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Pole-dominance model for nonleptonic charm decays

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Two- and three-body nonleptonic weak decays of D and F mesons are discussed. Current-field identities are used in the quantum-chromodynamics-corrected Hamiltonian. With a reasonable choice of strong Hamiltonian the resulting pole model gives rise to vector–axial-vector dominance in two-body D decays, vector-pseudoscalar dominance in two-body F decays and ρ dominance in $D^0 \to K^-\pi^+\pi^0$ decay. The K^* contribution to $D \to K\pi\pi$ decays is found to be small, which is consistent with experiment.

I. INTRODUCTION

The quantum-chromodynamics-corrected standard charm model[1] implies an exact $\Delta I = 1$ rule and an approximate $\Delta V = 0$ rule for Cabibbo-favored nonleptonic charm decays. The $\Delta I = 1$ rule by itself leads only to sum rules. It does not explain, for example, the experimentally observed large ratio $\Gamma(D^0 \to K^-\pi^+\pi^0)/\Gamma(D^+ \to K^-\pi^+\pi^+)$, although it does not exclude that possibility.[2] Soft-pion theorems, on the other hand, are doubtful to be applicable to charm decays since the pions are not soft. Inclusion of final-state interactions may be quite significant[3] but a comprehensive analysis in this direction is yet to be done.

As another approach, quark-line rules[4] emphasize the approximate $\Delta V = 0$ rule aspect of the Hamiltonian. It is assumed that charm decays are the decays of the charmed quark into ordinary quarks or the annihilation of the charmed quark with the accompanying ordinary quark. This approach leads to inclusive semileptonic and nonleptonic decay rates. However, as far as individual decay modes are concerned it neglects symmetry-breaking effects and gives only relative magnitudes of two-body decay rates in terms of quantum-chromodynamics (QCD) factors $\chi_{\mp}$.[4] It provides even less information on three- and more-body decays since they involve more than one type of unknown amplitudes.[5]

In order to obtain more detailed information on charm decays, for example, understand some exclusive decay modes, we consider a phenomenological model in which the current-field identities are used in the QCD-corrected nonleptonic weak Hamiltonian. The QCD correction implies a very specific pattern in charm decays such that $\{D^0\}_{S,P,V,A} \to \{\bar{K}^0\}_{S,P,V,A}$ transitions are suppressed whereas $\{F^+\}_{S,P,V,A} \to \{\pi^+\}_{S,P,V,A}$ transitions are enhanced. We also consider a strong Hamiltonian which describes vertices containing only three particles. With reasonable assumptions for the symmetry-breaking parameters, calculations of the two- and three-body charm decays show that two-body D decays are dominated by vector and axial-vector poles whereas two-body F decays are dominated by vector and pseudoscalar poles; the ρ-meson contribution accounts for a significant percentage of the observed rate for $D^0 \to K^-\pi^+\pi^0$ whereas the K^* contribution, in spite of being the dominant pole contribution in $D^0 \to \bar{K}^0\pi^+\pi^-$ and $D^+ \to K^-\pi^+\pi^+$ decays, accounts only for 15–20% of the observed rates, which implies that the direct-channel contribution is important for these modes. These results are considerably different from those obtained in Ref. 7 and are consistent with the experiment.[6]

In the following we discuss two- and three-body decays of charmed mesons. At the end we examine the application of the present model to ordinary meson decays.

II. MODEL

A similar model to the one we shall describe below was considered by Borchardt *et al.*[7] They have taken the short-distance enhancement factor for the $\underline{20}$-plets part of the nonleptonic weak Hamiltonian for charm decays to be the same as the one for the octet part of the Hamiltonian for ordinary meson decays, and they have calculated pseudoscalar (P) and vector (V) poles for two-body decay modes of charmed mesons. They have found that this model predicts large values for PV contributions to PPP three-body decays. However, there is strong evidence that the enhancement factor for charm decays is much smaller than the one for ordinary meson decays.[8] We take this fact into consideration; also, we consider the axial-vector contributions as well, and calculate two- and three-body charmed-meson decays.[9] The nonleptonic weak Hamiltonian is given by

$$\mathcal{H}^{\rm NL}_{\Delta C=\Delta S=-1}=\frac{G_F\cos^2\theta_C}{\sqrt{2}}\left[\left(\frac{f_++f_-}{2}\right)J^4_{\mu 3}J^{\dagger 2}_{\mu 1}+\left(\frac{f_+-f_-}{2}\right)J^2_{\mu 3}J^{\dagger 4}_{\mu 1}+\text{H.c.}\right], \quad (1)$$

where $J^\beta_{\mu\alpha}=V^\beta_{\mu\alpha}+A^\beta_{\mu\alpha}$ ($\alpha,\beta=1,\ldots,4$) and $f_\mp$ are short-distance factors. $V^\beta_{\mu\alpha}$ and $A^\beta_{\mu\alpha}$ are $\underline{15}$-plets of vector and axial-vector currents. α,β denote the quark flavors. Following Cabibbo *et al.*[4] we adopt the values

$$f_+\approx 2.15\,,\qquad f_-\approx 0.68\,. \quad (2)$$

We use the current-field identities given by

$$V^\beta_{\mu\alpha}=\sqrt{2}\,\frac{m_V{}^2}{f_V}\phi^\beta_{\mu\alpha}+\sqrt{2}f_S\partial_\mu S\,,\qquad A^\beta_{\mu\alpha}=\sqrt{2}\,\frac{m_A{}^2}{f_A}\psi^\beta_{\mu\alpha}+\sqrt{2}f_P\partial_\mu P\,, \quad (3)$$

where ϕ, S, ψ, P are $1^-, 0^+, 1^+, 0^-$ $\underline{15}$-plets of SU(4) respectively and $f_{V,S,A,P}$ are the proper coupling constants. From the $\pi\to l\nu$ decay,

$$f_\pi\approx 93\text{ MeV}\,. \quad (4)$$

Following Kandaswamy *et al.*[10] we use

$$f_K/f_\pi\approx 1.2\,,\qquad f_D/f_\pi\approx 1.6\,,\qquad F_F/f_\pi\approx 1.7\,. \quad (5)$$

We will postpone the discussion of the scalars for the moment. As far as $f_{V,A}$ is concerned, although f_ρ, f_{K^*} appear in the amplitudes, $(f_{V,A})_{\rm charm}$ appear only as g_V/f_V where $g_{V,A}$ are strong-interaction coupling constants which we will discuss later. For f_ρ and f_{K^*} we will use the KSRF relation[11]

$$f_V=\frac{m_V}{\sqrt{2}f_P}\,. \quad (6)$$

f_ρ obtained from this relation agrees fairly well with experiment. Now considering that a certain final state may be created by the Fierz-transformed form of the Hamiltonian, from Eq. (1) we obtain

$$(\mathcal{H}^{\rm NL}_{\Delta C=\Delta S=-1})_{\rm eff}=\frac{G_F}{\sqrt{2}}\cos^2\theta_C(\chi_+J^4_{\mu 3}J^{\dagger 2}_{\mu 1}+\chi_-J^2_{\mu 3}J^{\dagger 4}_{\mu 1}+\text{colored part}+\text{H.c.})\,, \quad (7)$$

where $\chi_\mp=(2f_+\mp f_-)/3$. Since we are interested in color-singlet final states we will not need the colored part. As to the choice of the strong Hamiltonian, that will certainly affect the results. We choose the following Hamiltonian which describes strong vertices containing three particles[7,12].

$$\mathcal{H}_{\rm str}=ig_{VPP}\,\text{Tr}\phi_\mu P\overleftrightarrow{\partial}_\mu P+ig_{AVP}\,\text{Tr}(\psi_\mu\phi_\mu P-P\phi_\mu\psi_\mu)+g_{SPP}\,\text{Tr}SPP+ig_{SVV}\,\text{Tr}S\phi_\mu\phi_\mu+ig_{APS}\,\text{Tr}\psi_\mu P\overleftrightarrow{\partial}_\mu S-\tfrac{2}{3}ig_{VVV}\,\text{Tr}F_{\mu\nu}\phi_\mu\phi_\nu\,, \quad (8)$$

where we have ignored interactions with dimension higher than four for simplicity. The fourth and sixth terms contribute only to VV modes and that contribution is proportional to χ_-, therefore we neglect them. On the other hand, the third term contributes to PP and PPP modes and the fifth term contributes to PPP modes. However, the dominant contribution (i.e., $\alpha\chi_+$) to PP modes is proportional to $m_\pi{}^2/(m_F{}^2-m_\pi{}^2)$, therefore it can be easily neglected, and the contribution to PPP modes which involve two pions is only through κ pole (Fig. 1). The status of κ is somewhat doubtful. It is not a Breit-Wigner resonance and it is very broad. We prefer to ignore this contribution. It can be shown that this will amount to ignoring the diagrams of the type shown in Fig. 1. We may expect future experiments to clarify this question. Thus, the only relevant terms in Eq. (7) are the first two terms which imply pseudoscalar, vector, and axial-vector poles for the processes under consideration. As far as the strong coupling constants are concerned, $g_{\rho\pi\pi}$ and $g_{K^*K\pi}$ appear explicitly in the amplitudes which we will consider, whereas $g_{A_1VP,\,Q_1VP}$ and $(g_{VPP,\,AVP})_{\rm charm}$ appear only in the ratios $\lambda_V\equiv g_{VPP}/f_V$ and $\lambda_{V_A}\equiv g_{AVP}/f_A$. For $g_{\rho\pi\pi}$ and $g_{K^*K\pi}$, using the observed widths and Eq. (8) we obtain

$$g\rho\pi\pi\approx 6.0\,,\qquad gK^*K\pi\approx 5.6\,. \quad (9)$$

For the other needed coupling constants we use the first Weinberg sum rule (saturated with vector, axial-vector, and pseudoscalar mesons) and the KSRF relation in the formulas obtained in Ref. 13 by using current algebra and pole dominance to obtain

$$\lambda_V=1-\frac{m_V{}^2}{2m_A{}^2}\,,\qquad \lambda_{AV}=\frac{m_V}{\sqrt{2}}\left(1-\frac{m_V{}^2}{m_A{}^2}\right)\,, \quad (10)$$

where V and A have the same quantum numbers except the parity. λ_{A_1P} and $\lambda_{Q_1K^*}$ obtained from Eq. (10) agree with experiment within 15%; however, λ_ρ and λ_{K^*} are off by 25%. For more accuracy in our numerical evaluation we will use Eq. (9) which with Eq. (6) implies $\lambda_\rho\approx\lambda_{K^*}\approx 1$ and hope that Eq. (10) is reasonable for the other coupling

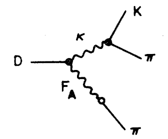

FIG. 1. Scalar pole contribution to $D\to K\pi\pi$ decay mode. κ has $I(J^P)=\frac{1}{2}(0^+)$.

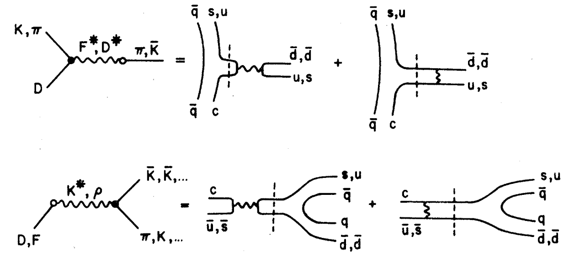

FIG. 2. Vector pole contributions to some PP decay modes. The dashed lines denote the poles.

constants within 15–25%. As we will see, even this level of accuracy will be enough to draw certain conclusions about charm decays. Finally, for the masses, following Fakirov *et al.*[14] we use $m_D \approx 1867$, $m_F \approx 2030$, $m_{D_A} \approx 2200$, $m_{F_A} \approx 2300$, $m_{D^*} \approx 2010$, $m_{F^*} \approx 2140$, all numbers in MeV.

III. DECAY RATES

(*a*) *PP modes.* The relevant diagrams are shown in Fig. 2. For example, $A(D^0\to K^-\pi^+)$ and $A(D^+\to \bar{K}^0\pi^+)$ are given by

$$A^{-+}=2\chi_+\lambda_{F^*}f_\pi(m_D^2-m_\pi^2)-2\chi_-\lambda_{K^*}f_D(m_K^2-m_\pi^2)\,,$$
$$A^{0+}=2\chi_+\lambda_{F^*}f_\pi(m_D^2-m_K^2)+2\chi_-\lambda_{D^*}f_K(m_D^2-m_\pi^2)\,, \tag{11}$$

from which we obtain

$$\Gamma^{-+}/\Gamma^{0+}\approx 2.0\,. \tag{12}$$

This is consistent with experiment if $\Gamma(D^+)\approx\Gamma(D^0)$. However, if we take the central values of the observed branching ratios, Eq. (12) suggests $\Gamma(D^+)<\Gamma(D^0)$ which is reasonable since $D^+\to K^{+*}$ transitions are doubly Cabbibo suppressed whereas $D^0\to\bar{K}^{0*}$ transitions are not. Dominant PP modes are listed in Table I. The suppressed modes are those which are proportional to $(\chi_-)^2$. Comparing our results with those of Ref. 15 we observe that although the suppressed modes are the same, we predict larger $\Gamma(D^0\to K^-\pi^+)/\Gamma(D^+\to\bar{K}^0\pi^+)$ and $\Gamma(F^+\to\eta\pi)$.

(*b*) *PV modes.* The relevant diagrams for $A(D^0\to K^-\rho^+)$ and $A(D^+\to\bar{K}^{0*}\pi^+)$ are shown in Fig. 3. For example, neglecting $m_\pi^2/(m_F^2-m_\pi^2)$ terms we obtain

$$A(D^0\to K^-\rho^+)=[2\chi-\lambda_{Q_1}f_D-4\chi_+\lambda_{F^*}m_\rho^2/f_\rho(1-m_\rho^2/m_F^2)+2\chi_-g_{\rho KK}f_Df_K/(1-m_K^2/m_D^2)]\,q_D\cdot\epsilon\,,$$
$$A(D^+\to\bar{K}^{0*}\pi^+)=[2\chi_+\lambda_{F_AK^*}f_\pi-4\chi_-\lambda_{D^*}m_{K^*}^2/f_{K^*}(1-m_{K^*}^2/m_{D^*}^2)]\,q_D\cdot\epsilon\,,$$
$$A(F^+\to\bar{K}^{0*}K^+)=[2\chi_+g_{K^*K\pi}f_Ff_\pi-4\chi_-\lambda_{D^*}m_{K^*}^2/f_{K^*}(1-m_{K^*}^2/m_{D^*}^2)+2\chi_+\lambda_{Q_1K^*}]\,q_F\cdot\epsilon\,. \tag{13}$$

We note that the 0^- contribution is small for D decays yet it can be dominant for F decays because the $F^+\to\phi\pi^+$ transition is enhanced. Comparing our results (Table II) with those obtained by Maiani in Ref. 15 we observe significant differences. For example, in our model $F^+\to\phi\pi^+$ is suppressed whereas $F^+\to\eta\rho^+$, $K^{*+}\bar{K}^0$, $K^{0*}K^+$ are not, which is opposite to his result. Also, in our

TABLE I. *PP* decay rates. The suppressed modes are proportional to $(\chi_-)^2$.

Decay mode	Decay rate (10^{10} sec^{-1})
$D^0\to K^-\pi^+$	17.8
$D^0\to\bar{K}^0\pi^+$	8.6
$D^0\to\bar{K}^0\pi^0$	suppressed
$D^0\to\bar{K}^0\eta$	suppressed
$F^+\to\eta\pi^+$	21.2
$F^+\to K^+\bar{K}^0$	suppressed

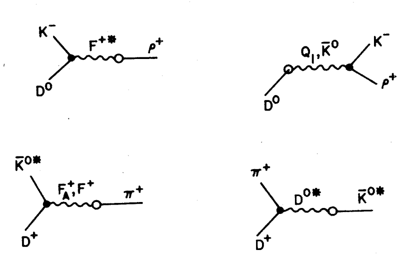

FIG. 3. Vector and axial-vector pole contributions to some PV decay modes.

TABLE II. PV decay rates. The suppressed modes are proportional to $(\chi_-)^2$.

Decay mode	Decay rate (10^{10} sec^{-1})
$D^+ \to \bar{K}^0\rho^+$	24.6(F^*contr $\approx$ 21.1)
$D^+ \to \bar{K}^{*0}\pi^+$	8.0
$D^0 \to K^{*-}\pi^+$	5.6
$D^0 \to K^-\rho^+$	32.6(F^*contr $\approx$ 21.1)
$D^0 \to \bar{K}^{*0}\pi^0$	suppressed
$D^0 \to \bar{K}^0\rho^0$	suppressed
$D^0 \to \bar{K}^{0*}\eta$	suppressed
$D^0 \to \bar{K}^0\omega$	suppressed
$F^+ \to \phi\pi^+$	2.2
$F^+ \to \eta\rho^+$	18.6
$F^+ \to K^{*+}\bar{K}^0$	19.5
$F^+ \to \bar{K}^{0*}K^+$	24.0

model $B(D^0 \to K^-\rho^+)/B(D^+ \to \bar{K}^{*0}\pi^+) \approx 2.46$ and $B(D^0 \to K^-\rho^+) \approx 5.82$ whereas he finds 1.35 and 0.85 respectively. We note that we have used the experimental input $B(K^-\pi^+) \approx 2.2\%$ and $B(\bar{K}^0\pi^+) \approx 1.6\%$ here. On the other hand, all of the PV rates we have obtained are much smaller than those of Ref. 7.

(*c*) *VV modes.* Pseudoscalars do not contribute to VV modes. Also our simple calculations show that vector contributions are negligible. Therefore we have axial-vector dominance in VV modes.

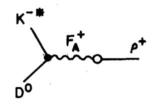

FIG. 4. Axial-vector pole contribution to a VV mode.

For example, $A(D^0 \to K^{-*}\rho^+)$ is given by (Fig. 4)

$$A(D^0 \to K^{-*}\rho^+) = [2\chi_+ \lambda_{F_A K^*} m_\rho^2 / f_\rho (1 - m_\rho^2/m_{F_A}^2)] \epsilon^{(K^*)} \cdot \epsilon^{(\rho)} . \tag{14}$$

Results are tabulated in Table III. They are smaller than the predictions of Ref. 15. $D^0 \to \bar{K}^0{}^*\rho^0$ is suppressed in both models and this is consistent with experiments.

(*d*) *PPP modes.* We consider $A(D^0 \to K^-\pi^+\pi^0)$ first. χ_+-enhanced pseudoscalar contributions are proportional to $m_\pi^2/(m_F^2 - m_\pi^2)$, therefore they can be neglected. Vector and axial-vector contributions are shown in Fig. 5. For simplicity, we shall neglect the graphs which are proportional to χ_- from here on [Fig. 5(d)]. The contribution from Fig. 5(a) has been calculated to be an order of magnitude smaller than the other contributions. Therefore the relevant diagrams are only Figs. 5(b) and 5(c). We obtain

$$A(D^0 \to K^-\pi^+\pi^0) = \frac{\sqrt{2}\chi_+ g_{KK^*\pi} \lambda_{F_A K^*} f_\pi}{(q_D - q_+)^2 - m_{K^*}^2 + i m_{K^*}\Gamma_{K^*}} [(1-x)\, q_+ \cdot q_K - (1+x)\, q_+ \cdot q_0]_{[\text{Fig. 5(b)}]} + \frac{2\sqrt{2}\chi_+ \lambda_{F^*} \lambda_\rho m_{F^*}^2 m_\rho^2}{(q^2 - m_{F^*}^2)(q^2 - m_\rho^2 + i m_\rho \Gamma_\rho)} [(q_D + q_K) \cdot (q_+ - q_0)]_{[\text{Fig. 5(c)}]} , \tag{15}$$

where $x = (m_K^2 - m_\pi^2)/m_{K^*}^2$ and $q = q_D - q_K$. This amplitude must be numerically integrated over the Dalitz plot. For F decays pseudoscalar contributions are again appreciable. Results for some of the PPP modes are given in Table IV. We observe that the K^* contribution to $\Gamma(D^+ \to K^-\pi^+\pi^+)$ and $\Gamma(D^0 \to \bar{K}^0\pi^+\pi^-)$ is 15–20% of the observed rate whereas the ρ contribution is practically absent. This is consistent with experiments. On the other hand, the ρ contribution to $\Gamma(D^0 \to K^-\pi^+\pi^0)$ is quite sizable. $B(D^0 \to K^-\pi^+\pi^0) = (9 \mp 5)\%$ implies

$$(\rho \text{ contribution})_{D^0 \to K^-\pi^+\pi^0} \approx 40^{+50}_{-25}\% . \tag{16}$$

This is approximately the ρ contribution alone, since

$$(K^* \text{ contr}/\rho \text{ contr})_{D^0 \to K^-\pi^+\pi^0} \approx 7\% \tag{17}$$

from Table IV. The idea of ρ dominance in $D^0 \to K^-\pi^+\pi^0$ is yet to be verified experimentally.

TABLE III. VV decay rates. The suppressed modes are proportional to $(\chi_-)^2$.

Decay mode	Decay rate (10^{10} sec^{-1})
$D^+ \to \bar{K}^{*0}\rho^+$	4.2
$D^0 \to K^{*-}\rho^+$	4.2
$D^0 \to \bar{K}^{0*}\rho^0$	suppressed
$D^0 \to \bar{K}^{0*}\omega$	suppressed
$F^+ \to \phi\rho^+$	4.8
$F^+ \to K^{*+}\bar{K}^{*0}$	suppressed

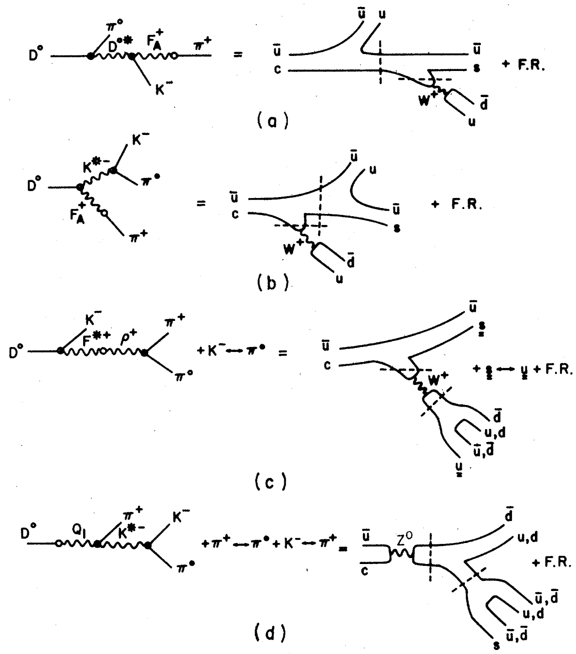

FIG. 5. Vector and axial-vector pole contributions to the $D^0 \to K^-\pi^+\pi^0$ decay mode. FR stands for Fierz rearrangement. The dashed lines denote the poles.

IV. DISCUSSION

The numerical results obtained above certainly depend on the choice of H_{str} and coupling constants such as g_{AVP}, etc. However, even the order of magnitudes as crude estimates show that $\Gamma(D^0 \to K^-\pi^+\pi^+)/\Gamma(D^+ \to K^-\pi^+\pi^+)$ may indeed be expected to be large since the ρ contribution is *much*

TABLE IV. *PPP* decay rates and branching ratios.

Decay mode	Decay rate (10^{10} sec^{-1})	Branching ratio
$D^+ \to K^-\pi^+\pi^+$	4.3	0.8%
$D^0 \to \bar{K}^0\pi^+\pi^-$	4.2	0.5%
$D^0 \to K^-\pi^+\pi^0$	29.6	3.7%
$D^+ \to \bar{K}^0\pi^+\pi^0$	29.6	5.4%
$F^+ \to \pi^+\pi^-\pi^+$	38.4	

larger than the K^* contribution. It has been discussed in Ref. 5 that ρ dominance may be responsible for large $\Gamma(D^0 \to K^-\pi^+\pi^0)/\Gamma(D^+ \to K^-\pi^+\pi^+)$. (See also Ref. 18.) Within a simple model we show that this is indeed the case. As far as the missing part of the branching ratios is concerned (see Table IV), a possible explanation might be that the type of diagrams shown in Fig. 6 are significant and they are approximately equal for each $K\pi\pi$ mode (within obvious symmetry factors). Then, for instance, a contribution of about 3% to each $K\pi\pi$ mode would bring our results to very good agreement with experiments. What is needed is, then, a more comprehensive phenomenological model which would preserve our results but also explain the direct-channel contributions through vertices containing more than three particles.[16]

As to ordinary-meson decays, it can be shown that the present model gives the correct $K \to \pi\pi$ rates provided that the QCD factor χ_+ is larger than theoretically expected.[7] However, it has been shown that[17] strong interactions not only modify various pieces of the bare Hamiltonian but also give rise to new structures of the type

$$\Delta\mathcal{H}^{NL}_{\substack{\Delta S=-1\\ \Delta C=0}} = C\,\frac{2G_F}{\sqrt{2}}\cos\theta_C\sin\theta_C\,[\bar{s}_L\gamma_\mu d_L(\bar{u}_R\gamma_\mu u_R + \bar{d}_R\gamma_\mu d_R + \bar{s}_R\gamma_\mu s_R) + \text{H.c.}]\,. \tag{18}$$

Now in the present model this would give rise to

$$(\mathcal{H}^{NL}_{\substack{\Delta S=-1\\ \Delta C=0}})_{\text{eff}} = a(K^{*0}_\mu - \partial_\mu K^0)(\phi_{0\mu} + \partial_\mu\eta') + b\,[(K^+ + K^+)(\delta^- - \pi^-) + (K^0 + K^0)(\delta^0 - \pi^0)]\,, \tag{19}$$

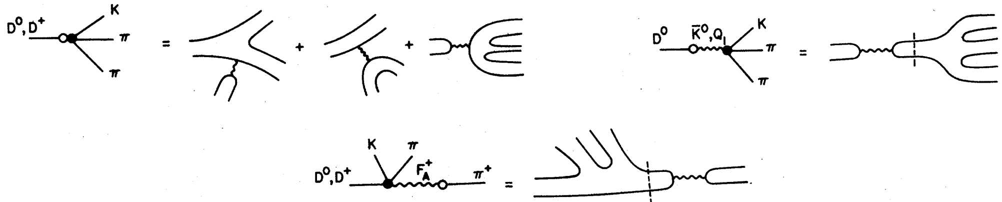

FIG. 6. Direct-channel and vector and axial-vector pole contributions to $D \to K\pi\pi$ decay modes where the strong vertices contain four particles. The dashed lines denote the poles.

where we have used the Fierz transformation on quark terms to obtain the effective Hamiltonian. We observe that the new terms do not give rise to $K \rightarrow \pi\pi$ decays through pole diagrams of the type shown in Fig. 2. However, let us note that they can give rise to different types of contributions such as $\langle \pi | J | 0 \rangle \langle \pi | J^+ | K \rangle$ and these have been argued to be significantly large.[17] Yet the calculation of the new contributions is a subtle problem and it requires further investigation in order to solve the octet enhancement problem completely in a satisfactory way.

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