Abstract

Determination of the characteristic $CP$-violating quantity $\sin(2\beta)$ should be the central goal of a $B$-meson factory in its first-round experiments. Except the gold-plated channels $B_d \to \psi K_S$ and $\psi K_L$, three other types of $B_d$ decays to $CP$ eigenstates can also serve for the extraction of $\sin(2\beta)$ in the standard model: (a) the $CP$-forbidden transitions $(B_d\bar{B}_d)_{\Upsilon(4S)} \to (X_c K_S)(X_c K_S)$ and $(X_c K_L)(X_c K_L)$, where $X_c = \psi, \psi', \eta_c$, etc; (b) the decay modes $B_d \to D^{(*)+}D^{(*)-}$ and $D^{(*)0}\bar{D}^{(*)0}$, whose amplitudes have simple isospin relations; and (c) the decay modes $B_d \to (f_{CP})D + (\pi^0, \rho^0, a_1^0, \text{etc})$, in which $f_{CP}$ is a $CP$ eigenstate (such as $\pi^+\pi^-$, $K^+K^-$ or $K_S\pi^0$) arising from either $D^0$ or $\bar{D}^0$ in the neglect of $D^0-\bar{D}^0$ mixing. We carry out an analysis of the $CP$-violating signals existing in these typical processes, without loss of the possibility that new physics might significantly affect $B_d^0-\bar{B}_d^0$ or $K^0-\bar{K}^0$ mixing. We also show that the magnitude of $\sin(2\beta)$ can be well determined, in terms of only $|V_{us}|$, $m_d/m_s$ and $m_u/m_c$, from a variety of quark mass ansätze.
1 Introduction

The origin of $CP$-violating phenomena, observed in neutral kaon decays, has been an intriguing puzzle of particle physics. Among various proposed mechanisms of $CP$ violation \[1\], the most natural and economical one is the Kobayashi-Maskawa (KM) picture which works within the standard electroweak model \[2\]. It is expected that large and theoretically clean signals of $CP$ violation, induced purely by the nontrivial phase of the KM matrix, may manifest themselves in some neutral $B$-meson decays to $CP$ eigenstates \[3, 4\]. This possibility has attracted a lot of phenomenological interest \[5\], leading experimentally to the $B$ factory programs at KEK, SLAC, DESY and LHC (as well as the upgrades of the existing facilities at Cornell and Fermilab).

The central goal of the first-round experiments at a $B$-meson factory should be to determine the $CP$-violating phase

$$\beta \equiv \arg\left(-\frac{V_{tb}^*V_{td}}{V_{cb}^*V_{cd}}\right),$$  \hspace{1cm} (1.1)

which represents one angle of the KM unitarity triangle $V_{ub}V_{ud} + V_{cb}V_{cd} + V_{tb}V_{td} = 0$ in the complex plane. The standard model predicts $CP$ asymmetries of the magnitude $\sin(2\beta)$, arising from the interference of decay and $B_d^0$-$\bar{B}_d^0$ mixing, in some $B_d$ decay modes such as $B_d^0$ vs $\bar{B}_d^0 \to \psi K_S$ and $\psi K_L$. The number of $B_d^0\bar{B}_d^0$ events needed for the pragmatic measurement of $\sin(2\beta)$ to three standard deviations can be estimated as follows:

$$N_{BB} = \left[\frac{3}{\sin(2\beta)}\right]^2 \frac{1}{B_{\text{eff}} \epsilon_{\text{com}}},$$  \hspace{1cm} (1.2)

where $B_{\text{eff}}$ is the effective branching fraction of $B_d^0$ or $\bar{B}_d^0$ decaying to a $CP$ eigenstate, and $\epsilon_{\text{com}}$ is the composite detection efficiency of the decay mode under consideration. An analysis of current experimental data on $|V_{ub}/V_{cb}|$, $B_d^0$-$\bar{B}_d^0$ mixing and $\epsilon_K$ yields the constraint $0.32 \leq \sin(2\beta) \leq 0.94$ \[3\]. If we assume $\epsilon_{\text{com}} = 10\%$ and $N_{BB} = 10^7$ (or $10^8$) in the first-round experiments of a $B$ factory, then the size of $B_{\text{eff}}$ is required to be $3.6 \times 10^{-5}$ (or $3.6 \times 10^{-6}$) for $\sin(2\beta) = 0.5$. Taking into account the fact that $B_{\text{eff}}$ should include the cost for flavor tagging of the parent $B_d^0$ and $\bar{B}_d^0$ mesons, one has to choose those $B_d$ decays of interest whose branching ratios are as large as possible.

The gold-plated decay modes for the extraction of $\sin(2\beta)$ are expected to be $B_d \to \psi K_S$ and $B_d \to \psi K_L$ \[3, 4\]. Their decay amplitudes are governed by

$$a_2 |V_{cb}V_{cs}| \approx A\lambda^2 a_2 \sim 8.9 \times 10^{-3}$$  \hspace{1cm} (1.3)

in the naive factorization approximation, where $A (\approx 0.8)$ and $\lambda (\approx 0.22)$ are the Wolfenstein parameters \[7\] and $a_2 (\approx 0.23)$ is the Bauer-Stech-Wirbel (BSW) factorization coefficient \[8, 9\]. There are other two types of $B_d$ decays to $CP$ eigenstates, which have branching ratios comparable in magnitude with that of $B_d^0 \to \psi K_S$. One typical example is $B_d \to D^+D^-$, whose decay amplitude is dominated by

$$a_1 |V_{cb}V_{cd}| \approx A\lambda^3 a_1 \sim 8.8 \times 10^{-3},$$  \hspace{1cm} (1.4)
where $a_1 \approx 1.03$ is the other BSW factorization coefficient [3]. Another typical example is $B_d \to (f_{CP})_D + \pi^0$, in which $f_{CP}$ is a $CP$ eigenstate coming from $D^0$ or $\bar{D}^0$ in the neglect of $D^0-\bar{D}^0$ mixing [4]. The primary transition amplitude of this decay mode is associated dominantly with

$$a_2 |V_{td}V_{ud}| \approx A\lambda^2 a_2 \sim 8.9 \times 10^{-3}.$$  \hfill (1.5)

$CP$ asymmetries in all three types of decays mentioned above are dominated by $\sin(2\beta)$ within the standard model.

In this talk we shall present a three-plus-one strategy to determine the $CP$-violating observable $\sin(2\beta)$. Starting from a variety of quark mass ansätze, we can calculate the KM matrix in terms of quark mass ratios and a $CP$-violating phase. The magnitude of $\sin(2\beta)$ is predictable, as shown in section 2, by use of the well-determined quantities $|V_{us}|$, $m_d/m_s$ and $m_u/m_c$. Nontrivially, section 3 is devoted to $CP$-forbidden decays of the type $(B_d\bar{B}_d)_{T(4S)} \to (\psi K_S)(\psi K_S)$ or $(\psi K_L)(\psi K_L)$, whose decay rates are proportional to $\sin^2(2\beta)$ in the standard model. In section 4, we carry out an isospin analysis of $CP$ violation in $B_d \to D^+D^-$ and $D^0\bar{D}^0$ to extract $\sin(2\beta)$ and probe the penguin-induced phase information. The possibility to determine $\sin(2\beta)$ in decay modes of the type $B_d \to (f_{CP})_D + \pi^0$ is discussed in section 5. A brief summary of our main results, together with some further discussions, is included in section 6.

It is worthwhile to point out that the explicit analyses in the subsequent sections allow the presence of new physics in $B_d^0-\bar{B}_d^0$ mixing and $K^0-\bar{K}^0$ mixing. The relevant $CP$-violating signals turn out to be $\sin(2\beta)$ if we adopt the standard model predictions for the mixing phases, i.e.,

$$\frac{q_B}{p_B} = \frac{V_{tb}^*V_{td}}{V_{td}^*V_{tb}}, \quad \frac{q_K}{p_K} = \frac{V_{cs}^*V_{cd}}{V_{cd}^*V_{cs}}.$$  \hfill (1.6)

Thus most of our results are also valid beyond the standard model, and they should be useful for the experimental studies to be carried out at the forthcoming $B$-meson factories.

## 2 Determination of $\sin(2\beta)$ from mass ansätze

It is expected that flavor mixing parameters can be completely predicted from fermion mass matrices in the framework of a theory beyond the standard model. Before the success in finding this more fundamental theory, the phenomenological approach is to look for the most proper pattern of quark mass matrices which are able to result in the experimentally favored relations between KM matrix elements and quark mass ratios [11, 12]. The relevant symmetries hidden in such quark mass ansätze may provide useful hints towards the dynamical details of fermion mass generation and $CP$ violation [13].

Here let us illustrate a variety of quark mass ansätze in order to predict the magnitude of $\sin(2\beta)$. In the standard model or its extensions which have no flavor-changing right-handed currents, we can choose the up and down quark mass matrices (denoted by $M_u$ and
we can draw the following points:

\[ M_a, \text{ respectively) to be Hermitian without loss of generality [14]. We also assume that } M_u \text{ and } M_d \text{ have the parallel structures (i.e., parallel hierarchies and texture zeros), coming naturally from the same dynamics. After the diagonalization of } M_{u,d} \text{ through the unitary transformation } O_{u,d}^\dagger M_{u,d} O_{u,d}, \text{ one obtains the mass eigenvalues. The KM matrix in the charged weak currents turns out to be } V \equiv O_u^\dagger O_d. \text{ Taking into account the facts}
\]

\[
m_u \ll m_c \ll m_t , \\
m_d \ll m_s \ll m_b ,
\]

and [13]

\[
|V_{tb}| > |V_{ud}| > |V_{cb}| > |V_{us}| > |V_{cd}| \\
\Rightarrow |V_{cb}| > |V_{ts}| \\
\Rightarrow |V_{td}| > |V_{ub}| > 0 ,
\]

we can draw the following points:

(a) \( M_{11}^u = M_{11}^d = 0 \) (or \( |M_{11}^{u,d}| \ll |M_{12}^{u,d}| \)) is a sufficient condition to get proper \( |V_{us}| \) and \( |V_{cd}| \) in leading order approximations:

\[
V_{us} \approx \sqrt{\frac{m_d}{m_s}} \exp(i\varphi_{12}) \sqrt{\frac{m_u}{m_c}} , \\
V_{cd} \approx \sqrt{\frac{m_u}{m_c}} - \exp(i\varphi_{12}) \sqrt{\frac{m_d}{m_s}} ,
\]

where \( \varphi_{12} \equiv \arg(M_{12}^u/M_{12}^d) \) is a phase parameter. These two relations form two congruent triangles in the complex plane, the so-called Cabibbo triangles [16].

(b) \( M_{13}^u = M_{13}^d = 0 \) (or \( |M_{13}^{u,d}| \ll |M_{23}^{u,d}| \)), together with condition (a), can approximately lead to

\[
|V_{ub}| \approx \sqrt{\frac{m_u}{m_c}} , \\
|V_{td}| \approx \sqrt{\frac{m_d}{m_s}} .
\]

We observe that \( |V_{ub}|/V_{cb} | < |V_{td}|/V_{ts} | \) due to the fact \( m_u/m_c < m_d/m_s \). It is worth mentioning that relations (2.3) and (2.4) are basically the results of the Fritzsch ansatz, which has texture zeros \( M_{11}^{u,d} = M_{22}^{u,d} = M_{13}^{u,d} = 0 \) [17].

(c) \( |V_{cb}| \) (or \( |V_{ts}| \)) depends upon the relative size between \( M_{22}^{u,d} \) and \( M_{23}^{u,d} \). If they are comparable in magnitude, we arrive at

\[
|V_{cb}| \approx |V_{ts}| \approx \sqrt{\frac{M_{23}^{d}}{M_{22}^{d}}} \frac{m_s}{m_b} - \exp(i\varphi_{23}) \sqrt{\frac{M_{23}^{u}}{M_{22}^{u}}} \frac{m_c}{m_t} \]

in leading order approximations, where \( \varphi_{23} \equiv \arg(M_{23}^u/M_{23}^d) \). It can be shown that the contribution of \( \varphi_{23} \) to \( CP \) violation in the KM matrix is negligibly small.

Indeed conditions (a) and (b) imply that the Hermitian mass matrices \( M_u \) and \( M_d \) should take the following generic form:

\[
\begin{pmatrix}
0 & \times & 0 \\
\times^* & \triangle & \nabla \\
0 & \nabla^* & \boxdot
\end{pmatrix} \quad .
\]
The key point of the above quark mass ansätze is that either of the two Cabibbo triangles can be rescaled by $V_{cb}^*$ or $V_{ts}^*$, and the resultant triangle is congruent approximately with the unitarity triangle $V_{ub}V_{ud} + V_{cb}V_{cd} + V_{tb}V_{td} = 0 \ [13, 16]$. Then three angles of the unitarity triangle are determinable from three sides of the Cabibbo triangle, which are nearly independent of the mass ratios $m_c/m_t$ and $m_s/m_b$ as well as the phase parameter $\varphi_{23}$. For our present purpose, we only write out the expression of \( \sin(2\beta) \):

\[
\sin(2\beta) \approx \frac{1}{2} \left[ \frac{m_s}{m_d} + \frac{1}{|V_{us}|^2} \left( 1 - \frac{m_u}{m_c} \cdot \frac{m_s}{m_d} \right) \right] \sqrt{\frac{4m_u}{m_c} \cdot \frac{m_d}{m_s} - \left( \frac{m_u}{m_c} + \frac{m_d}{m_s} - |V_{us}|^2 \right)^2}.
\tag{2.7}
\]

So far $|V_{us}|$ has been precisely measured \[17\]: $|V_{us}| = 0.2205 \pm 0.0018$. The latest result of the chiral perturbation theory yields $m_s/m_d = 19.3 \pm 0.9$ and $m_u = 5.1 \pm 0.9$ MeV at the scale 1 GeV \[18\]. The value of $m_c(1 \text{ GeV})$ is expected to be in the range $1.0 - 1.6$ GeV, or around 1.35 GeV \[17, 19\]. With these inputs, we calculate $\sin(2\beta)$ and plot the result in Fig. 1. One can see that the prediction of quark mass ansätze for $\sin(2\beta)$ is quite restrictive in spite of some errors associated with quark masses. It lies in the experimentally allowed region $0.32 \leq \sin(2\beta) \leq 0.94$, obtained from the analysis of current data on $|V_{ub}/V_{cb}|, B_d^0\bar{B}_d^0$ mixing and $\epsilon_K$ within the standard model \[6\].

In the above discussions, we did not assume any specific theory (or model) that can naturally guarantee conditions (a) and (b) for $M_u$ and $M_d$. It is very possible that such a theory exists at a superheavy energy scale (e.g., the scale of string theories or that of grand unification theories). Fortunately, the instructive relations (2.3) and (2.4) are independent of the renormalization-group effects to a good degree of accuracy \[20\]; in other words, they hold at both very high and very low energy scales. Thus the prediction (2.7) remains valid even if $M_u$ and $M_d$ are derived at a superheavy scale, and it can be confronted directly with the low-energy experimental data. In contrast, relation (2.5) will be spoiled by the renormalization-group effects, since $|V_{cb}|$ (or $|V_{ts}|$), $m_s/m_d$ and $m_c/m_t$ may have quite different evolution behaviors with energy scales (see, e.g., \[21\]).

### 3 Probing $\sin(2\beta)$ in $(B_d \bar{B}_d)\Upsilon(4S) \rightarrow (\psi K_S)(\psi K_S)$

It was first pointed out by Wolfenstein \[21\] that the search for $CP$-forbidden transitions of the type

\[
(B_d \bar{B}_d)_{\Upsilon(4S)} \rightarrow (f_a f_b)_{CP-even}, \tag{3.1}
\]

where $f_a$ and $f_b$ denote two $CP$ eigenstates with the same $CP$ parity (in contrast, the initial state has the $CP$-odd parity), would serve as a distinctive test of $CP$ violation in the $B_d^0\bar{B}_d^0$ system. For such a joint decay mode, the $CP$-violating signal can be established by measuring the decay rate itself other than the decay rate asymmetry. In practical experiments, this implies that neither flavor tagging (for the parent $B_d$ mesons) nor time-dependent measurements (for the whole decay chain) are necessary. The feasibility of detecting reaction (3.1) depends mainly upon the branching ratios $\mathcal{B}(B_d^0 \rightarrow f_a)$ and $\mathcal{B}(B_d^0 \rightarrow f_b)$. 
The most interesting $CP$-forbidden channels on the $\Upsilon(4S)$ resonance should be

\[
\begin{align*}
(B_d \bar{B}_d)_{\Upsilon(4S)} &\rightarrow (X_c K_S)_{B_d} (X_c K_S)_{\bar{B}_d}, \\
(B_d \bar{B}_d)_{\Upsilon(4S)} &\rightarrow (X_c K_L)_{B_d} (X_c K_L)_{\bar{B}_d},
\end{align*}
\]

(3.2)

in which $X_c$ stands for a set of possible charmonium states that can form $CP$ eigenstates with $K_S$ ($CP$-odd or $CP$-even) and $K_L$ ($CP$-even or $CP$-odd). The typical examples may include \[ X_c = \psi, \psi', \psi'' , \eta_c , \eta'_c , \text{ etc.} \] (3.3)

Since all the transitions $B^0_d \rightarrow X_c K_S$ occur through the same weak interactions, their branching ratios should be comparable in magnitude. Neglecting tiny $CP$ violation in the kaon system, we have $\mathcal{B}(B^0_d \rightarrow X_c K_S) = \mathcal{B}(B^0_d \rightarrow X_c K_L)$ to an excellent degree of accuracy. In contrast with (3.2), the transitions

\[
\begin{align*}
(B_d \bar{B}_d)_{\Upsilon(4S)} &\rightarrow (X_c K_S)_{B_d} (X_c K_L)_{\bar{B}_d}, \\
(B_d \bar{B}_d)_{\Upsilon(4S)} &\rightarrow (X_c K_L)_{B_d} (X_c K_S)_{\bar{B}_d}
\end{align*}
\]

(3.4)

are allowed by $CP$ symmetry. If we make use of $\mathcal{R}(K_S, K_S), \mathcal{R}(K_L, K_L), \mathcal{R}(K_S, K_L)$ and $\mathcal{R}(K_L, K_S)$ to respectively denote the rates of the above four types of joint decays, then

\[
\mathcal{S}_{CP} \equiv \frac{\mathcal{R}(K_S, K_S) + \mathcal{R}(K_L, K_L)}{\mathcal{R}(K_S, K_S) + \mathcal{R}(K_S, K_L) + \mathcal{R}(K_L, K_S) + \mathcal{R}(K_L, K_L)}
\]

(3.5)

is a clean signal of $CP$ violation independent of the ambiguity from hadronic matrix elements. Furthermore, a sum over all possible $X_c$ states as listed in (3.3) can enhance the statistics of a single mode (say, $X_c = \psi$) by several times (even one order) \[ \text{ii} \] without dilution of the $CP$-violating signal $\mathcal{S}_{CP}$. Only if the combined branching fraction of $(B_d \bar{B}_d)_{\Upsilon(4S)} \rightarrow (X_c K_S)(X_c K_S)$ can amount to $10^{-6}$ or so, a signal of the magnitude $\mathcal{S}_{CP} \sim 0.1$ should be explored in the first-round experiments of an $e^+ e^- B$-meson factory.

The generic formulas for coherent $B_d \bar{B}_d$ decays have been presented in the literature (see, e.g., \[ 3, 22, 23 \]). Explicitly, the time-independent decay rate of $(B_d \bar{B}_d)_{\Upsilon(4S)} \rightarrow (f_a f_b)$ can be written as

\[
\mathcal{R}(f_a, f_b) = N_f \left\{ \frac{x^2_d}{1 + x^2_d} \left[ 1 + |\rho_a|^2 |\rho_b|^2 - 2 \Re \left( \frac{q_B q_B^*}{p_B p_B^*} \rho_a \rho_b \right) \right] \right. \\
+ \left. \frac{2 + x^2_d}{1 + x^2_d} \left( |\rho_a|^2 + |\rho_b|^2 - 2 \Re (\rho_a \rho_b) \right) \right\},
\]

(3.6)

where $N_f$ is a normalization factor proportional to the product of the decay rates of $B^0_d \rightarrow f_a$ and $B^0_d \rightarrow f_b$; $x_d \equiv \Delta m / \Gamma \approx 0.73$ is the $B^0_d \rightarrow \bar{B}^0_d$ mixing parameter \[ 17 \]; $q_B / p_B$ signifies the phase information from $B^0_d \rightarrow \bar{B}^0_d$ mixing; and $\rho_{a,b}$ are ratios of the decay amplitudes $A(B^0_d \rightarrow f_a)$ to $A(B^0_d \rightarrow f_{a,b})$. In obtaining the above formula, we have neglected the tiny $CP$

\[ ^1 \text{Note that } \psi' \rightarrow \psi \pi \pi, \psi'' \rightarrow D \bar{D}, \text{ and } \eta'_c \rightarrow \eta_c \pi \pi. \]
violation induced purely by $B^0_d\bar B^0_d$ mixing (i.e., $|q_B/p_B| = 1$ is taken). Within the standard model, $q_B/p_B = (V_{ub}V_{td}^*)/(V_{tb}V_{bd}^*)$ given in (1.6) is a good approximation.

For the cases of $f_{a,b} = X_cK_{S,L}$, $\rho_{a,b}$ turn out to be

$$\rho_{X_cK_S} = -\rho_{X_cK_L} = \pm \frac{q_K^*}{p_K} \frac{V_{cb}V_{cs}^*}{V_{cb}^*V_{cs}},$$

where “±” is the $CP$ parity of $|X_cK_S|$ state, and $q_K/p_K$ stands for the phase information from $K^0\bar K^0$ mixing in the final state (here $|q_K/p_K| = 1$ is assumed). Defining the phase parameter

$$\phi_{\psi K} \equiv \frac{1}{2} \arg \left( \frac{q_B}{q_K^*} \frac{V_{cb}^*V_{cs}}{V_{cb}V_{cs}} \right),$$

then we obtain the following decay rates:

$$\mathcal{R}(K_S, K_S) = \mathcal{R}(K_L, K_L) = 4N_{X_cK} \left[ \frac{x_d^2}{1 + x_d^2} \sin^2(2\phi_{\psi K}) \right],$$

$$\mathcal{R}(K_S, K_L) = \mathcal{R}(K_L, K_S) = 4N_{X_cK} \left[ 2 - \frac{x_d^2}{1 + x_d^2} \sin^2(2\phi_{\psi K}) \right],$$

where the normalization factor $N_{X_cK}$ is proportional to the square of the decay rate of $B^0_d \to X_cK_S$. As a result, the $CP$-violating signal $S_{CP}$ defined in (3.5) reads

$$S_{CP} = \frac{1}{2} \frac{x_d^2}{1 + x_d^2} \sin^2(2\phi_{\psi K}),$$

purely determined by the $B^0_d\bar B^0_d$ mixing parameter $x_d$ and the combined weak phase $\phi_{\psi K}$. In the standard model, we get $\phi_{\psi K} = \beta$ to a good degree of accuracy. If new physics affects $B^0_d\bar B^0_d$ mixing and (or) $K^0\bar K^0$ mixing, however, $\phi_{\psi K}$ could significantly deviate from $\beta$.

For illustration, we plot the magnitude of $S_{CP}$ as a function of $\phi_{\psi K}$ in Fig. 2, with the input $x_d \approx 0.73$. Current constraint on $\beta$ is $9.3^\circ \leq \beta \leq 35^\circ$ [1], at 95% confidence level in the standard model. We see that there is large room for $\phi_{\psi K}$ or $S_{CP}$ to accommodate new physics. The maximal value of $S_{CP}$ ($\approx 0.17$) can be obtained when $\phi_{\psi K} = \pm 90^\circ$.

4 Probing $\sin(2\beta)$ in $B_d \to D^+D^-$ and $D^0\bar D^0$

The measurement of $CP$ asymmetries in $B_d \to D^+D^-$ and $D^0\bar D^0$ can not only cross-check the extraction of $\beta$ from decays of the type $B_d \to \psi K_S$, but also shed some light on the penguin effects and final-state interactions in nonleptonic $B$ decays to double charmed mesons. For this reason, it is worth studying $B_d \to D^+D^-$ and $D^0\bar D^0$ in a model-independent approach. The similar treatment is applicable to the processes $B_d \to D\bar D^*$, $D^*\bar D$, etc.

Let us carry out an isospin analysis of the decay modes $B \to D\bar D$, to relate their weak and strong phases to the relevant observables [24]. The effective weak Hamiltonians responsible for $B_u^- \to D^-D^0$, $B_d^0 \to D^+D^-$, $\bar B_d^0 \to D^0\bar D^0$ and their $CP$-conjugate processes have the
isospin structures \(|1/2, -1/2\rangle\) and \(|1/2, +1/2\rangle\) respectively. The decay amplitudes of these transitions can be written in terms of the \(I = 1\) and \(I = 0\) isospin amplitudes:

\[
A^{+-} \equiv \langle D^+ D^- | \mathcal{H}_{\text{eff}} | B^0_d \rangle = \frac{1}{2} (A_1 + A_0),
\]

\[
A^{00} \equiv \langle D^0 \bar{D}^0 | \mathcal{H}_{\text{eff}} | B^0_d \rangle = \frac{1}{2} (A_1 - A_0),
\]

\[
A^{+0} \equiv \langle D^+ \bar{D}^0 | \mathcal{H}_{\text{eff}} | B^+_u \rangle = A_1 ; (4.1)
\]

and

\[
\bar{A}^{+-} \equiv \langle D^+ D^- | \mathcal{H}_{\text{eff}} | \bar{B}^0_d \rangle = \frac{1}{2} (\bar{A}_1 + \bar{A}_0),
\]

\[
\bar{A}^{00} \equiv \langle D^0 \bar{D}^0 | \mathcal{H}_{\text{eff}} | \bar{B}^0_d \rangle = \frac{1}{2} (\bar{A}_1 - \bar{A}_0),
\]

\[
\bar{A}^{-0} \equiv \langle D^- D^0 | \mathcal{H}_{\text{eff}} | B^-_u \rangle = \bar{A}_1 . (4.2)
\]

The isospin relations (4.1) and (4.2) form two triangles in the complex plane:

\[
A^{+-} + A^{00} = A^{+0},
\]

\[
\bar{A}^{+-} + \bar{A}^{00} = \bar{A}^{-0}. (4.3)
\]

One is able to determine the relative size and phase difference of isospin amplitudes \(A_1\) (\(\bar{A}_1\)) and \(A_0\) (\(\bar{A}_0\)) from the above triangular relations. Denoting

\[
\frac{A_0}{A_1} \equiv z e^{i\theta}, \quad \frac{\bar{A}_0}{\bar{A}_1} \equiv \bar{z} e^{i\bar{\theta}}, (4.4)
\]

then we obtain

\[
z = \sqrt{\frac{2 \left( |A^{+-}|^2 + |A^{00}|^2 \right)}{|A^{+0}|^2}} - 1,
\]

\[
\theta = \arccos \left( \frac{|A^{+-}|^2 - |A^{00}|^2}{z |A^{+0}|^2} \right); (4.5)
\]

and

\[
\bar{z} = \sqrt{\frac{2 \left( |\bar{A}^{+-}|^2 + |\bar{A}^{00}|^2 \right)}{|\bar{A}^{-0}|^2}} - 1,
\]

\[
\bar{\theta} = \arccos \left( \frac{|\bar{A}^{+-}|^2 - |\bar{A}^{00}|^2}{\bar{z} |\bar{A}^{-0}|^2} \right). (4.6)
\]

If \(z = 1\) and \(\theta = 0\), e.g., we find that \(|A^{00}| = 0\) (i.e., the decay mode \(B^0_d \rightarrow D^0 \bar{D}^0\) is forbidden). Note that \(\theta\) (\(\bar{\theta}\)) is in general a mixture of the weak and strong phase shifts, since both \(A_0\) (\(\bar{A}_0\)) and \(A_1\) (\(\bar{A}_1\)) may contain the tree-level and penguin contributions.

It is worth pointing out that the same isospin relations hold for the decay modes \(B \rightarrow D \bar{D}^*\) and \(B \rightarrow D^* \bar{D}\). Of course, the isospin parameters \(z\) (\(\bar{z}\)) and \(\theta\) (\(\bar{\theta}\)) in \(B \rightarrow D \bar{D}, D \bar{D}^*\) and \(D^* \bar{D}\) may be different from one another due to their different final-state interactions.
As for $B \to D^* \bar{D}^*$, the same isospin relations hold separately for the decay amplitudes with helicity $\lambda = -1, 0, \text{or} +1$.

The quantities $|A^{+0}|$ and $|\bar{A}^{-0}|$ are obtainable from the time-independent measurements of decay rates of $B_u^+ \to D^+ D^0$ and $B_u^- \to D^- D^0$. A determination of $|A^{+}|$ ($|A^{00}|$) and $|\bar{A}^{+}|$ ($|\bar{A}^{00}|$) is possible through the time-integrated measurements of $B_d^0$ vs $\bar{B}_d^0 \to D^+ D^- (D^0 \bar{D}^0)$ on the $\Upsilon(4S)$ resonance, where the produced two $B_d$ mesons are in a coherent state (with odd charge-conjugation parity) until one of them decays. In practice, one can use the semileptonic transition of one $B_d$ meson to tag the flavor of the other meson decaying to $D^+ D^-$ or $D^0 \bar{D}^0$. To probe the $CP$ asymmetry induced by the interplay of direct decay and $B_d^0$-$\bar{B}_d^0$ mixing in $B_d \to D D$, the time-dependent measurements are necessary on the $\Upsilon(4S)$ resonance at asymmetric $B$ factories. In such an experimental scenario, the joint decay rates can be given as follows \cite{23, 24}:

$$
\mathcal{R}(l^\pm X^\pm, D^+ D^-; t) \propto |A_t|^2 e^{-\Gamma t} \left[ \frac{|A^{+}|^2 + |\bar{A}^{+}|^2}{2} \mp \frac{|A^{+}|^2 - |\bar{A}^{+}|^2}{2} \cos(x_d \Gamma t) \right] \\
\pm |A^{+}|^2 \text{Im} \left( \frac{q_B \bar{A}^{+}}{p_B A^{+}} \right) \sin(x_d \Gamma t) \tag{4.7}
$$

and

$$
\mathcal{R}(l^\pm X^\pm, D^0 \bar{D}^0; t) \propto |A_t|^2 e^{-\Gamma t} \left[ \frac{|A^{00}|^2 + |\bar{A}^{00}|^2}{2} \mp \frac{|A^{00}|^2 - |\bar{A}^{00}|^2}{2} \cos(x_d \Gamma t) \right] \\
\pm |A^{00}|^2 \text{Im} \left( \frac{q_B \bar{A}^{00}}{p_B A^{00}} \right) \sin(x_d \Gamma t) \tag{4.8}
$$

where $t$ is the proper time difference between the semileptonic and nonleptonic decays \cite{24}.

Denoting

$$
\phi_{DD} \equiv \frac{1}{2} \arg \left( \frac{q_B \bar{A}_1}{p_B A_1} \right) \tag{4.9}
$$

we express coefficients of the $\sin(x_d \Gamma t)$ term in (4.7) and (4.8) in terms of isospin parameters:

$$
\text{Im} \left( \frac{q_B \bar{A}^{+}}{p_B A^{+}} \right) = \frac{|A^{+0} \bar{A}^{-0}|}{4|A^{+0}|^2} \left[ \sin(2\phi_{DD}) - z \sin(\theta - 2\phi_{DD}) \\
+ \bar{z} \sin(\bar{\theta} + 2\phi_{DD}) + z \bar{z} \sin(\bar{\theta} - \theta + 2\phi_{DD}) \right] \tag{4.10}
$$

and

$$
\text{Im} \left( \frac{q_B \bar{A}^{00}}{p_B A^{00}} \right) = \frac{|A^{+0} \bar{A}^{-0}|}{4|A^{00}|^2} \left[ \sin(2\phi_{DD}) + z \sin(\theta - 2\phi_{DD}) \\
- \bar{z} \sin(\bar{\theta} + 2\phi_{DD}) + z \bar{z} \sin(\bar{\theta} - \theta + 2\phi_{DD}) \right] . \tag{4.11}
$$

All the quantities on the right-hand side of (4.10) or (4.11), except $\phi_{DD}$, can be determined through the time-independent measurements of $B \to D \bar{D}$ on the $\Upsilon(4S)$ resonance. Thus

\footnote{Note that the proper time sum of the semileptonic and nonleptonic decays has been integrated out, since it will not be measured at any $B$-meson factory.}
measuring the $CP$-violating observable on the left-hand side of (4.10) or (4.11) will allow a model-independent extraction of $\phi_{DD}$.

Two remarks about the results obtained above are in order:

(1) If the tree-level quark transition $\bar{b} \to (c\bar{c})\bar{d}$ is assumed to dominate the decay amplitude of $B_u^+ \to D^+ \bar{D}^0$, i.e., $\bar{A}_1/A_1 \approx (V_{cb}V^*_{cd})/(V_{cb}V^*_{cd})$, then we get $\phi_{DD} \approx \beta$ as a pure weak phase in the standard model. In general, $\phi_{DD}$ should be a mixture of both weak and strong phases due to the penguin effects [24]. We expect that a comparison of $\phi_{\psi K}$ (extracted from $B_d \to \psi K_S$ or $\psi K_L$) with $\phi_{DD}$ (extracted from $B_d \to D^+ D^-$ or $B_d \to D^0 \bar{D}^0$) would constrain the penguin-induced phase information in $B \to D \bar{D}$.

(2) A special but interesting case is $z = \bar{z} = 1$. It can be obtained if the decay modes $B \to D \bar{D}$ occur dominantly through the tree-level subprocess $b \to (c\bar{c})d$ or $\bar{b} \to (c\bar{c})\bar{d}$. In this case, $A_0$ ($\bar{A}_0$) and $A_1$ ($\bar{A}_1$) have a common KM factor; thus $\theta$ ($\bar{\theta}$) is a pure strong phase shift. This will lead, for arbitrary values of $\theta$ and $\bar{\theta}$, to the relations

\begin{align}
|A^+|^2 + |A^{00}|^2 &= |A^+|^2, \\
|\bar{A}^+|^2 + |\bar{A}^{00}|^2 &= |\bar{A}^-|^2; \quad (4.12)
\end{align}

i.e., the two isospin triangles in (4.3) become right-angled triangles. If $\theta = \bar{\theta}$ is further assumed, we obtain

\begin{align}
\text{Im} \left( \frac{q_B}{p_B} \frac{A^{++}}{A^{++}} \right) &= \frac{|A^+|^2 |A^{00}|^2}{|A^+|^2} \sin (2\phi_{DD}) \cos^2 \frac{\theta}{2}, \\
\text{Im} \left( \frac{q_B}{p_B} \frac{\bar{A}^{00}}{\bar{A}^{00}} \right) &= \frac{|A^+|^2 |A^{00}|^2}{|A^{00}|^2} \sin (2\phi_{DD}) \sin^2 \frac{\theta}{2}. \quad (4.13)
\end{align}

One can see that these two $CP$-violating quantities have the quasi-seesaw dependence on the isospin phase shift $\theta$. The magnitude of $\sin(2\phi_{DD})$ turns out to be

\begin{align}
\sin(2\phi_{DD}) = -\frac{1}{|A^+|^2 |A^{00}|^2} \left[ |A^+|^2 \text{Im} \left( \frac{q_B}{p_B} \frac{\bar{A}^{++}}{A^{++}} \right) + |A^{00}|^2 \text{Im} \left( \frac{q_B}{p_B} \frac{\bar{A}^{00}}{\bar{A}^{00}} \right) \right], \quad (4.14)
\end{align}

apparently independent of $\theta$.

5 Probing $\sin(2\beta)$ in $B_d \to (f_{CP})_D + (\pi^0, \rho^0, \phi^0)$

The third type of $B_d$ decays for the extraction of $\sin(2\beta)$ is expected to be [10]

\begin{align}
B_d \to (f_{CP})_D + (\pi^0, \rho^0, \phi^0, \phi^0, \text{etc}), \quad (5.1)
\end{align}

where the $CP$ eigenstate $f_{CP}$ may come from either $D^0$ or $\bar{D}^0$ in the neglect of non-trivial $D^0-\bar{D}^0$ mixing effects [24, 26]. Such transitions occur only through the tree-level quark diagrams, as illustrated by Fig. 3. We observe that the graph amplitudes of Fig. 3(a) are doubly KM-suppressed with respect to those of Fig. 3(b), and the ratio of their KM factors
is $|V_{cd}/V_{ud}| \cdot |V_{ub}/V_{cb}| \approx 2\%$ \cite{17} in the standard model. Therefore, the contribution from Fig. 3(a) can be safely neglected in discussing indirect CP violation induced by the interplay of decay and $B^0_d - \bar{B}^0_d$ mixing $^\dagger$.

We remark the assumption that possible effects induced by $D^0 - \bar{D}^0$ mixing are negligible in the $B_d$ decay modes under consideration. The latest constraint on the $D^0 - \bar{D}^0$ mixing rate is $r_D < 0.5\%$ \cite{27, 28}, which can be safely neglected for our present purpose. In case that the mixing phase $q_D/\bar{p}_D$ were nonvanishing, it would give rise to measurable CP asymmetries in some neutral $D$-meson decays to $CP$ eigenstates $f_{CP}$ (such as $f_{CP} = \pi^+\pi^-, K^+K^-$ and $K_S\pi^0$) \cite{29, 29}. The current limits on the asymmetries between the decay rates of $D^0 \to f_{CP}$ and $\bar{D}^0 \to f_{CP}$ show no CP violation at the percent level \cite{30}. If we further assume that the penguin amplitude of $D^0 \to f_{CP}$ is not enhanced by possible new physics \cite{10}, i.e., $D^0 \to f_{CP}$ occurs dominantly through the tree-level quark diagrams with a single KM factor, then the overall amplitudes of $B_d^0 \to (f_{CP})_D + \pi^0$ and $\bar{B}^0_d \to (f_{CP})_D + \pi^0$ can be written as follows:

$$
\langle (\pi^+\pi^-)_{D^0} \pi^0 | \mathcal{H}_{eff} | B^0_d \rangle = (V^*_{cb}V_{ud}) (V^*_{cd}V_{ud}) A_{D\pi} A_{\pi\pi},
\langle (K^+K^-)_{D^0} \pi^0 | \mathcal{H}_{eff} | B^0_d \rangle = (V^*_{cb}V_{ud}) (V^*_{cs}V_{us}) A_{D\pi} A_{KK},
\langle (K_S\pi^0)_{D^0} \pi^0 | \mathcal{H}_{eff} | B^0_d \rangle = (V^*_{cb}V_{ud}) (V^*_{cs}V_{us}) p^K_{B} A_{D\pi} A_{\pi\pi},
$$

and

$$
\langle (\pi^+\pi^-)_{D^0} \pi^0 | \mathcal{H}_{eff} | \bar{B}^0_d \rangle = -(V^*_{cb}V_{ud}) (V^*_{cd}V_{ud}) A_{D\pi} A_{\pi\pi},
\langle (K^+K^-)_{D^0} \pi^0 | \mathcal{H}_{eff} | \bar{B}^0_d \rangle = -(V^*_{cb}V_{ud}) (V^*_{cs}V_{us}) A_{D\pi} A_{KK},
\langle (K_S\pi^0)_{D^0} \pi^0 | \mathcal{H}_{eff} | \bar{B}^0_d \rangle = +(V^*_{cb}V_{ud}) (V^*_{cs}V_{us}) q^K_{B} A_{D\pi} A_{\pi\pi}.
$$

Here $A_{D\pi}, A_{\pi\pi}, A_{KK}$ and $A_{\pi\pi}$ denote the hadronic matrix elements containing strong interaction phases; $p^K_{B}$ and $q^K_{B}$ are the $K^0$-$\bar{K}^0$ mixing parameters; and the "±" sign arises from the $CP$-even or $CP$-odd final state. Let us define three phase observables:

$$
\phi_{\pi\pi} \equiv \frac{1}{2} \arg \left( \frac{q_{B} V^*_{cb}V_{cd}}{p_{B} V^*_{cd}V_{cd}} \right),
\phi_{KK} \equiv \frac{1}{2} \arg \left( \frac{q_{B} V^*_{cb}V^*_{cd}}{p_{B} V^*_{cd}V_{cd}} \right),
\phi_{\pi\pi} \equiv \frac{1}{2} \arg \left( \frac{q_{B} V^*_{cb}V_{cd}}{p_{B} V^*_{cd}V_{cd}} \right).
$$

At an asymmetric $B$ factory running on the $\Upsilon(4S)$ resonance, one can measure the following joint decay rates to extract $\phi_{\pi\pi}, \phi_{KK}$ and $\phi_{\pi\pi}$:

$$
\mathcal{R}(l^\pm X^\mp, (\pi^+\pi^-)_D \pi^0; t) \propto |A|^2 e^{-\Gamma t} \left[ 1 \mp \sin(2\phi_{\pi\pi}) \cdot \sin(x_d \Gamma t) \right],
$$

$^\dagger$Direct CP violation may appear due to interference between the graph amplitudes of Fig. 3(a) and Fig. 3(b). These two amplitudes have different isospin structures, hence a strong phase shift between them (denoted by $\Delta$) is possible as the necessary ingredient of a direct CP asymmetry (proportional to sin $\Delta$). However, there is no way to evaluate this strong phase theoretically. Even if $|\sin \Delta| \sim 1$, the CP asymmetry is at most of the percent level because of the KM suppression for the graph amplitudes in Fig. 3(a).
\[
\mathcal{R}(l^\pm X^\mp, (K^+K^-)_D \pi^0; t) \propto |A_t|^2 e^{-\Gamma[t]} [1 \mp \sin(2\phi_{KK}) \cdot \sin(x_d \Gamma t)] , \\
\mathcal{R}(l^\pm X^\mp, (K_S\pi^0)_D \pi^0; t) \propto |A_t|^2 e^{-\Gamma[t]} [1 \pm \sin(2\phi_{K\pi}) \cdot \sin(x_d \Gamma t)] ,
\]
where the semileptonic modes \((l^\pm X^\mp)\) serve for the flavor tagging of \(B_d\) mesons, and \(t\) is the proper time difference between the semileptonic and nonleptonic decays. Of course, these decay rates have different normalization factors.

The feasibility to measure the joint decay modes in (5.5) depends crucially upon the branching ratio of \(B_d^0 \rightarrow \bar{D}^0\pi^0\) and that of \(\bar{D}^0 \rightarrow f_{CP}\). The latter has been determined in experiments of charm physics \([17]\). Current data only yield the upper bound of the common KM factor \(V_{sd}^*V_{td}\). Neglecting tiny isospin-violating effects induced by the mass differences \(m_{D^0} - m_{D^+}\) and \(m_{\pi^0} - m_{\pi^+}\) as well as the life time difference \(\tau_{B_d} - \tau_{B_u}\), we get from (5.6) that \([31]\)

\[
\mathcal{B}(B_d^0 \rightarrow \bar{D}^0\pi^0) \geq \frac{1}{2} \left[ 1 - \left( \frac{\mathcal{B}(B_d^0 \rightarrow D^-\pi^+) \cdot \mathcal{B}(B_u^+ \rightarrow D^0\pi^+)}{\mathcal{B}(B_d^0 \rightarrow D^-\pi^+)} \right)^2 \right] \mathcal{B}(B_u^+ \rightarrow \bar{D}^0\pi^+) .
\]

Since \(\mathcal{B}(B_d^0 \rightarrow D^-\pi^+) = (3.0\pm0.4) \times 10^{-3}\) and \(\mathcal{B}(B_u^+ \rightarrow \bar{D}^0\pi^+) = (5.3\pm0.5) \times 10^{-3}\) have been measured \([17]\), we are able to obtain the lower bound of \(\mathcal{B}(B_d^0 \rightarrow \bar{D}^0\pi^0)\) model-independently, as numerically illustrated in Fig. 4. This result implies that the decay mode \(B_d^0 \rightarrow \bar{D}^0\pi^0\) should be detected soon.

In practice, it is necessary to sum over all possible decay modes of the same nature as \(B_d^0 \rightarrow \bar{D}^0\pi^0\), such as \(B_d^0 \rightarrow \bar{D}^0\rho^0\) and \(\bar{D}^0a_1^0\). These transitions are governed by the same weak interactions, thus their branching ratios are expected to be of the same order \([31, 32]\). It is also a good idea to sum over all possible \(\bar{D}^0 \rightarrow f_{CP}\) decays of the same nature, e.g., \(f_{CP} = K_S\pi^0, K_S\rho^0, K_S\omega, etc\) \([14]\). For a careful classification of \(CP\) parities in the final states of \(B_d \rightarrow (f_{CP})_D + (\pi^0, \rho^0, etc)\), we refer the reader to Ref. \([10]\).

### 6 Concluding remarks

We have discussed three different possibilities to determine the \(CP\)-violating quantity \(\sin(2\beta)\) in the first-round experiments of a \(B\)-meson factory. They should be supplementary to the gold-plated approach, where \(\sin(2\beta)\) is related to the \(CP\) asymmetry in \(B_d \rightarrow \psi K_S\) or \(B_d \rightarrow \psi K_L\) within the standard model. In addition, it has been pointed out that the
magnitude of \( \sin(2\beta) \) can be well constrained from a variety of quark mass ansätze. Some necessary remarks about the results obtained above are in order.

(a) The uncertainties associated with the approximate relations (2.3), (2.4) and (2.7) are expected to be less than 10\%, as a consequence of the significant hierarchy of quark mass values. This accuracy should be good enough to justify or rule out the relevant quark mass ansätze, if \( \sin(2\beta) \) can be measured to the similar extent of precision.

(b) The \( CP \)-violating signal \( S_{CP} \) in (3.10) is worth being pursued experimentally. If there exist some difficulties in detecting it within the first-round experiments of a \( B \) factory, further efforts should be made in the second-round experiments. Some other \( CP \)-forbidden channels of \( B_d\bar{B}_d \) decays on the \( \Upsilon(4S) \) resonance are also interesting for the study of \( CP \) violation.

(c) Within the standard model, we have \( \phi_{\pi\pi} \approx \phi_{KK} \approx \phi_{K\pi} \approx \beta \) to an excellent degree of accuracy. The deviation of \( \phi_{DD} \) from \( \phi_{\psi K} = \phi_{K\pi} \) might not be negligibly small, provided the penguin effects in \( B_d \to D\bar{D} \) were not as small as we naively expected. New physics in \( K^0$$\bar{K}^0 \) mixing could give rise to an observable difference between \( \phi_{\pi\pi} \) (\( \phi_{KK} \)) and \( \phi_{K\pi} \) or between \( \phi_{\psi K} \) and \( \phi_{DD} \). New physics in \( B_d^0$$\bar{B}_d^0 \) mixing would affect all the five \( CP \)-violating phases under discussion, but could not be isolated from the proposed measurements.

Of course, we have assumed unitarity of the \( 3 \times 3 \) KM matrix in the above discussions. Some of our results are indeed independent of this assumption. New physics, which can violate the KM unitarity and in turn affect \( CP \) asymmetries of some \( B_d \) decays, has been classified in Ref. [33].

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Figure 1: The $CP$-violating observable $\sin(2\beta)$ as a function of the quark mass $m_c$ at the scale $\mu = 1$ GeV, predicted by a variety of quark mass ansätze. The upper and lower bounds of $\sin(2\beta)$ (at 95% C.L.) are obtained, within the standard model (SM), from the analysis of current data on $|V_{ub}/V_{cb}|$, $B_d^0\bar{B}_d^0$ mixing and $\epsilon_K$.

Figure 2: The $CP$-violating signal $S_{CP}$ on the $Y(4S)$ resonance. The dark solid curve corresponds to the current constraint on $\phi_{\psi K} \approx \beta$ within the standard model.
Figure 3: The quark diagrams for $B_d^0 \rightarrow D^0 M^0$ and $\bar{D}^0 M^0$ with $M^0 = \pi^0, \rho^0, a_1^0$, etc.

Figure 4: The lower bound of $B(B_d^0 \rightarrow \bar{D}^0 \pi^0)$ as a function of the isospin phase shift $\delta \equiv \arg(A_{3/2}/A_{1/2})$. Here the dark solid curve corresponds to the central values of $B(B_d^0 \rightarrow D^- \pi^+)$ and $B(B_u^+ \rightarrow \bar{D}^0 \pi^+)$, while the dotted region arises from the errors associated with the relevant inputs.