Direct CP Violation and Final State Interactions in Hadronic $B$ Decays

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Final-state rescattering effects on the hadronic $B$ decays and their impact on direct CP violation are examined. The phenomenology of the polarization anomaly in $B \to \phi K^*$ decays is discussed.

1 Introduction

Although mixing-induced CP violation in $B$ decays has been observed in the golden mode $B^0 \to J/\psi K_S$ for several years, a first confirmed observation of direct CP asymmetry was not established until recently in the charmless $B$ decays $B^0(B^0) \to K^\pm \pi^\pm$ by both BaBar \cite{ref1} and Belle \cite{ref2}. Also the combined BaBar and Belle measurements of $B^0 \to \rho^\pm \pi^\mp$ imply a 3.6$\sigma$ direct CP asymmetry in the $\rho^+\pi^-$ mode \cite{ref3}. As for direct CP violation in $B^0 \to \pi^+\pi^-$, a 5.2$\sigma$ effect was claimed by Belle \cite{ref4}, but it has not been confirmed by BaBar \cite{ref5}.

Table 1 shows comparison of the model predictions of direct CP asymmetries with the world averages of experimental results \cite{ref3}. It appears that QCD factorization predictions \cite{ref6, ref7} for direct CP violation seem not consistent with experiment, whereas pQCD results \cite{ref8} are in the right ballpark. Since the observation of direct CP violation requires at least two different contributing amplitudes with distinct strong and weak phases, this means that one needs to explore the final-state interaction (FSI) rescattering phases seriously which are unlikely to be small possibly causing large compound CP-violating partial rate asymmetries in aforementioned charmless decay modes.

Besides the above-mentioned CP violation, there exist several other hints at large FSI effects in the $B$ sector \cite{ref9}. For example, the measured branching ratio $B(B^0 \to \pi^0\pi^0) = (1.5\pm0.3) \times 10^{-6}$ \cite{ref3} cannot be explained by either QCDF or pQCD and this may call for a possible rescattering effect to induce $\pi^0\pi^0$. The QCDF predictions for penguin-dominated modes such as $B \to K^*\pi$, $K\rho$, $K\phi$, $K^*\phi$ are consistently lower than the data by a factor of 2 to 3 \cite{ref7}. This large discrepancy between theory and experiment indicates the importance of subleading power corrections such as the annihilation topology and/or FSI effects.

The QCD factorization approach provides a systematic study of radiative corrections to naive factorization. While QCDF results in the limit $m_b \to \infty$ are model independent, power corrections always involve endpoint divergences. For example, the $1/m_b$ annihilation amplitude has endpoint divergences even at twist-2 level and the hard spectator scattering diagram at twist-3 order is power suppressed and possesses soft and collinear divergences arising from the
Table 1: Comparison of pQCD and QCD factorization (QCDF) predictions of direct CP asymmetries (in %) with experiment. Also shown are the QCDF results including large weak annihilation contributions denoted by QCDF(S4) and the FSI modifications to QCDF predictions [9] (see the main text for details). The pQCD results for $\rho\pi$ modes are taken from [10].

| Modes                  | Expt. | pQCD  | QCDF  | QCDF(S4)  | QCDF+FSI |
|------------------------|-------|-------|-------|-----------|----------|
| $\bar{B}^0 \to K^-\pi^+$ | $-11 \pm 2$ | $-17 \pm 5$ | $4.5^{+9.1}_{-9.9}$ | $-4.1$ | $-14^{+1}_{-3}$ |
| $\bar{B}^0 \to \rho^+\pi^-$ | $-47^{+13}_{-14}$ | $-7.1^{+0.1}_{-0.2}$ | $0.6^{+11.6}_{-11.8}$ | $-12.9$ | $-43 \pm 11$ |
| $\bar{B}^0 \to \pi^+\pi^-$ | $37 \pm 24$ | $23 \pm 7$ | $-6.5^{+13.7}_{-13.3}$ | $10.3$ | $64^{+3}_{-8}$ |
| $\bar{B}^0 \to \pi^0\pi^0$ | $28 \pm 39$ | $30 \pm 10$ | $45^{+52}_{-66}$ | $-19$ | $-30^{+1}_{-4}$ |
| $\bar{B}^0 \to \rho^-\pi^+$ | $-15 \pm 9$ | $12 \pm 2$ | $-1.5^{+8.6}_{-8.5}$ | $3.9$ | $-24 \pm 6$ |

soft spectator quark. Since the treatment of endpoint divergences is model dependent, subleading power corrections generally can be studied only in a phenomenological way. While the endpoint divergence is regulated in the pQCD approach by introducing the parton’s transverse momentum [8], it is parameterized in QCD factorization as

$$X_A \equiv \int_0^1 \frac{dy}{y} \ln \frac{m_B}{\Lambda_h} (1 + \rho_A e^{i\phi_A}),$$

with $\Lambda_h$ being a typical scale of order 500 MeV.

Just like the pQCD approach where the annihilation topology plays an essential role for producing sizable strong phases and for explaining the penguin-dominated $VP$ modes, it has been suggested in [7] that a favorable scenario (denoted as S4) for accommodating the observed penguin-dominated $B \to PV$ decays and the measured sign of direct CP asymmetry in $\bar{B}^0 \to K^-\pi^+$ is to have a large annihilation contribution by choosing $\rho_A = 1$, $\phi_A = -55^\circ$ for $PP$, $\phi_A = -20^\circ$ for $PV$ and $\phi_A = -70^\circ$ for $VP$ modes. The resultant direct CP asymmetries are shown in Table I. The sign of $\phi_A$ is chosen so that the direct CP violation $A_{K^-\pi^+}$ agrees with the data. However, the origin of these phases is unknown and their signs are not predicted. Moreover, the annihilation topologies do not help enhance the $\pi^0\pi^0$ and $\rho^0\pi^0$ modes. As stressed in [7], one would wish to have an explanation of the data without invoking weak annihilation. Therefore, it is of great importance to study final-state rescattering effects on decay rates and CP violation.

## 2 Final State Interactions in Hadronic $B$ decays

In QCDF there are two hard strong phases: one from the absorptive part of the penguin graph in $b \to s(d)$ transitions [11] and the other from the vertex corrections. However, these perturbative strong phases do not lead to the correct sign of direct CP asymmetries observed in $K^-\pi^+$, $\rho^+\pi^-$ and $\pi^+\pi^-$ modes. Therefore, one has to consider the nonperturbative strong phases induced from power suppressed contributions such as FSIs. Based on the Regge approach, Donoghue et
al. [12] have reached the interesting conclusion that FSIs do not disappear even in the heavy quark limit and soft FSI phases are dominated by inelastic scattering, contrary to the common wisdom. However, it was later pointed out by Beneke et al. [6] within the framework of QCD factorization that the above conclusion holds only for individual rescattering amplitudes. When summing over all possible intermediate states, there exist systematic cancellations in the heavy quark limit so that the strong phases must vanish in the limit of $m_b \to \infty$. Consequently, the FSI phase is generally of order $\mathcal{O}(\alpha_s, \Lambda_{QCD}/m_b)$. In reality, because the $b$ quark mass is not very large and far from being infinity, the aforementioned cancellation may not occur or may not be very effective for the finite $B$ mass. Hence, the strong phase arising from power corrections can be in principle very sizable.

At the quark level, final-state rescattering can occur through quark exchange and quark annihilation. In practice, it is extremely difficult to calculate the FSI effects, but it may become amenable at the hadron level where FSIs manifest as the rescattering processes with $s$-channel resonances and one particle exchange in the $t$-channel. In contrast to $D$ decays, the $s$-channel resonant FSIs in $B$ decays is expected to be suppressed relative to the rescattering effect arising from quark exchange owing to the lack of the existence of resonances at energies close to the $B$ meson mass. Therefore, we will model FSIs as rescattering processes of some intermediate two-body states with one particle exchange in the $t$-channel and compute the absorptive part via the optical theorem [9].

The approach of modelling FSIs as soft rescattering processes of intermediate two-body states has been criticized on several grounds [6]. For example, there are many more intermediate multi-body channels in $B$ decays and systematic cancellations among them are predicted to occur in the heavy quark limit. This effect of cancellation will be missed if only a few intermediate states are taken into account. As mentioned before, the cancellation may not occur or may not be very effective as the $B$ meson is not infinitely heavy. Hence, we may assume that two-body $\leftrightarrow n$-body rescatterings are negligible either justified from the $1/N_c$ argument or suppressed by large cancellations. Indeed, it has been even conjectured that the absorptive part of long-distance rescattering is dominated by two-body intermediate states, while the dispersive part is governed by multi-body states [13]. At any rate, we view our treatment of the two-body hadronic model for FSIs as a working tool. We work out the consequences of this tool to see if it is empirically working. If it turns out to be successful, then it will imply the possible dominance of intermediate two-body contributions.

The calculations of hadronic diagrams for FSIs involve many theoretical uncertainties. Since the particle exchanged in the $t$ channel is off shell and since final state particles are hard, form factors or cutoffs must be introduced to the strong vertices to render the calculation meaningful in perturbation theory. We shall parametrize the off-shell effect of the exchanged particle as

$$F(t, m) = \left( \frac{\Lambda^2 - m^2}{\Lambda^2 - t} \right)^n ,$$

normalized to unity at $t = m^2$ with $m$ being the mass of the exchanged particle. The monopole behavior of the form factor (i.e. $n = 1$) is preferred as it is consistent with the QCD sum rule expectation [14]. For the cutoff $\Lambda$, it should be not far from the physical mass of the exchanged particle. To be specific, we write $\Lambda = m_{\text{exc}} + r \Lambda_{QCD}$ where the parameter $r$ is expected to be of order unity and it depends not only on the exchanged particle but also on the external particles.
involved in the strong-interaction vertex. As we do not have first-principles calculations for form factors, we shall use the measured decay rates to fix the unknown cutoff parameters and then use them to predict direct CP violation. We discuss some applications below.

2.1 Penguin dominated modes

Penguin dominated modes such as $B \to K\pi$, $K^*\pi$, $K\rho$, $\phi K^{(*)}$ receive sizable contributions from rescattering of charm intermediate states (i.e. the so-called long-distance charming penguins). For example, the branching ratios of $B \to \phi K$ and $\phi K^*$ can be enhanced from $\sim 5 \times 10^{-6}$ predicted by QCDF to the level of $1 \times 10^{-5}$ by FSIs via rescattering of charm intermediate states [9].

2.2 Tree dominated modes

$B \to D\pi$ decays

The color-suppressed modes $\bar{B} \to D^{(*)}\pi^0$ have been measured by Belle, CLEO and BaBar. Their branching ratios are all significantly larger than theoretical expectations based on naive factorization. When combined with the color-allowed $\bar{B} \to D^{(*)}\pi$ decays, it indicates non-vanishing relative strong phases among various $B \to D^{(*)}\pi$ decay amplitudes. Neglecting the $W$-exchange contribution, a direct fit to the $D\pi$ data requires that $a_2/a_1 \approx (0.45 - 0.65)e^{\pm 0.06^o}$ [15]. The question is then why the magnitude and phase of $a_2/a_1$ are so different from the model expectation. We found that the rescattering from $B \to \{D\pi, D^*\rho\} \to D\pi$ contributes to the color-suppressed $W$-emission and $W$-exchange topologies and accounts for the observed enhancement of the $D^0\pi^0$ mode without arbitrarily assigning the ratio of $a_2/a_1$ a large magnitude and strong phase as done in many previous works. Note that the color-allowed $B \to D\pi$ decays are almost not affected by final-state rescattering.

$B \to \rho\pi$ decays

The color-suppressed $\rho^0\pi^0$ mode is slightly enhanced by rescattering effects to the order of $1.3 \times 10^{-6}$, which is consistent with the weighted average $(1.9 \pm 1.2) \times 10^{-6}$ of the experimental values. However, it is important to clarify the discrepancy between BaBar and Belle measurements for this mode, namely, $(1.4 \pm 0.7) \times 10^{-6}$ [16] vs $(5.1 \pm 1.8) \times 10^{-6}$ [17]. Note that the branching ratio of $\rho^0\pi^0$ is predicted to be of order $0.2 \times 10^{-6}$ in the pQCD approach [18], which is too small compared to experiment as the annihilation contribution does not help enhance its rate.

$B \to \pi\pi$ decays

There are some subtleties in describing $B^0 \to \pi\pi$ decays. The rescattering charming penguins in $\pi\pi$ are suppressed relative to that in $K\pi$ modes as the former are Cabibbo suppressed. Consequently, charming penguins are not adequate to explain the $\pi\pi$ data: the predicted $\pi^+\pi^- (\sim 9 \times 10^{-6})$ is too large whereas $\pi^0\pi^0 (\sim 0.4 \times 10^{-6})$ is too small. This means that a dispersive long-distance contribution is needed to interfere destructively with $\pi^+\pi^-$ so that $\pi^+\pi^-$ will be suppressed while $\pi^0\pi^0$ will get enhanced. This contribution cannot arise from the charming penguins or otherwise it will also contribute to $K\pi$ significantly and destroy all the nice predictions for $K\pi$. In the topological diagrammatic approach [19], this dispersive term comes from the so-called vertical $W$-loop diagram $V$ in which the meson annihilation topology such as $D^+D^- \to \pi\pi$ occurs.

It is often claimed in the literature that one needs large $C/T$ to accommodate the $\pi\pi$ data,
Table 2: Experimental data for CP averaged branching ratios (in units of $10^{-6}$) and polarization fractions for $B \to \phi K^*$ and $\rho K^*$ [20] [21].

| Mode       | BaBar            | Belle            | Average          |
|------------|------------------|------------------|------------------|
| $f_L(\phi K^{*0})$ | $0.52 \pm 0.05 \pm 0.02$ | $0.52 \pm 0.07 \pm 0.05$ | $0.52 \pm 0.05$ |
| $f_L(\phi K^{*0})$ | $0.22 \pm 0.05 \pm 0.02$ | $0.30 \pm 0.07 \pm 0.03$ | $0.25 \pm 0.04$ |
| $f_L(\phi K^{*+})$ | $0.46 \pm 0.12 \pm 0.03$ | $0.49 \pm 0.13 \pm 0.05$ | $0.47 \pm 0.09$ |
| $f_L(\phi K^{*+})$ | $0.12^{+0.11}_{-0.08} \pm 0.03$ | $0.12^{+0.11}_{-0.09}$ |
| $f_L(\rho^0 K^{*+})$ | $0.96^{+0.04}_{-0.15} \pm 0.04$ | $0.96^{+0.05}_{-0.16}$ |
| $f_L(\rho^+ K^{*0})$ | $0.79 \pm 0.08 \pm 0.04 \pm 0.02$ | $0.50 \pm 0.19^{+0.05}_{-0.07}$ | $0.74 \pm 0.08$ |

where $C$ and $T$ are color-allowed and color-suppressed topologies, respectively. In doing the fit to the $\pi\pi$ data, one actually redefines the quantities $T, C, P$ to absorb other topological amplitudes, e.g. $W$-exchange $E$, electroweak penguin $P_{\text{EW}}$. Hence, a large $C_{\text{eff}}/T_{\text{eff}}$ does not necessarily imply large $C/T$. In our calculation we found $|C_{\text{eff}}/T_{\text{eff}}| = 0.71$ for $C_{\text{eff}} = C - E - V$ and $T = T + E + V$, while $|C/T|_{SD} = 0.23$ and $|C/T|_{SD+LD} = 0.33$ [9]. The point is that one needs large $V/T = 0.56 \exp[\pi^2]$ to understand the $\pi\pi$ data.

2.3 Direct CP asymmetries

The strong phases in charmless $B$ decays are governed by final-state rescattering. We see from the last column of Table 1 that direct CP-violating partial rate asymmetries in $K^{-}\pi^{+}$, $\rho^{+}\pi^{-}$ and $\pi^{+}\pi^{-}$ modes are significantly affected by final-state rescattering and their signs are different from that predicted by the short-distance QCDF approach. The direct CP asymmetries in $B^0 \to \pi^0\pi^0, \rho^-\pi^+$ decays are also shown in Table 1 where we see that the predictions of pQCD and QCDF supplemented with FSIIs are opposite in sign. It will be interesting to measure direct CP violation in these two decays to test different models.

3 Polarization Anomaly in $B \to \phi K^*$

For $B \to V_1V_2$ decays with $V$ being a light vector meson, it is expected that they are dominated by longitudinal polarization states and respect the scaling law: $1 - f_L = O(m_V^2/m_B^2)$. However, a low value of the longitudinal fraction $f_L \approx 50\%$ in $\phi K^*$ decays was observed by both BaBar [20] and Belle [21] (see Table 2). This polarization anomaly poses an interesting challenge for any theoretical interpretation.

Working in the context of QCD factorization, Kagan [22] has argued that the lower value of the longitudinal polarization fraction and the large transverse rate can be accommodated by the $(S - P)(S + P)$ penguin-induced annihilation contributions. This is so because the transverse polarization amplitude induced from the above annihilation topologies is of the same $1/m_b$ order as the longitudinal one. Moreover, although the penguin-induced annihilation contribution is formally $1/m_b^2$ suppressed, it can be $O(1)$ numerically. Kagan showed that a fit to the data of $f_L(\phi K^{*0})$ and $f_L(\rho^+ K^{*0})$ favors $\rho_A \sim 0.5$. An alternative suggestion for the solution of the $\phi K^*$
anomaly was advocated in [23] that a energetic transverse gluon from the \(b \rightarrow sg\) chromodipole operator keeps most of its quantum numbers except color when it somehow penetrates through the \(B\) meson surface and descends to a transversely polarized \(\phi\) meson. Sizable transverse components of the \(B \rightarrow \phi K^*\) decay can be accommodated by having \(f_\parallel > f_\perp\). Since the gluon is a flavor singlet, this mechanism can distinguish \(\phi\) from \(\rho\), hence it affects \(B \rightarrow \phi K^*, \omega K^*\) but not \(B \rightarrow K^*\rho\).

Since the scaling law is valid only at short distances, one can also try to circumvent it by considering the long-distance rescattering contributions from intermediate states \(D(s)^*D_s^*\) [9, 24] [25]. The large transverse polarization induced from \(B \rightarrow D^*D_s^*\) will be propagated to \(\phi K^*\) via FSI rescattering. Furthermore, rescattering from \(B \rightarrow D^*D_s\) or \(B \rightarrow DD_s^*\) will contribute only to the \(A_\perp\) amplitude. Recently, we have studied FSI effects on \(B \rightarrow VV\). While the longitudinal polarization fraction can be reduced significantly from short-distance predictions due to such FSI effects, no sizable perpendicular polarization is found owing mainly to the large cancellations occurring in the processes \(B \rightarrow D_s^*D \rightarrow \phi K^*\) and \(B \rightarrow D_sD^* \rightarrow \phi K^*\) and this can be understood as a consequence of \(CP\) and \(SU(3)\) symmetry. Our result is different from a recent similar study in [24]. To fully account for the polarization anomaly (especially the perpendicular polarization) observed in \(B \rightarrow \phi K^*,\) FSI from other states or other mechanism, e.g. the aforementioned penguin-induced annihilation, may have to be invoked.

Li [26] pointed out an interesting observation that the polarization anomaly may be resolved in the pQCD approach provided that the form factor \(A_{BK^*0}(0)\) of order 0.30 is employed. This form factor is indeed very close to the result of 0.31 obtained in the covariant light-front quark model [27]. Since the pQCD approach tends to predict a large value of \(A_{BK^*0}(0)\) (0.46 and 0.41 in the previous pQCD calculations [18, 28]), the task is to see if a small value of \(A_0\) can be naturally produced in such an approach rather than put by hand artificially.

As for the transverse polarization in \(B \rightarrow \rho K^*\) decays, both final-state rescattering and large annihilation scenarios lead to \(f_L(\rho K^*) \sim 60\%\), whereas \(f_L(\rho K^*) \sim 1\) in the model of [23]. However, none of the aforementioned models can explain the observed disparity between \(f_L(\rho^+K^0)\) and \(f_L(\rho^0K^+)(\) see, however, [29] for a possible solution). This should be clarified both experimentally and theoretically.

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### References

[1] BaBar Collaboration, B. Aubert et al., [hep-ex/0407057](http://arxiv.org/abs/hep-ex/0407057).

[2] Belle Collaboration, Y. Chao et al., [hep-ex/0408100](http://arxiv.org/abs/hep-ex/0408100).

[3] Heavy Flavor Averaging Group, [http://www.slac.stanford.edu/xorg/hfag](http://www.slac.stanford.edu/xorg/hfag).

[4] Belle Collaboration, K. Abe et al., Phys. Rev. Lett. 93, 021601 (2004).
[5] BaBar Collaboration, B. Aubert et al., hep-ex/0408089.

[6] M. Beneke, G. Buchalla, M. Neubert, and C.T. Sachrajda, Phys. Rev. Lett. 83, 1914 (1999); Nucl. Phys. B 591, 313 (2000); ibid. 606, 245 (2001).

[7] M. Beneke and M. Neubert, Nucl. Phys. B 675, 333 (2003).

[8] Y.Y. Keum, H.n. Li, and A.I. Sanda, Phys. Rev. D 63, 054008 (2001); Y.Y. Keum and A.I. Sanda, ibid. 67, 054009 (2002).

[9] H.Y. Cheng, C.K. Chua, and A. Soni, hep-ph/0409317.

[10] Y.Y. Keum, private communication.

[11] M. Bander, D. Silverman, and A. Soni, Phys. Rev. Lett. 43, 242 (1979).

[12] J.F. Donoghue et al., Phys. Rev. Lett. 77, 2178 (1996).

[13] M. Suzuki and L. Wolfenstein, Phys. Rev. D 60, 074019 (1999).

[14] O. Grotchakov et al., Z. Phys. A 353, 447 (1996).

[15] H.Y. Cheng, Phys. Rev. D 65, 094012 (2002).

[16] BaBar Collaboration, B. Aubert et al., Phys. Rev. Lett. 93, 051802 (2004).

[17] Belle Collaboration, J. Dragic et al., hep-ex/0405068.

[18] C.D. Lü and M.Z. Yang, Eur. Phys. J. C 23, 275 (2002).

[19] L.L. Chau and H.Y. Cheng, Phys. Rev. Lett. 56, 1655 (1986).

[20] BaBar Collaboration, B. Aubert et al., Phys. Rev. Lett. 91, 171802 (2003); Phys. Rev. D 69, 031102 (2004); hep-ex/0408017 hep-ex/0408063 A.V. Gritsan, hep-ex/0409059.

[21] Belle Collaboration, K.F. Chen et al., Phys. Rev. Lett. 91, 201801 (2003); J. Zhang et al., ibid. 91, 221801 (2003); K. Abe et al., hep-ex/0408141.

[22] A.L. Kagan, Phys. Lett. B 601, 151 (2004); talk presented at 6th Workshop on Higher Luminosity B Factory, Nov. 16-18, 2004, KEK, Japan.

[23] W.S. Hou and M. Nagashima, hep-ph/0408007.

[24] P. Colangelo, F. De Fazio, and T. N. Pham, Phys. Lett. B 597, 291 (2004).

[25] M. Ladisa, V. Laporta, G. Nardulli, and P. Santorelli, hep-ph/0409286.

[26] H.n. Li, hep-ph/0411305.

[27] H.Y. Cheng, C.K. Chua, and C.W. Hwang, Phys. Rev. D 69, 074025 (2004).

[28] C.H. Chen and C.Q. Geng, Nucl. Phys. B 636, 338 (2002).

[29] H.n. Li and S. Mishima, hep-ph/0411146.