No final-state phase ambiguities in additional B_d decays with large CP violation

Isard Dunietz
Stanford Linear Accelerator Center, Stanford University, Stanford, California 94309

Art Snyder
Indiana University, Bloomington, Indiana 47401
(Received 23 April 1990)

The B_d decay modes $\bar{D}^0(\bar{D}^{*0})X^0$ generated by the quark process $\bar{b}\to\bar{c}+u\bar{d}$ have a large CP asymmetry within the Cabibbo-Kobayashi-Maskawa (CKM) model. Here X^0 denotes a light, neutral meson with zero strangeness, such as π^0, ρ^0, ω, etc. This asymmetry depends *only* on a ratio of CKM matrix elements and *not* on final-state phases. The CKM model predicts (up to a sign) the same asymmetries for the $\bar{D}^0(\bar{D}^{*0})X^0, \psi K_S$, and D^+D^- modes. Adding those modes might lower the required luminosity for observing CP violation within the CKM model. On the other hand, future high-statistics measurements could reveal violations of the CKM model simply by demonstrating that the modes ψK_S, D^+D^- and $\bar{D}^0(\bar{D}^{*0})X^0$ differ in their individual CP asymmetries.

I. INTRODUCTION

Recently there has been interest in searching for CP violation with neutral B meson decays. The classic decay mode is $B_d\to\psi K_S$.[1,2] The decay mode $B_d\to\pi^+\pi^-$ has also been studied in some detail.[2]

Here we suggest that decay modes, discussed earlier in the literature,[3,1,4,5] of the type $B_d\to\bar{D}^0X^0$ or $B_d\to\bar{D}^{*0}X^0$, generated by the quark process $\bar{b}\to\bar{c}+u\bar{d}$, may offer experimental sensitivity comparable to ψK_S when the $\bar{D}^0$ or $\bar{D}^{*0}$ particle decays into a CP eigenstate $(f)_{\mathcal{D}}$.[6] As in the ψK_S case, the large amplitude due to the direct decay $[B_d\to\bar{D}^0(\bar{D}^{*0})X^0\to(f)_{\mathcal{D}}X^0]$ interferes with the large amplitude due to B_d-$\bar{B}_d$ mixing $[B_{d,\text{phys}}\to\bar{B}_d\to D^0(D^{*0})X^0\to(f)_{\mathcal{D}}X^0]$ to yield a sizable asymmetry.[7]

Furthermore, the prediction for this asymmetry is theoretically clean. First, akin to the $B_d\to\psi K_S$ mode, Section II shows that uncertainties in hadronic matrix elements and final-state phases do not enter. Within the standard model of the electroweak interactions, in which CP violation arises from the complex Cabibbo-Kobayashi-Maskawa (CKM) matrix,[8] the interference term for $B_d\to(f)_{\mathcal{D}}X^0$ is the same as for $B_d\to\psi K_S$, namely, $|\text{Im}\lambda|=\sin(2\beta)$, and satisfies[9]

$$0.08\lesssim\sin(2\beta)\le 1 \ . \tag{1.1}$$

The angle β is the angle between $V_{tb}^*V_{td}$ and $V_{cb}^*V_{cd}$:

$$\beta=-\arg\left[-\frac{V_{tb}^*V_{td}}{V_{cb}^*V_{cd}}\right] .$$

Second, penguin contributions generally create uncertainties in the predictions of CP asymmetries. Well-known cases are the $B_d\to\pi^+\pi^-$ mode (and to a lesser extent $B_d\to D^+D^-$). In the modes advocated in this paper $[B_d\to\bar{D}^0(\bar{D}^{*0})X^0]$, however, penguin contributions are absent and cause no uncertainties.

To increase the data sample many different modes of the type $(f)_{\mathcal{D}}X^0$ can be summed over.[5] Since the asymmetries of final states with opposite CP parities differ by a minus sign, it will be crucial to distinguish between final states with different CP parities, so as not to dilute the signal.[10] Section III lists the relevant decay modes of the form $(f)_{\mathcal{D}}X^0$.

Whether the decays $\bar{B}_d\to D^0(D^{*0})X^0\to(f)_{\mathcal{D}}X^0$ are a realistic complement to the ψK_S mode depends on their branching ratios. Section IV estimates their branching ratios within the framework of the Bauer-Stech-Wirbel (BSW) model,[11] but cautions the reader that final-state interactions and other uncalculable effects render such estimates unreliable. Ultimately, measurements of the final states will determine whether those decays can substantially complement the ψK_S mode.

Even if the $(f)_{\mathcal{D}}X^0$ modes are not competitive with the ψK_S one, important information will be obtained by pursuing them. The standard model predicts the same CP-violating interference term $\text{Im}\lambda$ for the $(f)_{\mathcal{D}}X^0$, ψK_S, and D^+D^- modes. If new physics were to occur within the D^0-$\bar{D}^0$ complex, the interference term of the $(f)_{\mathcal{D}}X^0$ modes could differ from the ψK_S and D^+D^- ones,[12] which would be a violation of the standard model. Section V concludes.

II. DISCUSSION

A detailed discussion of the decay $B_d\to\bar{D}^0\pi^0\to(\pi^+\pi^-)_D\pi^0$ clarifies our idea. Unless otherwise indicated, the discussion will hold true for other decay modes, listed in Sec. III. Consider the process $B_d\to\bar{D}^0\pi^0\to(\pi^+\pi^-)_D\pi^0$. Here the CP asymmetry to measure is

$$A_{\text{sym}}\equiv\frac{\Gamma(B_{d,\text{phys}}\to(\pi^+\pi^-)_D\pi^0)-\Gamma(\bar{B}_{d,\text{phys}}\to(\pi^+\pi^-)_D\pi^0)}{\Gamma(B_{d,\text{phys}}\to(\pi^+\pi^-)_D\pi^0)+\Gamma(\bar{B}_{d,\text{phys}}\to(\pi^+\pi^-)_D\pi^0)} . \tag{2.1}$$

Four amplitudes give rise to the decay $B_{d,\mathrm{phys}} \to (\pi^+\pi^-)_D \pi^0$, when D^0-$\bar{D}^0$ mixing is neglected. Those amplitudes are

$$(1)\quad B_d \to \bar{D}^0\pi^0 \to (\pi^+\pi^-)_D\pi^0\ , \tag{2.2}$$

$$(2)\quad B_d \to D^0\pi^0 \to (\pi^+\pi^-)_D\pi^0\ , \tag{2.3}$$

$$(3)\quad B_{d,\mathrm{phys}} \to \bar{B}_d \to D^0\pi^0 \to (\pi^+\pi^-)_D\pi^0\ , \tag{2.4}$$

$$(4)\quad B_{d,\mathrm{phys}} \to \bar{B}_d \to \bar{D}^0\pi^0 \to (\pi^+\pi^-)_D\pi^0\ . \tag{2.5}$$

The amplitudes (2) and (4) are doubly Cabibbo-Kobayashi-Maskawa (CKM) suppressed in relation to (1) and (3), respectively.

A large interference between the main amplitudes (1) and (3) occurs, and the *CP*-violating interference term is given by

$$\mathrm{Im}\lambda = \mp \sin(2\beta)\ . \tag{2.6}$$

The sign of $\mathrm{Im}\lambda$ depends on whether the final state $(f)_{\mathcal{D}}X^0$ is *CP* even or odd. For *CP*-even [-odd] states, $\mathrm{Im}\lambda = -\sin(2\beta)$ [$\mathrm{Im}\lambda = +\sin(2\beta)$].[10] The final state of our example $(\pi^+\pi^-)_D\pi^0$ is *CP* odd.

Equation (2.6) is derived as follows. The pure B_d decay amplitude $B_d \to \bar{D}^0\pi^0$ can be parametrized as

$$A(B_d \to \bar{D}^0\pi^0) = V_{cb}^* V_{ud} |a| e^{i\delta}\ . \tag{2.7}$$

Because only one CKM combination $V_{cb}^* V_{ud}$ contributes to this decay, the hadronic matrix elements and final-state-interaction phases can be represented as a complex number $|a|e^{i\delta}$. Notice that this complex number includes final-state-interaction effects, such as the rescattering $B_d \to D^-\pi^+ \to \bar{D}^0\pi^0$. To be explicit, we note that the decay $B_d \to \bar{D}^0\pi^0$ has two different isospin amplitudes: $I = 1/2$ and $3/2$. The complex number is represented as

$$|a|e^{i\delta} = |a_{1/2}|e^{i\delta_{1/2}} + |a_{3/2}|e^{i\delta_{3/2}}\ .$$

No assumption whatsoever was made about the final-state phases δ_I and the magnitudes $|a_I|$. The crucial point is to observe that the amplitude for the process $B_d \to \bar{D}^0\pi^0$, although involving different isospin amplitudes and phases, involves only the one CKM combination $V_{cb}^* V_{ud}$.

The *CP*-conjugated mode leaves the final-state phases unchanged and complex conjugates the CKM elements:

$$A(\bar{B}_d \to D^0\pi^0) = V_{cb} V_{ud}^* |a| e^{i\delta}\ . \tag{2.8}$$

CP violation depends on the ratio of the two amplitudes:

$$\frac{A(\bar{B}_d \to D^0\pi^0)}{A(B_d \to \bar{D}^0\pi^0)} = \frac{V_{cb} V_{ud}^* |a| e^{i\delta}}{V_{cb}^* V_{ud} |a| e^{i\delta}} = \frac{V_{cb} V_{ud}^*}{V_{cb}^* V_{ud}}\ .$$

We see that the dependence on the final-state interactions—parametrized by $|a|e^{i\delta}$—factors out.

The doubly-CKM-suppressed amplitudes will in general have different final-state phases, but are negligible.[13] *The uncertainty due to final-state phases is removed, because those amplitudes where different final-state phases could occur have tiny CKM elements and are negligible.*

The amplitudes for the subsequent neutral D decays are given by

$$A(D^0 \to \pi^+\pi^-) \sim V_{cd}^* V_{ud}\ , \tag{2.9a}$$

$$A(\bar{D}^0 \to \pi^+\pi^-) \sim V_{cd} V_{ud}^*\ , \tag{2.9b}$$

where a common factor is suppressed (see the Appendix). Thus, the amplitudes for the decay chains are

$$A(B_d \to (\pi^+\pi^-)_D\pi^0) \sim V_{cb}^* V_{ud} V_{cd} V_{ud}^*\ , \tag{2.10a}$$

$$A(\bar{B}_d \to (\pi^+\pi^-)_D\pi^0) \sim -V_{cb} V_{ud}^* V_{cd}^* V_{ud}\ . \tag{2.10b}$$

The minus sign in Eq. (2.10b) arises because the final state is *CP* odd.[10] The interference term $\mathrm{Im}\lambda$ is[1-5]

$$\mathrm{Im}\lambda = \mathrm{Im}\frac{V_{tb}^* V_{td} A(\bar{B}_d \to (\pi^+\pi^-)_D\pi^0)}{V_{tb} V_{td}^* A(B_d \to (\pi^+\pi^-)_D\pi^0)}\ . \tag{2.11}$$

Henceforth we choose to work with the Wolfenstein parametrization,[14] where

$$\mathrm{Im}\lambda = -\mathrm{Im}\frac{V_{td}}{V_{td}^*}\ . \tag{2.12}$$

Since in the Wolfenstein parametrization the four CKM elements

$$V_{ud}, V_{us}, V_{cd}, V_{cs} \tag{2.13}$$

are real to excellent approximations, any $D^0 \to (f)_D$ decay satisfies[15]

$$\frac{A(D^0 \to (f)_D)}{A(\bar{D}^0 \to (f)_D)} \approx \pm 1\ . \tag{2.14}$$

The sign is plus (minus) for a *CP*-even (-odd) final state $(f)_D$.[10] For any hadronic *CP* eigenstate $(f)_D$, the interference term is

$$\mathrm{Im}\lambda = \pm \mathrm{Im}\frac{V_{td}}{V_{td}^*}\ . \tag{2.15}$$

The sign is determined as discussed in the paragraph following Eq. (2.6). Similar arguments hold for $D^0(\text{excited}) \to (f)_{\mathcal{D}}$.

A note about the neglect of D^0-$\bar{D}^0$ mixing is in order. If D^0-$\bar{D}^0$ were present, then for each of the decay chains (2.2)–(2.5) there exists another one, where the time-evolved $\bar{D}^0$ or D^0 mixes into its antiparticle. For instance, in addition to

$$B_d \to \bar{D}^0\pi^0 \to (\pi^+\pi^-)_D\pi^0\ , \tag{2.2}$$

D^0-$\bar{D}^0$ mixing yields

$$B_d \to \bar{D}^0\pi^0 \to D^0\pi^0 \to (\pi^+\pi^-)_D\pi^0\ . \tag{2.16}$$

Present limits on the square of the D^0-$\bar{D}^0$ mixing amplitude are $r_D < 0.37\%$ (90% C.L.),[16] and for our purposes can be neglected.[17]

To conclude the section, we recapitulate on the sign of $\mathrm{Im}\lambda$. Final states $(f)_{\mathcal{D}}X^0$ which are *CP*-even (-odd) eigenstates have an interference term $\mathrm{Im}\lambda = -\sin 2\beta$ ($+\sin 2\beta$).

III. RELEVANT DECAY MODES

We first discuss the decays $B_d \to (f)_{\mathcal{D}}X^0$, when the final state is a *CP* eigenstate. We concentrate on two possibilities.

(1) The B_d decays into the lowest-lying meson with $\bar{c}u$ quantum numbers, the $\bar{D}^0$, via $B_d \to \bar{D}^0 X^0 \to (f)_D X^0$. The meson X^0 can be one of the following:

(a) pseudoscalars ($J^{PC}=0^{-+}$)π^0, η, η' ,

(b) vectors ($J^{PC}=1^{--}$)ρ, ω, ϕ ,

(c) the $J^{PC}=1^{++}$ resonances

$a_1(1260), f_1(1285), f_1(1420)$,

(d) the $J^{PC}=0^{++}$ resonances

$a_0(980), f_0(975), f_0(1400)$,

(e) the $J^{PC}=1^{+-}$ resonances

$b_1(1235), h_1(1170)$,

(f) the $J^{PC}=2^{++}$ resonances

$a_2(1320), f_2'(1525), f_2(1270)$.

The above-mentioned $\bar{D}^0$ decays into a CP eigenstate, $\bar{D}^0 \to (f)_D$. The CP parity of the final state in the process

$$B_d \to \bar{D}^0 X^0 \to (f)_D X^0 , \tag{3.1}$$

is

$$CP[(f)_D X^0] = \begin{cases} -CP[(f)_D] & \text{if } X^0 \text{ belongs to set (a)–(c)} , \\ +CP[(f)_D] & \text{if } X^0 \text{ belongs to set (d)–(f)} . \end{cases}$$

Whereas X^0 from set (a)–(c) contributes with the same sign to the CP asymmetry, the X^0 from set (d)–(f) contributes with the opposite sign, for a given $(f)_D$. Note that all the particles and resonances X^0 below ~ 900 MeV enter with the same sign to the CP asymmetry.

(2) The B_d decays into $\bar{D}^{*0}$ or higher resonances with $\bar{c}u$ quantum numbers. Consider, thus, $B_d \to \bar{D}^0$(excited)$X^0 \to (f)_{\mathcal{D}} X^0$. For a spinless, excited $\bar{D}^0$, we allow the meson X^0 to be anything listed in (a)–(f). For excited $\bar{D}^0$ with higher spin, such as $\bar{D}^{*0}$, the meson X^0 must be spinless to allow a definite CP parity for the final state. Then the meson X^0 must be from (a) or (d) with $J=0$.

Since a definite CP parity is required for the final state of a B_d meson, by necessity the $(f)_{\mathcal{D}}$ mode, arising from $\bar{D}^0$(excited)$\to (f)_{\mathcal{D}}$, must have a definite CP parity also. This report focuses on the $\bar{D}^{*0}$,[18] which has two main decay modes:[19]

$$B(\bar{D}^{*0} \to \bar{D}^0 \pi^0) \approx B(\bar{D}^{*0} \to \bar{D}^0 \gamma) \approx 50\% . \tag{3.2}$$

Since the $\bar{D}^{*0}$ decays via parity-conserving interactions, both the $\bar{D}^0 \pi^0$ and the $\bar{D}^0 \gamma$ are in a P wave. Thus if the subsequent $\bar{D}^0$ decays into a CP eigenstate, the final state of $\bar{D}^{*0} \to (f)_{D^*}$ is a CP eigenstate. However the CP parity of this final state depends on whether the π^0 or the photon is involved, as they have opposite intrinsic CP parity.

Thus, consider the two decay chains, where X^0 is spinless:

(a) $\bar{B}_d \to D^{*0} X^0 \to [\gamma D^0]_{D^*} X^0 \to [\gamma (f)_D]_{D^*} X^0$,

(b) $\bar{B}_d \to D^{*0} X^0 \to [\pi^0 D^0]_{D^*} X^0 \to [\pi^0 (f)_D]_{D^*} X^0$.

We have shown that both final states have definite CP parities, and that the CP parity of the final state in decay chain (a) is opposite to the one in decay chain (b). Explicitly,

(a) $CP|[\gamma(f)_D]_{D^*} X^0\rangle = +\eta\{X^0\}\eta\{(f)_D\} \, |[\gamma(f)_D]_{D^*} X^0\rangle$,

(b) $CP|[\pi^0(f)_D]_{D^*} X^0\rangle = -\eta\{X^0\}\eta\{(f)_D\} \, |[\pi^0(f)_D]_{D^*} X^0\rangle$.

Here η denotes the intrinsic CP parity of the state in brackets. Thus it is crucial to distinguish the mode $D^{*0} \to \pi^0 D^0$ from $D^{*0} \to \gamma D^0$. This is possible if the detector has good calorimetry, such as the CsI detector recently installed by the CLEO Collaboration.[20]

To summarize, two possibilities (1) and (2) for large clean CP violation were discussed. The remainder of this section lists the D^0 modes $(f)_D$ into CP eigenstates.

(D1) Final states which include one neutral kaon, such as $\bar{K}^0\pi^0$, $\bar{K}^0\eta$, $\bar{K}^0\eta'$, $\bar{K}^0\rho$, $\bar{K}^0\omega$, $\bar{K}^0\phi$, $\bar{K}^0$ (c)–(f) (Ref. 21), $\bar{K}^{*0}\pi^0$, $\bar{K}^{*0}\eta$, $\bar{K}^{*0}\eta'$, $\bar{K}^{*0}(d)$.[21]

(D2) Final states which include an $\bar{s}s$ quark pair, such as $\phi\pi^0, \phi\eta, K^+K^-, \bar{K}^0K^0, \bar{K}^{*0}K^0 \, \bar{K}^0K^{*0}$.

(D3) Final states which include flavor-neutral decays, such as $\pi^+\pi^-$, $\{\pi^0, \eta, \eta', a_0, f_0\}$ $\{\pi^0, \eta, \eta', \rho, \omega$, (c)–(f)$\}$.[22,21]

The two vector modes of a D^0 might have a dominant CP parity as discussed in Ref. 23. If the modes in (D4) and (D5) below are experimentally found to be predominant CP eigenstates, then they could also be used for $(f)_D$.

(D4) Final states, such as, $\rho^0\phi, \omega\phi, \omega\bar{K}^{*0}, \rho\bar{K}^{*0}$.

(D5) Final states which are a particle-antiparticle system, such as, $\bar{K}^{*0}K^{*0}, K^{*-}K^{*+}, \rho^+\rho^-, \rho^0\rho^0, \omega\omega, \ldots$.

For cases (D1)–(D4) the neutral kaon K^0 is seen in its $\pi^+\pi^-$ decay, and the K^{*0} in its $K^0\pi^0$ mode. In contrast, case (D5) allows K^{*0} to be seen in its $K^0\pi^0$ and $K^+\pi^-$ modes, and K^{*+} in its $K^0\pi^+$ and $K^+\pi^0$ modes, because the final state $(f)_D$ is a particle-antiparticle system. In

TABLE I. The efficiencies of various final-state particles. Here ϵ_D is the D^0 branching fraction into visible CP eigenstates (see text).

Particle	ϵ, efficiency
$\pi^0,\pi^\pm,K^\pm$	100%
K^0	33%
K^{*0}	11%
ρ^0	100%
η	60%
ω	90%
η'	30%
ϕ	50%
D^0	ϵ_D
D^{*0}	$100\%\times\epsilon_D$

the next section, we estimate the combined B_d branching ratio to final states of the form (1) or (2).

IV. RATE ESTIMATES

We now wish to compare the statistical power of the $B_d\to(f)_{\mathcal{D}}X^0$ modes to that of the classic $B_d\to\psi K_S$ one. This requires estimates of branching ratios and efficiencies. For purposes of the paper efficiency refers to the branching fraction to visible modes. We do not discuss the ability to vertex those modes, which is an important requirement for an asymmetric $\Upsilon(4S)$ machine but not for a symmetric $\Upsilon(4S)^+$ or polarized Z^0 one. Specifically the efficiencies in Table I are obtained as follows. The neutral kaon is K_S half of the time, and K_S decays to $\pi^+\pi^-$ 2/3 of the time. The neutral K^{*0} decays to $K^0\pi^0$ 1/3 of the time. We assume the ρ^0 is seen in its $\pi^+\pi^-$ decay, the η in its 2γ and $\pi^+\pi^-\pi^0$ decays, the ω in its $\pi^+\pi^-\pi^0$ decay, the η' in its $\rho^0\gamma$ decay, and the ϕ in its K^+K^- mode. The D^{*0} is seen in its $D^0\pi^0$ and $D^0\gamma$ decays, and the D^0 in all its possible CP eigenstate modes.

We attempt to estimate the efficiency of D^0,ϵ_D. The measured rates for the definite CP eigenstates,

$$D^0\to\pi^+\pi^-,K^+K^-,\bar{K}^0\pi^0,\bar{K}^0\rho^0,\bar{K}^0\phi,\rho^0\bar{K}^{*0}\ , \quad (4.1)$$

yield a total branching rate of 5%.[19] However, when each of the decay modes is weighed by its efficiency (as given in table I), the visible rate is lowered to $\epsilon_D\approx 2\%$. We included the $\rho^0\bar{K}^{*0}$ mode, since it was shown to have a dominant CP parity.[24]

If experiment were to reveal that each of the V_1V_2 modes are dominated by a single CP parity, then we ought to include them in our analysis. As no experimental data is yet available, we use the BSW model (with $a_1=1.3$, $a_2=-0.55$ and a D^0 lifetime of $\tau_D=0.42$ ps) for the modes

$$D^0\to\omega\phi,\omega\bar{K}^{*0},\rho^0\rho^0,K^{*-}K^{*+},\rho^+\rho\ \ ,\omega\omega\ . \quad (4.2)$$

We estimate a total branching rate of 5%, and a visible rate of $\epsilon_D\approx 2\%$, for the V_1V_2 modes in Eq. (4.2). A note about efficiencies used in the calculation is in order. Whereas the efficiencies of the particles involved in modes from cases (D1)–(D4) are given in Table I, the efficiencies for the ones from (D5) differ, as explained in the previous section. For (D5), the K^{*+} is seen in its $K^0\pi^+$ mode 2/3 of the time and in its $K^+\pi^0$ mode 1/3 of the time, $\epsilon=5/9$. The K^{*0} is seen in its $K^0\pi^0$ mode 1/3 of the time and in its $K^+\pi^-$ mode 2/3 of the time, $\epsilon=7/9$.

Now we summarize our findings about the D^0 efficiencies. When definite CP eigenstates are used [Eq. (4.1)], $\epsilon_D\approx 2\%$. When the V_1V_2 modes of Eq. (4.2) are proven to have dominant CP parities and are added, the efficiency of a D^0 into CP eigenstates doubles to $\epsilon_D\approx 4\%$, but is still 3% if vertexing of the D^0 is required.

Bauer, Stech, and Wirbel[11] obtained the rates of decays of B_d into the modes from sets (1) and (2) shown in Table II. They calculated within the factorization approximation and obtained small branching ratios, neglecting annihilation diagrams and final-state interactions. The modes in Table II have a combined "branching ratio $\times$ efficiency" of 3.1×10^{-4}. Folding in the efficiency of the D^{*0} and D^0 of $\sim 4\%$, a visible rate of 1.2×10^{-5} results for the decays of the form $B_d\to(f)_{\mathcal{D}}X^0$. Since the visible rate of 1.2×10^{-5} for the modes $(f)_{\mathcal{D}}X^0$ arises by the summation of many channels, a careful analysis is required to determine the CP parity and backgrounds of

TABLE II. Branching rates for various CP eigenstate modes of a B_d. The theoretical decay widths (second column) are taken from BSW (Ref. 11). The theoretical branching rates (third column) use $a_2=-0.24$ and a B_d lifetime of 1.2 ps. The visible branching rates into CP eigenstates, $B_d\to(f)_{\mathcal{D}}X^0$, (fourth column) is obtained by multiplying the theoretical branching rate (third column) with the efficiencies given in Table I.

B_d decay	Width (BSW) (10^{10} sec^{-1})	Branching ratio (BSW)	Branching ratio visible ($\times\epsilon_D$)
$B_d\to\bar{D}^0\pi^0$	$0.11a_2^2$	7.6×10^{-5}	7.6×10^{-5}
$B_d\to\bar{D}^{*0}\pi^0$	$0.16a_2^2$	1.1×10^{-4}	1.1×10^{-4}
$B_d\to\bar{D}^0\rho^0$	$0.06a_2^2$	4.1×10^{-5}	4.1×10^{-5}
$B_d\to\bar{D}^{*0}\eta^0$	$0.07a_2^2$	4.8×10^{-5}	2.9×10^{-5}
$B_d\to\bar{D}^0\omega$	$0.06a_2^2$	4.1×10^{-5}	3.7×10^{-5}
$B_d\to\bar{D}^0\eta'$	$0.03a_2^2$	2.1×10^{-5}	6.2×10^{-6}
$B_d\to\bar{D}^{*0}\eta'$	$0.04a_2^2$	2.8×10^{-5}	8.3×10^{-6}

the individual final states $(f)_{\mathcal{D}}X^0$. In contrast, the ψK_S is a single mode with definite CP parity, simple topology, minimal background and a large "branching ratio $\times$ efficiency" of $(3\times10^{-4})\times0.14\times2/3=2.8\times10^{-5}$. It appears that the ψK_S mode is much favored over the $(f)_{\mathcal{D}}X^0$ ones.

Is the situation as hopeless for CP violation in the quark process $\bar{b}\rightarrow\bar{c}+u\bar{d}$? Not really, since the branching rates into modes of sets (1) and (2) could be enhanced over BSW estimates. Final-state interactions can play a role in making the $\bar{D}^0(\bar{D}^{*0})X^0$ modes. For instance, if at first mainly the $D^-\pi^+$ mode is made, final-state interactions can still mix the $D^-\pi^+\leftrightarrow\bar{D}^0\pi^0$. It is conceivable[25] that $B(B_d\rightarrow\bar{D}^0\pi^0)$ is of the same order as the measured[26]

$$B(B_d\rightarrow D^-\pi^+)\approx0.25\% \ . \tag{4.3}$$

It is thus possible that one of the many modes from sets (1) or (2) could be enhanced. Ultimately, experiments will answer which modes are enhanced and which ones can be used for the asymmetry measurements.

A final comment. We are intrigued by the possibility that some of the $\bar{D}^{*0}X^0$ $(J=1)$ modes, such as $\bar{D}^{*0}\rho^0$ or $\bar{D}^{*0}\omega$, could have large branching ratios. Those modes have no definite CP parity. However, the (0, 0) helicity component has a definite CP parity and can be extracted as advocated in Ref. 27. Even in the most general case, when no CP dominates, a detailed study of all the angular correlations, including the angle of the two decay planes, makes those decay modes competitive with definite CP eigenstate ones.[28]

V. CONCLUSION

Large CP-violating effects are predicted with the decays $\bar{B}_d\rightarrow(f)_{\mathcal{D}}X^0$—generated by the quark process $\bar{b}\rightarrow\bar{c}+u\bar{d}$. Similar to the ψK_S mode, Section II showed that uncertainties in hadronic matrix elements and final-state phases do not enter in calculating the CP-violating interference term Imλ. To increase statistics one could sum over many modes $(f)_{\mathcal{D}}X^0$, listed in Sec. III. However, this summation must be done carefully, because the asymmetries flip sign for final states with different CP parities.

Section IV compared the $(f)_{\mathcal{D}}X^0$ modes to the classic ψK_S one. Whereas a visible rate within the BSW model of $\sim10^{-5}$ for the $(f)_{\mathcal{D}}X^0$ modes arises by the summation of many channels, the ψK_S mode—a single mode with definite CP parity and simple topology—has a visible rate of $\sim3\times10^{-5}$. The BSW model predicts that the $(f)_{\mathcal{D}}X^0$ modes are not competitive with the ψK_S one. Final-state interactions may enhance some of the $D^0(D^{*0})X^0$ modes, and may make the $(f)_{\mathcal{D}}X^0$ modes more competitive. Only future experiments will tell. This paper focused on final states $(f)_{\mathcal{D}}X^0$ that are CP eigenstates. We can increase statistics by including final states that do not have a definite CP parity, but are mixtures of CP-even and -odd eigenstates. This increment can be achieved by utilizing all the information one could gain from angular correlations.[27,28]

Even if the $(f)_{\mathcal{D}}X^0$ modes are not competitive with the ψK_S one, important information will be obtained by pursuing them. The standard model predicts the same CP-violating interference term Imλ, for the $(f)_{\mathcal{D}}X^0$, ψK_S, and D^+D^- modes. If new physics were to occur within the D^0-$\bar{D}^0$ complex, the interference term of the $(f)_{\mathcal{D}}X^0$ modes could differ from the ψK_S and D^+D^- ones,[12] and would present a violation of the standard model. This idea will be pursued in the future.

ACKNOWLEDGMENTS

We thank H. Quinn for many illuminating discussions, and J. Rosner whose criticism on an earlier version required us to study the helicity formalism. We are grateful to F. DeJongh, D. Hitlin, B. Kayser, M. Peskin, C. Simopoulos, and B. Ward for discussions.

APPENDIX

In principle, three weak phases $\xi_i=V_{ci}^*V_{ui}$ contribute to the $D^0\rightarrow\pi^+\pi^-$ decay. Unitarity of the 3×3 CKM matrix, however, reduces the number of independent weak phases to two:

$$A(\bar{D}^0\rightarrow\pi^+\pi^-)=\xi_d^*|a_d|e^{i\delta_d}+\xi_s^*|a_s|e^{i\delta_s} \ , \tag{A1}$$

$$A(D^0\rightarrow\pi^+\pi^-)=\xi_d|a_d|e^{i\delta_d}+\xi_s|a_s|e^{i\delta_s} \ , \tag{A2}$$

where $|a_i|$, δ_i $(i=d,s)$ are the respective hadronic matrix elements and final-state-interaction phases. Experiment informs us that[14,19]

$$\xi_s/\xi_d=-1+iO(\lambda_c^4), \quad |\xi_s/\xi_d|\approx1 \ , \tag{A3}$$

where $\lambda_c=0.22$ is the Cabibbo angle.

The s term, $\xi_s|a_s|e^{i\delta_s}$, represents rescattering effects, such as $D^0\rightarrow K^+K^-\rightarrow\pi^+\pi^-$, and penguin diagrams, $c\rightarrow ug$. Clearly, the decay $D^0\rightarrow\pi^+\pi^-$ is dominated by the d term $\xi_d|a_d|e^{i\delta_d}$. Even without this physical insight, we show that[29]

$$\frac{A(\bar{D}^0\rightarrow\pi^+\pi^-)}{A(D^0\rightarrow\pi^+\pi^-)}\approx\frac{V_{cd}V_{ud}^*}{V_{cd}^*V_{ud}} \tag{A4}$$

holds, barring a fine-tuning, because the weak phases of the d and s terms are almost the same. The amplitudes for the $\pi^+\pi^-$ mode of the $\bar{D}^0$ and D^0 are

$$A(\bar{D}^0\rightarrow\pi^+\pi^-)=\xi_d^*\left[|a_d|e^{i\delta_d}+\left[\frac{\xi_s}{\xi_d}\right]^*|a_s|e^{i\delta_s}\right] , \tag{A5}$$

$$A(D^0\rightarrow\pi^+\pi^-)=\xi_d\left[|a_d|e^{i\delta_d}+\frac{\xi_s}{\xi_d}|a_s|e^{i\delta_s}\right] . \tag{A6}$$

Because of Eq. (A3), the large brackets in Eqs. (A5) and (A6) cancel to excellent approximation when the ratio $A(\bar{D}^0\rightarrow\pi^+\pi^-)/A(D^0\rightarrow\pi^+\pi^-)$ is taken. Thus, Eq. (A4) holds.

Since the asymmetries that we expect for $B_d \to (f)_{\mathcal{D}} X^0$ are of the order of $|\mathrm{Im}\lambda| = 10-100\,\%$, we neglect the small asymmetries in the D^0 modes advocated by Golden and Grinstein:[30]

$$|A(\bar{D}^0 \to \pi^+\pi^-)| \neq |A(D^0 \to \pi^+\pi^-)| \; . \tag{A7}$$

Those small D^0 asymmetries arise precisely because of the imaginary part of ξ_s/ξ_d.

[1]I. I. Bigi and A. I. Sanda, Nucl. Phys. **B193**, 85 (1981).

[2]Report No. SLAC-353, 1989, edited by D. Hitlin.

[3]A. B. Carter and A. I. Sanda, Phys. Rev. Lett. **45**, 952 (1980); Phys. Rev. D **23**, 1567 (1981).

[4]L. Wolfenstein, Nucl. Phys. **B246**, 45 (1984).

[5]I. I. Y. Bigi and A. I. Sanda, Nucl. Phys. **B281**, 41 (1987); S. Stone, in *Proceedings of the Workshop Towards Establishing a b Factory,* Syracuse, New York, 1989, edited by M. Goldberg and S. Stone (Syracuse Univ., Syracuse, 1989), p. 171.

[6]A remark on notation. The generic *CP* eigenstate, that a D^0 or an excited D^0 decays into, is denoted by $(f)_{\mathcal{D}}$. Here $\mathcal{D}$ denotes the generic parent D^0 or D^0(excited). On the other hand, the *CP* eigenstate $(f)_D$ denotes the final state of the D^0 itself. At times we denote the *CP* eigenstate from a D^{*0} as $(f)_{D^*}$.

[7]The convention is used that B^0 has $\bar{b}$ content and that B^0_{phys} refers to a time evolving B^0 meson and not to a B^0 mass eigenstate.

[8]M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1973).

[9]C. O. Dib, I. Dunietz, F. J. Gilman, and Y. Nir, Phys. Rev. D **41**, 1522 (1990).

[10]Bigi and Sanda (Refs. 1 and 5); Wolfenstein (Ref. 4); I. Dunietz and J. L. Rosner, Phys. Rev. D **34**, 1404 (1986), Appendix.

[11]M. Bauer, B. Stech, and M. Wirbel, Z. Phys. C **34**, 103 (1987).

[12]In analogy to Y. Nir and D. Silverman, Nucl. Phys. **B345**, 301 (1990), we would get additional relations by using the quark subprocess $\bar{b} \to \bar{c} + u\bar{d}$ with the B_d and B_s systems. The B_s modes governed by that quark subprocess, such as $B_s \to \bar{D}^0(\bar{D}^{*0})K_S \to (f)_{\mathcal{D}} K_S$ or $B_s \to \bar{D}^0\bar{K}^* \to (f)_D(\pi^0 K_S)_{K^*}$, are predicted to have a tiny *CP*-violating interference term $\mathrm{Im}\lambda$ within the standard model. Measuring a large $\mathrm{Im}\lambda$ with those modes would present a violation of the standard model.

[13]Similarly if we parametrize the highly CKM-suppressed mode as $A(B_d \to D^0\pi^0) = V^*_{ub} V_{cd} |b| e^{i\tau}$, then its *CP*-conjugated mode is given by $A(\bar{B}_d \to \bar{D}^0\pi^0) = V_{ub} V^*_{cd} |b| e^{i\tau}$. Even though $\tau \neq \delta$ is likely, and thus τ represents different final-state phases, those highly CKM-suppressed amplitudes are negligible to the process at hand.

[14]L. Wolfenstein, Phys. Rev. Lett. **51**, 1945 (1983).

[15]Barring, of course, extreme fine-tuning, as discussed for the $D^0 \to \pi^+\pi^-$ example in the Appendix.

[16]J. C. Anjos *et al.*, Phys. Rev. Lett. **60**, 1239 (1988).

[17]If large *CP*-violating effects were to occur within the D^0-$\bar{D}^0$ system, then we would not be allowed to neglect those effects and a careful analysis would be required. However, within the standard model such *CP*-violating effects are not expected, because long-distance contributions probably dominate over those short-distance ones that have a different weak phase. When long-distance effects are prevalent, the "real" CKM elements of Eq. (2.13) are involved. Thus, the dispersive M_{12} and absorptive Γ_{12} D^0 mass matrix have the same phase and negligible *CP* violation occurs in the D^0-$\bar{D}^0$ system.

[18]If other $\bar{D}^0$ resonances occur appreciably in B_d decays, one would want to investigate them also.

[19]Particle Data Group, G. P. Yost *et al.*, Phys. Lett. B **204**, 1 (1988).

[20]E. Blucher *et al.*, IEEE Trans. Nucl. Sci. **32**, 716 (1985).

[21]Where kinematically allowed.

[22]We mean any final state that is made from two particles—one being taken from the first set of brackets and one from the second.

[23]I. Dunietz and A. Snyder, Report No. SLAC-PUB-5234, 1990 (unpublished).

[24]Mark III Collaboration, J. Adler *et al.*, Phys. Rev. Lett. **64**, 2615 (1990), observe that the P wave is much smaller than the S and D waves for the $D^0 \to \bar{K}^{*0}\rho^0$ decay. Thus, the final state has a dominant *CP* parity when $\bar{K}^{*0}$ is observed in its *CP* eigenmode $K_S\pi^0$.

[25]We thank Helen Quinn for discussions.

[26]D. L. Kreinick, in *Proceedings of the XIVth International Symposium on Lepton and Photon Interactions,* Stanford, California, 1989, edited by M. Riordan (World Scientific, Singapore, 1990), p. 129.

[27]B. Kayser, M. Kuroda, R. D. Peccei, and A. I. Sanda, Phys. Lett. B **237**, 508 (1990).

[28]I. Dunietz, H. Quinn, A. Snyder, W. Toki, and H. J. Lipkin, Phys. Rev. D (to be published).

[29]I. I. Bigi, in *Proceedings of the Tau-Charm Factory Workshop,* Stanford, California, 1989, edited by L. V. Beers (SLAC Report No. 343, Stanford, 1989), p. 169.

[30]M. Golden and B. Grinstein, Phys. Lett. B **222**, 501 (1989).