Partial-decay-rate asymmetries of charged bottom mesons and CP violation

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The partial-decay-rate asymmetries of charged-bottom-meson decays $B_u^{\pm}, B_c^{\pm} \to PP$ and PV are systematically studied and estimated. More attention has been paid to penguin contributions. We find that $B_c^- \to \rho^- \bar{D}^0, K^0F^-$ and $B_u^- \to F^-D^{0*}, F^{-*}D^0$ are extremely promising because of their large asymmetry and small number of $b\bar{b}$ pairs needed for testing them.

I. INTRODUCTION

The observation of CP-violation effects has been confined to the K^0-$\bar{K}^0$ system for more than 20 years since its discovery in 1964 (Ref. 1). Recently, there seems to be some evidence for same-sign dimuon events for the B^0-$\bar{B}^0$ complex in the UA1 experiment.[2] If it is confirmed, it would be the first evidence for B^0-$\bar{B}^0$ mixing. As we know, there are many speculations and estimates in the literature, which claimed that there should be large CP-violation effects in bottom-meson systems.[3] The prediction of the standard model for the same-sign dilepton asymmetry is quite small[4] ($\lesssim 1\%$), while the prediction for asymmetries in nonleptonic decays of neutral bottom mesons is quite large[5] ($\gtrsim 10\%$). But a systematic study[6] shows that it is not easy to see such effects although asymmetries are large. Thus, it is natural to think about charged-bottom-meson decays. Actually, people have done so.[7,8] Because the bottom mesons have unexpected longer lifetime (~ 1 psec) and because it is easy to tag on charged bottom mesons, we might have a better chance to detect CP-violation effects by measuring the partial-decay-rate asymmetries in $B_u^{\pm}, B_c^{\pm}$ decays. Owing to the nonexistence of mixing in charged-bottom-meson decays, we actually measure "direct" CP violation (i.e., $\Delta B = 1$). But, the discussions about these asymmetries in the literature[7,8] are very limited. Many important decay channels are missing, and there are some ambiguities in the formulas used in the calculation of decay amplitudes. So, it is worthwhile to examine this problem again. In this article we shall search for all possible decay channels of $B_{u,c}^{\pm} \to PP$ and PV which have nonvanishing asymmetry. Here, P and V denote a pseudoscalar and a vector meson, respectively. We shall also clarify some ambiguities and confusions in the calculation of the decay amplitudes.

II. PARTIAL-DECAY-RATE ASYMMETRIES

We confine ourselves in the standard Kobayashi-Maskawa (KM) model[9] with three generations of quarks and leptons. The partial-decay-rate asymmetry is defined as

$$A^{CP} = \frac{|A|^2 - |\bar{A}|^2}{|A|^2 - |\bar{A}|^2}, \tag{1}$$

where

$$A \equiv K_1 A_1 + K_2 A_2 , \tag{2}$$

$$\bar{A} \equiv K_1^* A_1 + K_2^* A_2 \tag{3}$$

are the decay amplitudes of $B_{u,c}^-$ and its antiparticle $B_{u,c}^+$, respectively. In Eqs. (2) and (3), K_1, K_2 are KM factors and can be expressed as $V_{i\alpha}V_{j\beta}^*$. A_1, A_2 are partial amplitudes including strong-interaction phases. Substituting Eqs. (2) and (3) into (1), we have

$$A^{CP} = -4\,\mathrm{Im}(K_1K_2^*)\mathrm{Im}(A_1A_2^*)/[4\,\mathrm{Re}(K_1K_2^*)\mathrm{Re}(A_1A_2^*) + 2(|K_1A_1|^2 + |K_2A_2|^2)] \propto s_1{}^2 s_2 s_3 s_\delta . \tag{4}$$

Here a universal factor $s_1{}^2 s_2 s_3 s_\delta$ for CP violation appears. This is common both for charged- and neutral-bottom-meson decays. If there are more than two partial amplitudes, Eq. (1) is unchanged, only Eqs. (2) and (3) should be extended to include more partial amplitudes.

Now the question is how to calculate the decay amplitudes. We follow the method suggested in Refs. 7, 8, 10, and 11.

Because the weak force is short ranged, the weak decays involve mainly the short-distance behavior of quarks. Therefore, the hard-gluon corrections are expected to modify the structure and strength of the weak coupling. But for charm decays, the comparison of the theoretical prediction and the experimental data shows that[11] the light quarks and soft gluons inside the hadron also have significant effects on the decay properties. For example, the predicted color suppression of D^0-$\bar{K}^0\pi^0$ is not seen in experiments, and the hard-gluon effects are also obscure. But, in bottom decays, because the momentum transfer is much larger than in the charm decays, the situation might be better. Thus, we still only consider hard-gluon effects and use the valence-quark approximation for initial and final hadron states in the calculation of the decay amplitudes.

For later use we list all the useful formulas for various matrix elements below:

$$\langle 0 | A_\mu(0) | P(q) \rangle = i f_P q_\mu \mathrm{tr}(\Gamma P) , \tag{5}$$

$$\langle 0 | V_\mu(0) | V(q) \rangle = \lambda_V m_V{}^2 \epsilon_\mu \mathrm{tr}(\Gamma V) , \tag{6}$$

where ϵ_μ is the polarization vector, A_μ and V_μ are axial-vector and vector currents, respectively, and the pseudomeson decay constants f_P are taken as (in GeV units)

$$f_\pi = 0.13, \quad f_K = 0.16 , \\ f_D = 0.19, \quad f_B = 0.23 , \tag{7}$$

while the dimensionless vector-meson decay constants λ_V are taken to be[10]

$$\lambda_V(q_1\bar{q}_2) = \left[\frac{2M_u}{M_{q_1}+M_{q_2}} \right]^{1/2} \lambda_{\rho^+} . \tag{8}$$

In the above, $\lambda_{\rho^+}=0.24$ and $M_u=M_d=0.32$ GeV, $M_s=0.48$ GeV, $M_c=1.55$ GeV, $M_b=4.7$ GeV, are the corresponding constituent-quark mass. Note that we only use constituent-quark mass in Eq. (8); in all other places we only use current-quark masses. For transition matrix elements, we use

$$\langle P_2(q_2) | V_\mu(0) | P_1(q_1) \rangle = F_d^V(Q^2) \left[(q_1+q_2)_\mu - \frac{m_1-m_2}{m_1+m_2}(q_1-q_2)_\mu \right] \mathrm{tr}(P_2^\dagger \Gamma P_1 - P_1 \Gamma P_2^\dagger) , \tag{9}$$

where $Q=q_1-q_2$;

$$\langle P_2(q_2) P_1(q_1) | V_\mu(0) | 0 \rangle = F_a^V(Q^2) \left[(q_1-q_2)_\mu - \frac{m_1-m_2}{m_1+m_2}(q_1+q_2)_\mu \right] \mathrm{tr}(P_2^\dagger \Gamma P_1 - P_1 \Gamma P_2^\dagger) , \tag{10}$$

where $Q=q_1+q_2$;

$$\langle V(q_2) | A_\mu(0) | P(q_1) \rangle = (m_1+m_2) F_d^A(Q^2) \left[\epsilon_\mu - \frac{2(q_1\cdot\epsilon)}{(m_1+m_2)^2-Q^2} q_{2\mu} \right] \mathrm{tr}(V^\dagger \Gamma P - P \Gamma V^\dagger) , \tag{11}$$

where $Q=q_1-q_2$;

$$\langle V(q_2) P(q_1) | A_\mu(0) | \; \rangle = (m_1+m_2) F_a^A(Q^2) \left[\epsilon_\mu + \frac{2(q_1\cdot\epsilon)}{(m_1+m_2)-Q_2} q_{2\mu} \right] \mathrm{tr}(V^\dagger \Gamma P^\dagger - P^\dagger \Gamma V^\dagger) , \tag{12}$$

where $Q=q_1+q_2$. In Eqs. (5)–(12), the traces should be calculated in flavor space, P, P_1, P_2, V, Γ are corresponding flavor wave functions of pseudomesons P, P_1, P_2, vector meson V, and the currents A_μ, V_μ, respectively. In the standard KM model, we have 6 flavors. Thus the flavor space is of dimension 6. We can use 6×6 matrices or bra and ket vectors to describe these flavor wave functions. For example, when we label the six flavors u,d,c,s,t,b as 1,2,3,4,5,6, we have

$$B_u^- = b\bar{u} = \begin{bmatrix} 0 & 0 & \cdots & 0 \\ 0 & 0 & \cdots & 0 \\ 0 & 0 & \cdots & 0 \\ 0 & 0 & .. & 0 \\ 0 & 0 & .. & 0 \\ 1 & 0 & \cdots & 0 \end{bmatrix} \equiv |6\rangle\langle 1| . \tag{13}$$

In general, $|i\rangle\langle j|$ is equivalent to a 6×6 matrix (a) with only nonvanishing unit element a_{ij}. For π^0 or ρ^0, we have

$$\pi^0(\rho^0) = \frac{u\bar{u}-d\bar{d}}{\sqrt{2}} = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 0 & \cdots & 0 \\ 0 & -1 & \cdots & 0 \\ 0 & 0 & \cdots & 0 \\ 0 & 0 & \cdots & 0 \\ 0 & 0 & \cdots & 0 \\ 0 & 0 & \cdots & 0 \end{bmatrix}$$

$$\equiv \frac{1}{\sqrt{2}}(|1\rangle\langle 1| - |2\rangle\langle 2|) . \tag{14}$$

Special attention should be paid to the current flavor wave functions Γ. Because currents are expressed by means of the field operator, they are just opposite from the description by use of the creation operator. For example, the current $A_\mu = \bar{d}\gamma_\mu(1-\gamma_5)b$ includes the combination of the creation operators $a_d^* b_b^*$, so,

$$\underset{\text{field operator}}{\bar{d}\gamma_\mu(1-\gamma_5)b \,(\text{or } \bar{d}\gamma_\mu b)} \sim \underset{\text{quarks}}{d\bar{b}} \equiv |2\rangle\langle 6| . \tag{15}$$

For the form factors in Eqs. (9)–(12), we take[7]

$$F_d^V(Q^2), F_d^A(Q^2) \sim 1 , \tag{16}$$

$$F_a^A(Q^2=m_B{}^2) \sim -2.3+0.38i , \tag{17}$$

$$\mathrm{Im} F_a^V(Q^2) \sim 5\times10^{-3} . \tag{18}$$

Actually, in Eq. (9), if we change $q_1 \to -q_1$, $P_1 \to P_1^\dagger$, we get Eq. (10). Similarly, by changing $q_1 \to -q_1$, $P \to P^\dagger$, we obtain Eq. (12) from (11).

Notice that we shall apply the same form factor of Eq. (17) to all the processes

$$B_u^- \to F^- D^{0*}, F^{-*} D^0 , \tag{19}$$

$$B_u^- \to \rho^0 K^-, \pi^0 K^{-*} , \tag{20}$$

$$B_c^- \to \psi D^- , \tag{21}$$

$$B_c^- \to \rho^0 D^-, \rho^- \bar{D}^0, \pi^0 D^{-*}, F^- K^{0*}, K^{-*} \bar{D}^0 . \tag{22}$$

This does not mean that the form factor of Eq. (17) is universal. Instead, it reflects the lack of our knowledge of the corresponding form factors for different processes. We know that the form factor of Eq. (17) is postulated by the authors of Ref. 7 based on some experimental factors when they calculated the matrix element $\langle D^- D^{0*} | A_\mu(0) | 0 \rangle$. For the processes (19), Eq. (17) seems applicable because of SU(3) symmetry. For the processes (20)—(22), Eq. (17) would be applicable if SU(4) was a good symmetry. But the SU(4) is badly broken, so using this equation is a rough qualitative estimation. Note that the final state of the process (21) possesses charm quantum number; this is different from the processes (19).

Now we are in a better shape in the definitions of all the matrix elements in Eqs. (5)—(17). All the sign and symbol confusions in the literature are clarified.

The effective Hamiltonian including hard-gluon corrections to next to leading logarithm takes the form[11,12]

$$H_{\rm eff}^W = \frac{G_F}{\sqrt{2}}(C_+ O_+ + C_- O_-) + \text{H.c.} + H_{\rm penguin} , \quad (23)$$

where

$$O_\pm = \tfrac{1}{2}[(\bar q_{2/3} O_\mu V q_{1/3})(\bar q_{1/3} O^\mu V^+ q_{2/3}) \pm (\bar q_{2/3} O_\mu V q_{2/3})(\bar q_{1/3} O^\mu V^+ q_{1/3})] , \quad (24)$$

$$O_\mu = \gamma_\mu (1-\gamma_5) , \quad (25)$$

C_+ and C_- are numerical constants and assume 0.84 and 1.41 for bottom decays, respectively.

The penguin diagram contribution $H_{\rm penguin}$ takes the form, for $b \to s$ decay,

$$H_{\rm penguin}^{b\to s} = C_P \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \bar s \gamma_\mu (1-\gamma_5) \times \lambda_a b (\bar c \gamma^\mu \lambda_a c + \bar u \gamma^\mu \lambda_a u + \bar d \gamma^\mu \lambda_a d + \bar s \gamma^\mu \lambda_a s) + \text{H.c.} , \quad (26)$$

for $b \to d$ decay,

$$H_{\rm penguin}^{b\to d} = C_P \frac{G_F}{\sqrt{2}} V_{tb} V_{td}^* \bar d \gamma_\mu (1-\gamma_5) \times \lambda_a b (\bar c \gamma^\mu \lambda_a c + \bar u \gamma^\mu \lambda_a u + \bar d \gamma^\mu \lambda_a d + \bar s \gamma^\mu \lambda_a s) + \text{H.c.} \quad (27)$$

In the above, λ_a are Gell-Mann SU(3) generators:

$$C_p = \frac{\alpha_s(K^2)}{12\pi} \ln \left[\frac{m_t^2}{K^2} \right] \sim (2\text{--}5) \times 10^{-2} \quad (28)$$

for $K^2 \sim m_b^2$. We take $C_p \sim 0.03$ in later estimations.

In order to illustrate how to use these formulas, we calculate the decay amplitude of $B_u^- \to \rho^0 K^-$. We follow Chau's classification of quark diagrams in Ref. 13. Then for $B_u^- \to \rho^0 K^-$, only spectator diagrams a, b, annihilation diagram d, and penguin diagram $e(b \to s)$ can contribute.

According to Eqs. (23)—(28) the effective Hamiltonian responsible for these diagrams can be rewritten as

$$H_{\rm eff}^W = \frac{G_F}{\sqrt{2}} V_{ub} V_{us}^* \left[\frac{C_+ + C_-}{2} \bar s \gamma_\mu (1-\gamma_5) u \bar u \gamma^\mu (1-\gamma_5) b + \frac{C_+ - C_-}{2} \bar s \gamma_\mu (1-\gamma_5) b \bar u \gamma^\mu (1-\gamma^5) u \right] + \text{H.c.} + H_{\rm penguin}^{b\to s} . \quad (29)$$

When we consider the matrix elements $\langle \rho^0 K^- | H_{\rm eff}^W | B_u^- \rangle$, we should also include the contribution of the Fierz transformed $H_{\rm eff}^W$. For instance, the term $\bar s \gamma_\mu (1-\gamma_5) b \bar u \gamma^\mu (1-\gamma_5) u$ does not contribute to diagrams a,b,d before making Fierz transformation. But it does after Fierz transformation. Because of the color mismatch, its total contribution after the Fierz transformation should have an extra factor of $\frac{1}{3}$. Thus we have

$$\langle \rho^0 K^- | H_{\rm eff}^W | B_u^- \rangle^{(a+b)} = \frac{G_F}{\sqrt{2}} V_{ub} V_{us}^* \frac{2C_+ + C_-}{3} \frac{-i}{\sqrt{2}} (P_K \cdot \epsilon_\rho) f_K F_d^A(m_K^2) 2m_\rho \frac{(m_B + m_\rho)^2}{(m_B + m_\rho)^2 + m_K^2} , \quad (30)$$

$$\langle \rho^0 K^- | H_{\rm eff}^W | B_u^- \rangle^{(d)} = \frac{G_F}{\sqrt{2}} V_{ub} V_{us}^* \frac{2C_+ + C_-}{3} \frac{-i}{\sqrt{2}} (P_K \cdot \epsilon_\rho) f_B F_a{}^A(m_B^2) 2m_\rho \frac{(m_K + m_\rho)^2}{(m_K + m_\rho)^2 - m_B^2} . \quad (31)$$

Note that the neutral-current interaction $\bar s \gamma_\mu (1-\gamma_5) b \bar u \gamma^\mu (1-\gamma_5) u$ also contributes to the decay $B_u^- \to \rho^0 K^-$. $\bar s \gamma_\mu (1-\gamma_5) u \bar u \gamma^\mu (1-\gamma_5) b$ does after the Fierz transformation. Therefore, we have an extra contribution from the neutral-current interaction:

$$\langle \rho^0 K^- | H_{\rm eff}^W | B_u^- \rangle^{(N)} = \frac{G_F}{\sqrt{2}} V_{ub} V_{us}^* \frac{2C_+ - C_-}{3} \frac{1}{\sqrt{2}} 2\lambda_{\rho+} m_\rho{}^2 F_d^V(m_\rho{}^2)(P_K \cdot \epsilon_\rho) . \quad (32)$$

For the penguin contribution, we have

$$\langle \rho^0 K^- | H_{\rm penguin}^{b\to s} | B_u^- \rangle = \frac{G_F}{\sqrt{2}} V_{ub} V_{ts}^* \tfrac{16}{9} C_p \frac{-i}{\sqrt{2}} \left[\left[\frac{-m_B{}^2}{(m_s + m_u)(m_u + m_b)} + \frac{1}{2} \right] f_B F_a^A(m_B{}^2) 2m_\rho \frac{(m_K + m_\rho)^2}{(m_K + m_\rho)^2 - m_B{}^2} + \left[\frac{-m_K{}^2}{(m_s + m_u)(m_u + m_b)} + \frac{1}{2} \right] f_K f_d^A(m_K{}^2) 2m_\rho \frac{(m_B + m_\rho)^2}{(m_B + m_\rho)^2 - m_K{}^2} \right] . \quad (33)$$

TABLE I. Partial-decay-rate asymmetries of $B_u^{\pm}$.

Decay channel	Quark diagram	A_{tree}^{CP}	A_{penguin}^{CP}	B	$N_{b\bar{b}}$
F^-D^0	$a+d+e$	7.3×10^{-6}	6.9×10^{-6}	2.1×10^{-2}	1.0×10^{12}
F^-D^{0*}	$a+d+e$	3.2×10^{-2}	7.8×10^{-3}	1.3×10^{-1}	1.3×10^{5}
$F^{-*}D^0$	$a+d+e$	0.18	-4.6×10^{-3}	1.0×10^{-1}	4.7×10^{5}
ρ^0K^-	$a+b+d+e$	0	1.7×10^{-2}	1.4×10^{-6}	2.5×10^{9}
π^0K^{-*}	$a+b+d+e$	0	9.4×10^{-2}	4.9×10^{-5}	2.3×10^{6}

In deriving (33), the equations of motion have been used:

$$(m_1+m_2)\bar{\psi}_1\gamma_5\psi_2=-i\partial_\mu(\psi_1\gamma^\mu\gamma_5\psi_2)\ , \qquad (34)$$

$$(m_1-m_2)\bar{\psi}_1\psi_2=-i\partial_\mu(\bar{\psi}_1\gamma^\mu\psi_2)\ . \qquad (35)$$

Actually, in (33), we only use Eq. (34). But for other decays, Eq. (35) is also used.

It is argued that the equation of motion can only be used safely for one- and two-body matrix elements with the particles involved being on mass shell.[14] We are just in this safe case.

Adding all the matrix elements (30)—(33) up, we get the total decay amplitude just in the form of Eqs. (2) and (3). So, the partial-decay-rate asymmetry is easy to calculate. For KM parameters, we take $s_1=0.231$, $s_2=s_1^2$, $s_3=0.5s_2$, $c_\delta\sim0$. We list all asymmetries for the decays $B_{u,c}^{\pm}\to PP,PV$, in Tables I and II, where we have set $s_\delta\sim1$. For comparison, we also give the asymmetries for the case in which the penguin contribution is omitted.

For the branching ratios, we first calculate the decay rates by means of the amplitudes, then use the measured total lifetime $\tau_B\sim10^{-12}$ sec to compute the branching ratios. We also list them in Tables I and II. Note that the amplitudes calculated here are Lorentz invariant. We must recover the omitted factor $(2E_B2E_12E_2)^{-1/2}$ to the amplitudes when we calculate the decay rates. Here E_B,E_1,E_2 are the energies of the bottom meson and the final-state mesons, respectively.

In general, if the asymmetry

$$A^{CP}=\frac{n_+-n_-}{n_++n_-} \qquad (36)$$

we search for one-standard-deviation (1σ) signature, then we need

$$n_++n_-=\frac{1}{(A^{CP})^2} \qquad (37)$$

$b\bar{b}$ pairs. If we search for three-standard-deviation (3σ) signature, we need

$$n_++n_-=9\frac{1}{(A^{CP})^2} \qquad (38)$$

$b\bar{b}$ pairs. So, the total number of $b\bar{b}$ pairs needed for 1σ signature is

$$N_{b\bar{b}}=\frac{1}{B}\frac{1}{(A^{CP})^2}\ . \qquad (39)$$

We put all the numbers $N_{b\bar{b}}$ in Tables I and II.

III. DISCUSSION AND CONCLUSION

From Tables I and II we see that there are six decay modes which have vanishing A_{tree}^{CP}. That means that without the penguin contribution, there would be no asymmetry. The other eight decay modes have nonvanishing asymmetry even without the penguin contribution.

Furthermore, because $|V_{td}|^2\sim s_1^2s_2^2\ll|V_{ts}^*|^2\sim s_2^2$, the contribution of $H_{\text{penguin}}^{b\to s}\gg H_{\text{penguin}}^{b\to d}$. In $B_u^{\pm}$ decays only $H_{\text{penguin}}^{b\to s}$ contributes, while in $B_c^{\pm}$ decays only $H_{\text{penguin}}^{b\to d}$ contributes (except for $B_c^-\to K^-\bar{D}^{0*},K^{-*}\bar{D}^0$). So, the penguin effects on $B_u^{\pm}$ decays are larger than on $B_c^{\pm}$ decays.

We also see from the tables that some decay modes have quite large asymmetry and small number of $b\bar{b}$ pairs needed for testing them. The most promising ones are $B_c^-\to\rho^-\bar{D}^0$, $B_c^-\to K^0F^-$, which need 2.6×10^4 and 1.5×10^4 $b\bar{b}$ pairs for 1σ signature. If we want to have a very clear signal, we need $N_{b\bar{b}}$ to be one order of magnitude larger than the number of Tables I and II.

For $B_u^{\pm}$ decays, F^-D^{0*} and $F^{-*}D^0$ modes are also quite promising. They need $\sim10^5$ $b\bar{b}$ pairs for 1σ signature. Note that in the CERN collider LEP, the Stanford Linear Collider, and the Cornell Electron Storage Ring, we can have about 10^5—10^6 $b\bar{b}$ pairs per year running.

TABLE II. Partial-decay-rate asymmetries of $B_c^{\pm}$.

Decay mode	Quark diagram	A_{tree}^{CP}	A_{penguin}^{CP}	B	$N_{b\bar{b}}$
K^0F^-	$d+e$	0	0.15	2.9×10^{-3}	1.5×10^{4}
ψD^-	$a+b+d+e$	0	4.0×10^{-3}	6.9×10^{-3}	9.1×10^{6}
ρ^0D^-	$b+d+e$	0.55		1.5×10^{-6}	2.2×10^{6}
$\rho^-\bar{D}^0$	$a+d+e$	0.91	0.87	5.0×10^{-5}	2.6×10^{4}
π^0D^{-*}	$b+d+e$	2.0×10^{-2}	3.0×10^{-2}	2.8×10^{-8}	4.0×10^{10}
K^0F^{-*}	$d+e$	0	6.4×10^{-2}	7.3×10^{-6}	3.3×10^{7}
$K^{0*}F^-$	$d+e$	0	$-2{\cdot}2\times10^{-2}$	5.2×10^{-5}	4.0×10^{7}
$K^-\bar{D}^{0*}$	$a+d+e$	1.1×10^{-2}	9.8×10^{-3}	5.6×10^{-5}	1.9×10^{8}
$K^{-*}\bar{D}^0$	$a+d+e$	-6.6×10^{-2}	-5.6×10^{-2}	2.3×10^{-4}	1.4×10^{6}

Considering the rather easy tagging on the charged bottom mesons, it is quite hopeful to test the asymmetries of these bottom decays in the near future.

Finally, as we emphasized at the beginning of Sec. II, the predicted color suppression is not seen in charm-decay experiments, and the hard-gluon effects are also obscure. We really do not know what the situation will be for bottom decays. In addition, the fact that $F^{\pm}$ have shorter lifetimes than $D^{\pm}$ means that the annihilation amplitude d can be as large as the spectator amplitude a in charm decay. What will happen for bottom decay? We do not know. All of these need to be tested in future experiments.

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