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A Hybrid Model for Charmed Meson Decays

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Abstract

Two-body decays of the charmed mesons to both PP and PV final states are discussed, employing an effective Lagrangian formalism. We show that a hybrid model combining the pole model and the c-quark decay model provides a reasonable description of the amplitudes for these processes.

1. Introduction

Following the prediction of the existence of the charmed quark (Bjorken and Glashow 1964; Glashow *et al.* 1970) a fairly straightforward decay mechanism was envisaged for the charmed D and F mesons. The charmed quark was expected to decay weakly with the light antiquark remaining an uninvolved 'spectator' (Fig. 1), strong interaction corrections being smaller than in light meson decays due to the relatively large mass of the charmed quark (Gaillard *et al.* 1975). This would imply equal lifetimes for D^0, D^+ and F^+. However, there is now evidence (Bacino *et al.* 1980; Ushida *et al.* 1980*a*, 1980*b*) for D^+ having a significantly greater lifetime than D^0, so there is clearly a need to revise this simple picture.

In this paper we investigate decays of D^0 and D^+ to final states of two pseudoscalar mesons (PP) and of one pseudoscalar and one vector meson (PV), and attempt to describe the data by adding pole terms to the free quark decay terms of the original model. The pole contributions are calculated in such a way that the only free parameter is the strength of the quark decay terms. A reasonable fit to the data is obtained, which is an improvement on that given by the 'free quark' model or the pole model alone.

2. Direct Charmed Quark Decay

In the absence of strong interactions the effective four-fermion interaction contributing to nonleptonic charm decay has the form

$$L_{\Delta C=1} = \sqrt{\tfrac{1}{2}}\,G(\bar{s}'c)(\bar{u}d')\,, \tag{1}$$

using the abbreviated notation

$$\bar{s}'c = \sum_{\alpha} \bar{s}'_{\alpha}\gamma_{\mu}(1-\gamma_5)c^{\alpha}\,, \tag{2}$$

where α is a colour index, and (see Glashow *et al.* 1970)

$$d' = d\cos\theta_C + s\sin\theta_C\,, \qquad s' = d\sin\theta_C + s\cos\theta_C\,. \tag{3a, b}$$

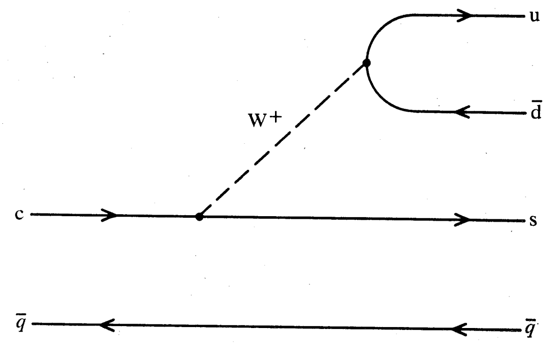

Fig. 1. Charmed-quark decay diagram.

If we now include QCD corrections, to leading log order equation (1) becomes (Gaillard and Lee 1974; Altarelli and Maiani 1974)

$$L_{\Delta C=1} = \sqrt{\tfrac{1}{2}}\,G\{\tfrac{1}{2}(f_+ + f_-)(\bar{s}'c)(\bar{u}d') + \tfrac{1}{2}(f_+ - f_-)(\bar{u}c)(\bar{s}'d')\}. \tag{4}$$

To leading log order we have (Ellis *et al.* 1975)

$$f_- = \{\alpha_s(m_c)/\alpha_s(m_W)\}^{0.48} \approx 2, \qquad f_+ = f_-^{-\frac{1}{2}} \approx 0{\cdot}7. \tag{5a, b}$$

where α_s is the quark–gluon coupling constant of QCD.

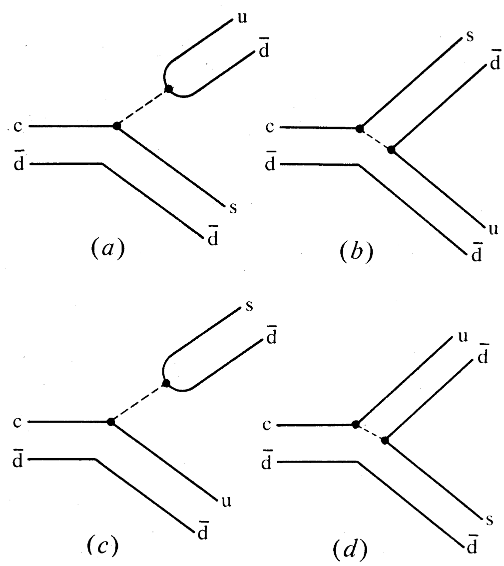

Fig. 2. Four c-quark decay diagrams contributing to $D^+ \to \bar{K}^0\pi^+$ (see Section 2).

Cabibbo and Maiani (1978) have discussed PP and PV decays in terms of the c-quark decay picture. We will use their results in our calculation and we reproduce here their arguments for the case $D^+ \to \bar{K}^0\pi^+$. The c-decay diagrams which must be considered in the effective Lagrangian formalism are shown in Fig. 2. Figs. 2*a* and 2*b* are associated with the $(\bar{s}'c)(\bar{u}d')$ term, and so will be proportional to $f_+ + f_-$, while Figs 2*c* and 2*d* will be proportional to $f_+ - f_-$. If we define the contribution of Fig. 2*a* as $(f_+ + f_-)A'$, with $A' = \sqrt{\tfrac{1}{2}}\,AG\cos^2\theta_C$, then the contribution of Fig. 2*c*

Table 1. **Amplitudes (in keV) for D → PP decays**

Decay	M_c	$M_{(a)}$	$M_{(b)}$	M_{tot} [A]	Experiment [B]
$D^0 \to \bar{K}^0\pi^0$	$-2\cdot640A$	$+0\cdot169$	$+1\cdot575$	$1\cdot03$	$2\cdot1^{+0\cdot7}_{-1\cdot0}$
$D^0 \to K^-\pi^+$	$+16\cdot752A$	$-0\cdot239$	$-7\cdot124$	$2\cdot83$	$2\cdot5^{+0\cdot6}_{-0\cdot7}$
$D^+ \to \bar{K}^0\pi^+$	$+13\cdot018A$	0	$-4\cdot896$	$1\cdot38$	$1\cdot4^{+0\cdot3}_{-0\cdot4}$
$D^0 \to K^+K^-$	$+3\cdot768A$	0	$-2\cdot102$	$1\cdot08$	
$D^0 \to \pi^+\pi^-$	$-3\cdot909A$	0	$+1\cdot778$	$0\cdot72$	
$D^+ \to \bar{K}^0K^+$	$+3\cdot768A$	0	$-2\cdot102$	$1\cdot08$	
$D^+ \to \pi^0\pi^+$	$+3\cdot635A$	0	$-1\cdot654$	$0\cdot67$	

[A] $M_{tot} = |M_c + M_{(a)} + M_{(b)}|$ is evaluated at $A = +0\cdot27$ GeV^3.
[B] Aguilar-Benitez *et al.* (1981).

will be $(f_+ - f_-)A'$, assuming SU(3) symmetry. To relate Fig. 2*b* to 2*c* we use the Fierz transformation (Altarelli *et al.* 1975)

$$(\bar{u}d')(\bar{s}'c) = \tfrac{1}{3}(\bar{s}'d')(\bar{u}c) + \tfrac{1}{2}\sum_{\alpha}(\bar{s}'\lambda^\alpha d')(\bar{u}\lambda^\alpha c)\,, \qquad (6)$$

where λ^α are the colour SU(3) Gell-Mann matrices. The second term in (6) produces colour octet states, so it is assumed not to contribute. Thus the contribution of Fig. 2*b* is $\frac{1}{3}(f_+ + f_-)A'$. Similarly the contribution of Fig. 2*d* is $\frac{1}{3}(f_+ - f_-)A'$, giving a total c-decay matrix element of $M_c = \frac{8}{3}f_+ A'$. The c-decay contributions to M_c for $D^+ \to \bar{K}^0\pi^+$ and six other PP decay modes are shown in Table 1, expressed in terms of the parameter A.

Taken alone the c-quark decay amplitudes above give a ratio of branching fractions

$$B(D^0 \to K^-\pi^+)/B(D^0 \to \bar{K}^0\pi^0) \approx 40\,.$$

Recent data from the Mark II collaboration include the following branching fractions (Schindler *et al.* 1981):

$$B(D^0 \to K^-\pi^+) = 3\cdot0 \pm 0\cdot6\%\,, \qquad (7a)$$

$$B(D^0 \to \bar{K}^0\pi^0) = 2\cdot2 \pm 1\cdot1\%\,, \qquad (7b)$$

$$B(D^+ \to \bar{K}^0\pi^+) = 2\cdot3 \pm 0\cdot7\%\,. \qquad (7c)$$

To attempt to redress this discrepancy in the ratios we look at other possible decay mechanisms.

3. W-exchange Generated Poles

So far we have assumed that the only contribution is from diagrams in which the light antiquark is a 'spectator'. If we allow the light antiquark to play an active role we need to consider processes in which the charm changing part of the decay occurs at a two-meson vertex. A $q_1\bar{q}_2$ pair may exchange a t-channel W boson (Fig. 3*a*) or may produce a virtual W^+ in the s channel (Fig. 3*b*). Originally these processes were expected to be negligible because of helicity suppression, however this suppression may be alleviated if gluons are emitted in the initial and/or final states (Bander *et al.* 1980; Bando *et al.* 1980). It becomes plausible that these exchange and annihilation processes contribute significantly, and it is appropriate to parametrize their contribution in some way.

For decays to two pseudoscalars we consider two types of vector pole terms, each of which involves a weak two-meson vertex and a strong three-meson vertex, as shown

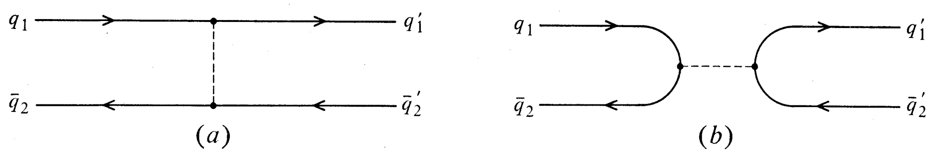

Fig. 3. Diagrams for (*a*) the exchange of a W boson and (*b*) annihilation.

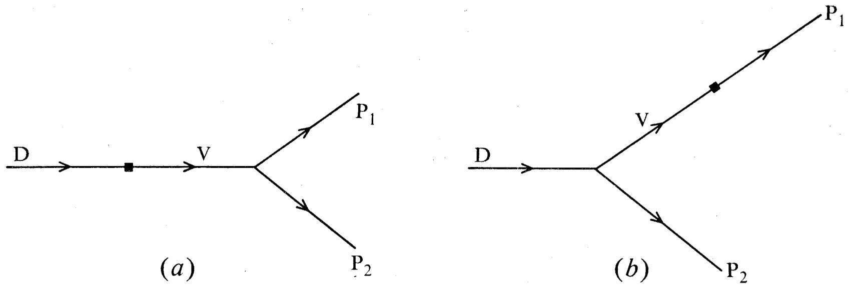

Fig. 4. Vector pole contributions to D → PP decay.

in Fig. 4. If we take $D^0 \to \bar{K}^{0*}$ as an example of the weak vertex, the two contributions are as shown in Fig. 3 with $q_1 = \mathrm{c}$, $\bar{q}_2 = \bar{\mathrm{u}}$, $q_1' = \mathrm{s}$, $\bar{q}_2' = \bar{\mathrm{d}}$. The effective Lagrangian after an appropriate Fierz rearrangement is

$$L^{\mathrm{eff}}_{\Delta C=1} = \surd\tfrac{1}{2}\, G X_-(\mathrm{c}\bar{\mathrm{u}})(\bar{\mathrm{s}}'\mathrm{d}')\,, \tag{8}$$

where

$$X_\pm = \tfrac{1}{3}(2f_+ \pm f_-)\,.$$

We now relate the quark currents $J^2_{\mu 1} = \bar{\mathrm{d}}\mathrm{u}'$ and $J^4_{\mu 3} = \bar{\mathrm{s}}'\mathrm{c}$ to meson fields by the current field identities (Sezgin 1979)

$$J^\beta_{\mu\alpha} = V^\beta_{\mu\alpha} + A^\beta_{\mu\alpha}\,, \tag{9a}$$

$$V^\beta_{\mu\alpha} = \surd 2(m_{\mathrm{v}}^2/f_{\mathrm{v}})\phi^\beta_{\mu\alpha} + \surd 2 f_S\,\partial_\mu S\,, \tag{9b}$$

$$A^\beta_{\mu\alpha} = \surd 2(m_{\mathrm{a}}^2/f_{\mathrm{a}})\psi^\beta_{\mu\alpha} + \surd 2 f_P\,\partial_\mu P\,, \tag{9c}$$

where $V^\beta_{\mu\alpha}$ and $A^\beta_{\mu\alpha}$ are vector and axial currents respectively, $\alpha, \beta = 1, \ldots, 4$ are flavour indices, and ϕ, ψ, S and P are $J^\pi = 1^-, 1^+, 0^+, 0^-$ fields respectively. Thus for the $D^0 \to \bar{K}^{0*}$ pole the coupling is of the form

$$L^{\mathrm{eff}}_{\mathrm{D^0\bar{K}^{0*}}} = \surd 2\, G X_-\, f_{\mathrm{D}}\,\partial_\mu P_{\mathrm{D}}(m_{\mathrm{v}}^2/f_{\mathrm{v}})\phi_\mu \cos^2\theta_{\mathrm{C}}\,. \tag{10}$$

For the strong three-body vertices we use the PPV coupling

$$H^{\mathrm{str}}_{\mathrm{PPV}} = g_{\mathrm{PPV}}\,\mathrm{Tr}\,\phi_\mu P \overleftrightarrow{\partial}_\mu P\,. \tag{11}$$

So for Fig. 4*a* the matrix element for a D^0 decay is

$$M_{(a)} = -\surd 2\, G \cos^2\theta_{\mathrm{C}}\, X_-\, f_{\mathrm{D}}(g_{\mathrm{P_1P_2V}}/f_{\mathrm{v}})(m_1^2 - m_2^2)\,. \tag{12}$$

For particular charge states of V, P_1 and P_2 we take $g_{P_1P_2V} = \gamma\tilde{g}_{P_1P_2V}$, where γ is an isospin factor, for example, $\gamma = -\sqrt{2}$ for $\overline{K}^{0*} \to K^-\pi^+$ and $\gamma = +1$ for $\overline{K}^{0*} \to \overline{K}^0\pi^0$, and we take $\tilde{g}_{P_1P_2V} \approx f_v$ (Sezgin 1979), so that both f_v and m_v drop out. For Fig. 4*b* the matrix element is, for a Cabibbo allowed process,

$$M_{(b)} = -\sqrt{2}\, G \cos^2\theta_C X_\pm f_{P_1}(g_{DP_1V}/f_v)(m_D^2 - m_2^2), \tag{13}$$

and again we take $\tilde{g}_{DP_1V} \approx f_v$. The pseudoscalar decay constants f_P are $f_\pi = 93$ MeV and $f_K = 122$ MeV where appropriate. For f_D the literature estimates vary widely, from about 150 to 800 MeV (Bander *et al.* 1980). Calculations have been done with $f_D = 200$ MeV (Bég 1981). Matrix elements for diagrams of the type shown in Figs 4*a* and 4*b* are given in Table 1 for several decay modes; $M_{(b)}$ is very much the dominant pole term in each case so that the uncertainty in f_D is not critical in our calculation.

It might be argued that we have 'double counted' by including direct decay diagrams (for example, Fig. 2) and pole diagrams such as Fig. 4*b* separately, the latter being the leading term of the total spectator amplitude which the former attempts to describe. Indeed, in the limit of exact SU(3) symmetry, the pole amplitudes would be directly proportional to the direct decay amplitudes. Thus, for the purposes of this paper, what we describe as the c-quark decay term is more correctly a parametrization of what remains of the spectator amplitude after the pole term has been extracted.

To obtain decay amplitudes from the observed branching ratios (7) we need to know the lifetimes of the D^0 and D^+. Statistics for these measurements remain fairly sparse but are improving. Comparison of semileptonic branching ratios for D^0 and D^+ observed on the DELCO detector at SPEAR (Bacino *et al.* 1980) indicate $\tau(D^+)/\tau(D^0) > 4 \cdot 3$ (95% CL), while the Mark II group (Schindler *et al.* 1981) obtained $\tau(D^+)/\tau(D^0) = 3 \cdot 1^{+4.6}_{-1.4}$ in a similar comparison. Direct determination of lifetimes is hindered by the short decay length $c\tau \leqslant 100\ \mu$m, which is below the range of conventional bubble or drift chambers; however, the high resolution hydrogen bubble chamber LEBC is now operating at CERN and initial results based on a total of 19 fully reconstructed decays indicate (Aguilar-Benitez *et al.* 1981)

$$\tau(D^\pm) = 8 \cdot 0^{+4.9}_{-2.4} \times 10^{-13}\ \mathrm{s}, \qquad \tau(D^0) = 3 \cdot 2^{+2.2}_{-1.0} \times 10^{-13}\ \mathrm{s}. \tag{14a, b}$$

The other detection technique used and until now the major source of lifetime data is nuclear emulsion spectrometry. The Fermilab group (see Stanton 1981) with improved statistics on its previous analysis (Ushida *et al.* 1980*a*, 1980*b*) now reports

$$\tau(D^+) = 10 \cdot 3^{+10.3}_{-4.2} \times 10^{-13}\ \mathrm{s}, \qquad \tau(D^0) = 3 \cdot 2^{+1.0}_{-0.7} \times 10^{-13}\ \mathrm{s}. \tag{15a, b}$$

For definiteness the CERN lifetimes (14) are used to convert observed branching ratios to amplitudes, which for PP modes are presented in the last column of Table 1.

The calculated total amplitudes M_{tot}, summing direct decay and pole terms, evaluated at the best fit value of $A = +0 \cdot 27\ \text{GeV}^3$ are also shown in Table 1, from which it can be seen that quite a satisfactory agreement is obtained. Mark II data are also available on ratios of branching fractions (Abrams *et al.* 1979; Kirkby 1979; Schindler *et al.* 1981), and these are compared with calculated values both with pole terms (at $A = 0 \cdot 27\ \text{GeV}^3$) and without (independent of A) in Table 2. This test has the advantage of not relying on lifetime measurements.

Inclusion of the pole terms goes a long way towards reducing the disparity between $\overline{K}^0\pi^0$ and $K^-\pi^+$ amplitudes seen in the pure c-quark decay picure, although it requires a significant cancellation between the two components to achieve this, as either contribution on its own is inadequate.

Table 2. Ratios of branching fractions for D → PP decays

Ratio	Without poles	With poles[A]	Experiment[B]
$B(D^0\to\pi^+\pi^-)/B(D^0\to K^-\pi^+)$	0·058	0·069	$0{\cdot}033\pm0{\cdot}015$
$B(D^0\to K^+K^-)/B(D^0\to K^-\pi^+)$	0·056	0·134	$0{\cdot}113\pm0{\cdot}030$
$B(D^0\to\overline{K}^0\pi^0)/B(D^0\to K^-K^+)$	0·025	0·132	$0{\cdot}73^{+1{\cdot}1}_{-0{\cdot}6}$
$B(D^+\to K^+\overline{K}^0)/B(D^+\to\overline{K}^0\pi^+)$	0·077	0·562	$0{\cdot}25\pm0{\cdot}15$
$B(D^+\to\pi^0\pi^+)/B(D^+\to\overline{K}^0\pi^+)$	0·084	0·237	$<0{\cdot}3$ (90% CL)

[A] Evaluated at $A = +0{\cdot}27\ \mathrm{GeV}^3$.
[B] Data from Abrams *et al.* (1979), Kirkby (1979) and Schindler *et al.* (1981).

4. Decays to a Vector and a Pseudoscalar

As for the PP case there are two types of pole terms in PV decays, both involving a pseudoscalar intermediate state. For diagrams of the type shown in Fig. 4*a* the matrix element is, for a Cabibbo favoured decay,

$$M_{(a)} = \sqrt{\tfrac{1}{2}}\,G\cos^2\theta_C X_\pm\, 2f_{P'}f_D\, g_{VPP'}\frac{m_D^2}{m_D^2-m_{P'}^2}\,\varepsilon^\mu_{(V)}(p_D+p_P)_\mu\,, \tag{16}$$

while for Fig. 4*b* the matrix element is

$$M_{(b)} = \sqrt{\tfrac{1}{2}}\,G\cos^2\theta_C X_\pm\, 2f_{P'}f_P\, g_{VDP'}\frac{m_P^2}{m_P^2-m_{P'}^2}\,\varepsilon^\mu_{(V)}(p_D+p_P)_\mu\,. \tag{17}$$

Because of the mass factors we have $|M_{(b)}|/|M_{(a)}| \leqslant 10$ for $D\to K\rho$ decays and $\leqslant 200$ for $D\to K^*\pi$, and so we choose to neglect $M_{(b)}$. The strong coupling $g_{VPP'}$ can be estimated from decay rate data. Writing $g_{VPP'} = \gamma\tilde{g}_{VPP'}$, where $\gamma = +1$ for $\overline{K}^{0*}\pi^0$ and $\overline{K}^0\rho^0$ final states and $+\sqrt{2}$ for charged final states, we find $\tilde{g}_{K^*K\pi} = 2{\cdot}0$. In the absence of direct measurement we assume that $\tilde{g}_{\rho KK} \approx \tilde{g}_{K^*K\pi}$.

Only decays to longitudinally polarized vector mesons have nonzero amplitudes as the form $M \propto \varepsilon^\mu_{(V)}(p_D+p_P)_\mu$ is the most general allowed, the polarization $\varepsilon^\mu_{(V)}$ being orthogonal to the vector meson four-momentum p_V. This form is used in the parametrization of the c-quark decay part of the amplitude so that, for instance,

$$M_c(D^+\to\overline{K}^{0*}\pi^+) = \tfrac{8}{3}f_+\,B\sqrt{\tfrac{1}{2}}\,G\cos^2\theta_C\,\varepsilon^\mu_{(V)}(p_D+p_\pi)_\mu\,. \tag{18}$$

The factor $\frac{8}{3}f_+$ is obtained by the same argument as in the $\overline{K}^0\pi^+$ case (see Section 2), however there is the reservation that in Fig. 2*a* the spectator quark is now in the vector meson while in Fig. 2*c* it goes into the pseudoscalar, so that in principle the two diagrams may not be so simply related as in the PP case. This distinction should not be important insofar as the four final state quarks can be considered as emerging from the one point, the D being in an s-wave state, so that the newly created $\bar{d}$ and the spectator $\bar{d}$ should be equivalent in the formation of new states (Cabibbo and Maiani 1978).

The c-quark decay amplitudes M_c and the pole amplitudes $M_{(a)}$ are presented in Table 3. The amplitudes in the last column are derived from Mark II branching ratio data (Lankford 1980), again using the assumed lifetimes (14). The value of the parameter B was chosen to give the best fit to this data both with and without pole terms included, and within the currently rather wide bounds of available evidence either fit is fairly adequate. Thus, the PV decays are not as exacting as the PP decays in testing the model.†

Table 3. Amplitudes (in keV) for D → PV decays

Decay	M_c	$M_{(a)}$	$M_{tot}{}^{A}$	$M_{tot}{}^{B}$	Experiment[C]
$D^0 \to K^{-*}\pi^+$	$+49\cdot81B$	$+1\cdot737$	$3\cdot19$	$2\cdot73$	$2\cdot9^{+0.8}_{-1.2}$
$D^0 \to \bar{K}^{0*}\pi^0$	$-7\cdot85B$	$-1\cdot228$	$0\cdot50$	$1\cdot39$	$1\cdot9^{+1.2}_{-1.9}$
$D^0 \to K^-\rho^+$	$+54\cdot32B$	$+1\cdot894$	$3\cdot48$	$2\cdot98$	$4\cdot4^{+1.3}_{-1.9}$
$D^0 \to \bar{K}^0\rho^0$	$-8\cdot56B$	$-1\cdot339$	$0\cdot55$	$1\cdot51$	$0\cdot5^{+0.6}_{-0.5}$
$D^+ \to \bar{K}^{0*}\pi^+$	$+38\cdot72B$	0	$2\cdot48$	$0\cdot77$	
$D^+ \to \bar{K}^0\rho^+$	$+42\cdot22B$	0	$2\cdot70$	$0\cdot84$	

[A] $M_{tot} = |M_c|$ evaluated at $B = 0\cdot064\ \mathrm{GeV}^2$.
[B] $M_{tot} = |M_c + M_{(a)}|$ evaluated at $B = 0\cdot020\ \mathrm{GeV}^2$.
[C] Data from Lankford (1980).

5. Conclusions

We find that it is not possible to fit the data with either pole diagrams or the direct decay term alone when one maintains the effective Lagrangian of equation (4). It is clear that to obtain agreement with the data additional freedom must be introduced into the parametrization. In this paper we have introduced this freedom by including both pole terms and direct decay terms in the amplitudes to provide a useful description of the data in terms of one free parameter.

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† Miller and McKellar (1982) have recently reached similar conclusions in that they show that a spectator based pole model gives a reasonable fit to PP and PV decays. The PP decays seem to be a much more sensitive test of the decay mechanism.

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