CP Violation and New Physics in $B_s$ Decays

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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 $B^0_s \bar{B}^0_s$ 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 $B_s \to D_s^{(\ast)} \pm K^\mp$ decays, which complement $B_d \to D^{(\ast)} \pm \pi^\mp$ modes. Since these strategies involve “tree” decays, the values of $\gamma$ thus extracted exhibit a small sensitivity on new physics.

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CP Violation and New Physics in $B_s$ Decays

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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 $B_0$-$\bar{B}_0$ 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 $B_s \rightarrow D(\ast)^\pm K^{\mp}$ 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

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

(1)

where we do not distinguish between initially present $B_s^0$ or $\bar{B}_s^0$ 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 $(\mathcal{C}\mathcal{P})|f\rangle = \pm |f\rangle$, we obtain the following time-dependent CP asymmetry [8]:

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

(2)

1
where

\[ A^\text{dir}_{CP} \equiv \frac{1 - |\xi_f^{(q)}|^2}{1 + |\xi_f^{(q)}|^2} \quad \text{and} \quad A^\text{mix}_{CP} \equiv \frac{2 \text{Im} \xi_f^{(q)}}{1 + |\xi_f^{(q)}|^2}, \]  

(3)

with

\[ \xi_f^{(q)} = -e^{-i\phi_q} \left[ \frac{A(B_q^0 \to f)}{A(B_q^0 \to f)} \right], \]  

(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

\[ \phi_q = 2 \arg(V_{tq}^* V_{tb}) = \begin{cases} +2\beta & (q = d) \\ -2\lambda^2 \eta & (q = s) \end{cases}, \]

(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 \( \Delta M_q \), as well as the \( \phi_q \). 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 \( A^\text{mix}_{CP} \). Scenarios of this kind were considered in several papers; for a selection, see [7]–[11]. Thanks to the “golden” mode \( B_d \to J/\psi K_S \), 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 \( B_0^0 - B_0^0 \) mixing. In order to explore these two solutions further, we may complement them with CP violation in \( B_d \to \pi^+\pi^- \) [12]. Following these lines [11], we obtain an allowed region in the \( \bar{\epsilon} - \eta \) 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 \( B_d \to \pi^+\pi^- \) 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 \( B_s \to J/\psi \phi \) channel plays the key rôle.

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

This decay is the counterpart of \( B_d \to J/\psi K_S \), and exhibits an analogous amplitude structure:

\[ A(B_s \to J/\psi \phi) \propto \left[ 1 + \lambda^2 a e^{i\theta} e^{i\gamma} \right]. \]

(6)

Here \( \gamma \) is the usual angle of the UT, and the hadronic parameter \( ae^{i\theta} \) measures the ratio of penguin to tree contributions, which is naïvely expected to be of \( \mathcal{O}(\lambda) \), where \( \lambda = \mathcal{O}(\lambda) = \ldots \).
allow the extraction of 

\( \mathcal{O}(0.2) \) is a “generic” expansion parameter \([10]\). In contrast to \( B_d \to J/\psi K_S \), the final state of \( B_s \to J/\psi \phi \) is an admixture of different CP eigenstates, which can be disentangled through an angular analysis of the \( J/\psi [\to \ell^+ \ell^-] \phi [\to K^+ K^-] \) decay products \([13]\). Their angular distribution exhibits tiny direct CP violation, whereas mixing-induced CP-violating effects allow the extraction of 

\[
\sin \phi_s + \mathcal{O}(X^0) = \sin \phi_s + \mathcal{O}(10^{-3}). \tag{7}
\]

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

Another interesting aspect of the \( B_s \to J/\psi \phi \) 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 \( \cos \delta_f \) through factorization, we may fix the sign of \( \cos \phi_s \), which allows an unambiguous determination of \( \phi_s \) \([15]\). In this context, \( B_s \to D_{\pm} \eta^{(0)} , D_{\pm} \phi , ... \) 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 \( B_s \to K^+ K^- \)

The decay \( B_s \to K^+ K^- \) is dominated by QCD penguins and complements \( B_d \to \pi^+ \pi^- \) 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:

\[
A(B_d^0 \to \pi^+ \pi^-) \propto \left[ e^{i\gamma} - de^{i\theta} \right], \quad A(B_s^0 \to K^+ K^-) \propto \left[ e^{i\gamma} + \left( \frac{1 - \lambda^2}{\lambda^2} \right) d'e^{i\theta} \right], \tag{8}
\]

where the hadronic parameters \( de^{i\theta} \) and \( d'e^{i\theta} \) measure the ratios of penguin to tree contributions to \( B_d^0 \to \pi^+ \pi^- \) and \( B_s^0 \to K^+ K^- \), respectively. Consequently, we obtain

\[
\mathcal{A}_{\text{CP}}^{\text{dir}}(B_d \to \pi^+ \pi^-) = \text{function}(d, \theta, \gamma), \quad \mathcal{A}_{\text{CP}}^{\text{mix}}(B_d \to \pi^+ \pi^-) = \text{function}(d, \theta, \gamma, \phi_d) \tag{9}
\]

\[
\mathcal{A}_{\text{CP}}^{\text{dir}}(B_s \to K^+ K^-) = \text{function}(d', \theta', \gamma), \quad \mathcal{A}_{\text{CP}}^{\text{mix}}(B_s \to K^+ K^-) = \text{function}(d', \theta', \gamma, \phi_s). \tag{10}
\]

As we saw above, \( \phi_d \) and \( \phi_s \) can “straightforwardly” be fixed, also if NP should contribute to \( B_d^0 \to B_s^0 \) mixing. Consequently, \( \mathcal{A}_{\text{CP}}^{\text{dir}}(B_d \to \pi^+ \pi^-) \) and \( \mathcal{A}_{\text{CP}}^{\text{mix}}(B_d \to \pi^+ \pi^-) \) 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 \( \mathcal{A}_{\text{CP}}^{\text{dir}}(B_s \to K^+ K^-) \) and \( \mathcal{A}_{\text{CP}}^{\text{mix}}(B_s \to K^+ K^-) \).

If we look at the corresponding Feynman diagrams, we observe that \( B_d \to \pi^+ \pi^- \) and \( B_s \to K^+ K^- \) 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

\[
d = d', \quad \theta = \theta'. \tag{11}
\]

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 $\mathcal{O}(10^0)$ and $\mathcal{O}(1^0)$, respectively, are expected \cite{11,21,20}. 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 \cite{3,19}, the result for $\gamma$ would be very robust under such corrections.

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

$$H = \left( \frac{1 - \lambda^2}{\lambda^2} \right) \left( \frac{f_K}{f_\pi} \right)^2 \left[ \frac{\text{BR}(B_d \rightarrow \pi^+\pi^-)}{\text{BR}(B_d \rightarrow \pi^+K^\pm)} \right] \sim 7.5.$$  \hspace{1cm} (12)

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

$$H = \text{function}(d, \theta, \gamma),$$ \hspace{1cm} (13)

which complements (9) and provides sufficient information to extract $\gamma$, $d$ and $\theta$ \cite{11,21}. This approach was applied in the UT analysis sketched at the end of Section 1, following \cite{11}. Interestingly, $H$ implies also a very narrow SM “target range” in the $\mathcal{A}_{\text{CP}}^{\text{mix}}(B_s \rightarrow K^+K^-)-\mathcal{A}_{\text{dir}}^{\text{dir}}(B_s \rightarrow K^+K^-)$ plane \cite{12}. The measurement of $\text{BR}(B_s \rightarrow K^+K^-)$, which is expected to be soon available from CDF-II \cite{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 $B_s \rightarrow K^+K^-$, $B_d \rightarrow \pi^+\pi^-$ system \cite{19}.

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

Let us finally turn to colour-allowed “tree” decays of the kind $B_s \rightarrow D_s^{(*)\pm} K^\mp$, which complement $B_d \rightarrow D^{(*)\pm}\pi^\mp$ transitions: these modes can be treated on the same theoretical basis, and provide new strategies to determine $\gamma$ \cite{23}. Following this paper, we may write these modes generically as $B_q \rightarrow D_q\bar{u}q$. Their characteristic feature is that both a $B_q^0$ and a $\bar{B}_q^0$ may decay into $D_q\bar{u}q$, thereby leading to interference between $B_q^0 \bar{B}_q^0$ mixing and decay processes, involving the weak phase $\phi_q + \gamma$. In the case of $q = s$, i.e. $D_s \in \{d^+, d^{**}, ...\}$ and $u_s \in \{K^+, K^{**}, ...\}$, these interference effects are governed by a hadronic parameter $x_se^{i\phi_s} \propto R_b \approx 0.4$, 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. $D_d \in \{d^+, d^{**}, ...\}$ and $u_d \in \{\pi^+, \rho^+, ...\}$, the interference effects are described by $x_de^{i\phi_d} \propto -\lambda^2 R_b \approx -0.02$, and hence are tiny. In the following, we shall only consider $B_q \rightarrow D_q\bar{u}q$ modes, where at least one of the $D_q$, $\bar{u}q$ states is a pseudoscalar; 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 \cite{24,25}: if we measure the observables $C(B_q \rightarrow D_q\bar{u}q) \equiv C_q$ and $C(B_q \rightarrow D_quq) \equiv \bar{C}_q$ provided by the $\cos(\Delta M_q t)$ pieces, we may determine the following quantities:

$$\langle C_q \rangle_+ \equiv (\bar{C}_q + C_q)/2 = 0, \quad \langle C_q \rangle_- \equiv (\bar{C}_q - C_q)/2 = (1 - x_q^2)/(1 + x_q^2),$$ \hspace{1cm} (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 \( x_s = \mathcal{O}(R_b) \), implying \( x_q^2 = \mathcal{O}(0.16) \), so that this may actually be possible, though challenging. On the other hand, \( x_d = \mathcal{O}(R_b) \) 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 \( x_q^2 = \mathcal{O}(0.0004) \) terms, so that other approaches to fix \( x_d \) are required [25]. In order to extract \( \phi_q + \gamma \), the mixing-induced observables \( S(B_q \rightarrow D_q \overline{\pi}_q) \equiv S_q \) and \( S(B_q \rightarrow D_q u_q) \equiv S_q \) 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

\[
\begin{align*}
    s_+ & \equiv (-1)^L \left[ \frac{1 + x_q^2}{2 x_q} \right] \langle S_q \rangle_+ = + \cos \delta_q \sin(\phi_q + \gamma) \quad (15) \\
    s_- & \equiv (-1)^L \left[ \frac{1 + x_q^2}{2 x_q} \right] \langle S_q \rangle_- = - \sin \delta_q \cos(\phi_q + \gamma), \quad (16)
\end{align*}
\]

which yield

\[
\sin^2(\phi_q + \gamma) = \frac{1}{2} \left[ (1 + s_+^2 - s_-^2) \pm \sqrt{(1 + s_+^2 - s_-^2)^2 - 4 s_+^2} \right]. \quad (17)
\]

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 \( D_q \overline{\pi}_q \) angular momentum, has to be properly taken into account [23]. This is a crucial issue for the extraction of the sign of \( \sin(\phi_d + \gamma) \) from \( B_d \rightarrow D^{\pm} \pi^\mp \) decays.

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

\[
\langle \Gamma(B_q(t) \rightarrow D_q \overline{\pi}_q) \rangle = \langle \Gamma(B_q \rightarrow D_q \overline{\pi}_q) \rangle \left[ \cosh(\Delta \Gamma_q t/2) - A_{\Delta \Gamma}(B_q \rightarrow D_q \overline{\pi}_q) \sinh(\Delta \Gamma_q t/2) \right] e^{-\Gamma_q t} \quad (18)
\]

and their CP conjugates provide \( A_{\Delta \Gamma}(B_s \rightarrow D_s \overline{\pi}_s) \equiv A_{\Delta \Gamma_s} \) and \( A_{\Delta \Gamma}(B_s \rightarrow D_s u_s) \equiv \overline{A}_{\Delta \Gamma_s} \).

Introducing, in analogy to (14), observable combinations \( \langle A_{\Delta \Gamma_s} \rangle_{\pm} \), we may derive the relations

\[
\tan(\phi_s + \gamma) = - \left[ \frac{\langle S_s \rangle_+}{\langle A_{\Delta \Gamma_s} \rangle_+} \right] = + \left[ \frac{\langle A_{\Delta \Gamma_s} \rangle_-}{\langle S_s \rangle_-} \right], \quad (19)
\]

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 \( \mathcal{O}(x_q^2) \) 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 \( x_q^2 = \mathcal{O}(0.0004) \), since the “untagged” untagged rates in (18) offer various interesting strategies to determine \( x_q \) from the ratio of \( \langle \Gamma(B_q \rightarrow D_q \overline{\pi}_q) \rangle + \langle \Gamma(B_q \rightarrow D_q u_q) \rangle \) 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

\[
| \sin(\phi_q + \gamma) | \geq | \langle S_q \rangle_+ |, \quad | \cos(\phi_q + \gamma) | \geq | \langle S_q \rangle_- |. \quad (20)
\]
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 $B^0_s \to D^{(*)^+}K^-$ and $B^0_d \to D^{(*)^+}\pi^-$. 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
\[ \tan \gamma = -\left[\frac{\sin \phi_d - S \sin \phi_s}{\cos \phi_d - S \cos \phi_s}\right]_{\phi_s=0^\circ} = -\left[\frac{\sin \phi_d}{\cos \phi_d - S}\right]. \] (22)

Here we have introduced
\[ S = -R \left[\frac{\langle S_d \rangle^+}{\langle S_s \rangle^+}\right], \] (23)

with
\[ R = \left(\frac{1 - \lambda^2}{\lambda^2}\right) \left[\frac{1}{1 + x_s^2}\right]. \] (24)

where $R$ can be fixed with the help of untagged $B_s$ rates through
\[ R = \left(\frac{f_K}{f_\pi}\right)^2 \left[\frac{\Gamma(B^0_s \to D^{(*)^+}K^-) + \Gamma(B^0_s \to D^{(*)^-}\pi^+)}{\Gamma(B_s \to D^{(*)^+}K^-) + \Gamma(B_s \to D^{(*)^-}\pi^+)}\right]. \] (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 $x^2_s$ 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 $B_q \to D_q \pi_q$ modes are also very interesting. In the case of the $B_d \to DK_{S(L)}$, $B_s \to D\eta(\prime), D\phi, ...$ modes, the interference effects between $B^0_q \to \overline{B^0_q}$ mixing and decay processes are governed by $x_{f_s} e^{i\delta_{f_s}} \propto R_b$. 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]:
\[ \Gamma^{f_s}_{+/-} \equiv \frac{\langle \Gamma(B_q \to D_{+}f_s) \rangle - \langle \Gamma(B_q \to D_{-}f_s) \rangle}{\langle \Gamma(B_q \to D_{+}f_s) \rangle + \langle \Gamma(B_q \to D_{-}f_s) \rangle}. \] (26)
which allows us to constrain $\gamma$ through $|\cos \gamma| \geq |\Gamma_{+_-}|$. Moreover, if we complement $\Gamma_{+_-}^{f_s}$ with 
\[
\langle S_{f_s} \rangle_{\pm} \equiv (S_{f_s}^{+} \pm S_{f_s}^{-})/2,
\]
where $S_{f_s}^{\pm} \equiv A_{\text{CP}}^{\text{mix}}(B_{q} \rightarrow D_{\pm} f_s)$, we may derive the following simple but exact relation:
\[
\tan \gamma \cos \phi_q = \left[ \frac{\eta_{f_s} \langle S_{f_s} \rangle_{+}}{\Gamma_{+_-}^{f_s}} \right] + \left[ \eta_{f_s} \langle S_{f_s} \rangle_{-} - \sin \phi_q \right],
\]
where $\eta_{f_s} \equiv (-1)^L \eta_{\text{CP}}$. This expression allows a conceptually simple, theoretically clean and essentially unambiguous determination of $\gamma$ \cite{16}; further applications, employing also $D$-meson decays into CP non-eigenstates, can be found in \cite{26}. Since the interference effects are governed by the tiny parameter $x_{f_d} e^{i\delta_{f_d}} \sim -\lambda^2 R_b$ in the case of $B_s \rightarrow D_{\pm} K_{S(L)}$, $B_d \rightarrow D_{\pm} \pi^0, D_{\pm} \rho^0, ...$, these modes are not as promising for the extraction of $\gamma$. However, they provide the relation
\[
\eta_{f_d} \langle S_{f_d} \rangle_{-} = \sin \phi_q + \mathcal{O}(x_{f_d}^2) = \sin \phi_q + \mathcal{O}(4 \times 10^{-4}),
\]
allowing very interesting determinations of $\phi_q$ with theoretical accuracies one order of magnitude higher than those of the conventional $B_d \rightarrow J/\psi K_S$, $B_s \rightarrow J/\psi \phi$ approaches (see Section \cite{22}). In particular, $\phi_{s}^{\text{SM}} = -2\lambda^2 \eta$ could be determined with only $\mathcal{O}(1\%)$ uncertainty \cite{16}.

5 Conclusions and Outlook

The most exciting question concerning $B_s \rightarrow J/\psi \phi$ is whether this mode will exhibit sizable mixing-induced CP-violating effects, thereby indicating NP contributions to $B_s^0 - \overline{B_s^0}$ mixing. As we have seen, the $B_s$-meson system offers interesting avenues to extract $\gamma$. For example, we may employ $B_s \rightarrow K^+ K^-$, which is governed by QCD penguin processes, to complement $B_d \rightarrow \pi^+ \pi^-$, or may complement pure “tree” decays of the kind $B_s \rightarrow D_s^{(*)\pm} K^\mp$ with their $B_d \rightarrow D^{(*)\pm} \pi^\mp$ 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

[1] K. Anikeev et al., hep-ph/0201071

[2] P. Ball et al., CERN-TH/2000-101 hep-ph/0003238, in CERN Report on Standard Model physics (and more) at the LHC (CERN, Geneva, 2000), p. 305.

[3] R. Fleischer, Phys. Rep. 370 (2002) 537 hep-ph/0207108.

[4] M. Battaglia, A. J. Buras, P. Gambino and A. Stocchi, eds., Proceedings of the First Workshop on the CKM Unitarity Triangle, CERN, February 2002, hep-ph/0304132.
For an overview, see M. Beneke and A. Lenz, J. Phys. G 27 (2001) 1219 [hep-ph/0012222].

I. Dunietz, Phys. Rev. D 52 (1995) 3048 [hep-ph/9501287];
R. Fleischer and I. Dunietz, Phys. Rev. D 55 (1997) 259 [hep-ph/9605220], Phys. Lett. B 387 (1996) 361 [hep-ph/9605221].

Y. Nir and D. J. Silverman, Nucl. Phys. B 345 (1990) 301.

Y. Grossman, Y. Nir and M. P. Worah, Phys. Lett. B 407 (1997) 307 [hep-ph/9704287].

J. P. Silva and L. Wolfenstein, Phys. Rev. D 62 (2000) 014018 [hep-ph/0002122].

R. Fleischer and T. Mannel, Phys. Lett. B 506 (2001) 311 [hep-ph/0101276].

R. Fleischer, G. Isidori and J. Matias, JHEP 0305 (2003) 053 [hep-ph/0302229].

R. Fleischer and J. Matias, Phys. Rev. D 66 (2002) 054009 [hep-ph/0204101].

A. S. Dighe et al., Phys. Lett. B 369 (1996) 144 [hep-ph/9511363];
A. S. Dighe, I. Dunietz and R. Fleischer, Eur. Phys. J. C 6 (1999) 647 [hep-ph/9804253].

R. Fleischer, Phys. Rev. D 60 (1999) 073008 [hep-ph/9903540].

I. Dunietz, R. Fleischer and U. Nierste, Phys. Rev. D 63 (2001) 114015 [hep-ph/0012219].

R. Fleischer, Phys. Lett. B 562 (2003) 234 [hep-ph/0301255].

M. Ciuchini et al., Phys. Rev. D 67 (2003) 075016 [hep-ph/0212397].

K. Kiers, J. Kolb, J. Lee, A. Soni and G. H. Wu, Phys. Rev. D 66 (2002) 095002 [hep-ph/0205082];
D. Silverman, W. K. Sze and H. Yao, UCI-TR-2003-15 [hep-ph/0305013].

R. Fleischer, Phys. Lett. B 459 (1999) 306 [hep-ph/9903456].

A. Golutvin, “Reaching for $\gamma$ (present and future)”, these proceedings.

R. Fleischer, Eur. Phys. J. C 16 (2000) 87 [hep-ph/0001253].

M. Martin, “Branching ratios from $B_s$ and $\Lambda_b$”, these proceedings.

R. Fleischer, CERN-TH/2003-084 [hep-ph/0304027].

R. Aleksan, I. Dunietz and B. Kayser, Z. Phys. C 54 (1992) 653.

I. Dunietz and R. G. Sachs, Phys. Rev. D 37 (1988) 3186 [E: D 39 (1989) 3515];
I. Dunietz, Phys. Lett. B 427 (1998) 179 [hep-ph/9712401];
M. Diehl and G. Hiller, Phys. Lett. B 517 (2001) 125 [hep-ph/0105213];
D. A. Suprun et al., Phys. Rev. D 65 (2002) 054025 [hep-ph/0110159].

R. Fleischer, Nucl. Phys. B 659 (2003) 321 [hep-ph/0301256].