Newly observed two-body decays of B mesons in a hybrid perspective

K. Terasaki
Yukawa Institute for Theoretical Physics, Kyoto University, Kyoto 606-8502, Japan

(Received 28 October 2002; published 14 May 2003)

In consistency with $\overline{B}\to D^{(*)}\pi$, $J/\psi\overline{K}$, and $J/\psi\pi$ decays, recently observed $B^0\to D_s^+\pi^-$ and $\overline{B}^0\to D_s^+K^-$ decays are studied in a hybrid perspective in which their amplitude is given by a sum of factorizable and nonfactorizable ones.

DOI: 10.1103/PhysRevD.67.097501 PACS number(s): 13.25.Hw

(Quasi-)two-body decays of B mesons have been studied extensively by using factorization [1,2]. However, recently measured rates [3] for the color mismatched spectator (CMS) decays, $\overline{B}_d^0\to D^{(*)0}\pi^0$, are much larger than the expectation of factorization. It suggests that nonfactorizable contributions can play an important role in these decays. In addition, very recently, $\overline{B}^0\to D_s^+K^-$ and $B^0\to D_s^+\pi^-$ have been observed [4]. The rate for the former is again much larger than the expectation of factorization; i.e., it is expected to be strongly suppressed (the helicity suppression) since it is described by an annihilation diagram in the weak boson mass $m_W\to\infty$ limit. It means that the nonfactorizable contribution is dominant in this decay. The latter is a pure spectator decay, $\overline{b}\to\overline{u}+(c\overline{s})$, but does not satisfy the kinematical condition of color transparency [5], so that it is not very clear if the factorization works well in this decay. Therefore, it is meaningful to study a possible role of nonfactorizable contributions in the newly observed $\overline{B}_d^0\to D^{(*)0}\pi^0$, $\overline{B}^0\to D_s^+K^-$, and $B^0\to D_s^+\pi^-$ decays in consistency with the $b\to c$ type of decays, $\overline{B}\to D^{(*)}\pi$, $J/\psi\overline{K}$, and $J/\psi\pi$.

We first review briefly our (hybrid) perspective (see Ref. [6] for more details). Our starting point is to assume that the amplitude can be decomposed into a sum of factorizable and nonfactorizable ones ($M_{\rm FA}$ and $M_{\rm NF}$, respectively). $M_{\rm FA}$ is estimated by using the factorization while $M_{\rm NF}$ is assumed to be dominated by dynamical contributions of various hadron states and calculated by using a hard pion (or kaon) approximation in the infinite momentum frame (IMF) [7,8] since, in the existing theories such as QCD sum rule [9], QCD factorization [10], p QCD [11], and soft-collinear effective theory [12] which treat hadronic weak interactions of heavy mesons, it is still too complicated to account for dynamical contributions of various hadron states at the B meson mass scale. In this approximation, $M_{\rm NF}$ is given by a sum of the surface term ($M_{\rm S}$) which is given by a sum of all possible pole amplitudes and the equal-time commutator term ($M_{\rm ETC}$) which arises from the contribution of nonresonant (multihadron) intermediate states [13]. Corresponding to the above decomposition of the amplitude, the effective weak Hamiltonian, $H_w\simeq(G_F/\sqrt{2})\{c_1O_1+c_2O_2\}+$H.c. (where c_1 and c_2 are the Wilson coefficients with QCD corrections [14]), is decomposed into a sum of the Bauer-Stech-Wirbel (BSW) Hamiltonian [1], $H_w^{\rm (BSW)}$, and an extra term, $\widetilde{H}_w$, i.e., $H_w\to H_w^{\rm (BSW)}+\widetilde{H}_w$, by using the Fierz reshuffling, where $H_w^{\rm (BSW)}$ is given by a sum of products of colorless currents and might provide the factorizable amplitude. However, the "external" hadron states which sandwich $H_w^{\rm (BSW)}$ might interact sometimes with each other through hadron dynamics (like a rescattering, etc.). In this case, the corresponding part of the amplitude is nonfactorizable and should be included in $M_{\rm NF}$, so that the values of the coefficients, a_1 and a_2, in $M_{\rm FA}$ arising from $H_w^{\rm (BSW)}$ might not be the same as the original $a_1^{\rm (BSW)}=c_1+c_2/N_c$ and $a_2^{\rm (BSW)}=c_2+c_1/N_c$ in $H_w^{\rm (BSW)}$, where N_c is the color degree of freedom. Since such a hadron dynamics cannot be controlled by the perturbative QCD, we will treat a_1 and a_2 as adjustable parameters later. The extra term $\widetilde{H}_w$ which is given by a color singlet sum of colored current products provides nonfactorizable amplitudes in the present perspective, although, in Ref. [2], contributions from $\widetilde{H}_w$ have been included in the factorized amplitudes by considering the effective colors.

Explicit expression of factorized and nonfactorizable amplitudes for the $\overline{B}\to D^{(*)}\pi$, $J/\psi\overline{K}$, and $J/\psi\pi$ decays have already been given in Ref. [6] in which $M_{\rm ETC}$ and $M_{\rm S}$ with contributions of low lying meson poles are taken into account. In the same way, we can calculate the amplitude for the $\overline{b}\to\overline{u}+(c\overline{s})$ decays, $B^0\to D_s^{(*)+}\pi^-$. These amplitudes, however, include many parameters, i.e., form factors, decay constants of heavy mesons, asymptotic matrix elements of $\widetilde{H}_w$ (matrix elements of $\widetilde{H}_w$ taken between single hadron states with infinite momentum), phases, $\widetilde{\delta}_I$, of $M_{\rm ETC}^{(I)}(\overline{B}\to D\pi)$, ($I=\frac{1}{2}$ and $\frac{3}{2}$), relative to $M_{\rm S}$ and the relative phase (Δ) between $M_{\rm FA}$ and $M_{\rm NF}$ which has not been considered in our previous studies. The other parameters involved are known or can be estimated by using related experimental data and asymptotic flavor symmetries [15].

To obtain improved values of the above amplitudes, we update values of parameters involved. Asymptotic matrix elements of axial charges are estimated as follows [8], i.e., $|\langle\rho^0|A_{\pi^+}|\pi^-\rangle|\simeq1.0$ by using the partially conserved axial-vector current (PCAC) and the measured decay rate $\Gamma(\rho\to\pi\pi)_{\rm exp}\simeq150$ MeV [16]. Here we take $\langle\rho^0|A_{\pi^+}|\pi^-\rangle=1.0$ and the other ones can be related to it by using related asymptotic flavor symmetries, for example, $\sqrt{2}\langle D^{*+}|A_{\pi^+}|D^0\rangle=-\langle\rho^0|A_{\pi^+}|\pi^-\rangle$, etc., as in our previous study [6]. As the values of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements [17] and the decay constants, we take $V_{cs}\simeq V_{ud}\simeq0.98$, $V_{cd}\simeq-0.22$, $V_{cb}\simeq0.040$, $|V_{ub}/V_{cb}|\simeq0.090$, and $f_\pi\simeq130.7$ MeV, $f_K\simeq160$ MeV from Ref. [16]. The decay constant, $f_{J/\psi}\simeq406$ MeV, can be

TABLE I. Factorized and nonfactorizable amplitudes for the $\overline{B}\to D^{(*)}\pi$, $J/\psi\overline{K}$, $J/\psi\pi$, and $B^0\to D_s^{(*)+}\pi^-$ decays. The CKM matrix elements are factored out.

Decay	$A_{\rm FA}(\times 10^{-5}$ GeV)	$A_{\rm NF}(\times 10^{-5}$ GeV)
$\overline{B^0}\to D^+\pi^-$	$1.94 a_1 e^{i\Delta}$	$-\left\{\frac{4}{3}e^{i\tilde{\delta}_{1/2}}-\frac{1}{3}e^{i\tilde{\delta}_{3/2}}\right\}B_H$
$\overline{B^0}\to D^0\pi^0$	$-1.14\left(\frac{f_D}{0.226\text{ GeV}}\right)a_2 e^{i\Delta}$	$-\left\{2\sqrt{\frac{2}{3}}e^{i\tilde{\delta}_{1/2}}+\sqrt{\frac{2}{3}}e^{i\tilde{\delta}_{3/2}}\right\}B_H$
$B^-\to D^0\pi^-$	$1.94 a_1\left\{1+0.48\left(\frac{f_D}{f_\pi}\right)\left(\frac{a_2}{a_1}\right)\right\}e^{i\Delta}$	$e^{i\tilde{\delta}_{3/2}}B_H$
$\overline{B^0}\to D^{*+}\pi^-$	$-1.68 a_1 e^{i\Delta}$	$-0.694 B_H$
$\overline{B^0}\to D^{*0}\pi^0$	$1.07\left(\frac{f_{D^*}}{0.226\text{ GeV}}\right)a_2 e^{i\Delta}$	$0.983 B_H$
$B^-\to D^{*0}\pi^-$	$-1.68\, a_1\left\{1+0.52\left(\frac{f_{D^*}}{f_\pi}\right)\left(\frac{a_2}{a_1}\right)\right\}e^{i\Delta}$	$-0.696 B_H$
$B^-\to J/\psi K^-$	$-3.60 a_2 e^{i\Delta}$	$-0.548 B_H$
$\overline{B^0}\to J/\psi\overline{K}^0$	$-3.60 a_2 e^{i\Delta}$	$-0.548 B_H$
$B^-\to J/\psi\pi^-$	$-3.08 a_2 e^{i\Delta}$	$-0.692 B_H$
$\overline{B^0}\to J/\psi\pi^0$	$2.18 a_2 e^{i\Delta}$	$0.489 B_H$
$B^0\to D_s^+\pi^-$	$1.95 a_1 e^{i\Delta}$	$e^{i\tilde{\delta}_1}B_H$
$B^0\to D_s^{*+}\pi^-$	$1.54 a_1 e^{i\Delta}$	$0.70 B_H$

obtained from $\Gamma(J/\psi\to e^+e^-)_{\rm exp}=5.26\pm 0.37$ keV [16]. The updated values of the decay constants of heavy mesons, $f_D\simeq 0.226$ GeV, $f_{D_s}\simeq 0.250$ GeV, and $f_B\simeq 0.198$ GeV, are taken from the lattice QCD [18], and $f_{D^*}\simeq f_D$ and $f_{D_s^*}\simeq f_{D_s}$ are assumed as expected by the heavy quark effective theory (HQET) [19]. The form factors, $F_0^{(D\overline{B})}(m_\pi^2)$ and $A_0^{(D^*\overline{B})}(m_\pi^2)$, are estimated by using the HQET and the data on the semi-leptonic decays of B mesons [16] as $F_0^{(D\overline{B})}(m_\pi^2)\simeq 0.74$ and $A_0^{(D^*\overline{B})}(m_\pi^2)\simeq 0.65$. The form factors, $F_0^{(\pi\overline{B})}(q^2)$ and $F_1^{(\pi\overline{B})}(q^2)$, are estimated by using extrapolation formulas based on the lattice QCD [20]. We here take $F_0^{(\pi B)}(m_D^2)\simeq 0.28$, $F_0^{(\pi B)}(m_{D_s}^2)\simeq 0.32$, $F_1^{(\pi B)}(m_{D^*}^2)\simeq 0.34$, $F_1^{(\pi B)}(m_\psi^2)\simeq 0.50$, and $F_1^{(KB)}(m_\psi^2)\simeq 0.59$. The annihilation amplitudes which contain $F_0^{(D\pi)}(m_B^2)$ and $A_0^{(D^*\pi)}(m_B^2)$ will be small and neglected because of the helicity suppression.

The asymptotic matrix element of $\tilde{H}_w$ is parametrized by

$$\frac{\langle D^0|\tilde{H}_w^{(ud;cb)}|\overline{B}_d^0\rangle}{V_{cb}V_{ud}f_\pi}=B_H\times 10^{-5}\ \ (\text{GeV}), \tag{1}$$

where $\tilde{H}_w^{(ud;cb)}$ is a component of $\tilde{H}_w$ which is given by a sum of $\tilde{O}_1^{(ud;cb)}=V_{ud}V_{cb}\{2\Sigma_a(\overline{d}t^a u)_L(\overline{c}t^a b)_L\}$ and $\tilde{O}_2^{(ud;cb)}=V_{ud}V_{cb}\{2\Sigma_a(\overline{c}t^a u)_L(\overline{d}t^a b)_L\}$ with the color $SU_c(3)$ generator t^a. To evaluate the $\overline{B}\to D^*\pi$ amplitudes, we assume

$$\langle D^{*0}|\tilde{H}_w^{(ud;cb)}|\overline{B}_d^{*0}\rangle=\langle D^0|\tilde{H}_w^{(ud;cb)}|\overline{B}_d^0\rangle \tag{2}$$

as expected by the HQET. All the other asymptotic matrix elements of $\tilde{H}_w$ involved in the nonfactorizable amplitudes are combined with the ones in Eq. (2), i.e.,

$$\begin{aligned}\langle J/\psi|\tilde{H}_w^{(cd;cb)}|\overline{B}_d^{*0}\rangle&=\left(\frac{V_{cd}}{V_{cs}}\right)\langle J/\psi|\tilde{H}_w^{(cs;cb)}|\overline{B}_s^{*0}\rangle\\&=\left(\frac{V_{cd}}{V_{ud}}\right)\langle D^{(*)0}|\tilde{H}_w^{(ud;cb)}|\overline{B}_d^{(*)0}\rangle\\&=-\left(\frac{V_{cd}V_{cb}}{V_{cs}V_{ub}}\right)\langle D_s^{(*)+}|\tilde{H}_w^{(cs;ub)}|B_u^{(*)+}\rangle,\end{aligned} \tag{3}$$

by inserting commutation relations, $[V_{K^0},\tilde{H}_w^{(cs;cb)}]=(V_{cs}/V_{cd})\tilde{H}_w^{(cd;cb)}$, $[V_{D^0},\tilde{H}_w^{(cd;cb)}]=(V_{cd}/V_{ud})\tilde{H}_w^{(ud;cb)}$, $[V_{\overline{D}^0},\tilde{H}_w^{(cs;cb)}]=(V_{cb}/V_{ub})\tilde{H}_w^{(cs;ub)}$, between related asymptotic states (single hadron states with infinite momentum) and using asymptotic $SU_f(3)$ and $SU_f(4)$ relations, $\langle\overline{B}_s^{*0}|V_{K^0}|\overline{B}_d^{*0}\rangle=-1$, $\langle D^{*0}|V_{D^0}|J/\psi\rangle=-1$, etc. To obtain the last equality in Eq. (3), we have used the CP invariance which is always assumed in this Brief Report and $\langle\{q\overline{q}\}_0|\tilde{O}_+|\{q\overline{q}\}_0\rangle=0$ from a quark counting [21], where $\tilde{O}_\pm=\tilde{O}_1\pm\tilde{O}_2$. The $\{q\overline{q}\}_0$'s denote the low lying mesons. In this way, we can obtain $M_{\rm FA}$ and $M_{\rm NF}$ in the second and third columns, respectively, of Table I, where we have neglected small contributions of annihilation terms in $M_{\rm FA}$ and excited meson poles in $M_{\rm NF}$.

TABLE II. A typical result on the branching ratios ($\times 10^{-3}$) for $\bar{B}\to D^{(*)}\pi$, $J/\psi\bar{K}$, and $J/\psi\pi$ decays, where the values of the parameters involved are given in the text. $\mathcal{B}_{\rm FA}$ and $\mathcal{B}_{\rm tot}$ are given by $M_{\rm FA}$ and $M_{\rm tot}$, respectively. $\mathcal{B}_{\rm exp}$ are taken from Ref. [22].

Decays	$\mathcal{B}_{\rm FA}$	$\mathcal{B}_{\rm tot}$	$\mathcal{B}_{\rm exp}$
$\bar{B}^0\to D^+\pi^-$	4.0	3.1	3.0 ± 0.4
$\bar{B}^0\to D^0\pi^0$	0.10	0.24	0.27 ± 0.06
$B^-\to D^0\pi^-$	5.6	5.6	5.3 ± 0.5
$\bar{B}^0\to D^{*+}\pi^-$	3.1	2.6	2.76 ± 0.21
$\bar{B}^0\to D^{*0}\pi^0$	0.09	0.22	0.22 ± 0.10
$\bar{B}^0\to D^{*0}\pi^-$	4.1	4.7	4.6 ± 0.4
$B^-\to J/\psi K^-$	0.82	0.99	1.01 ± 0.05
$\bar{B}^0\to J/\psi\bar{K}^0$	0.75	0.91	0.87 ± 0.05
$B^-\to J/\psi\pi^-$	0.030	0.039	0.042 ± 0.007
$\bar{B}^0\to J/\psi\pi^0$	0.014	0.018	0.021 ± 0.005

We now look for values of parameters, a_1, a_2, Δ, $\tilde{\delta}_I$, ($I=1/2$, and 3/2), and B_H, which reproduce the measured branching ratios for the $\bar{B}\to D^{(*)}\pi$, $J/\psi\bar{K}$, and $J/\psi\pi$ decays. a_1 and a_2 are treated as adjustable parameters with values around $a_1^{\rm (BSW)}$ and $a_2^{\rm (BSW)}$. The phase $\tilde{\delta}_I$ is restricted in the region $|\tilde{\delta}_I|<90°$ since resonant contributions have already been extracted as pole amplitudes in $M_{\rm S}$ while Δ and B_H are treated as free parameters. The result is not very sensitive to $\tilde{\delta}_I$, and the coefficients, a_1 and a_2, favor values close to the ones taken in Ref. [2] which is based on the factorization. The above implies that the nonfactorizable contribution is not very important in the color favored decays. We can reproduce the experimental data ($\mathcal{B}_{\rm exp}$) compiled by the Particle Data Group 2002 [22] taking values of parameters in the range $1.00\lesssim a_1\lesssim 1.13$, $0.28\lesssim a_2\lesssim 0.31$, $24°\lesssim|\Delta|\lesssim 32°$, $|\delta_{1/2}|\lesssim 70°$, $10°\lesssim|\delta_{3/2}|\lesssim 90°$, and $0.09\lesssim B_H\lesssim 0.25$. To see more explicitly a role of the nonfactorizable contribution, we list a typical result on the branching ratios (near the best fit to $\mathcal{B}_{\rm exp}$) for $a_1=1.08$, $a_2=0.29$, $\tilde{\delta}_1=0.0°$, $\tilde{\delta}_3=\pm 90°$, $\Delta=\pm 28°$, and $B_H=0.19$ in Table II, where we have used $\tau(B^-)=1.67\times 10^{-12}$ s and $\tau(\bar{B}^0)=1.54\times 10^{-12}$ s from Ref. [22]. $\mathcal{B}_{\rm FA}$ and $\mathcal{B}_{\rm tot}$ are given by $M_{\rm FA}$ and $M_{\rm tot}=M_{\rm FA}+M_{\rm NF}$, respectively. As seen in Table II, $\mathcal{B}_{\rm FA}$ in which $M_{\rm NF}$ is discarded is hard to reproduce the data on the CMS decays, $\bar{B}\to D^{(*)0}\pi^0$. If we add $M_{\rm NF}$, however, we can get a much better fit to the data including the CMS decays. In the color favored $\bar{B}\to D^{(*)}\pi$ decays, $M_{\rm NF}$ is rather small (but it can interfere efficiently with the main amplitude, $M_{\rm FA}$). In the $B^-\to D^0\pi^-$ decay, however, it is very small. In the $\bar{B}\to J/\psi\bar{K}$ and $J/\psi\pi$ decays, the color suppression does not work so well that $M_{\rm NF}$ is not dominant in contrast with the $\bar{B}\to D^{(*)0}\pi^0$ although all of them are the CMS decays.

Next, we study the $B^0\to D_s^{(*)+}\pi^-$ decays comparing with the $B^-\to D^0\pi^-$ which has been studied above. Using the same values of parameters as the above, i.e., $a_1=1.08$, $a_2=0.29$, $B_H=0.19$, we obtain

$$|M_{\rm NF}(B^0\to D_s^+\pi^-)|\simeq 0.09|M_{\rm FA}(B^0\to D_s^+\pi^-)|, \quad (4)$$

$$|M_{\rm NF}(B^0\to D_s^{*+}\pi^-)|\simeq 0.08|M_{\rm FA}(B^0\to D_s^{*+}\pi^-)|, \quad (5)$$

which imply that the factorization works considerably well in these decays although they do not satisfy the condition of the color transparency. Neglecting the rather small $M_{\rm NF}$ in the $B^0\to D_s^+\pi^-$ and using the same values of parameters as the above, we obtain

$$|M(B^0\to D_s^+\pi^-)|\simeq 0.074|M(B^-\to D^0\pi^-)|, \quad (6)$$

where we have used $|V_{ub}/V_{cb}|_{\rm exp}\simeq 0.090$ [16]. The measured branching ratio for the $B^-\to D^0\pi^-$ decay [22] leads us to

$$\mathcal{B}(B^0\to D_s^+\pi^-)\simeq 2.7\times 10^{-5}, \quad (7)$$

which reproduces well the recent measurements [4],

$$\mathcal{B}(B^0\to D_s^+\pi^-)_{\rm BABAR}=(3.1\pm 2.0)\times 10^{-5},$$

$$\mathcal{B}(B^0\to D_s^+\pi^-)_{\rm BELLE}=(2.4^{+1.0}_{-0.8}\pm 0.7)\times 10^{-5}.$$

In the same way, we obtain $\mathcal{B}(B^0\to D_s^{*+}\pi^-)\simeq 1.7\times 10^{-5}$, which is again compatible with the experimental upper limits [4].

In the $\bar{B}^0\to D_s^+K^-$ decay, $M_{\rm FA}$ is strongly suppressed because of the helicity suppression, so that $M_{\rm NF}$ dominates the decay in the present perspective, i.e.,

$$M(\bar{B}^0\to D_s^+K^-)\simeq M_{\rm NF}(\bar{B}^0\to D_s^+K^-)$$
$$\simeq -iV_{cb}V_{ud}\left(\frac{f_\pi}{f_K}\right)\frac{\langle D^0|\tilde{H}_w|\bar{B}_d^0\rangle}{V_{cb}V_{ud}f_\pi}e^{i\tilde{\delta}_1}. \quad (8)$$

The same value of parameters as the above leads to

$$\mathcal{B}(\bar{B}^0\to D_s^+K^-)\simeq 2.8\times 10^{-5}, \quad (9)$$

which should be compared with the measured values [4]

$$\mathcal{B}(\bar{B}^0\to D_s^+K^-)_{\rm BABAR}=(3.2\pm 2.0)\times 10^{-5},$$

$$\mathcal{B}(\bar{B}^0\to D_s^+K^-)_{\rm BELLE}=(4.6^{+1.2}_{-1.1}\pm 1.3)\times 10^{-5}.$$

In summary, we have studied the recently observed decays, $\bar{B}\to D^{(*)0}\pi^0$, $B^0\to D_s^+\pi^-$, and $\bar{B}^0\to D_s^+K^-$, in consistency with the $b\to c$ type of decays, $\bar{B}\to D^{(*)}\pi$, $J/\psi\bar{K}$, and $J/\psi\pi$, providing their amplitude by a sum of factorized and nonfactorizable ones. To study the nonfactorizable amplitudes, we have used the asymptotic $SU_f(3)$ and $SU_f(4)$ symmetries which may be broken. The size of the symmetry breaking can be estimated from the value of the form factor, $f_+(0)$, in the matrix element of the related vector current, where $f_+(0)=1$ in the symmetry limit. From the measured values of the form factors, $f_+^{(\pi D)}(0)=0.71\pm 0.06$ [23] and $|f_+^{(\pi D)}(0)/f_+^{(\bar{K}D)}(0)|=1.00\pm 0.13$ [24], the asymptotic $SU_f(4)$ symmetry seems to be broken to the extent of 30%

while the asymptotic $SU_f(3)$ still works well. However, such a large symmetry breaking has not caused any serious problem in the present study since $M_{\rm NF}$ is much smaller than $M_{\rm FA}$ except for some decays in which $M_{\rm FA}$ is strongly suppressed and whose experimental errors are still large. For more precise studies, of course, more detailed information of the symmetry breaking will be needed.

The amplitude with final state interactions has been included in the nonfactorizable one. For the color favored $\bar{B} \to D^{(*)}\pi$ decays, $M_{\rm NF}$ has been rather small and, therefore, the final state interactions seem to be not very important (but not necessarily negligible) in these decays. In the $\bar{B} \to D^{(*)0}\pi^0$ which are the CMS decays, $M_{\rm NF}$ has been dominant since $M_{\rm FA}$ is suppressed because of the color suppression. In the $\bar{B} \to J/\psi \bar{K}$ and $J/\psi\pi$, which also are the CMS decays, however, the color suppression has not worked so well that $M_{\rm NF}$ has not been dominant in contrast with the $\bar{B} \to D^{(*)0}\pi^0$ decays. In the $B^0 \to D_s^{(*)+}\pi^-$ decays which are the color favored $\bar{b} \to \bar{u} + (c\bar{s})$ type of spectator decays, $M_{\rm NF}$ has been small. The values of parameters which reproduce the measured branching ratios for the $\bar{B} \to D^{(*)}\pi$, $J/\psi\bar{K}$, and $J/\psi\pi$ decays have led to $\mathcal{B}(B^0 \to D_s^{(*)+}\pi^-)$ consistent with the very recent measurements. It means that the factorization works considerably well in these decays although they do not satisfy the condition of color transparency. In the $\bar{B}^0 \to D_s^+ K^-$ decay which is the annihilation decay, $M_{\rm NF}$ has been dominant and reproduced the very recent measurements within their large errors. All the above suggest that dynamical contributions of hadrons should be carefully treated in hadronic weak decays of B mesons.

In the CMS decays, $\bar{B}^0 \to D^{(*)0}\pi^0$, $J/\psi\bar{K}$, and $J/\psi\pi$, the annihilation decay, $\bar{B}^0 \to D_s^+ K^-$, and the $\bar{b} \to \bar{u} + (c\bar{s})$ type of spectator decay, $B^0 \to D_s^+\pi^-$, both of the theoretical and experimental ambiguities are still large although their measured rates have been reproduced considerably well by taking account of the nonfactorizable contributions. More theoretical and experimental studies on these decays will be needed.

The author would like to acknowledge Professor T. Onogi for discussions. He also thanks the Yukawa Institute for Theoretical Physics at Kyoto University. Discussions during the YITP workshop, YITP-W-02-08, ''TEA (Theoretical-Experimental-Astronomical Particle Physics) 02,'' were useful to complete this work.

[1] M. Bauer, B. Stech, and M. Wirbel, Z. Phys. C **34**, 103 (1987); M. Neubert, V. Rickert, B. Stech, and Q. P. Xu, in *Heavy Flavours*, edited by A. J. Buras and M. Lindner (World Scientific, Singapore, 1992), p. 286.
[2] M. Neubert, *Heavy Flavours II*, edited by A. J. Buras and M. Lindner (World Scientific, Singapore, 1998), p. 239.
[3] BELLE Collaboration, K. Abe *et al.*, Phys. Rev. Lett. **88**, 052002 (2002); CLEO Collaboration, T. E. Coan *et al.*, *ibid.* **88**, 062001 (2002).
[4] BABAR Collaboration, B. Aubert *et al.*, hep-ex/0207053; BELLE Collaboration, P. Krokovny, talk given at the 31st International Conference on High Energy Physics, Amsterdam, 2002; BELLE Collaboration, K. Abe *et al.*, BELLE-CONF-0259 (2002).
[5] J. D. Bjorken, Nucl. Phys. B (Proc. Suppl.) **11**, 325 (1989); M. J. Dugan and B. Grinstein, Phys. Lett. B **255**, 583 (1991).
[6] K. Terasaki, Int. J. Theor. Phys., Group Theor. Nonlinear Opt. **8**, 55 (2002), and references therein.
[7] K. Terasaki, S. Oneda, and T. Tanuma, Phys. Rev. D **29**, 456 (1984).
[8] S. Oneda and K. Terasaki, Prog. Theor. Phys. Suppl. **82**, 1 (1985).
[9] B. Block and M. Shifman, Nucl. Phys. **B389**, 534 (1993).
[10] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Nucl. Phys. **B591**, 313 (2000).
[11] Y.-Y. Keum, H.-N. Li, and A. I. Sanda, Phys. Rev. D **63**, 074006 (2001).
[12] C. W. Bauer, D. Pirjor, and I. W. Stewart, Phys. Rev. D **65**, 054022 (2002).
[13] V. S. Mathur and L. K. Pandit, in *Advances in Particle Physics*, edited by R. L. Cool and R. E. Marshak (Interscience, New York, 1968), Vol. 2, p. 383.
[14] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[15] Asymptotic flavor symmetry is the flavor symmetry of asymptotic matrix elements—matrix elements of charges, currents, and some other operators (like the weak Hamiltonian constructed by them) taken between single hadron states with infinite momentum. A comprehensive review on the asymptotic flavor symmetry and its fruitful results has been provided in Ref. [8].
[16] Particle Data Group, D. E. Groom *et al.*, Eur. Phys. J. C **15**, 1 (2000).
[17] N. Cabibbo, Phys. Rev. Lett. **10**, 531 (1963); M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1973).
[18] C. T. Sachrajda, Nucl. Instrum. Methods Phys. Res. A **462**, 23 (2001).
[19] N. Isgur and M. B. Wise, Phys. Lett. B **232**, 113 (1989); **237**, 527 (1990); H. Georgi, *ibid.* **240**, 447 (1990).
[20] A. Abada *et al.*, Nucl. Phys. **B619**, 565 (2001).
[21] K. Terasaki, Phys. Rev. D **59**, 114001 (1999).
[22] Particle Data Group, K. Hagiwara *et al.*, Phys. Rev. D **66**, 010001 (2002).
[23] Fermilab E687 Collaboration, P. L. Frabetti *et al.*, Phys. Lett. B **364**, 127 (1995).
[24] Fermilab E687 Collaboration, M. S. Nehring *et al.*, Nucl. Phys. B (Proc. Suppl.) **55A**, 131 (1997).