ASYMPTOTIC FREEDOM, LIGHT QUARKS AND THE ORIGIN OF THE 
$\Delta T = 1/2$ RULE IN THE NON-LEPTONIC DECAYS OF STRANGE PARTICLES

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A dynamical mechanism for the $\Delta T = 1/2$ rule in the non-leptonic decays of strange particles is considered. The weak interactions are described within the Weinberg-Salam model while the strong interactions are assumed to be mediated by the exchange of an octet of colour vector gluons. It is shown that the effect of the strong interactions gives rise to new operators in the effective Hamiltonian of weak interactions which contain both left- and right-handed fermions. These operators satisfy the $\Delta T = 1/2$ rule, and estimates within the relativistic quark model indicate that their contribution dominates the physical amplitudes of the $K \to 2\pi, 3\pi$ decays.

1. Introduction

In this paper we will consider a possible dynamical mechanism for the $\Delta T = 1/2$ rule in the non-leptonic decays of strange particles. The paper consists of two parts. First, we will derive the effective Hamiltonian of the weak interactions arising from the effect of strong interactions at short distances, and second, we will estimate the matrix elements of this Hamiltonian.

We will consider the "standard" model according to which the strong interactions are mediated by an octet of colour massless gluons while the weak interactions are described by the Weinberg-Salam model [1] with four quarks $p, n, \lambda,$ and $c$. Within this framework a number of interesting results were recently obtained concerning the structure of weak interactions [2,3]. The idea is to calculate the log $\mu_W$ terms by means of the renormalization group where $\mu_W$ is the intermediate W-boson mass. Such logarithmic terms arise from the graphs with gluon exchange and it was shown in refs. [2,3] that the effective Hamiltonian obtained in this way differs from the usual charged currents product by the fact that some pieces of the product are enhanced and some are suppressed.

We will show that the effect of strong interactions not only modifies the relative strengths of various pieces of the bare Hamiltonian but also gives rise to new structures of the type

$$\Delta H_W = c \frac{2G_F}{\sqrt{2}} \sin \theta \cos \theta \left\{ \bar{\lambda}_L \gamma_\mu n_L (\bar{p}_R \gamma_\mu p_R + \bar{n}_R \gamma_\mu n_R + \bar{\lambda}_R \gamma_\mu \lambda_R) \right\} + \text{h.c.} ,$$

$$316$$
where $G_F$ is the Fermi coupling constant, $\theta$ is the Cabibbo angle and the constant $c$ is related to the hadronic, W-boson masses and gluon coupling constant $g$; $p, n, \lambda$ are the operators of the corresponding quark fields.

Unlike the bare Hamiltonian, the new piece $\Delta H_W$ contains along with the left-handed particles $\psi_L$, the right-handed fermions $\psi_R$, where $\psi_{L,R} = \frac{1}{2}(1 \pm \gamma_5)\psi$. In what follows it is essential for us that $\Delta H_W$ satisfies the $\Delta T = \frac{1}{2}$ rule and commutes with the (V-A) generators of the chiral SU(2) × SU(2) group.

We will estimate the matrix elements from $\Delta H_W$ for kaonic decays $K \rightarrow 2\pi, 3\pi$, within the quark model with light quarks [4,5]. It turns out that the matrix elements from the Hamiltonian (1) are greatly enhanced as compared to the matrix elements from the Hamiltonian containing the left-handed particles alone. The reason is that $\Delta H_W$ can annihilate both left- and right-handed quarks from which the mesons are constructed.

However, explicit calculation shows that the coefficient $c$ in eq. (1) although it contains $\ln a W$ in positive powers, is rather small numerically so that the balance between the dynamical enhancement of the matrix elements and the smallness of the coefficient $c$ is a rather delicate one, though estimates still indicate that $\Delta H_W$ dominates the physical amplitudes.

If the $\Delta T = \frac{1}{2}$ rule is due to the piece of the Hamiltonian (1), then this could have important implications for the decays of charmed particles. The point is, that there is no analogous term for the decays of charmed particles, and for this reason the estimates of the relative probability of leptonic and non-leptonic decays of charmed particles extrapolated from the data on K-meson decays [6], may change, with more space given to leptonic modes. Let us note that in papers [6] the enhancement of non-leptonic decays of charmed particles is predicted based on an alternative explanation of the $\Delta T = \frac{1}{2}$ rule [2]. In refs. [2], the $\Delta T = \frac{1}{2}$ rule is explained through the dynamical enhancement of the corresponding piece of the Hamiltonian of the left-handed particles' interaction.

2. Summation of the leading logarithmic terms

In this section we will consider the Hamiltonian of weak interactions in the leading approximation in $g^2 \ln M/\mu$ where $\mu$ is some hadronic mass and $M$ is some large mass scale. Usually $M$ stands for the W-boson mass $\mu_W$, but in the case considered $M$ is sometimes $\mu_W$ and sometimes the mass of the charged quark $\mu_c$. Formally, we assume that $\mu \ll \mu_c \ll \mu_W$. Since $\mu_c$ can in fact not be too large, say, 2–3 GeV, the correction from the neglected terms is large, generally speaking. Still, it seems useful to sum up all the leading log terms and find the amplitude in a self-consistent way.

Let us start by considering the contribution of first order in $g^2$. (Hereafter we define the quark-gluon interaction vertex as $\frac{1}{2} g \bar{\psi} \gamma_\mu t^a \psi G^a_\mu$ where $t^a$ ($a = 1, \ldots, 8$) are the SU(3) matrices in the colour space, $G^a_\mu$ is the gluon field, and matrices $t^a$ are
normalized by the condition $\text{Tr}(a^a b^b) = 2\delta^{ab}$. In fig. 1 all the graphs containing the factor $g^2 \ln \mu_W^2/\mu^2$ where $\mu$ is some characteristic mass, are represented (hereafter we use the Landau gauge for the gluon field). The graphs of figs. 1a and 1b are in fact studied in refs. [2] and we will concentrate on the graph of fig. 1c. The explicit expression for it is as follows

\begin{equation}
M = \frac{2G_F}{\sqrt{2}} \sin \theta \cos \theta \frac{g^2}{3} \frac{1}{16\pi^2} \left( \ln \frac{\mu_W^2}{\mu^2} - \ln \frac{\mu_W^2}{\mu_c^2} \right)
\end{equation}

\begin{equation}
\times \bar{\lambda}_L \gamma_\rho t^a n_L (\bar{p} \gamma_\rho t^a p + \bar{n} \gamma_\rho t^a n + \bar{\lambda} \gamma_\rho t^a \lambda),
\end{equation}

where we have suppressed the colour indices. Eq. (2) is valid for all the outer momenta scalar products of the order $\mu^2$ and we neglected the quark masses as compared to the parameter $\mu$ of the infrared cut-off.

Due to GIM cancellation, the matrix element (2) vanishes in the limit of exact SU(4) symmetry, or when the masses of the $p$- and $c$-quarks coincide. However, it tends to zero only as $\ln \mu_c^2/\mu^2$ (or as $\ln \mu_W^2/\mu_c^2$ if $\mu < \mu_p$), and for the real case of $\mu^2 >> \mu_c^2$ the matrix element is not suppressed. Let us remark that at this point our conclusions are at variance with those of ref. [2] where it is virtually assumed that all the graphs except those of the type represented in figs. 1a, 1b have extra factors of the order $(\mu^2 - \mu_p^2)/\mu_W^2$.

For the physical applications considered below, it is sufficient to use the calculations of first order in $g^2$. To be more precise, however, we will sum all the terms of the order $(g^2 \ln \mu_W^2/\mu_c^2)^n$ or $(g^2 \ln \mu_W^2/\mu^2)^m$. The summation of such terms can be performed by solving the renormalization group equations as was suggested in refs. [7,8] and applied to the weak interactions in refs. [2,3].

The effect of the virtual gluons with momenta squared ranging from $p_1^2$ to $p_2^2$...
leads, in general, to the appearance of the factor
\[ [g^2(p_0^2)/g^2(p_0^2)]^\epsilon \]
where \( g^2(p_0^2) \) is the effective charge of the gluon interaction with momentum squared \( p_0^2 \). The value of constant \( \epsilon \) in eq. (3) can be inferred from the lowest-order calculation in perturbation theory.

Let us divide the whole range of integration over the virtual momenta into two intervals, one extending from \( \mu_c \) to \( \mu_W \) and the other covering the range from \( \mu \) to \( \mu_c \). In the former interval the graphs of the type represented in fig. 1c with p- and c-quarks in the intermediate state cancel each other in the leading logarithmic approximation. The summation of the graphs of the type represented in figs. 1a and 1b leads to the effective Hamiltonian obtained in refs. [2]
\[ H_{\text{eff}} = \frac{2G_F}{\sqrt{2}} \sin \theta \cos \theta \left\{ K_1^{12/25}(-O_1) + K_1^{-6/25}(\frac{1}{3}O_2 + \frac{1}{3}O_3 + \frac{2}{3}O_4) \right\}, \] (4)
where
\[ K_1 = g^2(\mu_c^2)/g^2(\mu_W^2) = 1 + \frac{25}{3} \frac{g^2(\mu_c^2)}{16 \pi^2} \ln \frac{\mu_W^2}{\mu_c^2}, \] (5)
and \( O_i \) are the local four-fermion operators with various selection rules in isotopic and unitary spin:
\[ O_1 = \bar{\lambda}_L \gamma \mu n_L \bar{p}_L \gamma \mu p_L - \bar{\lambda}_L \gamma \mu p_L \bar{p}_L \gamma \mu n_L, \quad (\{8_f\}, \Delta T = \frac{1}{2}), \]
\[ O_2 = \bar{\lambda}_L \gamma \mu n_L \bar{p}_L \gamma \mu p_L + \bar{\lambda}_L \gamma \mu p_L \bar{p}_L \gamma \mu n_L + 2\bar{\lambda}_L \gamma \mu n_L \bar{n}_L \gamma \mu n_L + 2\bar{\lambda}_L \gamma \mu n_L \bar{n}_L \gamma \mu n_L, \quad (\{8_d\}, \Delta T = \frac{1}{2}), \]
\[ O_3 = \bar{\lambda}_L \gamma \mu n_L \bar{p}_L \gamma \mu p_L + \bar{\lambda}_L \gamma \mu p_L \bar{p}_L \gamma \mu n_L + 2\bar{\lambda}_L \gamma \mu n_L \bar{n}_L \gamma \mu n_L - 3\bar{\lambda}_L \gamma \mu n_L \bar{n}_L \gamma \mu n_L + \bar{\lambda}_L \gamma \mu n_L \bar{n}_L \gamma \mu n_L, \quad (\{27\}, \Delta T = \frac{1}{2}), \]
\[ O_4 = \bar{\lambda}_L \gamma \mu n_L \bar{p}_L \gamma \mu p_L + \bar{\lambda}_L \gamma \mu p_L \bar{p}_L \gamma \mu n_L - \bar{\lambda}_L \gamma \mu n_L \bar{n}_L \gamma \mu n_L, \quad (\{27\}, \Delta T = \frac{3}{2}), \] (6)
where we have indicated in parentheses the properties of the corresponding operators with respect to the unitary (SU(3)) and isotopic (SU(2)) groups. The colour indices are suppressed and, e.g., \( \bar{\lambda}_L \gamma \mu n_L \bar{p}_L \gamma \mu p_L \) stands for \( (\bar{\lambda}_L \gamma \mu n_L') (\bar{p}_L \gamma \mu p_L') \) \((i, k = 1, 2, 3)\).

Let us consider now the virtual momenta ranging from \( \mu^2 \) to \( \mu_c^2 \). Apart from the graphs of figs. 1a and 1b the graph of fig. 1c is essential in this case. The effective
Hamiltonian now contains both right- and left-handed fermions since the strong coupling of gluons to quarks is involved. For this reason two new operators are added to the set (6):

\[
O_5 = \bar{\chi}_L \gamma_\mu t^d n_L (\bar{p}_R \gamma_\mu t^d p_R + \bar{n}_R \gamma_\mu t^d n_R + \bar{\lambda}_R \gamma_\mu t^d \lambda_R), \quad \{8\}, \Delta T = \frac{1}{2},
\]

\[
O_6 = \bar{\chi}_L \gamma_\mu n_L (\bar{p}_R \gamma_\mu p_R + \bar{n}_R \gamma_\mu n_R + \bar{\lambda}_R \gamma_\mu \lambda_R), \quad \{8\}, \Delta T = \frac{1}{2}. \tag{7}
\]

The higher-order corrections in gluon exchange mix operators \(O_1, O_2, O_5, O_6\) while leaving \(O_3\) and \(O_4\) diagonal since strong interactions induce transitions only among operators with the same selection rules in isotopic and unitary spin.

To find the linear combinations of operators \(O_{1,2,5,6}\) which are renormalized multiplicatively we need to calculate the mixing of the operators in first order in \(g^2\). Let us define the matrix \(A\) in such a way that the correction to operators \(O_{1,2,5,6}\) in first order in \(g^2\ln \frac{\mu^2}{\mu^2}\) is given by

\[
\delta O = \frac{g^2}{16\pi^2} \ln \frac{\mu^2}{\mu^2} \delta O \cdot A,
\]

where \(O\) is understood as a matrix with one row \(\{O_1, O_2, O_5, O_6\}\).

The straightforward calculation gives

\[
A = \begin{bmatrix}
4 - 2/9 & 10/9 & 4/3 & 0 \\
1/9 & -2 - 5/9 & -2/3 & 0 \\
1/6 & -5/6 & 6 & 3/2 \\
0 & 0 & 16/3 & 0
\end{bmatrix} \tag{8}
\]

The characteristic values of this matrix are as follows:

\[
x_1 = 3.756, \quad x_2 = -2.680, \quad x_3 = 7.222, \quad x_4 = -1.076, \tag{9}
\]

while the eigenvectors are given by

\[
U_1 = O_1 + 0.021 O_2 - 0.034 O_5 - 0.048 O_6,
\]

\[
U_2 = -0.204 O_1 + O_2 + 0.152 O_5 - 0.303 O_6,
\]

\[
U_3 = 0.366 O_1 - 0.064 O_2 + O_5 + 0.738 O_6,
\]

\[
U_4 = 0.034 O_1 + 0.093 O_2 - 0.202 O_5 + O_6.
\]

The effect of strong interactions at distances \(\mu_0^{-1} < \mu < \mu^{-1}\) reduces, for these operators, to the multiplicative renormalization by some power of the effective charge:

\[
U_i \rightarrow U_i \kappa^{2i/9}, \quad i = 1, 2, 3, 4. \tag{11}
\]
Let us stress that the replacement of the coefficient \( \frac{2\pi}{3} \) in the log term in eq. (5) by the coefficient 9 in eq. (12) is due to the fact that the contribution of the \( \bar{c}c \) pair into the vacuum polarization is not logarithmic for the momenta considered.

For the diagonal operators \( U_5 \equiv O_3 \) and \( U_6 \equiv O_4 \) the calculation of the lowest order graphs leads to

\[
x_5 = x_6 = -2, \quad U_i \rightarrow U_i \kappa_2^{x_i/9}, \quad i = 5, 6.
\]

To find the final answer for the effective Hamiltonian one has to represent the Hamiltonian (4) as a sum of diagonal operators \( U_i \) \((i = 1, \ldots, 6)\) and multiply each operator by the factor \( \kappa_2^{x_i/9} \). In this way we have

\[
H_2^{\Delta S, \bar{S}} = -\rho \frac{2G_F}{\sqrt{2}} \sin \theta \cos \theta \left( -\kappa_1^{0.48} O_1 (0.98 \kappa_2^{0.42} + 0.01 \kappa_2^{0.80}) 
-0.04 \kappa_1^{-0.72} \kappa_2^{0.42} + 0.04 \kappa_1^{-0.72} \kappa_2^{-0.30} \right) 
+ \frac{1}{3} \kappa_1^{0.24} O_2 (0.96 \kappa_2^{-0.30} + 0.03 \kappa_2^{-0.12} - 0.11 \kappa_1^{0.72} \kappa_2^{0.42}) 
+ 0.10 \kappa_1^{0.72} \kappa_2^{-0.30} 
+ \frac{2}{15} \kappa_1^{-0.24} \kappa_2^{-0.22} O_3 + \frac{2}{3} \kappa_1^{-0.24} \kappa_2^{-0.22} O_4 
+ 10^{-2} O_5 [\kappa_1^{0.48} (3.3 \kappa_2^{0.42} + 0.3 \kappa_2^{-0.30} - 3.9 \kappa_2^{0.80} + 0.3 \kappa_2^{-0.12}) 
+ \kappa_1^{-0.24} (-0.1 \kappa_2^{0.42} + 2.9 \kappa_2^{-0.30} - 1.4 \kappa_2^{0.80} - 1.4 \kappa_2^{-0.12})] 
+ 10^{-2} O_6 [\kappa_1^{0.48} (4.8 \kappa_2^{0.42} - 0.6 \kappa_2^{-0.30} - 2.9 \kappa_2^{0.80} - 1.3 \kappa_2^{-0.12}) 
+ \kappa_1^{-0.24} (-0.2 \kappa_2^{0.42} - 5.8 \kappa_2^{-0.30} - 1.0 \kappa_2^{0.80} + 7.0 \kappa_2^{-0.12})],
\]

The effect of the virtual momenta ranging from \( \mu \) to \( \mu_c \) does not change essentially the Hamiltonian of the left-handed particles' interaction found in refs. [2] where the annihilation graphs of the type represented in fig. 1c were not considered. The emergence of the new structures (7) seems to be more important. As follows from the estimates presented in sect. 3, the corresponding matrix elements can be large.
3. Estimates of the matrix elements

In this section we will consider the matrix element of the $K\pi$ transition

$$M_I \equiv \langle \pi^+ | O_I | K^+ \rangle,$$  \hspace{1cm} (15)

within a simple quark model. As is well-known all the matrix elements of $K \to 2\pi, 3\pi$ decays reduce to that of the $K\pi$ transition by application of the standard soft-pion technique. This circumstance as well as the simple quark structure in (15) serves as a motivation of our choice to consider the $K\pi$ transition.

We will assume that the $\pi$ ($K$) consists of a quark-antiquark pair. The matrix elements (15) can then be represented in the form (for example, for $M_1$)

$$M_1 = \langle \pi^+ | \bar{p}_{Lk} \gamma_\mu n^I_\mu | 0 \rangle \langle 0 | \bar{\chi}_L \gamma_\mu p^I_\mu | K^+ \rangle - \langle \pi^+ | \bar{p}_{Lk} \gamma_\mu n^I_\mu | 0 \rangle \langle 0 | \lambda_L \gamma_\mu p^I_L | K^+ \rangle,$$  \hspace{1cm} (16)

where $i, k$ are the colour indices.

Only the contribution of the vacuum intermediate state is kept in eq. (16) since the many quark admixture in the meson wave functions is neglected. It is worth emphasizing that eq. (16) accounts for the contribution of the vacuum state both in the $s$- and $t$- ($u$-) channels (to calculate the latter we use the Fierz transformation of the four-fermion operator). In this respect estimate (16) differs from a more familiar one which keeps only the contribution of the $s$-channel vacuum state. From the phenomenological point of view the vacuum state in $t$- ($u$-) channel corresponds to the pion and higher-mass states in the $s$-channel.

Since $\pi$ ($K$)-mesons are colour singlets we have

$$\langle \pi^+ | \bar{p}_{Lk} \gamma_\mu n^I_\mu | 0 \rangle = \frac{1}{3} \delta_i^k \frac{i}{2} f_\pi k_\mu,$$

$$\langle 0 | \bar{\chi}_L \gamma_\mu p^I_\mu | K^+ \rangle = \frac{1}{3} \delta_i^k \frac{i}{2} f_K k_\mu,$$  \hspace{1cm} (17)

where $f_\pi$ and $f_K$ are the constants of the $\pi \to \mu\nu, K \to \mu\nu$ decays and $k_\mu$ is the four-momentum of meson. Thus, we get for $M_1$

$$M_1 = -\frac{2}{3} \frac{1}{2} f_\pi f_K k^2.$$  \hspace{1cm} (18)

Let us estimate now the matrix elements from $O_5, O_6$. To this end we use the Fierz transformation

$$\bar{\psi}_{1L} \gamma_\mu \psi_{2L} \bar{\psi}_{3R} \psi_{4R} = -2 \bar{\psi}_{1L} \psi_{4R} \bar{\psi}_{3R} \psi_{2L}$$  \hspace{1cm} (19)

and the divergence equation

$$\bar{\psi}_1 \gamma_5 \psi_2 = \frac{-i}{m_1 + m_2} \partial_\mu (\bar{\psi}_1 \gamma_\mu \gamma_5 \psi_2),$$  \hspace{1cm} (20)

which is valid both for the bare and the interacting quark in the model of strong interactions considered.
As a result we have

\[ M_5 = \frac{16}{3} M_6 = - \frac{32}{9} \frac{1}{2\mu_\lambda + \mu_p} \frac{1}{2}(f_\pi m_\pi^2) (f_K m_K^2) . \]  

(21)

If the K-meson is on-mass-shell then the ratio of the matrix elements is equal to

\[ \frac{M_5}{M_1} = \frac{16 M_6}{3 M_1} = \frac{8 m_\pi^2}{3 \mu_\lambda (\mu_\lambda + \mu_p)} \frac{m_K^2}{k^2} \approx 70 , \]

(22)

where we used for the quark masses the values obtained in ref. [4] (see also ref. [5]): \( \mu_\lambda \approx 5.4 \text{ MeV}, \mu_p \approx 150 \text{ MeV} \). Any reference to the quark mass could be avoided by using the SU(6)-type relations between the matrix elements \( \langle \pi^+ | [\bar{\psi}_5 \gamma_5 n] | 0 \rangle \) and \( \langle \rho^0 | J_{\mu \nu}^\text{em} | 0 \rangle \).

In this way one gets for example [4]

\[ \mu_p = (m_\pi^2/3m_\rho) (F_\pi/F_\rho) , \]

(23)

where \( F_\pi = \sqrt{2} f_\pi \approx 0.68 \text{ m}_\pi \) and \( F_\rho \) is defined by relation \( \langle \rho^0 | J_{\mu \nu}^\text{em} | 0 \rangle = e_\mu F_\rho m_\rho \); numerically \( F_\rho \approx 1.0 \text{ m}_\pi \).

It is worth emphasizing that the ratio (22) turns out to be large not because the matrix elements from \( O_5, O_6 \) are anomalously large but rather because the matrix element from the standard Hamiltonian \( M_1 \) is small. The reason is that mesons contain both left- and right-handed quarks while \( O_1 \) can annihilate only left-handed particles (the two-particle operator \( \lambda^2 \delta^2 \delta n \) is not important in the Hamiltonian because of the small (of order \( (\mu_\lambda^2 - \mu_p^2)/\mu_W^2 \) coefficient). For light quarks this amounts to a suppression of the corresponding matrix elements. This suppression manifests itself in the smallness of the constants \( f_\pi, f_K \) which are of the order of \( m_\pi \).

Thus the ratio of the matrix elements \( M_5 \) and \( M_1 \) is quite large. However the corresponding operator \( O_5 \) enters the expansion of the effective Hamiltonian with a rather small coefficient, \( \approx 6.5 \cdot 10^{-2} \). Moreover, the contribution of \( O_1 \) is enhanced dynamically by a factor of 2.4 due to the summation of the log terms. In this way we get a factor of \( \frac{1}{37} \) less, and if all the estimates are taken literally the contribution of \( O_5 \) to the physical amplitudes is larger by a factor of 2 compared to that of the operator \( O_1 \). Numerical estimates are made above under the following choice of parameters:

\[ \mu = 0.7 \text{ GeV}, \quad \mu_c = 2 \text{ GeV}, \quad \mu_W = 70 \text{ GeV}, \quad g^2 (\mu^2) / 4\pi = 1 . \]

(24)

If other values of the parameters \( \mu, \mu_c, \mu_W \) were accepted, the final answer for the ratio of the contributions of \( O_5, O_1 \) would be slightly different.

If one considers decays of charmed particles the corresponding piece of the Hamiltonian containing the analog of \( O_5 \) vanishes in the limit of SU(3), and therefore its contribution is greatly suppressed. Thus if non-leptonic decays of K-mesons are due to new structures in \( H_W \) (that is \( O_5, O_6 \)) then there is no reason to believe that the non-leptonic decays of charmed particles are enhanced. Since the balance
between contributions of $O_1$ and $O_5$ is rather delicate, as follows from the estimates presented above, and depends on the effect of relatively large distances, it could be that only study of the decays of charmed particles will clarify the origin of the $\Delta T = \frac{1}{2}$ rule.

**Note added in proof**

Further consideration of the model proposed can be found in preprints ITEP-63, 64, May 1976, Moscow (JETP, in press).

**References**