CERN-TH/2003-149
hep-ph/0307033

# 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 $B_d \to D^{(*)\pm}\pi^{\mp}$ modes. Since these strategies involve "tree" decays, the values of $\gamma$ thus extracted exhibit a small sensitivity on new physics.

*Invited talk at FPCP 2003 – Flavour Physics & CP Violation*
*Ecole Polytechnique, Paris, France, 3–6 June 2003*
*To appear in the Proceedings*

CERN-TH/2003-149
July 2003

# CP Violation and New Physics in $B_s$ Decays

*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 $B_s \to J/\psi\phi$, the counterpart of $B_d \to J/\psi K_{\rm S}$, 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 $B_d \to D^{(*)\pm}\pi^\mp$ 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

$$\langle\Gamma(B_s(t) \to f)\rangle \equiv \Gamma(B_s^0(t) \to f) + \Gamma(\overline{B_s^0}(t) \to f), \tag{1}$$

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 $(\mathcal{CP})|f\rangle = \pm|f\rangle$, we obtain the following time-dependent CP asymmetry [3]:

$$\frac{\Gamma(B_q^0(t) \to f) - \Gamma(\overline{B_q^0}(t) \to f)}{\Gamma(B_q^0(t) \to f) + \Gamma(\overline{B_q^0}(t) \to f)} = \left[\frac{\mathcal{A}_{\rm CP}^{\rm dir}\cos(\Delta M_q t) + \mathcal{A}_{\rm CP}^{\rm mix}\sin(\Delta M_q t)}{\cosh(\Delta\Gamma_q t/2) - \mathcal{A}_{\Delta\Gamma}\sinh(\Delta\Gamma_q t/2)}\right], \tag{2}$$

where

$$\mathcal{A}_{\rm CP}^{\rm dir} \equiv \frac{1-|\xi_f^{(q)}|^2}{1+|\xi_f^{(q)}|^2} \quad \text{and} \quad \mathcal{A}_{\rm CP}^{\rm mix} \equiv \frac{2\,{\rm Im}\,\xi_f^{(q)}}{1+|\xi_f^{(q)}|^2}, \tag{3}$$

with

$$\xi_f^{(q)} = -e^{-i\phi_q}\left[\frac{A(\overline{B^0_q}\to f)}{A(B^0_q\to f)}\right], \tag{4}$$

describe the "direct" and "mixing-induced" CP-violating observables, respectively. In the SM, the CP-violating weak $B^0_q$–$\overline{B^0_q}$ mixing phase $\phi_q$ is associated with the well-known box diagrams, and is given by

$$\phi_q = 2\arg(V_{tq}^*V_{tb}) = \left\{\begin{array}{cc} +2\beta & (q=d) \\ -2\lambda^2\eta & (q=s), \end{array}\right. \tag{5}$$

where $\beta$ is the usual angle of the UT. Looking at (2), we observe that $\Delta\Gamma_q$ provides another observable $\mathcal{A}_{\Delta\Gamma}$, 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 $B^0_q$–$\overline{B^0_q}$ 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 $\mathcal{A}_{\rm CP}^{\rm mix}$. Scenarios of this kind were considered in several papers; for a selection, see [7]–[11]. Thanks to the "golden" mode $B_d\to J/\psi K_{\rm S}$, 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 $\phi_d\sim 47^\circ \vee 133^\circ$, where the former solution would be consistent with the SM, while the second would require NP contributions to $B^0_d$–$\overline{B^0_d}$ 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 $\overline{\rho}$–$\overline{\eta}$ 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 $B_d\to J/\psi K_{\rm S}$, and exhibits an analogous amplitude structure:

$$A(B_s\to J/\psi\,\phi) \propto \left[1+\lambda^2 a e^{i\vartheta}e^{i\gamma}\right]. \tag{6}$$

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

$\mathcal{O}(0.2)$ 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

$$\sin\phi_s + \mathcal{O}(\overline{\lambda}^3) = \sin\phi_s + \mathcal{O}(10^{-3}). \tag{7}$$

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, $B_s \to D_\pm\eta^{(\prime)}$, $D_\pm\phi$, ... 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

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

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

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

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 $\mathcal{O}(10^\circ)$ and $\mathcal{O}(1^\circ)$, 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 $B_d \to \pi^\mp K^\pm$ through a replacement of the down spectator quark through a strange quark, we have $\mathrm{BR}(B_s \to K^+K^-) \approx \mathrm{BR}(B_d \to \pi^\mp K^\pm)$. In order to play with the $B$-factory data, we may then consider

$$H = \left(\frac{1-\lambda^2}{\lambda^2}\right)\left(\frac{f_K}{f_\pi}\right)^2\left[\frac{\mathrm{BR}(B_d \to \pi^+\pi^-)}{\mathrm{BR}(B_d \to \pi^\mp K^\pm)}\right] \sim 7.5. \tag{12}$$

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

$$H = \mathrm{function}(d, \theta, \gamma), \tag{13}$$

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 $\mathcal{A}_{\mathrm{CP}}^{\mathrm{mix}}(B_s \to K^+K^-)$–$\mathcal{A}_{\mathrm{CP}}^{\mathrm{dir}}(B_s \to K^+K^-)$ plane [12]. The measurement of $\mathrm{BR}(B_s \to K^+K^-)$, 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 $B_d \to D^{(*)\pm}\pi^\mp$ 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 $B_q \to D_q \overline{u}_q$. Their characteristic feature is that both a $B_q^0$ and a $\overline{B_q^0}$ may decay into $D_q \overline{u}_q$, thereby leading to interference between $B_q^0$–$\overline{B_q^0}$ mixing and decay processes, involving the weak phase $\phi_q + \gamma$. In the case of $q = s$, i.e. $D_s \in \{D_s^+, D_s^{*+}, ...\}$ and $u_s \in \{K^+, K^{*+}, ...\}$, these interference effects are governed by a hadronic parameter $x_s e^{i\delta_s} \propto R_b \approx 0.4$, 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 $x_d e^{i\delta_d} \propto -\lambda^2 R_b \approx -0.02$, and hence are tiny. In the following, we shall only consider $B_q \to D_q \overline{u}_q$ modes, where at least one of the $D_q$, $\overline{u}_q$ 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 $C(B_q \to D_q \overline{u}_q) \equiv C_q$ 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:

$$\langle C_q \rangle_+ \equiv (\overline{C}_q + C_q)/2 = 0, \quad \langle C_q \rangle_- \equiv (\overline{C}_q - C_q)/2 = (1 - x_q^2)/(1 + x_q^2), \tag{14}$$

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 $x_s^2 = \mathcal{O}(0.16)$, so that this may actually be possible, though challenging. On the other hand, $x_d = \mathcal{O}(-\lambda^2 R_b)$ 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 $x_d^2 = \mathcal{O}(0.0004)$ terms, so that other approaches to fix $x_d$ are required [25]. In order to extract $\phi_q+\gamma$, the mixing-induced observables $S(B_q \to D_q \overline{u}_q) \equiv S_q$ and $S(B_q \to \overline{D}_q u_q) \equiv \overline{S}_q$ 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

$$s_+ \equiv (-1)^L \left[\frac{1+x_q^2}{2x_q}\right] \langle S_q \rangle_+ = +\cos\delta_q \sin(\phi_q+\gamma) \tag{15}$$

$$s_- \equiv (-1)^L \left[\frac{1+x_q^2}{2x_q}\right] \langle S_q \rangle_- = -\sin\delta_q \cos(\phi_q+\gamma), \tag{16}$$

which yield

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

This expression implies an eightfold solution for $\phi_q+\gamma$. If we assume that $\mathrm{sgn}(\cos\delta_q)>0$, 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 $D_q \overline{u}_q$ angular momentum, has to be properly taken into account [23]. This is a crucial issue for the extraction of the sign of $\sin(\phi_d+\gamma)$ from $B_d \to D^{*\pm}\pi^\mp$ decays.

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

$$\langle \Gamma(B_q(t) \to D_q \overline{u}_q) \rangle = \langle \Gamma(B_q \to D_q \overline{u}_q) \rangle \left[\cosh(\Delta\Gamma_q t/2) - \mathcal{A}_{\Delta\Gamma}(B_q \to D_q \overline{u}_q) \sinh(\Delta\Gamma_q t/2)\right] e^{-\Gamma_q t} \tag{18}$$

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

$$\tan(\phi_s+\gamma) = -\left[\frac{\langle S_s \rangle_+}{\langle \mathcal{A}_{\Delta\Gamma_s} \rangle_+}\right] = +\left[\frac{\langle \mathcal{A}_{\Delta\Gamma_s} \rangle_-}{\langle S_s \rangle_-}\right], \tag{19}$$

which allow an *unambiguous* extraction of $\phi_s+\gamma$ if we assume, in addition, that $\mathrm{sgn}(\cos\delta_q)>0$. 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 $\langle \mathcal{A}_{\Delta\Gamma_s} \rangle_\pm$ 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 $\langle \Gamma(B_q \to D_q \overline{u}_q) \rangle + \langle \Gamma(B_q \to \overline{D}_q u_q) \rangle$ 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

$$|\sin(\phi_q+\gamma)| \geq |\langle S_q \rangle_+|, \quad |\cos(\phi_q+\gamma)| \geq |\langle S_q \rangle_-|. \tag{20}$$

If $x_q$ is known, stronger constraints are implied by

$$|\sin(\phi_q+\gamma)| \geq |s_+|, \quad |\cos(\phi_q+\gamma)| \geq |s_-|. \tag{21}$$

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 $B_s^0 \to D_s^{(*)+}K^-$ and $B_d^0 \to D^{(*)+}\pi^-$. 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

$$\tan\gamma = -\left[\frac{\sin\phi_d - S\sin\phi_s}{\cos\phi_d - S\cos\phi_s}\right] \stackrel{\phi_s=0^\circ}{=} -\left[\frac{\sin\phi_d}{\cos\phi_d - S}\right]. \tag{22}$$

Here we have introduced

$$S = -R\left[\frac{\langle S_d\rangle_+}{\langle S_s\rangle_+}\right] \tag{23}$$

with

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

where $R$ can be fixed with the help of untagged $B_s$ rates through

$$R = \left(\frac{f_K}{f_\pi}\right)^2\left[\frac{\Gamma(\overline{B_s^0} \to D_s^{(*)+}\pi^-) + \Gamma(B_s^0 \to D_s^{(*)-}\pi^+)}{\langle\Gamma(B_s \to D_s^{(*)+}K^-)\rangle + \langle\Gamma(B_s \to D_s^{(*)-}K^+)\rangle}\right]. \tag{25}$$

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 $B_q \to D_q\overline{u}_q$ modes are also very interesting. In the case of the $B_d \to DK_{\mathrm{S(L)}}$, $B_s \to D\eta^{(')}$, $D\phi$, ... modes, the interference effects between $B_q^0$–$\overline{B_q^0}$ 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]:

$$\Gamma_{+-}^{f_s} \equiv \frac{\langle\Gamma(B_q \to D_+f_s)\rangle - \langle\Gamma(B_q \to D_-f_s)\rangle}{\langle\Gamma(B_q \to D_+f_s)\rangle + \langle\Gamma(B_q \to D_-f_s)\rangle}, \tag{26}$$

which allows us to constrain $\gamma$ through $|\cos\gamma| \geq |\Gamma^{f_s}_{+-}|$. Moreover, if we complement $\Gamma^{f_s}_{+-}$ with

$$\langle S_{f_s}\rangle_\pm \equiv (S^{f_s}_+ \pm S^{f_s}_-)/2, \tag{27}$$

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

$$\tan\gamma\cos\phi_q = \left[\frac{\eta_{f_s}\langle S_{f_s}\rangle_+}{\Gamma^{f_s}_{+-}}\right] + \left[\eta_{f_s}\langle S_{f_s}\rangle_- - \sin\phi_q\right], \tag{28}$$

where $\eta_{f_s} \equiv (-1)^L \eta^{f_s}_{\rm CP}$. 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 $B_s \to D_\pm K_{\rm S(L)}$, $B_d \to D_\pm \pi^0, D_\pm \rho^0, ...$, these modes are not as promising for the extraction of $\gamma$. However, they provide the relation

$$\eta_{f_d}\langle S_{f_d}\rangle_- = \sin\phi_q + \mathcal{O}(x^2_{f_d}) = \sin\phi_q + \mathcal{O}(4\times 10^{-4}), \tag{29}$$

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


[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.

[5] For an overview, see M. Beneke and A. Lenz, J. Phys. G **27** (2001) 1219 [hep-ph/0012222].

[6] 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].

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

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

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

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

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

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

[13] 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].

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

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

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

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

[18] 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].

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

[20] A. Golutvin, "Reaching for $\gamma$ (present and future)", *these proceedings*.

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

[22] M. Martin, "Branching ratios from $B_s$ and $\Lambda_b$", *these proceedings*.

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

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

[25] 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].

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