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Significance of Single Pion Exchange Inelastic Final State Interaction for $D \to VP$ Processes

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Abstract

We evaluate the contribution of the final state interaction (FSI) due to single pion exchange inelastic scattering for $D^+ \to \bar{K}^{0*}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$ processes. The effects are found to be very significant. The hadronic matrix elements of the weak transition are calculated in terms of the heavy quark effective theory (HQET), so less model-dependent and more reliable.

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I. Introduction

To precisely understand the mechanism governing weak transition process where the fundamental physics and possible new physics apply, one needs to face a synthesis problem including evaluation of the hadronic matrix elements and final state interaction (FSI). Unless one can fully understand the side effects, such as FSI, he can hardly extract information about new physics correctly. At least one needs to know the order of magnitude of FSI and determine its significance to the concerned problem and determine whether it is necessary to be taken into account.

Due to the success of the Standard Model, the hamiltonian for the weak transition $c \to s + u + \bar{d}$ is well understood [1], and thanks to the heavy quark effective theory (HQET) [2], we have some more reliable ways to handle the hadronic matrix elements for the $b \to c$ transition. The developments on chiral lagrangian [3] enables us to make feasible estimation of matrix elements from heavy mesons to light pseudoscalars. According to the recent work by Roberts and Ledroit [4], the transition matrix elements from B, D to $K^{(*)}$, become calculable in a unique theoretical framework.

On other side as discussed above, to really testify the theory concerning the matrix elements evaluation, as well as HQET, one important issue is to deprive of the FSI effects which sometimes become very significant and sometime can be negligible.

In the early work, Buccella et al. [5] found that the calculated branching ratios for $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$ deviated from the data quite apart and even using the recent data, the difference is still obvious. We use the HQET to recalculate (see below for details) and find that the discrepancy with data is not so acute, but still exists. One could suspect whether this declination is caused by the final state interaction (FSI). In fact, many authors have studied the problem of FSI in some D and B decays [5, 6, 7, 8]. The FSI due to s-channel resonances were found to be very important for some final channels[5, 6]. However for the processes $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$, the final states have isospin I=3/2, and therefore have no s-channel resonance FSI. Zheng [7] calculated the elastic FSI effects in $B \to DK$ caused by t-channel meson exchange, and obtained very small phase, so it does not make a substantial effect to the measured rate. Very recently, Donoghue et al. indicate that the inelastic scattering may dominate the FSI [8]. We concur with this. Our previous work in Ψ decays and $\bar{p}p$ annihilations[9] indicate that when a single$-\pi$ exchange at the t-channel can be realized, the corresponding mechanism would make significant contribution to the FSI and very probably, dominant.

Based on our previous studies on the FSI [9], we can conceive that the inelastic scattering is mainly caused by the t-channel single-pion exchange at the final states of the B and D decays. For $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$, the dominant FSI should be the inelastic rescattering between $\bar{K}^{*0}\pi^+$ and $\bar{K}^0\rho^+$, as shown in Fig.1. In this work we estimate this inelastic FSI effect. The hadronic matrix elements are evaluated in terms of the method given in ref.[4] which can alleviate the model-dependence of the calculations.

In Sec.II, we give the formulation of the transition amplitude with and without considering the inelastic scattering of $\bar{K}^{*0}\pi^+ \leftrightarrow \bar{K}^0\rho^+$, and in Sec.III, we present our numerical results, while the last section is devoted to our discussion and conclusion.

II. The formulation

In the weak interaction, the isospin is not conserved. There are four possible VP decay modes for D^+, i.e., $\bar{K}^{*0}\pi^+$, $\bar{K}\rho^+$, $K^{*+}\pi^0$ and $K^+\rho^0$; among them $I_3(\bar{K}^{*0}\pi^+) = I_3(\bar{K}^0\rho^+) = 3/2$ and $I_3(K^{*+}\rho^0) = I_3(K^{*+}\pi^0) = 1/2$. However, from the quark diagrams[10], one can find $D^+ \to K^{*+}\pi^0$ or $D^+ \to K^+\rho^0$ can only be realized via a Cabibbo double suppressed channel, so must be very small and can be neglected. For the FSI, the interaction is the strong interaction which conserves the isospin. The $\bar{K}^{*0}\pi^+$ and $\bar{K}^0\rho^+$ cannot rescatter into $K^{*+}\rho^0$ and $K^{*+}\pi^0$; while $\bar{K}^{*0}\pi^+ \leftrightarrow \bar{K}^0\rho^+$ can be realized by the t-channel single pion exchange diagrams as shown in Fig.1 and can be very important. Therefore, we only need to consider production of $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^{*0}\rho^+$, as well as their mutual conversion through the inelastic scattering.

(i) Without the FSI.

The weak interaction hamiltonian for non-leptonic decay is given as

$$H_W = \frac{G_F}{\sqrt{2}} V_{ud} V_{cs}^* [c_1 \bar{s}\gamma_\mu(1-\gamma_5)c\bar{u}\gamma^\mu(1-\gamma_5)d + c_2 \bar{s}\gamma_\mu(1-\gamma_5)d\bar{u}\gamma^\mu(1-\gamma_5)c] \quad (1)$$

where V_{ud} and V_{cs} are the Cabibbo-Kobayashi-Maskawa entries. It is noted that the concerned reactions are Cabibbo favored processes. The color indices are dropped out as well understood, and the coefficients c_1 and c_2 are obtained from the renormalization group equation [1].

In the calculations, we use the vacuum saturation and ignore the W-exchange (annihilation) contributions [10]. The transition matrix elements for $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$ are given below respectively:

$$\begin{aligned}
T_1(D^+ \to \bar{K}^{*0}\pi^+) &= (c_1 + \frac{1}{N_c}(1+\delta)c_2) < \pi^+|\bar{u}\gamma^\mu(1-\gamma_5)d|0 >< \bar{K}^{*0}|\bar{s}\gamma_\mu(1-\gamma_5)c|D^+ > + \\
&+ (c_2 + \frac{1}{N_c}(1+\delta)c_1) < \bar{K}^{*0}|\bar{s}\gamma^\mu(1-\gamma_5)d|0 >< \pi^+|\bar{u}\gamma_\mu(1-\gamma_5)c|D^+ > \\
&= -(c_1 + \frac{1}{N_c}(1+\delta)c_2) f_\pi p_\pi^\mu Tr\{[(\xi_3 + \not{p}_{K^*}\xi_4)\epsilon_{K^*}\cdot v \\
&+ \not{\epsilon}_{K^*}^*(\xi_5 + \not{p}_{K^*}\xi_6)]\gamma_\mu(1-\gamma_5)\cdot\frac{1}{2}\sqrt{M_D}\gamma_5(1-\not{v})\} \\
&+ (c_2 + \frac{1}{N_c}(1+\delta)c_1) f_{K^*}\epsilon_{K^*}^{\mu *} m_{K^*} Tr\{(\xi_1 + \not{p}_\pi\xi_2)\gamma_5\gamma_\mu\cdot\frac{1}{2}\sqrt{M_D}\gamma_5(1-\not{v})\},
\end{aligned} \quad (2)$$

and

$$\begin{aligned}
T_2(D^+ \to \bar{K}^0\rho^+) &= (c_1 + \frac{1}{N_c}(1+\delta)c_2) < \rho^+|\bar{u}\gamma^\mu(1-\gamma_5)d|0 >< \bar{K}^0|\bar{s}\gamma_\mu(1-\gamma_5)c|D^+ > + \\
&+ (c_2 + \frac{1}{N_c}(1+\delta)c_1) < \bar{K}^0|\bar{s}\gamma^\mu(1-\gamma_5)d|0 >< \rho^+|\bar{u}\gamma_\mu(1-\gamma_5)c|D^+ > \\
&= (c_1 + \frac{1}{N_c}(1+\delta)c_2) f_\rho m_\rho \epsilon_\rho^{*\mu} Tr\{(\xi_1 + \not{p}_{(K)}\xi_2)\gamma_5\gamma_\mu\cdot\frac{1}{2}\sqrt{M_D}\gamma_5(1-\not{v})\} \\
&- (c_2 + \frac{1}{N_c}(1+\delta)c_1) f_K p_K^\mu Tr\{[(\xi_3 + \not{p}_\rho\xi_4)\epsilon_\rho\cdot v \\
&+ \not{\epsilon}_\rho^*(\xi_5 + \not{p}_\rho\xi_6)]\gamma_\mu(1-\gamma_5)\cdot\frac{1}{2}\sqrt{M_D}\gamma_5(1-\not{v})\},
\end{aligned} \quad (3)$$

where v is the four-velocity of D^+ as $p_D^\mu = M_D v^\mu$ and ξ_i (i=1,...,6) are functions of momenta given in ref.[4]. δ is a non-factorization factor and cannot be evaluated in perturbative QCD [11, 12]. Recently, Sharma et al.[13] investigated the non-factorization effects in $D \to PV$ decays. Blok and Shifman[14] give a more theoretical estimation of the factor as

$$\delta = -N_c \frac{x m_{\sigma H}^2}{4\pi^2 f_\pi^2}, \tag{4}$$

where $x \sim 1$, $m_{\sigma H}$ is a numerical factor. Admittedly[14], one cannot take the number very seriously, so we keep it as a parameter in the region of -0.5$\sim$ -1.0. In fact, in our calculations, we take $\delta = -0.5$. Since we are mainly discussing the significance of FSI, the choice of δ does not influence our qualitative conclusion at all.

(ii) With the final state interaction.

Here as discussed in our previous work[9], to estimate the FSI, one only needs to calculate the absorptive part of the diagram. According to the Cutosky rule, for getting the absorptive part of the loop shown in Fig.1, there are two ways to make cuts, i.e. (1) let the V and P be on mass-shell while retain the t-channel pion off-shell; (2) let the pion and P on shell while leaving V off-shell. The second way refers to a three-body decay process and numerical computation for similar triangle diagrams in $\bar{p}p$ annihilation[15] shows that is much smaller than (1), so we omit the possibility in our later formulation.

(a) For $D^+ \to \bar{K}^{0*}\pi^+ \to \bar{K}^0\rho^+$.

In the CM frame of D^+ where $v = (1, \vec{0})$, the calculation can be greatly simplified. To obtain the absorptive part of the loop, for example T_3, one can just start from eq.(2), replace $\epsilon_{K^*}^\mu$ by $\frac{1}{2}(2\pi)^2\delta(p_1^2 - m_1^2)\delta(p_3^2 - m_3^2)(-g_{\mu\mu'} + \frac{p_{1\mu}p_{1\mu'}}{m_1^2})$ and add the effective vertices of strong interaction as well as the propagator of the t-channel pion.

$$\begin{aligned} T_3 \;=\; & \int \frac{d^4p_1}{(2\pi)^4}(p_2 - p_{\pi^0})^\mu \epsilon_\rho^* \cdot (p_3 - p_{\pi^0}) \frac{F(p_{\pi^0}^2)}{p_{\pi^0}^2 - m_{\pi^0}^2} \frac{1}{2}(2\pi)^2 \delta(p_1^2 - m_1^2)\delta(p_3^2 - m_3^2) g_{\rho\pi\pi} g_{K^*K\pi} \\ \times\; & \{-2(c_1 + \frac{1}{N_c}(1+\delta)c_2) f_\pi \sqrt{M_D}[(\xi_3 E_3 - \frac{1}{2}(\xi_4 - \xi_6)(M_D^2 - m_1^2 - m_3^2))(-g_{0\mu} + \frac{p_{10}p_{1\mu}}{m_1^2}) \\ -\; & (\xi_5 + \xi_6 E_1) p_3^{\mu'}(-g_{\mu\mu'} + \frac{p_{1\mu}p_{1\mu'}}{m_1^2}) + i\xi_6 \epsilon^{ijk} p_{1j} p_{3k}(-g_{i\mu} + \frac{p_{1i}p_{1\mu}}{m_1^2})] \end{aligned}$$

$$+ \quad [2(c_2 + \frac{1}{N_c}(1+\delta)c_1)f_{K^*}m_1\sqrt{M_D}(\xi_1(-g_{0\mu} + \frac{p_{10}p_{1\mu}}{m_1^2}) - \xi_2 p_3^{\mu'}(-g_{\mu\mu'} + \frac{p_{1\mu}p_{1\mu'}}{m_1^2}))]\}. \tag{5}$$

(b) For $D^+ \to \bar{K}^0\rho^+ \to \bar{K}^{*0}\pi^+$.

The amplitude T_4 has a similar form as T_3 with some changes, and can be easy to obtain from 3, For saving space, we do not give the explicit expression here.

In the expressions, for simplicity of bookkeeping we set $p_1 \equiv p_{K^*}$, $p_2 \equiv p_K$, $p_3 \equiv p_{\pi^+}$, $p_4 \equiv p_\rho$ and $m_1 \equiv m_{K^*}$, $m_2 \equiv m_K$, $m_3 \equiv m_{\pi^+}$, $m_4 \equiv m_\rho$.

$F(p_{\pi^0}^2)$ is an off-shell form factor for the vertices $\rho\pi\pi$ and $\bar{K}^{*0}\bar{K}^0\pi$. Because we use the experimental data where all the three particles are real and on mass-shell to fix the effective coupling at the vertices, in the case of the pion being off-shell, a compensation form factor is needed and it is

$$F(p_{\pi^0}^2) = (\frac{\Lambda^2 - m_{\pi^0}^2}{p_{\pi^0}^2 - m_{\pi^0}^2})^2, \tag{6}$$

with Λ in the range of $1.2 \sim 2.0$ GeV.

From above equations and employing the helicity-coupling amplitude formalism given by Chung [16], the whole calculation is straightforward though tedious, and a direct comparison at the amplitude level is feasible. We present our numerical results in the next section.

III. The numerical results.

The two strong coupling constants $g_{\rho\pi\pi}$ and $g_{K^*K\pi}$ can be obtained from the ρ and K^* decay width, respectively. From the newest PDG data [17], we have $g_{\rho\pi\pi} = 6.1$ and $g_{K^*K\pi} = 5.8$. The values of c_1 and c_2 are taken from Ref.[13], i.e., $c_1 = 1.26$ and $c_2 = -0.51$.

As aforementioned, we take the non-factorization factor δ to be -0.5. The ξ_i's (i=1,...,6) have simple Gaussian forms or polynomials. Their explicit forms and parameters can be found in ref.[4]. There are three sets of parameters of a_i and b_i (notation in ref.[4]), which look very disperse. We substitute all the three sets into our expressions to carry out the calculations and compare their results.

(i) First, we calculate the decay rate without taking into account the FSI, i.e. we only use eqs. (2) and (3). The results are listed in Table I.

	Fit 1	Fit 2	Fit 3	Exp.
$\Gamma(D^+ \to \bar{K}^0\rho^+)$ ($\times 10^{-13}$ GeV)	1.0	1.0	0.65	0.41 ± 0.15
$\Gamma(D^+ \to \bar{K}^{*0}\pi^+)$ ($\times 10^{-13}$ GeV)	0.09	0.049	0.079	0.124 ± 0.025

Table I. Results without considering FSI and with three sets of parameters from [4].

Here Fit1, Fit2 and Fit3 correspond to three different sets of parameters for the functions ξ_i (i=1,...,6) from Roberts et al.[4] fitting $D \to K^{(*)}l\nu$; $D \to K^{(*)}l\nu$, $B \to K^{(*)}J/\psi$ and $D \to K^{(*)}l\nu$, $B \to KJ/\psi$, $B \to K\psi'$, respectively. One can notice that even though the parameters in the three fits are quite apart (their signs can change either), the obtained results are rather close to each other. These results are still qualitatively consistent with that obtained by authors of ref.[5] without the FSI, namely the calculated rate for $D^+ \to \bar{K}^{*0}\pi^+$ is lower than the experimental value by 1.4~1.6 times, while for $D^+ \to \bar{K}^0\rho^+$, the calculated number is $1.6 \sim 2.5$ times larger than data.

(ii) The FSI contribution.

Since we only consider the absorptive part of the loop, we can get the imaginary part of the FSI amplitude. The dispersive part can be obtained by the dispersive relation with some cut-off parameters and is supposed to hold the same order as the absorptive one [15]. The absorptive part of the amplitude which gives rise to the lower bound of the magnitude caused by FSI.

For a clean comparison and avoid ambiguity, we use the helicity amplitudes. In the helicity picture, for the CM frame of D^+ all momenta of the outgoing P and V are aligned along $\hat{z}$ axis (or oppositely). It is noted that even the mesons in the loop are real, but their momenta can deviate from $\hat{z}$ axis by an angle θ. So in this scenario, for a reaction $D \to PV$, only $\epsilon^*(0)$ polarization of the final outgoing vector meson contributes. In Table II, we present the numerical value for the absorptive part of the FSI loop with $\Lambda = 1.6GeV$, and as a comparison also the helicity amplitude without FSI evaluated in the same theory.

	Fit1	Fit2	Fit3
$T(D^+ \to \bar{K}^0\rho^+ \to \bar{K}^{*0}\pi^+)$	-i0.220	-i0.223	-i0.178
$T(D^+ \to \bar{K}^{*0}\pi^+ \to \bar{K}^0\rho^+)$	-i0.071	-i0.051	-i0.062
$T(D^+ \to \bar{K}^0\rho^+)$	0.52	0.52	0.42
$T(D^+ \to \bar{K}^{*0}\pi^+)$	0.15	0.11	0.14

Table II.

In the table, we only keep the relative values of the helicity amplitudes for with and without the FSI, while dropping out the common factor such as $(G_F/\sqrt{2})V_{cs}V_{ud}^*$ etc. The values given in Table II are calculated with $\Lambda = 1.6$ GeV. In fact, our numerical computation show that as Λ varies from 1.2 to 2.0 GeV, the corresponding results in the first two rows of Table II can change by a factor 2. As one takes a more restrained region for Λ, the results are not very sensitive to the Λ value.

Obviously, the FSI effect of $T(D^+ \to \bar{K}^0\rho^+ \to \bar{K}^{*0}\pi^+)$ is stronger than that of $T(D^+ \to \bar{K}^{*0}\pi^+ \to \bar{K}^0\rho^+)$, we will discuss the results below.

IV. Conclusion and discussion

In above, we calculate the decay width of $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$ in terms of the HQET and consider the contribution of the t-channel single pion exchange to the inelastic scattering $\bar{K}^{*0}\pi^+ \leftrightarrow \bar{K}^0\rho^+$.

From Table I, we notice that the directly calculated value without considering FSI for $D^+ \to \bar{K}^{*0}\pi^+$ is lower than the experimental value about 1.5 times while for $D^+ \to \bar{K}^0\rho^+$ it is 1.6~2.5 times larger. In our calculations, three different sets of parameters from Roberts et al.[4] are used for the functions ξ_i (i=1,...,6) in the Gaussian forms (or polynomials). The three sets of parameters are obtained by fitting (1) $D \to K^{(*)}l\nu$; (2) $D \to K^{(*)}l\nu$, $B \to K^{(*)}J/\psi$; and (3) $D \to K^{(*)}l\nu$, $B \to KJ/\psi$, $B \to K\psi'$, respectively. Obviously, the third set of parameters fits more sets of data, and happens to give the best fit (Fit3) to the data in our cases. But there is still discrepancy.

Table II shows that the FSI effects are of similar order of magnitude of the direct production rates. Especially, the amplitude $T(D^+ \to \bar{K}^0\rho^+ \to \bar{K}^{*0}\pi^+)$ is quite large. This will reduce the rate for $\bar{K}^0\rho^+$ final state and raise the rate for $\bar{K}^{*0}\pi^+$ final state, therefore will bring the theoretical results to coincide with the experimental data.

In this work we are not trying to fit the data numerically. We do not introduce any free parameter. We just check what is the theoretical prediction for the processes $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \bar{K}^0\rho^+$, based on the more reliable and widely accepted theoretical framework[4] without considering the FSI. We find that the calculated values without considering the FSI obviously deviate from the experimental data. Then we investigate the magnitude of the FSI effects due to single pion exchange inelastic scattering with the unitarity approximation which has been testified in many practical processes and proved to be reasonable [9, 15]. Our results indicate that the single pion exchange process $\bar{K}^0\rho^+ \leftrightarrow \bar{K}^{*0}\pi^+$ has significant effects to the process $D^+ \to \bar{K}^0\rho^+$ and $D^+ \to \bar{K}^{*0}\pi^+$, and may be the reason for the discrepancy between the experimental data and the theoretical prediction without considering the FSI effects.

We cannot give more accurate results so far, since the disperse part of the inelastic scattering amplitude is very model-dependent, so in this work we are not going to adjust the cut-off parameters for fitting data to cause some mess and uncertainty at this stage.

Our discovery of the significance of the single pion exchange inelastic scattering in FSI may have important applications to many other D and B decays. For example, for $B \to DK$, Zheng [7] concluded that the FSI from the elastic scattering is negligible. It is noted that for the elastic scattering the lightest exchanged mesons are σ or ρ, while the inelastic scattering $D^*K^* \leftrightarrow DK$ can be realized by exchanging pions and may give more significant contributions. The single pion exchange inelastic FSI may play important role in the channels which are relevant to evaluating CP violation in some channels and precise measurement of the CKM matrix entries.

The investigation of those effects is in progress.

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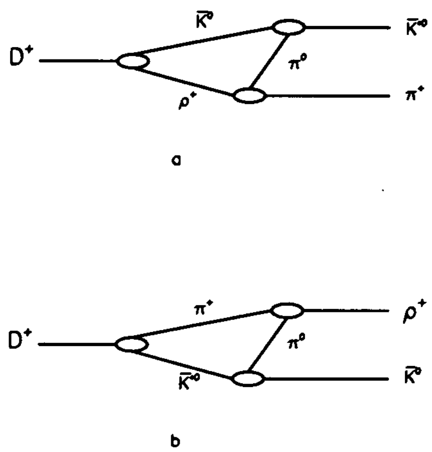

Figure 1: The single pion exchange inelastic FSI loop for $D^+ \to \bar{K}^{*0}\pi^+$ and $D^+ \to \rho^+\bar{K}^0$