Right handed currents and FSI phases in $B^0 \rightarrow \phi K^{*0}$

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Abstract

We consider possible effects of New Physics (NP) on the angular distributions of the decay $B^0 \rightarrow \phi K^{*0}$, showing how these effects depend on the nature of nonstandard interactions. In a general framework based on factorization, we show that triple products can be used to probe the chirality of NP currents. In this analysis we take into account the presence of non-vanishing strong phases, which is motivated by recent experimental evidence. It is seen that the observability of right-handed NP is strongly dependent on the relation between the relative magnitude of these phases and the ratio of Standard Model and NP scales. As an application we estimate the expected values of relevant observables in a particular Left Right Symmetric Model.

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I Introduction

The theoretical and experimental study of B meson physics offers a good opportunity to get new insight about the origin of CP violation. Taking into account the rich variety of decays channels, one can look for many different observables, providing stringent tests for the consistency and stoutness of the CKM mechanism of CP violation proposed by the Standard Model (SM). These analyses may definitely unveil the presence of New Physics (NP) beyond the SM, or provide hints for future searches. In the last few years, an important goal has been achieved through the measurements of the time dependent CP asymmetry in \( B \rightarrow J/\psi K_S \) decays by BABAR \cite{1} and Belle \cite{2}, which have firmly established the violation of CP in the \( B \) system. Within the SM, these experiments allow to get a “clean” measurement of the value of \( \sin(2\beta) \), where \( \beta \) is one of the angles of the so-called unitarity triangle. The present world average value for this quantity is \cite{3}

\[
\sin(2\beta)_{J/\psi K_S} = 0.731 \pm 0.056 ,
\]

which is in agreement with SM expectations. However, the success of the SM is not clear in the case of CP asymmetries in the channel \( B \rightarrow \phi K_S \), where recent measurements of \( \sin(2\beta) \) lead to

\[
\begin{align*}
\sin(2\beta)_{\phi K_S} &= 0.47 \pm 0.34 \quad \text{BABAR} \cite{4} \\
\sin(2\beta)_{\phi K_S} &= -0.96 \pm 0.51 \quad \text{Belle} \cite{5} ,
\end{align*}
\]

showing an apparent disagreement with the previous result. It is not unreasonable to expect this discrepancy to be originated by the presence of NP contributions. Indeed, NP effects are likely to be more important in the case of \( B \rightarrow \phi K_S \) than \( B \rightarrow J/\psi K_S \) since in the SM the latter is governed by a tree level amplitude, while the former is shown to occur dominantly through penguin-like processes.

In this work we propose to explore the NP hypothesis by considering a different channel, namely the decay \( B^0 \rightarrow \phi K^{*0} \), and its CP conjugate, \( \bar{B}^0 \rightarrow \phi \bar{K}^{*0} \). These decays, which are driven by the same quark level processes as \( B \rightarrow \phi K_S \), are found to be attractive for several reasons. Among the various charmless \( B \rightarrow VV \) channels, the neutral and charged \( B \rightarrow \phi K^* \) decays are the first ones that have been experimentally observed \cite{6}. In addition, \( B \rightarrow VV \) processes offer the possibility of measuring many different observables, taking into account the angular distributions of the final outgoing states. In fact, even if \( B^0 \rightarrow \phi K^{*0} \) is a neutral decay, various asymmetries that are relevant for the study of CP violation can be determined with no need of either flavor tagging or time-dependent measurements, hence the experimental analyses are considerably simplified. Detailed studies of the angular distributions have been performed recently by BABAR \cite{7} and Belle \cite{8} collaborations, leading to some puzzling results that deserve significant theoretical interest \cite{9}. Although the corresponding experimental errors are still relatively large so as to unveil the presence of NP, the channel is a promising one to encourage the search. In this sense, it is worth to notice that within the SM the \( B^0 \rightarrow \phi K^{*0} \) decay amplitude has the feature of being dominated by penguin contributions carrying approximately a common weak phase. Therefore, CP-violating observables
are expected to be suppressed within the SM, enhancing the possibility of finding signals of NP.

If the observation of NP effects is confirmed, the various observables to be measured at $B$ factories should be used to distinguish between possible extensions of the SM and to determine the corresponding parameters $[10]$. In this article we point out that the angular distributions of decay products in $B^0 \rightarrow \phi K^{*0}$ could provide not only evidences of physics beyond the SM $[11]$ but also relevant information about the nature of this NP. In particular, using the factorization approach $[12]$, we show that it is possible to perform a chirality test $[13]$ for the structure of nonstandard effective current-current operators through the measurement of the so-called triple products $[14]$. In our analysis we take into account the presence of non-vanishing strong FSI phases, which is stimulated by recent experimental results reported by BABAR and Belle $[7,8]$. In this regard, we point out that the expected order of magnitude of the relevant CP-odd quantities depends on the interplay between the scale of NP and the size of strong phases. The proposed chirality test can be used to distinguish between NP effects arising from right-handed currents (e.g. those predicted by left-right symmetric models) from those which come from SM-like operators, as one would find for instance in the case of the Minimal Supersymmetric Standard Model (MSSM), or multi-Higgs models with $SU(2)_L \otimes U(1)_Y$ electroweak gauge symmetry. Finally, in order to illustrate the potential significance of our analysis, we give an estimation of the size of the expected effects in the case of a Left-Right Symmetric Model (LRSM).

This paper is organized as follows: in Sec. II we write the effective Hamiltonian in the presence of NP right-handed currents, we identify different contributions to the amplitude in the helicity basis and we set the hypotheses to be used in later analysis. In Sec. III we explain how the angular analysis of the decay is used, within factorization, to build the appropriate observables to test the chirality of NP currents. Some of these observables are estimated in Sec. IV within a particular left-right symmetric model. Finally, in Sec. V we state our conclusions.

II Amplitudes and phases in $B \rightarrow VV$ decays

From the theoretical point of view, the standard way of dealing with nonleptonic $B$ decays is based on the effective Hamiltonian approach, which makes use of the Operator Product Expansion as a fundamental tool. Within the SM, this program has been developed in detail, including next-to-leading order calculations $[15]$. After integrating out the degrees of freedom corresponding to all particles with masses above the $b$ quark scale, the low-energy effective Hamiltonian responsible for $b \rightarrow s$ transitions can be written as

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} \left[ V_{cb}^* V_{cs} \sum_{i=1,2} C_i O_i + V_{tb}^* V_{ts} \sum_{i=3\ldots10,\gamma} C_i O_i \right].$$

(3)

where $O_i$ are local current-current quark operators, and $C_i$ are the corresponding Wilson coefficients evaluated at a renormalization scale $\mu \approx m_b$. Following the notation of Ref. $[10]$,...
here $O_{1,2}$ are standard current-current operators, $O_{3...6}$ and $O_{7...10}$ stand for QCD and EW penguin operators, respectively, and $O_{8}$ ($O_{9}$) denotes the gluonic (photonic) magnetic dipole operator. Explicit expressions for both the effective operators and the Wilson coefficients within the SM can be found in several articles and reports (see e.g. Refs. [16, 17]) and we will not repeat them here. Clearly, this effective Hamiltonian would be in general modified if one allows for the presence of NP. It could happen, for instance, that NP effects only show up effectively through SM-like operators, hence the new contributions would just add some new terms to the coefficients $C_i$. This is indeed the situation in the context of some popular schemes, such as the Two-Higgs doublet model and the MSSM. Conversely, it might happen that NP contributions give rise to new effective operators, to be added to the $O_i$‘s in Eq. (3). Here we will address the particular case in which NP manifests effectively through right handed currents, as e.g. in Left-Right (LR) symmetric models [18]. Within these models one would find new operators $O_i'$, obtained from the $O_i$‘s through the interchange $1 \pm \gamma_5 \rightarrow 1 \mp \gamma_5$. Moreover, in this case one should add to $H_{\text{eff}}$ four new tree operators, namely

$$O_{11} = \frac{m_b}{m_c} \langle \bar{s}_\alpha \gamma_\mu (1 - \gamma_5) c^\beta \rangle \langle \bar{c}_\beta \gamma^\mu (1 + \gamma_5) b^\alpha \rangle, \quad (4)$$

$$O_{12} = \frac{m_b}{m_c} \langle \bar{s}_\alpha \gamma_\mu (1 - \gamma_5) c^\alpha \rangle \langle \bar{c}_\beta \gamma^\mu (1 + \gamma_5) b^\beta \rangle, \quad (5)$$

plus $O_{11}'$, $O_{12}'$, which are obtained from the previous two by interchanging $1 \pm \gamma_5 \rightarrow 1 \mp \gamma_5$. Thus the effective $\Delta B = 1$ Hamiltonian will be given by

$$H_{\text{eff}} = H^{(\text{SM})} + H^{(\text{NP,L})} + H^{(\text{NP,R})}$$

$$= \sum_{i=1...10,g,\gamma} \tilde{C}_i O_i + \sum_{i=1...12,g,\gamma} \delta \tilde{C}_i O_i + \sum_{i=1...12,g,\gamma} \delta \tilde{C}_i' O_i', \quad (7)$$

where the NP coefficients $\delta \tilde{C}_i$ and $\delta \tilde{C}_i'$ can be calculated as before using renormalization group equations in the extended LR symmetric theory. In Eq. (7) we have added a tilde to the Wilson coefficients, which for simplicity have been redefined absorbing the common factor $G_F/\sqrt{2}$ and the corresponding CKM matrix elements. Notice that $\delta \tilde{C}_i$ correspond to contributions driven by both SM-like operators ($i = 1...10, g, \gamma$) and nonstandard tree operators ($O_{11,12}$).

Given the relevant effective Hamiltonian, the decay amplitude of a $B$ meson into two vector mesons $V_1$ and $V_2$ in a definite helicity state is written as $A_\lambda = \langle V_1(\lambda) V_2(\lambda) | H_{\text{eff}} | B \rangle$, where $\lambda = 0, \pm 1$ is the helicity of both $V_1$ and $V_2$. According to its Lorentz structure, this amplitude can be parameterized as [19]

$$A_\lambda = \epsilon^*_\mu_1(\lambda) \epsilon^*_{\mu_2}(\lambda) \left[ \frac{a g^{\mu\nu}}{m_1 m_2} p^\mu p'^\nu + \frac{i c}{m_1 m_2} \epsilon^{\mu\nu\alpha\beta} p_{1\alpha} p_{2\beta} \right], \quad (8)$$

where $p$ is the four-momentum of the $B$, while $m_i$, $p_i$ and $\epsilon_i$ stand for the masses, momenta and polarization vectors of the $V_i$ mesons respectively.

In general, the parameters $a$, $b$ and $c$ are complex numbers which arise from the sum of various interfering contributions. Each of these contributions will carry in principle different
strong (CP-conserving) and weak (CP-violating) phases, which we will denote as usual by $\delta_i$ and $\phi_i$ respectively. Now, for the sake of clarity, let us assume that each term in the decomposition of $H_{\text{eff}}$ in Eq. (6) provides a single dominant contribution (we recall that within the SM this is the situation in the case of $B^0 \to \phi K^{*0}$). In general, this does not need to be the case, but the introduction of more terms will not alter our conclusions qualitatively.

Labelling the (SM), (NP,L) and (NP,R) contributions in Eq. (6) with 1, 2 and 3 respectively, we write $a, b$ and $c$ as

$$a = (a_1 e^{i(\tilde{\delta}_1 + \phi_1)} + a_2 e^{i(\tilde{\delta}_2 + \phi_2)} + a_3 e^{i(\tilde{\delta}_3 + \phi_3)}) e^{i\delta_A}.$$  

$$b = (b_1 e^{i(\tilde{\delta}_1 + \phi_1)} + b_2 e^{i(\tilde{\delta}_2 + \phi_2)} + b_3 e^{i(\tilde{\delta}_3 + \phi_3)}) e^{i\delta_B}.$$  

$$c = (c_1 e^{i(\tilde{\delta}_1 + \phi_1)} + c_2 e^{i(\tilde{\delta}_2 + \phi_2)} + c_3 e^{i(\tilde{\delta}_3 + \phi_3)}) e^{i\delta_C}. \tag{9}$$  

Here we have distinguished between strong phases originated from high energy absorptive contributions, denoted by $\tilde{\delta}_i$, and those arising from low energy final state interactions (FSI), $\delta_{a,b,c}$. It is reasonable to assume that the former come only through the Wilson coefficients, hence they are common to $a$, $b$ and $c$ for each contribution 1, 2 and 3. The same argument holds also for the weak phases $\phi_i$. On the contrary, while FSI interaction phases may be different for $a$, $b$ and $c$, we expect them to appear in each case as global factors, since low energy FSI should be blind to the decomposition into standard and nonstandard quark-level operators. In Eq. (9) there is still an overall phase ambiguity, which allows e.g. to fix the SM phase $\tilde{\delta}_1 + \phi_1 = 0$ without loss of generality. Here we have kept all phases as nonvanishing for the sake of clarity in the forthcoming results. It is important to notice that the contributions 2 and 3, arising from NP, are expected to be suppressed with respect to the SM contribution 1. The order of magnitude of the corresponding suppression factor, say $\xi$, will be given by the ratio between typical standard and new physics energy scales.

A key point to be remarked is the fact that in Eq. (9) the Lorentz tensors associated with coefficients $a$ and $b$ have an opposite behavior under parity in comparison with the pseudotensor associated with $c$. We will see below that this can be used to disentangle the different NP contributions to the effective Hamiltonian, in particular, allowing to probe the magnitude of the contribution of right-handed currents in left-right symmetric extensions of the SM. It is crucial to find adequate observables to achieve this disentanglement, and this is the subject of the following Section.

### III Strong phases and observables sensitive to right-handed NP

Let us consider a decay of a $B$ into two vector mesons, $B \to V_1 V_2$, followed by the decays $V_1 \to P_1 P'_1$ and $V_2 \to P_2 P'_2$, where $P_i, P'_i$ are pseudoscalars. The normalized differential
angular distribution can be written as [20]

\[
\frac{1}{\Gamma_0} \frac{d^3\Gamma}{d\cos\theta_1 \ d\cos\theta_2 \ d\psi} = \frac{9}{8\pi N} \left\{ |A_0|^2 \cos^2\theta_1 \cos^2\theta_2 + \frac{|A_\parallel|^2}{2} \sin^2\theta_1 \sin^2\theta_2 \cos^2\psi \right.
\]
\[
+ \frac{|A_\perp|^2}{2} \sin^2\theta_1 \sin^2\theta_2 \sin^2\psi + \frac{\text{Re}(A_\parallel A_\parallel^*)}{2\sqrt{2}} \sin 2\theta_1 \sin 2\theta_2 \cos\psi \right.
\]
\[- \frac{\text{Im}(A_\perp A_\perp^*)}{2\sqrt{2}} \sin 2\theta_1 \sin 2\theta_2 \sin\psi - \frac{\text{Im}(A_\parallel A_\parallel^*)}{2} \sin^2\theta_1 \sin^2\theta_2 \sin 2\psi \right\}
\]

(10)

where \( \theta_1 (\theta_2) \) is the angle between the three-momentum of \( P_1 (P_2) \) in the \( V_1 (V_2) \) rest frame and the three-momentum of \( V_1 (V_2) \) in the \( B \) rest frame, \( \psi \) is the angle between the planes defined by the \( P_1 P_1' \) and \( P_2 P_2' \) three-momenta in the \( B \) rest frame, and we have defined \( N \equiv |A_0|^2 + |A_\perp|^2 + |A_\parallel|^2 \). The amplitudes \( A_\perp \) and \( A_\parallel \) are related to the helicity amplitudes introduced in Eq. (8) by

\[
A_\perp = \frac{A_{+1} - A_{-1}}{\sqrt{2}}, \quad A_\parallel = \frac{A_{+1} + A_{-1}}{\sqrt{2}},
\]

(11)

while \( A_0 \) is common to both bases. In addition, the observables in Eq. (10) can be written in terms of the previously introduced complex parameters \( a, b \) and \( c \) as

\[
|A_0|^2 = |x a + (x^2 - 1) b|^2 \quad \text{Re}(A_\parallel A_\parallel^*) = -\sqrt{2} \left[ x |a|^2 + (x^2 - 1) \text{Re}(a^* b) \right]
\]
\[
|A_\parallel|^2 = 2 |a|^2 \quad \text{Im}(A_\perp A_\parallel^*) = \sqrt{2} \left[ x \text{Im}(a c^*) + (x^2 - 1) \text{Im}(b c^*) \right]
\]
\[
|A_\perp|^2 = 2 (x^2 - 1) |c|^2 \quad \text{Im}(A_\perp A_\perp^*) = 2\sqrt{x^2 - 1} \text{Im}(c a^*),
\]

(12)

where \( x \equiv (m_B^2 - m_1^2 - m_2^2)/(2m_1 m_2) \). Now as usual we denote the amplitudes corresponding to the CP conjugated process by \( \bar{A}_i \). This comes together with the replacements \( a \rightarrow \bar{a}, \ b \rightarrow \bar{b}, \ c \rightarrow -\bar{c} \) in the parameterization of Eq. (8). In addition, following the notation of Ref. [21], we define the combined observables

\[
\Lambda_{ii} = \frac{1}{2} (|A_i|^2 + |\bar{A}_i|^2) , \quad \Sigma_{ii} = \frac{1}{2} (|A_i|^2 - |\bar{A}_i|^2) , \quad \Lambda_{ij} = -\text{Im}(A_i A_j^* - \bar{A}_i \bar{A}_j^*) , \quad \Sigma_{ij} = -\text{Im}(A_i A_j^* + \bar{A}_i \bar{A}_j^*) , \quad \Lambda_{||} = -\text{Re}(A_0 A_0^* + \bar{A}_0 \bar{A}_0^*) , \quad \Sigma_{||} = -\text{Re}(A_0 A_0^* - \bar{A}_0 \bar{A}_0^*) ,
\]

(13)

where \( i = 0, ||, \perp, \) and \( j = 0, || \).

Among the kinematic observables in the decay of a \( B \) into two vector mesons, one has the so-called “triple products” \( \vec{q}_i \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2) \), where \( \vec{\epsilon}_1, \vec{\epsilon}_2 \) are the polarization vectors of the two vector mesons, and \( \vec{q}_i \) is the momentum of one of them. Since the T-transformation reverses both momentum and spin, triple products are T-odd quantities. Regarding the observables defined above, one has that \( \text{Im}(A_\perp A_0^*) \) is proportional to the asymmetry

\[
A_T = \frac{\Gamma(\vec{q}_0 \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2) > 0) - \Gamma(\vec{q}_0 \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2) < 0)}{\Gamma(\vec{q}_0 \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2) > 0) + \Gamma(\vec{q}_0 \cdot (\vec{\epsilon}_1 \times \vec{\epsilon}_2) < 0)} ,
\]

(14)
which is also T-odd. Note that this asymmetry may be not null even in the absence of CP violation, due to the presence of strong phases \([14]\). If, as in the case of the combined observables \(\Lambda_{\perp(\|,0)}\), we define the CP-asymmetry \(A_T = \Lambda_T - \Lambda_T\), where \(\Lambda_T\) is the triple-product asymmetry corresponding to the CP-conjugated decay, then \(A_T\) will be both a T-odd and CP-odd quantity. Since \(A_T\) is an observable of CP violation, it will be also an observable of T violation if CPT invariance is assumed. On the other hand, the combined observable \(\Sigma_{\perp(\|,0)}\) is proportional to \(A_T + \bar{A}_T\), therefore it is a T-odd but CP-even quantity.

Let us focus our attention in the neutral decay \(B^0 \to \phi K^*\) (and the corresponding CP conjugated process, \(\bar{B}^0 \to \phi \bar{K}^*\)) in the context of a model including left- and right-handed quark currents arising from new physics, as discussed in Sec. II. This decay proceeds through the quark level process \(\bar{b} \to \bar{s}s\bar{s}\), which is mainly driven by penguin operators \(O_i\) and \(O'_i\), with \(i = 3, \ldots, 10\) in the notation introduced above. In order to deal with the hadronic matrix element, we will use the well-known factorization approximation, in which one neglects the contribution of color-octet quark transitions. However, we will allow for the presence of nonzero final state interaction phases. In this way, the standard and nonstandard contributions to the amplitude can be factorized as

\[
A = \left[ \tilde{C} + \delta\tilde{C} \right] \langle K^*|\bar{b}\gamma_\mu(1 - \gamma_5)s|B^0\rangle \langle \phi|\bar{s}\gamma_\mu s|0\rangle + \left[ \delta\tilde{C}' \right] \langle K^*|\bar{b}\gamma_\mu(1 + \gamma_5)s|B^0\rangle \langle \phi|\bar{s}\gamma_\mu s|0\rangle ,
\]

(15)

where the factors in square brackets involve combinations of the coefficients \(\tilde{C}_i, \delta\tilde{C}_i\) and \(\delta\tilde{C}'_i\) in Eq. (7) (explicit expressions, which are not important at this point, will be given in the next Section). The information on high energy contributions to the decay process— including SM and NP interactions— is carried by the Wilson coefficients, whereas low energy interactions are taken into account in the hadronic matrix elements. From Eq. (15), it is seen that the factorized matrix elements \(\langle K^*|J_\mu|B^0\rangle\) can be separated into two terms, namely those driven by vector and by axial-vector currents. Taking into account the Lorentz decomposition in Eq. (8), the latter are found to contribute to coefficients \(a\) and \(b\), and the former to \(c\). In this way, one arrives to the following simple relations between the various parameters in Eq. (9):

\[
\frac{a_1}{c_1} = \frac{a_2}{c_2} = -\frac{a_3}{c_3} .
\]

(16)

We recall that the decay \(B^0 \to \phi K^*\) is mainly driven by penguin-like operators. Since \(O_{11,12}\) and \(O'_{11,12}\) do not contribute, the operators corresponding to contributions 2 and 3 are nothing but the SM operators and their parity-conjugated (i.e. those obtained by changing \(\pm \gamma_5 \to \mp \gamma_5\)), respectively. Hence the sign difference in Eq. (16) will allow in principle to distinguish between NP contributions arising from effective left-handed (i.e. SM-like) and right-handed currents.

Regarding Eqs. (12) and (13), it is clear that the above relations between the contributions to \(a, b\) and \(c\) will become particularly relevant for the triple products \(\Lambda_{\perp j}\) and \(\Sigma_{\perp j}\), with \(j = ||, 0\). Let us focus first on \(\Lambda_{\perp ||}\) and \(\Sigma_{\perp ||}\), and then comment on the other two. As stated, with good approximation the process under study carries only one global weak phase within
the SM, therefore CP-odd observables will vanish in the absence of NP. This is the case of $\Lambda_{\perp \parallel}$. From Eq. (9), using the relations in Eq. (16) one has

$$
\Lambda_{\perp \parallel} = -8\sqrt{x^2-1} \left[ a_1 c_2 \sin(\delta_a - \delta_c) \sin(\delta_2 - \delta_1) \sin(\phi_2 - \phi_1) + 
\quad a_1 c_3 \cos(\delta_a - \delta_c) \cos(\delta_3 - \delta_1) \sin(\phi_3 - \phi_1) - 
\quad a_2 c_3 \cos(\delta_a - \delta_c) \cos(\delta_3 - \delta_2) \sin(\phi_3 - \phi_2) \right].
$$

(17)

In the framework of factorization, one expects strong FSI phases to be relatively small. However, notice that $\Lambda_{\perp \parallel}$ may be still non-vanishing even in the limit where these are zero. This is a particular feature of triple products, and an equivalent observable is not to be found in $PP$ or $VP$ decay channels of $B$ mesons. As stated, another relevant observable (in this case CP-even) is $\Sigma_{\perp \parallel}$, which in our framework is given by

$$
\Sigma_{\perp \parallel} = 8\sqrt{x^2-1} \left[ \frac{1}{2} (a_1 c_1 + a_2 c_2 + a_3 c_3) \sin(\delta_a - \delta_c) + 
\quad a_1 c_2 \sin(\delta_a - \delta_c) \cos(\delta_2 - \delta_1) \cos(\phi_2 - \phi_1) - 
\quad a_1 c_3 \cos(\delta_a - \delta_c) \sin(\delta_3 - \delta_1) \cos(\phi_3 - \phi_1) + 
\quad a_2 c_3 \cos(\delta_a - \delta_c) \sin(\delta_2 - \delta_3) \cos(\phi_2 - \phi_3) \right].
$$

(18)

Notice that this observable is different from zero only in the presence of non-vanishing strong phases. In particular, within the SM one needs FSI phases $\delta_a$ or $\delta_c$ to be nonzero—and to be different from each other—, therefore the value of $\Sigma_{\perp \parallel}$ is regarded as a measure of FSI effects [7, 8].

As stated, within factorization it is natural to assume that strong FSI phases are small, i.e. $\delta \ll 1$. Moreover, we will also assume $\tilde{\delta} \ll 1$, since these phases arise in general from absorptive contributions which are higher order in perturbation theory. Then from Eq. (17) it can be seen that, depending on the nature of NP contributions, one obtains sizable different behaviors in $\Lambda_{\perp \parallel}$. Indeed, the contribution to $\Lambda_{\perp \parallel}$ arising from a combination of coefficients of the same chirality [first term on the r.h.s. of Eq. (17)] is suppressed by a factor $\sim \delta \tilde{\delta}$ with respect to the remaining two (chirality-mixing) terms. In this sense, it could be said that $\Lambda_{\perp \parallel}$ behaves as a filter of right-handed NP. On the other hand, in the case of $\Sigma_{\perp \parallel}$ it is seen that all contributions are of order $\delta$ or $\tilde{\delta}$. For this observable it is also necessary to take into account the relative magnitude between the various coefficients, the SM term (carrying the combination $a_1 c_1$) being the dominant one. As stated above, NP contributions are expected to be suppressed by a factor $\xi \sim \delta C/C \sim \delta C'/C'$, which leads to $\Sigma_{\perp \parallel} \sim a_1 c_1 [\delta + \mathcal{O}(\delta \xi, \tilde{\delta} \xi)]$. Notice that the leading term in the square brackets is independent of the nature of NP.

In order to make use of the chirality-filtering property of $\Lambda_{\perp \parallel}$ when comparing with experiment, it is desirable to get rid of the coefficients $a_i c_j$ which in general will be hardly determined from the theory. In this way, we find convenient to consider the ratio $r_{\parallel} \equiv \Lambda_{\perp \parallel}/\Sigma_{\perp \parallel}$, which has two main advantageous features: (i) it is sensitive to different NP chiral
structures, and depends on the coefficients through the ratios \(c_2/c_1 \sim c_3/c_1 \sim \xi\) [see Eqs. (19-21)]; (ii) since the order of magnitude of the denominator does not change with the chirality of NP, it does not spoil the filtering properties of \(\Lambda_{\perp\parallel}\), and the ratio \(r_{\parallel}\) becomes enhanced in the case of right-handed NP currents. Indeed, one could have three different situations upon the NP scenario considered:

- **SM** \(\rightarrow r_{\parallel} \simeq 0\) (19)
- only left-handed NP \(\rightarrow r_{\parallel} \simeq -2(\tilde{\delta}_2 - \tilde{\delta}_1) \frac{c_a}{c_1} \sin(\phi_2 - \phi_1) \sim \xi \tilde{\delta}\) (20)
- right-handed NP \(\rightarrow r_{\parallel} \simeq -\frac{2}{\delta_a - \delta_c} \frac{c_3}{c_1} \sin(\phi_3 - \phi_1) \sim \xi / \delta\). (21)

Two main conclusions come out from these relations: first, a non-vanishing measurable value of \(r_{\parallel}\) represents a signature of NP. Second, this observable is particularly sensitive to the presence of NP if nonstandard interactions lead to effective right-handed quark currents. Notice the interplay between the characteristic scale of NP and the strong phases involved in the decay: if FSI phases \(\delta\) are small, but still measurable, the effect of right-handed currents should be observable when \(\xi \gtrsim \delta\). In contrast, SM-like contributions to \(r_{\parallel}\) arising from NP at this same scale would be additionally suppressed by strong phases originated in the high energy region. In this way, the possible detection of right-handed NP through this observable is intimately associated with the relation between the NP scale and the FSI phases, \(\xi \leftrightarrow \delta\). We notice at this point that a comparison of the predictions for \(r_{\parallel}\) with experimental data is not possible yet, since present measurements \([7, 8]\) of \(\Sigma_{\perp\parallel}\) give values which are still compatible with zero. However, given the potential power of \(B\) factories, we expect that the experimental analyses will be able to set bounds for \(r_{\parallel}\), providing a clue about the nature of NP.

Let us now comment on the other two triple products, \(\Lambda_{\perp 0}\) and \(\Sigma_{\perp 0}\), and also the ratio \(r_0 \equiv \Lambda_{\perp 0}/\Sigma_{\perp 0}\). Regarding the definitions in Eqs. (12-13), it is seen that a similar behavior to their cousins \(\Lambda_{\perp\parallel}\), \(\Sigma_{\perp\parallel}\) should be expected, even if the results will be more involved since the amplitude \(A_0\) receives contributions from both Lorentz invariant parameters \(a\) and \(b\). In fact, it is easy to see that \(\Lambda_{\perp 0}\) acts as a filter of right-handed NP currents, in a similar way as \(\Lambda_{\perp\parallel}\), in the limit of small strong phases. The ratio \(r_0\), upon the nature of the NP, behaves as follows:

- **SM** \(\rightarrow r_0 \simeq 0\) (22)
- only left-handed NP \(\rightarrow r_0 \simeq -2(\tilde{\delta}_2 - \tilde{\delta}_1) \frac{c_a}{c_1} \sin(\phi_2 - \phi_1) \sim \xi \tilde{\delta}\) (23)
- right-handed NP \(\rightarrow r_0 \simeq \frac{2[xa_1 + (x^2 - 1)b_1]}{x(\delta_c - \delta_a)a_1 + (x^2 - 1)(\delta_c - \delta_b)b_1} \frac{c_3}{c_1} \sin(\phi_3 - \phi_1)
\sim \xi / \delta\). (24)

Thus, the same qualitative behaviour found for \(r_{\parallel}\), c.f. Eqs. (19-21), holds for \(r_0\) as well. Notice, however, that the combinations of the parameters \(a_1\) and \(b_1\) and the FSI phases in
Eq. (24) may lead to cancellations which modify the naively expected orders of magnitude. Thus, in this sense, \( r_0 \) (which represents a better observable from the experimental point of view) will be not as conclusive as \( r_1 \). We recall that present experiments \([8,9]\) have measured a nonzero value for \( \Sigma_{\perp 0} \), allowing to establish first experimental bounds on the value of \( r_0 \), namely \( r_0 = -0.14 \pm 0.21 \) \([8]\).

To conclude this section, some final remarks are in order. First, notice that the chirality-mixing suppression in the \( \Lambda_{\perp j} \)'s could be modified by the effect of large nonfactorizable contributions, which have been neglected in our analysis. The consistency of the factorization approximation can be studied by performing a global analysis of the various observables that can be measured in this decay channel, as well as in related processes such as \( B^\pm \to \phi K^{*\pm} \) \([22]\). Second, we point out that a key feature of the neutral channel analyzed here is that it has a unique way to factorize the hadronic currents in the factorization approximation, and this fact makes possible the chirality test. As a matter of fact, this same analysis could be applied to the channels \( B^0 \to K^0 K^{*0} \), \( B^0 \to \phi \phi \), \( B^0 \to \rho^0 \phi \) and \( B^0 \to \omega \phi \), which have not been experimentally observed yet. Finally, as mentioned in Sect. II, we stress that the existence of other left- or right-handed NP contributions may be still worked out within this framework leaving our conclusions unchanged, since the property of filtering right-handed currents relies on the structure of the observables \( \Lambda_{\perp j} \) and does not depend on the number of contributions considered.

### IV Numerical estimations in a LR symmetric model

Theories based upon the electroweak gauge group \( SU(2)_L \times SU(2)_R \times U(1) \) represent well-known extensions of the Standard Model. We analyze here the so-called Left-Right Symmetric Model (LRSM) \([18]\), in which the elements of left and right quark mixing matrices are equal in modulus, i.e. \( |V_{ij}^L| = |V_{ij}^R| \). This model has been largely discussed in connection with penguin-dominated decays \( \bar{b} \to \bar{s}q\bar{q} \) \([23]\), and serves as a good example to illustrate the special features of the observables \( \Lambda_{\perp l} \) and \( \Sigma_{\perp l} \) (and their ratio \( r_l \)) considered in the preceding Sections.

The effective Hamiltonian for \( \bar{b} \to \bar{s}q\bar{q} \) decays within the LRSM has been calculated in LL precision by Cho and Misiak \([24]\). Keeping only the top and bottom quark masses as non-vanishing, the matching conditions at the \( W \) scale lead to

\[
\begin{align*}
C_2(M_1) &= 1, & C'_2(M_1) &= 0, \\
C_7(M_1) &= D'_0(x) + A^*_t \bar{D}'_0(x), & C'_7(M_1) &= A^*_t \bar{D}'_0(x), \\
C_9(M_1) &= E'_0(x) + A^*_t \bar{E}'_0(x), & C'_9(M_1) &= A^*_t \bar{E}'_0(x),
\end{align*}
\]

(25)

whereas the remaining coefficients \( C_i \) and \( C'_i \) are equal to zero. Here \( M_1 \) is the mass of the \( W^\pm \) gauge bosons (equivalent to the standard \( W^\pm \)), \( E'_0(x) \) and \( D'_0(x) \) are SM Inami-Lim

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\( ^1 \)Notice that the experimental values reported in Ref. \([8]\) result from the average between the data obtained from \( B^0 \to \phi K^{*0} \) and \( B^\pm \to \phi K^{*\pm} \), which is justified if the annihilation contribution to \( B^\pm \to \phi K^{*\pm} \) can be safely neglected. If this is not the case \([22]\), one should take into account for our purposes only the data corresponding to the neutral channel.
functions, and their left-right analogues are denoted by $\tilde{E}_0'(x)$ and $\tilde{D}_0'(x)$ (see Ref. [21]). These functions are evaluated at $x = m_t^2/M_W^2$, while the coefficients $A_{tb}$ and $A_{ts}$ are given in terms of LRSM parameters,

$$A_{tb} = \xi \frac{m_t}{m_b} \frac{V_{tb}^R}{V_{tb}^L} e^{i\eta} \equiv \xi \frac{m_t}{m_b} e^{i\sigma_1}$$  \hspace{1cm} (26)

$$A_{ts} = \xi \frac{m_t}{m_b} \frac{V_{ts}^R}{V_{ts}^L} e^{i\omega} \equiv \xi \frac{m_t}{m_b} e^{i\sigma_2} .$$  \hspace{1cm} (27)

Here $\xi$ is the $W_L$-$W_R$ mixing angle times the ratio between right and left coupling constants, while $\sigma_{1,2}$ are unknown CP-violating phases that take values in the range $[0,2\pi]$ (for simplicity we take the SM amplitude to be real). The mass of the gauge bosons $W^+\gamma$ does not appear here explicitly since its lower bound, in the multi-hundred GeV region, leads to negligible contributions to the effective Hamiltonian.

The renormalization group mixing is governed by a $20 \times 20$ anomalous dimension matrix $\gamma$, which decomposes into two identical $10 \times 10$ submatrices. The expression for the SM $8 \times 8$ submatrix can be found in Ref. [25] and the rest of the entries have been computed in Ref. [24]. The running of the Wilson coefficients to the $m_b$ scale in the LL approximation for five active flavors yields

$$C_i(\mu = m_b) = \sum_{k,l} (S^{-1})_{ik} \eta \frac{3\lambda_k}{4} S_{kl} C_i(M_1) ,$$  \hspace{1cm} (28)

where the $\lambda_k$’s in the exponent of $\eta = \alpha_s(M_1)/\alpha_2(m_b)$ are the eigenvalues of the anomalous dimension matrix, and $S$ is a matrix containing the corresponding eigenvectors. The explicit computation of these Wilson coefficients at the $m_b$ scale has been performed in Ref. [26].

In order to evaluate the relevant hadronic matrix elements in the case of the decay $B^0 \rightarrow \phi K^{*0}$, once again we make use of the factorization approximation. Then, using the equations of motion, it is seen that the matrix elements of magnetic dipole operators can be written in terms of the penguin ones as

$$\langle O_8^G \rangle = -\frac{\alpha_s}{4\pi} \frac{m_b}{\sqrt{q^2}} \left[ \langle O_4 \rangle + \langle O_6 \rangle - \frac{1}{N_c} \left( \langle O_3 \rangle + \langle O_5 \rangle \right) \right] ,$$  \hspace{1cm} (29)

$$\langle O_7^G \rangle = -\frac{\alpha}{3\pi} \frac{m_b}{\sqrt{q^2}} \left[ \langle O_7 \rangle + \langle O_9 \rangle \right] ,$$  \hspace{1cm} (30)

and similar equations hold for their primed counterparts. In this way, one can get rid of the magnetic dipole operators when writing the $B^0 \rightarrow \phi K^{*0}$ decay amplitude by including their contributions into the Wilson coefficients $C_i^{\text{eff}}$ and $C_i'^{\text{eff}}$, with $i = 3, \ldots, 10$. The amplitude can be written as [16]

$$A(B^0 \rightarrow \phi K^{*0}) = \langle H_{\text{eff}} \rangle$$

$$= -\frac{G_F}{\sqrt{2}} V_{tb}^L V_{ts}^L \left[ a_3 + a_4 + a_5 - \frac{1}{2}(a_7 + a_9 + a_{10}) \right] X_{L}^{(BK^*,\phi)}$$

$$-\frac{G_F}{\sqrt{2}} V_{tb}^L V_{ts}^L \left[ a_3' + a_4' + a_5' - \frac{1}{2}(a_7' + a_9' + a_{10}') \right] X_{R}^{(BK^*,\phi)} ,$$  \hspace{1cm} (31)
where \( X_{L,R}^{(BK,\phi)} = \langle K^{*0} | \bar{b} \gamma^\mu (1 \mp \gamma_5) s | B^0 \rangle \langle \phi | \bar{s} \gamma_\mu s | 0 \rangle \),
\[
a_{2i-1} = C_{2i-1}^{\text{eff}} + \frac{1}{N_c} C_{2i}^{\text{eff}}, \quad a_{2i} = C_{2i}^{\text{eff}} + \frac{1}{N_c} C_{2i-1}^{\text{eff}},
\]
and similar relations are satisfied by the coefficients \( a'_i \). Taking now \( \sqrt{\langle q^2 \rangle} = m_b/\sqrt{2} \) in Eqs. (30), together with \( m_t/m_b = 60 \) and \( N_C = 3 \), the amplitude reads
\[
A(B^0 \to \phi K^{*0}) \simeq - \frac{G_F}{\sqrt{2}} V_{ts}^L V_{tb}^* \left( -0.019 + 0.12 \xi e^{-i\sigma_1} \right) X_{L}^{(BK,\phi)}
- \frac{G_F}{\sqrt{2}} V_{ts}^L V_{tb}^* \left( 0.12 \xi e^{i\sigma_2} \right) X_{R}^{(BK,\phi)}.
\]
(33)

Notice that, aside from the scale suppression \( \xi \), the NP contributions turn out to be enhanced. This is mainly given by a factor \( m_t/m_b \), which arises in the case of \( V + A \) interactions since the usual helicity flip is not needed in penguin amplitudes.

In order to extract from Eq. (33) the different coefficients and phases in Eq. (9), we write the factorized matrix elements in terms of form factors \( f_V, V^{B \to V}(q^2) \), and \( A_i^{B \to V}(q^2) \), \( i = 0, 1, 2 \):
\[
\langle V(\varepsilon,p)|V_\mu|0\rangle = f_V m_V \varepsilon^*_\mu,
\]
\[
\langle V(\varepsilon,p')|V_\mu|B(p)\rangle = - \frac{2}{m_V + m_B} \epsilon_{\mu\nu\alpha\beta} \varepsilon^*_{\nu} p^\alpha p'^\beta V^{B \to V}(q^2),
\]
\[
\langle V(\varepsilon,p')|A_\mu|B(p)\rangle = i \frac{2 m_V (\varepsilon^* \cdot q)}{q^2} q_\mu A_0^{B \to V}(q^2) + i (m_V + m_B) \left[ \varepsilon^*_\mu - \frac{(\varepsilon^* \cdot q)}{q^2} q_\mu \right] A_1^{B \to V}(q^2)
- i \left[ (p + p')_\mu - \frac{(m_B^2 - m_V^2)}{q^2} q_\mu \right] \frac{(\varepsilon^* \cdot q)}{m_V + m_B} A_2^{B \to V}(q^2).
\]
(34)

Here \( V(\varepsilon,p') \) stands for the outgoing vector mesons \( \phi \) or \( K^{*0} \), \( V_\mu \) and \( A_\mu \) are the corresponding vector and axial-vector quark currents and \( q = p - p' \) is the transferred momentum. The vector and axial-vector form factors can be estimated from the analysis of semi-leptonic \( B \) decays, using the ansatz of pole dominance to account for the momentum dependencies in the region of interest. Taking into account Eqs. (33) and (34) and comparing with Eqs. (8) and (9) one can obtain explicit expressions for the coefficients \( a_i \) and \( c_i \), needed to estimate the values of the observables \( \Lambda_{||} \) and \( \Sigma_{||} \). We find
\[
a_1 = -0.019 \kappa (m_B + m_{K^*}) A_1^{B \to K^*}(m_\phi^2),
\]
\[
c_1 = 0.019 \kappa \left( \frac{2 m_{K^*} m_\phi}{m_B + m_{K^*}} \right) V^{B \to K^*}(m_\phi^2),
\]
\[
a_2 = 0.12 \xi \kappa (m_B + m_{K^*}) A_1^{B \to K^*}(m_\phi^2),
\]
\[
c_2 = -0.12 \xi \kappa \left( \frac{2 m_{K^*} m_\phi}{m_B + m_{K^*}} \right) V^{B \to K^*}(m_\phi^2),
\]
\[
a_3 = -a_2,
\]
\[
c_3 = c_2,
\]
(35)
where \( \kappa = i \frac{G_F}{\sqrt{2}} V_{ts} V_{tb}^* m_\phi f_\phi \). The weak phases in Eq. (9) are given by \( \phi_2 = -\sigma_1 \) and \( \phi_3 = \sigma_2 \), the absorptive phases \( \tilde{\delta}_i \) have been neglected in this approximation, and the strong FSI phases \( \delta \), which are expected to be small, are left as unknown parameters. As expected, the coefficients in Eqs. (35) satisfy the relations in Eq. (16). Finally, the observables \( \Lambda_{\perp\parallel} \) and \( \Sigma_{\perp\parallel} \) can be estimated in the LRSM by making use of Eqs. (17) and (18). This leads to

\[
\Lambda_{\perp\parallel} = -16\sqrt{x^2 - 1} \kappa^2 m_{K^*} m_\phi V_{B^0 \rightarrow K^*} (m_\phi^2) A_{1}^{B \rightarrow K} (m_\phi^2) \times
\left[ 0.019 \times 0.12 \xi \sin \sigma_2 - (0.12)^2 \xi^2 \sin(\sigma_1 + \sigma_2) \right],
\]

\[
\Sigma_{\perp\parallel} = 8\sqrt{x^2 - 1} \kappa^2 m_{K^*} m_\phi V_{B^0 \rightarrow K^*} (m_\phi^2) A_{1}^{B \rightarrow K} (m_\phi^2) (0.019)^2 (\delta_c - \delta_a) + O(\xi \delta).
\]

Eq. (36) deserves some discussion. As it can be easily seen, the dependence of \( \Lambda_{\perp\parallel} \) on \( \sigma_1 \) is suppressed by a factor \( \xi \) with respect to that on \( \sigma_2 \). In fact, this could be expected since the right-handed NP is incorporated in the effective Hamiltonian through the dipolar operators, which are obtained by a mass insertion in the \( b \) quark line, whereas the contribution of the \( s \) quark is neglected. Therefore, although the SM-like NP operators and the right-handed ones are both proportional to \( \xi \), the former depend only on \( \sigma_1 \) and the latter only on \( \sigma_2 \). Since the structure of \( \Lambda_{\perp\parallel} \) is such that suppresses in \( \tilde{\delta}(=0) \) the mixture of terms having the same chirality, the only combined contributions that remain are those given by SM \( \times \) Right-NP \( (\sim \xi) \) and Left-NP \( \times \) Right-NP \( (\sim \xi^2) \).

Using Eqs. (36,37) we obtain the value for \( r_{\parallel} = \Lambda_{\perp\parallel}/\Sigma_{\perp\parallel} \) within the LRSM to leading order in \( \xi \),

\[
r_{\parallel} = \frac{0.12 \xi}{0.019} \frac{2}{\delta_a - \delta_c} \sin \sigma_2.
\]

Notice that in the ratio \( r_{\parallel} \) one gets rid of form factors, which are in general theoretically uncertain. In Eq. (38), \( \sigma_2 \) is a free parameter, hence it is natural to assume \( \sin \sigma_2 \approx O(1) \). The scale of NP is given basically by the mass of the gauge bosons \( W^\pm_2 \); present limits lead to an upper bound for \( \xi \) of about 0.04 [3]. Clearly, the value of \( r_{\parallel} \) turns out to be enhanced by the FSI phases (assumed to be small) in the denominator. As discussed in the previous section, in order to distinguish right-NP from left-NP by means of \( r_{\parallel} \), it is essential to consider the relation between the NP scale and the FSI phases \( \delta \).

As mentioned, present data are not precise enough to place experimental bounds on \( r_{\parallel} \). However, in view of the positive result in the case of \( \Sigma_{\perp0} \), it is natural to expect that future measurements provide a nonzero value for the combination \( \delta_a - \delta_c \), and hence a nonzero \( \Sigma_{\perp\parallel} \). We expect that forthcoming experimental data will be able to place bounds on \( r_{\parallel} \), and in this way quest for the existence of right-handed NP currents as those proposed by the LRSM.

V Conclusions

In this work we have studied possible signatures of New Physics through the angular analysis of the decay \( B^0 \rightarrow \phi K^{*0} \), which is shown to be an optimal channel to search for nonstandard
effects. In particular, we have studied how the effects of left and/or right-handed NP currents show up in some selected observables.

In the construction of the effective Hamiltonian we have considered, besides the SM physics, possible NP contributions to SM-like operators, as well as contributions driven by new operators arising from the presence of right-handed NP currents. We have worked within the framework of factorization, considering for simplicity the case in which one has only one NP contribution of each chirality. Since the decay under study has only one possible way to factorize, the nonstandard contributions are subject to kinematic cancellations in some observables, allowing to obtain a chirality filter for the NP.

The coefficients in the angular distributions of the decay of a $B$ into two vector mesons, such as $B^0 \to \phi K^{*0}$, give a rich variety of observables. We have found that if both the absorptive and FSI strong phases ($\delta$ and $\tilde{\delta}$, respectively) as well as the ratio between SM and NP scales ($\xi$) are assumed to be small quantities, then one can define ratios which show a strongly different behavior upon the nature of the NP. In particular, we have concentrated on the T-odd triple products $\Lambda_{\perp i}$ ($i = 0, \parallel$), which can be written as sums of various terms including contributions from different invariant amplitudes. The key is to realize that those terms that contain a mixture of two contributions coming from operators with the same chiral structure are suppressed by an order of magnitude of $\sim \delta \tilde{\delta}$ with respect to those which arise from the mixture of operators with different chiralities. In this way, they effectively serve as filters of right-handed NP. We have also considered the CP-conserving observables $\Sigma_{\perp i}$, which at leading order behave as $\sim \delta$ independently of the nature of the NP. As explained in the text, we propose to measure the ratio $r_i = \Lambda_{\perp i}/\Sigma_{\perp i}$, which shows a different behavior upon the nature of the NP, namely

$$r_i \sim \begin{cases} 0 & \text{SM} \\ \xi\tilde{\delta} & \text{only Left-NP} \\ \xi/\delta & \text{Right-NP} \end{cases}$$  (39)

Although $r_i$ should be handled with care, since it is a ratio between two small quantities, from Eq. (39) it is seen that this represents an optimal observable for the search of effects of right-handed NP, provided that strong FSI phases $\delta$ are found to be nonzero but sufficiently small. On the other hand, if NP effects are observed in any other process, a measurement of a tiny upper bound for $r_i$ should be taken as a strong indication in favor that the underlying theory leads to effective operators of the same type of the SM. In addition, Eq. (39) shows that the perceptiveness of right-handed NP is strongly dependent on the relation between the relative scale of NP and the strong FSI phases. We notice that although the result in Eq. (39) is clean for the ratio $r_\parallel$, it should be taken with some care for $r_0$, since a fine tuning of the parameters could lead to cancellations modifying the naively expected orders of magnitude. On the experimental side, present measurements of both $\Sigma_{\perp\parallel}$ and $\Lambda_{\perp\parallel}$ are still compatible with zero, and only a positive result has been achieved in the case of $\Sigma_{\perp0}$. We expect forthcoming experiments to establish precise bounds for these observables in the near future, providing clues about the nature of new physics beyond the SM. Our analysis has been complemented with a numerical estimation of the ratio $r_\parallel$ in a LRSM, which shows
that the allowed parameter space for these models is compatible with the presence of large NP effects in penguin-dominated modes like $B^0 \to \phi K^{*0}$.

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