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GLUON EFFECTS IN CHARMED MESON WEAK DECAYS

by

V.L.CHERNYAK, A.R.ZHITNITSKY

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V.L.CHERNYAK, A.R.ZHITNITSKY

630090, Novosibirsk, USSR

ABSTRACT

The method for computation of the annihilation contributions into the two-particle decay amplitudes of the D,F-mesons is presented. It is shown that the annihilation contributions play the essential role in these decays. The estimates of the Cabibbo-allowed and the Cabibbo-suppressed decay widths are presented ($D \to K\pi$, $K\rho$, $K^*\pi$, KK, $\pi\pi$...). The results are in correspondence with the available experimental data.

I. Introduction.

The convincing experimental evidence has been obtained at the last years /1/ against the applicability of the standard theoretical scheme to the description of the D,F-meson decays (see the reviews /2-4/). The standard scheme /5/ supposes that the dominant contribution into the charmed particle decays gives the diagram in Fig.1. (Indeed, the contribution of the annihilation diagram in Fig.2 into the D^o -meson decay amplitude includes the additional suppression $\sim \frac{m_s}{m_D} \sim 10^{-1}, m_s \sim 150 MeV$). This supposition leads to the equality of the D^o - and D^+ lifetimes, and this contradicts to the experiment. Besides, the measured branching ratios for D-meson two-particle decays can not be explained within the standard scheme as well /5,6/.

There were a number of suggestions to explain the situation /7-9/. The supposition about the strengthening of the annihilation contribution due to one or few gluons emission (in other words, due to a many-particle component of the D-meson wave function) /7/ is the most fruitful, from our viewpoint. The corresponding estimates within the nonrelativistic quark model show that Fig.3 contribution is indeed of the same order as that of Fig.1 /7a/. The exact calculation of Fig.3 contribution is not possible at present, because the large distance interaction plays an essential role here even in the limit $M_c \to \infty$ and one needs to know the U-quark wave function in D^o-meson, which is unknown. At the same time, only the diagram in Fig.4 gives the main contribution into the D^o-meson two-particle decay amplitudes at $M_c \to \infty$, and the large distance interaction enters through the π (K)-meson wave function only. The main properties of this wave function are known to us /11,12/.

The main goal of this paper is to calculate the two-particle decay widths of the D,F-mesons. The experimental data we have at present there have, from our point of view, no adequate theoretical explanation.

One may think that the annihilation contributions are obviously essential in the two-particle D^0-meson decays if they are essential for the total decay width. We want to stress that this is not the case in general, and it can not be excluded beforehand that the annihilation contributions play no essential role in the two-particle decays.* Indeed, the naive estimate leads just to such conclusion. Using the dimensional considerations one obtains for the Fig.1 contribution into the $D^0 \to K^-\pi^+$ decay amplitude: $M_{dir} \sim G f_\pi m_D^2$, where G is the Fermi constant, $m_D \simeq 1.86\ GeV$ is the D-meson mass and $f_\pi \simeq 133\ MeV$ characterizes the value of the π meson wave function. The analogous estimate of the annihilation contribution includes the wave functions of all three mesons and has the form: $M_{ann} \sim G f_\pi f_K f_D$; $f_K \simeq 1.25 f_\pi$; $f_D \sim 1.5 f_\pi$. Therefore, $M_{ann}/M_{dir} \sim f_D f_K/m_D^2 \sim 10^{-2}$, and it seems that the annihilation contribution plays no role in the two-particle decays.

The main result of the present paper is that we show that the annihilation contributions are not small really as compared with the direct ones. The naive estimate $M_{ann}/M_{dir} = C_0 \frac{f_D f_K}{m_D^2}$, $C_0 \sim 1$ is wrong. The coefficient $C_0 \sim 100$ and compensates

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* This corresponds to the case when the "$\bar{u}c\bar{d}$"-state in Fig.3 transforms mainly into the hadron states with more than two particles.

for the smallness of the factor $f_D f_K/m_D^2 \sim 10^{-2}$. The reason is as follows. The dimensional estimate implies that the virtualities of the propagators in Fig.4 are of the order $\bar{q}_i^2 \sim m_D^2$. (When each quark in the π (K)-meson carries one half of the meson momentum, then $\bar{q}^2 \simeq m_D^2/4$). This is not the case really. The π-meson wave function $\varphi_\pi(\xi)$ which determine the distribution of the quark longitudinal momentum is such that π-meson momentum is devided between the quarks in proportion ~85%:15% /11,12/. Therefore, $\bar{q}^2 \ll m_D^2/4$ really. As a result, more accurate calculation of the annihilation contributions shows that they are not small (see Sect.III)*.

Moreover, the sign of the annihilation contribution relative to the direct one is determined unambigiously in each given decay. The determination of the relative sign of these two contributions is an important element of the calculation because they are of the same order. We show that in all cases for which the experimental data are available this interference works in the right direction.

In Section II we discuss the questions connected with the calculation of the diagram in Fig.5 contribution. The decay $D^+ \to \bar{K}^0\pi^+$ is chosen as the standard one because the annihilation contributions are absent here. In Section III we describe the calculation of the annihilation contributions (Fig.4). The method of calculation we apply here is analogous to this used for the description of the strong and electro-

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* At $M_c \to \infty$ the annihilation contributions die off, of course, but at real value $M_c \simeq 1.5\ GeV$ they are still important.

magnetic decays of the charmonium /10,12/. The annihilation amplitude contribution is expressed through the π - and K-meson wave functions and the constant f_D (which is analogous to f_π). The main properties of the π (K)-meson wave function were found by us in the previous paper /11/ and allow to obtain the good description of the charmonium decays: $\chi_0(3415), \chi_2(3555), \psi(3100), \psi'(3685) \to \pi^+\pi^-, K^+K^-$ /12/. So we expect that our approach allows to obtain the reliable estimate of the annihilation contributions into the charmed meson decays as well.

In Section IV we consider the Cabibbo-suppressed decays of the D-mesons. In particular, we argue that $D^+ \to \bar{K}^0 K^+ / D^+ \to \bar{K}^0 \pi^+ > tg^2\theta$ mainly due to the annihilation contribution. The D-meson decays into the vector and pseudoscalar particles are considered in Sect.V. In Sect.VI we discuss the F-meson decays. Our results are presented in Tables I,II.

II. $D \to PP$ decays. Direct contributions.

We consider here the decays into two pseudoscalar particles. The weak Hamiltonian has the form: /13,14/

$$H = \frac{G}{\sqrt{2}}\left[c_1\ \bar{s}'\gamma_\mu(1+\gamma_5)c \cdot \bar{u}\gamma_\mu(1+\gamma_5)d' + c_2\ \bar{s}'\gamma_\mu(1+\gamma_5)d' \cdot \bar{u}\gamma_\mu(1+\gamma_5)c\right], \tag{1}$$

$$d' = d\cos\theta + s\cdot\sin\theta \quad , \quad s' = s\cdot\cos\theta - d\sin\theta$$

$$c_1 = \tfrac{1}{2}(c_+ + c_-) \quad , \quad c_2 = \tfrac{1}{2}(c_+ - c_-).$$

The color summation is implied in (1). We use:

$$\alpha_s(\mu^2) = \frac{4\pi}{b}\frac{1}{\ln \mu^2/\Lambda^2} \quad , \quad \Lambda = 100\ MeV \, , \quad b = 11 - \tfrac{2}{3}N_f . \tag{2}$$

So:

$$c_- = \left[\frac{\alpha_s(M_c)}{\alpha_s(M_W)}\right]^{4/b} \simeq 1.48 \quad , c_+ = 1/\sqrt{c_-} \simeq 0.82, \ c_1 \simeq 1.15, \quad c_2 \simeq -0.33. \tag{3}$$

It is worth noting that the values $c_- \simeq 2.2$ and $c_+ \simeq 0.7$ corresponding to $\Lambda \simeq 550\ MeV$ are used frequently when considering the weak decays. The value $\Lambda \simeq 100\ MeV$ is preferable, from our point of view, as dictated by the charmonium sum rules /15a /. In any case, small changes in the $c_\pm$ -values are not important for the most decay modes.

When calculating the direct contributions like those shown in Fig.5 we follow the standard scheme /15b / in which the factorization of the matrix elements according to the valence quark model is performed.

The $D^+ \to \bar{K}^0\pi^+$ decay amplitude has the form (see Fig.5):

$$\begin{aligned} M(D^+ \to \bar{K}^0\pi^+) = \frac{G}{\sqrt{2}}\cos^2\theta \Big\{ & c_1 \langle \bar{K}^0 | \bar{s}^i\gamma_\mu(1+\gamma_5)c^i | D^+ \rangle \langle \pi^+ | \bar{u}^i\gamma_\mu(1+\gamma_5)d^i | 0 \rangle + \\ + & c_1 \langle \pi^+ | \bar{u}^i\gamma_\mu(1+\gamma_5)c^i | D^+ \rangle \langle \bar{K}^0 | \bar{s}^i\gamma_\mu(1+\gamma_5)d^i | 0 \rangle + \\ + & c_2 \langle \bar{K}^0 | \bar{s}^i\gamma_\mu(1+\gamma_5)c^j | D^+ \rangle \langle \pi^+ | \bar{u}^j\gamma_\mu(1+\gamma_5)d^i | 0 \rangle + \\ + & c_2 \langle \pi^+ | \bar{u}^j\gamma_\mu(1+\gamma_5)c^i | D^+ \rangle \langle \bar{K}^0 | \bar{s}^i\gamma_\mu(1+\gamma_5)d^j | 0 \rangle \Big\} = \\ = \frac{G}{\sqrt{2}}\cos^2\theta \Big\{ & (c_1 + \tfrac{1}{3}c_2) f_+^K \cdot f_\pi + (c_2 + \tfrac{1}{3}c_1) f_+^\pi \cdot f_K \Big\} (-i m_D^2) , \end{aligned} \tag{4}$$

$$\langle \pi^+(p^\pi) | \bar{u}^i\gamma_\mu\gamma_5 d^j | 0 \rangle = -\tfrac{i}{3}\delta^{ij} f_\pi\, p_\mu^\pi$$

$$\langle \bar{K}^0(p^K) | \bar{s}^i\gamma_\mu c^j | D^+(p) \rangle = \tfrac{1}{3}\delta^{ij}\left[f_+^K (p+p^K)_\mu + f_-^K (p-p^K)_\mu \right].$$

Here: $f_\pi \simeq 133\,MeV$, $f_K = 1.23 f_\pi$. Everywhere we neglect the SU(3)-symmetry breaking effects and small corrections due to a light hadron masses. So:

$$M(D^+ \to \bar{K}^0\pi^+) \simeq -i m_D^2 \frac{6}{\sqrt{2}} \cos^2\theta\, f_+(0) f_\pi \cdot \frac{4}{3}(c_1 + c_2). \tag{5}$$

The form factor $f_+(0) = 1$ in the exact SU(4)-symmetry limit. The SU(4)-symmetry is badly broken in reality, and so we expect $f_+(0)$ to be considerably smaller than unity. The estimates indicate that $f_+(0) \simeq 0.5$ / 16 /. We give the rough estimate here which corfirms that $f_+(0)$ can be considerably smaller than unity. Let us consider the form factor (see (4)):

$$(P-P^K)_\mu \langle \bar{K}^0(P^K) | \bar{s}\gamma_\mu c | D^+(P)\rangle \equiv m_D^2 \Phi(q^2),\ q = P - P^K,\ \Phi(0) = f_+(0).$$

In the soft K-meson limit we have:

$$\Phi(q^2 = m_D^2) = \frac{f_D}{f_K}, \quad \langle 0 | \bar{d}\gamma_\mu\gamma_5 c | D^+(P)\rangle = i P_\mu f_D \tag{6}$$

(The value f_D obtained by using the QCD-sum rules / 17 / is: $f_D \simeq 200\,MeV$; we use this value everywhere below). Let us approximate the form factor $\Phi(q^2)$ by the nearest scalar resonance contribution:

$$\Phi(q^2) = \frac{C_0}{M_S^2 - q^2}, \quad C_0 \simeq \frac{f_D}{f_K}(M_S^2 - M_D^2). \tag{7}$$

The mass difference between the K-meson and the $K(1420)$-scalar meson is ~ 900MeV. We expect the corresponding difference in the charmed meson case to be somewhat smaller (K-meson is too light): $M_S \simeq M_D + 700\,MeV \simeq 2.6\,GeV$. Then:

$$f_+(0) = \Phi(0) \simeq 0.6. \tag{8}$$

If we use the experimental data / 1 / for the branching ratio: $Br(D^+ \to \bar{K}^0\pi^+) \simeq 2\%$ and for the total lifetime: $\tau_{D^+} \simeq (4 \div 10) \cdot 10^{-13} sec$ then we obtain from (5),(3):

$$f_+(0) \simeq 0.25 \div 0.4. \tag{9}$$

Therefore, all considerations show that $f_+(0)$ is smaller then unity by a factor ~ 2-3. In what follows we use the value $f_+(0) \simeq 0.4$ in the numerical estimates.

III. $D \to PP$. Annihilation contributions.

The corresponding diagrams are shown in Figs. 4a-d. The method for the calculation of diagrams like those in Fig. 4 is described in / 10 / and in our previous paper / 12 /. We discuss only the results here.* The Figs. 4c,d contributions have the form:

$$M_c + M_d \sim \langle 0 | A_\mu(0) | D^0(P)\rangle \langle \pi^+(P_1) K^-(P_2) | V_\mu(0) | 0\rangle \sim$$

$$\sim f_D (m_s - m_u) \langle \pi^+(P_1) K^-(P_2) | S(0) = \bar{s}(0) d(0) | 0\rangle \sim f_D (m_s - m_u) \mu^3 / m_D^2$$

(The form factor $\langle \pi^+ K^- | S | 0\rangle \sim \mu^3/m_D^2$, where μ is some characteristic mass). These contributions are small, and we neglect them in what follows. The Fig. 4a contribution is:

$$M_{ann}(D^0 \to K^-\pi^+) = -i \frac{32}{27} \pi \alpha_s \cos^2\theta \frac{6}{\sqrt{2}} f_D f_\pi f_K \cdot c_1 \cdot I_a, \tag{10}$$

$$I_a \equiv I_a^{(1)} \cdot I_a^{(2)} = \int_{-1}^{1} \frac{d\xi_1 \varphi(\xi_1)}{(1-\xi_1)^2} \cdot \int_{-1}^{1} \frac{d\xi_2 \varphi(\xi_2)}{(1-\xi_2)}.$$

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* Let us point out only that virtualities of all the propagators in Fig. 4 are $\sim O(M_c^2)$ at $M_c \to \infty$. We neglect the light quark momentum in comparison with M_c (in the D-meson rest frame).

The Fig.4b contribution has the form (10) with the replacement $I_a \to I_b$ *:

$$I_b = \int_{-1}^{1} \frac{d\xi_2\, \varphi(\xi_2)}{(1-\xi_2)} \cdot \int_{-1}^{1} \frac{d\xi_1\, \varphi(\xi_1)}{[4-(1+\xi_1)(1+\xi_2)]} \tag{11}$$

Here $\varphi(\xi, \mu^2)$ is the π (K)-meson wave function normalized by the condition: $\int_{-1}^{1} d\xi\, \varphi(\xi, \mu^2) = 1.$ The wave function $\varphi(\xi, \mu^2)$ has the meaning of the decay amplitude into two quarks with the longitudinal momentum fractions X and 1-X, $\xi = x-(1-x) = 2x-1$ and with virtualities up to μ^2 (for more details see /11 /).

We estimate now the annihilation contributions using the model wave function proposed in the paper /11 /:

$$\varphi_\pi(\xi, \mu_0^2 \simeq (500\,MeV)^2) \simeq \varphi_K(\xi, \mu_0^2 \simeq (500\,MeV)^2) = \frac{15}{4}\xi^2(1-\xi^2). \tag{12}$$

The wave function (12) satisfies the QCD-sum rules /11 / and leads to the predictions for the charmonium $\pi^+\pi^-$ -decay widths in agreement with the experimental data /12 /. The value $\mu_0^2 \simeq (500\,MeV)^2$ was the characteristic one for the charmonium decays ($M \simeq 3 \div 3.5\,GeV$) /12 /. As the D-meson mass m_D is: $m_D \simeq 1865\,MeV$, the characteristic virtuality of the π (K)-meson wave functions and the propagators in Fig.4 are smaller (at $|\xi_{1,2}| \simeq 0.7$, $q^2 \simeq \frac{1-\xi_1}{2} \cdot \frac{1-\xi_2}{2} m_D^2 \simeq (300\,MeV)^2$). So: $\alpha_s = \alpha_s(\mu^2 \simeq (300\,MeV)^2) \simeq 0.6$ (see (2)) in formulae (10), (11). For the wave function (12) the integral $I_a^{(2)} \simeq 2.5$, the integral $I_b \simeq 1.5$ and the integral $I_a^{(1)}$ is logarithmi-

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* The leading at $M_c \to \infty$ contributions are given here only. We neglect temporarily power corrections (see Sect.VII).

cally divergent. We, therefore, estimate with the logarithmic accuracy*:

$$I_a^{(1)} \simeq \frac{15}{2} \ln\left(\frac{1}{1-\xi_{max}}\right) \simeq 10. \tag{13}$$

It is seen that the Fig.4b contribution is small as compared with that of Fig.4a. This can be expected beforehand as the gluon is radiated from the heavy C-quark line.

Taking the expressions (5),(10) as the characteristic values of the direct and annihilation contributions respectively, we have for their ratio ($f_D \simeq 200\,MeV$, $f_K \simeq 165\,MeV$, $f_+ \simeq 0.4$, $I_a \simeq 25$, $\alpha_s \simeq 0.6$):

$$\rho^{DD} = \frac{(annihilation)_{DD}}{(direct)_{DD}} = \frac{32}{27}\pi\alpha_s \frac{f_D f_K}{m_D^2 f_+(0)} I_a \simeq 1.3. \tag{14}$$

Therefore, it is seen, that the annihilation contribution is not small as compared with the direct one, inspite of the presence of the suppression factor $\sim f_D f_K / m_D^2 \sim 10^{-2}$. The reason is clear from the previous considerations: the π (K)-meson wave function is wide enough (see (12)) and so the propagator virtualities in Fig.4 are much smaller than m_D^2 (i.e. the integral I_a is large). For the same reason the coupling constant α_s in (10) is not very small.

We realize that the above calculation of the annihilation contribution is not very precize. (We expect it is accurate within a factor $\sim$ 2). The D-meson mass is not very large really, the π (K)-meson wave function is wide, and so the pro-

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* Really, the values of the integrals I_i are somewhat larger because the wave function (12) should be renormalized to the point $\mu^2 \simeq (300\,MeV)^2$. We neglect this effect.

pagator virtualities in Fig.4 are not large. So,it is natural to expect that the power corrections are more important here as compared with the case of the charmonium decays. We are certain however,that the wave function (12) reproduces correctly the main characteristic properties of the π (K)-meson wave function. In particular,the integral I appearing in the charmonium decays is equal for this wave function /12 /:

$$I = \int_{-1}^{1} d\xi_1\, d\xi_2\, \frac{\varphi_\pi(\xi_1,\mu_0^2)}{1-\xi_1^2} \cdot \frac{\varphi_\pi(\xi_2,\mu_0^2)}{1-\xi_2^2}\, \frac{1}{1-\xi_1^2\xi_2^2} \simeq 13, \tag{15}$$

and this value agrees well with the experimental data for the charmonium decays /12 /. Comparison of the value of the integral (15) with the above given estimate

$$I_a = \int_{-1}^{1} \frac{d\xi_1 (1+\xi_1^2)\varphi_\pi(\xi_1)}{(1-\xi_1^2)^2} \cdot \int_{-1}^{1} \frac{d\xi_2\, \varphi_\pi(\xi_2)}{1-\xi_2^2} \simeq 25 \simeq 2 \cdot I \tag{16}$$

shows,from our viewpoint,that the estimate (16) is reasonable, because I_a is certainly larger than I when the wave function is wide.

On the whole,we think that we have given strong enough arguments that

$$\rho^{PP} = \text{(annihilation)/(direct)} \simeq 1 \tag{17}$$

(within a factor ~ 2). This conclusion is one of the main results of this paper. It shows that the annihilation contributions play an important role in the charmed meson exclusive decays.

What remains to be done is the account for the final state interaction. At $M_c \to \infty$ the final state interaction plays

no role due to a factorization (analogously to the case of form factor asymptotic behaviour /10 /). The D-meson mass is, however,not very large,and so there are the (scalar) resonances close to it. With the final state interaction taken into account,the $D^0 \to K^-\pi^+$, $D^0 \to \bar{K}^0\pi^0$ decay amplitudes have the form:

$$M(D^0 \to K^-\pi^+) = M_0 \left[\tfrac{4}{9}(c_1+c_2) e^{i\delta_{3/2}} + \left(\tfrac{5}{9}c_1 - \tfrac{4}{9}c_2 + c_1\rho^{PP}\right) e^{i\delta_{1/2}} \right], \tag{18}$$

$$M(D^0 \to \bar{K}^0\pi^0) = \frac{M_0}{\sqrt{2}} \left[\tfrac{8}{9}(c_1+c_2) e^{i\delta_{3/2}} + \left(-\tfrac{5}{9}c_1 + \tfrac{4}{9}c_2 - c_1\rho^{PP}\right) e^{i\delta_{1/2}} \right], \tag{19}$$

where

$$\rho^{PP} = M_{ann}/M_0 \quad , \quad M_0 = -i\frac{G}{\sqrt{2}} \cos^2\theta\, m_D^2\, f_\pi\, f_+(0) \tag{20}$$

and $\delta_{1/2}$, $\delta_{3/2}$ are the S-wave $K\pi$ -scattering phases in states with the total isospin I=1/2, I=3/2 respectively. In this case these phases are well known /18,19/: $\delta_{1/2} \simeq 90^\circ$, $\delta_{3/2} \simeq -30^\circ$. (There is the scalar resonance $K'(1900)$ close to the D-meson,/ 19 /). So,for the values: $\rho^{PP} \simeq 1.3$ (see (14)), $c_1 \simeq 1.15$, $c_2 \simeq -0.53$ (see (3)) we have

$$\frac{\Gamma(\bar{K}^0\pi^0)}{\Gamma(K^-\pi^+)} \simeq 0.9 \quad \left(0.7^{+0.5}_{-0.3} \quad /11/ \right)$$

$$\frac{\Gamma(D^+ \to \bar{K}^0\pi^+)}{\Gamma(D^0 \to K^-\pi^+) + \Gamma(D^0 \to \bar{K}^0\pi^0)} \simeq 0.16 \quad \left((0.4 \pm 0.2)\frac{\tau_0}{\tau_+} \quad /11/ \right). \tag{21}$$

The corresponding experimental values are given in parenthesis. The fact that the final state interaction can play the essential role in given decay channels has been pointed out earlier /8/. However,it seems impossible to obtain simultaneously $\Gamma(\bar K^0\pi^+)/\Gamma(K^-\pi^+)\sim 1$ and $\Gamma(\bar K^0\pi^0)/[\Gamma(\bar K^0\pi^0)+\Gamma(K^-\pi^+)]\ll 1$ by phases choice if the annihilation amplitude $\rho^{PP}=0$. (Besides, at $\rho^{PP}=0$ the ratio is very sensitive to the precise phase values, while in our case the sensitivity is small).

Let us point in conclusion of this Section that the decay mode $D^0\to\bar K^0\eta'$ is strengthened significantly due to the annihilation contribution. We expect: $\Gamma(D^0\to\bar K^0\eta')/\Gamma(K^-\pi^+)\approx 0.7$.

IV. $D\to PV$ decays.

In this Section we consider the D-meson decays into pseudoscalar and vector particles. The vector meson is longitudinally polarized in the $D\to PV$ decays, i.e. its helicity $\lambda=0$. With the good enough accuracy one can replace: $m_V e_\mu^{\lambda=0}(p)\simeq p_\mu$, where e_μ is the polarization vector, m_V is the vector particle mass. So, the kinematics of the $D\to PV$ and $D\to PP$ decays become identical.

The new additional form factor $g_+(0)$ appears in the direct contributions here: $q_\mu\langle V(p_V)|A_\mu(0)|D(p)\rangle = m_D^2 g_+(0)$, $q=p-p_V$. The matrix element $\langle\rho(p_V)|V_\mu(0)|0\rangle = f_\rho m_\rho e_\mu^{\lambda=0}(p_V)\simeq f_\rho p_\mu^V$, $f_\rho\approx 200\,\mathrm{MeV}$ is analogous to that

$\langle\pi^+(p_2)|A_\mu(0)|0\rangle$.

(For other particles in the vector nonet the matrix elements are determined by the SU(3)-symmetry). Applying PCAC to the D-meson we can determine the sign: $g_+(0)>0$. Two types of integrals appear now in the annihilation contributions (see (10),(14)),Fig.4a:

$$\rho^{VP}\to\int d\xi_1 d\xi_2\,\frac{\varphi_V(\xi_1)}{(1-\xi_1)^2}\cdot\frac{\varphi_P(\xi_2)}{(1-\xi_2)}\,;\quad \rho^{PV}\to\int d\xi_1 d\xi_2\,\frac{\varphi_P(\xi_1)}{(1-\xi_1)^2}\cdot\frac{\varphi_V(\xi_2)}{(1-\xi_2)}\,.$$

The Figs.4b,c,d contributions are small for the same reasons as in Sect.III).

We have at present no detailed information about the ρ-meson wave function $\varphi_\rho(\xi)$. We have no doubts, however, that its properties are similar to those of $\varphi_\pi(\xi)$. The preliminary investigation of the QCD-sum rules for the $\varphi_\rho(\xi)$, $\int d\xi\,\varphi_\rho(\xi)=1$ (analogous to those for $\varphi_\pi(\xi)$, see /11/) shows that $\varphi_\rho(\xi)$ is a little more narrow than $\varphi_\pi(\xi)$. So, the integral like that of I_a in (10) will be somewhat smaller than I_a (16), and this will work in an opposite direction as compared with the effect due to $f_\rho>f_\pi$.

In order to get semiquantitative impression about the decay properties, one can put all the corresponding quantities for the P and V-mesons to be the same: $\rho^{PV}\simeq\rho^{VP}\simeq\rho^{PP}$, $f_\rho f_+(0)\simeq f_\pi g_+(0)$. It is easy to see that even such very rough approach does not contradict to the experimental data available as the latter include large uncertainties.

The value $g_+(0)$ can be obtained from the $D^+\to\bar K^{0*}\pi^+$ decay. The latter is not measured, however, at present. The only limitation is that the $\bar K^{0*}\pi^+$ mode constitutes less than 15% of the decay $D^+\to K^-\pi^+\pi^+$ /20/. This gives the limitation: $\beta = f_\pi g_+(0)/f_\rho f_+(0)\lesssim 0.7$. Let us note also that the final state interaction effects can be neglected in the $D\to PV$-decays because there are no known pseudoscalar resonances. For illustration in the Table II the ratios $\Gamma(D\to PV)/\Gamma(D^+\to\bar K^0\rho^+)$ * are given for the case: $\beta=0.7$, $\rho_{PV}\simeq\rho_{VP}\simeq 0.8$. In particular,

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* The $D^+\to\bar K^0\rho^+$ is chosen as the normalizing one because the annihilation contribution is absent here, the coefficient β enters with the small factor $c_2+\frac{1}{3}c_1$ and, moreover, the final state isospin is equal to 3/2.

we have the following relations:

$$\frac{\Gamma(D^0\to \bar{K}^{0*}\pi^0)}{\Gamma(D^0\to K^{-*}\pi^+)} \simeq 0.2\left(0.4^{+0.7}_{-0.2}\ /11/\right), \quad \frac{\Gamma(D^0\to K^{-*}\pi^+)}{\Gamma(D^0\to K^-\rho^+)} \simeq 0.7\left(0.44\pm0.29\ /11/\right)$$

$$\frac{\Gamma(D^0\to \bar{K}^0\rho^0)}{\Gamma(D^0\to K^-\rho^+)} \simeq 0.1\left(0.01^{+0.06}_{-0.01}\ /11/\right), \quad \frac{\Gamma(D^+\to \bar{K}^0\rho^+)}{\Gamma(D^0\to K^-\rho^+)+\Gamma(D^0\to \bar{K}^0\rho^0)} \simeq 0.25.$$

Let us also point out that we expect the branching ratio $\sim 1\%$ for the $D^0\to \bar{K}^0\varphi$ decay. This decay is interesting because it is due to the annihilation contribution only.

V. Cabibbo-suppressed decays.

Let us reminde that in this case (i.e. $\sim \sin\theta\cos\theta$) the annihilation contributions are present not only in the D^0 decays, but in the D^+-decays as well.* Let us first consider the latter case because the final state interaction is of no importance there.

The annihilation contributions increase somewhat $D^+\to \bar{K}^0K^+$, KK^* decays. In particular, we have (see Tables I, II):

$$\frac{\Gamma(D^+\to \bar{K}^0K^+)}{\Gamma(D^+\to \bar{K}^0\pi^+)} = 1.8\, tg^2\theta\left(0.25\pm0.15\ /21/\right), \quad \frac{\Gamma(D^+\to \pi^+\pi^0)}{\Gamma(D^+\to \bar{K}^0\pi^+)} \simeq 0.5\, tg^2\theta$$

$$\frac{\Gamma(D^+\to \bar{K}^0K^{+*})}{\Gamma(D^+\to \bar{K}^0\rho^+)} \simeq 1.6\, tg^2\theta, \qquad \frac{\Gamma(D^+\to \bar{K}^{0*}K^+)}{\Gamma(D^+\to \bar{K}^{0*}\pi^+)} \simeq 1.8\, tg^2\theta \tag{22a}$$

---

* We suppose that the "penguin" contribution can be neglected for the decays under considerations.

As for the D^0 -decays, the annihilation contributions are present in both Cabibbo-suppressed and allowed decays, so the ratios are

$$\frac{\Gamma(D^0\to K^{-*}K^+)}{\Gamma(D^0\to K^{-*}\pi^+)} \simeq \frac{\Gamma(D^0\to \rho^-\pi^+)}{\Gamma(D^0\to K^{-*}\pi^+)} \simeq tg^2\theta \tag{22b}$$

$$\frac{\Gamma(D^0\to K^-K^{+*})}{\Gamma(D^0\to K^-\rho^+)} \simeq \frac{\Gamma(D^0\to \pi^-\rho^+)}{\Gamma(D^0\to K^-\rho^+)} \simeq tg^2\theta.$$

The final state interactions do not influence (22a) and, it seems, can be neglected in (22b). This is not the case for the $D^0\to\pi\pi$, $K\bar{K}$ decays, where the S-wave resonance structure is present. Note also that the annihilation contribution is absent in the $D^+\to\pi^+\pi^0$ decay as there $\Delta I = 3/2$. Therefore, the prediction is the standard one here.

In conclusion of this Section let us note the following. We do not expect that the SU(3)-symmetry breaking effects lead to the significant increase of the ratio $D\to KK/D\to\pi\pi$ in contrast with that proposed in /2 /. The reason is that the symmetry breaking not only increases $f_\pi\to f_K\simeq 1.25 f_\pi$, but simultaneously changes the wave function shape, $\varphi_\pi(\xi)\neq\varphi_K(\xi)$. These effects tend to compensate each other, as can be seen from the experimental data for the charmonium decays /21/:

$$Br(\chi_0(3415)\to\pi^+\pi^-) = Br(\chi_0(3415)\to K^+K^-) = (1.0\pm 0.3)\%.$$

(The effect entirely due to $f_K\neq f_\pi$ give:

$$(\chi_0\to K^+K^-/\chi_0\to\pi^+\pi^-) = (f_K/f_\pi)^4 \simeq 2.5.$$

Therefore, we expect that the role of the symmetry-breaking effects is reasonably small in the ratios (22) and the main effects are due to annihilation contributions.

VI. F-meson decays.

The annihilation contributions in the F-meson decays are proportional to the coefficient C_2 and so are ~ 3.5 times smaller as compared with the effects due to C_1 in D^0-decays. (The coefficient C_2 is equal to zero in the absence of the QCD logarithmic corrections, see (3)). Therefore, the interference of the direct and annihilation contributions is smaller here as well as compared with that of the D-decays. The most clear manifestation of the annihilation contribution is the existence of the $F^+ \to \pi^+ \omega$ decay. We expect (see TableII):

$$\frac{\Gamma(F^+ \to \pi^+ \omega)}{\Gamma(F^+ \to \pi^+ \varphi)} \gtrsim 0.3 .$$

Let us point also some increase in the widths of $\pi^+ \xi,\ K^+ \bar{K}^0,\ \rho^+ \xi$ decays due to the annihilation contributions.

We expect the following ratios for the Cabibbo-suppessed decays:

$$\frac{\Gamma(F^+ \to K^+ \pi^0)}{\Gamma(F^+ \to K^0 \pi^+)} \simeq 1.3 \quad ; \quad \frac{\Gamma(F^+ \to \bar{K}^{0*} \pi^+)}{\Gamma(F^+ \to K^+ \bar{K}^{0*})} \sim 15\, tg^2\theta \qquad (23)$$

$$\frac{\Gamma(F^+ \to K^0 \rho^+)}{\Gamma(F^+ \to K^+ \bar{K}^{0*})} \sim 26\, tg^2\theta \sim 1 ; \quad \frac{\Gamma(F^+ \to K^+ \pi^0)}{\Gamma(F^+ \to K^+ \bar{K}^0)} \sim 6\, tg^2\theta .$$

The first ratio in (23) is analogous to $\Gamma(D^0 \to \bar{K}^0 \pi^0)/\Gamma(D^0 \to K^- \pi^+)$. The large ratios for the rest terms in (23) are due to the suppression of the Cabibbo-allowed decays.

VII. Phenomenology.

In order to reduce the role of uncertainties in the values of the direct and annihilation contributions to $D \to PP$

decays*, the following method is used. We suppose that the flavor structure of the annihilation amplitude is the same as that in Fig.4 and use the SU(3)-symmetry for calculation of the matrix elements in the direct and annihilation contributions. The quantity ρ^{PP} is considered now as a free parameter (with the limitation: $\rho^{PP} \sim 1$ within a factor ~ 2.). In this way we take into account the greater part of neglected before power corrections present in Fig.4, Fig.6 and many other power corrections. The above described suggestion is, from our point of view, weak enough and should be valid with a good accuracy. In the preceding Sections we have given rough estimates of the decay widths by choosing $\rho^{PP} \simeq 1.3$ (see (14)). According to the considerations described above in this Section, there is an optimal value of ρ^{PP} (not differing greatly from $\rho^{PP} \sim 1$) which allows much more precise predictions for the decay widths. (We expect the accuracy $\sim$ 20-30% corresponding to the SU(3)-symmetry accuracy). We do not try to find this optimal ρ^{PP} value here as the experimental data have large errors at present.

Analogously, there are the optimal values of the parameters $\beta, \rho^{PV}, \rho^{VP}$ (all O(1)) in the VP-decays which describe the experimental data with much better accuracy as compared with our rough estimates given in the text.

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* The contributions of the type shown in Fig.6 which are connected with the 3-particle light meson wave function, can also be important in the annihilation amplitude. Their contribution is suppressed by a factor $\sim \mu^2/m_c^2$ at $m_c \to \infty$ as compared with that in Fig.4. However, the C-quark mass is not very large really.

The expressions for the decay amplitudes in terms of the parameters $\rho^{PP}, \rho^{PV}, \rho^{VP}, \beta$ are presented in Tables I,II.

VIII. Conclusions.

As has been shown in the preceding Sections, the naive dimensional estimate $\rho = M_{ann}/M_{dir} \sim f_D f_K/m_D^2 \sim 10^{-2}$ is wrong, and in fact $\rho \simeq 1$. So, the annihilation contributions are not small in comparison with the direct ones and play an essential role in the nonleptonic decays of the charmed particles. The account of the annihilation contributions leads in many cases to the predictions which differ greatly from those of the standard scheme in which only the direct contributions are taken into account. In allcases our results are in better agreement with the experimental data as compared with the standard scheme. In particular, we have:

1. $$\frac{\Gamma(D^+ \to \bar{K}^0\pi^+)}{\Gamma(D^0 \to K^-\pi^+) + \Gamma(D^0 \to \bar{K}^0\pi^0)} \simeq 0.16 \left[(0.4 \pm 0.2)\frac{\tau_0}{\tau_+}\ /11/\right]$$

2. $$\frac{\Gamma(D^0 \to \bar{K}^0\pi^0)}{\Gamma(D^0 \to K^-\pi^+)} \simeq 0.9 \left[0.7^{+0.5}_{-0.3}\ /11/\right]$$

3. $$\frac{\Gamma(D^0 \to \bar{K}^{0*}\pi^0)}{\Gamma(D^0 \to K^{-*}\pi^+)} \simeq 0.2 \left[0.4^{+0.7}_{-0.2}\ /11/\right]$$

4. $$\frac{\Gamma(D^+ \to \bar{K}^{0*}\pi^+)}{\Gamma(D^0 \to K^{-*}\pi^+) + \Gamma(D^0 \to \bar{K}^{0*}\pi^0)} \simeq 0.17 \left[< 0.8\frac{\tau_0}{\tau_+}\ /11/\right]$$

5. $$\frac{\Gamma(D^0 \to K^{-*}\pi^+)}{\Gamma(D^0 \to K^-\rho^+)} \simeq 0.7 \left[0.44 \pm 0.29\ /11/\right]$$

6. $$\frac{\Gamma(D^+ \to \bar{K}^0K^+)}{\Gamma(D^0 \to \bar{K}^0\pi^+)} \simeq 1.8\, tg^2\theta \sim 0.1 \left[0.24 \pm 0.16\ /11/\right].$$

As a whole, we have at present, from our point of view, the clear enough understanding of all the main properties of the charmed meson two-particle decays.

To verify the scheme it is important to measure the branching ratios for the $D^0 \to \bar{K}^0\varphi$, $F^+ \to \pi^+\omega$ decays which are entirely due to the annihilation contributions. Besides, it is useful to measure the $D^+ \to \bar{K}^{0*}\pi^+$ decay width as it allows determination of the value of the parameter β.

In conclusion let us note that the annihilation contributions like those described above can play the role in the K-meson nonleptonic decays as well. Unfortunately, the reliable calculation can not be done here at present as all the masses are small.

We are indebted to A.I.Vainshtein for useful discussions.

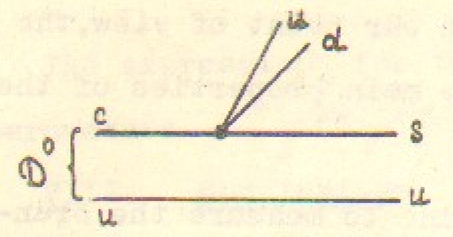

Fig.1

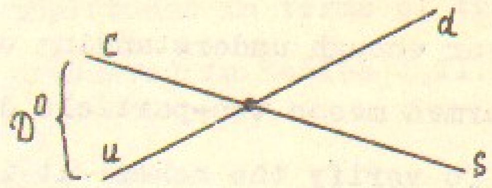

Fig.2

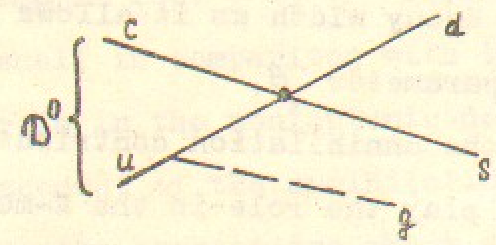

Fig.3

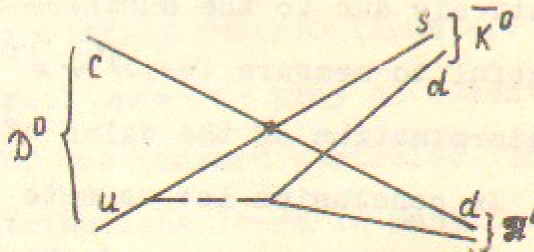

Fig.4a

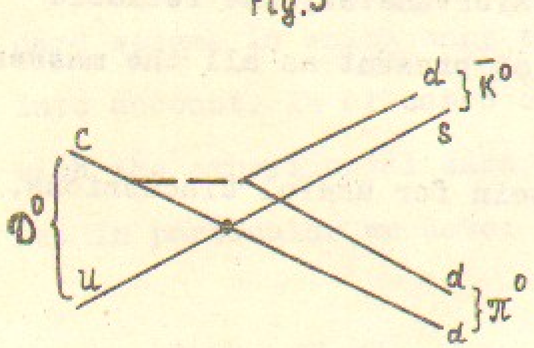

Fig.4в

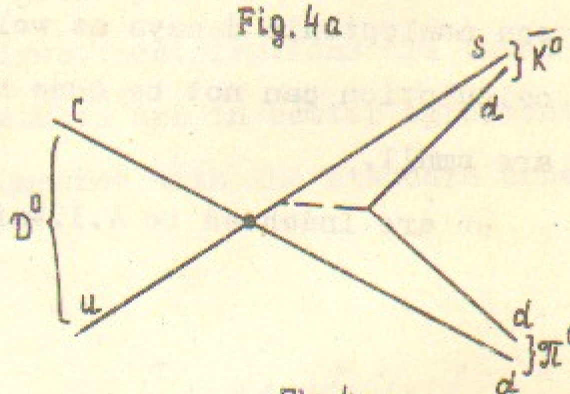

Fig.4c

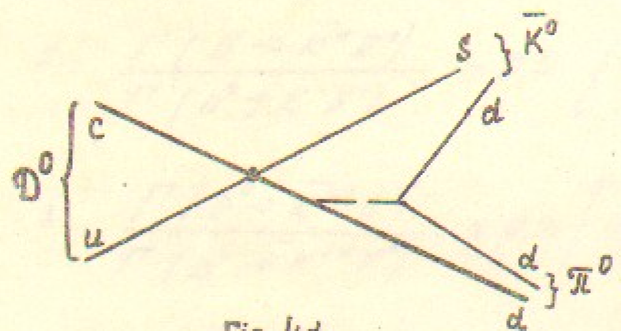

Fig.4d

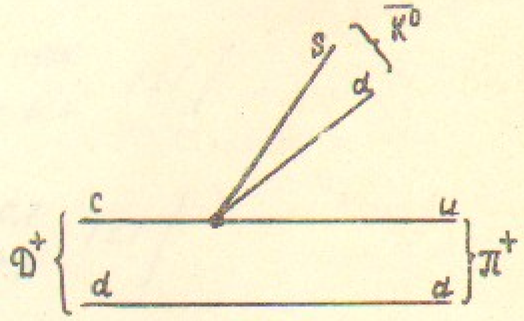

Fig.5

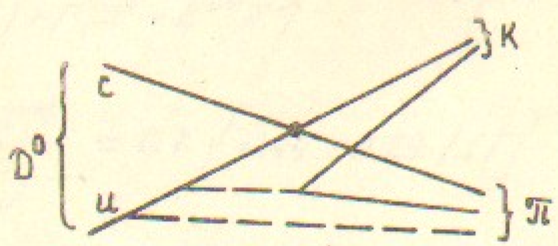

Fig.6

TABLE I

MODE	AMPLITUDE	$\frac{\Gamma(DF \to PP)}{\Gamma(D^+ \to \bar{K}^0\pi^+)}$ $c_1 = 1.15$, $c_2 = -0.33$	
	$-i\frac{G_F}{\sqrt{2}} m_D^2 f_\pi f_+ \cdot \begin{cases} \cos^2\theta \\ \sin\theta\cos\theta \end{cases}$	$\rho^{PP} = 0$	$\rho^{PP} = 1.3$
$D^+ \to \bar{K}^0\pi^+$	$\frac{4}{3}(c_1 + c_2)$	1	1
$D^0 \to \bar{K}^0\pi^0$	$\frac{1}{\sqrt{2}}\left[\frac{8}{9}e^{i\delta_{3/2}}(c_1+c_2) + e^{i\delta_{1/2}}\left(-\frac{5}{9}c_1 + \frac{1}{9}c_2 - \rho^{PP}\cdot c_1\right)\right]$	0.61 $\delta_{1/2} - \delta_{3/2} =$	2.9 120°
$D^0 \to K^-\pi^+$	$\frac{4}{9}e^{i\delta_{3/2}}(c_1+c_2) + e^{i\delta_{1/2}}\left(\frac{5}{9}c_1 - \frac{1}{9}c_2 + \rho^{PP}c_1\right)$	0.29 $\delta_{1/2} - \delta_{3/2}$	3.4 $= 120^\circ$
$D^0 \to \bar{K}^0\eta$	$\frac{1}{\sqrt{6}}(c_2 + \frac{1}{3}c_1) - \frac{1}{\sqrt{6}}\rho^{PP}\cdot c_1$	$\sim 10^{-3}$	0.3
$D^0 \to \bar{K}^0\eta'$	$\frac{1}{\sqrt{3}}(c_2 + \frac{1}{3}c_1) + \frac{2}{\sqrt{3}}\rho^{PP}\cdot c_1$	$\sim 10^{-3}$	2.5
$D^+ \to \bar{K}^0K^+$	$(c_1 + \frac{1}{3}c_2) - \rho^{PP}\cdot c_2$	$0.9\,tg^2\theta$	$1.8\,tg^2\theta$
$D^+ \to \pi^+\pi^0$	$\frac{4}{3\sqrt{2}}(c_1 + c_2)$	$0.5\,tg^2\theta$	$0.6\,tg^2\theta$
$D^0 \to K^+K^-$	$\frac{1}{2}(c_1 + \frac{1}{3}c_2 + \rho^{PP}\cdot c_1)(e^{i\delta_0} + e^{i\delta_1})$	$0.9\,tg^2\theta$ $\delta_0 = \delta_1$	$5.4\,tg^2\theta$
$D^0 \to \bar{K}^0K^0$	$\frac{1}{2}(c_1 + \frac{1}{3}c_2 + \rho^{PP}\cdot c_1)(e^{i\delta_0} - e^{i\delta_1})$	0 $\delta_0 = \delta_1$	0
$D^0 \to \pi^+\pi^-$	$-\frac{4}{9}e^{i\delta_2}(c_1+c_2) + \left(-\frac{5}{9}c_1 + \frac{1}{9}c_2 - c_1\cdot\rho^{PP}\right)e^{i\delta_0}$	$0.1\,tg^2\theta$ $\delta_0 - \delta_2 =$	$2.7\,tg^2\theta$ 180°
$D^0 \to \pi^0\pi^0$	$\frac{8}{9}e^{i\delta_2}(c_1+c_2) + \left(-\frac{5}{9}c_1 + \frac{1}{9}c_2 - c_1\cdot\rho^{PP}\right)e^{i\delta_0}$	$1.7\,tg^2\theta$ $\delta_0 - \delta_2 =$	$3.5\,tg^2\theta$ 180°
$D^0 \to \pi^0\eta$	$-\frac{1}{\sqrt{3}}\left[(c_2 + \frac{1}{3}c_1) - c_1\cdot\rho^{PP}\right]$	$\sim 10^{-3}/tg^2\theta$	$0.6\,tg^2\theta$
$D^0 \to \pi^0\eta'$	$\frac{1}{\sqrt{5}}\left[(c_2 + \frac{1}{3}c_1) + 2c_1\rho^{PP}\right]$	$\sim 10^{-3}/tg^2\theta$	$1.2\,tg^2\theta$
$D^0 \to \eta\eta$	$-(c_2 + \frac{1}{3}c_1) + c_1\cdot\rho^{PP}$	$\sim 10^{-3}/tg^2\theta$	$0.9\,tg^2\theta$
$F^+ \to \pi^+\eta$	$-\sqrt{\frac{2}{3}}(c_1 + \frac{1}{3}c_2) + \sqrt{\frac{2}{3}}c_2\cdot\rho^{PP}$	0.6	1.2
$F^+ \to \pi^+\eta'$	$\frac{1}{\sqrt{3}}(c_1 + \frac{1}{3}c_2) + \frac{2}{\sqrt{3}}c_2\cdot\rho^{PP}$	0.3	$\sim 10^{-2}$
$F^+ \to K^+\bar{K}^0$	$(c_2 + \frac{1}{3}c_1) + c_2\cdot\rho^{PP}$	$\sim 10^{-3}$	0.15
$F^+ \to K^+\pi^0$	$\frac{1}{\sqrt{2}}\left[\frac{8}{9}e^{i\delta_{3/2}}(c_1+c_2) + e^{i\delta_{1/2}}\left(-\frac{5}{9}c_1 + \frac{1}{9}c_2 + c_2\rho^{PP}\right)\right]$	$0.6\,tg^2\theta$ $\delta_{1/2} - \delta_{3/2}$	$tg^2\theta$ $= 120^\circ$
$F^+ \to K^0\pi^+$	$-\frac{4}{9}e^{i\delta_{3/2}}(c_1+c_2) + e^{i\delta_{1/2}}\left(-\frac{5}{9}c_1 + \frac{1}{9}c_2 - c_2\rho^{PP}\right)$	$0.3\,tg^2\theta$ $\delta_{1/2} - \delta_{3/2}$	$0.8\,tg^2\theta$ $= 120^\circ$
$F^+ \to K^+\eta$	$-\frac{1}{\sqrt{6}}\left[\frac{11}{3}c_2 + 3c_1 + c_2\rho^{PP}\right]$	$0.8\,tg^2\theta$	$0.5\,tg^2\theta$

TABLE II

MODE	AMPLITUDE $\frac{G_F}{\sqrt{2}} m_D^2 f_\rho f_+ \cdot \begin{cases} \cos^2\theta \\ \cos\theta\sin\theta \end{cases}$	$\frac{\Gamma(D,F \to PV)}{\Gamma(D^+ \to \bar K^0 \rho^+)}$ $c_1=1.15,\ c_2=-0.33\ \beta=0.7$	
		$\rho^{PV}=\rho^{VP}=0$	$\rho^{PV}=\rho^{VP}=0.8$
$D^+ \to \bar K^0 \rho^+$	$(c_1+\frac{1}{3}c_2)+\beta(c_2+\frac{1}{3}c_1)$	1	1
$D^+ \to \bar K^{0*}\pi^+$	$(c_2+\frac{1}{3}c_1)+\beta(c_1+\frac{1}{3}c_2)$	0.5	0.5
$D^0 \to \bar K^0 \rho^0$	$\frac{1}{\sqrt{2}}\beta(c_2+\frac{1}{3}c_1)-\frac{1}{\sqrt{2}}c_1\cdot\rho^{PV}$	$\sim 10^{-3}$	0.4
$D^0 \to K^- \rho^+$	$(c_1+\frac{1}{3}c_2)+c_1\rho^{PV}$	1	3.6
$D^0 \to K^{-*}\pi^+$	$(c_1+\frac{1}{3}c_2)\beta+c_1\cdot\rho^{VP}$	0.5	2.5
$D^0 \to \bar K^{0*}\pi^0$	$\frac{1}{\sqrt{2}}(c_2+\frac{1}{3}c_1)-\frac{1}{\sqrt{2}}c_1\rho^{VP}$	$\sim 10^{-3}$	0.4
$D^0 \to \bar K^0 \varphi$	$c_1\cdot\rho^{VP}$	0	0.7
$D^0 \to \bar K^0 \omega$	$\frac{1}{\sqrt{2}}\beta(c_2+\frac{1}{3}c_1)+\frac{1}{\sqrt{2}}c_1\rho^{PV}$	$\sim 10^{-3}$	0.4
$D^+ \to \bar K^0 K^{+*}$	$(c_1+\frac{1}{3}c_2)-c_2\rho^{VP}$	$tg^2\theta$	$1.6\,tg^2\theta$
$D^+ \to \bar K^{0*}K^+$	$(c_1+\frac{1}{3}c_2)\beta-c_2\rho^{PV}$	$0.5\,tg^2\theta$	$0.9\,tg^2\theta$
$D^0 \to K^{-*}K^+$	$(c_1+\frac{1}{3}c_2)\beta+c_1\rho^{VP}$	$0.5\,tg^2\theta$	$2.5\,tg^2\theta$
$D^0 \to K^- K^{+*}$	$(c_1+\frac{1}{3}c_2)+c_1\rho^{PV}$	$tg^2\theta$	$3.6\,tg^2\theta$
$D^0 \to \rho^-\pi^+$	$-(c_1+\frac{1}{3}c_2)\beta-c_1\rho^{VP}$	$0.5\,tg^2\theta$	$2.5\,tg^2\theta$
$D^0 \to \rho^+\pi^-$	$-(c_1+\frac{1}{3}c_2)-c_1\rho^{PV}$	$tg^2\theta$	$3.6\,tg^2\theta$
$D^0 \to \rho^0\pi^0$	$\frac{1}{2}(c_2+\frac{1}{3}c_1)(1+\beta)-\frac{1}{2}c_1(\rho^{PV}+\rho^{VP})$	$\sim 10^{-3}\,tg^2\theta$	$0.8\,tg^2\theta$
$D^0 \to \pi^0\omega$	$-\frac{1}{2}(c_2+\frac{1}{3}c_1)(1-\beta)+\frac{1}{2}c_1(\rho^{PV}+\rho^{VP})$	$\sim 10^{-3}\,tg^2\theta$	$0.8\,tg^2\theta$
$F^+ \to \pi^+\omega$	$\frac{1}{\sqrt{2}}c_2(\rho^{PV}+\rho^{VP})$	0	0.13
$F^+ \to \rho^+\eta$	$-\frac{2}{\sqrt{6}}[c_1+\frac{1}{3}c_2+\beta(c_2+\frac{1}{3}c_1)]+\frac{c_2}{\sqrt{6}}(\rho^{PV}+\rho^{VP})$	0.66	1.1
$F^+ \to K^0\rho^+$	$-(c_1+\frac{1}{3}c_2)+c_2\rho^{VP}$	$tg^2\theta$	$1.6\,tg^2\theta$
$F^+ \to K^{0*}\pi^+$	$-(c_1+\frac{1}{3}c_2)\beta+c_2\rho^{PV}$	$0.5\,tg^2\theta$	$0.9\,tg^2\theta$

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