Looking for novel *CP*-violating effects in $\overline{B} \rightarrow K^* l^+ l^-$

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The *CP*-violating asymmetries in the exclusive decay $\overline{B} \to K^* l^+ l^ (l=e,\mu,\tau)$ are predicted to be exceedingly small in the standard model (SM), thereby offering an opportunity to assess various new-physics scenarios. We derive quantitative predictions for various integrated observables in $\overline{B} \to K^* \mu^+ \mu^-$ decay in the presence of physics beyond the SM with additional *CP* phases and an extended operator basis. In particular, a model-independent analysis of *CP* asymmetries that require the presence of unitarity phases, in addition to *CP* violation, is performed. We find that in the low dimuon invariant mass region $2m_{\mu} \leq M_{\mu^+\mu^-} < M_{J/\psi}$, the *CP* asymmetries are highly suppressed by small dynamical phases, assuming that new physics is unlikely to significantly alter the Wilson coefficients of the operators governing two-body hadronic *B* decays. Taking into account current experimental data on the measured $b \rightarrow s \gamma$ rate and the upper limit on $\mathcal{B}(B^0 \rightarrow K^{*0}\mu^+\mu^-)$, *CP*-violating effects of a few percent are estimated, even in the presence of new physics with *CP* phases of O(1). By contrast, in the high dimuon invariant mass region $M_{\psi'} < M_{\mu^+\mu^-} \leq (M_B - M_{K^*})$ significant *CP*-violating effects are possible. Given a branching ratio of 1.8×10^{-6} , the *CP* asymmetries can be quite substantial (~20% or more), and thus may serve as a means of discovering physics transcending the SM.

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I. INTRODUCTION

CP violation has been observed so far only in the neutral kaon system. Within the standard model (SM), the experimental results on indirect (ϵ_K) and direct (ϵ'/ϵ_K) *CP* violation can be explained by the complex phase of the Cabibbo-Kobayashi-Maskawa (CKM) matrix if one takes into account the large theoretical uncertainties associated with the hadronic matrix elements that enter the analysis of ϵ'/ϵ_K [1]. It therefore remains an open question whether the CKM mechanism of *CP* violation can account for the new experimental result on ϵ'/ϵ_K .

A great deal of effort is given to the study of CP violation in the B system, which will provide invaluable information on the pattern of *CP* violation and open up the possibility to look for new physics. In this paper we are concerned with the exclusive decay $\overline{B} \rightarrow K^* l^+ l^-$ which is of special interest because (i) it probes the underlying effective Hamiltonian describing flavor-changing neutral current (FCNC) processes in B decay, (ii) the analysis of CP-violating effects in decays governed by $b \rightarrow s l^+ l^-$ may offer a deeper insight into the mechanism of CP violation since the SM prediction for CP asymmetries is extremely small, typically $\leq 10^{-3}$ [2], and (iii) the process $\overline{B} \rightarrow K^* \mu^+ \mu^-$ is a very promising decay mode since it is likely to be observed in the next round of Bphysics experiments. The most stringent limit has been set by the Collider Detector at Fermilab (CDF) of $\mathcal{B}(B^0)$ $\rightarrow K^{*0}\mu^+\mu^-) < 4.0 \times 10^{-6}$ at the 90% C.L. [3] to be compared with the SM prediction of $(1.9\pm0.7)\times10^{-6}$ [4].

The object of the present work is to explore the possibility of sizable *CP* asymmetries in $\overline{B} \rightarrow K^* l^+ l^-$ decay, whose observation would clearly indicate the presence of physics beyond the SM. We perform a largely model-independent analysis by considering a new-physics scenario with additional *CP* phases and an extended operator basis, taking into account existing experimental data on $\overline{B} \rightarrow X_s \gamma$ and the upper bound on $B^0 \rightarrow K^{*0} \mu^+ \mu^-$.

Our paper is organized as follows. Section II contains the SM operator basis and short-distance matrix element for the process $\overline{B} \rightarrow K^* l^+ l^-$. A formalism for dealing with real $c\overline{c}$ intermediate states such as $J/\psi, \psi'$ entering via the decay chain $\overline{B} \to K^* V_{cc} \to K^* l^+ l^-$ is described. In Sec. III, we briefly discuss possible extensions of the SM including those with an extended operator basis. Section IV is concerned with the parametrization of the hadronic matrix elements and gives the differential decay spectrum as well as the forwardbackward asymmetry. The corresponding CP-violating observables are discussed in Sec. V. Numerical estimates for integrated observables in $\overline{B} \rightarrow K^* \mu^+ \mu^-$ decay in the presence of new physics are given in Sec. VI. Particular attention is paid to the CP-violating partial-rate asymmetry and the *CP*-violating effect related to the angular distribution of μ^{-} in B and \overline{B} decays. Our conclusions are contained in Sec. VII.

II. EFFECTIVE HAMILTONIAN

A. Short-distance contributions

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The effective Hamiltonian for the decay $\overline{B} \rightarrow K^* l^+ l^-$ in the standard model is given by [5,6]

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TABLE I. Numerical values of the Wilson coefficients c_1, \ldots, c_{10} at $\mu = m_b$ within the SM.

<i>c</i> ₁	<i>c</i> ₂	<i>c</i> ₃	<i>C</i> ₄	<i>c</i> ₅	<i>C</i> ₆	$c_7^{\rm eff}$	C 9	<i>c</i> ₁₀
-0.249	+1.108	+0.011	-0.026	+0.007	-0.031	-0.314	+4.216	-4.582

$$H_{\rm eff} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \Biggl\{ \sum_{i=1}^{10} c_i(\mu) \mathcal{O}_i(\mu) + \lambda_u \{c_1(\mu) [\mathcal{O}_1^u(\mu) - \mathcal{O}_1(\mu)] + c_2(\mu) [\mathcal{O}_2^u(\mu) - \mathcal{O}_2(\mu)] \Biggr\} \Biggr\},$$
(2.1)

where we have used the unitarity of the CKM matrix, $\lambda_u \equiv V_{ub} V_{us}^* / V_{tb} V_{ts}^*$, and the operator basis is defined as follows:

$$\begin{split} \mathcal{O}_{1} &= (s_{\alpha}\gamma_{\mu}P_{L}c_{\beta})(c_{\beta}\gamma^{\mu}P_{L}b_{\alpha}), \\ \mathcal{O}_{1}^{u} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}u_{\beta})(\bar{u}_{\beta}\gamma^{\mu}P_{L}b_{\alpha}), \\ \mathcal{O}_{2} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}c_{\alpha})(\bar{c}_{\beta}\gamma^{\mu}P_{L}b_{\beta}), \\ \mathcal{O}_{2}^{u} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}u_{\alpha})(\bar{u}_{\beta}\gamma^{\mu}P_{L}b_{\beta}), \\ \mathcal{O}_{3} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}b_{\alpha}) \sum_{q=u,d,s,c,b} (\bar{q}_{\beta}\gamma^{\mu}P_{L}q_{\beta}), \\ \mathcal{O}_{4} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}b_{\beta}) \sum_{q=u,d,s,c,b} (\bar{q}_{\beta}\gamma^{\mu}P_{R}q_{\beta}), \\ \mathcal{O}_{5} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}b_{\beta}) \sum_{q=u,d,s,c,b} (\bar{q}_{\beta}\gamma^{\mu}P_{R}q_{\beta}), \\ \mathcal{O}_{6} &= (\bar{s}_{\alpha}\gamma_{\mu}P_{L}b_{\beta}) \sum_{q=u,d,s,c,b} (\bar{q}_{\beta}\gamma^{\mu}P_{R}q_{\alpha}), \\ \mathcal{O}_{7} &= \frac{e}{16\pi^{2}} \bar{s}_{\alpha}\sigma_{\mu\nu}(m_{b}P_{R}+m_{s}P_{L})b_{\alpha}F^{\mu\nu}, \\ \mathcal{O}_{8} &= \frac{g_{s}}{16\pi^{2}} \bar{s}_{\alpha}\gamma^{\mu}P_{L}b_{\alpha}\bar{l}\gamma_{\mu}l, \\ \mathcal{O}_{10} &= \frac{e^{2}}{16\pi^{2}} \bar{s}_{\alpha}\gamma^{\mu}P_{L}b_{\alpha}\bar{l}\gamma_{\mu}\gamma_{5}l, \end{split}$$

where α , β are color indices, *a* labels the SU(3) generators, and $P_{L,R} = (1 \mp \gamma_5)/2$. Evolution of the Wilson coefficients $c_i(\mu)$ in Eq. (2.1) from the weak scale $\mu = M_W$ down to the low energy scale $\mu = m_b$ by means of the renormalization group equations (RGE's) then leads to the QCD-corrected matrix element in the next-to-leading logarithmic approximation [5,6]

$$\mathcal{M} = \frac{G_F \alpha}{\sqrt{2} \pi} V_{tb} V_{ts}^* \left\{ \left[(c_9^{\text{eff}} - c_{10}) \langle K^*(k) | \bar{s} \gamma_\mu P_L b | \bar{B}(p) \right] - \frac{2 c_7^{\text{eff}}}{s} \langle K^*(k) | \bar{s} i \sigma_{\mu\nu} q^\nu (m_b P_R + m_s P_L) b | \bar{B}(p) \rangle \right] \times \bar{l} \gamma^\mu P_L l + (c_{10} \rightarrow -c_{10}) \bar{l} \gamma^\mu P_R l \right\}.$$
(2.3)

Here q=p-k, $s=q^2$ is the invariant mass of the lepton pair, and the effective Wilson coefficient c_9^{eff} has the form

$$c_9^{\text{eff}} = c_9 + Y(s),$$
 (2.4)

with

$$Y(s) = g(m_c, s)(3c_1 + c_2 + 3c_3 + c_4 + 3c_5 + c_6) + \lambda_u [g(m_c, s) - g(m_u, s)](3c_1 + c_2) - \frac{1}{2}g(m_s, s)(c_3 + 3c_4) - \frac{1}{2}g(m_b, s)(4c_3 + 4c_4) + 3c_5 + c_6) + \frac{2}{9}(3c_3 + c_4 + 3c_5 + c_6) + \cdots, \quad (2.5)$$

where the ellipsis represents the order α_s correction to the matrix element of the operator \mathcal{O}_9 , which can be regarded as a contribution to the form factors [7], and hence will be omitted in the calculations that follow. Table I summarizes our results for the Wilson coefficients $c_i(m_b)$. Observe that c_7^{eff} , c_9 , and c_{10} are real in the framework of the SM. The function $g(m_i, s)$ in the above formula arises from the one-loop contributions of the four-quark operators $\mathcal{O}_1 - \mathcal{O}_6$, and is given by (at $\mu = m_b$)

$$g(m_i, s) = -\frac{8}{9} \ln(m_i/m_b) + \frac{8}{27} + \frac{4}{9} y_i - \frac{2}{9} (2+y_i) \sqrt{|1-y_i|} \\ \times \left\{ \Theta(1-y_i) \left[\ln \left(\frac{1+\sqrt{1-y_i}}{1-\sqrt{1-y_i}} \right) - i \pi \right] \\ + \Theta(y_i - 1) 2 \arctan \frac{1}{\sqrt{y_i - 1}} \right\},$$
(2.6)

with $y_i = 4m_i^2/s$. This expression reduces to

$$g(0,s) = \frac{8}{27} - \frac{4}{9}\ln(s/m_b^2) + \frac{4}{9}i\pi, \qquad (2.7)$$

(2.2)

in the limit $m_i \rightarrow 0$. A few remarks are in order here.

(i) The Wilson coefficient c_9^{eff} , Eq. (2.4), has absorptive parts for $s > 4m_u^2$ and $s > 4m_c^2$, and thus contains dynamical (unitarity) phases. As will become clear, these are prerequisites, besides a *CP*-violating phase, for observing *CP* asymmetries in partial rates. Since the remaining coefficients c_7^{eff} and c_{10} do not contain any strong phases, the unitarity phases below the $c\bar{c}$ threshold are generated by light quark contributions, whereas for $s > 4m_c^2$ they arise mainly from $c\bar{c}$ intermediate states.

(ii) Within the SM, *CP* violation in $b \rightarrow sl^+l^-$ transition is caused by the ratio of CKM factors appearing in c_9^{eff} , namely $\lambda_u \sim \lambda^2$ ($\lambda \equiv V_{us} \approx 0.22$), which is further reduced by a factor of order $(3c_1+c_2)/c_9 \approx 0.085$. As a result, the effective Hamiltonian for $b \rightarrow sl^+l^-$ essentially involves only one independent CKM factor $V_{lb}V_{ls}^*$, so that *CP* violation in this channel is unobservably small. Thus, the numerical effect of λ_u in Eq. (2.5) is negligible for decays based on the transition $b \rightarrow sl^+l^-$.

B. Resonance contributions

In addition to the short-distance contributions discussed so far, there are possible quark antiquark resonant intermediate states like $\phi, J/\psi, \psi'$ etc. Since the *s*-quark contributions in Eq. (2.5) are suppressed by small Wilson coefficients, and terms proportional to λ_u may be dropped in the case of $b \rightarrow s$ transition, we are left with the J/ψ family.

Following the procedure in Ref. [8], we implement the charmonium resonances utilizing e^+e^- annihilation data. The absorptive part of the one-loop function g, Eq. (2.6), can be related to the experimentally accessible quantity

$$R(s) \equiv \sigma_{\text{tot}}(e^+e^- \rightarrow \text{hadrons}) / \sigma(e^+e^- \rightarrow \mu^+\mu^-)$$
(2.8)

by virtue of the optical theorem. Specifically, it is found that [8]

Im
$$g(m_c,s) = \frac{\pi}{3} R^{J/\psi}(s),$$
 (2.9)

with $R^{J/\psi}(s) \equiv R_{\text{cont}}^{c\bar{c}}(s) + R_{\text{res}}^{J/\psi}(s)$, whereas the dispersive part Re $g(m_c, s)$ may be obtained via a once-subtracted dispersion relation

$$g(m_c,s) = g(m_c,0) + \frac{s}{3} \int_{4M_{\pi}^2}^{\infty} \frac{R^{J/\psi}(s')}{s'(s'-s-i\epsilon)} ds', \quad \epsilon \to +0,$$
(2.10)

with $g(m_c,0) = -8/9 \ln(m_c/m_b) - 4/9$. The continuum contributions, $R_{\text{cont}}^{c\bar{c}}$, can be determined using experimental data from Ref. [9], while the narrow resonances are well described by a relativistic Breit-Wigner distribution [9,10]

$$R_{\rm res}^{J/\psi}(s) = \sum_{V=J/\psi,\psi',\dots} \frac{9s}{\alpha^2} \frac{\mathcal{B}(V \to l^+ l^-) \Gamma_{\rm tot}^V \Gamma_{\rm had}^V}{(s - M_V^2)^2 + M_V^2 \Gamma_{\rm tot}^{V^2}},$$
(2.11)

with the properties of the vector mesons summarized in Ref. [11].

To account for experimental data on direct J/ψ production via the relation

$$\mathcal{B}(\overline{B} \to K^* V_{cc} \to K^* l^+ l^-) = \mathcal{B}(\overline{B} \to K^* V_{cc}) \mathcal{B}(V_{cc} \to l^+ l^-),$$
(2.12)

where $V_{cc} = J/\psi$, ψ' , ..., one usually modifies the Breit-Wigner distribution in Eq. (2.11) by introducing an *ad hoc* factor κ_V [7,12]. This suggests that the factorization ansatz inherent in the approaches that have been advocated to incorporate resonance effects in $b \rightarrow sl^+l^-$ (see, e.g., Ref. [13]) is inadequate for two-body hadronic decays of *B* mesons. Using the set of form factors that will be discussed below, one finds $\kappa_{J/\psi}=1.7$, $\kappa_{\psi'}=2.4$, whereas for the remaining resonances we shall take (as in Ref. [4]) $\kappa=2$. In Fig. 1, we show the real and imaginary parts of $g(m_c, \hat{s})$ based on Eq. (2.10), as a function of $\hat{s}=s/M_B^2$.

III. EFFECTIVE HAMILTONIAN AND NEW PHYSICS

The new-physics contribution to the decay mode $b \rightarrow sl^+l^-$ can manifest itself in two distinct ways: (i) The absolute values and phases of the Wilson coefficients at the electroweak scale are modified. (ii) New operators in addition to the ones defined in Eq. (2.2) arise. We consider them in turn.

A. Wilson coefficients and new physics

The non-standard contributions can be divided into three main groups [14]: first, models with tree-level contributions to the four-quark operators such as supersymmetry (SUSY) with broken R parity or models with Z-mediated flavorchanging neutral currents; second, models with contributions at the one-loop level with new particles running in the loop like SUSY with conserved R parity or models with four quark generations; third, models with no significant effect on the above Wilson coefficients including multi-Higgs-doublet models with natural flavor conservation and left-right symmetric models.

Let us begin with the Wilson coefficients c_3, \ldots, c_6 of the QCD penguin operators. It follows from the RGE analysis that their numerical values at $\mu = m_b$ are essentially determined by the Wilson coefficient c_2 of the four-quark operator \mathcal{O}_2 at the electroweak scale. Thus, in order to have considerable deviations from the SM predictions for the coefficients c_3-c_6 , one needs large new-physics contributions to the short-distance coefficient c_2 , which in turn would affect the theoretical branching ratio for two-body non-leptonic *B* decays. Recent studies [15] suggest, however, that the short-distance coefficients of the SM can account for existing data if one allows departures from the naive factorization prescription.

For the numerical analysis here, we adopt the SM values for the Wilson coefficients c_1, \ldots, c_6 summarized in Table I. Furthermore, it is convenient for later use to parametrize the new-physics contributions to the remaining coefficients



FIG. 1. The predicted \hat{s} dependence of $g(m_c, \hat{s})$ using experimental data on $e^+e^- \rightarrow$ hadrons via a dispersion relation, Eq. (2.10), including the effects of $c\bar{c}$ resonances (solid curve). Note that Im $g(m_c, \hat{s}) \neq 0$ is entirely due to unitarity phases. Also shown is the result of the perturbative calculation according to Eq. (2.6) (dashed curve).

 c_7^{eff} , c_9^{eff} , and c_{10} by the following ratios (defined at the scale $\mu = m_b$):

$$R_7 = c_7^{\text{eff}} / c_7^{\text{eff},\text{SM}} \equiv |R_7| e^{i\phi_7},$$
 (3.1a)

$$R_9 = c_9 / c_9^{\rm SM} \equiv |R_9| e^{i\phi_9}, \qquad (3.1b)$$

$$R_{10} = c_{10} / c_{10}^{\rm SM} \equiv |R_{10}| e^{i\phi_{10}}, \qquad (3.1c)$$

where we recall Eq. (2.4), and

$$c_i = c_i^{\rm SM} + c_i^{\rm new}, \qquad (3.2)$$

with c_i^{SM} shown in Table I. Note that the Wilson coefficient c_{10} in the above does not depend on the renormalization scale, and thus $c_{10} \equiv c_{10}(M_W)$.

B. Extended operator basis

It is conceivable that physics beyond the SM induces new operator structures containing scalar, pseudoscalar, and tensor interactions, in addition to the SM operator basis, Eq. (2.2).

The Wilson coefficients of the scalar operators

$$\mathcal{O}_{9}^{S} = s_{\alpha} P_{L} b_{\alpha} l l,$$

$$\mathcal{O}_{10}^{S} = \overline{s}_{\alpha} P_{L} b_{\alpha} \overline{l} \gamma_{5} l,$$

$$\mathcal{O}_{9}^{S'} = \overline{s}_{\alpha} P_{R} b_{\alpha} \overline{l} l,$$

$$\mathcal{O}_{10}^{S'} = \overline{s}_{\alpha} P_{R} b_{\alpha} \overline{l} \gamma_{5} l \qquad (3.3)$$

involve the lepton mass in most extensions of the SM, and hence will give only small contributions in the case of l=eor μ . Moreover, possible tensor-type operators $\bar{s}\sigma_{\mu\nu}b\bar{l}\sigma^{\mu\nu}l$ and $\bar{s}i\sigma_{\mu\nu}b\bar{l}\sigma_{\alpha\beta}l\epsilon^{\mu\nu\alpha\beta}$ can be safely neglected since their numerical contributions have been found to be small [16].

This leaves an extended operator basis which consists of the SM operators and the opposite-chirality operators [16,17]

$$\mathcal{O}_{7}^{\prime} = \frac{e}{16\pi^{2}} \overline{s}_{\alpha} \sigma_{\mu\nu} (m_{b}P_{L} + m_{s}P_{R}) b_{\alpha}F^{\mu\nu},$$

$$\mathcal{O}_{8}^{\prime} = \frac{g_{s}}{16\pi^{2}} \overline{s}_{\alpha} T^{a}_{\alpha\beta} \sigma_{\mu\nu} (m_{b}P_{L} + m_{s}P_{R}) b_{\beta}G^{a\mu\nu},$$

$$\mathcal{O}_{9}^{\prime} = \frac{e^{2}}{16\pi^{2}} \overline{s}_{\alpha} \gamma^{\mu} P_{R} b_{\alpha} \overline{l} \gamma_{\mu} l,$$

$$\mathcal{O}_{10}^{\prime} = \frac{e^{2}}{16\pi^{2}} \overline{s}_{\alpha} \gamma^{\mu} P_{R} b_{\alpha} \overline{l} \gamma_{\mu} \gamma_{5} l.$$
(3.4)

TABLE II. LCSR predictions for the $B \rightarrow K^*$ form factors with $f(\hat{s}) = f(0)\exp(c_1\hat{s} + c_2\hat{s}^2)$, $\hat{s} = s/M_B^2$ [4]. Recall that A_3 is given in terms of A_1 and A_2 via Eq. (4.3).

	V	A_1	A_2	A_0	T_1	T_2	T_3
f(0)	0.457	0.337	0.282	0.471	0.379	0.379	0.260
c_1	1.482	0.602	1.172	1.505	1.519	0.517	1.129
c_2	1.015	0.258	0.567	0.710	1.030	0.426	1.128

IV. DECAY DISTRIBUTION FOR $\overline{B} \rightarrow K^* l^+ l^-$

A. Form factors

The hadronic matrix elements appearing in Eq. (2.3) can be expressed in terms of *s*-dependent form factors (recall $s \equiv q^2$), namely [18]

$$\begin{split} \langle K^*(k) | \overline{s} \gamma_{\mu} P_{L,R} b | \overline{B}(p) \rangle \\ &= i \epsilon_{\mu\nu\alpha\beta} \epsilon^{\nu*} p^{\alpha} q^{\beta} \frac{V(s)}{M_B + M_{K*}} \\ &= \frac{1}{2} \Biggl\{ \epsilon^*_{\mu} (M_B + M_{K*}) A_1(s) - (\epsilon^* \cdot q) (2p - q)_{\mu} \\ &\times \frac{A_2(s)}{M_B + M_{K*}} - \frac{2M_{K*}}{s} (\epsilon^* \cdot q) [A_3(s) - A_0(s)] q_{\mu} \Biggr\}, \end{split}$$

$$(4.1)$$

with the convention $\epsilon_{0123} = +1$, q = p - k, and

$$\langle K^*(k) | \overline{s} i \sigma_{\mu\nu} q^{\nu} P_{R,L} b | \overline{B}(p) \rangle$$

$$= -i \epsilon_{\mu\nu\alpha\beta} \epsilon^{\nu*} p^{\alpha} q^{\beta} T_1(s) \pm \frac{1}{2} \Biggl\{ \left[\epsilon^*_{\mu} (M_B^2 - M_{K^*}^2) - (\epsilon^* \cdot q) (2p - q)_{\mu} \right] T_2(s) + (\epsilon^* \cdot q) \Biggr\}$$

$$\times \Biggl[q_{\mu} - \frac{s}{M_B^2 - M_{K^*}^2} (2p - q)_{\mu} \Biggr] T_3(s) \Biggr\},$$

$$(4.2)$$

where $T_1(0) = T_2(0)$, and ϵ^{μ} is the K^* polarization vector. The form factor A_3 can be written in terms of A_1 and A_2 , i.e.

$$A_3(s) = \frac{M_B + M_{K^*}}{2M_{K^*}} A_1(s) - \frac{M_B - M_{K^*}}{2M_{K^*}} A_2(s), \quad (4.3)$$

with the relation $A_3(0) = A_0(0)$. The terms proportional to q_{μ} in Eqs. (4.1) and (4.2) do not contribute to the differential decay rate when the final leptons are massless. For the results presented below, we adopt the $B \rightarrow K^*$ form factors of Ref. [4] which have been obtained using light cone sum rule (LCSR) results, and are displayed in Table II. Throughout our discussion, we assume the above form factors to be real, in the absence of final-state interactions.

B. Differential decay spectrum

Squaring the matrix element (2.3), summing over spins, and introducing the shorthand notation

$$\lambda(a,b,c) = a^2 + b^2 + c^2 - 2(ab + bc + ac), \qquad (4.4)$$

$$\hat{M}_i = M_i / M_B, \quad \hat{m}_i = m_i / M_B, \quad \hat{s} = s / M_B^2,$$
(4.5)

$$X = \frac{1}{2} \lambda^{1/2} (1, \hat{s}, \hat{M}_{K^*}^2), \qquad (4.6)$$

the spectrum of $\overline{B} \rightarrow K^* l^+ l^-$ decay with respect to \hat{s} and θ_l , the angle between l^- and the outgoing hadron in the dilepton center-of-mass system, is [19]

$$\frac{d\Gamma(\bar{B} \to K^* l^+ l^-)}{d\hat{s} \, d \cos \theta_l} = \frac{G_F^2 \alpha^2 M_B^5}{2^9 \pi^5} |V_{lb} V_{ls}^*|^2 X \sqrt{1 - \frac{4\hat{m}_l^2}{\hat{s}}} \times [A(\hat{s}) + B(\hat{s}) \cos \theta_l + C(\hat{s}) \cos^2 \theta_l].$$
(4.7)

The quantities A, B, C are defined as follows:

$$A(\hat{s}) = \frac{2X^2}{\hat{M}_{K^*}^2} \left[\hat{s} \hat{M}_{K^*}^2 f_1(\hat{s}) + \frac{1}{4} \left(1 + \frac{2\hat{s} \hat{M}_{K^*}^2}{X^2} \right) \right. \\ \left. \times f_2(\hat{s}) + X^2 f_3(\hat{s}) + f_4(\hat{s}) \right] + 2\hat{m}_l^2 I(\hat{s}), \quad (4.8)$$

$$B(\hat{s}) = 8X \sqrt{1 - \frac{4\hat{m}_l^2}{\hat{s}}} \operatorname{Re} \{ c_{10}^* [c_9^{\text{eff}} \hat{s} A_x A_y - c_7^{\text{eff}} (A_x B_y + A_y B_x)] \}, \qquad (4.9)$$

$$C(\hat{s}) = \frac{2X^2}{\hat{M}_{K^*}^2} \left(1 - \frac{4\hat{m}_l^2}{\hat{s}}\right) \left[\hat{s}\hat{M}_{K^*}^2 f_1(\hat{s}) - \frac{1}{4}f_2(\hat{s}) - X^2 f_3(\hat{s}) - f_4(\hat{s})\right], \qquad (4.10)$$

with the auxiliary functions

$$I(\hat{s}) = 4X^2 f_1(\hat{s}) + f_2(\hat{s}) + f_5(\hat{s}), \qquad (4.11)$$

$$f_{1}(\hat{s}) = (|c_{9}^{\text{eff}}|^{2} + |c_{10}|^{2})A_{x}^{2} + \frac{4|c_{7}^{\text{eff}}|^{2}}{\hat{s}^{2}}B_{x}^{2} - \frac{4\operatorname{Re}(c_{7}^{\text{eff}}c_{9}^{\text{eff}*})}{\hat{s}}A_{x}B_{x}, \qquad (4.12)$$

$$f_2(\hat{s}) = f_1(\hat{s})_{x \to y}, \quad f_3(\hat{s}) = f_1(\hat{s})_{x \to z},$$
 (4.13)

$$f_{4}(\hat{s}) = \frac{1}{2} (1 - \hat{s} - \hat{M}_{K^{*}}^{2}) \left[(|c_{9}^{\text{eff}}|^{2} + |c_{10}|^{2}) A_{y} A_{z} + \frac{4|c_{7}^{\text{eff}}|^{2}}{\hat{s}^{2}} B_{y} B_{z} - \frac{2 \operatorname{Re} (c_{7}^{\text{eff}} c_{9}^{\text{eff}})}{\hat{s}} (A_{y} B_{z} + A_{z} B_{y}) \right],$$

$$(4.14)$$

$$f_{5}(\hat{s}) = |c_{10}|^{2} \left[-8X^{2}A_{x}^{2} - 3A_{y}^{2} + \frac{X^{2}}{\hat{M}_{K^{*}}^{2}} \{ [2(1 + \hat{M}_{K^{*}}^{2}) - \hat{s}]A_{z} + 2A_{y} \}A_{z} + \frac{4X^{2}}{\hat{s}\hat{M}_{K^{*}}} \{ \hat{M}_{K^{*}}(A_{3} - A_{0}) - [(1 - \hat{M}_{K^{*}}^{2})A_{z} + A_{y}] \}(A_{3} - A_{0}) \right].$$
(4.15)

In these equations,

$$A_{x} = \frac{V(\hat{s})}{1 + \hat{M}_{K^{*}}}, \quad A_{y} = (1 + \hat{M}_{K^{*}})A_{1}(\hat{s}), \quad A_{z} = -\frac{A_{2}(\hat{s})}{1 + \hat{M}_{K^{*}}},$$
(4.16)

ŀ

$$B_{x} = -T_{1}(\hat{s})(\hat{m}_{b} + \hat{m}_{s}),$$

$$B_{y} = -(1 - \hat{M}_{K*}^{2})T_{2}(\hat{s})(\hat{m}_{b} - \hat{m}_{s}),$$
(4.17)

$$B_{z} = \left[T_{2}(\hat{s}) + \frac{\hat{s}}{1 - \hat{M}_{K^{*}}^{2}} T_{3}(\hat{s}) \right] (\hat{m}_{b} - \hat{m}_{s}), \qquad (4.18)$$

with the form factors listed in Table II.

A further observable of interest is the forward-backward (FB) asymmetry of l^- , defined as

$$A_{\rm FB}(\hat{s}) = \frac{\int_0^1 d\cos\theta_l \frac{d\Gamma}{d\hat{s}\,d\cos\theta_l} - \int_{-1}^0 d\cos\theta_l \frac{d\Gamma}{d\hat{s}\,d\cos\theta_l}}{\int_0^1 d\cos\theta_l \frac{d\Gamma}{d\hat{s}\,d\cos\theta_l} + \int_{-1}^0 d\cos\theta_l \frac{d\Gamma}{d\hat{s}\,d\cos\theta_l}},$$
(4.19)

and we obtain

$$A_{\rm FB}(\hat{s}) = 12X \sqrt{1 - \frac{4\hat{m}_l^2}{\hat{s}}} \frac{\operatorname{Re}\left\{c_{10}^* \left[c_9^{\rm eff} \hat{s} A_x A_y - c_7^{\rm eff} (A_x B_y + A_y B_x)\right]\right\}}{[3A(\hat{s}) + C(\hat{s})]}.$$
(4.20)

Note that the Wilson coefficients in the above expressions are defined through Eq. (3.2). The decay distributions in the presence of new physics with additional operators are discussed further in Sec. VI. We now proceed to a discussion of *CP*-violating observables in the process $\overline{B} \rightarrow K^* l^+ l^-$.

V. CP-VIOLATING OBSERVABLES

Suppose the decay amplitude for $\overline{B} {\rightarrow} F$ has the general form

$$\mathcal{A}(\overline{B} \to F) = e^{i\phi_1}A_1e^{i\delta_1} + e^{i\phi_2}A_2e^{i\delta_2}, \qquad (5.1)$$

where $A_{1,2}$ are real matrix elements, δ_i and ϕ_i are the strong phases (*CP*-conserving) and weak phases (*CP*-violating) respectively. Using *CPT* invariance, which requires that the total decay rate for particle and antiparticle be equal, the decay amplitude for the conjugate process takes the form

$$\overline{\mathcal{A}}(B \to \overline{F}) = e^{-i\phi_1} A_1 e^{i\delta_1} + e^{-i\phi_2} A_2 e^{i\delta_2}, \qquad (5.2)$$

giving rise to the *CP* asymmetry

$$A_{CP} \equiv \frac{|\mathcal{A}|^2 - |\mathcal{\bar{A}}|^2}{|\mathcal{A}|^2 + |\mathcal{\bar{A}}|^2} = \frac{-2r\sin\phi\sin\delta}{1 + 2r\cos\phi\cos\delta + r^2}, \quad (5.3)$$

with $r = A_2/A_1$, $\phi = \phi_1 - \phi_2$, and $\delta = \delta_1 - \delta_2$. Notice that in the limit $r \ll 1$ the asymmetry is approximately $A_{CP} \approx$ $-2r \sin \phi \sin \delta$. Inspection of Eq. (5.3) reveals that a nonzero partial-rate asymmetry requires *CP* violation ($\phi \neq 0$) and the presence of dynamical phases ($\delta \neq 0$), the latter being provided by the one-loop function $g(m_i, s)$ present in the Wilson coefficient c_9^{eff} [Eq. (2.4)].

To determine the impact of the strong phases on the *CP* asymmetries, it is useful to separate those contributions to the decay amplitude that generate absorptive parts. To this end, we choose $\delta_1 = \phi_2 = 0$ in Eq. (5.1) and require that $A_2 e^{i\delta_2}$ vanish when $Y(s) \rightarrow 0$. Moreover, as discussed in Sec. III, the corrections to the short-distance coefficients of the SM multiplying the absorptive parts of the decay amplitude are not expected to be large in many extensions of the SM or required by current data on two-body hadronic *B* decays [15]. Consequently, the strong phases in $b \rightarrow sl^+l^-$ decay are essentially unaffected by possible new interactions transcending the SM, so that the non-standard effects on *CP* asymmetries can be described by the two parameters *r* and ϕ .



FIG. 2. The quantity $R_Y \equiv -\text{Im}[Y(\hat{s})]/|Y(\hat{s})|$ in $b \rightarrow sl^+l^-$ decay vs \hat{s} for the low- \hat{s} (a) and high- \hat{s} (b) regions, with $\hat{s} = s/M_B^2$.

In order to get a quantitative idea of the magnitude of unitarity phases in $b \rightarrow sl^+l^-$ transition, we plot in Fig. 2 the parameter $R_Y \equiv -\text{Im}[Y(\hat{s})]/|Y(\hat{s})|$ for two different regions of the dilepton invariant mass, namely

$$4\hat{m}_{l}^{2} \leq \hat{s} \leq (\hat{M}_{J/\psi} - \epsilon_{\rm cut})^{2}, \quad (\hat{M}_{\psi'} + \epsilon_{\rm cut}')^{2} \leq \hat{s} \leq (1 - \hat{M}_{K^{*}})^{2},$$
(5.4)

which we refer to as the low- \hat{s} and high- \hat{s} region respectively.¹ As we alluded to earlier, in the low- \hat{s} region the strong phase is suppressed by small Wilson coefficients of the QCD penguin operators, whereas in the high- \hat{s} region it can be large.

Figure 3 shows the dependence of A_{CP} on r for different choices of the weak phase $\sin \phi$ in the low- \hat{s} and high- \hat{s} regions. Observe that the asymmetry is maximized when r = 1 (i.e. the two interfering amplitudes are of comparable size) but is strongly suppressed if either $r \ge 1$ or $r \le 1$.

Using the two-dimensional decay distribution $d\Gamma/d\hat{s} d \cos \theta_l$ derived in the preceding section, we may define the following average *CP* asymmetries:

$$\langle A_{CP}^{D} \rangle = \frac{\int_{\hat{s}_{0}}^{\hat{s}_{1}} d\hat{s} \int_{D} d\cos \theta_{l} \frac{d\Gamma_{\text{sum}}}{d\hat{s} \, d\cos \theta_{l}}}{\int_{\hat{s}_{0}}^{\hat{s}_{1}} d\hat{s} \int_{-1}^{+1} d\cos \theta_{l} \frac{d\Gamma_{\text{sum}}}{d\hat{s} \, d\cos \theta_{l}}},$$

$$\langle A_{CP}^{S} \rangle = \frac{\int_{\hat{s}_{0}}^{\hat{s}_{1}} d\hat{s} \int_{S} d\cos \theta_{l} \frac{d\Gamma_{\text{diff}}}{d\hat{s} \, d\cos \theta_{l}}}{\int_{\hat{s}_{0}}^{\hat{s}_{1}} d\hat{s} \int_{-1}^{+1} d\cos \theta_{l} \frac{d\Gamma_{\text{sum}}}{d\hat{s} \, d\cos \theta_{l}}},$$
(5.5)

where

$$\int_{D,S} \equiv \int_{0}^{1} \mp \int_{-1}^{0},$$
(5.6)

$$\Gamma_{\rm sum} = \Gamma(\bar{B} \to K^* l^+ l^-) + \bar{\Gamma}(B \to \bar{K}^* l^+ l^-), \qquad (5.7)$$

$$\Gamma_{\text{diff}} = \Gamma(\bar{B} \to K^* l^+ l^-) - \bar{\Gamma}(B \to \bar{K}^* l^+ l^-).$$
(5.8)

Notice that the asymmetry A_{CP}^{D} is a *CP*-violating effect in the angular distribution of l^{-} (or equivalently the Dalitz-plot distribution) in *B* and \overline{B} decays while A_{CP}^{S} represents the asymmetry in the partial widths of these decays. The special feature of the former is that it can be determined even for an untagged equal mixture of *B* and \overline{B} events, i.e. without flavor identification. We emphasize that both asymmetries are odd

¹We use $\epsilon_{cut} = 0.2$ GeV/ M_B and $\epsilon'_{cut} = 0.1$ GeV/ M_B in the case of $\mu^+ \mu^-$ in the final state.



FIG. 3. *CP*-violating asymmetry A_{CP} , Eq. (5.3), as a function of *r* for (a) the low- \hat{s} region with sin $\delta = -0.1$ and (b) the high- \hat{s} region with sin $\delta = -0.5$.

under *CP* but even under "naive" T, and thus vanish in the limit that there are no strong phases.²

VI. NUMERICAL ANALYSIS

In the remainder of this paper, we wish to focus on the $\overline{B} \rightarrow K^* \mu^+ \mu^-$ mode. We start this section with a brief discussion of the experimental constraints that we shall be using in our subsequent calculations.

A. Experimental constraints

Constraints on the parameters R_i , Eqs. (3.1), are provided by the following experimental results: (i) The CDF Collaboration has recently obtained the upper bound

$$\mathcal{B}(B^0 \to K^{*0} \mu^+ \mu^-) < 4.0 \times 10^{-6} \tag{6.1}$$

at 90% C.L. [3]. (ii) The measurement of the inclusive branching ratio $\mathcal{B}(\overline{B} \rightarrow X_s \gamma)$ yields [21]

$$2.0 \times 10^{-4} < \mathcal{B}(\bar{B} \to X_s \gamma) < 4.5 \times 10^{-4} \quad (95\% \quad \text{C.L.}).$$
(6.2)

Employing the leading-order result for $\mathcal{B}(\overline{B} \to X_s \gamma)$ (see, e.g., Ref. [22]), useful bounds can be placed on the absolute value of R_7 : namely,

$$0.881 < |R_7| < 1.321. \tag{6.3}$$

B. Integrated observables in $\overline{B} \rightarrow K^* \mu^+ \mu^-$

Below we present our predictions (i) for the non-resonant invariant mass spectrum of the muon pair and (ii) for the low- \hat{s} and high- \hat{s} region, as defined in Eq. (5.4). The resonance effects are taken into account by employing Eqs. (2.9) and (2.10).

(i) The branching ratio resulting from the SM operator basis given in Eq. (2.2) can be conveniently written as

$$\mathcal{B} = [a_0 + a_1 |R_{10}|^2 + a_2 |R_7|^2 + a_3 |R_9|^2 + a_4 \text{Im } R_7 + a_5 \text{Im } R_9 + a_6 \text{Re} (R_7 R_9^*) + a_7 \text{Re} R_7 + a_8 \text{Re} R_9] \times 10^{-7}.$$
(6.4)

Here we have introduced the coefficients $a_i \equiv \alpha_i + \beta_i$ which allow us to incorporate the effects of new operators \mathcal{O}'_i [Eq. (3.4)] into our computation of the branching ratio by simply replacing

$$(\alpha_i + \beta_i) f(R_i) \rightarrow \alpha_i f(R_i + R'_i) + \beta_i f(R_i - R'_i), \quad (6.5)$$

²For a review of *CP*-odd observables, we refer the reader to Ref. [20].

TABLE III. Numerical estimate of the coefficients $a_i \equiv \alpha_i + \beta_i$ entering the expressions for the integrated branching ratio \mathcal{B} and the average FB asymmetry $\langle A_{\text{FB}} \rangle$ in $\overline{B} \rightarrow K^* \mu^+ \mu^-$ decay, as described in the text.

Coefficients	$\mathcal{B}^{\mathrm{nr}}$	${\cal B}^{ m low}$	${\cal B}^{ m high}$
(α_0,β_0)	(0.03,0.12)	(0.02,0.08)	(0.02,0.13)
(α_1,β_1)	(1.89,8.64)	(0.62,3.84)	(0.36,1.76)
(α_2,β_2)	(1.05,1.06)	(0.95,0.90)	(0.01,0.04)
(α_3, β_3)	(1.62,7.34)	(0.54,3.26)	(0.30,1.50)
(α_4, β_4)	-(0.07, 0.18)	-(0.01, 0.02)	-(0.02,0.10)
(α_5, β_5)	(0.23,0.93)	(0.01,0.07)	(0.11,0.60)
(α_6, β_6)	-(1.52,3.12)	-(0.90, 1.59)	-(0.13,0.52)
(α_7, β_7)	-(0.15,0.30)	-(0.12,0.21)	(0.02,0.07)
(α_8,β_8)	(0.33,1.41)	(0.18,0.91)	-(0.08, 0.44)
Coefficients	$\langle A_{\rm FB} \rangle^{\rm nr}$	$\langle A_{\rm FB} \rangle^{\rm low}$	$\langle A_{\rm FB} \rangle^{\rm high}$
(α_0,β_0)	(0.33,0.33)	(0.16,0.16)	-(0.05,0.05)
(α_1, β_1)	-(1.27, 1.42)	-(0.69, 0.77)	-(0.16,0.18)
(α_2,β_2)	(3.29,3.29)	(0.94,0.94)	(0.82,0.82)
(α_3,β_3)	(0.25,0.25)	(0.01,0.01)	(0.16,0.16)

where we have defined the quantities R'_i analogous to those in Eqs. (3.1). Our numerical results for the coefficients α_i and β_i are listed in Table III.

(ii) Similarly, we parametrize the average FB asymmetry as

$$\langle A_{\rm FB} \rangle = - \left[\operatorname{Re} R_{10}^*(a_0 + a_1 R_7 + a_2 R_9) + a_3 \operatorname{Im} R_{10} \right] \left(\frac{10^{-7}}{\mathcal{B}} \right),$$

(6.6)

with the values of a_i tabulated in Table III. As before, the average FB asymmetry in the presence of chirality-flipped operators can be obtained from Eq. (6.6) by the following replacements:

$$(\alpha_{i} + \beta_{i})f(R_{7}, R_{9}, R_{10}) \rightarrow \alpha_{i}f(R_{7} - R_{7}', R_{9} - R_{9}', R_{10} + R_{10}') + \beta_{i}f(R_{7} + R_{7}', R_{9} + R_{9}', R_{10} - R_{10}').$$
(6.7)

(iii) Our results for the *CP* asymmetries are as follows:

$$\langle A_{CP}^{S} \rangle^{\text{low}} = (a_0 \text{Im } R_7 + a_1 \text{Im } R_9) \left(\frac{10^{-9}}{\hat{\mathcal{B}}^{\text{low}}} \right), \quad (6.8)$$

$$\langle A_{CP}^{S} \rangle^{\text{high}} = (a_0 \text{Im } R_7 + a_1 \text{Im } R_9) \left(\frac{10^{-8}}{\hat{\mathcal{B}}^{\text{high}}} \right),$$

(6.9)

and

$$\langle A_{CP}^D \rangle^{\text{low}} = -a_0 \sin \phi_{10} |R_{10}| \left(\frac{10^{-9}}{\hat{\mathcal{B}}^{\text{low}}} \right),$$
 (6.10)

$$\langle A_{CP}^{D} \rangle^{\text{high}} = -a_0 \sin \phi_{10} |R_{10}| \left(\frac{10^{-8}}{\hat{\mathcal{B}}^{\text{high}}} \right),$$

(6.11)

 $\hat{\mathcal{B}} = (\mathcal{B} + \overline{\mathcal{B}})/2$ being the *CP*-averaged branching ratio [see Eqs. (5.1) and (5.2)]. It is obvious that $\langle A_{CP}^S \rangle$ is sensitive to the *CP*-violating phases ϕ_7 and ϕ_9 , while $\langle A_{CP}^S \rangle$ also probes the phase ϕ_{10} . The predictions for the coefficients a_i are reported in Table IV.

C. Numerical results and predictions

We first analyze the *CP* asymmetries in the context of the SM operator basis, Eq. (2.2). To determine the implications of new physics for the CP asymmetries, we adopt the following procedure. The absolute value of R_7 is chosen such that it satisfies the constraints implied by the measured b $\rightarrow s \gamma$ rate, Eq. (6.3), while the remaining parameters R_9 and R_{10} are required to be consistent with the current experimental upper limit on the non-resonant branching ratio $\mathcal{B}(B^0)$ $\rightarrow K^{*0}\mu^+\mu^-$) given in Eq. (6.1). To gain predictivity we take $|R_7|$ to be unity, so that we end up with a set of free parameters comprising ϕ_7 , ϕ_9 , and $|R_9|$. In fact, imposing the requirement that the non-resonant branching ratio coincide with the experimental upper bound or, alternatively, with the SM prediction, the quantity R_{10} is computed for any given set of the parameters $\phi_7, \phi_9, |R_9|$.³ This enables us to consider a scenario where the predicted $b \rightarrow s \gamma$ fraction coincides with the SM expectation while the non-resonant branching ratio of $\overline{B} \rightarrow K^* \mu^+ \mu^-$ may well be accessible in the next round of *B* experiments.

TABLE IV. Values of the coefficients $a_i \equiv \alpha_i + \beta_i$ for *CP*-violating observables in $\overline{B} \rightarrow K^* \mu^+ \mu^-$.

Coefficients	$\langle A_{CP}^S \rangle^{\text{low}}$	$\langle A^{S}_{CP} \rangle^{ m high}$	$\langle A^{D}_{CP} \rangle^{\text{low}}$	$\langle A^{D}_{CP} \rangle^{\text{high}}$
$(lpha_0,eta_0) \ (lpha_1,eta_1)$	-(0.99,1.76) (1.20,7.22)	-(0.25,1.03) (1.16,5.98)	(1.05,1.05)	(1.62,1.62)

³As far as new CP-violating phases are concerned, we note that in the context of a specific model one has also to take into account the severe constraints on the electric dipole moments of electron and neutron.



FIG. 4. The parameters *r* and sin ϕ [see Eq. (5.3)] vs ϕ_7 and ϕ_9 in the low dimuon invariant mass region. For simplicity, we have taken $|R_7| = |R_9| = 1$ while R_{10} is chosen to coincide with the experimental upper limit on the non-resonant branching ratio, i.e. $\mathcal{B}^{nr} = 4.0 \times 10^{-6}$.

As mentioned above, the value of the strong phase entering the amplitude of the $\overline{B} \rightarrow K^* \mu^+ \mu^-$ process depends on the Wilson coefficients c_1, \ldots, c_6 which we have assumed to be unaffected by new-physics contributions. Hence, the value of the strong phase sin δ is fixed, and estimated to be -0.07 and -0.51 in the low- \hat{s} and high- \hat{s} domain respectively.

Last, using the parametrization for the amplitude given in Eq. (5.1), and employing the integrated expressions above, we determine r and $\sin \phi$ numerically as a function of ϕ_7, ϕ_9 , and $|R_9|$.

1. CP asymmetry in the partial widths

Figures 4 and 5 show the parameters r, sin ϕ , and A_{CP}^S in the low dimuon invariant mass region as a function of the phases ϕ_7 and ϕ_9 , taking $|R_9|=1$ and requiring that $\mathcal{B}^{nr} = 4.0 \times 10^{-6.4}$ It may be noted that the predictions for the average *CP* asymmetry depend very little on the phase ϕ_7 and its absolute value is no greater than 1%, regardless of the size of the *CP*-violating phases. If we allow for devia-

⁴The impact of new phases on the partial-rate asymmetry in $\overline{B} \rightarrow K^* l^+ l^-$ decay for the case of $|R_i| = 1$ has also been studied in Ref. [23].



FIG. 5. *CP*-violating average asymmetry $\langle A_{CP}^{S} \rangle$ in the low dimuon invariant mass region vs ϕ_7 and ϕ_9 , with the parameters $|R_i|$ (*i*=7,9,10) as in Fig. 4.



FIG. 6. The parameters *r* and sin ϕ appearing in A_{CP} , Eq. (5.3), as a function of ϕ_7 and ϕ_9 for the high dimuon invariant mass region, with $|R_7| = |R_9| = 1$. The ratio R_{10} is chosen to be consistent with the experimental upper limit on the non-resonant branching ratio given in Eq. (6.1).

tions from $|R_9|=1$, the magnitude of the *CP* asymmetry does not change significantly. We may therefore conclude that the partial-rate asymmetry, indicative of *CP* violation, in the low- \hat{s} region is too small to be observable (assuming that the indispensable strong phase does not receive any substantial non-standard contributions).

We now turn our attention to the high- \hat{s} region. It is clear from the above discussion that physics beyond the SM can give rise to sizable *CP* asymmetries above the ψ' resonance where we expect to have an appreciable strong phase but also a lower branching ratio. In fact, in the high- \hat{s} region we find the approximate relation $r^{\text{high}} \simeq 2r^{\text{low}}$, whereas the numerical value of the weak phase $\sin \phi$ is of the same order of magnitude in both the low and the high dimuon invariant mass region. This can be seen in Figs. 6 and 7, where we show our results for r, sin ϕ , and A_{CP}^S . For certain values of the phases ϕ_7 and ϕ_9 , the *CP* asymmetry can be of order $\pm 10\%$. As far as theoretical uncertainties are concerned, we merely remark that the numerical estimates for average CP asymmetries are largely independent of the parametrizations of form factors, so that the theoretical uncertainty associated with real $c\bar{c}$ intermediate states discussed in Sec. II gives by far the largest uncertainty in the predicted CP asymmetry.

Next we consider the case where we allow for higher values of $|R_9|$. We begin by noting that for $|R_9| > 1.75$ some



FIG. 7. *CP*-violating average asymmetry $\langle A_{CP}^{S} \rangle$ in the high invariant mass region of the muon pair, as a function of ϕ_7 and ϕ_9 . The parameters $|R_i|$ (*i*=7,9,10) have been chosen as in Fig. 6.



FIG. 8. The quantity $R_{CP} \equiv \langle A_{CP}^D \rangle^{\text{high}/\text{sin}} \phi_{10}$ [cf. Eq. (6.11)] as a function of ϕ_7 and ϕ_9 with $|R_7| = 1$. (a) $|R_{10}| = 1.9$ and $|R_9|$ is chosen such that it coincides with the upper limit on $\mathcal{B}(B^0 \to K^{*0} \mu^+ \mu^-)$. (b) $|R_{10}| = 1.2$ and the SM branching fraction of 1.8×10^{-6} has been adopted.

part of the (ϕ_7, ϕ_9) parameter space is already excluded by the experimental upper bound on the non-resonant branching ratio, Eq. (6.1). Exploiting the fact that A_{CP}^S depends only weakly on the phase ϕ_7 , we may take $\phi_7=4.8$ (see Fig. 6). Then, setting $|R_9|=1$ we find a magnitude of the average *CP* asymmetry varying from -0.05 to 0.12, whereas for $|R_9|=1.75$ we predict the range $-0.15 \le \langle A_{CP}^S \rangle \le 0.20$. In this latter case, numerical values of *r* between 0.26 and 0.33 are estimated while sin ϕ can be maximal. In particular, for $|R_9|=1.75$ the weak phase sin ϕ is nearly unity if we demand the non-resonant branching fraction to be 4.0×10^{-6} .

Let us now discuss a scenario in which we abandon the assumption of having a non-resonant branching ratio of 4.0 $\times 10^{-6}$. We focus here on the high- \hat{s} region as we do not expect any significant deviation from our results obtained in the region below the charmonium resonances. We first consider the case where the branching ratio is still fixed to the SM value of $\mathcal{B}^{nr}=1.8\times 10^{-6}$. Further, we take $\phi_7=4.8$ and $|R_9|=0.9$. (Note that larger values of $|R_9|$ are not consistent with the SM branching ratio.) As a result, the weak phase sin ϕ exhibits almost the same ϕ_9 -dependent behavior as in the previous case with maximum branching ratio presently allowed by experiment. In addition, a smaller value for \mathcal{B}^{nr} leads to a wider range of r, namely $0.47 \leq r \leq 0.70$. As for *CP* violation, an average asymmetry of anything between -0.20 and 0.30 is predicted.

It is also interesting to analyze the case in which the branching ratio is not fixed to any particular value but is still compatible with the experimental results described above. Remembering that R_{10} contributes only to the branching ratio, we set $R_{10}=0$ in order to get the highest possible value for the *CP* asymmetry (see Sec. V). Consequently, the interference of the terms R_7 and R_9 now plays an essential role as the branching ratio diminishes for certain values of ϕ_9 , and so *r* can be unity (which corresponds to the maximal size of the *CP* asymmetry). In this case, the *CP* asymmetry $\langle A_{CP}^S \rangle$ varies considerably and can take on any value between -0.5 and 0.4 for \mathcal{B}^{nr} ranging from 5.0×10^{-7} to 1.6×10^{-6} .

2. CP asymmetry in the angular distribution of μ^-

A similar analysis has been carried out for the asymmetry A_{CP}^{D} , Eq. (5.5), which is of considerable interest since it probes the phase of R_{10} . Here again the *CP* asymmetry in the low- \hat{s} domain is fairly small, typically a few percent, and we therefore concentrate on the high- \hat{s} region where *CP*-violating effects are not suppressed by small unitarity phases. Results for the ratio $R_{CP} \equiv \langle A_{CP}^D \rangle^{\text{high}/\sin\phi_{10}}$ as a function of ϕ_7 and ϕ_9 are displayed in Fig. 8. Demanding that $|R_9|$ reproduce the experimental upper bound on the non-resonant branching ratio leads to the predictions shown in Fig. 8(a). Moreover, the quantity R_{10} obeys the constraint $|R_{10}| \leq 1.9$. On the other hand, assuming the SM prediction of $\mathcal{B}^{\text{nr}} = 1.8 \times 10^{-6}$, we obtain the results shown in Fig. 8(b).

In this case, present experimental data on $\mathcal{B}(B^0 \to K^{*0}\mu^+\mu^-)$ lead to the upper bound $|R_{10}| \le 1.2$. In both cases, the *CP* asymmetry in the angular distribution of μ^- in *B* and \overline{B} decays turns out to be about -10%.⁵

Finally, taking $|R_9|=0$ we find that the *CP* asymmetry can be as large as -25% for a rather low $|R_{10}|$, which corresponds to a non-resonant branching fraction of $O(10^{-7})$. If we consider instead a branching ratio of $\mathcal{B}^{nr} \approx 10^{-6}$, the asymmetry can reach to values of -15%.

3. Comment on CP violation and additional operators

As seen in the preceding, within the framework of the SM operator basis it is possible to account for the maximum values of the *CP*-violating asymmetries. Hence our quantitative results for *CP* violation in the decay $\overline{B} \rightarrow K^* \mu^+ \mu^-$ are not affected by the chirality-flipped operators [Eq. (3.4)] once existing experimental constraints are taken into account. In other words, the observation of an appreciable *CP* asymmetry alone does not provide a test of the chirality structure of operators that enter the effective Hamiltonian, and thus is not sufficient to disentangle different new-physics scenarios.

VII. SUMMARY AND CONCLUSIONS

We have performed a largely model-independent analysis of the exclusive decay $\overline{B} \rightarrow K^* l^+ l^-$ in the presence of physics transcending the SM. In particular, we have investigated the implications of new *CP*-violating phases for the decay $\overline{B} \rightarrow K^* \mu^+ \mu^-$, and derived analytic expressions for the branching ratio, the FB asymmetry, and certain CP-violating observables. The formalism presented is applicable to any effective Hamiltonian containing the SM operator basis as well as operators with a different chirality structure. We have studied in some detail the CP asymmetries in the partial rates and the angular distribution of μ^- in $\overline{B} \rightarrow K^* \mu^+ \mu^$ and $B \rightarrow \overline{K}^* \mu^+ \mu^-$ decays, which require the simultaneous presence of weak and strong phases. Adopting the SM operator basis and assuming that new-physics contributions to two-body non-leptonic B decays are unlikely to be significant, the CP-violating effects in the $2m_{\mu} \leq M_{\mu^+\mu^-} < M_{J/\psi}$ domain are estimated to be small (up to at most a few percent). Ultimately, this result is a consequence of the smallness of the dynamical phase associated with the absorptive part of the penguin diagram; thus, the presence of large nonstandard CP-violating phases does not necessarily imply sizable CP-violating effects in the lower part of the decay spectrum. Even so, studies of *CP*-violating effects in the low- \hat{s} region will provide a crucial test of the SM, since any indication of CP violation would represent new physics. On the other hand, in the high dimuon invariant mass region $M_{\mu'}$ $< M_{\mu^+\mu^-} \leq (M_B - M_{K^*})$ appreciable *CP*-violating effects can show up, which are consistent with current experimental data on the $b \rightarrow s \gamma$ branching fraction and the upper bound on $\mathcal{B}(B^0 \to K^{*0} \mu^+ \mu^-)$. We find that *CP* asymmetries up to 30% can arise for a non-resonant branching ratio of 1.8 $\times 10^{-6}$. That is, new physics gives the same rate as in the SM while large CP-violating effects may occur. It should be kept in mind that our numerical results for the CP asymmetries in the high dimuon invariant mass region are plagued with theoretical uncertainties due mainly to real $c\bar{c}$ intermediate states, so that precise predictions are more difficult in this case. Nevertheless, the CP asymmetries provide a particularly useful tool for discovering new physics and their study may gain insight into the mechanism of CP violation. Given an asymmetry of 20% and a branching ratio of 3.0 $\times 10^{-7}$ in the high-s region, a measurement at 3σ level will necessitate at least $7.5 \times 10^8 \ b\bar{b}$ pairs, which seems to be achievable in the B-factory era (see, e.g., Ref. [25]). In this connection it is worth pointing out that the asymmetry A_{CP}^{D} , which is a CP-violating effect related to the angular distribution of μ^- in B and \overline{B} decays, has the piquant feature that it does not require flavor identification and can be obtained from a measurement of the sum of B and \overline{B} events.

As far as new operators are concerned, we have argued that in the case of massless leptons, i.e. l = e or μ , the dominant contributions may come from operators with a non-SM chirality structure. However, since the SM operator basis can accommodate maximum *CP* asymmetries, the inclusion of new operator structures does not affect the main conclusions of our analysis, but it is worth considering once sufficient data are accumulated. We are thus eagerly awaiting the upcoming *B* experiments which will provide useful information on the short-distance coefficients governing FCNC processes like $b \rightarrow s \gamma$ and $b \rightarrow s l^+ l^-$.

We conclude that the *CP* asymmetries in the exclusive decay $\overline{B} \rightarrow K^* l^+ l^-$ can serve as an important test of the SM mechanism of *CP* violation and hence provide a testing ground for new physics. A measurement of these asymmetries in forthcoming *B* experiments would signal the presence of non-standard physics and rule out the SM as a primary source of *CP* violation.

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⁵An analysis of such a *CP*-violating effect in the presence of non-standard *Z* couplings has recently been performed in Ref. [24], which estimates an asymmetry of about 10% in the high dimuon invariant mass region.

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