Eur. Phys. J. C 4, 283–287 (1998)
DOI 10.1007/s100529800757

THE EUROPEAN
PHYSICAL JOURNAL C

Towards isolation of new physics in B_s^0-$\bar{B}_s^0$ mixing

Zhi-zhong Xing

Department of Physics, Nagoya University, Chikusa-ku, Nagoya 464-01, Japan (e-mail: xing@eken.phys.nagoya-u.ac.jp)

Received: 24 July 1997 / Published online: 20 February 1998

Abstract. In many extensions of the standard model, new physics is possible to contribute significantly to B_s^0-$\bar{B}_s^0$ mixing. We show that the effect of new physics, both its phase information and its magnitude, can be isolated from measurements of CP asymmetries in the semileptonic B_s transitions and in some nonleptonic B_s decays into hadronic CP eigenstates. We also find that the rates of CP-forbidden decays at the $\Upsilon(5S)$ resonance are not suppressed by the large B_s mass difference, thus they can be used to extract the CP-violating phase of new physics in either B_s^0-$\bar{B}_s^0$ or K^0-$\bar{K}^0$ mixing.

1 Introduction

Recently some experimental efforts have been made to measure the B_s^0-$\bar{B}_s^0$ mixing parameter x_s (corresponding to the mass difference between two B_s mass eigenstates), which is anticipated to be very large in the standard model (SM) [1]. This stimulates some theoretical or phenomenological interest for a more delicate study of the problem of B_s^0-$\bar{B}_s^0$ mixing and CP violation (see, e.g., [2–4]). In particular, the other B_s^0-$\bar{B}_s^0$ mixing parameter y_s (corresponding to the width difference between two B_s mass eigenstates) is predicted to be about 0.1 within the SM, a value making the measurement of untagged B_s data samples meaningful in the future experiments [2,5].

The B_s^0-$\bar{B}_s^0$ mixing system is a good place for the exploration of new physics (NP) beyond the SM [6]. For instance, observation of a sufficiently small x_s, an asymmetry between the semileptonic B_s^0 and $\bar{B}_s^0$ decays, or a significant CP-violating signal in B_s^0 vs $\bar{B}_s^0 \to \psi\phi$ modes should imply the existence of NP in B_s^0-$\bar{B}_s^0$ mixing. In most extensions of the SM, NP is expected to contribute only to the mass difference of B_s mesons [7]. Any NP does not significantly affect the direct decays of B_s mesons via the tree-level W-mediated channles, thus the total decay width Γ remains to amount to its SM value as a good approximation. However, the effect of NP can appear in both x_s and y_s, and it may also give rise to a significant CP-violating phase, which is primarily vanishing in the SM. Consequently it is possible to isolate NP in B_s^0-$\bar{B}_s^0$ mixing through the measurement of x_s, y_s as well as CP violation in both semileptonic and nonleptonic B_s transitions.

Up to now, the possibility to extract the *phase* of NP from CP asymmetries in some B_s decays has been considered (see, e.g., [5,8]). In contrast, little attention has been paid to determining the *magnitude* of NP in B_s^0-$\bar{B}_s^0$ mixing, a crucial quantity for one to get at the specific non-standard electroweak model. In this work we shall discuss a few instructive approaches towards the isolation of NP in B_s^0-$\bar{B}_s^0$ mixing, both its phase information and its magnitude.

The remaining parts of this paper are organized as follows: in Sect. 2 we describe a parametrization of NP in different B_s^0-$\bar{B}_s^0$ mixing quantities. The NP effects on CP asymmetries in semileptonic and nonleptonic decays of B_s mesons are discussed in Sects. 3 and 4, respectively. Section 5 is devoted to extracting the phase information of NP from the CP-forbidden $B_s^0\bar{B}_s^0$ transitions. Finally some concluding remarks are given in Sect. 6.

2 NP in B_s^0-$\bar{B}_s^0$ mixing

In the assumption of CPT invariance, the mass eigenstates of B_s^0 and $\bar{B}_s^0$ mesons can be written as

$$
\begin{aligned}
|B_1\rangle &= p|B_s^0\rangle \ + \ q|\bar{B}_s^0\rangle \, , \\
|B_2\rangle &= p|B_s^0\rangle \ - \ q|\bar{B}_s^0\rangle \, ,
\end{aligned}
\tag{1}
$$

where p and q are complex mixing parameters. In terms of off-diagonal elements of the 2×2 B_s^0-$\bar{B}_s^0$ mixing Hamiltonian $\mathbf{M} - \mathrm{i}\mathbf{\Gamma}/2$, we express the ratio q/p as

$$
\frac{q}{p} \ = \ \sqrt{\frac{M_{12}^* \ - \ \mathrm{i}\Gamma_{12}^*/2}{M_{12} \ - \ \mathrm{i}\Gamma_{12}/2}} \, . \tag{2}
$$

The mass and width differences of B_1 and B_2 are defined by $\Delta m \equiv |m_1 - m_2|$ and $\Delta\Gamma \equiv |\Gamma_1 - \Gamma_2|$, respectively. Furthermore, we define $m \equiv (m_1+m_2)/2$ and $\Gamma \equiv (\Gamma_1+\Gamma_2)/2$. It is also convenient to use the following dimensionless parameters to describe B_s^0-$\bar{B}_s^0$ mixing:

$$
x_s \ \equiv \ \frac{\Delta m}{\Gamma} \, , \qquad y_s \ \equiv \ \frac{\Delta\Gamma}{2\Gamma} \, . \tag{3}
$$

The present experimental lower bound for x_s is $x_s \geq 15$ [1], implying that $\Delta m \gg \Delta\Gamma$ or $|M_{12}| \gg |\Gamma_{12}|$ generally holds. Consequently one gets

$$\Delta m \;=\; 2|M_{12}| \tag{4}$$

to a high degree of accuracy. Then the width difference $\Delta\Gamma$ reads [5]

$$\begin{aligned}\Delta\Gamma &= \frac{4|\mathrm{Re}(M_{12}\Gamma_{12}^*)|}{\Delta m} \\ &= 2|\Gamma_{12}| \cdot |\cos\phi_m| \,,\end{aligned} \tag{5}$$

where $\phi_m \equiv \arg(-M_{12}/\Gamma_{12})$. In addition, calculations based on the SM yields $y_s \approx 0.08$ [3]. This value will always be reduced if B_s^0-$\bar{B}_s^0$ mixing receives CP-violating contributions from NP (i.e., $\phi_m \neq 0$). Thus $x_s/y_s > 10^2$ is true both within and beyond the SM.

Now we assume that NP contributes to Δm through M_{12}, which can be parametrized as

$$M_{12} \;=\; M_{12}^{\rm SM} \;+\; M_{12}^{\rm NP} \;=\; M_{12}^{\rm SM}\left(1 + z e^{\mathrm{i}\theta}\right) \tag{6}$$

with $z \equiv |M_{12}^{\rm NP}/M_{12}^{\rm SM}|$ and $\theta \equiv \arg(M_{12}^{\rm NP}/M_{12}^{\rm SM})$. The effect of NP on Γ_{12} is anticipated to be negligibly small, hence $\Gamma_{12} = \Gamma_{12}^{\rm SM}$ and $\Gamma = \Gamma^{\rm SM}$ hold as a good approximation [6]. Since $M_{12}^{\rm SM}$ and $\Gamma_{12}^{\rm SM}$ are dominantly proportional to the Cabibbo-Kobayashi-Maskawa matrix elements $V_{tb}V_{ts}^*$ and $V_{cb}V_{cs}^*$, respectively, $\phi_m = 0$ turns out to be an excellent approximation in the Wolfenstein phase convention. The presence of NP in M_{12} modifies the SM value of ϕ_m in the following manner:

$$\begin{aligned}\phi_m &= \arg\left(1 + z e^{\mathrm{i}\theta}\right) \\ &= \arctan\left[\frac{z\sin\theta}{1 + z\cos\theta}\right] \,.\end{aligned} \tag{7}$$

The B_s^0-$\bar{B}_s^0$ mixing observables x_s and y_s, in the existence of NP, read

$$x_s \;=\; x_s^{\rm SM}\sqrt{1 + z^2 + 2z\cos\theta} \tag{8}$$

and

$$y_s \;=\; y_s^{\rm SM}\sqrt{1 - \frac{z^2\sin^2\theta}{1 + z^2 + 2z\cos\theta}} \,, \tag{9}$$

where $x_s^{\rm SM}$ and $y_s^{\rm SM}$ are (in principle) calculable from the SM. Finally, the ratio q/p can be expressed as

$$\begin{aligned}\frac{q}{p} &= \sqrt{\frac{(1 + z e^{-\mathrm{i}\theta}) \;+\; \mathrm{i}y_s^{\rm SM}/x_s^{\rm SM}}{(1 + z e^{+\mathrm{i}\theta}) \;+\; \mathrm{i}y_s^{\rm SM}/x_s^{\rm SM}}} \\ &\approx e^{-\mathrm{i}2\phi_m}\left[1 \;-\; \frac{z\sin\theta}{1 + z^2 + 2z\cos\theta}\cdot\frac{y_s^{\rm SM}}{x_s^{\rm SM}}\right]\end{aligned} \tag{10}$$

in the next-to-leading order approximation. Note that ϕ_m is indeed a function of z and θ, as given in (7). Clearly $\arg(q/p) = -2\phi_m$ is a good approximation when we discuss CP violation in nonleptonic B_s decays. The small correction term in q/p, which is proportional to $z\sin\theta$ and $y_s^{\rm SM}/x_s^{\rm SM}$, is crucial for an asymmetry between the semileptonic B_s^0 and $\bar{B}_s^0$ transitions (see Sect. 3).

One can see, from (9), that $y_s \leq y_s^{\rm SM}$ is always true. In contrast, $x_s \leq x_s^{\rm SM}$ takes place only if $\theta \in [-\pi, -\pi/2]$ or $[\pi/2, \pi]$, corresponding to the destructive interference between $M_{12}^{\rm NP}$ and $M_{12}^{\rm SM}$. Provided $z \gg 1$ (i.e., NP dominates B_s^0-$\bar{B}_s^0$ mixing), we obtain $x_s \approx z x_s^{\rm SM}$ and $y_s \approx y_s^{\rm SM}\cos\theta$. In this case, $q/p = e^{-\mathrm{i}2\phi_m}$ holds to a higher degree of accuracy.

3 NP effect in semileptonic B_s decays

Now we examine the NP effect on CP violation in semileptonic transitions of B_s mesons. Assuming both the $\Delta Q = \Delta B$ rule and CPT invariance, one is able to get the time-integrated wrong-sign events of semileptonic B_s decays as follows [9,10]:

$$\begin{aligned}N^-(B_s^0) &= n_0\left[\frac{1}{1 - y_s^2} \;-\; \frac{1}{1 + x_s^2}\right]\left|\frac{q}{p}\right|^2 \,, \\ N^+(\bar{B}_s^0) &= n_0\left[\frac{1}{1 - y_s^2} \;-\; \frac{1}{1 + x_s^2}\right]\left|\frac{p}{q}\right|^2 \,,\end{aligned} \tag{11}$$

where n_0 is a normalization factor proportional to the rate of the right-sign semileptonic B_s decay. The CP asymmetry between the above two processes turns out to be

$$\mathcal{A}_{\rm SL} \;\equiv\; \frac{N^+(\bar{B}_s^0) \;-\; N^-(B_s^0)}{N^+(\bar{B}_s^0) \;+\; N^-(B_s^0)} \;=\; \frac{|p|^4 \;-\; |q|^4}{|p|^4 \;+\; |q|^4} \,, \tag{12}$$

apparently irrelevant to the mixing parameters x_s and y_s.

Such an asymmetry can also be obtained from the same-sign dilepton events of coherent $B_s^0\bar{B}_s^0$ decays at the $\Upsilon(5S)$ resonance, whose rates are given by [10,11]

$$\begin{aligned}N_C^{--}(B_s^0\bar{B}_s^0) &= n_C\left[\frac{1 + C y_s^2}{(1 - y_s^2)^2} \;-\; \frac{1 - C x_s^2}{(1 + x_s^2)^2}\right]\left|\frac{q}{p}\right|^2 \,, \\ N_C^{++}(B_s^0\bar{B}_s^0) &= n_C\left[\frac{1 + C y_s^2}{(1 - y_s^2)^2} \;-\; \frac{1 - C x_s^2}{(1 + x_s^2)^2}\right]\left|\frac{p}{q}\right|^2 ,\end{aligned} \tag{13}$$

in which C ($= \pm 1$) is the charge-conjugation parity of the $B_s^0\bar{B}_s^0$ pair, and n_C denotes the normalization factor proportional to the rates of right-sign semileptonic decays of B_s^0 and $\bar{B}_s^0$ mesons. Obviously, we arrive at the same CP asymmetry $\mathcal{A}_{\rm SL}$:

$$\frac{N_C^{++}(B_s^0\bar{B}_s^0) \;-\; N_C^{--}(B_s^0\bar{B}_s^0)}{N_C^{++}(B_s^0\bar{B}_s^0) \;+\; N_C^{--}(B_s^0\bar{B}_s^0)} \;=\; \frac{|p|^4 \;-\; |q|^4}{|p|^4 \;+\; |q|^4} \,, \tag{14}$$

for both $C = -1$ and $C = +1$ cases.

Taking $|M_{12}| \gg |\Gamma_{12}|$ into account, one can obtain $\mathcal{A}_{\rm SL} \approx \mathrm{Im}(\Gamma_{12}/M_{12})$ as a good approximation. By use of (10), we find that $\mathcal{A}_{\rm SL}$ explicitly reads

$$\mathcal{A}_{\rm SL} \;\approx\; \frac{2z\sin\theta}{1 + z^2 + 2z\cos\theta}\cdot\frac{y_s^{\rm SM}}{x_s^{\rm SM}} \,. \tag{15}$$

Within the SM (i.e., $z = 0$ and $\theta = 0$), $\mathcal{A}_{\rm SL}$ is vanishingly small. In view of the fact $y_s^{\rm SM}/x_s^{\rm SM} \sim O(10^{-2})$, we believe

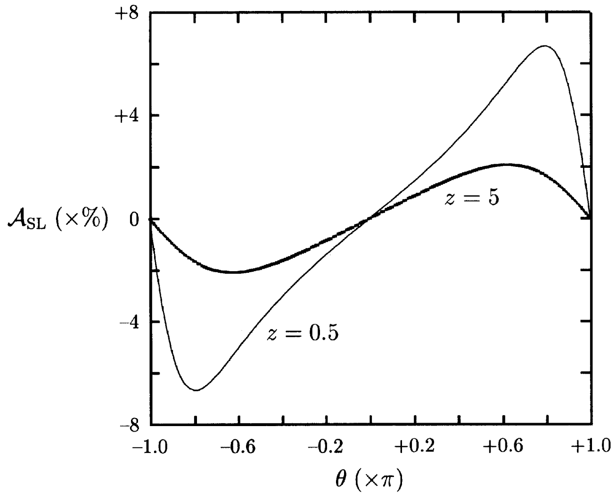

Fig. 1. Illustrative plot for changes of the semileptonic CP asymmetry $\mathcal{A}_{\rm SL}$ with θ, where $y_s^{\rm SM}/x_s^{\rm SM} = 0.05$ and $z = 0.5$ or $z = 5$ have been taken

that the presence of NP in B_s^0-$\bar{B}_s^0$ mixing could enhance the asymmetry $\mathcal{A}_{\rm SL}$ to the level of $10^{-3} - 10^{-2}$. For the purpose of illustration, we plot changes of $\mathcal{A}_{\rm SL}$ with the CP-violating phase θ by assuming $y_s^{\rm SM}/x_s^{\rm SM} = 0.05$ and $z = 0.5$ (or $z = 5$) in Fig. 1. It is clear that $\mathcal{A}_{\rm SL}$ may reach the percent level for favorable values of θ and z. Thus the measurement of $\mathcal{A}_{\rm SL}$ at the forthcoming B-meson factories can provide some useful information or constraints on NP in B_s^0-$\bar{B}_s^0$ mixing.

Finally it is worth emphasizing that reliable calculations of $y_s^{\rm SM}/x_s^{\rm SM}$ are crucial for probing any NP effect through $\mathcal{A}_{\rm SL}$. A detailed analysis of $y_s^{\rm SM}/x_s^{\rm SM}$ with updated data can be found in [12].

4 NP effect in nonleptonic B_s decays

Let us turn attention to the NP effect on CP asymmetries in some nonleptonic B_s decays into CP eigenstates, such as $D_s^+ D_s^-$, $D_s^{*+} D_s^- \oplus D_s^+ D_s^{*-}$, $(\psi\phi)_+$ (i.e., the CP-even state of $\psi\phi$) [4], ψK_S and ψK_L. These transitions can be used to extract the CP-violating phase ϕ_m, in most cases, without ambiguities from hadronic matrix elements. They are also favorable in view of the nearest experimental detectability for weak decays of B_s mesons at e^+e^- colliders and (or) hadron machines. Obviously, $q/p = e^{-{\rm i}2\phi_m}$ should be an excellent approximation.

Due to B_s^0-$\bar{B}_s^0$ mixing, the time-dependent rates of B_s mesons decaying into a CP eigenstate f can be written as [10]

$$
\begin{aligned}
&\mathcal{R}[B^0_{s,\rm phys}(t) \to f] \\
&= n_f e^{-\Gamma t} \left[\frac{1+|\lambda_f|^2}{2} \cosh(y_s \Gamma t) \; - \; {\rm Re}\lambda_f \sinh(y_s \Gamma t) \right. \\
&\left. + \; \frac{1-|\lambda_f|^2}{2} \cos(x_s \Gamma t) \; - \; {\rm Im}\lambda_f \sin(x_s \Gamma t) \right] , \\
&\mathcal{R}[\bar{B}^0_{s,\rm phys}(t) \to f] \\
&= n_f e^{-\Gamma t} \left[\frac{1+|\lambda_f|^2}{2} \cosh(y_s \Gamma t) \; - \; {\rm Re}\lambda_f \sinh(y_s \Gamma t) \right. \\
&\left. - \; \frac{1-|\lambda_f|^2}{2} \cos(x_s \Gamma t) \; + \; {\rm Im}\lambda_f \sin(x_s \Gamma t) \right] , \qquad (16)
\end{aligned}
$$

where n_f denotes the normalization factor proportional to $|\langle f|\mathcal{H}_{\rm eff}|B_s^0\rangle|^2$, and λ_f stands for a rephasing-invariant quantity defined by

$$
\lambda_f \; \equiv \; e^{-{\rm i}2\phi_m} \frac{\langle f|\mathcal{H}_{\rm eff}|\bar{B}_s^0\rangle}{\langle f|\mathcal{H}_{\rm eff}|B_s^0\rangle} \; . \qquad (17)
$$

We see that the x_s-induced oscillations, anticipated to be rapid in the SM, cancel in the rate of the flavor-untagged data samples:

$$
\begin{aligned}
\mathcal{S}(t) &\equiv \mathcal{R}[B^0_{s,\rm phys}(t) \to f] \; + \; \mathcal{R}[\bar{B}^0_{s,\rm phys}(t) \to f] \; , \\
&= n_f e^{-\Gamma t} \left[\left(1+|\lambda_f|^2\right) \cosh(y_s \Gamma t) \right. \\
&\quad \left. - \; 2{\rm Re}\lambda_f \sinh(y_s \Gamma t) \right] \; . \qquad (18)
\end{aligned}
$$

In contrast, the y_s-induced oscillations disappear in the rate asymmetry between the flavor-tagged data sample:

$$
\begin{aligned}
\mathcal{A}(t) &\equiv \mathcal{R}[B^0_{s,\rm phys}(t) \to f] \; - \; \mathcal{R}[\bar{B}^0_{s,\rm phys}(t) \to f] \; , \\
&= n_f e^{-\Gamma t} \left[\left(1-|\lambda_f|^2\right) \cos(x_s \Gamma t) \right. \\
&\quad \left. - \; 2{\rm Im}\lambda_f \sin(x_s \Gamma t) \right] \; . \qquad (19)
\end{aligned}
$$

Therefore, measurements of $\mathcal{S}(t)$ and (or) $\mathcal{A}(t)$ allow one to determine ${\rm Re}\lambda_f$ versus $(1+|\lambda_f|^2)$ and (or) ${\rm Im}\lambda_f$ versus $(1-|\lambda_f|^2)$ in most cases, except the special situation where y_s is vanishingly small (e.g., $y_s \leq 10^{-2}$) due to NP and x_s is considerably large (e.g., $x_s > 50$).

Within the SM, $\lambda_{D_s^+ D_s^-} \approx 1$, $\lambda_{(\psi\phi)_+} \approx 1$, $\lambda_{\psi K_S} \approx -1$ and $\lambda_{\psi K_L} \approx 1$ hold as a good approximation. For these decay modes, the strong or electroweak penguin contributions (dominated by the top quark in the loop) are significantly suppressed in magnitude and do not give rise to additional CP-violating phases [13]. Beyond the SM, the NP appearing in B_s^0-$\bar{B}_s^0$ mixing is likely to simultaneously manifest itself in the loop-induced penguin amplitudes. Generally we do not expect that the NP-enhanced penguin contributions can be significant enough in those B_s transitions mentioned above [6], and a test of this assumption is possible by detecting the deviation of $|\lambda_f|$ from unity in $\mathcal{A}(t)$. Here we assume $|\lambda_f| \approx 1$ to hold even in the presence of NP, for $B_s \to D_s^+ D_s^-$, $(\psi\phi)_+$, ψK_S and ψK_L. Then we obtain

$$
\begin{aligned}
{\rm Re}\lambda_{D_s^+ D_s^-} &\approx {\rm Re}\lambda_{(\psi\phi)_+} \; \approx \; -{\rm Re}\lambda_{\psi K_S} \\
&\approx {\rm Re}\lambda_{\psi K_L} \; \approx \; \cos(2\phi_m) \; , \\
{\rm Im}\lambda_{D_s^+ D_s^-} &\approx {\rm Im}\lambda_{(\psi\phi)_+} \; \approx \; -{\rm Im}\lambda_{\psi K_S} \\
&\approx {\rm Im}\lambda_{\psi K_L} \; \approx \; -\sin(2\phi_m) \; , \qquad (20)
\end{aligned}
$$

where

$$
\begin{aligned}
\cos(2\phi_m) \; &= \; \frac{1 + 2z\cos\theta + z^2\cos(2\theta)}{1 + z^2 + 2z\cos\theta} \; , \\
\sin(2\phi_m) \; &= \; \frac{2z\sin\theta\,(1 + z\cos\theta)}{1 + z^2 + 2z\cos\theta} \; . \qquad (21)
\end{aligned}
$$

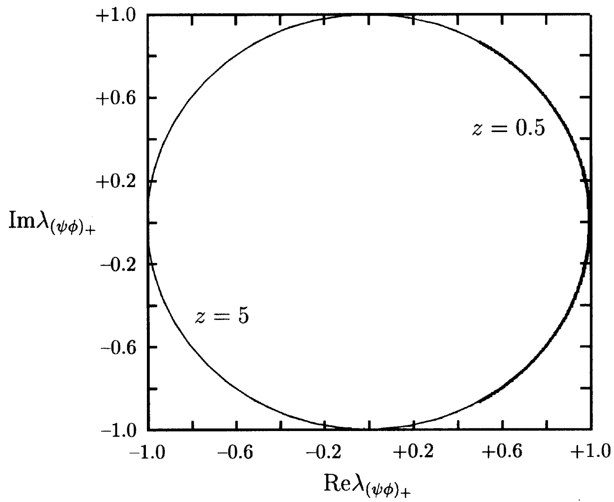

Fig. 2. Illustrative plot for the correlation between two observables $\mathrm{Re}\lambda_{(\psi\phi)_+}$ and $\mathrm{Im}\lambda_{(\psi\phi)_+}$, where $\theta \in [-\pi, +\pi]$ and $z = 0.5$ (*dark solid curve*) or $z = 5$ (*solid circle*) have been taken

Through this way, we can determine the phase information of NP in B_s^0-$\bar{B}_s^0$ mixing. To illustrate the correlation between two observables, e.g., $\mathrm{Re}\lambda_{(\psi\phi)_+}$ and $\mathrm{Im}\lambda_{(\psi\phi)_+}$, we plot them in Fig. 2 by taking $\theta \in [-\pi, +\pi]$ and $z = 0.5$ or $z = 5$. Clearly the two-fold ambiguity in extracting ϕ_m from $\mathrm{Re}\lambda_{(\psi\phi)_+} \approx \cos(2\phi_m)$ can be removed if $\mathrm{Im}\lambda_{(\psi\phi)_+} \approx -\sin(2\phi_m)$ is also measured.

It is obvious that a determination of ϕ_m itself cannot isolate the magnitude (z) and phase (θ) of NP effects in B_s^0-$\bar{B}_s^0$ mixing. For this reason, combining the measurements of $\mathcal{A}_{\mathrm{SL}}$, $\mathcal{S}(t)$ and (or) $\mathcal{A}(t)$ is useful to determine or constrain z and θ separately.

5 NP effect in CP-forbidden $B_s^0\bar{B}_s^0$ decays

If the expected large size of x_s^{SM} cannot be remarkably reduced by the contribution from NP in B_s^0-$\bar{B}_s^0$ mixing, then a study of CP-forbidden $B_s^0\bar{B}_s^0$ decays at the $\Upsilon(5S)$ resonance will become interesting. The straightforward reason is that the rates of such CP-forbidden transitions are not suppressed by the largeness of x_s, as we shall show subsequently. The CP-violating signal may be established directly from the observation of a CP-forbidden decay mode itself other than the decay rate asymmetry, thus neither flavor tagging nor time-dependent measurements are necessary in practical experiments [9]. For a B-meson factory running at the $\Upsilon(5S)$ resonance, one will have a chance to detect the typical CP-forbidden channels like $(B_s^0\bar{B}_s^0)_{C=+1} \rightarrow (\psi K_S)(\psi K_L)$ and $(B_s^0\bar{B}_s^0)_{C=-1} \rightarrow (\psi K_S)(\psi K_S)$ or $(\psi K_L)(\psi K_L)$, where CP parities of the initial and final states are opposite. For simplicity and illustration, we concentrate only on the coherent decays of $(B_s^0\bar{B}_s^0)_{C=-1}$ events later on.

The time-independent rate of a $(B_s^0\bar{B}_s^0)_{C=-1}$ pair decaying into the final state $(f_1 f_2)$ can be given by [9,10]

$$\mathcal{R}(f_1, f_2) = n_{12}\left\{\frac{x_s^2}{1+x_s^2}\left[1+|\lambda_{f_1}|^2|\lambda_{f_2}|^2 - 2\mathrm{Re}\left(\lambda_{f_1}\lambda_{f_2}\right)\right] + \frac{2+x_s^2}{1+x_s^2}\left[|\lambda_{f_1}|^2+|\lambda_{f_2}|^2 - 2\mathrm{Re}\left(\lambda_{f_1}\lambda^*_{f_2}\right)\right]\right\} , \qquad (22)$$

where n_{12} is a normalization factor proportional to the product of $\mathcal{R}(B_s^0 \rightarrow f_1)$ and $\mathcal{R}(B_s^0 \rightarrow f_2)$; and λ_{f_i} ($i = 1$ or 2) is defined like λ_f in (17). In obtaining the above formula, we have made a safe approximation: $1/(1-y_s^2) \approx 1$ due to $y_s \leq y_s^{\mathrm{SM}} \sim 0.1$. For those B_s decays into hadronic CP eigenstates via quark subprocesses $b \rightarrow (c\bar{c})s$ and $b \rightarrow (c\bar{c})d$, $|\lambda_{f_1}| \approx |\lambda_{f_2}| \approx 1$ is expected to hold to a good degree of accuracy. With the help of $\mathcal{R}(B_s^0 \rightarrow \psi K_S) \approx \mathcal{R}(B_s^0 \rightarrow \psi K_L)$, one can calculate the following ratios of two joint decay rates:

$$\xi_{1K} \equiv \frac{\mathcal{R}(D_s^+D_s^-, \psi K_L)}{\mathcal{R}(D_s^+D_s^-, \psi K_S)} , \qquad \xi_{2K} \equiv \frac{\mathcal{R}(\psi K_S, \psi K_S)}{\mathcal{R}(\psi K_S, \psi K_L)} . \qquad (23)$$

The results explicitly read

$$\xi_{1K} = \frac{x_s^2\left[1-\mathrm{Re}(\lambda_{D_s^+D_s^-}\lambda_{\psi K_L})\right]+(2+x_s^2)\left[1-\mathrm{Re}(\lambda_{D_s^+D_s^-}\lambda^*_{\psi K_L})\right]}{x_s^2\left[1-\mathrm{Re}(\lambda_{D_s^+D_s^-}\lambda_{\psi K_S})\right]+(2+x_s^2)\left[1-\mathrm{Re}(\lambda_{D_s^+D_s^-}\lambda^*_{\psi K_S})\right]} , \quad = \frac{x_s^2\sin^2(2\phi_m-\phi_K)+(2+x_s^2)\sin^2\phi_K}{x_s^2\cos^2(2\phi_m-\phi_K)+(2+x_s^2)\cos^2\phi_K} ; \qquad (24)$$

and

$$\xi_{2K} = \frac{x_s^2\left[1-\mathrm{Re}(\lambda^2_{\psi K_S})\right]+(2+x_s^2)\left[1-\mathrm{Re}(|\lambda_{\psi K_S}|^2)\right]}{x_s^2\left[1-\mathrm{Re}(\lambda_{\psi K_S}\lambda_{\psi K_L})\right]+(2+x_s^2)\left[1-\mathrm{Re}(\lambda_{\psi K_S}\lambda^*_{\psi K_L})\right]} , \quad = \frac{x_s^2\sin^2[2(\phi_m-\phi_K)]}{x_s^2\cos^2[2(\phi_m-\phi_K)]+(2+x_s^2)} , \qquad (25)$$

where the phase parameter ϕ_K comes from K^0-$\bar{K}^0$ mixing in the final states of $B_s \rightarrow \psi K_S$ and ψK_L: $2\phi_K \equiv -\arg(q_K/p_K)$. Within the SM, $\phi_K \approx 0$ is an excellent approximation. However, the NP appearing in B_s^0-$\bar{B}_s^0$ mixing is also likely to manifest itself in K^0-$\bar{K}^0$ mixing, leading probably to a small CP-violating phase ϕ_K (see, e.g., [8]). One can see that the above two CP-violating signals are not suppressed by the largeness of x_s. For $x_s \geq 15$ [1], we approximate (24) and (25) to

$$\xi_{1K} \approx \frac{\sin^2(2\phi_m-\phi_K)+\sin^2\phi_K}{\cos^2(2\phi_m-\phi_K)+\cos^2\phi_K} , \qquad \xi_{2K} \approx \frac{\sin^2[2(\phi_m-\phi_K)]}{\cos^2[2(\phi_m-\phi_K)]+1} . \qquad (26)$$

Clearly the information on ϕ_m and ϕ_K is extractable from ξ_{1K} and ξ_{2K}. If $\phi_K \approx 0$ is further taken [8], then (26) is simplified as

$$\xi_{1K} \approx \xi_{2K} \approx \frac{\sin^2(2\phi_m)}{1+\cos^2(2\phi_m)} , \qquad (27)$$

which is a pure function of ϕ_m.

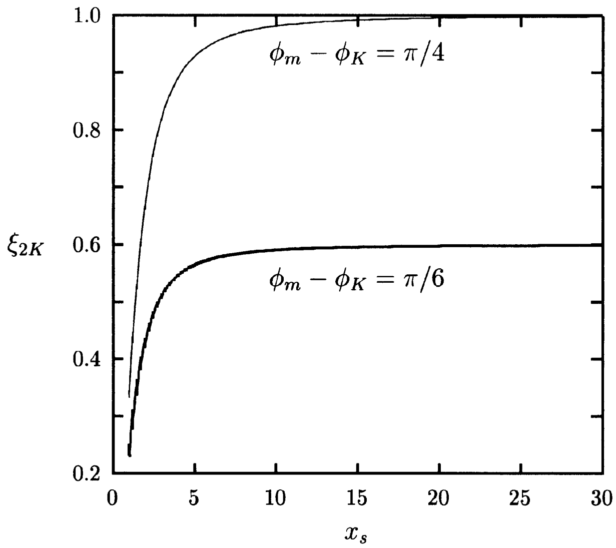

Fig. 3. Illustrative plot for the CP-violating signal ξ_{2K} changing with $x_s \in [1, 30]$, where $(\phi_m - \phi_K) = \pi/6$ or $\pi/4$ has been taken

For the purpose of illustration, we assume $\phi_m - \phi_K = \pi/6$ or $\pi/4$ to plot the change of ξ_{2K} with $x_s \in [1, 30]$ in Fig. 3. We find that ξ_{2K} becomes close to its maximal value only if $x_s \geq 10$, for a definite input of $\phi_m - \phi_K$. Hence such a CP-violating signal will be of particular experimental interest, provided NP introduces a significant phase into B_s^0-$\bar{B}_s^0$ (or K^0-$\bar{K}^0$) mixing.

6 Concluding remarks

We have studied some phenomenological possibilities to isolate the NP effect in B_s^0-$\bar{B}_s^0$ mixing. If $x_s^{\rm SM}$ (or $y_s^{\rm SM}$) can be reliably evaluated and x_s (or y_s) can be experimentally determined, then significant deviation of $x_s/x_s^{\rm SM}$ (or $y_s/y_s^{\rm SM}$) from unity will be a clear signal of NP. In addition, NP in B_s^0-$\bar{B}_s^0$ mixing is likely to enhance the semileptonic CP asymmetry $\mathcal{A}_{\rm SL}$ to the level of $10^{-3} - 10^{-2}$. The phase information of NP can be extracted from either flavor-tagged or flavor-untagged data samples of B_s decays into hadronic CP eigenstates, or both of them, depending upon the experimental sensitivity to the x_s- and (or) y_s-induced oscillations. It is remarkable that the rates of CP-forbidden $B_s^0\bar{B}_s^0$ transitions on the $\Upsilon(5S)$ resonance are not suppressed by the largeness of x_s, and this feature will be useful for measuring the CP-violating phase in B_s^0-$\bar{B}_s^0$ (or K^0-$\bar{K}^0$) mixing.

Of course, the experimental feasibility of the above-proposed measurements has to be studied in some detail [12]. In particular, the complication to observe coherent $B_s^0\bar{B}_s^0$ decays at the $\Upsilon(5S)$ resonance, where coherent $B_d^0\bar{B}_d^0$ events can also be produced [12,14], should be taken seriously. We hope that a systematic search for various CP-violating signals in the B_s^0-$\bar{B}_s^0$ system will be available in the future experiments of B-meson physics.

Acknowledgements. The author likes to thank A.I. Sanda for some useful discussions. This work was supported by the Japan Society for the Promotion of Science.

References

1. L. Gibbons (CLEO Collaboration), invited talk given at the Int. Conf. on HEP, Warsaw, Poland (July 1996)
2. I. Dunietz, Phys. Rev. **D 52** (1995) 3048
3. M. Beneke, G. Buchalla, and I. Dunietz, Phys. Rev. **D 54** (1996) 4419
4. A.S. Dighe, I. Dunietz, H.J. Lipkin, and J.L. Rosner, Phys. Lett. **B 369** (1996) 144
5. Y. Grossman, Phys. Lett. **B 380** (1996) 99
6. I.I. Bigi, V.A. Khoze, N.G. Uraltsev, and A.I. Sanda, in *CP Violation*, edited by C. Jarlskog (World Scientific, Singapore, 1988), p. 175;
   J.L. Hewett, T. Takeuchi, and S. Thomas, SLAC-PUB-7088 or CERN-TH/96-56;
   Y. Grossman, Y. Nir, and R. Rattazzi, SLAC-PUB-7379 or CERN-TH-96-368;
   M. Gronau and D. London, Phys. Rev. **D 55** (1997) 2845
7. Y. Nir, Phys. Lett. **B 327** (1994) 85
8. A.G. Cohen, D.B. Kaplan, F. Lepeintre, and A.E. Nelson, Phys. Rev. Lett. **78** (1997) 2300;
   J.P. Silva and L. Wolfenstein, Phys. Rev. **D 55** (1997) 5331
9. Z.Z. Xing, Phys. Rev. **D 53** (1996) 204
10. Z.Z. Xing, Phys. Rev. **D 55** (1997) 196
11. Z.Z. Xing, Phys. Lett. **B 379** (1996) 257
12. A comprehensive study of this problem is underway
13. Y. Nir and H.R. Quinn, Annu. Rev. Nucl. Part. Sci. **42** (1992) 211
14. P. Krawczyk, D. London, and H. Steger, Nucl. Phys. **B 321** (1989) 1