ON THE ROLE OF FSI IN $K \to 2\pi$ DECAY

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Abstract

Contrary to widely-spread opinion that the final state interaction (FSI) enlarges the amplitude $< 2\pi; I = 0 | K^0 >$, we argue that FSI is not able to increase the absolute value of this amplitude.

1 Introduction

The great progress in understanding the nature of the $\Delta I = 1/2$ rule in $K \to 2\pi$ decays was achieved in the paper [1], where the authors had found a considerable enlargement of contribution of the operators containing a product of left-handed and right-handed quark currents generated by the diagrams called later the penguin ones. But for a quantitative agreement with the experimental data, a search for some additional enlargement produced by long-distance effects was utterly desirable.

In the literature, two possible mechanisms of such increase were discussed. First one was based on assumption that the strengthening of $s$-wave $2\pi$ amplitude with isospin $I = 0$ arises due to small mass of the intermediate scalar $\sigma$ meson. And the calculations in the framework of chiral theory to leading order in momentum expansion of $< 2\pi; I = 0 | K >$ amplitude confirmed such a possibility [2], [3].

The second mechanism of strengthening the $< 2\pi; I = 0 | K >$ amplitude was ascribed to final state interaction of the pions [4]-[12].

In the papers [4]-[7], the calculations were based on the dispersion relation

$$\text{Re} A_I(s) = \frac{1}{\pi} \int_{4\mu^2}^{\infty} dz \frac{\text{Im} A_I^I(z)}{z - s + i\varepsilon} + \text{subtractions}$$

(1)
But $\text{Im}A_j^I(z)$ itself is unknown at $z > m_K^2$, the channels $KK$, $4\pi$ and so on contribute at the values $z$ larger than threshold of elasticity and, in addition, even at $z$ smaller than threshold of elasticity, the resonances of the type of $\sigma$ mesons are possible [13] changing considerably the expression for $\text{Im}A_j^I(z)$ calculated in the framework of ChPT and used in eq.(1).

All these circumstances are the cause to reexamine the role of FSI, using some other approach. It will be done in present paper basing on linear $U(3)_L \otimes U(3)_R \sigma$ model with broken symmetry.

# 2 The technique of calculation

Our calculations are based on employment of the effective lagrangian of non-leptonic, strangeness-changing weak interaction represented in [1]

$$L^{\text{weak}} = \sqrt{2}G_F \sin \theta_C \cos \theta_C \sum_i c_iO_i,$$

where

$$\begin{align*}
O_1 &= \bar{s}\gamma_{\mu}d_L \cdot \bar{u}L\gamma_{\mu}u_L - \bar{s}_L\gamma_{\mu}u_L \cdot \bar{u}_L\gamma_{\mu}d_L \quad (\{8\}, \Delta I = 1/2), \\
O_2 &= \bar{s}_L\gamma_{\mu}d_L \cdot \bar{u}\gamma_{\mu}u_L + \bar{s}_L\gamma_{\mu}u_L \cdot \bar{u}_L\gamma_{\mu}d_L + 2\bar{s}_L\gamma_{\mu}d_L \cdot \bar{d}_L\gamma_{\mu}d_L \\
+ 2\bar{s}_L\gamma_{\mu}d_L \cdot \bar{s}_L\gamma_{\mu}s_L \quad (\{8\}, \Delta I = 1/2) \\
O_3 &= \bar{s}_L\gamma_{\mu}d_L \cdot \bar{u}_L\gamma_{\mu}u_L + \bar{s}_L\gamma_{\mu}u_L \cdot \bar{u}_L\gamma_{\mu}d_L + 2\bar{s}_L\gamma_{\mu}d_L \cdot \bar{d}_L\gamma_{\mu}d_L \\
- 3\bar{s}_L\gamma_{\mu}d_L \cdot \bar{s}_L\gamma_{\mu}s_L \quad (\{27\}, \Delta I = 1/2) \\
O_4 &= \bar{s}_L\gamma_{\mu}d_L \cdot \bar{u}\gamma_{\mu}u_L + \bar{s}_L\gamma_{\mu}u_L \cdot \bar{u}_L\gamma_{\mu}d_L \\
- \bar{s}_L\gamma_{\mu}d_L \cdot \bar{d}_L\gamma_{\mu}d_L \quad (\{27\}, \Delta I = 3/2), \\
O_5 &= \bar{s}_L\gamma_{\mu}\lambda^a d_L \sum_{q=u,d,s} \bar{q}_R\gamma_{\mu}\lambda^a q_R \quad (\{8\}, \Delta I = 1/2), \\
O_6 &= \bar{s}_L\gamma_{\mu}d_L \sum_{q=u,d,s} \bar{q}_R\gamma_{\mu}q_R \quad (\{8\}, \Delta I = 1/2).
\end{align*}$$

This set is sufficient for calculation of the CP-even part of $K \to 2\pi$ amplitude.

The bosonization of the operators $O_i$ in linear $U(3)_L \otimes U(3)_R \sigma$ model is carried out using the relations

$$\begin{align*}
\bar{q}\gamma_{\mu}(1 + \gamma_5)q_j &= i \left( [\partial_{\mu}U]^\dagger - U(\partial_{\mu}U^\dagger) \right)_{ji} \\
\bar{q}_i(1 + \gamma_5)q_j &= -\frac{\sqrt{2}F_m m^2}{m_u + m_d} U_{ji}, \quad F_{\pi} = 93 \text{ MeV}
\end{align*}$$

where

$$U = \hat{\sigma} + i\hat{\pi}, \quad <\sigma_0 + \sigma_8/\sqrt{2}>_0 = \sqrt{\frac{3}{2}}F_{\pi},$$

$$F_{\pi} = 93 \text{ MeV}.$$
and $\hat{\sigma}$ is $3 \times 3$ matrix of the scalar partners of the nonet of pseudoscalar mesons:

$$
\begin{pmatrix}
\frac{\pi_0}{\sqrt{3}} + \frac{\pi_8}{\sqrt{6}} + \frac{\pi_3}{\sqrt{2}} & \pi^+ & K^+
\frac{\pi^-}{\sqrt{3}} + \frac{\pi_8}{\sqrt{6}} - \frac{\pi_3}{\sqrt{2}} & K^0
\frac{\pi^-}{\sqrt{3}} & K^0
\frac{\pi_0}{\sqrt{3}} - \frac{2\pi_8}{\sqrt{6}} &
\end{pmatrix}.
$$

The operators $O_{1-3}$ produce the following part of $<2\pi; I = 0|K^0>$ amplitude:

$$
<\pi^+(p_+)\pi^-(p_-)|\sum_{i=1}^{3} c_i O_i|K^0(q) > = \frac{1}{4}\sqrt{2}F_{\pi}(c_1 - c_2 - c_3)(q^2 - p^2) \quad (7)
$$

The combination $c_5 O_5 + c_6 O_6$, after reordering of quarks in the spinor and colour spaces, turns into

$$
c_5 O_5 + c_6 O_6 = \frac{8}{9}\tilde{c}_5 \tilde{s}(1 - \gamma_5)q \cdot \bar{q}(1 + \gamma_5)d, \quad \tilde{c}_5 = c_5 + \frac{3}{16}c_6. \quad (8)
$$

After bosonization

$$
c_5 O_5 + c_6 O_6 = \frac{8}{9}\tilde{c}_5 \tilde{s}2\sqrt{2}m_\pi^2 [\sigma_0 K^- + \sqrt{\frac{3}{2}}(\sigma_8 K^0 - \pi_8 \sigma K^0) - \sqrt{\frac{1}{2}}(\sigma_3 K^0 - \pi_3 \sigma K^0)]
$$

A breakdown of the $SU(3)$ symmetry together with mixing of $\sigma_0$ with gluonium originates a mixing between $\sigma_0$ and $\sigma_8$, so that the physical states become of the forms

$$
\sigma_{\eta'} = \sigma_0 \cos \theta_S + \sigma_8 \sin \theta_S, \quad \sigma_{\eta} = -\sigma_0 \sin \theta_S + \sigma_8 \cos \theta_S. \quad (10)
$$

To find the $K \to \pi^+\pi^-$ amplitude generated by $O_5$ operator, it is necessary to know the amplitudes of $K^0 \to \sigma K^+\pi^-$ and $\sigma_8 \to \pi^+\pi^-$ transitions. They are determined by the lagrangian [14], [15]:

$$
L_{mesons} = \frac{1}{2} Tr \{ \partial_\mu U \partial_\mu U^\dagger \} - c Tr \{ UU^\dagger - A^2 l_0 \}^2
- c_i (Tr \{ UU^\dagger - A^2 l_0 \})^2 + \frac{F_{\pi}}{2\sqrt{2}} Tr \{ M(U + U^\dagger) \} + \Delta L_{PS} \quad (11)
$$

This lagrangian is considered here as the phenomenological effective lagrangian where the baryonic degrees of freedom are integrated out. In addition, the masses and coupling constants are such that their renormalization generated
by mesonic loops is not needed. However, the renormalization procedure

does not remove the corrections which are the finite functions of the external

momenta. They have to be taken into account and will be considered later.

The parameter \( \xi \) in eq.(11) characterizes a degree of mixing between \( \sigma_0 \)

and gluonium. Such a mixing increases \(< \pi^+; I = 0|K^0 > \) amplitude [2], but

in present paper, devoted to investigation of FSI role, we shall consider, for

simplicity, the case \( \xi = 0 \) in which \( \sin \theta_S = 1/\sqrt{3} \). Then (see [15])

\[
\begin{align*}
g_{\sigma_0 \pi^+} &= -(m_{\sigma_0}^2 - m_\pi^2)/F_\pi, \quad g_{\sigma_0 \pi^-} = 0, \\
g_{\sigma_K K^0 \pi^+} &= -(m_{\sigma_0}^2 - m_\pi^2)(2R - 1)/\sqrt{2}F_\pi, \\
m_{\sigma_0}^2 - m_\pi^2 &= (m_K^2 - m_\pi^2)/(R - 1)(2R - 1), \\
m_{\sigma_K}^2 - m_\pi^2 &= (m_{\sigma_0}^2 - m_\pi^2)(2R - 1)R
\end{align*}
\]

(12)

where \( R = F_K/F_\pi \). In our theory

\[
R = 1 + (m_K^2 - m_\pi^2)/(m_{\sigma_0}^2 - m_\pi^2)....
\]

(13)

The amplitude generated by operators \( O_{5,6} \) is [2]:

\[
\begin{align*}
< \pi^+ \pi^- |c_5 O_5 + c_6 O_6|K^0(q) > &= \frac{8}{9} \sqrt{2} F_\pi m_\pi^4 \left[ m_{\sigma_0}^2 - m_\pi^2 \right] \left[ \frac{1}{1/R} \right].
\end{align*}
\]

(14)

Neglecting the higher-order corrections, we obtain the total amplitude

\[
\begin{align*}
< \pi^+(p_+)\pi^-(p_-)| \sum_{i=1}^{6} c_i O_i|K^0(q) > &= \frac{1}{2} \sqrt{2} F_\pi \left[ (c_1 - c_2 - c_3)(q^2 - p_+^2) + \frac{2\beta}{m_{\pi}^2} \right].
\end{align*}
\]

(15)

where

\[
\beta = \frac{2m_\pi^4}{(m_u + m_d)^2(m_{\sigma_0}^2 - m_\pi^2)}
\]

(16)

Such is the amplitude \(< \pi^+; I = 0|K^0 > \) produced by weak interaction and

incorporating the corrections produced by strong quark-gluon interactions at

short distances. Our next step is to find the changes arising due to rescattering

of pions at large distances where the real intermediate hadrons play

role. We begin from a consideration of the elastic \( \pi \pi \) scattering itself.

4
3 The elastic $\pi\pi$ scattering

An amplitude of elastic $\pi\pi$ scattering depends on two variables: initial energy of two pions and their scattering angle in c.m. system. When this amplitude is expressed in terms of the phase shifts, the partial amplitudes depending only on $s = (p_1 + p_2)^2$ are the objects of investigation. Let’s the $S$-wave partial amplitude (the only one participating in $K \to 2\pi$ decay ) calculated in the tree approximation is

$$A^{\text{tree}} = f(s).$$  \hspace{1cm} (17)

To one-loop it turns into

$$A^{\text{one-loop}} = f(s)[1 + \text{Re}\Pi_R(s) + i\text{Im}\Pi(s)] = f(s)[1 + \text{Re}\Pi_R(s) + i \frac{f(s)\sqrt{1 - 4m^2/s}}{16\pi}]$$  \hspace{1cm} (18)

where Re$\Pi_R(s)$ means the finite part depending on external momenta remaining after regularization of Re$\Pi(s)$.

The unitarization, representing the summing up of the chains built from the loop diagrams corresponding to rescattering of pions, leads to

$$A^{\text{unitar}} = \frac{f(s)}{1 - \text{Re}\Pi_R(s) - i\text{Im}\Pi_R(s)} = \frac{f(s)}{1 - \text{Re}\Pi_R(s)} \cdot \frac{1}{1 - i \frac{f(s)\sqrt{1 - 4m^2/s}}{16\pi(1 - \text{Re}\Pi_R(s))}}$$ \hspace{1cm} (19)

Introducing the notation

$$\frac{f(s)\sqrt{1 - 4m^2/s}}{16\pi(1 - \text{Re}\Pi_R(s))} = \tan \delta,$$ \hspace{1cm} (20)

we come to the known from the non-relativistic quantum mechanics expression of the scattering amplitude through the phase shifts:

$$A^{\text{unitar}} = \frac{16\pi \sin \delta e^{i\delta}}{\sqrt{1 - 4m^2/s}}$$ \hspace{1cm} (21)

leading to the cross-section:

$$\sigma = \frac{4\pi \sin^2 \delta}{k^2}, \quad k = \frac{\sqrt{s}}{2} \sqrt{1 - 4m^2/s}$$ \hspace{1cm} (22)
From eq. (20) and representation of $A_{unitar}$ in the form (19)

$$A_{unitar} = \frac{f(s)}{1 - \text{Re}\Pi_R(s)} \cos \delta e^{i\delta},$$

one concludes that at $\text{Re}\Pi_R(s) < 0$ (more exactly, $\text{Re}\Pi_R(s) < 1 - \cos \delta$) FSI diminishes the absolute value of the initial amplitude (17).

It should be noted in addition that according to eq. (20), $\text{Re}\Pi_R(s)$ can be found using the data on the phase shifts and the theoretical value of $f(s)$ obtained in one or another model of $\pi\pi$ interaction. And this is extremely important because the direct calculation of the contribution from the closed hadron loops in the lowest order of perturbation theory does not allow to obtain the right conclusion on a value of $\text{Re}\Pi_R(s)$. A reason is that such a contribution turns out to be of order 1. In such case the perturbation theory is not applicable.

In the literature devoted to study of the role of the resonances in mesonic theory, a necessity of incorporation of $\text{Re}\Pi_R$ in analysis was not considered, as a rule [16]. An exception are the papers [17] where $\text{Re}\Pi_R$ was calculated in the lowest order of perturbation theory, which does not work in the case of strong coupling, as it was mentioned above. Nevertheless, it is possible to obtain the reliable estimate of $\text{Re}\Pi_R$. As the function $f(s)$ in the chiral theory is proportional to $g_0^2$, the effect of $\text{Re}\Pi_R$ can be transferred into $g^2(s)$:

$$g_0^2 / (1 - \text{Re}\Pi_R(s)) \implies g^2(s) = g_0^2 F(s),$$

where $F(s)$ is the phenomenological form factor chosen so that the theoretical results for the set of phase shifts $\delta_0^0, \delta_0^2, \delta_1^0, \delta_2^0$ and $\delta_2^2$ should coincide with the experimental data in some broad range of values of $s$. This method, simplifying an analysis of the effects of $\text{Re}\Pi_R(s)$ was used in [18] where the Chiral Resonance Theory of $\pi\pi$ Scattering in the range $4m_{\pi}^2 \leq s \leq 1$ GeV$^2$ was elaborated. It follows from [18] that the main contribution into $\text{Re}\Pi_R(s = m_{K^*}^2)$ is generated by $\sigma\pi\pi$ interaction and

$$F(s = m_{K^*}^2) \approx \exp(-2m_{K^*}^2/m_{\pi}^2/2\text{GeV}^2) = 0.894,$$

that is

$$\text{Re}\Pi_R(s = m_{K^*}^2) \approx -0.12.$$
The sign of $\text{Re}\Pi_R(s)$ coincides with the one following from the calculation in the lowest order of perturbation theory [17], but the absolute value turns out to be considerably smaller.

Now, let us pass to study of the FSI influence on the amplitude of $K \to 2\pi$ decay.

4 FSI in $K^0 \to 2\pi$ decay

To one loop, the amplitude (15) transforms into

$$<\pi^+(p_+)\pi^-|p_-)|\sum_{i=1}^6 c_i O_i |K^0(q) >^{\text{one-loop}} = \frac{F_\pi}{2\sqrt{2}} \{ (c_1 - c_2 - c_3)(q^2 - p_-) + \frac{f(s)}{(2\pi)^4} \int \frac{d^4p}{[(p-q)^2-m_\pi^2][p^2-m_\pi^2]} + i \frac{f(s)}{16\pi} \sqrt{1 - 4m_\pi^2/q^2(q^2 - p_-^2)} \}$$

$$+ \frac{32}{9} \beta \tilde{c}_5 (m_K^2 - m_\pi^2)[1 + \frac{f(s)}{(2\pi)^4} \int \frac{d^4p}{[(p-q)^2-m_\pi^2][p^2-m_\pi^2]} + i \frac{f(s)}{16\pi} \sqrt{1 - 4m_\pi^2/q^2}]$$

Taking into account that the part of the first integral proportional to finite function of the external momenta is equal to (see Appendix)

$$(q^2 - m_\pi^2) \frac{f(s)}{16\pi^2} \sqrt{1 - 4m_\pi^2/q^2} \ln \frac{1 - \sqrt{1 - 4m_\pi^2/q^2}}{1 + \sqrt{1 - 4m_\pi^2/q^2}} = (q^2 - m_\pi^2) \text{Re}\Pi_R(q^2)$$

(27)

and the analogous part of the second integral is

$$(q^2 - m_\pi^2) \text{Re}\Pi_R(q^2)$$

and accomplishing the unitarization according to the prescription (19) we come to

$$<\pi\pi; I = 0|K^0(q^2 = m_K^2) > = \frac{F_\pi}{2\sqrt{2}} (m_K^2 - m_\pi^2)[c_1 - c_2 - c_3 + \frac{32}{9} \beta \tilde{c}_5] \frac{\cos \delta e^{i\delta}}{1 - \text{Re}\Pi_R(m_K^2)}$$

(29)

Remembering that the integrals in eq.(27) calculated in the leading order of perturbation theory do not give a reliable estimate of $\text{Re}\Pi_R$ and using the estimate (26) together with $\delta_0^0 \approx 37^\circ$ [18] we come to conclusion that FSI diminishes the tree amplitude (15) by 30%.

The analogous analysis of the influence of FSI on the amplitude $<\pi\pi; I = 2|K^0 >$ leads to conclusion that FSI enlarges this amplitude by 5%.
5 Conclusion

We did not find an enlargement of the amplitude $<\pi\pi; I = 0|K^0>$ due to final state interaction of pions. On the contrary, our analysis showed that FSI diminishes this amplitude. The amplitude $<\pi\pi; I = 2|K^0>$ is increased a little by FSI. So that, FSI is not at all the mechanism bringing us nearer to explanation of the $\Delta I = 1/2$ rule in $K \to 2\pi$ decay. But our result, of course, does not mean that this rule can not be explained. In particular, an increase of the coefficient $\tilde{c}_5$ calculated in theory with large uncertainty (see [1]) can compensate the negative influence of FSI.

Appendix

The loop integral

$$I^{(0)} = \frac{1}{(2\pi)^4} \int \frac{d^4p}{[(p-q)^2 - m^2][p^2 - m^2]}$$

(30)

calculated using the t’Hooft-Veltman dimensional regularization [19] is equal to

$$I^{(0)} = \frac{i}{16\pi^2} \left( \ln \frac{M_0^2}{m^2} + 2 + \sqrt{1 - 4m^2/q^2} \ln \frac{1 - \sqrt{1 - 4m^2/q^2}}{1 + \sqrt{1 - 4m^2/q^2}} \right)$$

(31)

where

$$\ln \frac{M_0^2}{\mu^2} = 1/\epsilon - \gamma_E + \ln(4\pi), \quad \epsilon \to 0$$

and $\mu$ is some arbitrary mass disappearing in the above expression for $I^{(0)}$. After regularization removing from $I^{(0)}$ the terms independent on the external momenta we obtain

$$I_R^{(0)} = \frac{i}{16\pi^2} \sqrt{1 - 4m^2/q^2} \ln \frac{1 - \sqrt{1 - 4m^2/q^2}}{1 + \sqrt{1 - 4m^2/q^2}}$$

(32)
The first loop integral in eq.(27) is

\[
I^{(1)} = \frac{1}{(2\pi)^2} \int \frac{(q^2-p^2)d^4p}{[(p-q)^2-m^2][p^2-m^2]} = \{ q^2I^{(0)} - \frac{im^2}{16\pi^2} \left( 2 \ln \frac{M_0^2}{m^2} + 3 + \sqrt{1 - 4m^2/q^2} \ln \frac{1 - \sqrt{1 - 4m^2/q^2}}{1 + \sqrt{1 - 4m^2/q^2}} \right) \}
\]

(33)

After regularization it converts into

\[
I^{(1)}_R = (q^2 - m^2)I^{(0)}_R.
\]

(34)

Just this part of \(I^{(1)}\) turns out to be proportional on mass shell to the combination \((m_K^2 - m^2)\) that vanishes in the limit of exact SU(3) symmetry in accordance with Cabibbo-Gell-Mann theorem for \(K \to 2\pi\) amplitude [20].

For \(\Pi_R(s)\) in eq.(18), it can be obtained in the lowest order of perturbation theory (see [17]):

\[
\begin{align*}
\Pi_R(s) &= \frac{f(s)}{16\pi} \sqrt{1 - 4m^2/s} \ln \frac{1 - \sqrt{1 - 4m^2/s}}{1 + \sqrt{1 - 4m^2/s}} + i \frac{f(s)}{16\pi} \sqrt{1 - 4m^2/s}, \quad s > 4m^2 \\
\Pi_R(s) &= -\frac{f(s)}{16\pi} |\sqrt{1 - 4m^2/s}| \cdot (1 - \frac{2}{\pi} \arctan |\sqrt{1 - 4m^2/s}|), \quad s < 4m^2.
\end{align*}
\]

(35)

But outside the perturbation theory, a value of \(\text{Re}\Pi_R(s)\) depends on the used model of \(\pi\pi\) interaction. In our case this value is given by eq.(26).

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