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The pattern of CP asymmetries in $b \to s$ transitions

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Abstract: New CP violating physics in $b \to s$ transitions will modify the CP asymmetries in B decays into final CP eigenstates (ϕK_S, $\eta' K_S$, $\pi^0 K_S$, ωK_S, $\rho^0 K_S$ and ηK_S) from their Standard Model values. In a model independent analysis, the pattern of deviations can be used to probe which Wilson coefficients get a significant contribution from the new physics. We demonstrate this idea using several well-motivated models of new physics, and apply it to current data.

Keywords: Rare Decays, Beyond Standard Model, B-Physics, CP violation.

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Contents

1. Introduction

CP asymmetries in the decays of neutral B mesons into final CP eigenstates f exhibit a time-dependent behavior (for a review of CP violation in meson decays, see [1]),

$$\mathcal{A}_f(t) = S_f \sin(\Delta m t) - C_f \cos(\Delta m t). \tag{1.1}$$

The Standard Model predicts that for most of the decays that proceed via the quark transitions $\bar{b} \to \bar{q}' q' \bar{s}$ ($q' = c, s, d, u$), the following relations hold to a good approximation:

$$-\eta_f S_f \simeq \sin 2\beta\,, \qquad C_f \simeq 0\,, \tag{1.2}$$

where $\eta_f = \pm 1$ is the CP eigenvalue for the final state f, and $\beta \equiv \phi_1 \equiv \arg[-\frac{(V_{cd}V_{cb}^*)}{(V_{td}V_{tb}^*)}]$. New physics effects can appear in two ways. First, new physics in $\mathcal{H}_{\Delta B=2}$ contributes to the $B^0 - \overline{B}^0$ mixing amplitude. Such a contribution shifts all S_f's in a universal way: The S_f asymmetries remain equal to each other, though different from $\sin 2\beta$. The C_f's still vanish. Second, new physics in $\mathcal{H}_{\Delta B=1}$ can contribute to the $\bar{b} \to \bar{q} q \bar{s}$ transitions ($q = s, d, u$) [2]. (The tree level transition $\bar{b} \to \bar{c} c \bar{s}$ is unlikely to get a significant contribution from new physics, and consequently the asymmetries $S_{\psi K}$ and $C_{\psi K}$ would not be modified by $\mathcal{H}_{\Delta B=1}$.) Such a contribution would lead to interesting consequences:

1. The S_f's would be different from each other and from $S_{\psi K}$.
2. The C_f's would be different from each other and from zero.

Thus, in the presence of new physics, a pattern of deviations, $\delta S_f = -\eta_f S_f - S_{\psi K} \neq 0$ and $C_f \neq 0$, will arise.

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f	$-\eta_{\rm CP} S_f$	C_f	refs.
ψK_S	$+0.73 \pm 0.04$	$+0.03 \pm 0.03$	[6, 7]
ϕK_S	$+0.35 \pm 0.21$	-0.04 ± 0.17	[8, 9]
$\eta' K_S$	$+0.43 \pm 0.11$	-0.04 ± 0.08	[10, 9]
$\pi^0 K_S$	$+0.34 \pm 0.29$	$+0.08 \pm 0.14$	[11, 9]
ωK_S	$+0.56 \pm 0.32$	-0.49 ± 0.25	[9, 12]

Table 1: World averages of the experimental results for the CP asymmetries in neutral B decays via the quark transitions $\bar{b} \to \bar{q}' q' \bar{s}$ modes, $q' = c, s, d, u$.

Given a set of experimental ranges for δS_f and C_f, one would like to interpret these data in terms of new physics. There are two ways to proceed. First, one can choose a model of new physics and analyze whether the model can accommodate the data. The second way is model independent. The effects of new physics can be described by a modification of the Wilson coefficients in the operator product expansion for $\Delta B = 1$ interactions. Thus, one can fit a set of Wilson coefficients to the data and learn which operators can account for the observed deviations. In this paper, we study this second, model independent, method.

While our analysis has the advantage of being model independent, it suffers from two limitations:

- The number of Wilson coefficients is larger than the number of measured CP asymmetries. Consequently, an analysis that is entirely generic is impossible to carry out at present. Hence, one has to assume that only a subset of all possible operators are modified. In this work we consider some simple cases, where the new physics is parameterized by a single complex parameter. The cases that we study are motivated by actual models of new physics. Extensions to a larger number of new physics parameters is left for future work. It is also possible to include in the analysis additional data, beyond the CP asymmetries in decays into final CP eigenstates. This extension is also left for future work.

- To interpret the data in terms of Wilson coefficients, one has to know the values of the (mode-dependent) hadronic matrix elements of the operators. At present, there is no first principle calculation of the matrix elements that has been tested to a high level of precision. Thus, hadronic uncertainties prevent a clean theoretical interpretation. We use factorization [3 − 5] for our analysis.[1]

A number of relevant measurements exists already. The experimental situation is summarized in table 1.

[1] In this work we perform our calculations to leading order in α_s and drop $1/m_b$ corrections, except for the chirally enhanced terms related to a_6 and a_8. In this approximation, QCD- and naive-factorization are identical. We note, however, differences in the expressions for the $f = \eta^{(\prime)} K$ modes between refs. [4] and [5]: A term $\propto f^u_{\eta^{(\prime)}}(a_6 - \frac{1}{2}a_8)$ should be omitted from the first, while a term $\propto f^c_{\eta^{(\prime)}}(a_7 - a_9)$ should be added to the second.

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The plan of this paper is as follows. In section 2 we introduce our formalism and evaluate, to leading log approximation and within the QCD factorization approach, the Standard Model (SM) predictions for the CP asymmetries. In section 3 we analyze the effects of new physics. We start with a generic, model independent analysis, using the operator product expansion for $\Delta B = 1$ operators. We then focus on scenarios where the effects of new physics depend on a single complex parameter. We give three explicit examples of such scenarios. In section 4 we apply our approach to present data. We conclude in section 5.

2. Formalism and standard model predictions

We follow the notations of ref. [3]. We consider the following effective hamiltonian for $\Delta B = \pm 1$ decays:

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} V_{ps}^* V_{pb} \left(C_1 O_1^p + C_2 O_2^p + \sum_{i=3}^{10} C_i O_i + C_{7\gamma} O_{7\gamma} + C_{8g} O_{8g} \right) + \text{h.c.}\,, \tag{2.1}$$

with

$$\begin{aligned}
O_1^p &= (\bar{p}b)_{V-A}(\bar{s}p)_{V-A}\,, & O_2^p &= (\bar{p}_\beta b_\alpha)_{V-A}(\bar{s}_\alpha p_\beta)_{V-A}\,,\\
O_3 &= (\bar{s}b)_{V-A}\sum_q (\bar{q}q)_{V-A}\,, & O_4 &= (\bar{s}_\alpha b_\beta)_{V-A}\sum_q (\bar{q}_\beta q_\alpha)_{V-A}\,,\\
O_5 &= (\bar{s}b)_{V-A}\sum_q (\bar{q}q)_{V+A}\,, & O_6 &= (\bar{s}_\alpha b_\beta)_{V-A}\sum_q (\bar{q}_\beta q_\alpha)_{V+A}\,,\\
O_7 &= \frac{3}{2}(\bar{s}b)_{V-A}\sum_q e_q (\bar{q}q)_{V+A}\,, & O_8 &= \frac{3}{2}(\bar{s}_\alpha b_\beta)_{V-A}\sum_q e_q (\bar{q}_\beta q_\alpha)_{V+A}\,,\\
O_9 &= \frac{3}{2}(\bar{s}b)_{V-A}\sum_q e_q (\bar{q}q)_{V-A}\,, & O_{10} &= \frac{3}{2}(\bar{s}_\alpha b_\beta)_{V-A}\sum_q e_q (\bar{q}_\beta q_\alpha)_{V-A}\,,\\
O_{7\gamma} &= -\frac{e m_b}{8\pi^2}\bar{s}\sigma^{\mu\nu}(1+\gamma_5)F_{\mu\nu}b\,, & O_{8g} &= -\frac{g_s m_b}{8\pi^2}\bar{s}\sigma^{\mu\nu}(1+\gamma_5)G_{\mu\nu}b\,,
\end{aligned} \tag{2.2}$$

where $(\bar{q}_1 q_2)_{V\pm A} = \bar{q}_1\gamma_\mu(1\pm\gamma_5)q_2$, the sum is over active quarks, with e_q denoting their electric charge in fractions of $|e|$ and α, β are color indices.

The CP asymmetries in $B \to f$ decays are calculated as follows. One defines a complex quantity λ_f,

$$\lambda_f = e^{-i\phi_B}\left(\frac{\overline{A}_f}{A_f}\right), \tag{2.3}$$

where ϕ_B is the phase of M_{12}, the $B^0 - \overline{B}^0$ mixing amplitude, and A_f ($\overline{A}_f$) is the decay amplitude for $B^0(\overline{B}^0) \to f$. We have

$$S_f = \frac{2\mathcal{I}m(\lambda_f)}{1+|\lambda_f|^2}\,, \qquad C_f = \frac{1-|\lambda_f|^2}{1+|\lambda_f|^2}\,. \tag{2.4}$$

The decay amplitudes can be calculated from the effective hamiltonian of eq. (2.1) [4]:

$$A_f = \langle f|\mathcal{H}_{\rm eff}|B^0\rangle\,, \qquad \overline{A}_f = \langle f|\mathcal{H}_{\rm eff}|\overline{B}^0\rangle\,. \tag{2.5}$$

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-2	6	$-\frac{2}{9}$	$\frac{2}{3}$	$-\frac{2}{9}$	$\frac{2}{3}$	0	0	0	0	$\frac{416}{81}$	$\frac{70}{27}$
6	-2	0	0	0	0	0	0	0	0	0	3
0	0	$-\frac{22}{9}$	$\frac{22}{3}$	$-\frac{4}{9}$	$\frac{4}{3}$	0	0	0	0	$-\frac{464}{81}$	$\frac{545}{27}$
0	0	$\frac{44}{9}$	$\frac{4}{3}$	$-\frac{10}{9}$	$\frac{10}{3}$	0	0	0	0	$\frac{136}{81}$	$\frac{512}{27}$
0	0	0	0	2	-6	0	0	0	0	$\frac{32}{9}$	$-\frac{59}{3}$
0	0	$-\frac{10}{9}$	$\frac{10}{3}$	$-\frac{10}{9}$	$-\frac{38}{3}$	0	0	0	0	$-\frac{296}{81}$	$-\frac{703}{27}$
0	0	0	0	0	0	2	-6	0	0	$-\frac{16}{9}$	$\frac{5}{6}$
0	0	$-\frac{1}{9}$	$\frac{1}{3}$	$-\frac{1}{9}$	$\frac{1}{3}$	0	-16	0	0	$-\frac{1196}{81}$	$-\frac{11}{54}$
0	0	$\frac{2}{9}$	$-\frac{2}{3}$	$\frac{2}{9}$	$-\frac{2}{3}$	0	0	-2	6	$\frac{232}{81}$	$-\frac{59}{54}$
0	0	$-\frac{1}{9}$	$\frac{1}{3}$	$-\frac{1}{9}$	$\frac{1}{3}$	0	0	6	-2	$\frac{1180}{81}$	$-\frac{46}{27}$
0	0	0	0	0	0	0	0	0	0	$\frac{32}{3}$	0
0	0	0	0	0	0	0	0	0	0	$-\frac{32}{9}$	$\frac{28}{3}$

Table 2: The anomalous dimension matrix at leading log. The rows and columns correspond to the twelve operators $O_{1,\ldots,10,7\gamma,8g}$ defined in eq. (2.2).

The electroweak model determines the Wilson coefficients while QCD (or, more practically, a calculational method such as QCD factorization) determines the matrix elements $\langle f|O_i|B^0(\overline{B}^0)\rangle$.

We perform our calculations to leading-log approximation. In particular, to run the Wilson coefficients from the weak scale m_W to a low scale of order m_b, we use the 12-dimensional leading-log anomalous dimension matrix γ [13] that is given in table 2. The α_s mixing of electroweak penguins onto the dipole operators is deduced from [13–15].

Within the SM, we have the following set of Wilson coefficients at leading order:

$$C_1^{\rm SM}(m_W) = 1\,, \qquad C_{i\neq 1}^{\rm SM}(m_W) = 0\,. \tag{2.6}$$

(Strictly speaking, $C_{7\gamma}^{\rm SM}(m_W)$ and $C_{8g}^{\rm SM}(m_W)$ are also different from zero. However, their contributions to the decay processes of interest occur at next-to-leading order which we neglect in the present work.) The SM contribution to the decay amplitudes, related to $\bar b \to \bar q' q' \bar s$ transitions, can always be written as a sum of two terms, $A_f^{\rm SM} = A_f^c + A_f^u$, with $A_f^c \propto V_{cb}^* V_{cs}$ and $A_f^u \propto V_{ub}^* V_{us}$. Defining the ratio $a_f^u \equiv e^{-i\gamma}(A_f^u/A_f^c)$, where $\gamma \equiv \phi_3 \equiv \arg\left[-(V_{ud}V_{ub}^*)/(V_{cd}V_{cb}^*)\right]$, we have

$$A_f^{\rm SM} = A_f^c\left(1 + a_f^u e^{i\gamma}\right). \tag{2.7}$$

For $B \to \psi K_S$, the a_f^u term can be safely neglected (its effects are below the percent level) and, consequently,

$$\lambda_{\psi K_S}^{\rm SM} = -e^{-2i\beta}\,. \tag{2.8}$$

For charmless modes, the effects of the a_f^u terms (often called 'the SM pollution') are at least of order $|(V_{ub}V_{us}^*)/(V_{cb}V_{cs}^*)| \sim$ a few percent. They can lead to a deviation of S_f from $S_{\psi K}$ and of C_f from zero. In table 3 we give the values of the a_f^u parameter (obtained, as explained above, by using factorization [3–5]) for all relevant modes.

f	a_f^u	$-\eta_{\rm CP} S_f$
ψK_S	0	0.73
ϕK_S	0.019	0.75
$\pi^0 K_S$	$0.052\ [0.094, 0.021]$	$0.79\ [0.83, 0.76]$
ηK_S	$0.08\ [0.16, 0.02]$	$0.82\ [0.88, 0.76]$
$\eta' K_S$	$0.007\ [-0.006, 0.019]$	$0.74\ [0.72, 0.75]$
ωK_S	$0.22\ [0.37, 0.04]$	$0.92\ [0.98, 0.78]$
$\rho^0 K_S$	$-0.16\ [-0.32, 0.005]$	$0.49\ [0.19, 0.74]$

Table 3: The a_f^u parameters, calculated in QCD factorization at leading log and to zeroth order in Λ/m_b (except for chirally enhanced corrections), and the SM values of S_f for $\mu = m_b$ and in parentheses the respective values for $\mu = 2m_b$ (first) and $\mu = m_b/2$ (second) if different from the central one.

Decay constants (MeV)	
f_π	131
f_K	160
f_η^q	108
f_η^s	-111
$f_{\eta'}^q$	89
$f_{\eta'}^s$	136
f_ω	187
f_ρ	209

Form Factors	
$F^{B\to\pi}(0)$	0.28
$F^{B\to K}(0)$	0.34
$F^{B\to\eta}(0)$	0.23
$F^{B\to\eta'}(0)$	0.19
$A^{B\to\omega}(0)$	0.33
$A^{B\to\rho}(0)$	0.37

r_χ-factors, quark masses	
$r_\chi(m_b)$	1.170
$r_\chi^{\eta_s}(m_b)$	1.230
$r_\chi^{\eta'_s}(m_b)$	1.241
$m_b(m_b)$	4.2 GeV
$m_s(2\text{GeV})$	110 MeV
$(m_u + m_d)(2\text{GeV})$	9.1 MeV

Table 4: Input parameters used in the calculations.

Since we perform our calculations to leading log approximation, we neglect, for consistency, also power corrections within the QCD factorization approach, except for the dominant chirally-enhanced ones related to a_6 and a_8. This means that the strong phases vanish, and that $C_f = 0$ for all f in this approximation. The SM predictions for S_f (with $S_{\psi K_S} = 0.73$ and $\gamma = 62^o$ [16] taken as input) are also given in table 3. The values of the input parameters used are specified in table 4. To estimate our uncertainties, we give the values for $\mu = m_b = 4.2\,\text{GeV}$ and for a varied renormalization scale of $\mu = m_b/2$ and $\mu = 2m_b$. Note that the coefficients of the chirally enchanced terms, defined as $r_\chi = 2m_K^2/(m_b m_s)$ and $r_\chi^{S_s} = h_S^s/(f_S^s m_b m_s)$ for the singlets $S = \eta, \eta'$ (see [5] for details), evolve with the renormalization scale through the running ($\overline{\text{MS}}$-bar) quark masses.

An examination of table 3 shows that the SM pollution is small (that is, at the naively expected level of $|(V_{ub}V_{us}^*)/(V_{cb}V_{cs}^*)| \sim$ a few percent) for $f = \phi K_S$, $\eta' K_S$ and $\pi^0 K_S$. It is larger for $f = \eta K_S$, ωK_S and $\rho^0 K_S$. In these modes, a_f^u is enhanced because, within the QCD factorization approach, there is an accidental cancellation between the leading contributions to A_f^c. The reason for the suppression of the leading A_f^c piece in $f = \rho K$, ωK versus $f = \pi^0 K$ is that the dominant QCD-penguin coefficients a_4 and a_6 appear in $A^c_{(\rho,\omega)K}$ as $(a_4 - r_\chi a_6)$ and in $A^c_{\pi^0 K}$ as $(a_4 + r_\chi a_6)$. Since $r_\chi \simeq 1$ and, within the Standard Model,

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$a_4 \sim a_6$, there is a cancellation in $A^c_{(\rho,\omega)K}$ while there isn't one in $A^c_{\pi^0 K}$. The suppression for $A^c_{\eta K}$ with respect to $A^c_{\eta' K}$ has a different reason: it is due to the octet-singlet mixing, which causes destructive (constructive) interference in the $\eta(\eta')K$ penguin amplitude [17].

We stress that the numerical results presented above are often sensitive to the approximations that we make and to the values of the input parameters. We will refine them in the future by going beyond the leading log approximation and taking into account uncertainties in the input parameters other than the scale μ. At present they should be taken as indication to how close to $S_{\psi K_S}$ one should expect the various S_f's to be, but not as accurate predictions. Our findings are compatible with the next to leading order results in the SM given by [18].

We conclude this section by adding some general considerations concerning the accuracy of our approximation and the stability of our results. We discuss two main points. First, the validity of our approximation for branching ratios and direct CP asymmetries. Second, the presence of large power corrections in our analysis.

Regarding the first issue, there is a substantial effect from NLO corrections in factorization as far as the branching ratios are concerned. However, the impact of such NLO corrections is very moderate for the asymmetries S_f, which are the main focus of our study. This can be demonstrated explicitly in the SM by comparing our results (table 3) with the more complete, full-fledged NLO analysis of [18]. In fact, even the estimate of theoretical uncertainties, which we obtained by a variation of the renormalization scale, is in good agreement with [18], where also other sources of uncertainty are included. This clearly demonstrates that our approach gives a reasonable picture for the asymmetries S_f, including the issue of uncertainties and the stability against subleading effects.

In addition, our framework is still consistent with the observed direct CP asymmetries C_f, which are compatible with zero. This is the case in general for direct CP asymmetries in all B decays, where $A_{CP}(K^+\pi^-)$ is so far the only exception. Even in this case the observed direct CP asymmetry is not very large (about 10%).

Regarding the second issue, we stress again in this context that the dominant power corrections related to Q_6 and Q_8 are included in our analysis (as already mentioned in footnote 1). Further power corrections, in particular those from annihilation topologies can be estimated [3] to be typically of the order 10 - 20% with respect to the factorizable prediction. Hence the impact of such effects is subdominant. This is also confirmed by the comparison with [18] mentioned above.

3. New physics

3.1 Generic analysis

For the purpose of discussing new physics, it is convenient to define a phase $\beta_{\rm eff}$:

$$\lambda_{\psi K_S} \equiv -e^{-2i\beta_{\rm eff}}\,. \tag{3.1}$$

In writing down eq. (3.1), we are making the very plausible assumption that the quark transition $b \to c\bar{c}s$ is dominated by a single weak phase. This is clearly the case within

f	$\pi^0 K$	ηK	$\eta' K$	ϕK	ωK	$\rho^0 K$
b^c_{f1}	1	1	1	1	1	1
b^u_{f1}	0.052	0.08	0.0072	0.019	0.22	-0.16
b^c_{f2}	-0.12	-0.11	-0.13	-0.13	-0.096	-0.093
b^u_{f2}	0.32	0.59	-0.12	-0.0024	1.9	-1.8
b^c_{f3}	-0.34	17	-26	-46	210	9.3
b^c_{f4}	-12	-7.4	-17	-43	110	81
b^c_{f5}	-12	-28	12	-43	140	-55
b^c_{f6}	-31	-39	-21	-20	-51	-160
b^c_{f7}	-19	-32	3.8	22	80	-110
b^c_{f8}	3.6	-5.1	20	14	120	-1.6
b^c_{f9}	25	37	3.5	23	45	-140
b^c_{f10}	12	11	11	24	-36	-51

Table 5: The b^c_{fi} $(i = 1, \ldots, 10)$ and $b^u_{f1,2}$ parameters calculated in QCD factorization to leading log approximation at $\mu = m_b$ and to zeroth order in Λ/m_b (except for chirally enhanced corrections).

the SM, where the phase of $V_{cb}V^*_{cs}$ dominates to better than one percent. But even with new physics we do not expect significant new contributions to SM tree level processes. The approximation of a single weak phase may be not as good as in the SM, but it is still expected to be a very good approximation. On the other hand, eq. (3.1) allows for new physics in $B^0 - \overline{B}^0$ mixing, in which case $\beta_{\rm eff} \neq \beta$.

Contributions of new physics to $\mathcal{H}_{\rm eff}$ of eq. (2.1) modify the Wilson coefficients C_i. (New physics can also introduce additional operators. We do not consider this possibility here, but our analysis can be generalized to this case in a straightforward manner.) This modification can be parametrized as follows:

$$C_i(m_W) = C_i^{\rm SM}(m_W) + \varepsilon_i e^{i\theta_i}\,, \tag{3.2}$$

where $C_i^{\rm SM}$ is the Standard Model value of the Wilson coefficient, ε_i is real and positive, and θ_i is a phase in the range $[0, 2\pi]$.

The modification of A_f from the SM expression [eq. (2.7)] due to the new physics contribution [eq. (3.2)] can always be written as follows:

$$A_f = A^c_f \left[1 + a^u_f e^{i\gamma} + \sum_i \left(b^c_{fi} + b^u_{fi} e^{i\gamma}\right) \varepsilon_i e^{i\theta_i}\right]. \tag{3.3}$$

In table 5, we give the values of the b^c_{fi} parameters within the QCD factorization approach [3–5]. The b^u_{fi} for $i = 3, \ldots, 10$ are related to the b^c_{fi} through $b^u_{fi} = (\frac{|V_{ub}V^*_{us}|}{|V_{cb}V^*_{cs}|})b^c_{fi}$. Note that further $b^c_{f1} = 1$ and $b^u_{f1} = a^u_f$.

With new physics, the deviation of S_f of charmless modes from $S_{\psi K}$ depends on b^c_{fi}, b^u_{fi}, ε_i and θ_i. If enough measurements of S_f and C_f asymmetries become available, one will be able to find the values of the ε_i- and θ_i- parameters that account for the data. Given a small set of measurements, one can still find ε_i and θ_i if only a small number of operators is affected by the new physics in a significant way. In the next section, we focus on a simple (but well motivated) case in which the contribution of the new physics can be parametrized in terms of a single complex parameter.

3.2 A single complex parameter

Suppose that all the $\varepsilon_i e^{i\theta_i}$ parameters of eq. (3.2) are related to a single complex parameter, that is,

$$\varepsilon_i e^{i\theta_i} = x_i \varepsilon e^{i\theta} \qquad \text{for all } i\,, \tag{3.4}$$

where, in a given model of new physics, the x_i-constants are computable (see section 3.3 for three specific examples). In such cases, eq. (3.3) takes the following form:

$$A_f = A_f^c \left[1 + a_f^u e^{i\gamma} + \left(b_f^c + b_f^u e^{i\gamma}\right) \varepsilon e^{i\theta}\right]. \tag{3.5}$$

The analysis is simpler if $|b_f^c| \gg |b_f^u|$. This is the case for all generic combinations of operators O_i in eq. (2.2) that do not involve $i = 1, 2$. We proceed therefore in discussing the case where a combination of operators within the set $i = 3, \ldots, 10, 7\gamma, 8g$ is affected by new physics. For two classes of interesting scenarios, simple analytical expressions can be obtained for the modification of S_f and C_f.

(i) The new physics contribution is dominant. Explicitly,

$$\varepsilon |b_f^c| \gg |a_f^u|, 1\,. \tag{3.6}$$

The shift in all modes where eq. (3.6) holds is universal and depends only on θ:

$$\begin{aligned} -\eta_f S_f &\simeq \sin 2\beta_{\text{eff}} \cos 2\theta + \cos 2\beta_{\text{eff}} \sin 2\theta\,, \\ C_f &\simeq 0\,. \end{aligned} \tag{3.7}$$

(ii) The new physics contribution and the SM pollution are small. Explicitly,

$$\varepsilon |b_f^c| \ll 1 \qquad \text{and} \qquad |a_f^u| \ll 1\,. \tag{3.8}$$

The shift is mode dependent and depends on $\varepsilon \sin\theta$:

$$\begin{aligned} -\eta_f S_f &\simeq \sin 2\beta_{\text{eff}} + 2\cos 2\beta_{\text{eff}} \left[\mathcal{R}e(b_f^c)\varepsilon \sin\theta + \mathcal{R}e(a_f^u)\sin\gamma\right], \\ C_f &\simeq -2\mathcal{I}m(b_f^c)\varepsilon \sin\theta - 2\mathcal{I}m(a_f^u)\sin\gamma\,. \end{aligned} \tag{3.9}$$

Since the b_f^c values are known (see table 5), the pattern of deviations is predicted and can be used, by comparing with the data, to close in on the set of operators that is responsible for the deviations and to extract $\varepsilon \sin\theta$.

In case that none of these approximations applies, one can still find the pattern of deviations by numerical evaluation for any values for ε and θ. By comparing these theoretical δS_f and C_f to the data, allowed regions in the $\varepsilon - \theta$ plane (for each scenario) can be determined, and a best fit point (among all scenarios) can be found. We demonstrate this procedure in the next section by using present data.

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3.3 Specific scenarios

We now consider three explicit examples of new physics models. All three examples are well motivated and fulfill eq. (3.4).

- The dominant effect of the new physics is through Z-penguins [19]. Then, the Wilson coefficients are modified as follows:

$$\begin{aligned} C_3^{\rm NP}(m_W) &= \frac{1}{6}\varepsilon_z e^{i\theta_z}\,, \\ C_7^{\rm NP}(m_W) &= \frac{2}{3}\sin^2\theta_W \varepsilon_z e^{i\theta_z}\,, \\ C_9^{\rm NP}(m_W) &= -\frac{2}{3}(1-\sin^2\theta_W)\varepsilon_z e^{i\theta_z}\,. \end{aligned} \tag{3.10}$$

  Here

$$\varepsilon_z e^{i\theta_z} = \frac{g^2}{4\pi^2}\frac{Z_{sb}}{V_{tb}V_{ts}^*}\,, \tag{3.11}$$

  where Z_{sb} parametrizes the coupling of the Z^0 boson to the left handed strange and bottom quarks. The consistency of the measured $\mathrm{BR}(b \to s\ell^+\ell^-)$ with the SM prediction requires that $|Z_{sb}| \leq 0.08$ [19] which gives

$$\varepsilon_z \lesssim 0.02\,. \tag{3.12}$$

- Kaluza-Klein excitations of gluons in RS models [20] couple mainly to the third generation quarks and contribute to all gluonic penguin operators:[2]

$$\begin{aligned} C_4^{\rm NP}(m_W) &= C_6^{\rm NP}(m_W) = \varepsilon_k e^{i\theta_k}\,, \\ C_3^{\rm NP}(m_W) &= C_5^{\rm NP}(m_W) = -\frac{1}{3}\varepsilon_k e^{i\theta_k}\,. \end{aligned} \tag{3.13}$$

  Here

$$\varepsilon_k e^{i\theta_k} = -\pi\alpha_s(M_G)\left(\frac{v}{M_G}\right)^2 \frac{D_{sb}}{V_{tb}V_{ts}^*}\chi\,, \tag{3.14}$$

  where D_{sb} is the rotation matrix from interaction to mass basis for the left-handed down quarks, $\chi = \mathcal{O}(1)$ is a model dependent parameter (related to bulk fermion masses), and M_G is the mass of the lightest excitation. With $M_G \gtrsim 1\ TeV$ and $|D_{sb}| \sim |V_{tb}V_{ts}^*|$, we expect

$$\varepsilon_k \lesssim 0.02\,. \tag{3.15}$$

- Enhanced chromomagnetic operator [23] can be parametrized as follows:

$$C_{8g}^{\rm NP}(m_W) = \varepsilon_g e^{i\theta_g}\,. \tag{3.16}$$

[2] In ref. [21] it is argued that, quite generically in the RS framework, the set of operators in eq. (3.13) is subdominant and constrained to give only a small contribution to the decays in question. However, this issue is model dependent, and eq. (3.13) also represents a viable scenario.

f	$\pi^0 K$	ηK	$\eta' K$	ϕK	ωK	$\rho^0 K$
b^c_{fz}	-16	-21	-5.4	-16	24	57
	$[-12,-22]$	$[-16,-30]$	$[-4.0,-7.6]$	$[-10,-24]$	$[29,28]$	$[60,65]$
b^c_{fk}	-39	-43	-34	-33	-61	-64
	$[-28,-55]$	$[-33,-59]$	$[-22,-51]$	$[-24,-48]$	$[-63,-74]$	$[-66,-75]$
b^c_{fg}	0.75	0.64	0.86	1.4	-1.7	-1.5
	$[0.61,0.99]$	$[0.48,0.87]$	$[0.76,1.07]$	$[1.1,1.8]$	$[-2.8,-1.4]$	$[-2.0,-1.3]$

Table 6: The b^c_f parameters calculated in QCD factorization to leading log approximation and to zeroth order in Λ/m_b (chirally enhanced corrections are included) in the three scenarios: (i) Z-penguins [eq. (3.10)], denoted with subindex z; (ii) KK-gluons [eq. (3.13)], denoted with subindex k; (iii) Chromomagnetic operator [eq. (3.16)], denoted with subindex g. The first row in each case is obtained for $\mu = m_b$, whereas the second row gives the values for $\mu = m_b/2$ (first) and $2m_b$ (second).

(We still make the approximation that $C_{8g}^{\rm SM}(m_W) = 0$ and neglect further $C_{8g}^{\rm SM}(m_b)$ induced by O_1.) From $\text{BR}(B \to X_s g) < 0.09$ [22, 23], we have

$$\varepsilon_g \lesssim 1 \,. \tag{3.17}$$

Note that in this case the new physics contribution enters at next to leading order. Yet, the modifications of the CP asymmetries can be substantial. An experimental improvement by a factor of a few in the upper bound on $\text{BR}(B \to X_s g)$ will restrict the potential modifications in this scenario in a significant way.

We have calculated the b^c_f and b^u_f parameters of eq. (3.5) for each of the three scenarios. We use the factorized decay amplitudes from [3–5]. (We use asymptotic distribution amplitudes in the analysis of the scenario with an enhanced chromomagnetic operator.) The results are presented in table 6. The b^c_f parameters are quite stable under variation of the scale. As anticipated in section 3.2, in all three scenarios we have $b^u_f = (|V_{ub}V^*_{us}|/|V_{cb}V^*_{cs}|)\, b^c_f$.

Given the upper bounds on ε for the three scenarios — $\varepsilon_z \lesssim 0.02$, $\varepsilon_k \lesssim 0.02$ and $\varepsilon_g \lesssim 1$ — we learn that in none of the three examples can the new physics strongly dominate over the SM contributions, that is, $\varepsilon b^c_f \ggg a^u_f, 1$. Thus, if the data makes a convincing case for a universal pattern of deviations, it would mean that either this pattern is accidental or that the new physics is different from the three specific scenarios discussed here.

On the other hand, the interpretation of the results presented in table 6 is straightforward in the other limit discussed in the previous subsection. For the four modes $f = \pi^0 K_S$, ηK_S, $\eta' K_S$ and ϕK_S, we have, say, $|b^c_f \varepsilon| \lesssim 0.3$ for $\varepsilon_z \lesssim 0.01$, $\varepsilon_k \lesssim 0.005$, and $\varepsilon_g \lesssim 0.17$. In any of these cases, the deviation from the SM prediction is proportional to b^c_f [see eq. (3.9)]. Consequently, there is a distinct pattern of deviations for each of the three scenarios. For example, the deviation of $S_{\eta' K_S}$ from the SM prediction is a factor of four smaller than that of $S_{\eta K_S}$ in the first scenario ($\varepsilon_z \neq 0$), a factor ~ 1.3 smaller in the second ($\varepsilon_k \neq 0$) and a factor of ~ 1.3 larger in the third ($\varepsilon_g \neq 0$). The deviation of $S_{\phi K_S}$ is similar or somewhat smaller than that of $S_{\pi K_S}$ in the first two scenarios, but a factor of ~ 1.9 larger in the third. The second scenario has the characteristic feature that

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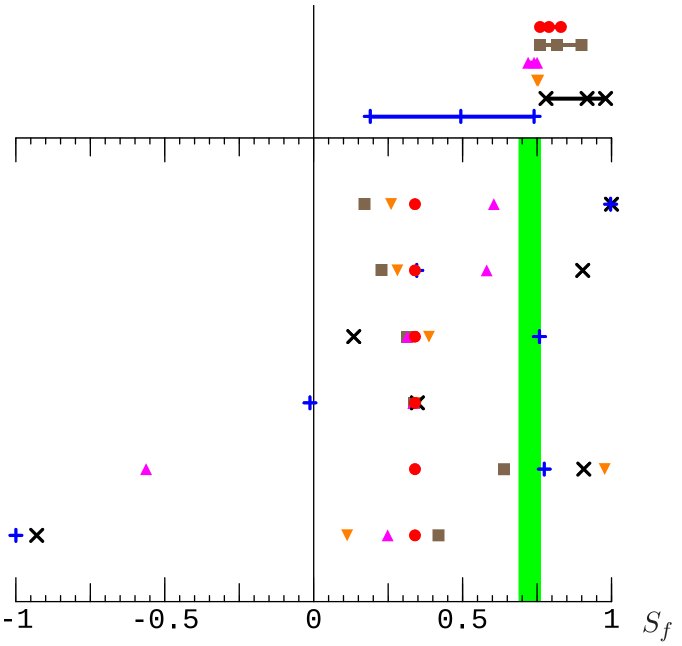

Figure 1: Pattern of S_f asymmetries for the three scenarios presented in section 3.3 with fixed $S_{\pi K_S} = 0.34$. The symbols represent $\bullet = \pi$, $\blacksquare = \eta$, $\blacktriangle = \eta'$, $\blacktriangledown = \phi$, $\times = \omega$ and $+ = \rho$. The upmost pattern gives the asymmetries in the SM for $\mu = m_b/2, m_b$ and $2m_b$. The other six patterns, from top to bottom, correspond to $(\varepsilon_z, \theta_z) = (0.02, 0.85)$, $(\varepsilon_z, \theta_z) = (0.02, 1.79)$, $(\varepsilon_k, \theta_k) = (0.02, 0.12)$, $(\varepsilon_k, \theta_k) = (0.02, 2.52)$, $(\varepsilon_g, \theta_g) = (1, 3.27)$, and $(\varepsilon_g, \theta_g) = (1, 5.64)$. The vertical bar represents the experimental value for $S_{\psi K_S}$.

the SM deviations of $S_{\phi K_S}$ and $S_{\eta' K_S}$ are very similar, whereas in scenario i (iii) $S_{\phi K_S}$ is a factor of 3 (1.6) larger than $S_{\eta' K_S}$. It is interesting to note that in all three new physics examples, the deviations for $b_f^c \varepsilon \ll 1$ in the modes that have been most accurately measured, $f = \pi^0 K_S$, $\eta' K_S$ and ϕK_S, and also ηK_S, are in the same direction, that is, the four S_f asymmetries are either all larger or all smaller [depending on sign$(\sin\theta)$] than the SM prediction.

To demonstrate how the pattern of deviations probes new physics, we perform the following exercise. For each of the three scenarios discussed in this section, we take ε_i at the phenomenological upper bound, that is $\varepsilon_z = 0.02$, $\varepsilon_k = 0.02$ and $\varepsilon_g = 1$. Since the new physics contribution is now as large as it can get, the patterns of deviations are expected to be more significant. We then pick the two values of θ_i for which $S_{\pi K_S} = 0.34$, the experimental central value. (There is no particular reason for selecting $S_{\pi K_S}$. Our aim is just to demonstrate the idea that the pattern of deviations is sensitive to new physics.) We present the results in figure 1. One can clearly see that the six patterns are different from each other, so that each scenario can in principle be easily tested.

4. Present data

At present, the B factories provide constraints on eight relevant CP asymmetries. The data are presented in table 1. We would like to test various scenarios of new physics by per-

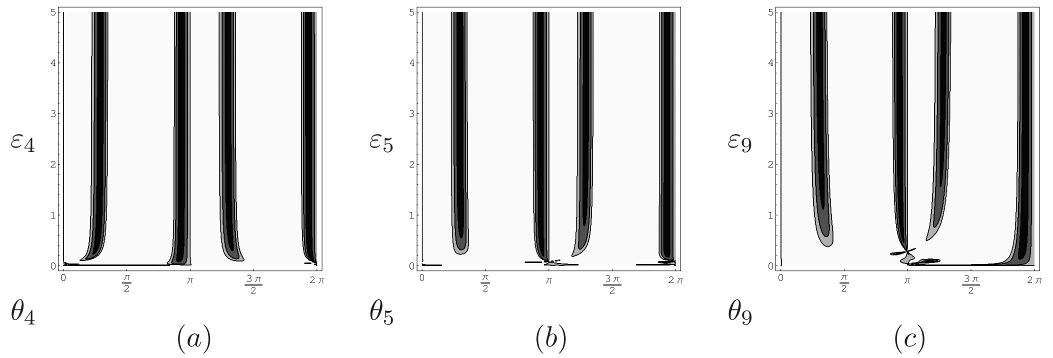

Figure 2: Allowed region in the $\varepsilon_i - \theta_i$ plane for new physics in a single operator: (a) O_4, (b) O_5 and (c) O_9. The black, dark grey, and light grey regions correspond to probability higher than 0.32, 0.046, and 0.0027, respectively.

forming a χ^2 fit to the data. Since we are mainly interested in demonstrating the potential power of the data to probe new physics in the future, we make several simplifications:

1. We work within the framework defined in section 3.2, that is, new physics contributions that depend on only a single complex parameter.

2. Since the measurement of $S_{\psi K_S}$ is very accurate, we fix the value of $\sin 2\beta_{\text{eff}}$ to its central value, and do not include it as a fit parameter. Furthermore, since experimental data [24] disfavor $\cos 2\beta_{\text{eff}} < 0$, we allow only positive $\cos 2\beta_{\text{eff}}$ values.

We thus calculate, for various scenarios,

$$\chi^2(\varepsilon, \theta) \equiv \sum_f \left[\frac{X_f^{\text{exp}} - X_f^{\text{QF}}(\varepsilon, \theta)}{\sigma_f^{\text{exp}}} \right]^2 , \tag{4.1}$$

where X_f^{exp} and σ_f^{exp} are, respectively, the central value and the one sigma range of the experimental measurements of S_f and C_f for $f = \pi^0 K_S$, $\eta' K_S$, ϕK_S and ωK_S, and $X_f^{\text{QF}}(\varepsilon, \theta)$ are the (central) theoretical values of S_f and C_f calculated within the QCD factorization approach for given values of ε and θ. Only experimental uncertainties are taken into account. Since there are eight observables and two free parameters, χ^2 has a six degrees of freedom distribution.

We consider eleven cases. Firstly, we investigate the possibility that the new physics contributes significantly via only a single O_i, where $i = 3, \ldots, 10$. (Specific models of new physics do not provide a strong motivation for such scenarios. Yet, this investigation is useful in the two limits discussed in section 3.2. If the new physics is dominant and depends on a single weak phase, we can see the universality of its effects [eq. (3.7)]. If the new physics contribution is small, the effects are well approximated by a sum over the shifts from the various operators [eq. (3.9)].) In each case, we perform our analysis for $\varepsilon \leq 5$ and all $\theta \in [0, 2\pi]$. Secondly, we investigate the three specific examples presented in section 3.3. We do this for ε within the phenomenologically allowed ranges and for all θ values.

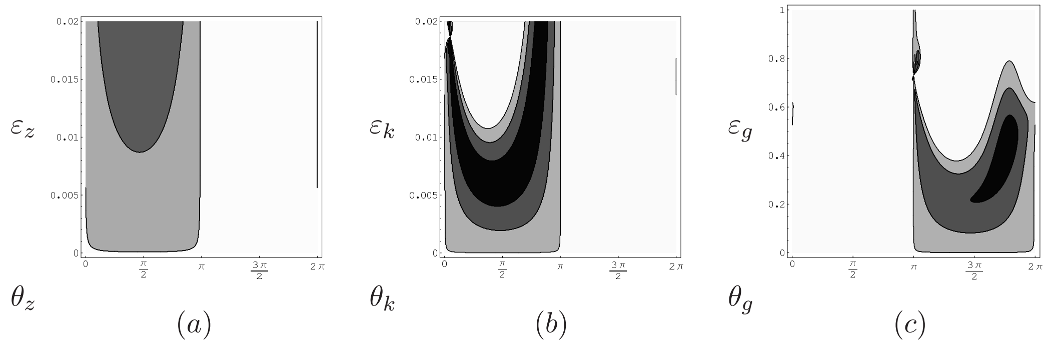

Figure 3: Allowed regions in the $\varepsilon - \theta$ plane for the three specific scenarios, (a) enhanced Z-penguins, (b) KK gluons in RS models, (c) enhanced chromomagnetic operator. The black, dark grey, and light grey regions correspond to probability higher than 0.32, 0.046, and 0.0027, respectively.

We use the χ^2 fit in two ways. For each case we find the favored regions in the $\varepsilon - \theta$ plane. The results are shown in figures 2 and 3. Since the fits to individual operators $O_{3\ldots10}$ give very similar results for large ϵ, we only show the ones for O_4, O_5 and O_9. Then we test the overall goodness of the fit for each case and learn whether the possibility that the data are accounted for by a new physics contribution to this (set of) operator(s) is disfavored. We draw the following conclusions:

1. For almost all the Wilson coefficients, a very good fit is obtained for a high ε value, such that eq. (3.6) holds and the shift is universal [eq. (3.7)]. The reason is that the central values of the three best measured S_f asymmetries ($f = \eta' K_S, \phi K_S, \pi^0 K_S$) show an approximately universal pattern: $\delta S_f \sim -0.4$.

2. For operators where eq. (3.6) gives the best fit, the preferred value of θ corresponds to $\sin 2(\beta + \theta) \sim 0.4$. Note that a four-fold ambiguity arises, giving $\theta_i \sim 55^o,\ 170^o,\ 235^o,\ 350^o$, which can be seen in figure 2. Had we allowed for $\cos 2\beta_{\rm eff} < 0$, there would be eight possible solutions for high ε.

3. For each of the three specific scenarios that we considered, there exists a region in the $\varepsilon - \theta$ plane where a good fit is obtained. Note that within the operator basis we use, in the first two scenarios with modified four-Fermi operators a phase smaller than π is preferred, whereas $\pi \leq \theta_g < 2\pi$ is favored.

4. Naively, the probability that there is no new physics contribution affecting the measured S_f's is small. (The SM point, $\varepsilon = 0$, lies in the white area.) One has to bear in mind, however, that we used a rather crude approximation, and that there are large uncertainties in our evaluation of the matrix elements. Furthermore, the discrepancy between $S^{\rm exp}_{\eta' K_S} = 0.43 \pm 0.11$ and $S^{\rm QF}_{\eta' K_S} = 0.74$ plays a significant role in this result. Note however that the Belle and Babar results for this observable are not quite consistent. Inflating the error according to the PDG prescription would yield $S^{\rm exp}_{\eta' K_S} = 0.43 \pm 0.17$ and a better SM fit. (Now the SM point lies in the light grey area.)

5. Conclusions

B-meson decays that proceed via the quark transition $b \to q\bar{q}s$, for $q = s, d, u$, open a window to new physics because the SM contribution is either loop- or CKM-suppressed. In particular, the CP violating asymmetries S_f and C_f in these modes may reveal the existence of new sources of CP violation beyond the KM phase.

With new physics, these decays may get significant contributions that depend on a phase that is different from the SM phase. Then, not only will these asymmetries be different from the SM prediction but also, in general, they will differ from each other. The pattern of deviations, $\delta S_f = -\eta_f S_f - S_{\psi K}$ and C_f, allows us to probe in detail the nature of the new physics that accounts for the effect.

Factorization schemes predict that there is an accidental cancellation between the leading Standard Model contributions to the ηK_S, ωK_S and ρK_S modes (see section 2). The suppression of the Standard Model amplitudes makes the corresponding CP asymmetries more sensitive to New Physics. At the same time, the sensitivity to hadronic uncertainties is also stronger.

We presented a way to analyze the data in a (new physics) model independent way. In a low energy effective theory, the effects of the new physics are expressed as modifications of Wilson coefficients in the $\Delta B = 1$ effective hamiltonian. The data can be used to find which Wilson coefficients are modified and by what (complex) factor.

The analysis is simplified when the modification of the Wilson coefficients depends on a single complex parameter, as is often the case in specific models. We gave analytic expressions for the modifications in some interesting limits. In particular, for $C_{3,\ldots,10,8g}$, when the new physics dominates, a universal shift of the asymmetries arises [see eq. (3.7)]. If the new physics contribution is small, then the deviations (for $f = \eta' K_S, \phi K_S, \pi^0 K_S, \eta K_S$) have a pattern that depends in a simple way on the b_f^c-parameters [see eq. (3.9)].

We applied this method to present data, using lowest order QCD factorization to calculate the hadronic matrix elements. We find that, since the best three measured S_f's have similar central values (~ 0.4) which are, however, different from $S_{\psi K}$, and all measured C_f's are consistent with zero, a good fit can be achieved for almost all operators if they are enhanced by the new physics to a level where they dominate over the SM contribution. In three specific scenarios that we considered, such new physics dominance cannot be realized in Nature because of phenomenological constraints. Still, a good fit can be obtained for each of these scenarios.

Our analysis can be improved in several ways. In particular, the calculation should go beyond the leading log approximation and incorporate power corrections. The combination of improved calculations, more accurate measurements and additional modes explored experimentally, is likely to make the pattern of deviations from the SM (or their absence) a powerful probe of flavor and CP violating new physics.

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