CP Violation and New Physics in $B_s$ Decays

Robert Fleischer

Theory Division, CERN, CH-1211 Geneva 23, Switzerland

Abstract

The $B_s$-meson system is a key element in the $B$-physics programme of hadron colliders, offering various avenues to explore CP violation and to search for new physics. One of the most prominent decays is $B_s \to J/\psi \phi$, the counterpart of $B_d \to J/\psi K_S$, providing a powerful tool to search for new-physics contributions to $B_s^0 - \overline{B}_s^0$ mixing. Another benchmark mode is $B_s \to K^+ K^-$, which complements $B_d \to \pi^+ \pi^-$, thereby allowing an extraction of the angle $\gamma$ of the unitarity triangle that is sensitive to new-physics effects in the QCD penguin sector. Finally, we discuss new methods to constrain and determine $\gamma$ with the help of $B_s \to D_s^{(*) \pm} K^\mp$ decays, which complement $B_d \to D_s^{(*) \pm} \pi^\mp$ modes. Since these strategies involve “tree” decays, the values of $\gamma$ thus extracted exhibit a small sensitivity on new physics.

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Robert Fleischer
Theory Division, CERN
CH-1211 Geneva 23, Switzerland

The $B_s$-meson system is a key element in the $B$-physics programme of hadron colliders, offering various avenues to explore CP violation and to search for new physics. One of the most prominent decays is $B_s \rightarrow J/\psi \phi$, the counterpart of $B_d \rightarrow J/\psi K_S$, providing a powerful tool to search for new-physics contributions to $B_s^0$–$\bar{B}_s^0$ mixing. Another benchmark mode is $B_s \rightarrow K^+K^-$, which complements $B_d \rightarrow \pi^+\pi^-$, thereby allowing an extraction of the angle $\gamma$ of the unitarity triangle that is sensitive to new-physics effects in the QCD penguin sector. Finally, we discuss new methods to constrain and determine $\gamma$ with the help of $B_s \rightarrow D_s(\ast)^\pm K_\mp$ decays, which complement $B_d \rightarrow D(\ast)^\pm \pi^\mp$ modes. Since these strategies involve “tree” decays, the values of $\gamma$ thus extracted exhibit a small sensitivity on new physics.

1 Setting the Stage

At the $e^+e^-$ $B$ factories operating at the $\Upsilon(4S)$ resonance, $B_s$ mesons are not accessible, i.e. their decays cannot be explored by the BaBar, Belle and CLEO collaborations. On the other hand, plenty of $B_s$ mesons will be produced at hadron colliders. Consequently, these particles are the “El Dorado” for $B$-decay studies at run II of the Tevatron [1], and later on at the LHC [2]. A detailed overview of the physics potential of $B_s$ mesons can be found in [3].

An important aspect of $B_s$ physics is the mass difference $\Delta M_s$, which can be complemented with $\Delta M_d$ to determine the side $R_t \propto |V_{td}/V_{cb}|$ of the unitarity triangle (UT). To this end, we use that $|V_{cb}| = |V_{ts}|$ to a good accuracy in the Standard Model (SM), and require an $SU(3)$-breaking parameter, which can be determined, e.g. on the lattice. At the moment, only experimental lower bounds on $\Delta M_s$ are available, which can be converted into upper bounds on $R_t$, implying $\gamma \lesssim 90^\circ$ [4]. Once $\Delta M_s$ is measured, more stringent constraints on $\gamma$ will emerge.

Another interesting quantity is the width difference $\Delta \Gamma_s$. While $\Delta \Gamma_d/\Gamma_d$ is negligibly small, where $\Gamma_d$ is the average decay width of the $B_d$ mass eigenstates, $\Delta \Gamma_s/\Gamma_s$ may well be as large as $\mathcal{O}(10\%)$ [5], thereby allowing interesting studies with “untagged” $B_s$ decay rates of the kind

$$\langle \Gamma(B_s(t) \rightarrow f) \rangle \equiv \Gamma(B_s^0(t) \rightarrow f) + \Gamma(\overline{B}_s^0(t) \rightarrow f),$$

where we do not distinguish between initially present $B_s^0$ or $\overline{B}_s^0$ mesons [6].

The focus of the following discussion will be CP violation. If we consider the decay of a neutral $B_q$ meson ($q \in \{d,s\}$) into a final state $|f\rangle$, which is an eigenstate of the CP operator satisfying $\langle \mathcal{CP} | f \rangle = \pm | f \rangle$, we obtain the following time-dependent CP asymmetry [3]:

$$\frac{\Gamma(B_q^0(t) \rightarrow f) - \Gamma(\overline{B}_q^0(t) \rightarrow f)}{\Gamma(B_q^0(t) \rightarrow f) + \Gamma(\overline{B}_q^0(t) \rightarrow f)} = \frac{\mathcal{A}_{\text{dir}} \cos(\Delta M_q t) + \mathcal{A}_{\text{mix}}^{\text{dir}} \sin(\Delta M_q t)}{\cosh(\Delta \Gamma_q t/2) - \mathcal{A}_{\Delta \Gamma} \sinh(\Delta \Gamma_q t/2)},$$

1
where
\[ A_{\text{CP}}^{\text{dir}} = \frac{1 - |\xi_f^{(q)}|^2}{1 + |\xi_f^{(q)}|^2} \quad \text{and} \quad A_{\text{CP}}^{\text{mix}} = \frac{2 \text{Im} \xi_f^{(q)}}{1 + |\xi_f^{(q)}|^2}, \] (3)
with
\[ \xi_f^{(q)} = -e^{-i\phi_q} \left[ \frac{A(B_0 \rightarrow f)}{A(B_0 \rightarrow f)} \right], \] (4)
describe the “direct” and “mixing-induced” CP-violating observables, respectively. In the SM, the CP-violating weak \( B_0^0 \)–\( B_0^0 \) mixing phase \( \phi_q \) is associated with the well-known box diagrams, and is given by
\[ \phi_q = 2 \text{arg}(V_{tq}^* V_{tb}) = \begin{cases} +2\beta & (q = d) \\ -2\lambda^2 \eta & (q = s), \end{cases} \] (5)
where \( \beta \) is the usual angle of the UT. Looking at (2), we observe that \( \Delta \Gamma_q \) provides another observable \( A_{\Delta \Gamma} \), which is, however, not independent from those in (3).

The preferred mechanism for new physics (NP) to manifest itself in (2) is through contributions to \( B_0^0 \)–\( B_0^0 \) mixing, which is a CKM-suppressed, loop-induced, fourth-order weak-interaction process within the SM. Simple dimensional arguments suggest that NP in the TeV regime may well affect the \( \Delta M_{q}^a \) sw e l la st h e \( \phi_q \). If NP enters differently in \( \Delta M_d \) and \( \Delta M_s \), the determination of \( R_t \) from \( \Delta M_d / \Delta M_s \) would be affected. On the other hand, NP contributions to \( \phi_q \) would affect the mixing-induced CP asymmetries \( A_{\text{CP}}^{\text{mix}} \). Scenarios of this kind were considered in several papers; for a selection, see [7]–[11]. Thanks to the “golden” mode \( B_d \rightarrow J/\psi K_S \), direct measurements of \( \sin \phi_d \) are already available. The current world average is given by \( \sin \phi_d \sim 0.734 \), which is in accordance with the indirect range following from the “CKM fits” [4]. Despite this remarkable feature, NP may still hide in the experimental value for \( \sin \phi_d \), since it implies \( \phi_d \sim 47^\circ \lor 133^\circ \), where the former solution would be consistent with the SM, while the second would require NP contributions to \( B^0_d \)–\( \bar{B}^0_d \) mixing. In order to explore these two solutions further, we may complement them with CP violation in \( B_d \rightarrow \pi^+\pi^- \) [12]. Following these lines [11], we obtain an allowed region in the \( \bar{\eta} \)–\( \eta \) plane that is consistent with the SM for \( \phi_d \sim 47^\circ \). In the case of \( \phi_d \sim 133^\circ \), we arrive at a range in the second quadrant, which corresponds to \( \gamma > 90 \) and is consistent with the \( \varepsilon_K \) hyperbola. Interestingly, also this exciting possibility cannot be discarded. The current \( B_d \rightarrow \pi^+\pi^- \) data do not yet allow us to draw definite conclusions, but the situation will significantly improve in the future. As far as \( B_s \) decays are concerned, the burning question in this context is whether \( \phi_s \), which is tiny in the SM, as can be seen in (5), is made sizeable through NP effects. In order to address this issue, the \( B_s \rightarrow J/\psi \phi \) channel plays the key rôle.

2 \( B_s \rightarrow J/\psi \phi \)

This decay is the counterpart of \( B_d \rightarrow J/\psi K_S \), and exhibits an analogous amplitude structure:
\[ A(B_s \rightarrow J/\psi \phi) \propto \left[ 1 + \lambda^2 a e^{i\vartheta} e^{i\gamma} \right]. \] (6)
Here \( \gamma \) is the usual angle of the UT, and the hadronic parameter \( a e^{i\vartheta} \) measures the ratio of penguin to tree contributions, which is naïvely expected to be of \( \mathcal{O}(\lambda) \), where \( \lambda = \mathcal{O}(\lambda) = \)
$O(0.2)$ is a “generic” expansion parameter [10]. In contrast to $B_d \rightarrow J/\psi K_S$, the final state of $B_s \rightarrow J/\psi \phi$ is an admixture of different CP eigenstates, which can be disentangled through an angular analysis of the $J/\psi \rightarrow \ell^+\ell^- \phi \rightarrow K^+K^-$ decay products [13]. Their angular distribution exhibits tiny direct CP violation, whereas mixing-induced CP-violating effects allow the extraction of

$$\sin \phi_s + O(\lambda^3) = \sin \phi_s + O(10^{-3}).$$

Since we have $\phi_s = -2\lambda^2\eta = O(10^{-2})$ in the SM, the determination of this phase from (7) is affected by generic hadronic uncertainties of $O(10\%)$, which may become important for the LHC era. These uncertainties can be controlled with the help of $B_s \rightarrow J/\psi \rho^0$ [14].

Another interesting aspect of the $B_s \rightarrow J/\psi \phi$ angular distribution is that it allows also the determination of $\cos \delta_f \cos \phi_s$, where $\delta_f$ is a CP-conserving strong phase. If we fix the sign of $\cos \phi_s$, which allows an unambiguous determination of $\phi_s$ [15]. In this context, $B_s \rightarrow D_{\pm}\eta^0$, $D_{\pm} \phi$, ... decays are also interesting [16].

The big hope is that experiments will find a sizeable value of $\sin \phi_s$, which would immediately signal NP. There are recent NP analyses where such a feature actually emerges, for example, within SUSY [17], or specific left–right-symmetric models [18].

## 3 $B_s \rightarrow K^+K^-$

The decay $B_s \rightarrow K^+K^-$ is dominated by QCD penguins and complements $B_d \rightarrow \pi^+\pi^-$ nicely, thereby allowing a determination of $\gamma$ with the help of $U$-spin flavour-symmetry arguments [19]. Within the SM, we may write the corresponding decay amplitudes as follows:

$$A(B^0_d \rightarrow \pi^+\pi^-) \propto \left[e^{i\gamma} - d e^{i\theta}\right], \quad A(B^0_s \rightarrow K^+K^-) \propto \left[e^{i\gamma} + \left(1 - \frac{\lambda^2}{\lambda^2}\right) d' e^{i\theta}\right],$$

where the hadronic parameters $de^{i\theta}$ and $d' e^{i\theta'}$ measure the ratios of penguin to tree contributions to $B^0_d \rightarrow \pi^+\pi^-$ and $B^0_s \rightarrow K^+K^-$, respectively. Consequently, we obtain

$$A_{\text{CP}}^\text{dir}(B_d \rightarrow \pi^+\pi^-) = \text{function}(d, \theta, \gamma), \quad A_{\text{CP}}^\text{mix}(B_d \rightarrow \pi^+\pi^-) = \text{function}(d, \theta, \gamma, \phi_d)$$

$$A_{\text{CP}}^\text{dir}(B_s \rightarrow K^+K^-) = \text{function}(d', \theta', \gamma), \quad A_{\text{CP}}^\text{mix}(B_s \rightarrow K^+K^-) = \text{function}(d', \theta', \gamma, \phi_s).$$

As we saw above, $\phi_d$ and $\phi_s$ can “straightforwardly” be fixed, also if NP should contribute to $B^0_q \rightarrow B^0_q$ mixing. Consequently, $A_{\text{CP}}^\text{dir}(B_d \rightarrow \pi^+\pi^-)$ and $A_{\text{CP}}^\text{mix}(B_d \rightarrow \pi^+\pi^-)$ allow us to eliminate $\theta$, thereby yielding $d$ as a function of $\gamma$ in a theoretically clean way. Analogously, we may fix $d'$ as a function of $\gamma$ with the help of $A_{\text{CP}}^\text{dir}(B_s \rightarrow K^+K^-)$ and $A_{\text{CP}}^\text{mix}(B_s \rightarrow K^+K^-)$.

If we look at the corresponding Feynman diagrams, we observe that $B_d \rightarrow \pi^+\pi^-$ and $B_s \rightarrow K^+K^-$ are related to each other through an interchange of all down and strange quarks. Because of this feature, the $U$-spin flavour symmetry of strong interactions implies

$$d = d', \quad \theta = \theta'.$$

Applying the former relation, we may extract $\gamma$ and $d$ from the clean $\gamma-d$ and $\gamma-d'$ contours. Moreover, we may also determine $\theta$ and $\theta'$, allowing an interesting check of the second relation.
This strategy is very promising from an experimental point of view: at CDF-II and LHCb, experimental accuracies for $\gamma$ of $\mathcal{O}(10^{5})$ and $\mathcal{O}(1^{5})$, respectively, are expected [1, 2, 20]. As far as $U$-spin-breaking corrections are concerned, they enter the determination of $\gamma$ through a relative shift of the $\gamma-d$ and $\gamma-d'$ contours; their impact on the extracted value of $\gamma$ depends on the form of these curves, which is fixed through the measured observables. In the examples discussed in [3, 19], the result for $\gamma$ would be very robust under such corrections.

As we have already noted, $B_s \to K^+K^-$ is not accessible at the BaBar and Belle detectors. However, since we obtain $B_s \to K^+K^-$ from $B_d \to \pi^+K^\pm$ through a replacement of the down spectator quark through a strange quark, we have $\text{BR}(B_s \to K^+K^-) \approx \text{BR}(B_d \to \pi^+K^\pm)$. In order to play with the $B$-factory data, we may then consider

$$H = \left(\frac{1 - \lambda^2}{\lambda^2}\right) \left(\frac{f_K}{f_\pi}\right)^2 \frac{\text{BR}(B_d \to \pi^+\pi^-)}{\text{BR}(B_d \to \pi^\pm K^\mp)} \approx 7.5.$$  

(12)

If we use (8) and (11), we may write

$$H = \text{function}(d, \theta, \gamma),$$  

(13)

which complements (9) and provides sufficient information to extract $\gamma$, $d$, and $\theta$ [19, 21]. This approach was applied in the UT analysis sketched at the end of Section 1, following [11]. Interestingly, $H$ implies also a very narrow SM “target range” in the $\mathcal{A}_{\text{mix}}(B_s \to K^+K^-) - \mathcal{A}_{\text{dir}}(B_s \to K^+K^-)$ plane [12]. The measurement of $\text{BR}(B_s \to K^+K^-)$, which is expected to be soon available from CDF-II [22], will already be an important achievement, allowing a better determination of $H$. Once also the CP asymmetries of this channel have been measured, we may fully exploit the physics potential of the $B_s \to K^+K^-$, $B_d \to \pi^+\pi^-$ system [19].

4 $B_s \to D_s^{(*)\pm} K^\mp$

Let us finally turn to colour-allowed “tree” decays of the kind $B_s \to D_s^{(*)\pm} K^\mp$, which complement $B_d \to D_s^{(*)\pm} K^\mp$ transitions: these modes can be treated on the same theoretical basis, and provide new strategies to determine $\gamma$ [23]. Following this paper, we may write these modes generically as $B_q \to D_q \pi_q$. Their characteristic feature is that both a $B^0_q$ and a $\overline{B}^0_q$ may decay into $D_q \pi_q$, thereby leading to interference between $B^0_q - \overline{B}^0_q$ mixing and decay processes, involving the weak phase $\phi_q + \gamma$. In the case of $q = s$, i.e. $D_s \in \{D_s^+, D_s^{*+}, \ldots\}$ and $u_s \in \{K^+, K^{*+}, \ldots\}$, these interference effects are governed by a hadronic parameter $x_s e^{i\phi_s} \propto R_b \approx 0.4$, where $R_b \propto |V_{ub}/V_{cb}|$ is the usual UT side, and hence are large. On the other hand, for $q = d$, i.e. $D_d \in \{D^+, D^{*+}, \ldots\}$ and $u_d \in \{\pi^+, \rho^+, \ldots\}$, the interference effects are described by $x_d e^{i\phi_d} \propto -\lambda^2 R_b \approx -0.02$, and hence are tiny. In the following, we shall only consider $B_q \to D_q \pi_q$ modes, where at least one of the $D_q$, $\pi_q$ states is a pseudoscalar meson; otherwise a complicated angular analysis has to be performed.

The time-dependent rate asymmetries of these decays take the same form as (2). It is well known that they allow a determination of $\phi_q + \gamma$, where the “conventional” approach works as follows [24, 25]: if we measure the observables $C(B_q \to D_q \pi_q) \equiv C_q$ and $C(B_q \to \overline{D}_q u_q) \equiv \overline{C}_q$ provided by the $\cos(\Delta M_q t)$ pieces, we may determine the following quantities:

$$\langle C_q \rangle_+ \equiv \frac{\langle \overline{C}_q + C_q \rangle}{2} = 0, \quad \langle C_q \rangle_- \equiv \frac{\langle \overline{C}_q - C_q \rangle}{2} = \frac{(1 - x_q^2)/(1 + x_q^2)}{2}.$$  

(14)
where \( \langle C_q \rangle - \) allows us to extract \( x_q \). However, to this end we have to resolve terms entering at the \( x_q^2 \) level. In the case of \( q = s \), we have \( x_s = \mathcal{O}(R_0) \), implying \( x_s^2 = \mathcal{O}(0.16) \), so that this may actually be possible, though challenging. On the other hand, \( x_d = \mathcal{O}(-\lambda^2 R_0) \) is doubly Cabibbo-suppressed. Although it should be possible to resolve terms of \( \mathcal{O}(x_d) \), this will be impossible for the vanishingly small \( x_d^2 = \mathcal{O}(0.0004) \) terms, so that other approaches to fix \( x_d \) are required [25]. In order to extract \( \phi_q + \gamma \), the mixing-induced observables \( S(B_q \to D_q \bar{u}_q) \equiv S_q \) and \( S(B_q \to \bar{D}_q u_q) \equiv S_q \) associated with the \( \sin(\Delta M_q t) \) terms of the time-dependent rate asymmetry must be measured. In analogy to (14), it is convenient to introduce observable combinations \( \langle S_q \rangle \pm \). If we assume that \( x_q \) is known, we may consider the quantities

\[
\begin{align*}
  s_+ & \equiv (-1)^L \left[ \frac{1 + x_q^2}{2 x_q} \right] \langle S_q \rangle_+ = + \cos \delta_q \sin \phi_q + \gamma \\
  s_- & \equiv (-1)^L \left[ \frac{1 + x_q^2}{2 x_q} \right] \langle S_q \rangle_- = - \sin \delta_q \cos \phi_q + \gamma,
\end{align*}
\]

which yield

\[
\sin^2(\phi_q + \gamma) = \frac{1}{2} \left[ 1 + s_+^2 - s_-^2 \right] \pm \sqrt{(1 + s_+^2 - s_-^2)^2 - 4 s_+^2 s_-^2}.
\]

This expression implies an eightfold solution for \( \phi_q + \gamma \). If we assume that \( \text{sgn}(\cos \delta_q) > 0 \), as suggested by factorization, a fourfold discrete ambiguity emerges. Note that this assumption allows us also to fix the sign of \( \sin(\phi_q + \gamma) \) through \( \langle S_q \rangle_+ \). To this end, the factor \((−1)^L\), where \( L \) is the \( D_q \bar{u}_q \) angular momentum, has to be properly taken into account [23]. This is a crucial issue for the extraction of the sign of \( \sin(\phi_q + \gamma) \) from \( B_d \to D^{\pm} \pi^\pm \) decays.

Let us now discuss new strategies to explore CP violation through \( B_q \to D_q \bar{u}_q \) modes, following [23]. If the width difference \( \Delta \Gamma_s \) is sizeable, the “untagged” rates (see (1))

\[
\Gamma(B_q(t) \to D_q \pi_q) = \langle \Gamma(B_q \to D_q \pi_q) \rangle \left[ \cosh(\Delta \Gamma_q t/2) - A_{\Delta \Gamma}(B_q \to D_q \pi_q) \sinh(\Delta \Gamma_q t/2) \right] e^{-\Gamma_q t}
\]

and their CP conjugates provide \( A_{\Delta \Gamma}(B_s \to D_s \pi_s) \equiv A_{\Delta \Gamma_s} \) and \( A_{\Delta \Gamma}(B_s \to \bar{D}_s u_s) \equiv \bar{A}_{\Delta \Gamma_s} \). Introducing, in analogy to (14), observable combinations \( \langle A_{\Delta \Gamma_s} \rangle \pm \), we may derive the relations

\[
\tan(\phi_s + \gamma) = - \left[ \frac{\langle S_s \rangle_+}{\langle A_{\Delta \Gamma_s} \rangle_+} \right] = + \left[ \frac{\langle A_{\Delta \Gamma_s} \rangle_-}{\langle S_s \rangle_-} \right],
\]

which allow an unambiguous extraction of \( \phi_s + \gamma \) if we assume, in addition, that \( \text{sgn}(\cos \delta_q) > 0 \). Another important advantage of (19) is that we do not have to rely on \( \mathcal{O}(x_s^2) \) terms, as \( \langle S_s \rangle \pm \) and \( \langle A_{\Delta \Gamma_s} \rangle \pm \) are proportional to \( x_s \). On the other hand, we need a sizeable value of \( \Delta \Gamma_s \). Measurements of untagged rates are also very useful in the case of vanishingly small \( \Delta \Gamma_q \), since the “unevolved” untagged rates in (18) offer various interesting strategies to determine \( x_q \) from the ratio of \( \Gamma(B_q \to D_q \pi_q) \) and \( \Gamma(B_q \to \bar{D}_q u_q) \) and CP-averaged rates of appropriate \( B^\pm \) or flavour-specific \( B_q \) decays.

If we keep the hadronic parameter \( x_q \) and the associated strong phase \( \delta_q \) as “unknown”, free parameters in the expressions for the \( \langle S_q \rangle \pm \), we may obtain bounds on \( \phi_q + \gamma \) from

\[
|\sin(\phi_q + \gamma)| \geq |\langle S_q \rangle_+|, \quad |\cos(\phi_q + \gamma)| \geq |\langle S_q \rangle_-|.
\]

5
If \( x_q \) is known, stronger constraints are implied by
\[
|\sin(\phi_q + \gamma)| \geq |s_+|, \quad |\cos(\phi_q + \gamma)| \geq |s_-|.
\] (21)

Once \( s_+ \) and \( s_- \) are known, we may of course determine \( \phi_q + \gamma \) through the “conventional” approach, using (17). However, the bounds following from (21) provide essentially the same information and are much simpler to implement. Moreover, as discussed in detail in [23] for several examples within the SM, the bounds following from \( B_s \) and \( B_d \) modes may be highly complementary, thereby providing particularly narrow, theoretically clean ranges for \( \gamma \).

Let us now further exploit the complementarity between the processes \( B^0_s \to D_s^{(*)+}K^- \) and \( B^0_d \to D_s^{(*)+}\pi^- \). If we look at the corresponding decay topologies, we observe that these channels are related to each other through an interchange of all down and strange quarks. Consequently, the \( U \)-spin symmetry implies \( a_s = a_d \) and \( \delta_s = \delta_d \), where \( a_s = x_s/R_b \) and \( a_d = -x_d/(\lambda^2 R_b) \) are the ratios of hadronic matrix elements entering \( x_s \) and \( x_d \), respectively. There are various possibilities to implement these relations [23]. A particularly simple picture emerges if we assume that \( a_s = a_d \) and \( \delta_s = \delta_d \), which yields
\[
\tan \gamma = - \left[ \sin \phi_d - S \sin \phi_s \over \cos \phi_d - S \cos \phi_s \right] \phi_s = 0^\circ = - \left[ \sin \phi_d \over \cos \phi_d - S \right].
\] (22)

Here we have introduced
\[
S = - R \left[ \langle S_d \rangle_+ \right]
\] (23)
with
\[
R = \left( {1 - \lambda^2 \over \lambda^2} \right) \left[ {1 \over 1 + x_s^2} \right],
\] (24)

where \( R \) can be fixed with the help of untagged \( B_s \) rates through
\[
R = \left( {f_K \over f_\pi} \right)^2 \left[ \Gamma(B^0_d \to D_s^{(*)+}\pi^-) + \Gamma(B^0_s \to D_s^{(*)-}\pi^+) \over \Gamma(B_s \to D_s^{(*)+}K^-) + \Gamma(B_s \to D_s^{(*)-}K^+) \right].
\] (25)

Alternatively, we may only assume that \( \delta_s = \delta_d \) or that \( a_s = a_d \), as discussed in detail in [23]. Apart from features related to multiple discrete ambiguities, the most important advantage with respect to the “conventional” approach is that the experimental resolution of the \( x_q^2 \) terms is not required. In particular, \( x_d \) does not have to be fixed, and \( x_s \) may only enter through a \( 1 + x_s^2 \) correction, which can straightforwardly be determined through untagged \( B_s \) rate measurements. In the most refined implementation of this strategy, the measurement of \( x_d/x_s \) would only be interesting for the inclusion of \( U \)-spin-breaking effects in \( a_d/a_s \). Moreover, we may obtain interesting insights into hadron dynamics and \( U \)-spin-breaking effects.

In order to explore CP violation, the colour-suppressed counterparts of the \( B_q \to D_q \eta_q \) modes are also very interesting. In the case of the \( B_d \to D\bar{K}_{S(L)} \), \( B_s \to D\eta^{(*)} \), \( D\phi \), \ldots modes, the interference effects between \( B_q^0 - B_q^0 \) mixing and decay processes are governed by \( x_f e^{i\delta_f} \propto R_b \). If we consider the CP eigenstates \( D_{\pm} \), we obtain additional interference effects at the amplitude level, which involve \( \gamma \), and may introduce the following “untagged” rate asymmetry [16]:
\[
\Gamma_{+/-}^L = \frac{\langle \Gamma(B_q \to D_+ f_s) \rangle - \langle \Gamma(B_q \to D_- f_s) \rangle}{\langle \Gamma(B_q \to D_+ f_s) \rangle + \langle \Gamma(B_q \to D_- f_s) \rangle},
\] (26)
which allows us to constrain $\gamma$ through $|\cos \gamma| \geq |\Gamma^f_\pm|$. Moreover, if we complement $\Gamma^f_\pm$ with

$$\langle S^f_\pm \rangle_\pm \equiv (S^f_+ \pm S^f_-)/2,$$

(27)

where $S^f_\pm \equiv A^\text{mix}_{\text{CP}}(B_q \to D_{\pm}f_s)$, we may derive the following simple but exact relation:

$$\tan \gamma \cos \phi_q = \left[ \frac{\eta_f \langle S^f_+ \rangle}{\Gamma^f_+} \right] + \left[ \eta_f \langle S^f_- \rangle - \sin \phi_q \right],$$

(28)

where $\eta_f \equiv (-1)^L \eta^f_{\text{CP}}$. This expression allows a conceptually simple, theoretically clean and essentially unambiguous determination of $\gamma$ [16]; further applications, employing also $D$-meson decays into CP non-eigenstates, can be found in [26]. Since the interference effects are governed by the tiny parameter $x^{f_d} e^{i\delta^{f_d}} \propto -\lambda^2 R_b$ in the case of $B_s \to D_{\pm}K_{S(L)}$, $B_d \to D_{\pm}\pi^0, D_{\pm}\rho^0, ...$, these modes are not as promising for the extraction of $\gamma$. However, they provide the relation

$$\eta_f \langle S^f_\pm \rangle = \sin \phi_q + O(x^{f_d}_b^2) = \sin \phi_q + O(4 \times 10^{-4}),$$

(29)

allowing very interesting determinations of $\phi_q$ with theoretical accuracies one order of magnitude higher than those of the conventional $B_d \to J/\psi K_S, B_s \to J/\psi$ approaches (see Section 2). In particular, $\phi^\text{SM}_s = -2\lambda^2\eta$ could be determined with only $O(1\%)$ uncertainty [16].

5 Conclusions and Outlook

The most exciting question concerning $B_s \to J/\psi \phi$ is whether this mode will exhibit sizeable mixing-induced CP-violating effects, thereby indicating NP contributions to $B^0_s \overline{B}^0_s$ mixing. As we have seen, the $B_s$-meson system offers interesting avenues to extract $\gamma$. For example, we may employ $B_s \to K^+K^-$, which is governed by QCD penguin processes, to complement $B_d \to \pi^+\pi^-$, or may complement pure “tree” decays of the kind $B_s \to D^{(s)\pm}K^{\mp}$ with their $B_d \to D^{(*)\pm}\pi^{\mp}$ counterparts. Here the burning question is whether the corresponding results for $\gamma$ will show discrepancies, which could indicate NP effects in the penguin sector. The exploration of $B_s$ decays is the “El Dorado” for $B$-physics studies at hadron colliders. Important first steps are already expected in the near future at run II of the Tevatron, whereas the rich physics potential of the $B_s$-meson system can be fully exploited by LHCb and BTeV.

References


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