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Analysis of two-body decays of charm mesons using the quark-diagram scheme

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Two-body decays of charm mesons are analyzed in the quark-diagram scheme. Effects of SU(3) breaking and final-state interactions are included in the formulation. Various theoretical models are reviewed and tested. In this scheme, the experimental observation of $D^0 \to \bar{K}^0 \phi$ gives a clear signal of the W-exchange diagram. The difference in the lifetimes of D^+, D^0, and F^+ is studied in terms of two-body decays.

I. INTRODUCTION

In this paper we analyze the experimental results for exclusive two-body decays of charmed mesons using the model-independent quark-diagram scheme. We show that the recent measurements of two-body exclusive decays of charm mesons D^+, D^0, and F^+ from the Mark III,[1] CLEO,[2] ARGUS,[3] HRS,[4] and TASSO Collaborations,[5] incorporating lifetime measurements, can allow us to determine the magnitudes and even the signs of some of the quark-diagram amplitudes for $P_c \to VP$ decays.[6] (Here, P_c represents D^+,D^0,F^+; V is a vector meson; and P is a pseudoscalar meson.) For $P_c \to PP$, we can also derive relations among various quark-diagram amplitudes. Using these experimentally determined quark amplitudes, we are able to make predictions for other charm-decay channels and test various theoretical models.

This paper is organized as follows. The quark-diagram scheme is presented in Sec. II. In Sec. III we incorporate the effects of SU(3) breaking and final-state interactions. In Sec. IV the status of the quark mixing matrix is briefly reviewed and its implications on charm-decay amplitudes is discussed. In Sec. V we analyze the experimental data of two-body decays of charmed mesons. Section VI is devoted to the discussions of various theoretical models and their tests. It is shown that the reaction $D^0 \to \bar{K}^0 \phi$ is a clear signal of W exchange. The difference in the lifetimes of D^+, D^0, and F^+ is studied in Sec. VII within two-body decays. Section VIII contains our conclusions. Some of the results in this paper have been reported by us in Refs. 6–8.

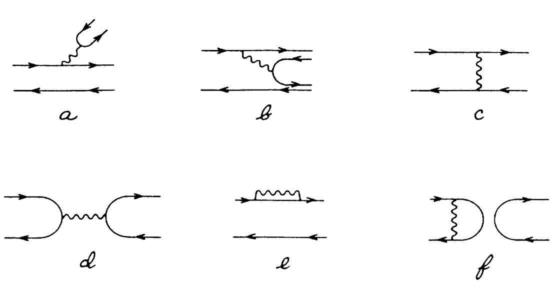

FIG. 1. The six quark diagrams for inclusive meson decay.

II. THE QUARK-DIAGRAM SCHEME

It has been established[9,10] that all weak decays of meson states can be classified according to six quark diagrams (see Figs. 1 and 2), the external W-emission diagram (a), the internal W-emission diagram (b), the W-exchange diagram (c), the W-annihilation diagram (d), the horizontal W-loop diagram (e), and the vertical W-loop diagram (f). This classification is independent of the strong-interaction schemes, and can incorporate any specific strong-interaction-model calculations. Thus all the strange, charm, bottom, and top particle decays can be expressed in terms of these six types of quark diagrams and the quark mixing matrix.[11,12] These quark-diagram amplitudes and the quark mixing matrix can be determined by comparing them with nonleptonic-decay experiments.

For the decay of a charmed meson into a pseudoscalar and a vector meson (PV), there are two different amplitudes for each quark diagram depending on whether or not the vector meson comes from the charmed-quark decay. We denote the primed amplitudes for the case that the vector meson arises from the decay of the c quark. It is obvious that this convention does not apply to the W-annihilation diagram since the final state $q_1\bar{q}_2$ comes from W annihilation. In this case we define d' (d) for the W-annihilation amplitude if the vector meson comes from q_1 ($\bar{q}_2$).

In the third column of Tables I and II we list quark-diagram amplitudes for all two-body decays of D^+, D^0,

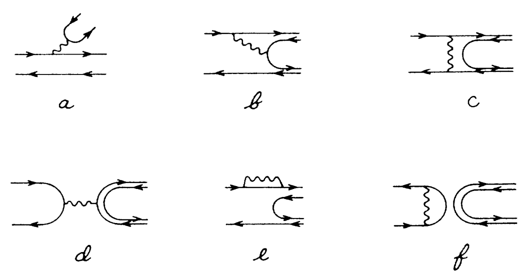

FIG. 2. The six quark diagrams for a meson decaying to two mesons.

TABLE I. Charm-meson decays into a vector boson and a pseudoscalar meson. (a) D^+ decays, (b) D^0 decays, and (c) F^+ decays.

Decays	Experimental branching ratio (%)	Amplitudes with SU(3) symmetry ($V_{us}V_{cs}^* \simeq -V_{ud}V_{cd}^* \simeq s_1c_1$ used)	Amplitudes with SU(3) breaking and final-state interaction
		(a)	
$\bar{K}^{*0}\pi^+$	3.0±1.9±1.7 (Ref. 1)	$(c_1)^2(a'+b') \rightarrow$	$(a'+b')e^{i\delta_{3/2}^{\bar{K}^*\pi}}$
$\rho^+\bar{K}^0$	12.2±2.8±1.9 (Ref. 1)	$(c_1)^2(a+b) \rightarrow$	$(a+b)e^{i\delta_{3/2}^{\rho\bar{K}}}$
$\phi\pi^+$	0.93±0.26±0.17 (Ref. 1)	$(s_1c_1)(b') \rightarrow$	$(b')e^{i\delta^{\phi\pi}}$,
$\omega\pi^+$		$-(1/\sqrt{2})(s_1c_1)(a'+b'+d+d') \rightarrow$	$[a'+b'+d+d'+(\delta e-\delta e')]e^{i\delta^{\omega\pi}}$
$\bar{K}^{*0}K^+$	0.53±0.24±0.14 (Ref. 1)	$(s_1c_1)(a'-\tilde{d}) \rightarrow$	$(a'-\tilde{d}+\delta e)e^{i\delta_1^{\bar{K}^*K}}$
$K^{*+}\bar{K}^0$		$(s_1c_1)(a-\tilde{d}') \rightarrow$	$(a-\tilde{d}'+\delta e')e^{i\delta_1^{\bar{K}^*K}}$
$\rho^+\pi^0$		$(1/\sqrt{2})(s_1c_1)(a+b-d+d')$ } $\rightarrow$	$[(a+b)(e^{i\delta_0^{\rho\pi}} \pm e^{i\delta_2^{\rho\pi}})/2+(a'+b')(e^{i\delta_0^{\rho\pi}} \mp e^{i\delta_2^{\rho\pi}})/2$
$\rho^0\pi^+$		$(1/\sqrt{2})(s_1c_1)(a'+b'+d-d')$ }	$\mp(d-d')e^{i\delta_1^{\rho\pi}} \pm 2(\delta e-\delta e')e^{i\delta_1^{\rho\pi}}]$
$\rho^+\eta_8$		$-(1/\sqrt{6})(s_1c_1)(a+3b+d+d') \rightarrow$	$[a+3b+d+d'+(\delta e+\delta e')]e^{i\delta_1^{\rho\eta_8}}$
$\rho^+\eta_0$		$-(1/\sqrt{3})(s_1c_1)(a+d+d') \rightarrow$	$[a+d+d'+(\delta e+\delta e')]e^{i\delta_1^{\rho\eta_0}}$
ϕK^+		$-(s_1)^2(\tilde{d}) \rightarrow$	$(\tilde{d})e^{i\delta^{\phi K}}$
ωK^+		$-(1/\sqrt{2})(s_1)^2(a'+d') \rightarrow$	$(a'+d')e^{i\delta^{\omega K}}$
$K^{*+}\eta_8$		$-(1/\sqrt{6})(s_1)^2(a+d-2d') \rightarrow$	$(a+d-2d')e^{i\delta^{K^*\eta_8}}$
$K^{*+}\eta_0$		$-(1/\sqrt{3})(s_1)^2(a+d+d') \rightarrow$	$(a+d+d')e^{i\delta^{K^*\eta_0}}$
ρ^+K^0		$-(s_1)^2(b+d') \rightarrow$	$[(b+d')+(a'-b-2d')\frac{1}{3}(1-e^{-i\Delta_{\rho K}})]e^{i\delta_{3/2}^{\rho K}}$
ρ^0K^+		$(1/\sqrt{2})(s_1)^2(a'-d') \rightarrow$	$[(a'-d')+(a'-b-2d')\frac{2}{3}(1-e^{-i\Delta_{\rho K}})]e^{i\delta_{3/2}^{\rho K}}$
$K^{*0}\pi^+$		Same as ρ^+K^0 but with primed and unprimed amplitudes interchanged.	
$K^{*+}\pi^0$		Same as ρ^0K^+ but with primed and unprimed amplitudes interchanged.	

TABLE I. (*Continued*).

Decays	Experimental branching ratio (%)	Amplitudes with SU(3) symmetry ($V_{us}V_{cs}^{*}\simeq -V_{ud}V_{cd}^{*}\simeq s_1c_1$ used)	Amplitudes with SU(3) breaking and final-state interaction
		(b)	
$\phi\bar{K}^0$	1.18±0.25±0.17 (Ref. 3) $1.1^{+0.7+0.4}_{-0.5-0.2}$ (Ref. 1) 1.18±0.40±0.17 (Ref. 2)	$(c_1)^2(c')\rightarrow$	$(\tilde{c}')e^{i\delta^{\phi\bar{K}}}$
$\omega\bar{K}^0$	3.8±1.5±1.0 (Ref. 1)	$(1/\sqrt{2})(c_1)^2(b+c)\rightarrow$	$(b+c)e^{i\delta^{\omega\bar{K}}}$
$K^{*-}\pi^+$	7.8±1.2±0.9 (Ref. 1) 7.1±1.6±1.3 (Ref. 1)	$(c_1)^2(a'+c')\rightarrow$	$[(a'+c')-\frac{1}{3}(a'+b')(1-e^{i\Delta_{\bar{K}^*\pi}})]e^{i\delta_{1/2}^{\bar{K}^*\pi}}$
$\bar{K}^{*0}\pi^0$	2.1±0.9±0.6 (Ref. 1)	$(1/\sqrt{2})(c_1)^2(b'-c')\rightarrow$	$[(b'-c')-\frac{2}{3}(a'+b')(1-e^{i\Delta_{\bar{K}^*\pi}})]e^{i\delta_{1/2}^{\bar{K}^*\pi}}$
$\bar{K}^{*0}\eta_8$		$(1/\sqrt{6})(c_1)^2(b'+c'-2c)\rightarrow$	$(b'+c'-2\tilde{c})e^{i\delta^{\bar{K}^*\eta_8}}$
$\bar{K}^{*0}\eta_0$		$(1/\sqrt{3})(c_1)^2(b'+c'+c)\rightarrow$	$(b'+c'+\tilde{c})e^{i\delta^{\bar{K}^*\eta_0}}$
ρ^+K^-	13.7±1.3±1.5 (Ref. 1)	$(c_1)^2(a+c)\rightarrow$	$[(a+c)-\frac{1}{3}(a+b)(1-e^{i\Delta_{\rho\bar{K}}})]e^{i\delta_{1/2}^{\rho\bar{K}}}$
$\rho^0\bar{K}^0$	1.3±0.4±0.3 (Ref. 1)	$(1/\sqrt{2})(c_1)^2(b-c)\rightarrow$	$[(b-c)-\frac{2}{3}(a+b)(1-e^{i\Delta_{\rho\bar{K}}})]e^{i\delta_{1/2}^{\rho\bar{K}}}$
$\phi\pi^0$		$(1/\sqrt{2})(s_1c_1)(b')\rightarrow$	$(b')e^{i\delta^{\phi\pi}}$
$\phi\eta_8$		$(1/\sqrt{6})(s_1c_1)(b'-2\tilde{c}'-2\tilde{c})\rightarrow$	$[b'-2\tilde{c}'-2\tilde{c}-2(\delta f+\delta f')]e^{i\delta^{\phi\eta_8}}$
$\phi\eta_0$		$(1/\sqrt{3})(s_1c_1)(b'+\tilde{c}'+\tilde{c})\rightarrow$	$[b'+\tilde{c}'+\tilde{c}+(\delta f+\delta f')]e^{i\delta^{\phi\eta_0}}$
$\omega\pi^0$		$\frac{1}{2}(s_1c_1)(b-b'+c+c')\rightarrow$	$(b-b'+c+c'+\delta e+\delta e')e^{i\delta^{\omega\pi}}$
$\omega\eta_8$		$(1/\sqrt{12})(s_1c_1)(-3b-b'-c-c')\rightarrow$	$(-3b-b'-c-c'+\delta e+\delta e'+\delta f+\delta f')e^{i\delta^{\omega\eta_8}}$
$\omega\eta_0$		$(1/\sqrt{6})(s_1c_1)(b+c+c')\rightarrow$	$[b+c+c'+(\delta e+\delta e'+2\delta f+2\delta f')]e^{i\delta^{\omega\eta_0}}$
$\bar{K}^{*0}K^0$		$-(s_1c_1)(c-c')\rightarrow$	$\{(c-c')-\delta f-\delta f'$ $+[(a'+c')+(\tilde{c}-\tilde{c}')-\delta e](1-e^{i\Delta_{\bar{K}^*K}})/2\}e^{i\delta_0^{\bar{K}^*K}}$

TABLE I. (*Continued*).

Decays	Experimental branching ratio (%)	Amplitudes with SU(3) symmetry ($V_{us}V_{cs}^* \simeq -V_{ud}V_{cd}^* \simeq s_1c_1$ used)	Amplitudes with SU(3) breaking and final-state interaction
$K^{*0}\overline{K}^0$	≤ 0.83 (Ref. 1)	Same as $\overline{K}^{*0}K^0$ but with primed and unprimed amplitudes interchanged.	
$K^{*-}K^+$	$1.02\pm0.47\pm0.21$ (Ref. 1)	$(s_1c_1)(a'+c')\rightarrow$	$\{(a'+c')+\delta e+\delta f+\delta f'$ $-[(a'+c')+(\tilde{c}-c')-\delta e](1-e^{i\Delta_{\overline{K}^*K}})/2\}e^{i\delta_0^{\overline{K}^*K}}$
$K^{*+}K^-$		Same as $K^{*-}K^+$ but with primed and unprimed amplitudes interchanged.	
$\rho^-\pi^+$ $\rho^+\pi^-$		$(s_1c_1)[-(a'+c');-(a+c)]\rightarrow$	$[-a'(2e^{i\delta_0^{\rho\pi}}\pm 3e^{i\delta_1^{\rho\pi}}+e^{i\delta_2^{\rho\pi}})/3$ $-a(2e^{i\delta_0^{\rho\pi}}\mp 3e^{i\delta_1^{\rho\pi}}+e^{i\delta_2^{\rho\pi}})/3$ $-(b+b')(e^{i\delta_0^{\rho\pi}}-e^{i\delta_2^{\rho\pi}})/3$ $-(c'-\delta e)(e^{i\delta_0^{\rho\pi}}\pm 3e^{i\delta_1^{\rho\pi}}+2e^{i\delta_2^{\rho\pi}})/3$ $-(c-\delta e')(e^{i\delta_0^{\rho\pi}}\mp 3ei^{i\delta_1^{\rho\pi}}+2e^{i\delta_2^{\rho\pi}})/3$ $-(\delta_f+\delta_{f'})e^{i\delta_0^{\rho\pi}}]$
$\rho^0\pi^0$		$\frac{1}{2}(s_1c_1)(b+b'-c-c')\rightarrow$	$[(a+a')(e^{i\delta_0^{\rho\pi}}-e^{i\delta_2^{\rho\pi}})/3$ $+(b+b')(e^{i\delta_0^{\rho\pi}}+2e^{i\delta_2^{\rho\pi}})/6$ $+(c+c')(e^{i\delta_0^{\rho\pi}}-4e^{i\delta_2^{\rho\pi}})/6$ $-(\delta e+\delta e')(e^{i\delta_0^{\rho\pi}}-4e^{i\delta_2^{\rho\pi}})/6$ $+(\delta f+\delta f')e^{i\delta_0^{\rho\pi}}]$
$\rho^0\eta_8$		$(1/\sqrt{12})(s_1c_1)(-3b+b'+c+c')\rightarrow$	$(-3b+b'+c+c'+\delta e+\delta e')e^{i\delta^{\rho\eta_8}}$
$\rho^0\eta_0$		$(1/\sqrt{6})(s_1c_1)(b'+c+c')\rightarrow$	$(b'+c+c'+\delta e+\delta e')e^{i\delta^{\rho\eta_8}}$
ϕK^0		$-(s_1)^2(\tilde{c})\rightarrow$	$(\tilde{c})e^{i\delta^{\phi K}}$
ωK^0		$-(1/\sqrt{2})(s_1)^2(b+c')$	$(b+c')e^{i\delta^{\omega\overline{K}}}$
$K^{*+}\pi^-$		$-(s_1)^2$ (same as $\overline{K}^*\pi^+$ but with primed and unprimed amplitudes interchange)	
$K^{*0}\pi^0$		$-(1/\sqrt{2})(s_1)^2$ (same as $\overline{K}^{*0}\pi^0$ but with primed and unprimed amplitudes interchange)	
$K^{*0}\eta_8$		$-(1/\sqrt{6})(s_1)^2(b'+c-2\tilde{c}')\rightarrow$	$(b'+c-2\tilde{c}')e^{i\delta_{1/2}^{K^*\eta_8}}$
$K^{*0}\eta_0$		$-(1/\sqrt{3})(s_1)^2(b'+c+\tilde{c}')\rightarrow$	$(b'+c+\tilde{c}')e^{i\delta^{K^*\eta_0}}$
ρ^-K^+		$-(s_1)^2[(a'+c')]\rightarrow$	$[(a'+c')-\frac{2}{3}(a'+b)(1-e^{i\Delta_{\rho K}})]e^{i\delta_{1/2}^{\rho K}}$
ρ^0K^0		$-(1/\sqrt{2})(s_1)^2[(b-c')]\rightarrow$	$[(b-c')-\frac{1}{3}(a'+b)(1-e^{i\Delta_{\rho K}})]e^{i\delta_{1/2}^{\rho K}}$

TABLE I. (*Continued*).

Decays	Experimental branching ratio (%)	Amplitudes with SU(3) symmetry ($V_{us}V_{cs}^* \simeq -V_{ud}V_{cd}^* \simeq s_1c_1$ used)	Amplitudes with SU(3) breaking and final-state interaction
		(c)	
$\phi\pi^+$	3.3±1.1 (Ref. 4) 4.4 (Ref. 2) 13.0±3.0±4.0 (Ref. 5)	$(c_1)^2(a')\rightarrow$	$(a')e^{i\delta^{\phi\pi}}$
$\omega\pi^+$		$(1/\sqrt{2})(c_1)^2(d+d')\rightarrow$	$(d+d')e^{i\delta_1^{\omega\pi}}$
$K^{*+}\bar{K}^0$		$(c_1)^2(b+d')\rightarrow$	$(b+\tilde{d}')e^{i\delta_1^{K^*K}}$
$\bar{K}^{*0}K^+$		$(c_1)^2(b'+d)\rightarrow$	$(b'+\tilde{d})e^{i\delta_1^{K^*K}}$
$\rho^+\pi^0$		$(1/\sqrt{2})(c_1)^2(d-d')\rightarrow$	$(d-d')e^{i\delta_1^{\rho\pi}}$
$\rho^0\pi^+$		$(1/\sqrt{2})(c_1)^2(d'-d)\rightarrow$	$(d'-d)e^{i\delta_1^{\rho\pi}}$
$\rho^+\eta_8$		$(1/\sqrt{6})(c_1)^2(-2a+d+d')\rightarrow$	$(-2a+d+d')e^{i\delta^{\rho\eta_8}}$
$\rho^+\eta_0$		$(1/\sqrt{3})(c_1)^2(a+d+d')\rightarrow$	$(a+d+d')e^{i\delta^{\rho\eta_0}}$
ϕK^+		$(s_1c_1)(a'+b'+\tilde{d})\rightarrow$	$(a'+b'+\tilde{d}+\delta e)e^{i\delta^{\phi K}}$
ωK^+		$-(1/\sqrt{2})(s_1c_1)(b'-d')\rightarrow$	$(b'-d'-\delta e')e^{i\delta^{\omega K}}$
$K^{*0}\pi^+$		$-(s_1c_1)(a'-d)\rightarrow$	$[(a'-d-\delta e)-(2a'+b-d-\delta e)\frac{1}{3}(1-e^{-i\Delta_{K^*\pi}})]e^{\delta_{3/2}^{K^*\pi}}$
$K^{*+}\pi^0$		$-(1/\sqrt{2})(s_1c_1)(b-d)\rightarrow$	$[(b-d+\delta e)-(2a'+b-d-\delta e)\frac{2}{3}(1-e^{-i\Delta_{K^*\pi}})]e^{\delta_{3/2}^{K^*\pi}}$
ρ^+K^0		$-(s_1c_1)(a-d')\rightarrow$	$[(a-d')-(2a+b'-d'-\delta e')\frac{1}{3}(1-e^{-i\Delta_{K^*\pi}})]e^{i\delta_{3/2}^{\rho K}}$
ρ^0K^+		$(1/\sqrt{2})(s_1c_1)(b'+d')\rightarrow$	$[(b'+d')-(2a+b'-d'-\delta e')\frac{2}{3}(1-e^{-i\Delta_{\rho K}})]e^{i\delta_{3/2}^{\rho K}}$
$K^{*+}\eta_8$		$(1/\sqrt{6})(s_1c_1)(-2a+3b+d-2\tilde{d}')\rightarrow$	$(-2a-3b+d+\tilde{d}'+\delta e-2\delta e')e^{i\delta^{K^*\eta_8}}$
$K^{*+}\eta_0$		$(1/\sqrt{3})(s_1c_1)(a+d+\tilde{d}')\rightarrow$	$[a+d+\tilde{d}'+(\delta e+\delta e')]e^{i\delta^{K^*\eta_8}}$
$K^{*+}K^0$		$-(s_1)^2(a+b)\rightarrow$	$[(a+b)+\frac{1}{2}(a'+b'-a-b)(1-e^{-i\Delta_{K^*K}})]e^{i\delta_1^{K^*K}}$
$K^{*0}K^+$		$-(s_1)^2(a'+b')\rightarrow$	$[(a'+b')-\frac{1}{2}(a'+b'-a-b)(1-e^{-i\Delta_{K^*K}})]e^{i\delta_1^{K^*K}}$

TABLE II. Charm-meson decays to two pseudoscalars. (a) D^+ decays, (b) D^0 decays, and (c) F^+ decays.

Decays	Experimental branching ratio (%)	Amplitudes with SU(3) symmetry ($V_{us}V_{cs}^* \simeq -V_{ud}V_{cd}^* = s_1c_1$ used)	Amplitudes with SU(3) breaking and final-state interaction
		(a)	
$\bar{K}^0\pi^+$	3.5±0.5±0.4 (Ref. 1)	$(c_1)^2(a+b) \to$	$(a+b)e^{i\delta_{1/2}^{\bar{K}\pi}}$
$\bar{K}^0K^+$	1.11±0.34±0.21 (Ref. 1)	$(s_1c_1)(a-d) \to$	$(a-\tilde{d}+\delta e)e^{i\delta_1^{KK}}$
$\pi^0\pi^+$	$\leq$0.53 (Ref. 1)	$(1/\sqrt{2})(s_1c_1)(a+b) \to$	$(a+b)e^{i\delta_2^{\pi\pi}}$
$\eta_8\pi^+$		$-(1/\sqrt{6})(s_1c_1)[(a+b)+2(b+d)] \to$	$[a+b+2(b+d)+2\delta e]e^{i\delta^{\eta\pi}}$
$\eta_0\pi^+$		$-(1/\sqrt{3})(s_1c_1)(a+2d) \to$	$[(a+2d)+2\delta e]e^{i\delta^{\eta\pi}}$
$K^0\pi^+$		$-(s_1)^2(b+d) \to$	$[(b+d)+(a-b-2d)\frac{1}{3}(1-e^{-i\Delta_{K\pi}})]e^{i\delta_{3/2}^{K\pi}}$
$K^+\pi^0$		$(1/\sqrt{2})(s_1)^2(a-d) \to$	$[(a-d)+(a-b-2d)\frac{2}{3}(1-e^{-i\Delta_{K\pi}})]e^{i\delta_{3/2}^{K\pi}}$
$K^+\eta_8$		$-(1/\sqrt{6})(s_1)^2(a-d) \to$	$(a-d)e^{i\delta^{K\eta_8}}$
$K^+\eta_0$		$-(1/\sqrt{3})(s_1)^2(a+2d) \to$	$(a+2d)e^{i\delta^{K\eta_0}}$
		(b)	
$K^-\pi^+$	4.9±0.4±0.4 (Ref. 1)	$(c_1)^2(a+c) \to$	$[(a+c)-(a+b)\frac{1}{3}(1-e^{i\Delta_{K\pi}})]e^{i\delta_{1/2}^{K\pi}}$
$\bar{K}^0\pi^0$	2.2±0.4±0.2 (Ref. 1)	$(1/\sqrt{2})(c_1)^2(b-c) \to$	$[(b-c)-(a+b)\frac{2}{3}(1-e^{i\Delta_{K\pi}})]e^{i\delta_{1/2}^{K\pi}}$
$\bar{K}^0\eta_8$		$(1/\sqrt{6})(c_1)^2(b-c) \to$	$(b+c-2\tilde{c})e^{i\delta^{K\eta_8}}$
$\bar{K}^0\eta_0$		$(1/\sqrt{3})(c_1)^2(b+2c) \to$	$(b+c+\tilde{c})e^{i\delta^{K\eta_0}}$
$\bar{K}^0\eta$	1.8±0.8±0.3 (Ref. 1)	$\cos\theta\, A(D^0\to\bar{K}^0\eta_8)+\sin\theta\, A(D^0\to\bar{K}^0\eta_0)$, $\theta\approx 20°$	
$K^0\bar{K}^0$	$\leq$0.62 (Ref. 1)	$0 \to$	$[(c-\tilde{c}-2\delta f)+(a+\tilde{c}-\delta e)\frac{1}{2}(1-e^{i\Delta_{K\bar{K}}})]e^{i\delta_0^{K\bar{K}}}$
K^-K^+	0.60±0.10±0.08 (Ref. 1)	$(s_1c_1)(a+c) \to$	$[(a+c)+(\delta e+2\delta f)-(a+\tilde{c}-\delta e)\frac{1}{2}(1-e^{i\Delta_{K\bar{K}}})]e^{i\delta_0^{K\bar{K}}}$
$\pi^+\pi^-$	0.16±0.05±0.03 (Ref. 1)	$-(s_1c_1)(a+c) \to$	$[(a+c)+(\delta e+2\delta f)-(a+b)\frac{1}{3}(1-e^{i\Delta_{\pi\pi}})]e^{i\delta_0^{\pi\pi}}$

TABLE II. (*Continued*).

Decays	Experimental branching ratio (%)	Amplitudes with SU(3) symmetry ($V_{us}V_{cs}^* \simeq -V_{ud}V_{cd}^* = s_1c_1$ used)	Amplitudes with SU(3) breaking and final-state interaction
		(b)	
$\pi^0\pi^0$		$(\sqrt{2})\frac{1}{2}(s_1c_1)(b-c)\rightarrow$	$[(b-c)+(\delta e+2\delta f)-(a+b)\frac{2}{3}(1-e^{i\Delta_{\pi\pi}})]e^{i\delta_0^{\pi\pi}}$
$\eta_8\eta_8$		$-(\sqrt{2})\frac{1}{2}(s_1c_1)(b-c)\rightarrow$	$[(b-c)+2(-\delta c/6-\delta e/6-\delta f)]e^{i\delta^{\eta_8\eta_8}}$
$\pi^0\eta_8$		$-(1/\sqrt{3})(s_1c_1)(b-c)\rightarrow$	$[(b-c-\delta e)]e^{i\delta^{\pi\eta_0}}$
$\pi^0\eta_0$		$(1/\sqrt{6})(s_1c_1)(b+2c)\rightarrow$	$(b+2c+2\delta e)^{i\delta^{\pi\eta_0}}$
$\eta_8\eta_0$		$(1/\sqrt{2})(s_1c_1)(b+2c)\rightarrow$	$[(b+2c+2\delta e)]e^{i\delta^{\eta_0\eta_8}}$
$\eta_0\eta_0$		$(\sqrt{2})0\rightarrow$	$(\delta c+\delta e+3\delta f)e^{i\delta^{\eta_0\eta_0}}$
$K^+\pi^-$		$-(s_1)^2(a+c)\rightarrow$	Same as for $K^-\pi^+$
$K^0\pi^0$		$-(1/\sqrt{2})(s_1)^2(b-c)\rightarrow$	Same as for $\bar{K}^0\pi^0$
$K^0\eta_8$		$-(1/\sqrt{6})(s_1)^2(b-c)\rightarrow$	Same as for $\bar{K}^0\eta_8$
$K^0\eta_0$		$-(1/\sqrt{3})(s_1)^2(b+2c)\rightarrow$	Same as for $\bar{K}^0\eta_0$
		(c)	
$\bar{K}^0K^+$		$(c_1)^2(b+d)\rightarrow$	$(b+d)e^{i\delta_1^{K\bar{K}}}$
$\pi^0\pi^+$		$0\rightarrow$	0
$\eta_8\pi^+$		$-(\sqrt{2}/\sqrt{3})(c_1)^2(a-d)\rightarrow$	$(a-d)e^{i\delta^{\pi\eta_8}}$
$\eta_0\pi^+$		$(1/\sqrt{3})(c_1)^2(a+2d)\rightarrow$	$(a+2d)e^{i\delta^{\pi\eta_0}}$
$K^0\pi^+$		$-(s_1c_1)(a-d)\rightarrow$	$[(a-d-\delta e)-(2a+b-d-\delta e)\frac{1}{3}(1-e^{-i\Delta_{K\pi}})]e^{i\delta_{3/2}^{K\pi}}$
$K^+\pi^0$		$(1/\sqrt{2})(s_1c_1)(b+d)\rightarrow$	$[(b+d+\delta e)-(2a+b-d-\delta e)\frac{2}{3}(1-e^{-i\Delta_{K\pi}})]e^{i\delta_{3/2}^{K\pi}}$
$K^+\eta_8$		$-(1/\sqrt{6})(s_1c_1)[2(a+b)+(b+d)]\rightarrow$	$[2(a+b)+b-d+2\tilde{d}+\delta e]e^{i\delta^{K\eta_8}}$
$K^+\eta_0$		$(1/\sqrt{3})(s_1c_1)(a+2d)\rightarrow$	$(a+d+\tilde{d}+2\delta e)e^{i\delta^{K\eta_0}}$
K^0K^+		$-(s_1)^2(a+b)\rightarrow$	$[(a+b)-(b+d)(1/\sqrt{2})(1-e^{-i\Delta_{KK}})]e^{i\delta_1^{KK}}$

and F^+ without SU(3) breaking and final-state interactions. It is of interest to note that in the absence of these effects $D^0 \to K^0\bar{K}^0, \eta_0\eta_0$ are forbidden. Also, as will be discussed in Sec. IV, the loop diagrams do not contribute to Cabibbo-suppressed decays. In Sec. III we will discuss how to incorporate the effects of the final-state interaction and SU(3) breaking in the quark-diagram scheme.

III. SU(3) BREAKING AND FINAL-STATE INTERACTIONS

As we shall see in Sec. V, a comparison of the quark-diagram analysis with the data shows the existence of SU(3)-breaking and final-state-interaction effects. Hence we incorporate the SU(3)-breaking effects. For instance, we distinguish the contributions from the strange quarks in the W-loop diagrams from those from down quarks. As a result, the W-loop contribution to charm decays is a manifestation of SU(3)-breaking effects. In Tables I and II the W-exchange or W-annihilation diagrams with strange quark-antiquark pair creations are denoted by a tilde to distinguish from those coming from up or down quark-antiquark creations. For simplicity, we do not include SU(3)-breaking effects in the external particles except in phase-space calculations. For example, we do not distinguish the amplitude a in $D^0 \to K^+K^-$ and $K^-\pi^+$.

When the strong interaction is turned on, the weak decay amplitudes are modified by the rescattering effects required by unitarity of the S matrix; the absorptive parts of the amplitudes arise from the final-state strong interaction. The final-state interaction is particularly troublesome for charm decays since some resonances are known to exist at energies close to the mass of the charmed meson.[13–15] Consequently, the inelastic scattering effects are crucial for understanding the pattern of charm weak decays. Hence not only phase shifts but also inelasticities ought to be taken into account in the quark-diagram analysis. To do this, we introduce an explicit factor $e^{i\delta}$ for each isospin partial-wave amplitudes. These phase shifts δ in general have both real and imaginary parts; the imaginary component indicates the inelastic effect. Because of the inelasticity, the phase of the weak amplitude is in general not the same as the strong-interaction phase shift.

In Table III we give the relations between quark-diagram amplitudes and isospin partial-wave amplitudes A_I for some decay modes. Introducing the phase factor $e^{i\delta_I}$ to each A_I and considering the SU(3)-breaking effects, we obtain the quark-diagram amplitudes in the last column of Tables I and II for all two-body decays of charmed mesons.

IV. STATUS OF QUARK MIXING MATRIX AND IMPLICATIONS ON CHARM-DECAY AMPLITUDES

Recently our knowledge of the quark mixing matrix[16] has dramatically improved owing to the b-lifetime measurement[17,18] and the $\Gamma(b \to u)/\Gamma(b \to c) < 0.05$ bound from the measurements of semileptonic decays of b-flavored particles.[19] Also these results show that the quark mixing matrix can be parametrized most conveniently in the form[16]

TABLE III. Relation between quark-diagram amplitudes and partial-wave isospin amplitudes.

$D^0 \to K^-\pi^+$	$\frac{1}{3}(c_1)^2(\sqrt{2}A_{1/2}+A_{3/2})$	$A_{1/2}=(1/\sqrt{2})(2a-b+3c)e^{i\delta^{K\pi}_{1/2}}$
$\to \bar{K}^0\pi^0$	$\frac{1}{3}(c_1)^2(-A_{1/2}+\sqrt{2}A_{3/2})$	$A_{3/2}=(a+b)e^{i\delta^{K\pi}_{3/2}}$
$D^+ \to \bar{K}^0\pi^+$	$(c_1)^2A_{3/2}$	
$D^+ \to K^0\pi^+$	$-(1/\sqrt{3})(s_1)^2(\sqrt{2}A'_{1/2}+A'_{3/2})$	$A'_{1/2}=(1/\sqrt{6})(a+2b+d)e^{i\delta^{K\pi}_{1/2}}$
$\to K^+\pi^0$	$-(1/\sqrt{3})(s_1)^2(-A'_{1/2}+\sqrt{2}A'_{3/2})$	$A'_{3/2}=(1/\sqrt{3})(-a+b+2d)e^{i\delta^{K\pi}_{3/2}}$
$D^0 \to \pi^+\pi^-$	$(1/\sqrt{3})s_1c_1(\sqrt{2}B_0+B_2)$	$B_0=-(3/\sqrt{6})(2a/3-b/3+c-\delta e-2\delta f)e^{i\delta^{\pi\pi}_0}$
$\to \pi^0\pi^0$	$(1/\sqrt{3})s_1c_1(B_0-\sqrt{2}B_2)$	
$D^+ \to \pi^0\pi^+$	$-(\sqrt{3}/\sqrt{2})s_1c_1B_2$	$B_2=-(1/\sqrt{3})(a+b)e^{i\delta^{\pi\pi}_2}$
$D^0 \to K^+K^-$	$(1/\sqrt{2})s_1c_1(E_0+E_1)$	$E_0=(1/\sqrt{2})(a+2c-\tilde{c}-\delta e-4\delta f)e^{i\delta^{K\bar{K}}_0}$
$\to K^0\bar{K}^0$	$(1/\sqrt{2})s_1c_1(E_0-E_1)$	$E_1=(1/\sqrt{2})(a+\tilde{c}'-\delta e)e^{i\delta^{KK}_1}$
$D^+ \to K^+\bar{K}^0$	$s_1c_1E'_1$	$E'_1=[a-d+(-\delta e)]e^{i\delta^{K\bar{K}}_1}$
$F^+ \to K^0\pi^+$	$-(1/\sqrt{3})s_1c_1(\sqrt{2}A''_{1/2}+A''_{3/2})$	$A''_{1/2}=(1/\sqrt{6})(2a+b-d-\delta e)e^{i\delta^{K\pi}_{1/2}}$
$\to K^+\pi^0$	$-(1/\sqrt{3})s_1c_1(-A''_{1/2}+\sqrt{2}A''_{3/2})$	$A''_{3/2}=(1/\sqrt{3})(a-b-2d-2\delta e)e^{i\delta^{K\pi}_{3/2}}$
$F^+ \to \bar{K}^0K^+$	$(c_1)^2E''_1$	$E''_1=(b+d)e^{i\delta^{KK}_1}$
$\to K^0K^+$	$-(1/\sqrt{2})(s_1)^2(E''_0+E''_1)$	$E''_0=[\sqrt{2}a+(\sqrt{2}-1)b-d]e^{i\delta^{KK}_0}$

$$V=\begin{bmatrix}1 & 0 & 0\\ 0 & c_y & s_y\\ 0 & -s_y & c_y\end{bmatrix}\begin{bmatrix}c_z & 0 & s_z e^{-i\phi}\\ 0 & 0 & 0\\ -s_z e^{+i\phi} & 0 & c_z\end{bmatrix}\begin{bmatrix}c_x & s_x & 0\\ -s_x & c_x & 0\\ 0 & 0 & 1\end{bmatrix}$$

$$\begin{array}{c} \quad d \qquad\qquad\qquad s \qquad\qquad\qquad b \end{array}$$

$$=\left[\begin{array}{ccc} c_x c_z & s_x c_z & s_z e^{-i\phi}\\ -s_x c_y - c_x s_y s_z e^{i\phi} & c_x c_y - s_x s_y s_z e^{i\phi} & s_y c_z\\ s_x s_y - c_x c_y s_z e^{i\phi} & -c_x s_y - s_x c_y s_z e^{i\phi} & c_y c_z \end{array}\right]\begin{array}{c} u\\ c\\ t\end{array}$$

$$\underset{(s_i\approx 0)}{\approx}\left[\begin{array}{ccc} 1 & s_x & s_z e^{-i\phi}\\ -s_x - s_y s_z e^{i\phi} & 1 - s_x s_y s_z e^{i\phi} & s_y\\ s_x s_y - s_z e^{i\phi} & -s_y - s_x s_z e^{+i\phi} & 1 \end{array}\right] . \tag{4.1}$$

An important feature of the three-generation quark mixing matrix is that all *CP* noninvariance in decays of those three-generation quark particles are described by one universal phase-convention-independent parameter:

$$X_{CP}=s_x s_y s_z s_\phi c_x c_y c_z{}^2 \approx 10^{-5}$$

$$\text{if } m_t \approx 50 \text{ GeV is used in fitting } \epsilon . \tag{4.2}$$

Because of this property we can put the phase factor at the most convenient place, i.e., where the matrix element is the smallest.[17] This has the advantage that for all practical purposes without involving *CP* noninvariance, the matrix can be considered to be real. Further it can be demonstrated that this parametrization is also the simplest in considering *CP* noninvariance of partial-decay-rate differences.

Another important feature of this parametrization is that it takes advantage of all the experimental information: each sine is directly related to one type of experiment:[18]

$$c_x \approx 0.9737 \text{ from nuclear } \beta \text{ decays ,} \tag{4.3a}$$

$$s_x \approx 0.22 \text{ determined from strange-particle decays ,} \tag{4.3b}$$

$$s_y \approx 0.05 \text{ from } b\text{-particle lifetime ,} \tag{4.3c}$$

$$s_z \lesssim 0.01 \text{ from bounds on } (b\to u)/(b\to c) . \tag{4.3d}$$

This is the inherited convenience from the original Maiani's parametrization,[19] though the rotation order was found originally in Ref. 16 and with final form in the real angles in agreement with Maiani's. The subscripts x,y,z are used because that was the order of experimental measurements and easiest to remember. Such parametrization via sequences of rotations and put phases at the furthest corners can be generalized to cases with higher generating of quarks.[20,21]

Following the convention that these matrix elements are positive and using the unitarity of the mixing matrix, assuming three generations of quarks, we obtain the following full quark mixing matrix (ignoring the imaginary part for this work which deals only with the absolute values of V_{ij}; for a convenient parametrization see Refs. 21 and 22):

$$V=\begin{array}{c}\begin{array}{ccc} d & s & b\end{array}\\ \left[\begin{array}{ccc} 0.9737 & 0.228 & \lesssim 0.0082\\ -0.23 & 0.98 & 0.059\\ 0.01 & -0.06 & 1\pm(\sim 10^{-2})\end{array}\right]\end{array}\begin{array}{c} u\\ c\ .\\ t\end{array} \tag{4.4}$$

It is interesting to note that owing to the severe suppression of the off-diagonal element in the third column and row, the upper left two-by-two matrix is almost unitary, i.e.,

$$V_{cd}/V_{cs} = -V_{us}/V_{ud} \approx -s_1/c_1 . \tag{4.5}$$

With such a well-determined mixing matrix, the unknown quantities remaining in the classification for all meson decays are the six amplitudes. An interesting result of this is the prediction:

$$\begin{aligned} |A(D^+\to\pi^0\pi^+)/A(D^+\to\bar{K}^0\pi^+)|^2 &= \tfrac{1}{2}|V_{cd}/V_{cs}|^2\\ &= \tfrac{1}{2}(s_1/c_1)^2\\ &= 0.028 . \end{aligned} \tag{4.6}$$

It is very important to measure this quantity. Note that Eqs. (4.5) and (4.6) are strictly true under the three-generation assumption. Deviation from them may have the very interesting implication of the existence of the fourth generation. In that case, all the analyses done in this paper have to be reassessed.

Equation (4.5) has another very striking implication for the charm-meson decays: in the SU(3) limit, contributions in the W loop from the strange and down quarks cancel each other.[23,24] This explains why we drop all the e and f amplitudes in the third column of Tables I and II. It also becomes clear that the penguin amplitude in charm decays does not play an essential role as in the kaon decays since it must vanish in the limit of SU(3) symmetry.

V. QUARK-DIAGRAM ANALYSIS

A. The decays $P_c \to VP$

Recently, many two-body decays of charm particles have been elegantly measured,[1–5] which we list in Tables I and II. Here we shall put the quark-diagram formalism to use, analyzing all existing charm two-body decay data and discussing their implication for various theoretical-model calculations.

We begin with the *PV* decays because of the relative simplicity in presenting the discussion, though the data of *PV* decays are not yet as good as those for some of the *PP* decays. The simplicity in discussing the *PV* decays comes from the purity of the quark contents in ϕ and ω. Many *PV* decays are given by one type of amplitude, as shown in Table I: e.g., $F^+\to\phi\pi^+$ ($\propto a'$), $D^+\to\phi\pi^+$ ($\propto b'$), $D^0\to\phi\bar{K}^0$ ($\propto \tilde{c}'$), respectively. The decay amplitude is of the form

$$M(P_c\to PV)=\epsilon^\mu p_\mu A , \tag{5.1}$$

where ϵ^μ is the polarization vector of the vector meson V, p_μ is the four-momentum of the charmed meson, and the amplitude A can be expressed in terms of the quark-diagram amplitudes as shown in Table I. The partial decay rate is then given by

$$\Gamma = |A^2| \frac{\{[m_c^2-(m_V+m_P)^2][m_c^2-(m_V-m_P)^2]\}^{3/2}}{64\pi m_c^3 m_V^2} . \tag{5.2}$$

From the decay rates of $F^+\to\phi\pi^+$, $D^+\to\phi\pi^+$, and $D^0\to\phi\bar{K}^0$ we can obtain the absolute values of the following quark-diagram amplitudes:

$$|a'| = (2.50\pm0.42)\times10^{-6} ,$$
$$|b'| = (3.67\pm0.51)\times10^{-6} , \tag{5.3}$$
$$|\tilde{c}'| = (1.92\pm0.29)\times10^{-6} ,$$

where we have assumed no inelasticity in these channels. (The discussion of the inelasticity in $D^0\to\bar{K}^0\phi$ is made in Sec. VI.) To obtain the decay widths from the measured branching ratios as given in Table I, the following charm lifetimes are used:

$$\tau(D^+) = (8.8^{+1.0}_{-0.8})\times10^{-13}\ \text{sec} ,$$
$$\tau(D^0) = (4.3^{+0.4}_{-0.3})\times10^{-13}\ \text{sec} , \tag{5.4}$$
$$\tau(F^+) = (2.8^{+1.4}_{-0.7})\times10^{-13}\ \text{sec} .$$

Among the three measurements of $F^+\to\phi\pi^+$ we have used the branching ratio $(3.3\pm1.1)\%$ (Ref. 4), which is consistent with various theoretical estimates[25,26] and the upper bound calculated from current algebra.[27] The $\tilde{c}'$ in Eq. (5.3) corresponds to $B(D^0\to\phi\bar{K}^0) = (1.18\pm0.25)\%$.

Since $D^+\to\bar{K}^{*0}\pi^+$ $[\propto(a'+b')]$ is an exotic channel, it is reasonable to neglect the inelasticity owing to the absence of isospin-$\frac{3}{2}$ resonances. Therefore, the rate gives

$$|a'+b'| = (1.16\pm0.37)\times10^{-6} . \tag{5.5}$$

From Eqs. (5.3) and (5.5), it is evident that a' and b' are of opposite signs, and the $|a'|$ and $|b'|$ obtained from $F^+\to\phi\pi^+, D^+\to\phi\pi^+$ are in excellent agreement with $a'+b'$ determined from $D^+\to\bar{K}^{*0}\pi^+$. Thus we have

$$a' = (2.50\pm0.42)\times10^{-6} ,$$
$$b' = -(3.67\pm0.51)\times10^{-6} . \tag{5.6}$$

Letting $\Delta_{\bar{K}^*\pi} = \delta_{1/2}^{\bar{K}^*\pi} - \delta_{3/2}^{\bar{K}^*\pi}$, and assuming $e^{i\delta_{1/2}^{\bar{K}^*\pi}} = 1$, we obtain the following two solutions from the measurements of $D^0\to K^{*-}\pi^+$ and $\bar{K}^{*0}\pi^0$:

$$(a'+c')/(a'+b') = 2.36\pm0.67,$$
$$c' = -(5.25\pm0.39)\times10^{-6} , \tag{5.7a}$$
$$\Delta_{\bar{K}^*\pi} = (52^{+30}_{-52})^\circ ;$$

$$(a'+c')/(a'+b') = -1.70\pm0.67,$$
$$c' = -(0.53\pm0.39)\times10^{-6} , \tag{5.7b}$$
$$\Delta_{\bar{K}^*\pi} = 180^\circ-(52^{+30}_{-52})^\circ .$$

We note that the errors on $\Delta_{\bar{K}^*\pi}$ are so large that the data can be accommodated by real amplitudes without final-state interactions. The amplitude c' is quite different from $\tilde{c}'$ in Eq. (5.3); this might be attributed to rescattering effects as we are going to discuss in Sec. VI.

We next proceed to determine the unprimed amplitudes from $D\to\rho\bar{K}$ and $D^0\to\omega\bar{K}$ decays. From Table I it follows that $D^+\to\rho^+\bar{K}^0, D^0\to\omega\bar{K}^0$ determine

$$|a+b| = (2.18\pm0.25)\times10^{-6} ,$$
$$|b+c| = (2.57\pm0.51)\times10^{-6} . \tag{5.8}$$

From the measurements of $D^0\to\rho^0\bar{K}^0,\rho^+K^-$, we find two solutions:

$$(a+c)/(a+b) = 1.55\pm0.16 ,$$
$$\Delta_{\rho\bar{K}} = (24^{+25}_{-24})^\circ ; \tag{5.9a}$$

$$(a+c)/(a+b) = -0.89\pm0.16 ,$$
$$\Delta_{\rho\bar{K}} = 180^\circ-(24^{+25}_{-24})^\circ . \tag{5.9b}$$

These give the following three possible solutions for amplitudes a,b,c:

$$a = (4.06\pm0.38)\times10^{-6},\quad b = -(1.88\pm0.28)\times10^{-6} ,$$
$$c = -(0.68\pm0.28)\times10^{-6},\quad \Delta_{\rho\bar{K}} = (24^{+25}_{-24})^\circ ; \tag{5.10a}$$

$$a = (1.50\pm0.38)\times10^{-6},\quad b = (0.68\pm0.28)\times10^{-6} ,$$
$$c = (1.89\pm0.28)\times10^{-6},\quad \Delta_{\rho\bar{K}} = (24^{+25}_{-24})^\circ ; \tag{5.10b}$$

$$a = (1.41\pm0.38)\times10^{-6},\quad b = (0.77\pm0.28)\times10^{-6} ,$$
$$c = -(3.34\pm0.28)\times10^{-6},\quad \Delta_{\rho\bar{K}} = 180^\circ-(24^{+25}_{-24})^\circ . \tag{5.10c}$$

Again, because of the large errors, the data are compatible with real amplitudes.

Since there is a severe destructive interference in the mode $D^+\to\bar{K}^{0*}\pi^+$, as indicated by Eq. (5.6), the first solution, Eq. (5.10a), for the unprimed amplitudes is preferred on the theoretical prejudice that it yields a destructive interference in $D^+\to\rho^+\bar{K}^0$. Also, the theoretical calculation in Sec. VI picks up the same solution. Of course, we need other measurements, as discussed shortly, to confirm this conjecture. If (5.10a) turns out to be the correct solution, then this will imply that (5.7b) is the solution for the primed amplitudes since the W-exchange diagram in $D\to\bar{K}^*\pi$ and $\bar{K}\rho$ decays is similar and there is no reason that W-exchange dominates in $D\to\bar{K}^*\pi$ but not so in $\bar{K}\rho$.

One nice prediction from this analysis is

$$B(D^0\to\phi\pi^0) = \tfrac{1}{2}B(D^+\to\phi\pi^+)\Gamma(D^+)/\Gamma(D^0) = 0.21\% . \tag{5.11}$$

This will be an important measurement if we are to test this scheme.

The measurement of $D^+\to\bar{K}^{*0}K^+$ gives us

$$|a'-\tilde{d}+\delta e| = (2.70\pm0.61)\times10^{-6} . \tag{5.12}$$

Future measurements of $\bar{K}^{*0}\eta_8$ $[\propto(b'+c'-2\tilde{c})]$ and $\bar{K}^{*0}\eta_0$ $[\propto(b'+c'+\tilde{c})]$ can help to determine amplitudes $\tilde{c}$ and $(b'+c')$, and then c', since b' is known. Useful information on the amplitude c can be obtained from the data of $D^0 \to \bar{K}^{*0}K^0, K^{*0}\bar{K}^0$ $[\propto(c-c')]$. Then we can check which solution of Eqs. (5.10) will be picked, and thus determine a and b individually. The measurements of $F^+ \to \rho^+\pi^0$ $[\propto(d-d')]$, $F^+ \to \omega\pi^+$ $[\propto(d+d')]$ can help the determination of d and d'. All the amplitudes a,b,c,d, and a',b',c',d', as well as their relative signs thus can be determined to a certain extent.

B. The decays $P_c \to PP$

Next we discuss the case of charm-meson decay into two pseudoscalars, $P_c \to PP$. (Note that the amplitudes a to f here for PP decays have no relation to those for the PV decays. When needed for clarity, we use subscript PP to denote the distinction.) Here, the data are of greater accuracy than for the $P_c \to PV$ case (see Table II). For the case without final-state interactions, one obtains from Table II

$$\frac{\Gamma(D^0 \to K^-\pi^+)}{\Gamma(D^+ \to \bar{K}^0\pi^+)} = \left|\frac{a+c}{a+b}\right|^2 , \tag{5.13a}$$

$$\frac{\Gamma(D^0 \to \bar{K}^0\pi^0)}{\Gamma(D^+ \to \bar{K}^0\pi^+)} = \left|1-\frac{a+c}{a+b}\right|^2 . \tag{5.13b}$$

If all the weak decay amplitudes a,b,c are real, it is easy to check that Eq. (5.13) cannot be fitted by the Mark III data in Table II. We need

$$\frac{a+c}{a+b} \simeq (1.69\pm0.28)e^{i40^\circ} \tag{5.14}$$

to fit the data. As discussed in Sec. IV, such a phase can come from the absorptive part in the weak decay amplitudes which is required by the unitarity of the S matrix once the strong interaction is turned on. Thereby, in order to have a complete agreement with experiment it is mandatory to include the effect of final-state interactions. From Table II(b) we see that $D^0 \to \bar{K}^0\pi^0$ receives a contribution proportional to $(a+b)$ in the presence of strong interactions. Physically, this means that $K^-\pi^+$ mode can be converted into $\bar{K}^0\pi^0$ through the final-state interaction. As a result, the color suppression of the channel $D^0 \to \bar{K}^0\pi^0$ is partially relieved.

From $D^+ \to \pi^+\bar{K}^0$, we obtain

$$|a+b|_{\bar{K}\pi} = (1.66\pm0.11)\times10^{-6}\ \text{GeV} . \tag{5.15}$$

Then from $D^0 \to K^-\pi^+, \bar{K}^0\pi^0$, we obtain the following two solutions:

$$[(a+c)/(a+b)]_{\bar{K}\pi} = 1.95\pm0.14 , \quad \Delta_{\bar{K}\pi} = (79^{+10}_{-14})^\circ , \tag{5.16a}$$

or

$$[(a+c)/(a+b)]_{\bar{K}\pi} = -1.28\pm0.14 , \quad \Delta_{\bar{K}\pi} = 180^\circ-(79^{+10}_{-14})^\circ , \tag{5.16b}$$

where we have assumed no inelasticity for both $I=\frac{1}{2}$ and $\frac{3}{2}$ $\bar{K}\pi$ scattering since we do not know how sizable the inelastic effect is. The model calculation in Sec. VI favors the first solution Eq. (5.16a). We want to caution about the interpretation of the phase-shift difference $\Delta_{\bar{K}\pi} \equiv \delta_{1/2}^{\bar{K}\pi} - \delta_{3/2}^{\bar{K}\pi}$ obtained here from charm decays. The final-state inelastic scattering can convert the state $(\bar{K}\pi)_{I=1/2}$ into the states $\bar{K}\eta$, $\bar{K}\eta'$, $(\bar{K}\pi\pi)_{I=1/2}, \ldots$. As already emphasized in Sec. IV, such inelastic effect would modify the phase shifts of the amplitudes. Hence, the phase difference $\Delta_{\bar{K}\pi}$ in the weak amplitudes is not necessarily the same as the one appearing in strong interactions.

Unlike the PV decays, it is much harder here to determine individual amplitudes, since none of the decays is given by a single amplitude. It is interesting to point out that the nonspectator-diagram amplitudes $\tilde{c}$ and d can be measured in a model-independent way by observing the following decay modes:

$$\frac{\Gamma(D^0 \to \bar{K}^0\eta_8)}{\Gamma(D^0 \to \bar{K}^0\eta)} = \frac{1}{2}\left|\frac{b+c-2\tilde{c}}{b+c+\tilde{c}}\right|^2_{PP} , \quad \frac{\Gamma(F^+ \to \eta_8\pi^+)}{\Gamma(F^+ \to \eta^+\pi)} = 2\left|\frac{a-d}{a+2d}\right|^2_{PP} . \tag{5.17}$$

We want to caution that because of SU(3) breaking (i.e., the large mass difference between π and η,η'), the amplitude c in the $\bar{K}\pi$ state may be substantially different from the same amplitude in the mode $\bar{K}^0\eta$ or $\bar{K}^0\eta'$. Consider the η-η' mixing and the phase-space correction, the current measurement of $D^0 \to \bar{K}^0\eta$ gives (i) $|1.23b-0.49c|_{\bar{K}\eta} = (4.57\pm1.01)\times10^{-6}$ GeV, if there is no SU(3) breaking, i.e., $\tilde{c}=c$, or (ii) $|b+c|_{\bar{K}\eta} = (3.71\pm0.83)\times10^{-6}$ GeV, if SU(3) breaking is maximal, i.e., $\tilde{c}=0$.

We next turn to Cabibbo-singly-suppressed channels. The measurement of $D^+ \to \bar{K}^0K^+$ yields

$$|a-\tilde{d}+\delta e|_{PP} = (4.29\pm0.66)\times10^{-6}\ \text{GeV} . \tag{5.18}$$

Next we discuss the implications of the two quark-mixing-matrix-suppressed decays measured by the Mark III Collaboration:[1]

$$\Gamma(D^0 \to K^-K^+)/\Gamma(D^0 \to K^-\pi^+) = 0.122\pm0.018\pm0.012 , \quad \Gamma(D^0 \to \pi^+\pi^-)/\Gamma(D^0 \to K^-\pi^+) = 0.033\pm0.010\pm0.006 . \tag{5.19}$$

From Table II(a), in the case without SU(3) breaking and final-state interaction one obtains

$$\bar{\Gamma}(D^0 \to K^-K^+)/\bar{\Gamma}(D^0 \to K^-\pi^+) = (s_1/c_1)^2 \approx 0.05 , \quad \bar{\Gamma}(D^0 \to \pi^+\pi^-)/\bar{\Gamma}(D^0 \to K^-\pi^+) = (s_1/c_1)^2 \approx 0.05 . \tag{5.20}$$

(Here $\bar{\Gamma}$ indicates the reduced width, i.e., phase-space factor has been factored out.) Obviously this is not right when comparing to the data Eq. (5.19), a well-known difficulty.[28] From Table II(b), we see that such differences might be attributed to SU(3)-breaking effect of a (i.e., $a_{\bar{K}K} \neq a_{\bar{K}\pi}$, for example), and of $(\delta e+2\delta f)$, which contributes with opposite sign to K^+K^- and $\pi^+\pi^-$, and/or to the final-state-interaction effect. As we shall see in Sec. VI, the theoretical calculation shows that the SU(3) breaking does not suffice to explain the data, Eq. (5.19); in particular, the contribution of $(\delta e+2\delta f)$ makes the situation even worse. To clarify these mechanisms it is thus very important to measure $D^0 \to \pi^0\pi^0$ [see Table II(b)], since the same unknowns $(\delta e + 2\delta f)$, $\delta_0^{\pi\pi}$ are present, but the rest of the amplitudes $(b-c)$ and $(a+b)$ are known [from Eq. (5.16)].

VI. MODEL CALCULATIONS

A. The decays $D \to PP$

Thus far we still do not have a reliable way to calculate the nonleptonic weak decay amplitudes because of the final-state interactions and soft-gluon nonperturbative effects. Indeed, after almost 30 years of the first observation of $\Delta I = \frac{1}{2}$ enhancement in $K \to 2\pi$, we have not yet a reasonable framework to compute the $\Delta I = \frac{1}{2}$ amplitude. In this section we first review the standard approach of calculating charm-decay amplitudes[22,26] and then we proceed to discuss the various models which have been proposed to modify the naive scheme.

The standard approach is based on the valence-quark assumption and vacuum-insertion approximation in which the matrix elements of two quark bilinear operators are saturated by the vacuum intermediate states in all possible ways. As an example we consider the $D \to \bar{K}\pi$ decays. The relevant Hamiltonian is

$$H_{\rm eff} = \frac{G_F}{2\sqrt{2}} V_{ud} V_{cs}^* (c_- O_- + c_+ O_+) ,\quad O_\pm = (\bar{u}d)(\bar{s}c) \pm (\bar{u}c)(\bar{s}d) , \tag{6.1}$$

where

$$(\bar{q}_1 q_2) = \bar{q}_1 \gamma_\mu (1-\gamma_5) q_2$$

and $c_- = (c_+)^{-2} = 1.80$–2.10 for $\Lambda_{\rm QCD} = 200$–500 MeV.

Using the vacuum-insertion method, one obtains

$$\begin{aligned} a &= \frac{G_F}{\sqrt{2}} \frac{2c_+ + c_-}{3} \langle \pi^+ | A_\mu | 0 \rangle \langle K^- | V^\mu | D^0 \rangle , \\ b &= \frac{G_F}{\sqrt{2}} \frac{2c_+ - c_-}{3} \langle \bar{K}^0 | A_\mu | 0 \rangle \langle \pi^+ | V^\mu | D^+ \rangle , \\ c &= \frac{-G_F}{\sqrt{2}} \frac{2c_+ - c_-}{3} \langle K^- \pi^+ | V_\mu | 0 \rangle \langle 0 | A^\mu | D^0 \rangle . \end{aligned} \tag{6.2}$$

Parametrizing the matrix elements as

$$\begin{aligned} \langle P(q) | A_\mu | 0 \rangle &= i f_P q_\mu , \\ \langle P^i(p_2) | V_\mu^j | P^k(P_1) \rangle &= i f^{ijk} [f_+ (p_1 + p_2)_\mu + f_- (p_1 - p_2)_\mu] , \end{aligned} \tag{6.3}$$

it follows that

$$\begin{aligned} a &= \frac{G_F}{\sqrt{2}} \frac{2c_+ + c_-}{3} f_\pi [f_+^{D^0K^-} (m_D{}^2 - m_K{}^2) + f_-^{D^0K^-} m_\pi{}^2] , \\ b &= \frac{G_F}{\sqrt{2}} \frac{2c_+ - c_-}{3} f_K [f_+^{D^+\pi^+} (m_D{}^2 - m_\pi{}^2) + f_-^{D^+\pi^+} m_K{}^2] , \\ c &= \frac{-G_F}{\sqrt{2}} \frac{2c_+ - c_-}{3} f_D [f_+^{K^-\pi^-} (m_K{}^2 - m_\pi{}^2) - f_-^{K^-\pi^-} m_D{}^2]. \end{aligned} \tag{6.4}$$

Form factors at $q^2 = 0$ can be evaluated in the Isgur model:[29]

$$\begin{aligned} & f_+^{D^0K^-}(0) = 1.15, \quad f_-^{D^0K^-}(0) = -0.60 , \\ & f_+^{D^+\pi^+}(0) = 1.33, \quad f_-^{D^+\pi^+}(0) = -0.94 , \\ & f_+^{K^-\pi^-}(0) = 1.03, \quad f_-^{K^-\pi^-}(0) = 0.13 . \end{aligned} \tag{6.5}$$

For the form factor $f_-^{K\pi}$ we take 0.13 as determined from the effective chiral Lagrangian[30] rather than -0.30 as calculated in Ref. 29. The negative sign for $f_-^{K\pi}$ is inconsistent with experiment.[31] Assuming the dipole form for the q^2 dependence of the form factors

$$f_+^{12}(q_{12}{}^2) = f_+^{12}(0) / [1 - (q_{12}{}^2 / m_{12}{}^2)] , \tag{6.6}$$

where $q_{12}{}^2$ is the invariant momentum squared of particles 1,2, and m_{12} is the mass of the contributing vector bosons, we obtain

$$\begin{aligned} & a = 4.4 \times 10^{-6}\ \text{GeV}, \quad b = -1.1 \times 10^{-6}\ \text{GeV} , \\ & c = -6.3 \times 10^{-8}\ \text{GeV} . \end{aligned} \tag{6.7}$$

Therefore, $(a+b) = 3.3 \times 10^{-6}$ GeV, which is to be compared with Eq. (5.15); the decay rate of $D^+ \to \bar{K}^0\pi^+$ is overestimated[32] by a factor of 4. On the other hand, the branching ratio of $D^0 \to \bar{K}^0\pi^0$ is underestimated by a factor of 5 or 6 in the absence of final-state interactions. From Eq. (6.7) it turns out

$$(b-c)/(a+b) = -0.29 \tag{6.8}$$

which is consistent with neither of solutions (5.16).

To improve the discrepancy between Eqs. (6.8) and (5.16), one may enhance the W-exchange amplitude c and/or the internal W-emission diagram (b). Historically, two extreme attempts have been made to improve the situation. In the first approach[33,34] the W-exchange diagram is enhanced by the perturbative gluon which is emitted from the initial light quark and creates the final quark-antiquark pair, or the diagram is enhanced by other mechanisms. An estimate in Ref. 33 gives

$$\frac{c}{a} = 0.8\text{–}1.0 , \tag{6.9}$$

where the amplitudes a and b stay the same as in Eq. (6.7). The result $(b-c)/(a+b) = -(1.37$–$1.63)$ is in agreement with the experimental result, Eq. (5.16a) if the corresponding final-state interaction is taken into account.

The second approach is to enhance the amplitude b, but keep c as negligible as before. Different mechanisms have been suggested in the literature,[35,36,15] for instance, the soft-gluon nonperturbative effects, to give

$$\frac{b}{a} = -0.57, \quad \frac{c}{a} = 0 \tag{6.10}$$

from Eq. (5.16a).

Although both models (6.9) and (6.10) are indistinguishable in $D \to \bar{K}\pi$ decays, they can be tested in other channels.[8]

(a) From the measurements of $D^+ \to \bar{K}^0 K^+$ and $\bar{K}^0\pi^+$ we obtain

$$R_1=\left|\frac{a-d+e}{a+b}\right|^2=6.7\pm1.8 . \tag{6.11}$$

Neglecting the penguin e for the moment, we find that R_1 is predicted to be 3.0 and 5.7, respectively, in models (i) and (ii). Following Ref. 37 we find that the penguin contribution is

$$e=-\frac{4}{9}\alpha_s\left[\frac{m_s^2-m_d^2}{m_c^2}\right]\left[1+2\frac{m_K^2}{(m_u+m_s)(m_c-m_s)}\right]a , \tag{6.12}$$

where use has been made of the fact that $V_{ub}V_{cb}^*$ is very small and α_s is the strong-interaction coupling constant. From (6.12) it is obvious that the penguin amplitude in charm decays is a manifestation of the SU(3)-breaking effect. Hence the penguin contribution to $D^+\to\bar{K}^0K^+$ is destructive and small.

(b) There are two different kinds of W-exchange diagrams in $D^0\to\bar{K}^0\eta$ decays; one is attributed from down-quark-antiquark pair production, the other comes from strange-quark pair creation. Introducing a suppression factor

$$\tilde{c}=\lambda c \tag{6.13}$$

and including η-η' mixing and phase-space corrections, we find

$$R_2=\frac{\Gamma(D^0\to K^0\eta)}{\Gamma(D^0\to\bar{K}^0\pi^0)}=\begin{cases}\text{model (i): } 0.12(\lambda=0);\ 0.03(\lambda=\tfrac{1}{3}) ,\\ \text{model (ii): } 0.4 .\end{cases} \tag{6.14}$$

Hence the W-exchange-enhanced model predicts a very small R_2 for a reasonable value of the SU(3)-breaking effect. Again, the experimental ratio $R_2=0.8\pm0.4\pm0.2$ disfavors the W-exchange-enhanced model. When we discuss $D\to\bar{K}^*\pi$ decays later, it becomes more clear that just enhancing the W-exchange does not suffice to consistently fit the data.

We next discuss the Cabibbo-suppressed decays $D^0\to K^+K^-$ and $\pi^+\pi^-$. It is well known that the SU(3) symmetry implies

$$R_3=\frac{\bar{\Gamma}(D^0\to K^+K^-)}{\bar{\Gamma}(D^0\to K^-\pi^+)}=\tan^2\theta_C ,\quad R_4=\frac{\bar{\Gamma}(D^0\to\pi^+\pi^-)}{\bar{\Gamma}(D^0\to K^-\pi^+)}=\tan^2\theta_C \tag{6.15}$$

(here $\bar{\Gamma}$ indicates the reduced width, i.e., the phase-space factor has been factored out), while experimentally[1]

$$R_3=(2.25\pm0.42)\tan^2\theta_C ,\quad R_4=(0.61\pm0.19)\tan^2\theta_C . \tag{6.16}$$

From Table II we can identify the following two sources responsible for the deviation from the SU(3) predictions, Eq. (6.15).

(a) SU(3) breaking. Because of $f_K>f_\pi$, it follows that $a_{K\bar{K}}>a_{\bar{K}\pi}=a_{\pi\pi}$ for the external W-emission amplitudes. Realistic calculation shows[38]

$$R_3\simeq1.5\tan^2\theta_C,\quad R_4\simeq\tan^2\theta_C . \tag{6.17}$$

Thus, the SU(3)-breaking effect in the spectator quark amplitudes is not sufficient to explain the data (6.16). It has been suggested that the penguin contributions may help to explain the discrepancy between theory and experiment.[39] However, it is evident from Eq. (6.12) that the inclusion of the penguin amplitude would reduce R_3 and enhance R_4. Furthermore, the penguin contribution is rather small, estimated to be $e/a=1\%$.

(b) Final-state interactions. The $\pi\pi$ mode can be converted into $\bar{K}K$ through the inelastic final-state interaction. In particular, there are two known 0^+ resonances $S(975)$ and $\delta(1300)$ which couple to both $\bar{K}K$ and $\pi\pi$ channels. A coupled-channel analysis for this final-state interaction has been performed by Kamal and Cooper.[13] Since the $\pi\pi$ scattering is highly inelastic and $\bar{K}K$ receives contributions from the final-state inelastic scattering, consequently, R_3 is enhanced and R_4 is suppressed. Unfortunately, a quantitative prediction cannot be made at present.

B. The decays $P_c\to VP$

As in the two-pseudoscalar case, we evaluate various quark-diagram amplitudes using the factorization with vacuum-insertion approximation

$$\begin{aligned}a',b&=\frac{G_F}{\sqrt{2}}\eta_{a,b}\langle P\mid A_\mu\mid0\rangle\langle V\mid A^\mu\mid P_c\rangle ,\\ a,b'&=\frac{G_F}{\sqrt{2}}\eta_{a,b}\langle V\mid V_\mu\mid0\rangle\langle P\mid V^\mu\mid P_c\rangle ,\\ c',d'=c,d&=\frac{G_F}{\sqrt{2}}\eta_{c,d}\langle PV\mid A_\mu\mid0\rangle\langle0\mid A^\mu\mid P_c\rangle ,\end{aligned} \tag{6.18}$$

where $\eta_a=\eta_d=(2c_++c_-)/3$, $\eta_b=\eta_c=(2c_+-c_-)/3$, V denotes a vector meson with 0 helicity. To evaluate the amplitudes we also need the matrix elements

$$\begin{aligned}&\langle V(q)\mid V_\mu\mid0\rangle=f_Vm_V^2\epsilon_\mu ,\\ &\langle V^i(q_2)\mid A_\mu^j\mid P^k(q_1)\rangle=-f^{ijk}\epsilon^\nu(F_1g_{\mu\nu}+F_2q_{1\mu}q_{1\nu}+F_3q_{2\mu}q_{1\nu}) ,\end{aligned} \tag{6.19}$$

where ϵ_μ is the polarization vector of the vector meson V. The coupling constant f_V can be inferred from the width of the leptonic decay $V\to e^+e^-$. From $\Gamma(\rho^0\to e^+e^-)=7.084$ keV (Ref. 31), we find $f_{\rho^0}=0.203$. SU(3) and isospin symmetry then give $f_{K^*}=f_{\rho^+}=\sqrt{2}f_{\rho0}=0.287$. For f_ϕ, we use the relation $\langle\phi\mid V_\mu\mid0\rangle=\langle\phi\mid\bar{s}\gamma_\mu s\mid0\rangle=-3\langle\phi\mid J_\mu^{\text{em}}\mid0\rangle$, so $\mid f_\phi\mid=0.228$ from $\Gamma(\phi\to e^+e^-)=1.31$ keV. For form factors we follow the ansatz given by Bernabeu and Jarlskog[40] from a constituent-quark model:

$$\begin{aligned}&F_-(q^2)=-\frac{m_1-m_2}{m_1+m_2}F_+(q^2),\quad F_+(0)=1 ,\\ &F_2(q^2)=0,\quad F_3(q^2)=-\frac{2}{(m_1+m_2)^2-q^2}F_1(q^2) ,\\ &F_1(0)=m_1+m_2 .\end{aligned} \tag{6.20}$$

These form factors are similar to the earlier one[41] except that the q^2 dependence of F_3/F_1 is kept. There are many other different parametrizations for the form factors F_1 and F_3 (Ref. 41). Fortunately, the spectator quark-diagram amplitudes are insensitive to the choice of the form factors.

To illustrate the calculation, let us take $D^+ \to \bar{K}^{0*}\pi^+ (\propto a'+b')$ as an example. The quark-diagram amplitudes are given by

$$a' = \frac{G_F}{\sqrt{2}} \frac{2c_+ + c_-}{3} \langle \pi^+ | (\bar{u}d) | 0 \rangle \langle \bar{K}^{0*} | (\bar{s}c) | D^+ \rangle ,$$
$$b' = \frac{G_F}{\sqrt{2}} \frac{2c_+ - c_-}{3} \langle \bar{K}^{0*} | (\bar{s}d) | 0 \rangle \langle \pi^+ | (\bar{u}c) | D^+ \rangle . \tag{6.21}$$

Applying Eqs. (6.19) and (6.20), and factoring out the polarization term $\epsilon \cdot P_D$, we obtain

$$a' = \frac{G_F}{\sqrt{2}} \frac{2c_+ + c_-}{3} 2m_{K^*} f_\pi \times \frac{(m_D + m_{K^*})^2}{(m_D + m_{K^*})^2 - m_\pi^2} F^{D^+ \bar{K}^*}(q^2 = m_\pi^2) ,$$
$$b' = \frac{G_F}{\sqrt{2}} \frac{2c_+ - c_-}{3} 2m_{K^*}^2 f_{K^*} f_+^{D^+\pi^+}(q^2 = m_{K^*}^2) , \tag{6.22}$$

where $F(q^2)$ is the form factor $F_1(q^2)$ normalized to unity at $q^2=0$. Assuming the dipole form for the q^2 dependence of the form factors, we find

$$a' = 2.13 \times 10^{-6}, \quad b' = -0.87 \times 10^{-6} . \tag{6.23}$$

The relative sign of a' and b' is very confusing in the literature. Without taking into account the structure constants f_{ijk} in Eqs. (6.3) and (6.19), as done by some authors, the relative sign will become positive. It has been argued in Ref. 42 that in the limit of V-spin symmetry, the decay $D^+ \to \bar{K}^{0*}\pi^+$ receives contributions only from the O_- operator. As a result, the matrix element $\langle \bar{K}^{0*} \cdots D^+ \rangle$ in b' must be the same as the matrix element $\langle \pi^+ \cdots D^+ \rangle$ in a' [see Eq. (6.21)] except with a sign difference. This means that b' is of the same sign as a'. However, this argument is not correct since $\bar{K}^{0*}$ and π^+ are in different V-spin doublets so they can form a p-wave V singlet as well as a V-spin triplet.[43] The experimental result (5.6) clearly indicates a large destructive interference in $D^+ \to \bar{K}^{0*}\pi^+$ decays.

By comparing Eq. (6.23) with Eq. (5.8), it is obvious that b' is underestimated by a factor of 4. This clearly indicates that the phenomenological model in which only the W exchange is being enhanced is ruled out. The reader may now appreciate the quark-diagram analysis in Sec. V which we just used to rule out the purely W-exchange-enhanced model.

C. *W*-exchange diagram

W exchange in $D \to \bar{K}\pi$ is suppressed due to the smallness of the form factors at large momentum transfer [see Eqs. (6.4) and (6.6)]. More precisely, m_{12} in Eq. (6.6) is the mass of $\kappa(890)$, so W exchange is form-factor suppressed. In terms of quark language, this is equivalent to say that W exchange in $D \to PP$ decays is helicity suppressed. For $P_c \to VP$ decays, the W exchange or annihilation is, however, no longer subject to form-factor suppression since m_{12} is now the mass of 1^+ state of the final two quarks; helicity suppression is no more active. The decay $D^0 \to \bar{K}^0\phi$, which was seen by ARGUS,[3] CLEO,[2] and Mark III[1] Collaborations, ought to be an unambiguous signal for W exchange.[6,44]

However, recently it was argued in Ref. 45 that rescattering effects required by unitarity can produce the reaction $D^0 \to \bar{K}^0\phi$, even when the W-exchange diagram is absent. An example of the rescattering graph is shown in Fig. 3(a). This problem was also examined within the framework of $1/N_c$ approach (N_c being the number of colors) by Baur, Buras, and Gerard.[46] They concluded that the leading contribution from the W-exchange diagram does not suffice to explain the experimental value of $B(D^0 \to \bar{K}^0\phi)$. Hence, the question is whether or not the W-exchange mechanism is the dominant contribution to $D^0 \to \bar{K}^0\phi$.

Here we wish to emphasize that in the quark-diagram scheme employed in this paper, all quark diagrams include strong-interaction gluon clouds. In other words, all quark graphs used in this approach are symbolic and meant to have all the strong interactions included, i.e., gluon lines are included in all possible ways;[9] hence, they are *not* Feynman graphs. The rescattering diagram, Fig. 3(a), is actually a W-exchange graph, which can be easily seen if we redraw it as Fig. 3(b). Therefore, $D^0 \to \bar{K}^0\phi$ is really dominated by the generic W-exchange graph. Of course, if one wishes to calculate the decay amplitude of $D \to \bar{K}\phi$ dynamically in terms of Feynman graphs, one would encounter the final-state rescattering effect and the nonleading nonfactorizable $1/N_c$ diagrams. Whether or not these contributions are important is another subject.

VII. LIFETIME DIFFERENCE BETWEEN CHARMED MESONS

This section is devoted to the understanding of the lifetime difference between D^+, D^0, and F^+ (Ref. 47). The large number of D branching ratios measured by the Mark III can now account for about 85% of the D^0 and D^+ decays.[1] Therefore, the difference between D^0 and

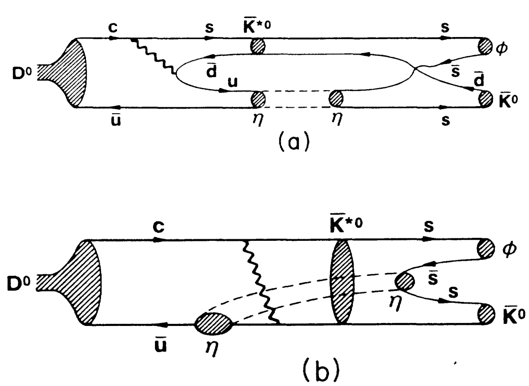

FIG. 3. (a) The final-state rescattering graph suggested in Ref. 45 as an important contribution to $D^0 \to \bar{K}^0\phi$. (b) Redrawing of (a).

D^+ lifetimes should be understandable at the level of two-body decays since three- and four-body decays are dominated by quasi-two-body channels.[48] The longer lifetime of D^+ compared to D^0 may be attributed to the suppression of D^+ rate and the enhancement of D^0 rate.[8]

(i) D^+ suppression. There are two reasons for the suppression of D^+ rate. (a) Owing to identical-particle effects in the Cabibbo-allowed decays of D^+, only a few channels are open to D^+. More precisely, there are 16 channels for Cabibbo-favored D^0 decays, but only 4 channels for D^+. (b) Neglecting the W-exchange diagram for the moment, then D^0 decay amplitude is proportional to a or b, whereas the Cabibbo-allowed D^+ amplitude is always of the form $(a+b)$ due to the identical-particle effect. Equations (5.6), (5.10a), and (6.10) imply that the interference between a and b is always destructive. Hence there is a severe destructive interference in all Cabibbo-favored D^+ decays.[49]

(ii) D^0 enhancement. The W-exchange diagram, which does not exist in D^+ decays, may enhance the D^0 decay rate. However, from the solutions (5.6), (5.10a), and the discussions in Sec. VI we know that W-exchange cannot be the dominant contribution. Furthermore, the contribution of W-exchange to the decay rates can be either constructive or *destructive;* it is channel dependent. For example, the partial rates of $D^0 \to \rho^+ K^-$ and $\rho^0 \bar{K}^0$ are suppressed in the presence of W-exchange effect. Although a quantitative prediction requires the details of final-state interactions, it is qualitatively clear that the enhancement of D^0 rate due to the presence of W exchange is not important.

From the above discussions, we see that the D^0 and D^+ lifetime difference is attributed mainly to the large suppression of the D^+ rate, namely, (a) fewer available channels and (b) large destructive interference. It should be emphasized that the difference of lifetimes of D^0 and D^+ cannot be solely due to the nonspectator contribution as proposed in models with the soft-gluon emission or with the presence of gluons in the initial D state.[50] First of all, this purely W-exchange enhanced model is ruled out when we discuss $P_c \to VP$ and PP decays. Second, the effect of D^+-rate suppression is model independent and has to be included.

Recently, there are several new measurements of the F^+ lifetime: $(3.2^{+3.0}_{-1.3})\times 10^{-13}$ sec by the NA11 experiment (Ref. 51), $(3.5^{+2.4}_{-1.8}\pm 0.9)\times 10^{-13}$ sec by the HRS group,[52] and $(2.6^{+1.6}_{-0.9})\times 10^{-13}$ sec by the new E531 experiment.[53] All the new measurements are substantially larger than the old E531 result (Ref. 54), $(1.9^{+1.3}_{-0.7})\times 10^{-13}$ sec, but still smaller than the world-average D^0 lifetime, Eq. (5.4).

The study of the lifetime of F^+ is of special interest since it can help to sort out the possible explanations for the D^0-D^+ lifetime difference. Indeed, if W exchange is the sole mechanism responsible for the large difference in the lifetimes of D^0 and D^+, F^+ would be expected to have a larger lifetime than D^0, which is not seen experimentally. As far as the principal spectator diagrams are concerned, F^+ decay is very similar to D^0. Hence it is plausible that the F^+-D^0 lifetime difference may be attributed to the nonspectator contributions. As the W-exchange diagram in D^0 decays, the W-annihilation diagram in $F^+ \to VP$ and VV decays is not subject to helicity or form-factor suppression. Moreover, W annihilation is not color suppressed and this is the crucial difference between W exchange and W annihilation. As a result, it is very likely that the difference of D^0 and F^+ lifetimes comes from the W-annihilation contributions in the VP and VV decays of the F^+ meson. In view of this, the future measurements of $F^+ \to \omega\pi^+$, $\rho^+\pi^0$, and $\rho\pi^+$, which proceed through the W annihilation, are of paramount importance to test the role of W annihilation in F^+ decays.

In most model calculations, $F^+ \to \pi^+\pi^0$ is forbidden by the isospin argument, and the amplitudes of $F^+ \to \rho^+\pi^0, \rho^0\pi^+$ vanish since d and d' are the same [see Table I(c)]. Consequently, it is difficult to understand why two-body decays of F^+ have larger decay rates compared with that of D^0 (Ref. 55). The experimental fact that $\tau(F^+) < \tau(D^0)$ may imply that the two W-annihilation graphs d and d' should not be the same when the spin of the final state is taken into consideration. It is thus quite important to measure $F^+ \to \rho^+\pi^0$ and $\rho^0\pi^+$ to see the role of W annihilation.

VIII. DISCUSSION AND CONCLUSIONS

In this paper we have analyzed exclusive two-body decays of charmed mesons within the framework of the quark-diagram scheme. We have shown that the Mark III data of charm decaying into two pseudoscalar $P_c \to PP$ cannot be fitted if all quark-diagram amplitudes are real and if SU(3) symmetry is assumed. We therefore incorporate the effects of SU(3) breaking and final-state interactions in our quark diagram to account for *inelastic* final-state interactions due to resonances and rescattering effects, we introduce an explicit factor $e^{i\delta}$ for each partial-wave isospin amplitudes. This phase shift δ has both real and imaginary parts; the imaginary component reflects the inelastic effect. Because of inelasticities and the communicating multichannels available in the decays it is not straightforward to compare the phase shifts in hadronic scattering to those determined in charm decays.

Owing to the purity of the quark contents in ϕ and ω, many PV decays are governed by one type of amplitude, as shown in Table I. This enables us to determine, for the first time, some of the quark diagrams from the data. Solutions (5.6), (5.7), (5.9), and (5.10) are our results. We also pointed out future measurements of $D^0 \to \bar{K}^{*0}\eta_8$, $\bar{K}^{*0}\eta_0$, $\bar{K}^{*0}K^0$, $K^{*0}\bar{K}^0$, and $F^+ \to \rho^+\pi^0$, $\omega\pi^+$ will give definite and model-independent results about individual amplitudes.

For $P_c \to PP$ decays, we obtain two solutions, Eq. (5.16). It is evident from these solutions that a nontrivial phase shift difference due to final-state interactions is indispensable to fit the data. Unlike the PV decays, it is much harder here to determine individual amplitudes, since none of the decays is given by a single amplitude. Nevertheless, it is pointed out that the future measurements of $D^0 \to \bar{K}^0\eta_8$, $\bar{K}^0\eta_0$, $F^+ \to \pi^+\eta_8$, $\pi^+\eta_0$ will help to sort out individual amplitudes. We also urge experimentalists to measure $D^0 \to \pi^0\pi^0$ to clarify which mechanism, SU(3) breaking and/or final-state interactions, is responsible for the ratio $\Gamma(D^0 \to K^+K^-)/\Gamma(D^0 \to \pi^+\pi^-) \approx 3.7$.

We then proceed to give theoretical model calculations of various quark-diagram amplitudes. It has been established that the vacuum-insertion calculations do not agree with the two-body exclusive decays. For example, the decay rate of $D^+ \to \bar{K}^0\pi^+$ is overestimated by a factor of 4, whereas the branching ratio of $D^0 \to \bar{K}^0\pi^0$ is underestimated by a factor of 5 or 6 in the absence of final-state interactions. Most perturbative QCD calculations give negligible W-exchange and W-annihilation amplitudes for $P_c \to PP$ decays. To improve the comparison with the data attempts have been made to enhance either the W-exchange diagram (c), or the internal W-emission diagram (b). From the Mark III data of $D^+ \to \bar{K}^0K^+$ and $D^0 \to \bar{K}^0\eta$, it appears that just enhancing W exchange does not suffice to consistently fit the data. Our model calculations reveal that SU(3) breaking alone is not sufficient to explain the observed ratios $\Gamma(D^0 \to K^+K^-)/\Gamma(D^0 \to K^-\pi^+)$ and $\Gamma(D^0 \to \pi^+\pi^-)/\Gamma(D^0 \to K^-\pi^+)$. Therefore, the two 0^+ resonances $S(975)$ and $\delta(1300)$ which couple to both $\bar{K}K$ and $\pi\pi$ channels may play an important role in this case. As mentioned above, the measurement of $D^0 \to \pi^0\pi^0$ is needed to clarify the situation.

For $P_c \to VP$ decays, we first clarify the sign confusion in the literature about amplitudes a' and b', a and b; the V-spin symmetry argument does not tell us what is the relative sign it should be. We then point out that the theoretical prediction of the internal W-emission amplitude b' is too small by a factor of 4. (Recall that b' can be uniquely determined from $D^+ \to \phi\pi^+$ and $D \to \bar{K}^*\pi$ data.) This clearly indicates that the phenomenological model in which only the W exchange is being enhanced is ruled out.

We then turn to the W-exchange diagram. Even in the naive model, W exchange or annihilation in $P_c \to VP$ decays is no longer subject to form-factor and helicity suppression. Indeed, the decay $D^0 \to \bar{K}^0\phi$, which proceeds through the W-exchange mechanism, is seen with large branching ratio $\sim 1\%$. It has been suggested that a substantial fraction of $D^0 \to \bar{K}^0\phi$ could come from the rescattering effects and nonleading $1/N_c$ contributions. We clarify that actually all these contributions are included in the generic W-exchange graph (which is not a Feynman graph).

Finally, we discuss the lifetime differences between D^+, D^0, and F^+. We argue that the difference in lifetimes should be understandable at the level of two-body decays since three- and four-body decays are dominated by quasi-two-body channels. We conclude that the difference of lifetimes of D^0 and D^+ is attributed mainly to the large suppression of the D^+ rate: namely, (i) fewer available channels and (ii) large destructive interference. We emphasize that the lifetime difference cannot be *solely* due to soft-gluon emission as proposed in some models. The study of the lifetime of F^+ can help to sort out the possible explanations for the D^0-D^+ lifetime difference. We contemplate that the difference between F^+-D^0 lifetimes is due to the nonspectator contributions. W-annihilation diagram in F^+ decays does not subject color suppression, whereas W exchange does. In view of this, future measurements of $F^+ \to \omega\pi^+$, $\rho^+\pi^0$, and $\rho\pi^+$ are of importance to test the role of W annihilation in F^+ decays, and hence the lifetime differences.

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