\mu\nu} u'_{Ra} \right) \left( \begin{array}{c} U^o_{\mu\nu} \\ U^{0}_{\mu\nu} \end{array} \right) + \text{h.c.} \quad (8)$$

We want to note here that if the right neutrino does not exist, the doublet of the tensor field $U_{\mu\nu}$ interacts only with the quark fields.

After the spontaneous symmetry breaking, mixing of the fields may result. We are not going to discuss here the spontaneous symmetry breaking mechanism in detail. We only need to note that the mass parameters appear from the interaction (7). In case of mixing, the propagator for the charged tensor fields $T_{\mu\nu}^\pm$ and $U_{\mu\nu}^\pm$, will have the form:

$$\mathcal{P}(q) = \left( \begin{array}{cc} T(T^+T^-) & (T(T^+U^-))_o \\ (T(U^+T^-))_o & (T(U^+U^-))_o \end{array} \right) = \frac{2i }{\Delta} \left( \begin{array}{cc} (q^2 - m^2)\Pi(q) \\ \mu^2 (q^2 - M^2)\Pi(q) \end{array} \right), \quad \Delta = (q^2 - m^2)(q^2 - M^2) - \mu^4. \quad (9)$$

Now we can write the effective Lagrangian for the lepton-quark interaction due to the exchange of the charged tensor fields:

$$\mathcal{L}_{\text{eff}} = t^a \left( \bar{e}_{Ra} \sigma_{\mu\nu} \nu_{La} \right) \frac{q^\mu q^\nu}{\Delta} \bar{u}_b \sigma^{\nu\lambda} \left[ (q^2 - m^2) t^a_{b}(1 + \gamma^5)U_{bc} - \mu^2 u^b_{0}(1 - \gamma^5)W_{bc} \right] d_c + \text{h.c.} \quad (10)$$

The main requirement put on the interaction (10) is that it should not contribute into the $\pi_{\text{e}_2}$ decay. Therefore, the pseudotensor terms $\bar{u}_b \sigma_{\mu\nu} \gamma^5 d_c$, contributing to $\pi_{\text{e}_2}$ decay, must cancel. Assuming for simplicity the coupling constant universality of the tensor field interaction, i.e. $t^t = t^q = u^d$, and the mixing matrices $U$ and $W$ rather close to
unity, in the static approximation $q^2 \ll \mu^2$, $m^2$, $M^2$, there will be no contribution to the $\pi_{e2}$ decay, provided that $\mu^2 = m^2$.

We shall discuss the interaction of the light quarks $u$, $d$ and $s$. In order to avoid the flavor changing neutral currents we suppose $S_{u,d} = T_{u,d}$. Then the mixing of the left and the right quarks $d'_{L,R} = \cos \theta_C d_{L,R} + \sin \theta_C s_{L,R}$, $s'_{L,R} = -\sin \theta_C d_{L,R} + \cos \theta_C s_{L,R}$ is parametrized by one Cabibbo angle $\theta_C$. The effective Lagrangian (10) for the conserving strangeness case will have the form

$$L_{\Delta S=0} = -\frac{G_F \cos \theta_C}{\sqrt{2}} f_t \bar{e} \sigma_{\mu\lambda}(1 - \gamma^5)\nu \frac{q^\mu q_\nu}{q^2} \bar{u} \sigma^{\nu\lambda} d + \text{h.c.},$$

while that for the changing strangeness case will have the form

$$L_{\Delta S=1} = -\frac{G_F \sin \theta_C}{\sqrt{2}} f_t \bar{e} \sigma_{\mu\lambda}(1 - \gamma^5)\nu \frac{q^\mu q_\nu}{q^2} \bar{u} \sigma^{\nu\lambda} s + \text{h.c.},$$

where $f_t G_F / \sqrt{2} = \mu^2 / (M^2 - m^2)$. The mass matrix in (9) will have definite sign in case $m^2 \leq M^2$. This condition fixes the sign of $f_t$: $f_t \geq 0$. As it will be shown below, this leads to the right sign for the interference amplitude of the $\pi^- \rightarrow e^- \bar{\nu} \gamma$ decay.

### 4 The $K^+ \rightarrow \pi^0 e^+ \nu$ Decay

The most general form of the matrix element for the $K_{e3}^+$ decay is\(^{14}\)

$$M = \frac{G_F \sin \theta_C}{\sqrt{2}} \bar{\nu}(1 + \gamma^5)\left\{M_K f_S - \frac{1}{2} \left[ (P_K + P_\pi) \mu f_+ + (P_K - P_\pi) \mu f_- \right] \gamma^\mu + \frac{i f_T}{M_K} \sigma_{\mu\nu} P_K^{\mu} P_\pi^{\nu} \right\} e,$$ \hspace{1cm} (13)

where the form factors $f_+$, $f_-$, $f_S$ and $f_T$ are functions only of the square of the four-momentum transfer to leptons $q^2 = (P_K - P_\pi)^2$.

In the relativistic quark model the form factor $f_T$ can be obtained from the diagram in Fig. 3. Substituting $i g \bar{u} \gamma^5 u \cdot \pi^0$ and $i \sqrt{2} g \bar{u} \gamma^5 s \cdot K^+$ into the quark-meson vertices of the pion and the kaon correspondingly, and (12) into the quark-lepton vertex we obtain

$$M_T = i G_F \sin \theta_C f_t g^2 N_c \bar{\nu} \sigma_{\mu\lambda}(1 + \gamma^5) e \frac{q^\mu q_\nu}{q^2}$$

$$\times \int \frac{d^4l}{(2\pi)^4} S_p \left[ (\hat{l} - \hat{P_\pi} - M_u)^{-1} \gamma^5 (\hat{l} - M_u)^{-1} \gamma_5 (\hat{l} - \hat{P_K} - M_s)^{-1} \sigma^{\nu\lambda} \right],$$ \hspace{1cm} (14)

where $M_u = 340$ MeV and $M_s = 490$ MeV are the constituent masses of $u$ and $s$ quarks. We want to note that the one-loop integral is a convergent one and can be directly calculated. Finally, we obtain a relation between the form factor $f_T$ and the constant $f_t$:

$$f_T = \frac{\sqrt{2}}{4\pi^2 N_c} \frac{M_u a + (M_s - M_u) b}{M_K} g^2 f_t,$$ \hspace{1cm} (15)
where \( N_c = 3 \) is the colour factor; \( a \) and \( b \) are defined by

\[
16\pi^2 M_K^2 i \int \frac{d^4l}{(2\pi)^4} \{ (l - P_\pi)^2 - M_u^2 \}^{-1} \{ l^2 - M_u^2 \}^{-1} \{ (l - P_K)^2 - M_s^2 \}^{-1} = a(q^2),
\]

\[
16\pi^2 M_K^2 i \int \frac{d^4l}{(2\pi)^4} l\mu \{ (l - P_\pi)^2 - M_u^2 \}^{-1} \{ l^2 - M_u^2 \}^{-1} \{ (l - P_K)^2 - M_s^2 \}^{-1} = b(q^2) P_\mu^\pi + c(q^2) P_\pi^\mu,
\]

and they depend weakly on \( q^2 \). Their average values are \( a \approx 0.97 \) and \( b \approx 0.29 \).

With the help of the Goldberger-Treiman relation, the constant \( g \) can be expressed either through the pion decay constant \( F_\pi = 131 \text{ MeV} \) or the kaon decay constant \( F_K = 160 \text{ MeV} \): 

\[
g \approx \frac{2 M_u}{\sqrt{2} F_\pi} \approx \frac{(M_u + M_s)}{\sqrt{2} F_K} \approx 3.67.
\]

Using the experimental data \( f_T/f_+(0) = 0.38 \pm 0.11 \) and taking \( f_+(0) = \sqrt{2} \) we obtain \( | f^K_+ | = 0.49 \pm 0.14 \).

The tensor interaction (12) cannot explain the value \( f_S/f_+ = 0.084 \pm 0.023 \) and the relative phase angle value \( \phi_{st} = 0.00 \pm 0.52 \) of the scalar form factor in (13). The account of the radiative corrections to (12) leads to an insufficient value for \( \frac{f_S}{f_+} \approx 10^{-3} \) and \( \phi_{st} = \pi \). Hence in our model two Higgs doublets exist, an interaction due to charged Higgs fields is possible:

\[
\mathcal{L}_S = - \frac{G_F \sin \theta_C}{\sqrt{2}} f_s \bar{e}(1 - \gamma^5)\nu \cdot \bar{u}s + \text{h.c.}
\]

Here as well, in order to avoid the contribution to the \( \pi_{e2} \) decay, the quark pseudoscalar term must be absent. Then

\[
M_S = iG_F \sin \theta_C f_s g^2 N_c \bar{\nu}(1 - \gamma^5) e
\]

\[
\times \int \frac{d^4l}{(2\pi)^4} S \{ (\hat{l} - \hat{P}_\pi - M_u)^{-1}\gamma^5(\hat{l} - M_u)^{-1}\gamma^5(\hat{l} - \hat{P}_K - M_s)^{-1} \}.
\]
The leading diverging part of the integral in (18) has been evaluated using the quark loop for the kaon decay $K \rightarrow e\nu$

$$2(M_u + M_s)\sqrt{2}gN_c\int \frac{d^4l}{(2\pi)^4} \left[ l^2 - M_u^2 \right]^{-1} \left[ (l - p)^2 - M_s^2 \right]^{-1} = iF_K. \quad (19)$$

The following relation is obtained: $|f_S/f_+| = f_sM_s/M_K$. Then the coupling constant of the scalar interaction (17) is $|f_s| = 0.086 \pm 0.023$.

5 The $\pi^- \rightarrow e^-\tilde{\nu}\gamma$ Decay

We shall use the same parametrization of the matrix element as the one in Ref. 16

$$M = M_{IB} + M_{SD}, \quad (20)$$

where

$$M_{IB} = -i\frac{eG_F\cos\theta_C}{\sqrt{2}}F_\pi m_e \epsilon^\mu \left[ \left( \frac{k}{kq} - \frac{p}{pq} \right)^\mu - \frac{i\sigma^{\mu\nu} q^\nu}{2kq} \right] (1 - \gamma^5)\nu_e \quad (21)$$

is a QED correction to the $\pi \rightarrow e\nu$ decay (inner bremsstrahlung) and

$$M_{SD} = -\frac{eG_F\cos\theta_C}{\sqrt{2}M_\pi} \epsilon^\mu [F_V e_{\mu\rho\sigma} p^\rho q^\sigma - iF_A (pq \cdot g_{\mu\nu} - p_\mu q_\nu)] \bar{\epsilon}\gamma^\nu (1 - \gamma^5)\nu_e \quad (22)$$

is a structure-dependent amplitude parametrized by two form factors $F_V$ and $F_A$; $\epsilon^\mu$ is the photon polarization vector; $p$, $k$ and $q$ are the pion, electron and photon four-momenta respectively.

Fig. 4.
Using the new interaction (11), the additional to (20) matrix element can be calculated in the framework of the relativistic quark model (Fig. 4)

\[ M_T = e G_F \cos \theta_C f_t g N_c (e_d + e_u) \epsilon^\mu \bar{\epsilon} \sigma^{\gamma \beta} (1 - \gamma^5) \nu_e \frac{(p - q)_\gamma (p - q)^\mu}{(p - q)^2} \]

\[ \times \int \frac{d^4 l}{(2\pi)^4} S p \left[ (\hat{l} - \hat{q} - m)^{-1} \gamma_\mu (\hat{l} - m)^{-1} \gamma^5 (\hat{l} - \hat{p} - m)^{-1} \sigma_{\alpha \beta} \right] \]

(23)

here \( m \equiv M_u \approx M_d \) is the constituent mass of u and d quarks. The leading diverging part of the integral in (23) has been evaluated using the quark loop for the pion decay \( \pi \rightarrow e \nu \)

\[ 4m \sqrt{2} g N_c \int \frac{d^4 l}{(2\pi)^4} \left[ l^2 - m^2 \right]^{-1} \left[ (l - p)^2 - m^2 \right]^{-1} = i F_\pi . \]

Therefore, the additional amplitude reads

\[ M_T = \frac{e G_F \cos \theta_C F_T}{\sqrt{2}} \frac{1}{2} \left[ \epsilon^\mu q^\nu + \frac{(\epsilon \nu) q^\mu - (p q) \epsilon^\mu}{(p - q)^2} \right] \bar{\epsilon} \sigma_{\mu \nu} (1 - \gamma^5) \nu_e. \]

(24)

The constant \( F_T \) is expressed by \( f_t \) in a similar way as in Ref. 4:

\[ F_T = (e_d + e_u) \frac{F_\pi}{M} f_t. \]

(25)

Neglecting the electron mass, the decay amplitude with the leading interference term has the form:

\[ \frac{d^2 \Gamma}{dx d\lambda} = \frac{\alpha}{2\pi} \Gamma_{\pi \rightarrow e\nu} \left\{ IB(x, \lambda) + a_{SD}^2 \left[ (F_V + F_A)^2 SD^+(x, \lambda) + (F_V - F_A)^2 SD^-(x, \lambda) \right] - a_{SD} F_T I(x, \lambda) \right\} \]

(26)

where \( a_{SD} = M_{\pi}^2 / 2 F_\pi m_e \),

\[ IB(x, \lambda) = \frac{1 - \lambda (1 - x)^2 + 1}{x}, \quad SD^+(x, \lambda) = \lambda^2 (1 - x) x^3, \]

\[ SD^-(x, \lambda) = (1 - \lambda)^2 (1 - x) x^3, \quad I(x, \lambda) = (1 - \lambda) x^2. \]

In the pion frame the variables \( x \) and \( \lambda \) are \( x = 2 E_\gamma / M_\pi, \lambda = 2 E_e / M_\pi \sin^2 \theta_{ee} / 2 \).

The theoretical branching ratio, calculated within the standard \( V - A \) model for the kinematical region \( 0.3 < x < 1.0, 0.2 < \lambda < 1.0 \) is \( B^{th} = (2.41 \pm 0.07) \times 10^{-7} \). The discrepancy between the experimental total decay probability \( B^{exp} = (1.61 \pm 0.23) \times 10^{-7} \) and the calculated one \( B^{th} \) is due to the negative value of \( SD^- \). As far as the distributions of the spectra for \( I(x, \lambda) \) and \( SD^-(x, \lambda) \) terms are approximately similar ones, this discrepancy can be avoided if \( F_T = (3.72 \pm 1.20) \times 10^{-2} \). From (26) we can obtain \( f_t^\pi = 0.29 \pm 0.09 \).
6  Conclusions

From the calculation of the constants $f_t^\pi$ and $f_t^K$, on the basis of the meson decay experiments, it shows that they may be equal. Therefore, we can conclude that a universal effective interaction may exist

$$\mathcal{L}_T = -\frac{G_F}{\sqrt{2}} f_t \bar{\epsilon} \sigma_{\mu\lambda} (1 - \gamma^5)\nu \frac{q^\mu q_\nu}{q^2} \bar{u} \sigma^{\nu\lambda} d' + \text{h.c.} \quad (28)$$

with an average coupling constant $f_t = 0.39 \pm 0.18$, where $d' = \cos \theta_C d + \sin \theta_C s$. 

The absence of the quark pseudotensor terms in (28) allows to avoid the possible anomalously great contribution of the radiative corrections to the $\pi e^2$ decay. The effective scalar interaction

$$\mathcal{L}_S = -\frac{G_F}{\sqrt{2}} f_s \bar{\epsilon} (1 - \gamma^5)\nu \bar{u} d' + \text{h.c.} \quad (29)$$

does not contribute to the $\pi e^2$ decay as well. As far as (29) gives no contribution in $\pi^- \to e^- \bar{\nu} \gamma$ decay, the value of the scalar constant $| f_s | = 0.086 \pm 0.023$ is obtained only from the $K^+ \to \pi^0 e^+\nu$ decay. Unfortunately, we cannot say anything about the values of the Yukawa coupling constants of the tensor fields without using higher symmetries. Therefore, we will abstain from conclusions concerning the tensor particle masses.

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