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Exploring New Physics in the $B \to \phi K$ System

Robert Fleischer*
Deutsches Elektronen-Synchrotron DESY, Notkestraße 85, D–22607 Hamburg, Germany

Thomas Mannel†
CERN Theory Division, CH–1211 Geneva 23, Switzerland
and
Institut für Theoretische Teilchenphysik, Universität Karlsruhe, D–76128 Karlsruhe, Germany

Abstract

Employing the $SU(2)$ isospin symmetry of strong interactions and estimates borrowed from effective field theory, we explore the impact of new physics on the decays $B^\pm \to \phi K^\pm$ and $B_d \to \phi K_S$ in a model-independent manner. To this end, we introduce – in addition to the usual mixing-induced CP asymmetry in $B_d \to \phi K_S$ – a set of three observables, which may not only provide "smoking-gun" signals for new-physics contributions to different isospin channels, but also valuable insights into hadron dynamics. Imposing dynamical hierarchies of amplitudes, we discuss various patterns of these observables, including also scenarios with small and large rescattering processes. Whereas the $B \to \phi K$ system provides, in general, a powerful tool to search for indications of new physics, there is also an unfortunate case, where such effects cannot be distinguished from those of the Standard Model.

*E-mail: Robert.Fleischer@desy.de
†E-mail: Thomas.Mannel@cern.ch

1 Introduction

The experimental data collected at the B factories will allow stringent tests of the Kobayashi–Maskawa picture of CP violation [1]. Among the various B decays that can be used to achieve this goal [2], $B \to \phi K$ decays play an important role. In these modes, which are governed by QCD penguin processes [3], also electroweak (EW) penguins are sizeable [4], and physics beyond the Standard Model may have an important impact [5]. In the summer of 2000, the observation of the $B^\pm \to \phi K^\pm$ channel was announced by the Belle and CLEO collaborations. The present results for the CP-averaged branching ratio are given as follows:

$$\mathrm{BR}(B^\pm \to \phi K^\pm) = \begin{cases} \left(1.39^{+0.37+0.14}_{-0.33-0.24}\right) \times 10^{-5} \text{ (Belle [6])} \\ \left(5.5^{+2.1}_{-1.8} \pm 0.6\right) \times 10^{-6} \text{ (CLEO [7]).} \end{cases} \tag{1}$$

The Belle and CLEO results are only marginally compatible with each other. Evidence for the neutral mode $B_d^0 \to \phi K^0$ at the 2.9σ level, corresponding to a branching ratio of $\left(5.4^{+3.7}_{-2.7} \pm 0.7\right) \times 10^{-6}$, was also reported by CLEO, whereas a significant signal for this decay has not yet been observed by the Belle collaboration.

In our discussion of the $B \to \phi K$ system, we follow closely our recent $B \to J/\psi K$ analysis [8], and make use of the $SU(2)$ isospin symmetry of strong interactions to derive a model-independent parametrization of the $B^+ \to \phi K^+$, $B_d^0 \to \phi K^0$ decay amplitudes. After recapitulating the structure of the Standard-Model amplitudes, we include new-physics effects in a general manner, and estimate their generic size with the help of arguments borrowed from the picture of effective field theory. In order to deal with hadronic matrix elements, we impose certain dynamical hierarchies of decay amplitudes, where we distinguish between small and large rescattering effects. Although we do not consider the latter case, which is also not favoured by the "QCD factorization" approach [9] and the present experimental upper bounds on $B \to KK$ branching ratios [10], as a very likely scenario,[1] it deserves careful attention to separate possible new-physics effects from those of the Standard Model. Moreover, following the strategies proposed below, we may not only obtain insights into new physics, but also into hadron dynamics. To this end, we introduce – in addition to the usual mixing-induced CP asymmetry in $B_d \to \phi K_{\mathrm{S}}$ – a set of three observables, providing "smoking-gun" signals for new-physics contributions to different isospin channels. Two of these new-physics observables may be significantly enhanced by large rescattering processes. In general, the $B \to \phi K$ system offers powerful tools to search for new physics. However, there is also an unfortunate case, where such effects cannot be disentangled from those of the Standard Model.

The outline of this paper is as follows: in Section 2, we employ a low-energy effective Hamiltonian and the isospin symmetry of strong interactions to parametrize the $B^\pm \to \phi K^\pm$, $B_d \to \phi K_{\mathrm{S}}$ decay amplitudes arising within the Standard Model. The impact of new physics on these amplitudes is explored in a model-independent way in Section 3, where we make use of dimensional estimates following from effective field theory, and introduce plausible dynamical hierarchies of amplitudes. The set of observables to search for "smoking-gun" signals of new-physics contributions to different isospin channels of the $B \to \phi K$ decay amplitudes is introduced in Section 4, and is discussed in further detail in Section 5. In Section 6, our conclusions are summarized.

[1] Arguments against this possibility, i.e. large rescattering effects, were also given in [11].

2 Phenomenology of $B \to \phi K$ Decays

The $B \to \phi K$ system is described by the following low-energy effective Hamiltonian:

$$\mathcal{H}_{\rm eff} = \frac{G_{\rm F}}{\sqrt{2}} \left[V_{cs} V_{cb}^* \left(\mathcal{Q}_{\rm CC}^c - \mathcal{Q}_{\rm QCD}^{\rm pen} - \mathcal{Q}_{\rm EW}^{\rm pen} \right) + V_{us} V_{ub}^* \left(\mathcal{Q}_{\rm CC}^u - \mathcal{Q}_{\rm QCD}^{\rm pen} - \mathcal{Q}_{\rm EW}^{\rm pen} \right) \right], \tag{2}$$

where the $\mathcal{Q}$ are linear combinations of perturbative Wilson coefficient functions and four-quark operators, consisting of current–current (CC), QCD penguin and EW penguin operators. As discussed in [8], this Hamiltonian is a combination of isospin $I=0$ and $I=1$ pieces:

$$\mathcal{H}_{\rm eff} = \mathcal{H}_{\rm eff}^{I=0} + \mathcal{H}_{\rm eff}^{I=1}, \tag{3}$$

where $\mathcal{H}_{\rm eff}^{I=0}$ receives contributions from all of the operators appearing in (2), whereas $\mathcal{H}_{\rm eff}^{I=1}$ is due to only $\mathcal{Q}_{\rm CC}^u$ and $\mathcal{Q}_{\rm EW}^{\rm pen}$. If we employ the $SU(2)$ isospin flavour symmetry of strong interactions, we obtain

$$\langle \phi K^+ | \mathcal{H}_{\rm eff}^{I=0} | B^+ \rangle = +\langle \phi K^0 | \mathcal{H}_{\rm eff}^{I=0} | B_d^0 \rangle \tag{4}$$

$$\langle \phi K^+ | \mathcal{H}_{\rm eff}^{I=1} | B^+ \rangle = -\langle \phi K^0 | \mathcal{H}_{\rm eff}^{I=1} | B_d^0 \rangle, \tag{5}$$

yielding

$$A(B^+ \to \phi K^+) = \frac{G_{\rm F}}{\sqrt{2}} \left[V_{cs} V_{cb}^* \left\{ \mathcal{A}_c^{(0)} - \mathcal{A}_c^{(1)} \right\} + V_{us} V_{ub}^* \left\{ \mathcal{A}_u^{(0)} - \mathcal{A}_u^{(1)} \right\} \right] \tag{6}$$

$$A(B_d^0 \to \phi K^0) = \frac{G_{\rm F}}{\sqrt{2}} \left[V_{cs} V_{cb}^* \left\{ \mathcal{A}_c^{(0)} + \mathcal{A}_c^{(1)} \right\} + V_{us} V_{ub}^* \left\{ \mathcal{A}_u^{(0)} + \mathcal{A}_u^{(1)} \right\} \right], \tag{7}$$

where the CP-conserving strong amplitudes[2]

$$\mathcal{A}_c^{(0)} = \mathcal{A}_{\rm CC}^c - \mathcal{A}_{\rm QCD}^{\rm pen} - \mathcal{A}_{\rm EW}^{(0)}, \quad \mathcal{A}_c^{(1)} = -\mathcal{A}_{\rm EW}^{(1)} \tag{8}$$

$$\mathcal{A}_u^{(0)} = \mathcal{A}_{\rm CC}^{u(0)} - \mathcal{A}_{\rm QCD}^{\rm pen} - \mathcal{A}_{\rm EW}^{(0)}, \quad \mathcal{A}_u^{(1)} = \mathcal{A}_{\rm CC}^{u(1)} - \mathcal{A}_{\rm EW}^{(1)} \tag{9}$$

can be expressed in terms of hadronic matrix elements $\langle \phi K | \mathcal{Q} | B \rangle$. Taking into account that

$$V_{cs} V_{cb}^* = \left(1 - \frac{\lambda^2}{2} \right) \lambda^2 A, \quad V_{us} V_{ub}^* = \lambda^4 A \, R_b \, e^{i\gamma}, \tag{10}$$

where γ is the usual angle of the unitarity triangle of the CKM matrix [2], and

$$\lambda \equiv |V_{us}| = 0.22, \quad A \equiv |V_{cb}|/\lambda^2 = 0.81 \pm 0.06, \quad R_b \equiv |V_{ub}/(\lambda V_{cb})| = 0.41 \pm 0.07, \tag{11}$$

we finally arrive at

$$A(B^+ \to \phi K^+) = \frac{G_{\rm F}}{\sqrt{2}} \left(1 - \frac{\lambda^2}{2} \right) \lambda^2 A \left\{ \mathcal{A}_c^{(0)} - \mathcal{A}_c^{(1)} \right\} \left[1 + \frac{\lambda^2 R_b}{1 - \lambda^2/2} \left\{ \frac{\mathcal{A}_u^{(0)} - \mathcal{A}_u^{(1)}}{\mathcal{A}_c^{(0)} - \mathcal{A}_c^{(1)}} \right\} e^{i\gamma} \right] \tag{12}$$

$$A(B_d^0 \to \phi K^0) = \frac{G_{\rm F}}{\sqrt{2}} \left(1 - \frac{\lambda^2}{2} \right) \lambda^2 A \left\{ \mathcal{A}_c^{(0)} + \mathcal{A}_c^{(1)} \right\} \left[1 + \frac{\lambda^2 R_b}{1 - \lambda^2/2} \left\{ \frac{\mathcal{A}_u^{(0)} + \mathcal{A}_u^{(1)}}{\mathcal{A}_c^{(0)} + \mathcal{A}_c^{(1)}} \right\} e^{i\gamma} \right]. \tag{13}$$

[2] The labels of $\mathcal{A}^{(0)}$ and $\mathcal{A}^{(1)}$ refer to the isospin channels $I=0$ and $I=1$, respectively.

At first sight, expressions (12) and (13) are completely analogous to the ones for the $B^+ \to J/\psi K^+$ and $B^0_d \to J/\psi K^0$ amplitudes given in [8]. However, the dynamics, which is encoded in the strong amplitudes $\mathcal{A}$, is very different. In particular, the current–current operators $\mathcal{Q}^c_{\rm CC}$ cannot contribute to $B \to \phi K$ decays, i.e. to $\mathcal{A}^c_{\rm CC}$, through tree-diagram-like topologies; they may only do so through penguin topologies with internal charm-quark exchanges, which include also

$$B^+ \to \{D_s^+ \overline{D^0}, ...\} \to \phi K^+, \quad B^0_d \to \{D_s^+ D^-, ...\} \to \phi K^0 \tag{14}$$

rescattering processes [12], and may actually play an important role [13]. On the other hand, the $\mathcal{A}^{u(0,1)}_{\rm CC}$ amplitudes receive contributions from penguin processes with internal up- and down-quark exchanges, as well as from annihilation topologies.[3] Such penguins may also be important, in particular in the presence of large rescattering processes [12, 14]; a similar comment applies to annihilation topologies. In the $B \to \phi K$ system, the relevant rescattering processes are

$$B^+ \to \{K^+ \pi^0, ...\} \to \phi K^+, \quad B^0_d \to \{K^+ \pi^-, ...\} \to \phi K^0, \tag{15}$$

containing – in addition to long-distance penguins – also annihilation processes (see Figs. 1 and 2). In contrast to (14), large rescattering effects of the kind described by (15) may affect the search for new physics with $B \to \phi K$ decays, since these processes are associated with the weak phase factor $e^{i\gamma}$. Moreover, they involve "light" intermediate states, and are hence expected to be enhanced more easily, dynamically, through long-distance effects than (14), which involve "heavy" intermediate states.

As is well known, the ϕ-meson is an almost pure $\overline{s}s$ state; the mixing angle with its isoscalar partner $\omega \sim (\overline{u}u + \overline{d}d)/\sqrt{2}$ is small, i.e. at the few per cent level. Whereas ω–ϕ mixing does not at all affect the isospin relations (4) and (5), which rely on the fact that the ϕ is an isospin singlet, it has an impact on the size of the amplitudes $\mathcal{A}^{u(0,1)}_{\rm CC}$, since an ω component of the ϕ state permits current–current operator contributions through tree-diagram-like topologies. However, the arguments given below are not modified by the small ω–ϕ mixing.

Let us now have a closer look at the structure of the $B \to \phi K$ decay amplitudes, focusing first on the case corresponding to small rescattering effects. Looking at (8) and (9), we expect

$$\left|\mathcal{A}^{(0,1)}_u / \mathcal{A}^{(0)}_c\right| = \mathcal{O}(1). \tag{16}$$

In the case of the amplitude $\mathcal{A}^{(1)}_c$, the situation is different. Here we have to deal with an amplitude that is essentially due to EW penguins. Moreover, the $B \to \phi K$ matrix elements of $I = 1$ operators, having the general flavour structure

$$\mathcal{Q}_{I=1} \sim (\overline{u}u - \overline{d}d)(\overline{b}s), \tag{17}$$

are expected to suffer from a dynamical suppression. In order to keep track of these features, we introduce, as in [8], a "generic" expansion parameter $\overline{\lambda} = \mathcal{O}(0.2)$ [15], which is of the same order as the Wolfenstein parameter $\lambda = 0.22$, and suggests

$$\left|\mathcal{A}^{(1)}_c / \mathcal{A}^{(0)}_c\right| = \underbrace{\mathcal{O}(\overline{\lambda})}_{\text{EW penguins}} \times \underbrace{\mathcal{O}(\overline{\lambda})}_{\text{Dynamics}} = \mathcal{O}(\overline{\lambda}^2). \tag{18}$$

[3]Note that the isospin projection operators $\mathcal{Q} \sim (\overline{u}u \pm \overline{d}d)(\overline{b}s)$ involve also $\overline{d}d$ quark currents.

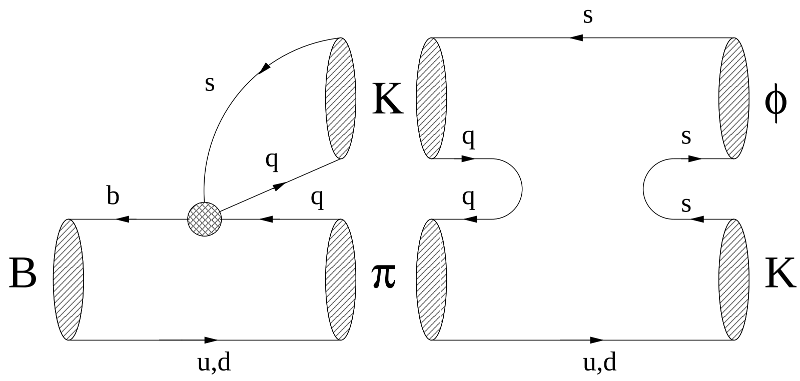

Figure 1: Rescattering processes contributing to $B \to \phi K$ through penguin-like topologies with internal q-quark exchanges ($q \in \{u, d\}$). The shaded circle represents insertions of the corresponding current–current operators.

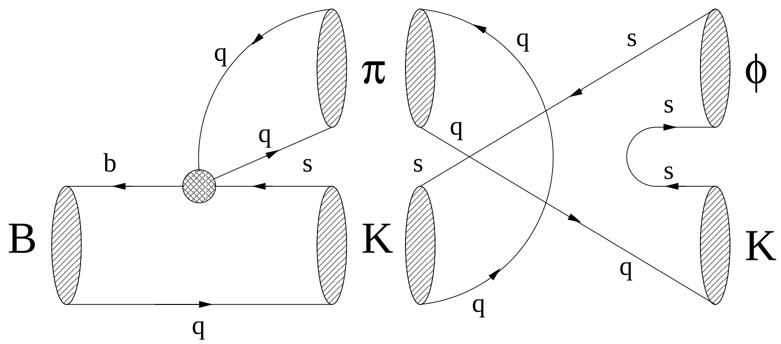

Figure 2: Rescattering processes contributing to $B \to \phi K$ through annihilation topologies. The shaded circle represents insertions of current–current operators ($q \in \{u, d\}$).

Consequently, we obtain

$$A(B^+ \to \phi K^+) = \mathcal{A}_{\rm SM}^{(0)}\left[1 + \mathcal{O}(\overline{\lambda}^2)\right] = A(B_d^0 \to \phi K^0), \tag{19}$$

with

$$\mathcal{A}_{\rm SM}^{(0)} \equiv \frac{G_{\rm F}}{\sqrt{2}}\lambda^2 A\, \mathcal{A}_c^{(0)}. \tag{20}$$

The terms entering (19) at the $\overline{\lambda}^2$ level contain also pieces that are proportional to the weak-phase factor $e^{i\gamma}$, thereby leading to direct CP violation in the $B \to \phi K$ system.

Let us now consider large rescattering effects of the kind given in (15). In the worst case, (16) would be dynamically enhanced as

$$\left|\mathcal{A}_u^{(0,1)}/\mathcal{A}_c^{(0)}\right| = \mathcal{O}(1/\overline{\lambda}), \tag{21}$$

and the dynamical suppression in (18) would no longer be effective, i.e.

$$\left|\mathcal{A}_c^{(1)}/\mathcal{A}_c^{(0)}\right| = \mathcal{O}(\overline{\lambda}). \tag{22}$$

In such a scenario, (19) would receive corrections of $\mathcal{O}(\overline{\lambda})$, involving also $e^{i\gamma}$. This feature may complicate the search for new physics with the help of CP-violating effects in $B \to \phi K$ decays. On the other hand, the rescattering processes described by (14) may only affect the amplitude $\mathcal{A}_{\rm SM}^{(0)}$ sizeably through its $\mathcal{A}_{\rm CC}^{c}$ piece, and are not related to a CP-violating weak-phase factor within the Standard Model.

Before turning to new physics, we would like to point out an interesting difference between the $B \to \phi K$ system and the $B \to J/\psi K$ decays discussed in [8]. In the $B \to J/\psi K$ case, the Standard-Model amplitudes corresponding to (19) receive corrections at the $\overline{\lambda}^3$ level, which may be enhanced to $\mathcal{O}(\overline{\lambda}^2)$ in the presence of large rescattering effects. Consequently, within the Standard Model, there may be direct CP-violating effects in $B \to J/\psi K$ transitions of at most $\mathcal{O}(\overline{\lambda}^2)$, whereas such asymmetries may already arise at the $\overline{\lambda}$ level in the $B \to \phi K$ system. On the other hand, as $B \to \phi K$ modes are governed by penguin processes, their decay amplitudes are more sensitive to new physics. In the next section, we have a closer look at the generic size of such effects, using dimensional arguments inspired by effective field theory.

3 Effects of Physics Beyond the Standard Model

The most general way of introducing new physics is to employ the picture of effective field theory, where we write down all possible dim-6 operators and construct the generalization of the Standard-Model effective Hamiltonian (2) at the scale of the b quark. The relevant operators are again dim-6 operators, where the Wilson coefficients of those already present in the Standard Model now contain also a possible piece of new physics. The problem with this generic point of view is that the list of possible dim-6 operators contains close to one hundred entries [16], not yet taking into account the flavour structure, making this general approach almost useless. However, as we are dealing with non-leptonic decays in which, because of our ignorance of the hadronic matrix elements, there is no senistivity neither to the helicity structure of the operators nor to their colour structure, only the flavour structure of the operators is relevant for us.

A discussion of the $\Delta B = \pm 2$ operators mediating B_d^0–$\overline{B_d^0}$ mixing within this framework was given in [8]. For a characteristic new-physics scale Λ in the TeV regime, new-physics contributions could in principle be as large as those of the Standard Model. This well-known feature was also found in several model-dependent studies of physics beyond the Standard Model [17]. As far as CP violation is concerned, new-physics contributions involving also new CP-violating phases are of particular interest. In this case, not only the "strength" of B_d^0–$\overline{B_d^0}$ mixing is affected, which would manifest itself as an inconsistency in the usual "standard analysis" of the unitarity triangle [18], but also "mixing-induced" CP-violating asymmetries [2], arising, for instance, in $B_d \to J/\psi K_{\rm S}$ or $B_d \to \phi K_{\rm S}$ decays.

Let us now turn to the $B \to \phi K$ amplitudes. As in (3), also in the presence of new physics, the corresponding low-energy effective Hamiltonian can be decomposed into $I = 0$ and $I = 1$ pieces, where the new physics may affect the Wilson coefficients and may introduce new dim-6 operators, thereby modifying (19) as follows:

$$A(B^+ \to \phi K^+) = \mathcal{A}_{\rm SM}^{(0)} \left[1 + \sum_k v_0^{(k)} e^{i\Delta_0^{(k)}} e^{i\Phi_0^{(k)}} - \sum_j v_1^{(j)} e^{i\Delta_1^{(j)}} e^{i\Phi_1^{(j)}} \right] \tag{23}$$

$$A(B_d^0 \to \phi K^0) = \mathcal{A}_{\rm SM}^{(0)} \left[1 + \sum_k v_0^{(k)} e^{i\Delta_0^{(k)}} e^{i\Phi_0^{(k)}} + \sum_j v_1^{(j)} e^{i\Delta_1^{(j)}} e^{i\Phi_1^{(j)}} \right]. \tag{24}$$

Here $v_0^{(k)}$ and $v_1^{(j)}$ correspond to the $I=0$ and $I=1$ pieces, respectively, $\Delta_0^{(k)}$ and $\Delta_1^{(j)}$ are CP-conserving strong phases, and $\Phi_0^{(k)}$ and $\Phi_1^{(j)}$ the corresponding CP-violating weak phases. The amplitudes for the CP-conjugate processes can be obtained straightforwardly from (23) and (24) by reversing the signs of the weak phases. The labels k and j distinguish between different new-physics contributions to the $I=0$ and $I=1$ sectors.

As we have already noted, within the Standard Model, $B \to \phi K$ decays are governed by QCD penguins. Neglecting, for simplicity, EW penguins and the proper renormalization-group evolution, we may write

$$\mathcal{A}_{\rm SM}^{(0)} \sim \frac{G_{\rm F}}{\sqrt{2}} \lambda^2 A \left[\frac{\alpha_s}{4\pi} \mathcal{C} \right] \langle P_{\rm QCD} \rangle, \tag{25}$$

where $\mathcal{C} = \mathcal{O}(1)$ is a perturbative short-distance coefficient, which is multiplied by the characteristic loop factor $\alpha_s/(4\pi)$, and $P_{\rm QCD}$ denotes an appropriate linear combination of QCD penguin operators [19]. Since (25) is a doubly Cabibbo-suppressed loop amplitude, new physics could well be of the same order of magnitude. If we assume that the physics beyond the Standard Model is associated with a scale Λ and impose that it yields contributions to the $B \to \phi K$ amplitudes of the same size as the Standard Model, we obtain

$$\frac{G_{\rm F}}{\sqrt{2}} \frac{M_W^2}{\Lambda^2} \sim \frac{G_{\rm F}}{\sqrt{2}} \lambda^2 A \left[\frac{\alpha_s}{4\pi} \mathcal{C} \right], \tag{26}$$

corresponding to $\Lambda \sim 3\,{\rm TeV}$.[4] Consequently, for a generic new-physics scale in the TeV regime, which was also considered in our $B \to J/\psi K$ analysis [8], we may have

$$v_0^{(k)} = \mathcal{O}(1). \tag{27}$$

In the case where the new-physics effects are less pronounced, it may be difficult to disentangle them from the Standard-Model contributions. We shall come back to this issue in Section 5, where we shall discuss various scenarios.

Concerning possible new-physics contributions to the $I=1$ sector, we assume a "generic strength" of the corresponding operators similar to (26). However, since the $I=1$ operators have the general flavour structure given in (17), they are expected to suffer from a dynamical suppression in $B \to \phi K$ decays. As in (18), we assume that this brings a factor of $\overline{\lambda}$ into the game, yielding

$$v_1^{(j)} = \mathcal{O}(\overline{\lambda}). \tag{28}$$

If we impose such a hierarchy of amplitudes, the new-physics contributions to the $I=1$ sector would be enhanced by a factor of $\mathcal{O}(\overline{\lambda})$ with respect to the $I=1$ Standard-Model pieces. This may actually be the case if new physics shows up, for example, in EW penguin processes.

Consequently, we finally arrive at

$$A(B \to \phi K) = \mathcal{A}_{\rm SM}^{(0)} \Big[1 + \underbrace{\mathcal{O}(1)}_{{\rm NP}_{I=0}} + \underbrace{\mathcal{O}(\overline{\lambda})}_{{\rm NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\rm SM} \Big]. \tag{29}$$

[4] In our numerical estimate, we have assumed $A \times \mathcal{C} \sim 1$ and $\alpha_s = \alpha_s(m_b) \sim 0.2$.

In deriving this expression, we have assumed that the $B \to \phi K$ decays are not affected by rescattering effects. On the other hand, in the presence of large rescattering processes of the kind described by (15), the dynamical suppression assumed in (28) would no longer be effective, thereby yielding $v_1^{(j)} = \mathcal{O}(1)$. Analogously, the $B \to \phi K$ matrix elements of $I = 0$ operators with flavour structure

$$\mathcal{Q}_{I=0}^{\overline{u}u,\overline{d}d} \sim (\overline{u}u + \overline{d}d)(\overline{b}s) \tag{30}$$

would no longer be suppressed with respect to those of the dynamically favoured $I = 0$ operators

$$\mathcal{Q}_{I=0}^{\overline{s}s} \sim (\overline{s}s)(\overline{b}s), \tag{31}$$

and would also contribute to $v_0^{(k)}$ at $\mathcal{O}(1)$. A similar comment applies to the matrix elements of the $I = 0$ operators with the following flavour content:

$$\mathcal{Q}_{I=0}^{\overline{c}c} \sim (\overline{c}c)(\overline{b}s), \tag{32}$$

whose dynamical suppression in $B \to \phi K$ decays may be reduced through rescattering effects of the kind given in (14), which may also affect the $\mathcal{A}_{\rm SM}^{(0)}$ amplitude, as we have noted above. Consequently, in the presence of large rescattering effects, the decay amplitude (29) is modified as follows:

$$A(B \to \phi K)\big|_{\rm res.} = \mathcal{A}_{\rm SM}^{(0)}\Big|_{\rm res.} \times \left[1 + \underbrace{\mathcal{O}(1)}_{{\rm NP}_{I=0}} + \underbrace{\mathcal{O}(1)}_{{\rm NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda})}_{\rm SM}\right]. \tag{33}$$

Let us emphasize once again that the rescattering contributions to the prefactor on the right-hand side of this equation are due to (14), whereas the hierarchy in square brackets is governed by large rescattering processes of the type described by (15).

In the following section, we introduce a set of observables, allowing us to separate the $I = 0$ contributions from the $I = 1$ sector. These observables play a key role to search for new physics and to obtain, moreover, valuable insights into the $B \to \phi K$ hadron dynamics.

4 Observables of $B \to \phi K$ Decays

The decays $B^+ \to \phi K^+$, $B_d^0 \to \phi K^0$ and their charge conjugates are described by four decay amplitudes A_i. If the corresponding rates are measured, the $|A_i|^2$ can be extracted. Since we are not interested in the overall normalization $\mathcal{A}_{\rm SM}^{(0)}$ of these amplitudes, we may construct the following three independent observables with the help of the $|A_i|^2$:

$$\mathcal{A}_{\rm CP}^{(+)} \equiv \frac{|A(B^+ \to \phi K^+)|^2 - |A(B^- \to \phi K^-)|^2}{|A(B^+ \to \phi K^+)|^2 + |A(B^- \to \phi K^-)|^2} \tag{34}$$

$$\mathcal{A}_{\rm CP}^{\rm dir} \equiv \frac{|A(B_d^0 \to \phi K^0)|^2 - |A(\overline{B_d^0} \to \phi \overline{K^0})|^2}{|A(B_d^0 \to \phi K^0)|^2 + |A(\overline{B_d^0} \to \phi \overline{K^0})|^2} \tag{35}$$

$$\mathcal{B} \equiv \frac{\langle |A(B_d \to \phi K)|^2\rangle - \langle |A(B^\pm \to \phi K^\pm)|^2\rangle}{\langle |A(B_d \to \phi K)|^2\rangle + \langle |A(B^\pm \to \phi K^\pm)|^2\rangle}, \tag{36}$$

where the "CP-averaged" amplitudes are defined in the usual way:

$$\langle|A(B_d \to \phi K)|^2\rangle \equiv \frac{1}{2}\left[|A(B_d^0 \to \phi K^0)|^2 + |A(\overline{B_d^0} \to \phi \overline{K^0})|^2\right] \tag{37}$$

$$\langle|A(B^\pm \to \phi K^\pm)|^2\rangle \equiv \frac{1}{2}\left[|A(B^+ \to \phi K^+)|^2 + |A(B^- \to \phi K^-)|^2\right]. \tag{38}$$

It should be noted that $\mathcal{B}$ is a CP-conserving quantity. In the case of the neutral decay $B_d \to \phi K_{\rm S}$, interference between B_d^0–$\overline{B_d^0}$ mixing and decay processes yields an additional observable [2]:

$$\frac{\Gamma(B_d^0(t) \to \phi K_{\rm S}) - \Gamma(\overline{B_d^0}(t) \to \phi K_{\rm S})}{\Gamma(B_d^0(t) \to \phi K_{\rm S}) + \Gamma(\overline{B_d^0}(t) \to \phi K_{\rm S})} = \mathcal{A}_{\rm CP}^{\rm dir}\cos(\Delta M_d t) + \mathcal{A}_{\rm CP}^{\rm mix}\sin(\Delta M_d t), \tag{39}$$

where the "direct" CP-violating contribution $\mathcal{A}_{\rm CP}^{\rm dir}$ was already introduced in (35), and the "mixing-induced" CP asymmetry is given by

$$\mathcal{A}_{\rm CP}^{\rm mix} = \frac{2\,{\rm Im}\,\xi}{1+|\xi|^2}, \tag{40}$$

with

$$\xi = e^{-i\phi}\left[\frac{1+\sum_k v_0^{(k)} e^{i\Delta_0^{(k)}} e^{-i\Phi_0^{(k)}} + \sum_j v_1^{(j)} e^{i\Delta_1^{(j)}} e^{-i\Phi_1^{(j)}}}{1+\sum_k v_0^{(k)} e^{i\Delta_0^{(k)}} e^{+i\Phi_0^{(k)}} + \sum_j v_1^{(j)} e^{i\Delta_1^{(j)}} e^{+i\Phi_1^{(j)}}}\right]. \tag{41}$$

Here (24) was used to parametrize the decay amplitudes. The CP-violating phase ϕ is given by $\phi = \phi_{\rm M} + \phi_K$, where $\phi_{\rm M}$ and ϕ_K are the weak B_d^0–$\overline{B_d^0}$ and K^0–$\overline{K^0}$ mixing phases, respectively.

As in our analysis of the $B \to J/\psi K$ system [8], it is useful to introduce the following combinations of the observables (34) and (35):

$$\mathcal{S} \equiv \frac{1}{2}\left[\mathcal{A}_{\rm CP}^{\rm dir} + \mathcal{A}_{\rm CP}^{(+)}\right], \quad \mathcal{D} \equiv \frac{1}{2}\left[\mathcal{A}_{\rm CP}^{\rm dir} - \mathcal{A}_{\rm CP}^{(+)}\right]. \tag{42}$$

The interesting feature of these combinations is that $\mathcal{S}$ is governed by the $I=0$ pieces of the $B \to \phi K$ amplitudes, whereas $\mathcal{D}$ and $\mathcal{B}$ are proportional to the $I=1$ amplitudes.

As the general expressions are very complicated and not very instructive, let us focus on the simplified case where the new-physics contributions to the $I=0$ and $I=1$ sectors involve either the same weak or strong phases:

$$A(B^+ \to \phi K^+) = \mathcal{A}_{\rm SM}^{(0)}\left[1 + v_0 e^{i\Delta_0} e^{i\Phi_0} - v_1 e^{i\Delta_1} e^{i\Phi_1}\right] \tag{43}$$

$$A(B_d^0 \to \phi K^0) = \mathcal{A}_{\rm SM}^{(0)}\left[1 + v_0 e^{i\Delta_0} e^{i\Phi_0} + v_1 e^{i\Delta_1} e^{i\Phi_1}\right]. \tag{44}$$

Then we obtain

$$\mathcal{S} = -2\left[\frac{ac - bd\,v_1^2}{c^2 - d^2\,v_1^2}\right] = -2\,v_0\left[\frac{\sin\Delta_0 \sin\Phi_0}{1 + 2\,v_0\cos\Delta_0\cos\Phi_0 + v_0^2}\right] + \mathcal{O}(v_1^2) \tag{45}$$

$$\mathcal{D} = -2\,v_1\left[\frac{bc - ad}{c^2 - d^2\,v_1^2}\right] \tag{46}$$

$$\mathcal{B} = v_1\,\frac{d}{c} = 2\,v_1\left[\frac{\cos\Delta_1\cos\Phi_1 + v_0\cos(\Delta_0 - \Delta_1)\cos(\Phi_0 - \Phi_1)}{1 + 2\,v_0\cos\Delta_0\cos\Phi_0 + v_0^2 + v_1^2}\right], \tag{47}$$

with

$$a = v_0 \sin \Delta_0 \sin \Phi_0 \tag{48}$$

$$b = \sin \Delta_1 \sin \Phi_1 + v_0 \sin(\Delta_0 - \Delta_1) \sin(\Phi_0 - \Phi_1) \tag{49}$$

$$c = 1 + 2\, v_0 \cos \Delta_0 \cos \Phi_0 + v_0^2 + v_1^2 \tag{50}$$

$$d = 2 \left[\cos \Delta_1 \cos \Phi_1 + v_0 \cos(\Delta_0 - \Delta_1) \cos(\Phi_0 - \Phi_1)\right]. \tag{51}$$

Let us finally give the expression for the mixing-induced CP asymmetry (40):

$$\mathcal{A}_{\rm CP}^{\rm mix} = -\sin \phi - 2\, \frac{z}{n}, \tag{52}$$

where

$$\begin{aligned} z = {} & [v_0 \cos \Delta_0 \sin \Phi_0 + v_1 \cos \Delta_1 \sin \Phi_1] \cos \phi \\ & + v_0 v_1 \left[\cos(\phi + \Phi_1) \sin \Phi_0 + \cos(\phi + \Phi_0) \sin \Phi_1\right] \cos(\Delta_0 - \Delta_1) \\ & + v_0^2 \cos(\phi + \Phi_0) \sin \Phi_0 + v_1^2 \cos(\phi + \Phi_1) \sin \Phi_1, \end{aligned} \tag{53}$$

and

$$\begin{aligned} n = {} & 1 + 2\, v_0 \cos \Delta_0 \cos \Phi_0 + 2\, v_1 \cos \Delta_1 \cos \Phi_1 \\ & + 2\, v_0 v_1 \cos(\Delta_0 - \Delta_1) \cos(\Phi_0 - \Phi_1) + v_0^2 + v_1^2. \end{aligned} \tag{54}$$

5 Discussion

Because of large uncertainties, the data on $B \to \phi K$ decays that are available at present from the Belle and CLEO collaborations (see (1)) do not allow us to speculate on new physics. However, the experimental situation should significantly improve in the next couple of years. In this section, we discuss various patterns of the observables $\mathcal{S}$, $\mathcal{D}$ and $\mathcal{B}$ introduced above that may shed light both on new physics and on the $B \to \phi K$ hadron dynamics.

An interesting probe for new physics is of course also provided by the mixing-induced CP asymmetry of the $B_d \to \phi K_{\rm S}$ channel, which can be compared with the one of the "gold-plated" mode $B_d \to J/\psi K_{\rm S}$ [5]. Using the expression for $\mathcal{A}_{\rm CP}^{\rm mix}(B_d \to J/\psi K_{\rm S})$ given in [8], the parametrization for $\mathcal{A}_{\rm CP}^{\rm mix}(B_d \to \phi K_{\rm S})$ given in (52), and the hierarchy arising in (29), we obtain

$$\mathcal{A}_{\rm CP}^{\rm mix}(B_d \to \phi K_{\rm S}) - \mathcal{A}_{\rm CP}^{\rm mix}(B_d \to J/\psi K_{\rm S}) = \underbrace{\mathcal{O}(1)}_{{\rm NP}_{I=0}} + \underbrace{\mathcal{O}(\overline{\lambda})}_{{\rm NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\rm SM}, \tag{55}$$

where the $\sin \phi$ terms, which may deviate from the Standard-Model expectation [18], cancel. The contributions entering at the $\overline{\lambda}$ and $\overline{\lambda}^2$ levels may also contain new-physics effects from $B_d \to J/\psi K_{\rm S}$, whereas the $\mathcal{O}(1)$ term would essentially be due to new physics in $B_d \to \phi K_{\rm S}$.

In order to disentangle the $I = 0$ and $I = 1$ isospin sectors, the observables $\mathcal{S}$, $\mathcal{D}$ and $\mathcal{B}$ play a key role. As can be seen in (45)–(47), $\mathcal{S}$ provides a "smoking-gun" signal for new-physics contributions to the $I = 0$ amplitude, whereas $\mathcal{D}$ and $\mathcal{B}$ probe new-physics effects in the $I = 1$ sector. Using the hierarchy arising in (29), we obtain

$$\mathcal{S} = \underbrace{\mathcal{O}(1)}_{\mathrm{NP}_{I=0}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{SM}}, \quad \mathcal{D} = \underbrace{\mathcal{O}(\overline{\lambda})}_{\mathrm{NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{SM}}, \quad \mathcal{B} = \underbrace{\mathcal{O}(\overline{\lambda})}_{\mathrm{NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{SM}}, \tag{56}$$

where the Standard-Model contributions are not under theoretical control. If the dynamical suppression of the $I = 1$ contributions were larger than $\mathcal{O}(\overline{\lambda})$, $\mathcal{D}$ and $\mathcal{B}$ would be further suppressed with respect to $\mathcal{S}$. On the other hand, if the rescattering effects described by (15) were very large – and not small, as assumed in (56) – *all* observables would be of $\mathcal{O}(1)$. In such a situation, we would not only have signals for physics beyond the Standard Model, but also for large rescattering processes.

The discussion given above corresponds to the most optimistic scenario concerning the generic strength of the new-physics effects in the $B \to \phi K$ system. Let us now consider a more pessimistic case, where the new-physics contributions are smaller by a factor of $\mathcal{O}(\overline{\lambda})$:

$$A(B \to \phi K) = \mathcal{A}_{\mathrm{SM}}^{(0)}\Big[1 + \underbrace{\mathcal{O}(\overline{\lambda})}_{\mathrm{NP}_{I=0}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{SM}}\Big]. \tag{57}$$

Now the new-physics contributions to the $I = 1$ sector can no longer be separated from the Standard-Model contributions. However, we would still get an interesting pattern for the $B \to \phi K$ observables, providing evidence for new physics: whereas (55) and $\mathcal{S}$ would both be sizeable, i.e. of $\mathcal{O}(10\%)$ and within reach of the B-factories, $\mathcal{D}$ and $\mathcal{B}$ would be strongly suppressed. However, if these two observables, in addition to (55) and $\mathcal{S}$, are found to be also at the 10% level, new physics cannot be distinguished from Standard-Model contributions, which could also be enhanced to the $\overline{\lambda}$ level by large rescattering effects. This would be the most unfortunate case for the search of new-physics contributions to the $B \to \phi K$ decay amplitudes. An analogous discussion of the $B \to J/\psi K$ system was given in [8].

6 Summary

The decays $B^{\pm} \to \phi K^{\pm}$ and $B_d \to \phi K_{\mathrm{S}}$ offer an interesting probe to search for physics beyond the Standard Model. Using estimates borrowed from effective field theory, we have explored the impact of new physics on the $B \to \phi K$ system in a model-independent manner, and have derived parametrizations of the corresponding decay amplitudes, which rely only on the $SU(2)$ isospin symmetry of strong interactions. We have introduced – in addition to the mixing-induced CP asymmetry of the $B_d \to \phi K_{\mathrm{S}}$ channel – a set of three observables, providing the full picture of possible new-physics effects in $B \to \phi K$ decays. In particular, these observables allow us to separate the new-physics contributions to the $I = 0$ and $I = 1$ isospin sectors, and offer, moreover, valuable insights into hadron dynamics. Imposing dynamical hierarchies of amplitudes, we have discussed various patterns of these observables, including also scenarios corresponding to small and large rescattering effects. We find that $B \to \phi K$ decays may provide, in general, powerful "smoking-gun" signals for new physics. However, there is also an unfortunate case, where such effects cannot be distinguished from those of the Standard Model. We strongly encourage our experimental colleagues to focus not only on the measurement of the mixing-induced CP asymmetries in $B_d \to \phi K_S$ and $B_d \to J/\psi K_S$, but also on the observables $\mathcal{S}$, $\mathcal{D}$ and $\mathcal{B}$. Hopefully, these will yield evidence for physics beyond the Standard Model.

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