Charmless Hadronic Two-body Decays of B Mesons

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A way of circumventing the gauge and infrared problems with effective Wilson coefficients is shown. Implications of experimentally measured charmless B decays are discussed.

I. INTRODUCTION

In past years we have witnessed remarkable progress in the study of exclusive charmless B decays. Experimentally, CLEO [1] has discovered many new two-body decay modes $B \to \eta' K^\pm$, $\eta' K^0$, $\pi^\pm K^0$, $\pi^\pm K^\mp$, $\pi^0 K^\pm$, $\rho^0\pi^\pm$, $\omega K^\pm$ and found a possible evidence for $B \to \phi K^*$. Moreover, CLEO has provided new improved upper limits for many other decay modes. While all the measured channels are penguin dominated, the most recently observed $\rho^0\pi^-$ mode is dominated by the tree diagram. In the meantime, updates and new results of many $B \to PV$ decays with $P = \eta, \eta', \pi, K$ and $V = \omega, \phi, \rho, K^*$ as well as $B \to PP$ decays will be available soon. With the B factories Babar and Belle starting to collect data, many exciting and harvest years in the arena of B physics and CP violation are expected to come. Theoretically, many significant improvements and developments have been achieved over past years. For example, a next-to-leading order effective Hamiltonian for current-current operators and QCD as well as electroweak penguin operators becomes available. The renormalization scheme and scale problems with the factorization approach for matrix elements can be circumvented by employing scale- and scheme-independent effective Wilson coefficients. Heavy-to-light form factors have been computed using QCD sum rules, lattice QCD and potential models.

II. EFFECTIVE WILSON COEFFICIENTS

Although the hadronic matrix element $\langle O(\mu)\rangle$ can be directly calculated in the lattice framework, it is conventionally evaluated under the factorization hypothesis so that $\langle O(\mu)\rangle$ is factorized into the product of two matrix elements of single currents, governed by decay constants and form factors. In spite of its tremendous simplicity, the naive factorization approach encounters two principal difficulties. One of them is that the hadronic matrix element under factorization is renormalization scale μ independent as the vector or axial-vector current is partially conserved. Consequently, the amplitude $c_i(\mu)\langle O\rangle_{\rm fact}$ is not truly physical as the scale dependence of Wilson coefficients does not get compensation from the matrix elements. A plausible solution to the aforementioned scale problem is to extract the μ dependence from the matrix element $\langle O(\mu)\rangle$, combine it with the μ-dependent Wilson coefficients to form μ-independent effective Wilson coefficients. However, it was pointed out recently in [2] that $c_i^{\rm eff}$ suffer from the gauge and infrared ambiguities since an off-shell external quark momentum, which is usually chosen to regulate the infrared divergence occurred in the radiative corrections to the local 4-quark operators, will introduce a gauge dependence. Therefore, this solution, though removes the scale and scheme dependence of a physical amplitude in the framework of the factorization hypothesis, often introduces the infrared cutoff and gauge dependence.

It was recently shown in [3] that the above-mentioned problems on gauge dependence and infrared singularity associated with the effective Wilson coefficients can be resolved by perturbative QCD (PQCD) factorization theorem. In this formalism, partons, *i.e.*, external quarks, are assumed to be on shell, and both ultraviolet and infrared divergences in radiative corrections are isolated using the dimensional regularization. Because external quarks are on shell, gauge invariance of the decay amplitude is maintained under radiative corrections to all orders. This statement is confirmed by an explicit one-loop calculation in [3]. The obtained ultraviolet poles are subtracted in a renormalization scheme, while the infrared poles are absorbed into universal nonperturbative bound-state wave functions. The remaining finite piece is grouped into a hard decay subamplitude. The decay rate is then factorized into the convolution of the hard subamplitude with the bound-state wave functions, both of which are well-defined and gauge invariant. Explicitly, the effective Wilson coefficient has the expression

$$c^{\rm eff} = c(\mu)g_1(\mu)g_2(\mu_f)\ , \tag{1}$$

where $g_1(\mu)$ is an evolution factor from the scale μ to m_b, whose anomalous dimension is the same as that of $c(\mu)$, and $g_2(\mu_f)$ describes the evolution from m_b to μ_f (μ_f being a factorization scale arising from the dimensional regularization of infrared divergences), whose anomalous dimension differs from that of $c(\mu)$ because of the inclusion of the dynamics associated with spectator quarks. Setting $\mu_f = M_b$, the effective Wilson coefficients obtained from the one-loop vertex corrections to the 4-quark operators O_i have the form:

$$c_i^{\rm eff}\Big|_{\mu_f=m_b} = c_i(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{ij} c_j(\mu) + \cdots, \tag{2}$$

where the matrix $\hat{r}$ gives momentum-independent constant terms which depend on the treatment of γ_5. The expression of $\hat{r}$ for $\Delta B = 1$ transition current-current operators O_1, O_2 is given in [3], while the complete result for QCD-penguin operators $O_3, \cdots, O_6$ and electroweak penguin operators $O_7, \cdots, O_{10}$ is given in [4]. It should be accentuated that, contrary to the previous work based on Landau gauge and off-shell regularization [5], the matrix $\hat{r}$ given in Eq. (2) is gauge invariant. Consequently, the effective Wilson coefficients (2) are not only scheme and scale independent but also free of gauge and infrared problems.

III. EFFECTIVE PARAMETERS AND NONFACTORIZABLE EFFECTS

It is known that the effective Wilson coefficients appear in the factorizable decay amplitudes in the combinations $a_{2i} = c_{2i}^{\rm eff} + \frac{1}{N_c}c_{2i-1}^{\rm eff}$ and $a_{2i-1} = c_{2i-1}^{\rm eff} + \frac{1}{N_c}c_{2i}^{\rm eff}$ ($i = 1, \cdots, 5$). Phenomenologically, the number of colors N_c is often treated as a free parameter to model the nonfactorizable contribution to hadronic matrix elements and its value can be extracted from the data of two-body nonleptonic decays. As shown in [6], nonfactorizable effects in the decay amplitudes of $B \to PP$, VP can be absorbed into the parameters $a_i^{\rm eff}$. This amounts to replacing N_c in $a_i^{\rm eff}$ by $(N_c^{\rm eff})_i$. Explicitly,

$$a_{2i}^{\rm eff} = c_{2i}^{\rm eff} + \frac{1}{(N_c^{\rm eff})_{2i}}c_{2i-1}^{\rm eff}, \qquad a_{2i-1}^{\rm eff} = c_{2i-1}^{\rm eff} + \frac{1}{(N_c^{\rm eff})_{2i-1}}c_{2i}^{\rm eff}, \qquad (i = 1, \cdots, 5), \tag{3}$$

where $(1/N_c^{\rm eff})_i \equiv (1/N_c) + \chi_i$ with χ_i being the nonfactorizable terms which receive main contributions from color-octet current operators [7]. In the absence of final-state interactions, we shall assume that χ_i and hence $N_c^{\rm eff}$ are real. If χ_i are universal (i.e. process independent) in charm or bottom decays, then we have a generalized factorization scheme in which the decay amplitude is expressed in terms of factorizable contributions multiplied by the universal effective parameters $a_i^{\rm eff}$. Phenomenological analyses of the two-body decay data of D and B mesons indicate that while the generalized factorization hypothesis in general works reasonably well, the effective parameters $a_{1,2}^{\rm eff}$ do show some variation from channel to channel, especially for the weak decays of charmed mesons. A recent updated analysis of $B \to D\pi$ data gives [8]

$$N_c^{\rm eff}(B \to D\pi) \sim (1.8 - 2.2), \qquad \chi_2(B \to D\pi) \sim (0.12 - 0.21). \tag{4}$$

It is customary to assume in the literature that $(N_c^{\rm eff})_1 \approx (N_c^{\rm eff})_2 \cdots \approx (N_c^{\rm eff})_{10}$ so that the subscript i can be dropped; that is, the nonfactorizable term is usually postulated to behave in the same way in penguin and tree decay amplitudes. A closer investigation shows that this is not the case. It has been argued in [9] that nonfactorizable effects in the matrix elements of $(V-A)(V+A)$ operators are *a priori* different from that of $(V-A)(V-A)$ operators. One primary reason is that the Fierz transformation of the $(V-A)(V+A)$ operators $O_{5,6,7,8}$ is quite different from that of $(V-A)(V-A)$ operators $O_{1,2,3,4}$ and $O_{9,10}$. As a result, contrary to the common assertion, $N_c^{\rm eff}(LR)$ induced by the $(V-A)(V+A)$ operators are theoretically different from $N_c^{\rm eff}(LL)$ generated by the $(V-A)(V-A)$ operators [9]. Therefore, we shall assume that

$$\begin{aligned} N_c^{\rm eff}(LL) &\equiv \left(N_c^{\rm eff}\right)_1 \approx \left(N_c^{\rm eff}\right)_2 \approx \left(N_c^{\rm eff}\right)_3 \approx \left(N_c^{\rm eff}\right)_4 \approx \left(N_c^{\rm eff}\right)_9 \approx \left(N_c^{\rm eff}\right)_{10}\,, \\ N_c^{\rm eff}(LR) &\equiv \left(N_c^{\rm eff}\right)_5 \approx \left(N_c^{\rm eff}\right)_6 \approx \left(N_c^{\rm eff}\right)_7 \approx \left(N_c^{\rm eff}\right)_8\,, \end{aligned} \tag{5}$$

and $N_c^{\rm eff}(LR) \neq N_c^{\rm eff}(LL)$ in general. Since the energy release in the energetic two-body charmless B decays is not less than that in $B \to D\pi$ decays, it is thus expected that $|\chi(\text{two} - \text{body rare B decay})| \lesssim |\chi(B \to D\pi)|$ and hence $N_c^{\rm eff}(LL) \approx N_c^{\rm eff}(B \to D\pi) \sim 2$. From the data analysis in the next section, we shall see that $N_c^{\rm eff}(LL) < 3$ and $N_c^{\rm eff}(LR) > 3$.

IV. ANALYSIS OF DATA

We have studied in detail the two-body charmless B decays for $B_{u,d}$ mesons in [4] and for B_s mesons in [10]. In what follows we show some highlights of the data analysis.

A. Spectator-dominated rare B decays

Very recently, CLEO has made the first observation of a hadronic $b \to u$ decay, namely $B^\pm \to \rho^0\pi^\pm$ [11,12]. The preliminary measurement yields:

$$\mathcal{B}(B^\pm \to \rho^0\pi^\pm) = (1.5 \pm 0.5 \pm 0.4) \times 10^{-5}. \tag{6}$$

The branching ratios of this mode decreases with $N_c^{\rm eff}(LL)$ as it involves interference between external and internal W-emission amplitudes. From Fig. 1 it is clear that this class-III decay is sensitive to $1/N_c^{\rm eff}$ if $N_c^{\rm eff}(LL)$ is treated as a free parameter, namely, $N_c^{\rm eff}(LR) = N_c^{\rm eff}(LL) = N_c^{\rm eff}$; it has the lowest value of order 1×10^{-6} and then grows with $1/N_c^{\rm eff}$. We see from Fig. 1 that $0.35 \le 1/N_c^{\rm eff} \le 0.92$. Since the tree diagrams make the dominant contributions, we then have

$$1.1 \le N_c^{\rm eff}(LL) \le 2.6 \quad \text{from } B^\pm \to \rho^0\pi^\pm. \tag{7}$$

Therefore, $N_c^{\rm eff}(LL)$ is favored to be less than 3, as expected.

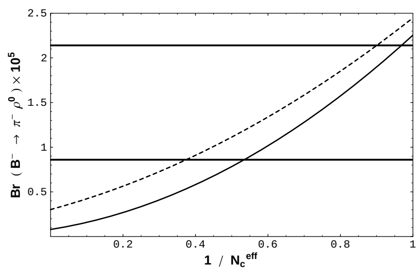

FIG. 1. The branching ratio of $B^- \to \rho^0\pi^-$ versus $1/N_c^{\rm eff}$. The solid (dotted) curve is calculated using the Baure-Stech-Wirbel (light-cone sum rule) model for form factors, while the solid thick lines are the CLEO measurements with one sigma errors.

There are two additional experimental hints that favor the choice $N_c^{\rm eff}(LL) \sim 2$. First is the class-III decay $B^\pm \to \pi^\pm\omega$. This mode is very similar to $\rho^0\pi^\pm$ as its decay amplitude differs from that of $\omega\pi^\pm$ only in the $2(a_3 + a_5)$ penguin term which is absent in the former. Since the coefficient $(a_3 + a_5)$ is small and it is further subject to the quark-mixing angle suppression, the decay rates of $\omega\pi^\pm$ and $\rho^0\pi^\pm$ are very similar. Although experimentally only the upper limit $\mathcal{B}(B^\pm \to \pi^\pm\omega) < 2.3 \times 10^{-5}$ is quoted by CLEO [13], the CLEO measurements $\mathcal{B}(B^\pm \to K^\pm\omega) = (1.5^{+0.7}_{-0.6} \pm 0.2) \times 10^{-5}$ and $\mathcal{B}(B^\pm \to h^\pm\omega) = (2.5^{+0.8}_{-0.7} \pm 0.3) \times 10^{-5}$ with $h = \pi,\ K$ indicate that the central value of $\mathcal{B}(B^\pm \to \pi^\pm\omega)$ is about 1×10^{-5}. This means $0.4 < 1/N_c^{\rm eff}(LL) < 0.6$ (see Fig. 3 of [4]) or $1.7 < N_c^{\rm eff}(LL) < 2.5$ is favored; the prediction for $N_c^{\rm eff}(LL) = 2$ is $\mathcal{B}(B^\pm \to \omega\pi^\pm) = 0.8 \times 10^{-5}$ and 1.1×10^{-5} in the Bauer-Stech-Wirbel (BSW) model [14] and the light-cone sum rule (LCSR) analysis [15] for heavy-to-light form factors, respectively. The second hint comes from the decay $B^0 \to \pi^+\pi^-$. A very recent CLEO analysis of $B^0 \to \pi^+\pi^-$ presents an improved upper limit, $\mathcal{B}(B^0 \to \pi^+\pi^-) < 0.84 \times 10^{-5}$ [16]. It also implies that $N_c^{\rm eff}(LL)$ is preferred to be smaller [4].

B. $B \to \phi K,\ \phi K^*$ decays

The decay amplitudes of the penguin-dominated modes $B \to \phi K$ and $B \to \phi K^*$ are governed by the effective coefficients $[a_3 + a_4 + a_5 - \frac{1}{2}(a_7 + a_9 + a_{10})]$. Note that the QCD penguin coefficients a_3 and a_5 are sensitive to $N_c^{\rm eff}(LL)$ and $N_c^{\rm eff}(LR)$, respectively. We see from Figs. 6 and 7 of [4] that the decay rates of $B \to \phi K^{(*)}$ increase with $1/N_c^{\rm eff}(LR)$ irrespective of $N_c^{\rm eff}(LL)$. The new CLEO upper limit $\mathcal{B}(B^\pm \to \phi K^\pm) < 0.59 \times 10^{-5}$ at 90% C.L. [12] implies that

$$N_c^{\rm eff}(LR) \geq \begin{cases} 4.2 & \text{BSW}, \\ 3.2 & \text{LCSR}, \end{cases} \tag{8}$$

with $N_c^{\rm eff}(LL)$ being fixed at the value of 2. Note that this constraint is subject to the corrections from spacelike penguin and W-annihilation contributions. At any rate, it is safe to conclude that $N_c^{\rm eff}(LR) > 3 > N_c^{\rm eff}(LL)$.

CLEO has seen a 3σ evidence for the decay $B \to \phi K^*$. Its branching ratio, the average of ϕK^{*-} and ϕK^{*0} modes, is reported to be [12]

$$\mathcal{B}(B \to \phi K^*) \equiv \frac{1}{2}\left[\mathcal{B}(B^\pm \to \phi K^{*\pm}) + \mathcal{B}(B^0 \to \phi K^{*0})\right] = \left(1.1^{+0.6}_{-0.5} \pm 0.2\right) \times 10^{-5}. \tag{9}$$

We find that the branching ratio of $B \to \phi K^*$ is in general larger (less) than that of $B \to \phi K$ in the LCSR (BSW) model. This is because $\Gamma(B \to \phi K^*)$ is very sensitive to the form factor ratio $x = A_2^{BK^*}(m_\phi^2)/A_1^{BK^*}(m_\phi^2)$, which is equal to 0.875 (1.03) in the LCSR (BSW) model. In particular, $\mathcal{B}(B \to \phi K^*) = 0.77 \times 10^{-5}$ is predicted by the LCSR for $N_c^{\rm eff}(LL) = 2$ and $N_c^{\rm eff}(LR) = 5$, which is in accordance with experiment. It is evident that the data of $B \to \phi K$ and $B \to \phi K^*$ can be simultaneously accommodated in the LCSR analysis (see Figs. 6 and 7 of [4]). Therefore, the non-observation of $B \to \phi K$ does not necessarily invalidate the factorization hypothesis; it could imply that the form-factor ratio A_2/A_1 is less than unity. Of course, it is also possible that the absence of $B \to \phi K$ events is a downward fluctuation of the experimental signal. At any rate, in order to clarify this issue and to pin down the effective number of colors $N_c^{\rm eff}(LR)$, measurements of $B \to \phi K$ and $B \to \phi K^*$ are urgently needed with sufficient accuracy.

C. $B \to \eta' K$ and ηK decays

The published CLEO results [17] on the decay $B \to \eta' K$: $\mathcal{B}(B^\pm \to \eta' K^\pm) = \left(6.5^{+1.5}_{-1.4} \pm 0.9\right) \times 10^{-5}$ are several times larger than earlier theoretical predictions [18–20] in the range of $(1-2) \times 10^{-5}$. It was pointed out in past two years by several authors [5,21] that the decay rate of $B \to \eta' K$ will get enhanced because of the small running strange quark mass at the scale m_b and sizable $SU(3)$ breaking in the decay constants f_8 and f_0.

As shown in [4], if $N_c^{\rm eff}(LL)$ is treated to be the same as $N_c^{\rm eff}(LR)$, the branching ratio of $(B^- \to \eta' K^-)$ is $\sim (2.7-4.7) \times 10^{-5}$ at $0 < 1/N_c^{\rm eff} < 0.5$ and it becomes $(3.5-3.8) \times 10^{-5}$ when the η' charm content contribution with $f_{\eta'}^c = -6.3$ MeV is taken into account. However, the discrepancy between theory and experiment is greatly improved by treating $N_c^{\rm eff}(LL)$ and $N_c^{\rm eff}(LR)$ differently. Setting $N_c^{\rm eff}(LL) = 2$, we find that (see Fig. 2) the decay rates of $B \to \eta' K$ are considerably enhanced especially at small $1/N_c^{\rm eff}(LR)$. Specifically, $\mathcal{B}(B^\pm \to \eta' K^\pm)$ at $1/N_c^{\rm eff}(LR) \leq 0.2$ is enhanced from $(3.6-3.8) \times 10^{-5}$ to $(4.6-6.1) \times 10^{-5}$ due to three enhancements. First, the η' charm content contribution $a_2 X_c^{(BK,\eta')}$ now always contributes in the right direction to the decay rate irrespective of the value of $N_c^{\rm eff}(LR)$. Second, the interference in the spectator amplitudes of $B^\pm \to \eta' K^\pm$ is constructive. Third, the term proportional to $2(a_3 - a_5)X_u^{(BK,\eta')} + (a_3 + a_4 - a_5)X_s^{(BK,\eta')}$ is enhanced when $(N_c^{\rm eff})_3 = (N_c^{\rm eff})_4 = 2$. Therefore, the data of $B \to K\eta'$ provide a strong support for $N_c^{\rm eff}(LL) \sim 2$ and the relation $N_c^{\rm eff}(LR) > N_c^{\rm eff}(LL)$. The mode $B \to \eta' K$ has the largest branching ratio in the two-body charmless B decays due mainly to the constructive interference between the penguin contributions arising from the $(\bar{u}u + \bar{d}d)$ and $\bar{s}s$ components of the η'. By contrast, the destructive interference for the η production leads to a much small decay rate for $B \to \eta K$.

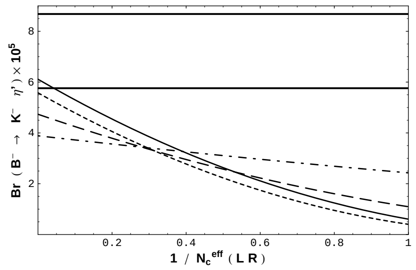

FIG. 2. The branching ratio of $B^{\pm} \to \eta' K^{\pm}$ as a function of $1/N_c^{\rm eff}(LR)$ with $N_c^{\rm eff}(LL)$ being fixed at the value of 2 and $\eta = 0.370$, $\rho = 0.175$, $m_s(m_b) = 90$ MeV. The calculation is done using the BSW model for form factors. The charm content of the η' with $f_{\eta'}^c = -6.3$ MeV contributes to the solid curve but not to the dotted curve. The anomaly contribution to $\langle \eta' | \bar{s}\gamma_5 s | 0 \rangle$ is included. For comparison, predictions for $N_c^{\rm eff}(LL) = N_c^{\rm eff}(LR)$ as depicted by the dashed curve with $f_{\eta'}^c = 0$ and dot-dashed curve with $f_{\eta'}^c = -6.3$ MeV are also shown. The solid thick lines are the preliminary updated CLEO measurements [11,12] with one sigma errors.

D. $B^{\pm} \to \omega K^{\pm}$ and $B^{\pm} \to \rho^0 K^{\pm}$ decays

The CLEO observation [13] of a large branching ratio for $B^{\pm} \to \omega K^{\pm}$, $\mathcal{B}(B^{\pm} \to \omega K^{\pm}) = \left(1.5^{+0.7}_{-0.6} \pm 0.2\right) \times 10^{-5}$, is rather difficult to explain at first sight. We showed in [4] that in the absence of FSI, the branching ratio of $B^+ \to \omega K^+$ is expected to be of the same order as $\mathcal{B}(B^+ \to \rho^0 K^+) \sim (0.5 - 1.0) \times 10^{-6}$, while experimentally it is of order 1.5×10^{-5}. We argued that $B^+ \to \omega K^+$ receives a sizable final-state rescattering contribution from the intermediate states $K^{*-}\pi^0, K^{*-}\eta, K^{*0}\pi^-, K^-\rho^0, K^0\rho^-$ which interfere constructively, whereas the analogous rescattering effect on $B^+ \to \rho^0 K^+$ is very suppressed. However, if the measured branching ratio $\rho^0 K^+$ is similar to ωK^+, then W-annihilation and spacelike penguins may play a prominent role. Likewise, the decay mode $B^0 \to K^+K^-$ is expected to be dominated by inelastic rescattering from $\rho^+\rho^-, \pi^+\pi^-$ intermediate states.

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