
QUARK-MUONIC CURRENTS AND VIOLATION OF CP INVARIANCE

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In [1] we postulated, from cosmological considerations, the existence of quark-muonic currents whose interaction constant \( g_\alpha \) with the fractional-charge vector field \( a_{i \alpha} \) of the order of \( 10^{-3} \). In this note we consider a hypothesis in which the violation of CP invariance in \( K_0 \to L \) decay (see [2]) is ascribed to the difference between the phase constants \( g_\alpha \exp(i \phi) \) for the ordinary and strange quarks.

We assume interaction Lagrangians with maximum phase difference for ordinary and strange quarks \(^*)\) and with P-parity conservation:

\[
\begin{align*}
L &= \sum_{\alpha, \mu} g_\alpha \left[ (\bar{\Psi}_{-q} a_{i \mu} \gamma^i \Psi_{\mu}) + \text{h.c.} \right], \\
L &= i \sum_{\alpha, \mu} g_\alpha \lambda \left[ (\bar{\Psi}_{-\lambda} a_{i \mu} \gamma^i \Psi_{\mu}) - \text{h.c.} \right],
\end{align*}
\]

\( \alpha, \mu, \lambda, \) and \( \mu \) are the indices of the electric charge and take on the values \( q = -1/3, +2/3; \lambda = -1/3; \mu = 0, 0; \alpha = -2/3, +1/3, +4/3. \)

Generally speaking, the constant \( g_\alpha \) can depend on the index \( \alpha \), but we shall not consider these variants.

Figure 1 shows the main diagrams for the transformation of \( K_0 = \bar{\Lambda}n \) into \( \bar{K}_0 = \Lambda\bar{n} \). The

\[
\begin{align*}
\text{Fig. 1}
\end{align*}
\]

matrix element of the transition, \( V_{12} = (K_0 | V | \bar{K}_0) \), is complex:

\[
V_{12} \sim 2 g_W^2 \delta_{\lambda \lambda} + i g_W \delta_{\lambda \lambda} \delta_{\alpha \lambda};
\]

\( g_W \) is the constant for the interaction of the weak current with the W boson,
$$\frac{4\pi G_W^2}{m_W^2} = \frac{G}{\sqrt{2}} = \frac{10^{-5}}{\sqrt{2} m_p^2},$$

where $m_W$ is the W-boson mass. In the expression for $V_{12}$ we neglected the possible differences between the masses $m_{\mu}$ and $m_{W}$, which enter in diverging expressions. The eigenfunction of the mass operator are proportional to

$$K_0 V_{12} = \tilde{K}_0 |V_{12}|,$$

and in particular

$$K_L = \cos \nu \cdot K_2 + i \sin \nu \cdot K_1,$$

where

$$\nu = \frac{1}{2} \frac{\text{Im} V_{12}}{\text{Re} V_{12}} = \frac{1}{4} \frac{\delta_{AA}}{\delta_{WW}}.$$

The difference between $K_L$ and $K_0$ determines completely the amplitude of the decay $K_L \rightarrow \pi^+ \pi^-$, since the "direct" decay $K_L \rightarrow \pi^+ \pi^-$ is forbidden by P-parity conservation in the a-interaction. Therefore

$$\nu = \frac{\Lambda(K_L \rightarrow \pi^+ \pi^-)}{\Lambda(K_s \rightarrow \pi^+ \pi^-)} = 2 \cdot 10^{-3}.$$  

Using the value $g^2_{\pi} = (137^{-2})$ from [1], we get $g^2_{W} = (137)^{-2}$ and

$$m_{\pi} \sim 1.0 m_{H} \sim 137 (m_{H} / 2) = 137^2 m_c.$$  

$m_{\pi}$ remains unknown.

The quark currents, as well as the quark-electronic currents which are possible in principle, should change the ratio of the yields of the two channels of $\pi^+$-meson decay

$$R = \frac{\Lambda(\pi^+ \rightarrow e^+ \nu)}{\Lambda(\pi^+ \rightarrow \mu^+ \nu)}.$$  

The experimental value is $R = (1.24 \pm 0.05) \times 10^{-4}$ (see [3]) and agrees within the limits of measurement accuracy both with the theoretical value of $R_W$ given by the V - A theory with electromagnetic corrections (see [4]), and with the possible value of $R_{W+a}$ measured as a result of the presence of quark currents $[(R_{W+a} - R_{W})/R_{W}] \sim 0.01 (m_{W}^2/m_{\pi}^2)$. We use the formula

$$R_{W} = \frac{m_{e}}{f_{\pi}} \left( \frac{m_{\pi}^2}{m_{e}^2} - \frac{m_{\pi}^2}{m_{\mu}^2} \right)^2 (1.21 \cdot 10^{-4}).$$

$\gamma$ takes into account the electromagnetic correction to the nonrelativistic approximation, as an effect that is due to the attraction of the charged particles in the neutral channel and depends on their relative velocity $v$. 

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\[ f = \frac{|\Psi(0)|}{|\Psi(\infty)|} = \left( \frac{2 \pi e^2}{\nu \left( 1 - e^{2 \pi e^2/\nu} \right)} \right)^{1/2}, \quad \frac{f^2}{f^2_{\mu}} = 0.945. \]

The change in the decay amplitude due to the quark currents is

\[ A_a = -A_w \frac{e^2_s}{2 g_w^2} \frac{m_w^2}{m_a^2} \frac{m_\pi}{m_\mu} e^{i \phi} \]

(see Fig. 2). The uncertainty in the phase is here a reflection of the uncertainty in the relative phases in the expression (2) for the weak current. Putting \( \varphi_1 = \varphi_2 \) we get

\[ R_{W +} = R_w (1 - \frac{2 m_\mu}{m_\pi} \frac{m_{\pi^0}}{m_\mu} \nu)^{-2}. \]

We note that the quark-electronic currents of comparable magnitude (for \( m_a \sim m_w \) and \( \varphi_2 \neq \pi/2 \)) would change \( R \) by several times ten per cent, which is completely excluded by the experiment.

In the case of \( K^+ \) decay, a similar effect of the change in the yield ratio of the two channels is very small (\( \sim \nu^2 \)) when the phases \( \varphi \) coincide, owing to the fact that the a-amplitude is imaginary. An effect on the order of \( \nu (m_w^2/m_a^2) \) (i.e., \( \sim 1 - 0.1\% \)) must be expected in the expression for the probability of \( K \) capture of \( \mu^- \) in hydrogen and \( \text{He}^3 \), and in effects of transverse polarization in three-particle decay of \( K_L \).

![Fig. 2](image)

Fig. 2

Particular interest attaches to effects of violation of C symmetry of the partial probabilities in decays with \( P \)-parity conservation and change of strangeness, for example, the deviation of the ratio of the partial probabilities

\[ \frac{K^+ \rightarrow \pi^+ + \pi^+ + \pi^-}{K^- \rightarrow \pi^- + \pi^- + \pi^+}, \quad \frac{\Sigma^+ \rightarrow N + \pi^+}{\Sigma^- \rightarrow N + \pi^-} \]

from unity (the effect of S. Okubo). These effects are also of the order of \( \nu (m_w^2/m_a^2) \), but since they depend on the phase differences \( \varphi_{1,2} \sim \left( A_w/A_a \right) \) for different values of the isospin \( I \) and on the phases of the strong interaction, their numerical value should be much lower than 0.1%.

As indicated by L. B. Okun' in the course of a discussion, processes of interest from the point of view of checking the theory are those in which pairs of charged mesons are produced, for example \( K^+ \rightarrow \pi^+ + \mu^+ + \mu^- \) (relative yield \( \sim \nu^2 \)). The processes \( K_L \rightarrow \mu^+ + \mu^- \) and \( K_L \rightarrow \mu^+ + \mu^- + \pi^0 \) are strongly forbidden, but the processes \( K_L \rightarrow \mu^+ + \mu^- + \gamma \) and \( \Sigma^+ \rightarrow \Sigma^+ + \mu^+ + \mu^- \) are possible (all are \( \sim \nu^2 \)).

The author takes this opportunity to thank L. B. Okun' for a discussion and advice.


*) The equivalent form of the theory - introduction of complex phases in the expression for the weak current

\[ j_w = \pi O \nu + \pi O \mu \alpha \phi_1 + \pi O \rho \phi_2 + \frac{\pi W L}{\rho W} \phi_3 + \ldots \]

\[ \phi_3 = \phi_2 - \pi \/ 2. \]

FRAGMENT ANGULAR DISTRIBUTION IN THE FISSION OF Th^{232} BY 1.6-MeV NEUTRONS

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Henkel and Brolley [1], in a study of the angular anisotropy of the fission of Th^{232} by neutrons, were in fact the first to obtain weighty evidence in favor of the hypothesis of A. Bohr [2], that the fissioning nucleus is strongly cooled in the transition state, as a result of which the fission reaction proceeds at low excitations via a small number of accessible channels. Willets and Chase [3], analyzing the angular distribution \( W(\theta) \) of the fragments, obtained in [1], established that the fission of Th^{232} by neutrons of energy \( E_n = 1.6 \) MeV proceeds predominantly via a band of fission channels with \( K = 3/2^- \) (\( K \) = projection of total angular momentum \( J \) of the compound nucleus on the fission axis). They represented the distribution \( W(\theta) \) in the form

\[ W(\theta) = a_0 + \sum_{J=3/2}^{7/2} a_J W_{KJ}(\theta), \]

where \( W_{KJ}(\theta) \) are the fragment angular distributions in fission via channels with characteristics \( K = 3/2^- \) and \( J > K \); the coefficients \( a_0 \) and \( a_J \) were determined by least squares. Later Strutinskii [4] and Kittmair [5] have shown that good agreement can be attained with experiments by using a more consistent calculation based on the probability of formation of a compound nucleus in the optical model. In particular, the authors of [4,5] have noted that the agreement with experiment does not become worse if the contribution of the band with \( K = 1/2 \) is assumed in lieu of the isotropic component in \( W(\theta) \). The predominance of the states of the transition Th^{232} nucleus with \( K = 1/2 \) at lower values of \( E_n \) was established experimentally [6]. On the basis of measurements of the angular anisotropy \( W(0^\circ)/W(90^\circ) \) of fragment emission, Lamphere [6] proposes the sequence \( 1/2^+, 3/2^- \), and \( 1/2^- \) for the fission channel bands \( K^2 \) that are possible in the Th^{232}(n, f) reaction.
considered by us yields for $1/T_1$ and $\alpha$ values which are approximately $10^4$ times larger in AF-EP than in AF-EA. A numerical estimate for hematite ($\alpha$-Fe$_2$O$_3$) gives a value of $1/T_1$ which agrees with experiment [4]. (For other AF-EP the parameter $\omega^{(1)}_{\text{ns}}$ is unknown, and furthermore there are no data on $1/T_1$.)

The preferred method of observing acoustic NMR is to suppress the ordinary NMR acoustically (using ultrasound at the NMR frequency). We therefore present also an estimate for the sound flux necessary for acoustic saturation of the nuclear spin system:

$$\Pi = \frac{\rho \omega^4 \gamma^3 \Delta \omega_n}{(\omega^{(1)}_{\text{ns}})^2 \omega_k^2 (\gamma_n H'_n)^2 T_1}.$$  \hspace{1cm} (5)

(Here $T_1$ must be obtained from experiment.) For hematite $\Pi \approx 10^{-7}$ W/cm$^2$.

For AF-EP for which the position of the antiferromagnetic resonance $\omega_1$ at low temperature depends essentially on the temperature of the nuclear spin system, we can propose one more method of observing acoustic NMR. The latter can be observed by determining the shift of the frequency $\omega_1$ when ultrasound of frequency $\omega_n$ is applied to the sample. The required sound-flux power must again be estimated from (5). It should be noted that the field $5H_2'$ being due to the coupling of the nuclear spins with the lattice via the low-frequency branch of the spin waves ($\omega_{1k}$), nevertheless excites that branch of the nuclear-spin oscillations, which interacts with the high-frequency branch of the spin waves ($\omega_{2k}$). The frequency $\omega_n$ of this NMR branch remains undisplaced (see, e.g., [5]).

The field $5H_2'$ gives analogous formulas for $1/T_1$ and $\alpha$, except that $\omega_{10}$ is replaced by $\omega_{20}$ and $\omega_{1n}^{(1)}$ is replaced by a combination of $\omega_{2n}^{(2)}$ and $\omega_{n}^{(2)}$ (which depends on the direction of $k$ and $e_\lambda$). This component of $5H$ excites the nuclear-spin oscillation branch interacting with $\omega_{2k}$. For such weakly-anisotropic ferromagnets as RbMnF$_3$ and K$_2$MnF$_3$, for which $\omega_{20} \approx \omega_{10}$, the effects due to $5H_2'$ and $5H_2'$ are comparable in magnitude.


**ERRATA**

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On page 24, line 12 from bottom, read $\vec{v}_\mu = \vec{v}_0$ instead of $v_\mu = v_0$

" " 24, " 9 " " , " $\vec{P}$ and $\vec{N}$ " " $P$ and $N$

" " 25, " 7 " top , " $\vec{E}_-$ " " $\Sigma_-$

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On page 25, line 14 from top, read [3] instead of [5].

" " 26, " 4 " top, " [5]  " " [7].
" " 26, " 16 " " , " Curl a_1  " " R_0 t a_1 .
" " 28, " 1 " " , " \sqrt{2} \frac{m_c}{p}  " " \sqrt{2} \frac{m_c}{p}.

" " 29, last equation  " \Sigma^- \rightarrow \bar{N} + \pi^-  " " \Sigma^- \rightarrow \bar{N} + \pi^- .