Factorization Approaches to $B$ Meson Decays

Hsiang-nan Li
Institute of Physics, Academia Sinica, Nankang, Taipei 115, Taiwan, Republic of China

Department of Physics, National Cheng-Kung University, Tainan 701, Taiwan, Republic of China

We compare the theoretical frameworks and the phenomenological applications of the factorization approaches to exclusive $B$ meson decays, which include QCD-improved factorization, perturbative QCD, and soft-collinear effective theory. Recent progress on two-body nonleptonic $B$ meson decays made in these approaches are reviewed.

1. Introduction

"Factorizations" in the naive factorization assumption and in factorization theorem have very different meanings. The former refers to the factorization of a process into subprocesses. For example, a $B$ meson decay amplitude $A(B \to M_1M_2)$ is written, in the factorization assumption, as the product $f_{M_2}F^{BM_1}$,

$$A(B \to M_1M_2) = f_{M_2}F^{BM_1}; \quad (1)$$

where the $m$ meson decay constant $f_{M_2}$ arises from the production of the $m$ meson $M_2$ from the vacuum, and the form factor $F^{BM_1}$ is associated with the $B \to M_1$ transition. The latter refers to the factorization of perturbative and nonperturbative dynamics in a QCD process. According to factorization theorem, the above amplitude is expressed as

$$A(B \to M_1M_2) = f_{M_2}H(x;Q^2); \quad (2)$$

where $H(x;Q^2)$ denotes the convolution over parton kinematic variables, the hard kernel $H$ absorbs perturbative dynamics, and the $B \to M_1M_2$ meson distribution amplitude $f_{M_2}$ absorbs nonperturbative dynamics in the $B \to M_1M_2$ decay. A piece of contribution to $B$ meson decays is factorizable, if it respects either Eq. (1) in the sense of the factorization assumption, or Eq. (2) in the sense of factorization theorem. Below we shall use the terms "factorizable" and "nonfactorizable" without specifying which sense it refers to.

2. Collinear vs. $k_T$ Factorization

Both collinear and $k_T$ factorizations are the fundamental tools of perturbative QCD, where $k_T$ denotes parton transverse momentum. We first explain these two types of theorems by considering the simplest scattering process $(P_1)$ ! $(P_2)$ as an example. The $m$ meson $P_1$ of the pion and the $m$ meson $P_2$ of the outgoing on-shell photon are chosen as

$$P_1 = (P_{1T}^0; 0; 0_T); \quad P_2 = (0; P_2; 0_T); \quad (3)$$

The leading-order (LO) quark diagram, in which the anti-quark $q$ carries the on-shell fraction $m$ in the on-shell pion and the internal quark carries $P_2 \to k_T$, leads to the amplitude,

$$G^{(0)}(x;Q^2) = \frac{1}{xQ^2} \left\{ \frac{\text{tr}}{\text{tr}}[6 \text{ P}_2 \text{ P}_1 \text{ s}] \right\}; \quad (4)$$

with the leading spin structure $\text{P}_1 \cdot \text{P}_s$ of the pion and the $m$ mesonium transfer squared $Q^2 = 2P_2 \cdot \text{P}_1$. We have suppressed other constant factors, such as the electric charge, the color number, and the pion decay constant, which are irrelevant in the following discussion.

The trivial collinear factorization of Eq. (4) reads,

$$G^{(0)}(x;Q^2) = \frac{1}{xQ^2} \left\{ \frac{\text{tr}}{\text{tr}}[6 \text{ P}_2 \text{ P}_1 \text{ s}] \right\}; \quad (5)$$

The zeroth-order distribution amplitude $G^{(0)}$ is proportional to $(x\cdot k_T^2)$, implying that the parton entering the LO hard kernel $H^{(0)}$ carries the same momentum as the parton entering the distribution amplitude does. The trivial $k_T$ factorization of Eq. (4) reads,

$$G^{(0)}(x;Q^2) = \frac{1}{xQ^2} \left\{ \frac{\text{tr}}{\text{tr}}[6 \text{ P}_2 \text{ P}_1 \text{ s}] \right\}; \quad (6)$$

Because of the zeroth-order wave function $(0)$ / $(k_T^2)$, $H^{(0)}$ does not depend on the parton transverse momentum actually.

The $O(\alpha_s)$ quark diagrams for Eq. (4) from full QCD are displayed in Fig. 1, in which the upper line represents the quark. The collinear factorization of...
these radiative corrections is given by

\[
G^{(1)}(x; Q^2) = \int x dx \, G^{(0)}(x; Q^2) H^{(1)}(x; Q^2) + H^{(1)}(x; Q^2);
\]

where the first-order distribution amplitude \((1)\) is defined by the effective diagrams in Fig. 3. Expressions from Figs. 2(c), 2(e), and 2(f) are proportional to \((x' x I = P_1^1)\), while the second-order term is proportional to \((x' x' I = P_1^1)\). In this case the collinear gluon exchange modifies both the longitudinal and transverse parton momentum fractions entering into the hard kernel.

\[
G^{(1)}(x; Q^2) = \int x dx k_T^{(1)}(x; Q^2) k_T^{(0)} + H^{(1)}(x; Q^2);
\]

where the first-order wave function \((1)\) from Figs. 2(c), 2(e), and 2(f) is proportional to \((x' x' I = P_1^1)\). In this case the collinear gluon exchange modifies both the longitudinal and transverse parton momentum fractions entering into the hard kernel.

It is observed that \(H^{(0)}\) in Eq. (5) depends on \(k_T\) nontrivially in the first-order \(k_T\) factorization. Being convoluted with \(H^{(0)}\), the partons entering the next-to-leading (NLO) hard kernel \(H^{(1)}\) are still on-shell. To acquire a nontrivial \(k_T\) dependence, \(H^{(1)}\) must be convoluted with the higher-order wave functions \((1)\) and \((1)\); the gluon exchanges in \((1)\) render the incoming partons of \(H^{(1)}\), i.e., the incoming partons of the quark diagram \(s^{(1)}\) and the effective diagrams \(s^{(1)}\) on-shell by \(k_T^2\). We thus derive \(H^{(1)}(x; Q^2; j k_T)\) according to the formula,

\[
H^{(1)}(x; Q^2; j k_T) = G^{(1)}(x; Q^2; j k_T)
\]

with the quark carrying the momentum \(k = (P_1^1; P_1^1; j k_T)\). This is the way to obtain a \(k_T\)-dependent hard kernel without breaking gauge invariance, since the gauge dependences cancel between \(G^{(1)}\) and \(H^{(1)}\). A physical quantity is written as a convolution of a hard kernel with the collinear wave function, which are determined by methods beyond a perturbation theory, such as lattice QCD and QCD sum rules, or extracted from experimental data. A gauge-invariant hard kernel then leads to gauge-invariant predictions from \(k_T\) factorization.

3. B Decays in QCDF, SCET, and PQCD

Factorization theorems have been applied to exclusive B meson decays, and different approaches have been developed. Below we compare the frameworks of perturbative QCD (PQCD) [1, 2, 3, 4], QCDF [5, 6, 7, 8], and soft-collinear effective theory (SCET) [9, 10]. The factorization approach is based on the semileptonic decay \(B \rightarrow \ell \nu + X\), where \(\ell\) is expressed in collinear factorization, as

\[
F_B = \int x_1 dx_2 \, (x_1 Q x_2) (x_2); \quad (10)
\]

with the LO hard kernel \(H^{(0)}\). The parton momentum fractions \(x_1\) and \(x_2\) are carried by the spectator quarks on the B meson and pion sides, respectively. Obviously, the above integral is logarithmically divergent for the asymptotic model.

An end-point singularity implies that exclusive B meson decays are dominated by soft dynamics. That is, a heavy-to-light form factor is not calculable in collinear factorization, and \(F_B\) should be treated as a soft object [11, 12]. This is the basic idea of QCDF, and subleading corrections are added systematically [12]. The above treatment has been further elucidated in the framework of SCET [13]; only the 1 term in \(H^{(0)}\) contains the end-point singularity, which leads to an \(O(1)\) factorization. \(F_B\) is calculated by the B meson...
mass $m_B$. Therefore, at leading power in $1-m_B$, the $B \to two-body$ nonleptonic form factor can be split into the nonfactorizable and factorizable components,

$$F^B = F^{NF} + F^F;$$

(11)

which have different power counting in the strong coupling constant $\alpha$: the former is of $O(\alpha)$ and the latter of $O(\alpha^2)$. The values of $F^{NF}$ and $F^F$ have been determined from $a$ to the $B$ data [14].

The formulation of the $B$ transition in $K$-factorization theorem is different. When the parton transverse momenta are included, $F^{NF}$ does not develop an end-point singularity, and both $F^{NF}$ and $F^F$ are factorizable. Hence, they are of the same order in $\alpha$ and can be combined into a single term, giving $\text{SCET} [15].$

$$Z = \int dx_1 dx_2 d^2 k_{1T} d^2 k_{2T} \int (x_1, k_{1T}) H (x_2, k_{1T}, k_{2T}) (x_2, k_{2T});$$

(12)

The end-point singularity is smeared into the large logarithm $\ln^2 (m_B - k_0^2)$, and absorbed into the pion wave function. Repeating this logarithm at all orders in the conjugate space $[14,17]$, we derived the Sudakov factor $S(m_B, b)$, which describes the parton distribution in $b$. Since $F^{NF}$ has been included, the large-energy symmetry is respected in $PQCD$. Recently, it was proposed that the end-point singularity is attributed to a double counting of soft degrees of freedom in collinear factorization [13]. The zero-bin subtraction removes the double counting, and leads to a modified SCET formalism for $F^{NF}$, labelled by SCET$_6$ hereafter. The power counting of SCET$_6$ in both $1-m_B$ and $s$ is then consistent with that of PQCD. The regularization of the end-point singularity introduces the logarithm $\ln^2 (m_B, b)$, whose treatment has not yet been explained.

When applying the above factorization approaches to two-body nonleptonic $B$ meson decays, further difference appears in the treatment of end-point singularities. The $O(\alpha)$ $m_B=m_B$ annihilation amplitudes from the penguin operators are divergent in collinear factorization, where $m_B$ is the chiral enhancement scale. Because of the end-point singularity, an annihilation amplitude has been parametrized as

$$m_B \ln 1 + A e^{iA};$$

(13)

in $\text{QCDF} [8]$, where $A$ is a hadronic scale and the free parameter $\Lambda$ is postulated to vary in the range $0 < \Lambda < 1$. It is not clear what mechanism generates the strong phase $\Lambda$. With the similar zero-bin subtraction, an annihilation amplitude is factorizable in SCET$_6$, but found to be absent in $\text{QCDF} [13]$. A strong phase can only be generated at loop level, e.g., at $O(\alpha^2)$. However, we notice that the residual

The scalar penguin annihilation amplitude also factorizable in $K$-factorization with the absence of the end-point singularity. Furthermore, it was almost in agreement in $PQCD [3]$, whose corresponding mechanism is similar to the Bander-Silveu-V-Gonion [21]; when the $u$ or $c$ quark in a loop goes on mass shell, a strong phase is produced. In the case of the annihilation topology for heavy-to-light decays, the loop is formed by the virtual particles in the LO PQCD hard kernel and the in nearly any Sudakov gluons exchanged between two partons in a light meson. The virtual particle acquires the transverse momenta of $K_1$ through the Sudakov gluon exchange. A sizable strong phase is then given, in terms of the principle-value prescription for the virtual particle propagator, by

$$\frac{1}{\beta^2 + i} = \frac{P}{(\beta^2 + k_1^2)^2 + i} (\beta^4 + k_1^4)$$

(14)

Therefore, the treatment and the effect of the scalar penguin annihilation amplitude are very different in $\text{QCDF, SCET}_6$, and $PQCD$.

Though the scalar penguin annihilation amplitude is factorizable in both $PQCD$ and $\text{SCET}_6$, it is almost in agreement in $\text{QCDF}$, but real in the latter. We argue that the above different opinions can be discriminated by comparing the direct CP asymmetries in the charged $B$ meson decays $B \to K^0$ and $B \to K^0$ decays. The $B \to K^0$ decays involve a $B$ meson transition to a pseudoscalar meson, so the penguin emission amplitude is proportional to the constructive combination of the W boson couplings $a_4 + 2(m_B^0 - m_B) a_5$, where $m_B^0$ is the chiral enhancement scale associated with the kaon. The $B \to K^0$ decays involve a $B$ meson transition to a vector meson, so the penguin emission amplitude is proportional to the destructive combination $a_4 - 2(m_B^0 - m_B) a_5$. The annihilation effect is then less in essential in the former than in the latter. If the scalar penguin annihilation is real, both decays will exhibit small direct CP asymmetries, e.g., $A_{CP} (B \to K^0)$, $A_{CP} (B \to K^0)$, $A_{CP} (B \to K^0)$, the current data $A_{CP} (B \to K^0)$, $A_{CP} (B \to K^0)$, and $A_{CP} (B \to K^0)$ seem to prefer an inaginary scalar penguin annihilation.
4. Recent Results

4.1. The B ! Puzzle

According to a naive estimate of the color-suppressed tree amplitude, the hierarchy of
the branching ratios \( B(B^0 \rightarrow 0 0) \) is expected. However, the data [24]
\[
B(B^0 \rightarrow 0 0) = (5.2 \pm 0.2) \times 10^{-6};
B(B^0 \rightarrow 0 0) = (1.31 \pm 0.21) \times 10^{-6}; (15)
\]
show \( B(B^0 \rightarrow 0 0) \), giving rise to the B ! puzzle. As observed in
[23], the NLO corrections, despite of increasing the color-suppressed tree amplitudes signifi-
cantly, are not enough to enhance the \( B(B^0 \rightarrow 0 0) \) branching ratio to the measured value. A much larger color-suppressed tree amplitude, about the same order as the color-suppressed tree amplitude, must be obtained in order to resolve the puzzle [24,25]. To make sure that the above NLO effects are reasonable, the PQCD formalism has been applied to the B ! decays [23], which also resolve the color-suppressed tree contribution. It was found that the NLO PQCD predictions are in agreement with the data \( B(B^0 \rightarrow 0 0) = (1.31 \pm 0.21) \times 10^{-6} [24] \). We conclude that it is unlikely to accommodate the measured \( B(B^0 \rightarrow 0 0) \) and \( B(B^0 \rightarrow 0 0) \) branching ratios simultaneously in PQCD, and that the B ! puzzle remains.

It has been claimed that the B ! puzzle is resolved in the QCD approach [8] with an input from SCET [26,27,28]; the inclusion of the NLO jet function, the hard core of SCET II, into the QCD formula for the color-suppressed tree amplitude gives a soft enhancement of the \( B(B^0 \rightarrow 0 0) \) branching ratio, if adopting the parametrization scenario "S4" [26]. It is necessary to investigate whether the proposed new mechanism deteriorates the consistency of theoretical results with other data. The formalism in [26] has been extended to the B ! decays as a check [23]. It was found that the NLO jet function overshoots the observed \( B(B^0 \rightarrow 0 0) \) branching ratio very much as adopting "S4". That is, it is also unlikely to accommodate the B ! data simultaneously in QCD.

4.2. The B ! K Puzzle

For penguin-dominated B ! K decays, as shown in Table I [22], the polarization fractions deviate from the naive counting rules based on kinematics [39]. This is the so-called the B ! K puzzle. Many attempts to resolve the B ! K polarizations have been made [31], which include new physics [32,33,34,35,36], the annihilation contribution [37,38] in the QCD approach, the charm penguin in SCET [14], the rescattering effects [33,40,41], and the b ! s (the magnetic penguin) [42] and b ! s [43] transitions. The annihilation contribution from the B ! K decays is expected to be a large uncertainty.

| Mode | Belle | BaBar |
|------|------|-------|
| K *  | 0.52 | 0.01  | 0.02 |
| K 0  | 0.05 | 0.05  | 0.02 |

4.3. The B ! K Puzzle

The B ! K decays depend on the tree amplitude, the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations.

Table I Polarization fractions in the penguin-dominated B ! K decays.

The B ! K decays depend on the tree amplitude, the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations. The data of the direct CP asymmetry \( A_{CP} (B^0 \rightarrow K^0) \) depend on the QCD penguin and the B ! K polarizations.
expected. However, this naive expectation is in conflict with the data.

\[ A_{CP}(B^0 \to K^0) = 0.093 \pm 0.015 \]
\[ A_{CP}(B^+ \to K^-) = 0.047 \pm 0.026 \]  \hspace{1cm} (16)

leading to one of the B \to K puzzles.

While LO PQCD gives a negligible C^0 \cite{3,7}, it is possible that this supposedly tiny amplitude receives a significant subleading correction. Note that the small C^0 is attributed to the accidental cancellation between the Wilson coefficients C_1 and C_2 = \pi, at the scale of the b quark mass m_b. In \cite{9} the important NLO contributions to the B \to K decays from the vertex corrections, the quark loops, and the magnetic penguins were calculated. It was observed that the vertex corrections increase C^0 by a factor of 3, and induce a large phase about 80° relative to T^0. The large and imaginary C^0 then renders the total tree amplitude T^0 + C^0 more or less parallel to the total penguin amplitude B^0 + P^0 in the B \to K decays, leading to nearly vanishing \( A_{CP}(B^0 \to K^0) = (1/3)\% \) at NLO (it is about -8% at LO). We conclude that the B \to K puzzle has been alleviated, but not yet gone away completely. Whether new physics effects \cite{4,43} are called for will become clear when the data are more detailed. More discussion on this subject can be found in \cite{50}.

4.4. Nonleptonic B_s Decays

Two-body nonleptonic B_s meson decays are interesting, since their study can test SU(3) or U-spin symmetry. The framework for these decays is basically identical to that for B_d meson decays. The results of two-body nonleptonic B_s meson decays from different factorization approaches can be found in \cite{23} for QCDF, in \cite{51} for SCET, and in \cite{23,52} for PQCD. Roughly speaking, the branching ratios predicted by QCDF and by PQCD are similar, but the predicted direct CP asymmetries are usually opposite in sign.

5. Conclusion

The factorization approaches are systematic theoretical tools for exclusive B meson decays, in which hadronic inputs are universal, and the hard kernels can be computed order by order. NLO corrections have been obtained for some B meson decay modes, and the consistency between the theoretical predictions and the experimental data is improved in general. More need to be done in order to pin down QCD uncertainty especially for those quantities exhibiting puzzling behaviors. Higher-power corrections are another important source of QCD uncertainty, which deserves a careful investigation. The recent development in SCET is encouraging, whose counting rules become consistent with those in PQCD. However, the arbitrary logarithm sinh resulting from the zero-bin subtraction needs to be handled (recall that the double logarithm ln^2(m_b = k_f) from the smearing of the end-point singularity has been resummed in PQCD). We did not review the progress on the S puzzle appearing in the extraction of the weak phase sin(2\theta_1) from penguin-dominated B meson decays. For the details, refer to \cite{54}.

Acknowledgments

The work was supported in part by the National Science Council of R.O.C. under Grant No. NSC-95-2112-M-050-M3, and by the National Center for Theoretical Sciences of R.O.C. .

References

[1] M. Bauer, B. Stech, and M. W. M"{u}zel, Z. Phys. C 29, 637 (1985); 34, 103 (1987).
[2] M. Nagashima and H.-n. Li, Phys. Rev. D 67, 034001 (2003).
[3] H.-n. Li, Phys. Rev. D 64, 014019 (2001); M. Nagashima and H.-n. Li, Eur. Phys. J. C 40, 395 (2005).
[4] S. Nandi and H.-n. Li, arXiv:0704.3790, to appear in Phys. Rev. D.
[5] H.-n. Li and H. L. Yu, Phys. Rev. Lett. 74, 4388 (1995); Phys. Lett. B 353, 301 (1995); Phys. Rev. D 53, 2480 (1996).
[6] Y. K. Keum, H.-n. Li, and A. I. Sanda, Phys. Lett. B 504, 6 (2001); Phys. Rev. D 63, 054008 (2001); Y. K. Keum and H.-n. Li, Phys. Rev. D 63, 074006 (2001).
[7] C. D. Lu, K. Ukai, and M. Z. Yang, Phys. Rev. D 63, 074009 (2001).
[8] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Phys. Rev. Lett. 83, 1914 (1999); Nucl. Phys. B 591, 313 (2000).
[9] C. W. Bauer, S. Fleming, and M. Luke, Phys. Rev. D 63, 014006 (2001).
[10] C. W. Bauer, S. Fleming, D. Pirjol, and I. Stewart, Phys. Rev. D 63, 114020 (2001).
[11] A. Szczepaniak, E. M. Henley, and S. Brodsky, Phys. Lett. B 243, 287 (1990).
[12] M. Beneke and T. Felkann, Nucl. Phys. B 592, 3 (2000).
[13] C. W. Bauer, D. Pirjol, and I. Stewart, Phys. Rev. D 67, 071502 (2003); R. J. Ilis, T. Becker, S. J. Lee, and M. Neubert, JHEP 0407, 081 (2004).
[14] C. W. Bauer, D. Pirjol, I.Z. Rothstein, and I. Stewart, Phys. Rev. D 70, 054015 (2004).

fpocp07.323
[15] T. Kurihoto, H.-n. Li, and A.I. Sanda, Phys. Rev. D 65, 014007 (2002).
[16] J. Botts and G. Sterman, Nucl. Phys. B 325, 62 (1989).
[17] H.-n. Li and G. Sterman, Nucl. Phys. B 381, 129 (1992).
[18] A.V. Manohar and I.W. Stewart, hep-ph/0605001.
[19] C.M. Amsler, Z. Ligeti, I.Z. Rothstein, I.W. Stewart, hep-ph/0607001.
[20] J. Chay, H.-n. Li, and S.M. Ishma, in preparation.
[21] M. Bander, D. Silverman and A. Soni, Phys. Rev. Lett. 43, 242 (1979).
[22] Heavy Flavor Averaging Group, http://www.slac.stanford.edu/xorg/hfag/.
[23] H.-n. Li and S.M. Ishma, Phys. Rev. D 73, 114014 (2006).
[24] Y.Y. Chang and H.-n. Li, Phys. Rev. D 71, 014036 (2005).
[25] T.N. Pham, hep-ph/0610063.
[26] M. Beneke and D. Yang, Nucl. Phys. B 736, 34 (2006).
[27] M. Beneke and S. Jager, hep-ph/0512101.
[28] M. Beneke and S. Jager, Nucl. Phys. B 751, 160 (2006).
[29] M. Beneke and M. Neubert, Nucl. Phys. B 675, 333 (2003).
[30] C.H. Chen, Y.Y. Keum, and H.-n. Li, Phys. Rev. D 66, 054013 (2002).
[31] H.-n. Li and S.M. Ishma, Phys. Rev. D 71, 054025 (2005).
[32] Y. Grossman, Int. J. Mod. Phys. A 19, 907 (2004).
[33] Y.D. Yang, R.M. Wang, and G.R. Lu, Phys. Rev. D 72, 015009 (2005).
[34] P.K. Das and K.C. Yang, Phys. Rev. D 71, 094002 (2005).
[35] C.H. Chen and C.Q. Geng, Phys. Rev. D 71, 115004 (2005).
[36] C.S. Huang, P. Ko, X.H. Wu, and Y.D. Yang, Phys. Rev. D 73, 034026 (2006).
[37] A.L. Kagan, Phys. Lett. B 601, 151 (2004); hep-ph/0407078.
[38] M. Beneke, J. Rohrer, and D. Yang, hep-ph/0612290.
[39] F. Colangelo, F. De Fazio, and T.N. Pham, Phys. Lett. B 597, 291 (2004).
[40] M. Ladisa, V. Laporta, G.Nardulli, and P. Santorelli, Phys. Rev. D 70, 114025 (2004).
[41] H.Y. Cheng, C.K. Chua, and A. Soni, Phys. Rev. D 71, 014030 (2005).
[42] W.S. Hou and M. Nagashima, hep-ph/0408007.
[43] M. Beneke, J. Rohrer, and D. Yang, Phys. Rev. Lett. 96, 141801 (2006).
[44] The BABAR Collaboration, B. Aubert, et. al., Phys. Rev. Lett. 97, 201801 (2006).
[45] H.-n. Li, Phys. Lett. B 622, 63 (2005).
[46] A. Datta et al., arXiv:0705.3913.
[47] H.-n. Li, S.M. Ishma, and A.I. Sanda, Phys. Rev. D 72, 114005 (2005).
[48] A.J. Buras, R. Fleischer, S. Recksiegel, and F. Schwab, Eur. Phys. J. C 45, 701 (2006); R. Fleischer, S. Recksiegel, and F. Schwab, hep-ph/0702273.
[49] S. Baek and D. London, arXiv:hep-ph/0701181.
[50] M. Gronau, arXiv:0706.2156, talk presented at Flavor Physics and CP Violation Conference, Bled, 2007.
[51] A.R.W. Williams and J. Zupan, Phys. Rev. D 74, 014003 (2006).
[52] Z.J. Xiao, X.F. Chen, and D.Q. Guo, hep-ph/0608222; X.F. Chen, D.Q. Guo, and Z.J. Xiao, hep-ph/0701146.
[53] A. Ali et al., hep-ph/0703162.
[54] J. Zupan, arXiv:0707.3323, talk presented at Flavor Physics and CP Violation Conference, Bled, 2007.