CP asymmetry in $B_d \to \phi K_S$: Standard model pollution

Yuval Grossman
Stanford Linear Accelerator Center, Stanford University, Stanford, California 94309

Gino Isidori
Stanford Linear Accelerator Center, Stanford University, Stanford, California 94309
and INFN, Laboratori Nazionali di Frascati, I-00044 Frascati, Italy

Mihir P. Worah
Stanford Linear Accelerator Center, Stanford University, Stanford, California 94309
and Department of Physics, University of Washington, Seattle, Washington 98185

(Received 9 September 1997; published 29 July 1998)

The difference in the time dependent CP asymmetries between the modes $B \to \psi K_S$ and $B \to \phi K_S$ is a clean signal for physics beyond the standard model. This interpretation could fail if there is a large enhancement of the matrix element of the $b \to u\bar{u}s$ operator between the B_d initial state and the ϕK_S final state. We argue against this possibility and propose some experimental tests that could shed light on the situation. [S0556-2821(98)04817-6]

PACS number(s): 13.25.Hw, 11.30.Er, 12.15.Hh

It is well known that in the standard model the time-dependent CP-violating asymmetry in $B_d \to \psi K_S$ $[a_{CP}(\psi K_S)]$ measures $\sin 2\beta$, where $\beta = \arg(-V_{cd}V_{cb}^*/V_{td}V_{tb}^*)$ and V_{ij} denote the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements [1,2]. Moreover, being dominated by the tree-level transition $b \to c\bar{c}s$, the decay amplitude of $B_d \to \psi K_S$ is unlikely to receive significant corrections from new physics.[1] Interestingly, within the standard model the CP asymmetry in $B_d \to \phi K_S$ $[a_{CP}(\phi K_S)]$ also measures $\sin 2\beta$ if, as naively expected, the decay amplitude is dominated by the short-distance penguin transition $b \to s\bar{s}s$ [4]. Since $B_d \to \phi K_S$ is a loop mediated process within the standard model, it is not unlikely that new physics could have a significant effect on it [3]. The expected branching ratio and the high identification efficiency for this decay suggests that $a_{CP}(\phi K_S)$ is experimentally accessible at the early stages of the asymmetric B factories. Thus, the search for a difference between $a_{CP}(\psi K_S)$ and $a_{CP}(\phi K_S)$ is a promising way to look for physics beyond the standard model [3,5–8].

If, indeed, it turns out that $a_{CP}(\psi K_S)$ is not equal to $a_{CP}(\phi K_S)$, it would be extremely important to know how precise the standard model prediction of them being equal is. In particular, one has to rule out the possibility of unexpected long distance effects altering the prediction that $a_{CP}(\phi K_S)$ measures $\sin 2\beta$ in the standard model.

The weak phases of the transition amplitudes are ruled by products of CKM matrix elements. In the $b \to s q\bar{q}$ case, relevant to both $B_d \to \psi K_S$ and $B_d \to \phi K_S$, we denote these by $\lambda_q^{(s)} = V_{qb}V_{qs}^*$. For the purpose of CP violation studies, it is instructive to use CKM unitarity and express any decay amplitude as a sum of two terms [9]. In particular, for $b \to s q\bar{q}$ we eliminate $\lambda_t^{(s)}$ and write

$$A_f = \lambda_c^{(s)} A_f^{cs} + \lambda_u^{(s)} A_f^{us}. \tag{1}$$

The unitarity and the experimental hierarchy of the CKM matrix imply [10] $\lambda_c^{(s)} \simeq -\lambda_t^{(s)} = A\lambda^2 + \mathcal{O}(\lambda^4)$ and $\lambda_u^{(s)} = AR_u\lambda^4 e^{i\gamma} + \mathcal{O}(\lambda^6)$, where $A \approx 0.8$, $\lambda = \sin\theta_c = 0.22$, $R_u \equiv |V_{ub}^* V_{ud}/V_{cb}^* V_{cd}| \sim 0.2–0.5$ [2] and γ is a phase of order one. Thus the first and dominant term is real (we work in the standard parametrization). The correction due to the second term, that is complex and doubly Cabibbo suppressed, is negligibly small unless $A_f^{us} \gg A_f^{cs}$.

The A_f^{qs} amplitudes cannot be calculated since they depend on hadronic matrix elements. However, in some cases we can reliably estimate their relative sizes. For $B \to \psi K_S$ the dominant term includes a tree level diagram while the CKM-suppressed term contains only one-loop (penguin) and higher order diagrams. This leads to $A_{\psi K_S}^{cs} \gg A_{\psi K_S}^{us}$, and thus insures that $a_{CP}(\psi K_S)$ measures $\sin 2\beta$ in the standard model. Since both terms for $B \to \phi K_S$ begin at one-loop order one expects $A_{\phi K_S}^{cs} \sim A_{\phi K_S}^{us}$. This is also supported by various model calculations [11]. However, since $A_{\phi K_S}^{cs}$ and $A_{\phi K_S}^{us}$ are hadronic matrix elements that depend on long distance dynamics, there is no firm theoretical argument that the two should be approximately equal. In the case that $A_{\phi K_S}^{cs} \sim A_{\phi K_S}^{us}$, $a_{CP}(\phi K_S)$ also measures $\sin 2\beta$ in the standard model up to corrections of $\mathcal{O}(\lambda^2)$. However, any unexpected enhancement of $A_{\phi K_S}^{us}$ would violate this result. In particular, an enhancement of $\mathcal{O}(\lambda^{-2}) \sim 25$ (analogous to the $\Delta I = 1/2$ rule in K decays) leads to $\mathcal{O}(1)$ violations, and subsequently to $a_{CP}(\psi K_S) \neq a_{CP}(\phi K_S)$ even in the standard model.

[1] There is, of course, a possible new contribution to the $B^0 - \bar{B}^0$ mixing amplitude. This does not affect the generality of our arguments or the conclusions [3].

In this note we argue against this possibility, presenting different arguments that suggest the pollution of $A^{us}_{\phi K_S}$ in $B_d \to \phi K_S$ is very small. Moreover, we will propose some experimental tests that in the near future could provide quantitative bounds on this pollution.

The natural tool to describe the B decays of interest is by means of an effective $b \to s\bar{q}q$ Hamiltonian. This can be generally written as

$$\mathcal{H}^{(s)}_{eff} = \frac{G_F}{\sqrt{2}} \left\{ \lambda^{(s)}_t \sum_{k=3..10} C_k(\mu) Q^s_k + \lambda^{(s)}_c \sum_{k=1,2} C_k(\mu) Q^{cs}_k + \lambda^{(s)}_u \sum_{k=1,2} C_k(\mu) Q^{us}_k \right\}, \quad (2)$$

where Q^i_k denote the local four fermion operators and $C_k(\mu)$ the corresponding Wilson coefficients, to be evaluated at a renormalization scale $\mu \sim \mathcal{O}(m_b)$. For our discussion it is useful to emphasize the flavor structure of the operators: $Q^{qs}_{1,2} \sim \overline{b}s\overline{q}q$ and $Q^s_{3..8} \sim \overline{b}s \Sigma_{q=u,d,s,c} \overline{q}q$, as well as the order of magnitude of their Wilson coefficients: $C_{1,2} \sim \mathcal{O}(1)$ and $C_{3..8} \sim \mathcal{O}(10^{-2})$. The estimates of the $C_k(\mu)$ beyond the leading logarithmic approximation and the definitions of the Q^i_k, can be found in [12]. To an accuracy of $O(\lambda^2)$ in the weak phases, $\mathcal{H}^{(s)}_{eff}$ can be rewritten as

$$\mathcal{H}^{(s)}_{eff} = \frac{G_F}{\sqrt{2}} \left\{ \lambda^{(s)}_c \left[\sum_{k=1,2} C_k(\mu) Q^{cs}_k - \sum_{k=3..10} C_k(\mu) Q^s_k \right] + \lambda^{(s)}_u \sum_{k=1,2} C_k(\mu) Q^{us}_k \right\}. \quad (3)$$

It is clear that, when sandwiched between the B_d initial state and the ϕK_S final state, the first term corresponds to $A^{cs}_{\phi K_S}$ and the second to $A^{us}_{\phi K_S}$ [cf. Eq. (1)]. The pollution is then generated by $Q^{us}_{1,2}$, corresponding to the $b \to s\bar{u}u$ transition.

Since the matrix elements of the Q^i_k have to be evaluated at $\mu \sim \mathcal{O}(m_b)$, a realistic estimate of their relative sizes can be obtained within perturbative QCD. We recall that the $|\phi\rangle$ is an almost pure $|\overline{s}s\rangle$ state. The $\omega - \phi$ mixing angle is estimated to be below 5% [13,2]. We neglect this small mixing in the following. Then, the matrix elements of $Q^{us}_{1,2}$ and $Q^{cs}_{1,2}$ evaluated at the leading order (LO) in the factorization approximation are identically zero. At LO only $Q_{3..8}$, i.e. the short-distance $b \to s\bar{s}s$ penguins, have a nonvanishing matrix element in $B_d \to \phi K_S$. As a consequence, the weak phase of the $B_d \to \phi K_S$ decay amplitude is essentially zero. Nonetheless, given the large Wilson coefficients of $Q^{qs}_{1,2}$, a more accurate estimate of their contribution is required.

At next to leading order (NLO), working in a modified factorization approximation, one obtains additional contributions from penguin-like matrix elements of the operators Q^{us}_2 and Q^{cs}_2 [14]. These have been reevaluated recently, and shown to be important in explaining the CLEO data on charmless two-body B decays [15–17]. However, even in this case the $b \to s\bar{u}u$ pollution in $B_d \to \phi K_S$ is very small. The reason is that, in the limit where we can neglect both the charm and the up quark masses with respect to m_b, the matrix elements of $Q^{us}_{1,2}$ and $Q^{cs}_{1,2}$ are identical from the point of view of perturbative QCD [up to corrections of $\mathcal{O}(m_c/m_b) \sim 0.3$]. However, the overall contribution of the charm operators $Q^{cs}_{1,2}$ is enhanced by a factor λ^{-2} with respect to the one of $Q^{us}_{1,2}$. Thus, either if the $B_d \to \phi K_S$ transition is dominated by Q^s_{3-10} (short-distance penguins) or if it is dominated by $Q^{cs}_{1,2}$ (long-distance charming penguins), the weak phase is vanishingly small.

Of course one could not exclude *a priori* a scenario where the contributions of $Q^s_{3..8}$ and $Q^{cs}_{1,2}$ cancel each other to an accuracy of $O(\lambda^2)$. However, this extremely unlikely possibility would result in an unobservably small $B(B_d \to \phi K_S)$, rendering this entire discussion moot.

As discussed above, any enhancement of $\langle \phi K_S | Q^{us}_{1,2} | B_d \rangle$, that could spoil the prediction that $a_{CP}(\phi K_S)$ measures $\sin 2\beta$ in the standard model should occur at low energies in order not to be compensated by a corresponding enhancement of $\langle \phi K_S | Q^{cs}_{1,2} | B_d \rangle$. This possibility is not only disfavored by the OZI rule [18],[2] but is also suppressed by the smallness of the energy range where the enhancement should occur with respect to the scale of the process. We are not aware of any dynamical mechanism that could favor this scenario. Inelastic rescattering effects in B decays due to Pomeron exchange have been argued not to be negligible and to violate the factorization limit [20]. However, even within this context violations of the OZI rule are likely to be suppressed [21].

There are experimental tests of our arguments that can be achieved in the sector of $b \to d$ transitions. These are described by an effective Hamiltonian $\mathcal{H}^{(d)}_{eff}$ completely similar to the one in Eq. (2) except for the substitution $s \to d$ in the flavor indices of both CKM factors and four-fermion operators. SU(3) flavor symmetry can be used to obtain relation among several matrix elements [22]. In particular

$$\sqrt{2} \langle \phi K_S | Q^{us}_{1,2} | B_d \rangle = \langle \phi \pi^+ | Q^{ud}_{1,2} | B^+ \rangle + \langle K^* K^+ | Q^{ud}_{1,2} | B^+ \rangle. \quad (4)$$

[SU(3) breaking effects, which are typically at the 30% level, are neglected here.] The coefficients of these matrix elements are, however, proportional to different CKM factors. This is illustrated in Table I, where we show the relevant B decay modes along with the Cabibbo factors corresponding to the leading and sub-leading contributions to the decay amplitudes. If our arguments hold, one expects $B(B_d \to \phi K_S) \sim \mathcal{O}(\lambda^4)$ and $B(B^+ \to K^* K^+)$, $B(B^+ \to \phi \pi^+) \sim \mathcal{O}(\lambda^6)$. Notice, however, that the overall contribution of $Q^{ud}_{1,2}$ in $B^+ \to K^* K^+$ and $B^+ \to \phi \pi^+$ is enhanced with respect to the one of $Q^{us}_{1,2}$ in $B_d \to \phi K_S$ by the corresponding CKM factors:

[2]This non-perturbative prescription has never been fully understood in the framework of perturbative QCD, but can be justified in the framework of the $1/N_c$ expansion, and is known to work well in most cases and particularly in the vector meson sector [19].

TABLE I. SU(3) related B decay modes that allow us to quantify the standard model pollution in $a_{CP}(\phi K_S)$.

	Operators and CKM factors		
Decay mode	penguins	c trees	u trees
$B_d \to \phi K_S$	$Q^s_{3..8}$	$Q^{cs}_{1,2}$	$Q^{us}_{1,2}$
	$\lambda_t^{(s)} \sim \lambda^2$	$\lambda_c^{(s)} \sim \lambda^2$	$\lambda_u^{(s)} \sim \lambda^4$
$B^+ \to \phi\pi^+$ and $B^+ \to K^*K^+$	$Q^d_{3..8}$	$Q^{cd}_{1,2}$	$Q^{ud}_{1,2}$
	$\lambda_t^{(d)} \sim \lambda^3$	$\lambda_c^{(d)} \sim \lambda^3$	$\lambda_u^{(d)} \sim \lambda^3$

$\lambda_u^{(d)}/\lambda_u^{(s)} = \mathcal{O}(\lambda^{-1})$. Thus, if $\langle \phi K_S | Q^{us}_{1,2} | B_d \rangle$ is enhanced by $\mathcal{O}(\lambda^{-2})$ in order to interfere with the dominant $\mathcal{O}(\lambda^2)$ contributions, then $B(B^+ \to \phi\pi^+)$ and/or $B(B^+ \to K^*K^+)$ would be dominated by the similarly enhanced matrix elements of $Q^{ud}_{1,2}$. This would result in an enhancement of the naively Cabibbo suppressed modes, i.e. we should observe $B(B^+ \to \phi\pi^+) \sim \mathcal{O}(\lambda^2)$ and/or $B(B^+ \to K^*K^+) \sim \mathcal{O}(\lambda^2)$ [while $B(B_d \to \phi K_S)$ is still $\sim \mathcal{O}(\lambda^4)$]. Similar arguments hold for the corresponding B_d decay modes, however in that case the SU(3) relation is not quite as precise.

To get a quantitative bound we define the ratios

$$R_1 = \frac{B(B^+ \to \phi\pi^+)}{B(B_d \to \phi K_S)}, \quad R_2 = \frac{B(B^+ \to K^*K^+)}{B(B_d \to \phi K_S)}, \qquad (5)$$

such that in the standard model the following inequality holds

$$|a_{CP}(\psi K_S) - a_{CP}(\phi K_S)| < \sqrt{2}\lambda(\sqrt{R_1} + \sqrt{R_2}) \times [1 + R_{SU(3)}] + \mathcal{O}(\lambda^2), \quad (6)$$

where $R_{SU(3)}$ represents the SU(3) breaking effects. While measuring $a_{CP}(\phi K_S)$ it should be possible to set limits at least of order one on R_1 and R_2 and thus to control by means of Eq. (6) the accuracy to which $a_{CP}(\phi K_S)$ measures $\sin 2\beta$ in the standard model. The limits $\sqrt{R_1}, \sqrt{R_2} \lesssim 0.25$ would reduce the theoretical uncertainty to the 10% level.

It may be possible to confirm that $B(B^+ \to \phi\pi^+)$ and $B(B^+ \to K^*K^+)$ are not drastically enhanced based just on the current CLEO data. The CLEO collaboration already has reported the bounds $B(B^+ \to \phi\pi^+) < 0.56 \times 10^{-5}$ [23] and $B(B^+ \to K^*\pi^+) < 4.1 \times 10^{-5}$ [24]. Given the similarity of energetic K's and π's in the CLEO environment, it is plausible that a bound similar to the latter can also be derived for the mode $B^+ \to K^*K^+$. Bounds on these branching ratios of $\mathcal{O}(10^{-5})$ would clearly imply that the rates are not $\mathcal{O}(\lambda^2)$ as they would be if the matrix elements of $Q^{ud}_{1,2}$ were enhanced by $\mathcal{O}(\lambda^{-2})$.

The above experimental test can only confirm that $a_{CP}(\phi K_S)$ measures $\sin 2\beta$ in the standard model. If it turns out that R_1 or R_2 is large, this may be either due to the failure of our conjectures or due to new physics. If, however, R_1 and R_2 are small, and $a_{CP}(\psi K_S) - a_{CP}(\phi K_S)$ violates the standard model prediction of Eq. (6), this would be an unambiguous sign of new physics.

Another possible check of our conjecture could be achieved through the measurement of the CP asymmetry in $B_d \to \eta' K_S$. Recently CLEO has measured a large branching ratio for the related decay $B^+ \to \eta' K^+$, suggesting these processes are penguin dominated and thus that $a_{CP}(\eta' K_S)$ also should measure $\sin 2\beta$ in the standard model [7]. Nonetheless, the $|\eta'\rangle$ has a non-negligible $|\bar{u}u\rangle$ component that could enhance the $b \to u\bar{u}s$ pollution and the η' mass is one of the few exception where the OZI rule is known to be badly broken. Thus, without fine tuning, a sufficient condition to support our claim on $a_{CP}(\phi K_S)$ could be obtained by an experimental evidence of $a_{CP}(\eta' K_S) = a_{CP}(\phi K_S)$. This would imply that the $b \to u\bar{u}s$ pollution is negligible in both cases.

To summarize, we have argued that the deviation from the prediction that $a_{CP}(\phi K_S)$ measures $\sin 2\beta$ in the standard model is of $\mathcal{O}(\lambda^2) \sim 5\%$. Moreover, we have shown how the accuracy of this prediction can be tested experimentally. While we concentrated on the time-dependent CP asymmetry it is clear that our arguments hold also for direct CP violation in charged and neutral $B \to \phi K$ decays. Namely, that in the standard model the direct CP asymmetry is $\mathcal{O}(\lambda^2)$. Experimentally, we can hope to get an accuracy for both the time dependent and the direct CP violation of about 10%. Therefore, any measurable direct CP violation in $B \to \phi K$ or an indication that $a_{CP}(\psi K_S) \neq a_{CP}(\phi K_S)$, combined with experimental evidence that the standard model pollution is of $\mathcal{O}(\lambda^2)$ will signal physics beyond the standard model.

We thank Ann Nelson for instigating this investigation. We have benefitted from useful conversations with Roy Briere, Gerhard Buchalla, Bob Cahn, Alex Kagan, David Kaplan, Marek Karliner, Sacha Kop and Yossi Nir. The research at SLAC supported by the Department of Energy under contract DE-AC03-76SF00515.

[1] For a review see, e.g., Y. Nir, Lectures presented in the 20th SLAC Summer Institute, SLAC-PUB-5874 (1992); Y. Nir and H. R. Quinn, Annu. Rev. Nucl. Part. Sci. **42**, 211 (1992).
[2] Particle Data Group, R. M. Barnett *et al.*, Phys. Rev. D **54**, 1 (1996).
[3] Y. Grossman and M. Worah, Phys. Lett. B **395**, 241 (1997), and in 2nd International Conference on B Physics and CP violation, hep-ph/9707280.
[4] D. London and R. D. Peccei, Phys. Lett. B **223**, 257 (1989).
[5] Y. Nir and H. Quinn, Phys. Rev. D **42**, 1473 (1990).
[6] M. Ciuchini *et al.*, Phys. Rev. Lett. **79**, 978 (1997).
[7] D. London and A. Soni, Phys. Lett. B **407**, 61 (1997).
[8] R. Barbieri and A. Strumia, Nucl. Phys. **B508**, 3 (1997).
[9] Y. Grossman and H. R. Quinn, Phys. Rev. D **56**, 7259 (1997).

[10] L. Wolfenstein, Phys. Rev. Lett. **51**, 1945 (1983).
[11] See, e.g., G. Kramer, W. F. Palmer, and H. Simma, Z. Phys. C **66**, 429 (1995).
[12] A. Buras and R. Fleischer, TUM-HEP-275-97, hep-ph/9704376.
[13] R. S. Chivukula and M. Flynn, Phys. Lett. B **186**, 127 (1986); P. Jain *et al.*, Phys. Rev. D **37**, 3252 (1988).
[14] R. Fleischer, Z. Phys. C **58**, 483 (1993).
[15] M. Ciuchini *et al.*, Nucl. Phys. **B501**, 271 (1997); **B512**, 3 (1998).
[16] R. Fleischer and T. Mannel, Phys. Rev. D **57**, 2752 (1998).
[17] A. Lenz, U. Nierste, and G. Ostermeyer; Phys. Rev. D **56**, 7228 (1997).
[18] S. Okubo, Phys. Lett. **5**, 165 (1963); G. Zweig, CERN Reports TH-401 and TH-412 (1964); J. Iiuzuka, Prog. Theor. Part. Sci. **37**, 21 (1996).
[19] P. Geiger and N. Isgur, Phys. Rev. Lett. **67**, 1066 (1991); Phys. Rev. D **47**, 5059 (1993).
[20] J. Donoghue *et al.*, Phys. Rev. Lett. **77**, 2178 (1996).
[21] See, e.g., G. Veneziano, in *Color Symmetry and Quark Confinement*, edited by T. Thanh Van, Proceedings of the 12th Rencontré de Moriond (1977).
[22] See, for example, A. Dighe, M. Gronau, and J. Rosner, Phys. Rev. D **57**, 1783 (1998).
[23] CLEO Collaboration, M. S. Alam *et al.*, CLEO CONF 97-23.
[24] CLEO Collaboration, D. Asner *et al.*, Phys. Rev. D **53**, 1039 (1996).