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# Nonleptonic $B_s$ to charmonium decays: Analysis in pursuit of determining the weak phase $\beta_s$

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We analyze nonleptonic $B_s$ decays to a charmonium state and a light meson, induced by the $b \to c\bar{c}s$ transition, which are useful to access the $B_s$-$\bar{B}_s$ mixing phase $\beta_s$. We use generalized factorization and $SU(3)_F$ symmetry to relate such modes to correspondent $B$ decay channels. We discuss the feasibility of the measurements in the various channels, stressing the importance of comparing different determinations of $\beta_s$ in view of the hints of new physics effects recently emerged in the $B_s$ sector. Finally, adopting a general parametrization of new physics contributions to the decay amplitudes, we discuss how to experimentally constrain new physics parameters.

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## I. INTRODUCTION

The detailed analysis of $CP$ violation in particle physics is a powerful tool to test the standard model (SM) of elementary interactions and unveil the effects of new interactions. The fundamental role in the SM description of $CP$ violation is played by the Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix, which is unitary and implies $CP$ violation if it is complex. The constraints stemming from unitarity can be represented as triangles, the lengths of whose sides are the moduli of products of CKM elements, while the angles represent relative phases between them. The most studied $bd$ unitarity triangle is defined by the relation $V_{ud}V_{ub}^* + V_{cd}V_{cb}^* + V_{td}V_{tb}^* = 0$, and has been probed mainly through the extensive analysis of $B_d$ phenomenology. As a result, the CKM parameters in the Wolfenstein parametrization have been fixed with small errors through the measurement of the sides and the angles of this triangle [1]. The next, already ongoing, effort is to look at processes in which to test the SM requires a greater experimental and theoretical precision. The $B_s$ sector is suitable for such a purpose. As in the $B_d$ case, one of the CKM unitarity constraints involves matrix elements related to $B_s$ decays: $V_{us}V_{ub}^* + V_{cs}V_{cb}^* + V_{ts}V_{tb}^* = 0$, and one of the angles of the corresponding $bs$ triangle is the phase of the $B_s$-$\bar{B}_s$ mixing: $\beta_s = \mathrm{Arg}[-\frac{V_{ts}V_{tb}^*}{V_{cs}V_{cb}^*}]$. In SM $\beta_s$ is expected to be tiny: $\beta_s \simeq 0.017$ rad.

$B_s$ is produced at the $B$ factories running at the peak of $\Upsilon(5S)$ and in hadron collisions. In particular, the experiments CDF and D0 at the Tevatron have obtained a number of remarkable results, such as the measurement of the mixing parameters: The mass difference of the two $B_s$ mass eigenstates has been fixed to $\Delta m_s = 17.77 \pm 0.10(\mathrm{stat}) \pm 0.07(\mathrm{syst})\ \mathrm{ps}^{-1}$ [2], while the value of their width difference $\Delta\Gamma_s$ depends on the constraints on $\beta_s$ adopted in the experimental analysis [3]; noticeably, $\Delta\Gamma_s$ is not small as in the $B_d$ case. Furthermore, these Collaborations have provided us with results which seem to signal new physics (NP) effects. The first one concerns the phase $\beta_s$, extracted from the angular analysis of the time-dependent differential decay width in the process $B_s \to J/\psi\phi$. The study is rather involved: an angular analysis is needed to disentangle the $CP$-even and $CP$-odd components, required since the final state of two vector mesons is not a $CP$ eigenstate. Moreover, the measurement can be carried out either considering flavor tagged or untagged decays. Another issue concerns the use or not of assumptions on the strong phases among the different helicity amplitudes in the considered process: this assumption has been once adopted by the D0 Collaboration in one study [4]. Different results have been obtained from the different analyses [4,5], and, averaging them, the Heavy Flavor Averaging Group has provided a value of $\beta_s$ consistent with SM only at $2.2\sigma$ level: $\phi_s^{J/\psi\phi} = -2\beta_s = -0.77\pm^{0.29}_{0.37}$ or $\phi_s^{J/\psi\phi} = -2\beta_s = -2.36\pm^{0.37}_{0.29}$ [6]. A new measurement announced by CDF: $\beta_s \in [0.0, 0.5] \cup [1.1, 1.5]$ (at 68% CL) [7], if confirmed, would reconcile the SM prediction with experiment.

Another signal of a possible inadequacy of the SM is the measurement of an anomalous like-sign dimuon charge asymmetry of semileptonic $b$-hadron decay, reported by the D0 Collaboration [8] (updating a previous measurement [9]):

$$A_{sl}^b = \frac{N_b^{++} - N_b^{--}}{N_b^{++} + N_b^{--}} = -(9.57 \pm 2.51 \pm 1.46) \times 10^{-3}. \tag{1}$$

Hence, there is a large excess of negatively charged muons over positively charged ones which would have been generated by the oscillation of one neutral $b$ meson into the other, at odds with the SM expectation

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$A^b_{sl} = (-0.310^{+0.083}_{-0.098}) \times 10^{-3}$ [10], a result which might imply an NP effect in the oscillation.

In this complex scenario it is worthwhile to further analyze the $B_s$ sector, trying to identify and reduce the uncertainties affecting the theoretical predictions, with the aim of improving the measurement of $\beta_s$, overcoming the difficulties in $B_s \to J/\psi\phi$. Notice that this channel is considered a golden mode, since it is induced by the $b \to c\bar{c}s$ transition in which, in SM, the only weak phase involved is that of the mixing, so that the indirect $CP$ asymmetry would be proportional to $\sin(2\beta_s)$, much in the same way as $B_d \to J/\psi K_S$ has provided a determination of the angle $\beta$. The feasibility in the reconstruction of the products of the subsequent decays $J/\psi \to \mu^+\mu^-$, $\phi \to K^+K^-$ makes this channel also experimentally appealing.

There are other modes that can be used to access $\beta_s$, namely $B_s \to M_{c\bar{c}} + L$, where $M_{c\bar{c}}$ is a charmonium state $J/\psi$, $\psi(2S)$, $\eta_c$, $\eta_c(2S)$, $\chi_{c0}$, $\chi_{c1}$, $\chi_{c2}$, $h_c$ and $L$ is a light scalar, pseudoscalar or vector meson, $f_0(980)$, $\eta$, $\eta'$ and $\phi$. Each of these channels presents specific features and advantages/difficulties which we want to discuss here. Standing the general theoretical difficulty in the calculation of nonleptonic decay amplitudes, in the next section we discuss approaches to afford the problem, and exploit the generalized factorization to determine the branching fractions in the SM for all the channels listed above, except for those involving $\chi_{c0,2}$ or $h_c$. Generalized factorization gives a vanishing result in these last cases, and hence we only exploit symmetry requirements to estimate their branching ratios using information from the $B_d$ sector. In this way, suitable processes to determine $\beta_s$ can be identified. In Sec. III we also consider the possible impact of new physics in these modes, and discuss how to exploit experimental data to constrain NP parameters. Conclusions are presented in the last section.

## II. $B_s \to M_{c\bar{c}}L$ DECAYS

In SM, assuming CKM unitarity and neglecting the tiny product $V_{ub}V^*_{us}$, the relation holds: $V_{tb}V^*_{ts} = -V_{cb}V^*_{cs}$, and the effective Hamiltonian governing the nonleptonic decays induced by the $b \to c\bar{c}s$ transition reads as [11]:

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}}\Big\{V_{cb}V^*_{cs}[C_1(\mu)O_1(\mu) + C_2(\mu)O_2(\mu)] - V_{tb}V^*_{ts}\Big[\sum_{i=3}^{10,7\gamma,8g} C_i(\mu)O_i(\mu)\Big]\Big\}. \tag{2}$$

$G_F$ is the Fermi constant, the $C_i$ are Wilson coefficients, and $O_i$ are

(i) current–current (tree) operators

$$O_1 = \bar{c}\gamma_\mu(1-\gamma_5)b\bar{s}\gamma^\mu(1-\gamma_5)c,\qquad O_2 = \bar{c}\gamma_\mu(1-\gamma_5)c\bar{s}\gamma^\mu(1-\gamma_5)b, \tag{3}$$

(ii) QCD penguin operators

$$\begin{aligned} O_3 &= (\bar{s}_\alpha b_\alpha)_{V-A}\sum_{q'}(\bar{q}'_\beta q'_\beta)_{V-A},\\ O_4 &= (\bar{s}_\beta b_\alpha)_{V-A}\sum_{q'}(\bar{q}'_\alpha q'_\beta)_{V-A},\\ O_5 &= (\bar{s}_\alpha b_\alpha)_{V-A}\sum_{q'}(\bar{q}'_\beta q'_\beta)_{V+A},\\ O_6 &= (\bar{s}_\beta b_\alpha)_{V-A}\sum_{q'}(\bar{q}'_\alpha q'_\beta)_{V+A}, \end{aligned} \tag{4}$$

(iii) electroweak penguin operators

$$\begin{aligned} O_7 &= \frac{3}{2}(\bar{s}_\alpha b_\alpha)_{V-A}\sum_{q'}e_{q'}(\bar{q}'_\beta q'_\beta)_{V+A},\\ O_8 &= \frac{3}{2}(\bar{s}_\beta b_\alpha)_{V-A}\sum_{q'}e_{q'}(\bar{q}'_\alpha q'_\beta)_{V+A},\\ O_9 &= \frac{3}{2}(\bar{s}_\alpha b_\alpha)_{V-A}\sum_{q'}e_{q'}(\bar{q}'_\beta q'_\beta)_{V-A},\\ O_{10} &= \frac{3}{2}(\bar{s}_\beta b_\alpha)_{V-A}\sum_{q'}e_{q'}(\bar{q}'_\alpha q'_\beta)_{V-A}, \end{aligned} \tag{5}$$

(iv) magnetic moment operators

$$\begin{aligned} O_{7\gamma} &= -\frac{e}{4\pi^2}\bar{s}_\alpha\sigma^{\mu\nu}(m_sP_L + m_bP_R)b_\alpha F_{\mu\nu},\\ O_{8g} &= -\frac{g}{4\pi^2}\bar{s}_\alpha\sigma^{\mu\nu}(m_sP_L + m_bP_R)T^a_{\alpha\beta}b_\beta G^a_{\mu\nu}. \end{aligned} \tag{6}$$

$\alpha$ and $\beta$ are color indices, and $q'$ are the active $q' = (u, d, s, c, b)$ quark fields at the scale $m_b$ with charge $e_{q'}$. The right (left) handed current is defined as $(\bar{q}'_\alpha q'_\beta)_{V\pm A} = \bar{q}'_\alpha\gamma_\nu(1\pm\gamma_5)q'_\beta$, with projection operators $P_{R,L} = \frac{1\pm\gamma_5}{2}$.

The Hamiltonian (2) induces the decays of $SU(3)_F$ related states, namely, the decays of $B_d$, $B^-$ and $B_s$; we consider the general case of the decay of a $B_a$ meson ($a = u, d, s$ being the light flavor index). The simplest approach to compute the matrix element of the Hamiltonian (2) between given initial and final hadronic states is the naive factorization approach. In such an approach, neglecting the magnetic moment operators in (6), the $B_a \to M_{c\bar{c}}L$ amplitude reads:

$$\begin{aligned} \mathcal{A}(\bar{B}_a \to M_{c\bar{c}}L) &= \frac{G_F}{\sqrt{2}}V_{cb}V^*_{cs}a_2^{\rm eff}(\mu)\langle M_{c\bar{c}}|\bar{c}\gamma^\mu(1-\gamma_5)c|0\rangle\\ &\quad\times\langle L|\bar{s}\gamma_\mu(1-\gamma_5)b|\bar{B}_a\rangle, \end{aligned} \tag{7}$$

where $a_2^{\rm eff}(\mu) = a_2(\mu) + a_3(\mu) + a_5(\mu)$ and $a_2 = C_2 + \frac{C_1}{N_c}$, $a_3 = C_3 + \frac{C_4}{N_c} + \frac{3}{2}e_c(C_9 + \frac{C_{10}}{N_c})$ and $a_5 = C_5 + \frac{C_6}{N_c} + \frac{3}{2}e_c(C_7 + \frac{C_8}{N_c})$. However, naive factorization predictions are not able to reproduce several branching ratios for which experimental data are available. Among these there are $B_d$ decays induced by the transition $b \to c\bar{c}s$: some of these modes, which are of interest for the present analysis, are listed in Table I together with the experimental branching fractions.

Several modifications of the naive factorization ansatz have been proposed. One possibility is to consider the Wilson coefficients as effective parameters to be determined from experiment [13]. In principle, this implies that such coefficients are channel-dependent. However, some channels could be related, namely, invoking flavor symmetries, so that universal values for the coefficients can be assumed within a certain class of modes. In our case, this *generalized* factorization approach consists in considering the quantity $a_2^{\rm eff}$ in (7) as a process-dependent parameter to be fixed from experiment. In particular, on the basis of $SU(3)_F$ symmetry, $B_q$ ($B_u$ or $B_d$) decays can be related to analogous $B_s$ decays induced by the same $b \to c\bar{c}s$ transition, so that experimental data concerning $B_q$ modes provide predictions for $B_s$ related ones. Also this method presents some drawbacks, for example, the issue of rescattering in the final state and of the strong phases in the various amplitudes cannot be faced [14]. Nevertheless, it is useful from a phenomenological point of view, at least to understand the size of nonleptonic branching ratios.

A different procedure to analyze nonleptonic decays is the hard-scattering approach, based on the assumption of the dominance of hard gluon exchange and of the suppression of soft mechanisms due to low energy gluon exchanges. A nonleptonic amplitude is expressed as a convolution of a hard kernel, computed in perturbation theory, with the light-cone wave functions of the hadrons involved in the decay. In this so-called perturbative QCD approach (pQCD) the suppression of the soft term is achieved by suitable Sudakov suppression factors, but the uncertainty in the wave functions limits the accuracy of the predictions [15].

A systematic improvement of naive factorization is QCD factorization (BBNS) [16]. In this approach, a factorization formula is written for a nonleptonic $B_a \to M_1 M_2$ decay amplitude ($M_1$ denotes the meson picking up the $B_a$ spectator quark), valid in the heavy quark limit (i.e. up to $\Lambda_{\rm QCD}/m_b$ corrections). This formula reproduces the naive factorization result at leading order in $\alpha_s$ and $\Lambda_{\rm QCD}/m_b$; however, it cannot be applied when the meson that does not pick up the $B_a$ spectator quark is heavy. The knowledge of the meson wave functions is required and represents therefore a limiting factor.

A particular case is represented by the decays $B \to M_{c\bar{c}}L$ considered here. Since the charmonium state is a heavy meson, the BBNS factorization formula does not hold. However, it has been pointed out that, due to the feature of a charmonium meson of being a state with small transverse extension, one can still adopt the factorization formula. Still a problem arises going beyond the leading twist for the wave functions, since the factorization formula contains convolution integrals of such wave functions, and higher twist wave functions do not vanish in the end point, developing divergences. Getting rid of such divergences requires the introduction of a cutoff, a parameter to be fixed from experiment.

Two-body nonleptonic $B$ decays have also been analyzed in a modified formulation of light-cone QCD sum rules originally proposed in [17] to calculate the $B \to \pi\pi$ matrix element, finding results in agreement with QCD factorization. Applying this method to $B$ to charmonium decays, one finds that nonfactorizable contributions are important, but that their inclusion does not allow to reproduce experimental data for $B \to J/\psi K$ [18].

Hence, no satisfactory treatment of nonleptonic $B$ to charmonium decays exists at present, each method having its own advantages/drawbacks. Motivated by the phenomenological importance of these modes, we afford a study based on generalized factorization, aiming at establishing at least the sizes of the branching ratios of these modes and their role for a measurement of $\beta_s$.

TABLE I. Experimental results for the branching fractions $\mathcal{B}(B \to M_{c\bar{c}}K^{(*)})$ ($\times 10^4$) [12]; results in parentheses are from [6].

| $B \to M_{c\bar{c}}K$ | $J/\psi$ | $\eta_c$ | $\psi(2S)$ | $\eta_c(2S)$ |
|---|---|---|---|---|
| $B^-$ | $10.07 \pm 0.35$ | $9.1 \pm 1.3$ | $6.48 \pm 0.35$ | $3.4 \pm 1.8$ |
| $B^0$ | $8.71 \pm 0.32$ | $8.9 \pm 1.6$ | $6.2 \pm 0.6$ | |
| $B \to M_{c\bar{c}}K^*$ | $J/\psi$ | $\eta_c$ | $\psi(2S)$ | $\eta_c(2S)$ |
| $B^-$ | $14.3 \pm 0.8$ | $12.0 \pm 7.0$ | $6.7 \pm 1.4$ | |
| $B^0$ | $13.3 \pm 0.6$ | $9.6 \pm 3.3$ | $7.2 \pm 0.8$ | $<3.9$ |
| $B \to M_{c\bar{c}}K$ | $\chi_{c0}$ | $\chi_{c1}$ | $\chi_{c2}$ | $h_c$ |
| $B^-$ | $1.43 \pm 0.21$ | $5.1 \pm 0.5$ | $<0.29$ | $<0.38$ |
| $B^0$ | $<1.13$ | $3.9 \pm 0.4$ | $<0.26$ | |
| $B \to M_{c\bar{c}}K^*$ | $\chi_{c0}$ | $\chi_{c1}$ | $\chi_{c2}$ | $h_c$ |
| $B^-$ | $<2.1$ | $3.6 \pm 0.9$ | $<0.12$ | |
| $B^0$ | $1.70 \pm 0.40$ | $2.0 \pm 0.6$ | $<0.36(0.66 \pm 0.19)$ | $(< 2.2)$ |

To apply Eq. (7) to the modes we are analyzing, we need the following hadronic quantities:

(i) charmonium decay constants:

$$\langle \eta_c(q)|\bar{c}\gamma_\mu\gamma_5 c|0\rangle = -if_{\eta_c} q_\mu,$$
$$\langle J/\psi(q,\epsilon)|\bar{c}\gamma_\mu c|0\rangle = f_\psi m_\psi \epsilon^*_\mu, \qquad (8)$$
$$\langle \chi_{c1}(q,\epsilon)|\bar{c}\gamma_\mu\gamma_5 c|0\rangle = f_{\chi_{c1}} m_{\chi_{c1}} \epsilon^*_\mu,$$

($\epsilon(\lambda)$ polarization vector); for $\chi_{c0,2}$ and $h_c$ one has $\langle \chi_{c0}(q)|\bar{c}\gamma^\mu c|0\rangle = \langle \chi_{c2}(q,\epsilon)|\bar{c}\gamma^\mu c|0\rangle = \langle h_c(q,\epsilon)|\bar{c}\gamma^\mu(1-\gamma_5)c|0\rangle = 0$;

(ii) $\bar{B}_a \to L$ form factors, with $L$ a pseudoscalar ($P$) or a scalar ($S$) meson:

$$\langle P(S)(p')|\bar{s}\gamma^\mu(\gamma_5)b|\bar{B}_a(p)\rangle = F_1(q^2)\left[(p+p')^\mu - \frac{m^2_{B_a} - m^2_{P(S)}}{q^2} q^\mu\right] + F_0(q^2)\frac{m^2_{B_a} - m^2_{P(S)}}{q^2} q^\mu, \qquad (9)$$

(iii) $\bar{B}_a \to L$ form factors, with $L$ a vector ($V$) meson:

$$\langle L(p',\epsilon)|\bar{s}\gamma_\mu(1-\gamma_5)b|\bar{B}_a(p)\rangle = \frac{2V(q^2)}{m_{B_a}+m_L}\epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}p^\alpha p'^\beta - i\left[\epsilon^*_\mu(m_{B_a}+m_L)A_1(q^2) - (\epsilon^*\cdot q)(p+p')_\mu \frac{A_2(q^2)}{m_{B_a}+m_L} - (\epsilon^*\cdot q)\frac{2m_L}{q^2}(A_3(q^2)-A_0(q^2))q_\mu\right], \qquad (10)$$

with $A_3(q^2) = \frac{m_{B_a}+m_L}{2m_L}A_1(q^2) - \frac{m_{B_a}-m_L}{2m_L}A_2(q^2)$. The first two equations in (8) also hold for $\eta_c(2S)$ and $\psi(2S)$, respectively, while the vanishing of the matrix elements $\langle \chi_{c0,2}(q)|\bar{c}\gamma^\mu c|0\rangle$ and $\langle h_c(q,\epsilon)|\bar{c}\gamma^\mu(1-\gamma_5)c|0\rangle$ implies that the $B_a \to \chi_{c0,2}L$ and $B_a \to h_c L$ amplitudes vanish in the factorization approximation.

By the factorization ansatz one has expressions for the decay widths. Moreover, for decays in two $J=1$ mesons, also the polarization fractions can be computed, namely $f_L$, the fraction of the decay width when both the final mesons are longitudinally polarized.[1] The results are the following:

(iv) modes where $M_{c\bar{c}}$ is either a $J^{PC}=1^{--}$ charmonium state ($J/\psi$ or $\psi(2S)$), or a $J^{PC}=1^{++}$ P-wave $\chi_{c1}$ meson:

$$\Gamma(B_a \to M_{c\bar{c}}L) = \frac{G_F^2|V_{cb}V^*_{cs}|^2(a_2^{\rm eff})^2 f^2_{M_{c\bar{c}}}}{32\pi m^3_{B_a}}[F_1^{B_a\to L}(m^2_{M_{c\bar{c}}})]^2\lambda^{3/2}(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_L), \qquad (11)$$

$$\begin{aligned}\Gamma(B_a \to M_{c\bar{c}}V) &= \frac{G_F^2|V_{cb}V^*_{cs}|^2(a_2^{\rm eff})^2 f^2_{M_{c\bar{c}}}}{16\pi m^3_{B_a}}\frac{\lambda^{1/2}(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V)}{8m_V^2}\Big\{(m_{B_a}+m_V)^2[A_1^{B_a\to V}(m^2_{M_{c\bar{c}}})]^2[\lambda(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V) + 12m^2_{M_{c\bar{c}}}m^2_V] \\ &+ \frac{[A_2^{B_a\to V}(m^2_{M_{c\bar{c}}})]^2}{(m_{B_a}+m_V)^2}\lambda^2(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V) - 2A_1^{B_a\to V}(m^2_{M_{c\bar{c}}})A_2^{B_a\to V}(m^2_{M_{c\bar{c}}})(m^2_{B_a} - m^2_{M_{c\bar{c}}} - m^2_V)\lambda(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V) \\ &+ 8m^2_{M_{c\bar{c}}}m^2_V\frac{[V^{B_a\to V}(m^2_{M_{c\bar{c}}})]^2}{(m_{B_a}+m_V)^2}\lambda(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V)\Big\}, \end{aligned} \qquad (12)$$

$$\begin{aligned}f_L(B_a \to M_{c\bar{c}}V) &= \frac{1}{\Gamma(B_a\to M_{c\bar{c}}V)}\frac{G_F^2|V_{cb}V^*_{cs}|^2(a_2^{\rm eff})^2 f^2_{M_{c\bar{c}}}}{16\pi m^3_{B_a}}\frac{\lambda^{1/2}(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V)}{8m_V^2}\Big\{(m_{B_a}+m_V)[A_1^{B_a\to V}(m^2_{M_{c\bar{c}}})](m^2_{B_a} - m^2_{M_{c\bar{c}}} - m^2_V) \\ &- \lambda(m^2_{B_a}, m^2_{M_{c\bar{c}}}, m^2_V)\frac{[A_2^{B_a\to V}(m^2_{M_{c\bar{c}}})]}{(m_{B_a}+m_V)}\Big\}^2; \end{aligned} \qquad (13)$$

(v) modes with a $J^{PC}=0^{-+}$ $P_{c\bar{c}}$ ($\eta_c$ or $\eta_c(2S)$) charmonium state:

[1] The two $J=1$ mesons in the final state have the same helicity since the decaying $B_a$ is spinless.

$$\Gamma(B_a \to P_{c\bar{c}}L) = \frac{G_F^2|V_{cb}V_{cs}^*|^2(a_2^{\rm eff})^2 f_{P_{c\bar{c}}}^2}{32\pi m_{B_a}^3}[F_0^{B_a\to L}(m_{P_{c\bar{c}}}^2)]^2(m_{B_a}^2 - m_L^2)^2\lambda^{1/2}(m_{B_a}^2, m_{P_{c\bar{c}}}^2, m_L^2), \tag{14}$$

$$\Gamma(B_a \to P_{c\bar{c}}V) = \frac{G_F^2|V_{cb}V_{cs}^*|^2(a_2^{\rm eff})^2 f_{P_{c\bar{c}}}^2}{32\pi m_{B_a}^3}[A_0^{B_a\to P}(m_{P_{c\bar{c}}}^2)]^2\lambda^{3/2}(m_{B_a}^2, m_{P_{c\bar{c}}}^2, m_V^2). \tag{15}$$

$\lambda$ is the triangular function, $\lambda(a, b, c) = (a - b - c)^2 - 4bc$. Equations (11) and (14) apply to both the cases in which the light meson is pseudoscalar or scalar, $L = P, S$.

Using in these expressions the coefficients $a_i(\mu)$ computed in renormalization group improved perturbation theory, the experimental data are badly reproduced: the $b \to c\bar{c}s$ induced modes under scrutiny are color suppressed, and the predictions of naive factorization undershoot the data. The most striking discrepancy is for the modes with $\chi_{c0}$ in the final state, which have sizable rates despite the fact that their amplitudes vanish in the factorization approach. Our strategy is to exploit the data in Table I to determine an effective parameter $a_2^{\rm eff}$ (generally channel-dependent) and, assuming $SU(3)_F$ symmetry, to use these values to predict the flavor related $B_s$ decays. Since the results depend on the form factors, to estimate this hadronic uncertainty we use two sets of form factors computed by variants of the QCD sum rule method [19], the set in [20] obtained using sum rules based on the short-distance expansion, and the set in [21] based on the light-cone expansion. In the case of $B_s \to \phi$ and $B_s \to f_0(980)$ we use form factors determined by light-cone sum rules [22,23]. It is noticeable that in the future, when enough experimental data will be available for both $B_d$ and $B_s$ modes, a comparison between the effective Wilson coefficients $a_2^{\rm eff}$ from the various $SU(3)_F$ related modes will be possible.

The numerical inputs $V_{cb} = 0.0412 \pm 0.0011$, $V_{cs} = 1.04 \pm 0.06$, $\tau(B^-) = (1.638 \pm 0.011)$ ps, $\tau(B^0) = (1.530 \pm 0.009)$ ps, $\tau(B_s) = (1.470^{+0.026}_{-0.027})$ ps, together with the values of the meson masses, are taken from the Particle Data Group [12]. Moreover, from $J/\psi(\psi(2S)) \to e^+e^-$ [12] we obtain: $f_{J/\psi} = (416.3 \pm 5.3)$ MeV and $f_{\psi(2S)} = (296.1 \pm 2.5)$ MeV. The decay constant of $\eta_c$ comes from $\eta_c \to 2\gamma$: $f_{\eta_c} = (380.0 \pm 87.1)$ MeV, while only the upper bound $f_{\eta_c(2S)} < 438.8$ MeV is known. In the heavy quark limit, the pseudoscalar and vector charmonia are collected in a doublet of states with degenerate masses and same decay constants. Therefore, $f_{\eta_c(2S)}$ can be obtained using the relation

$$f_{\eta_c(2S)} = \frac{f_{\eta_c}}{f_{J/\psi}} f_{\psi(2S)} = (270.3 \pm 62.0)\ \text{MeV}; \tag{16}$$

symmetry breaking terms, coming from removing the meson degeneracy, are expected to cancel in the ratio. Since the constant $f_{\chi_{c1}}$ is not known, for the modes involving $\chi_{c1}$ we determine the product $a_2^{\rm eff} f_{\chi_{c1}}$ from data. $SU(3)_F$ symmetry allows to relate $B_s$ decays to those listed in Table I: data on $B \to M_{c\bar{c}}K$ allow us to predict $B_s \to M_{c\bar{c}}\eta^{(\prime)}$, while information on $B \to M_{c\bar{c}}K^*$ is used to predict $B_s \to M_{c\bar{c}}\phi$. As for $B_s \to M_{c\bar{c}}f_0(980)$, they are obtained using the effective $|a_2|$ determined from $B \to M_{c\bar{c}}K$. The $B_s \to \eta^{(\prime)}$ form factors are related to the analogous $B \to K$ form factors: for a generic form factor $F$ we have $F^{B_s\to\eta} = -\sin\theta F^{B\to K}$ and $F^{B_s\to\eta'} = \cos\theta F^{B\to K}$ where $\theta$ is the mixing angle in the flavor basis [24]

$$\eta = \eta_q \cos\theta - \eta_s \sin\theta, \quad \eta' = \eta_q \sin\theta + \eta_s \cos\theta, \tag{17}$$

with $\eta_q = (\bar{u}u + \bar{d}d)/\sqrt{2}$ and $\eta_s = \bar{s}s$. The mixing angle between $\eta_q$ and $\eta_s$ can be fixed to the value measured by the KLOE Collaboration: $\theta = (41.5 \pm 0.3_{\rm stat} \pm 0.7_{\rm syst} \pm 0.6_{\rm th})^\circ$ [25], which agrees with the outcome of a QCD sum rule analysis of the radiative $\phi \to \eta^{(\prime)}\gamma$ modes [26]. The errors on the various $a_2^{\rm eff}$ correspond to the uncertainty of the form factors at zero momentum transfer and to the experimental errors of the branching ratios of the related modes reported in Table I. In the case of the transitions involving $\eta$ or $\eta'$, the uncertainty on the form factors is not included since, on the basis of $SU(3)_F$, the dependence on the form factors cancels when $B_s \to M_{c\bar{c}}\eta^{(\prime)}$ branching ratio is related to $B \to M_{c\bar{c}}K$, leaving only a dependence on the $\eta$-$\eta'$ mixing angle. The obtained values of $a_2^{\rm eff}$ are collected in Table II, and the predictions for $B_s$ branching

TABLE II. Effective Wilson coefficients $a_2^{\rm eff}$ and combination $a_2^{\rm eff} f_{\chi_{c1}}$ (in GeV) appearing in the decay amplitudes of the various $B_s \to M_{c\bar{c}}L$ modes, obtained from the $SU(3)_F$ related $B$ decay modes and using the form factors in Ref. [20] (CDSS) and Ref. [21] (BZ).

| mode | $\lvert a_2^{\rm CDSS}\rvert$ | $\lvert a_2^{\rm BZ}\rvert$ | mode | $\lvert a_2^{\rm CDSS}\rvert$ | $\lvert a_2^{\rm BZ}\rvert$ |
|---|---|---|---|---|---|
| $J/\psi\,\eta(\eta')$ | $0.40 \pm 0.007$ | $0.26 \pm 0.005$ | $\eta_c\,\eta(\eta')$ | $0.36 \pm 0.03$ | $0.25 \pm 0.02$ |
| $\psi(2S)\eta(\eta')$ | $0.50 \pm 0.02$ | $0.31 \pm 0.01$ | $\eta_c(2S)\eta(\eta')$ | $0.31 \pm 0.08$ | $0.21 \pm 0.06$ |
| mode | $\lvert a_2^{\rm CDSS} f_{\chi_{c1}}\rvert$ | $\lvert a_2^{\rm BZ} f_{\chi_{c1}}\rvert$ | mode | $\lvert a_2^{\rm CDSS} f_{\chi_{c1}}\rvert$ | $\lvert a_2^{\rm BZ} f_{\chi_{c1}}\rvert$ |
| $\chi_{c1}\,\eta(\eta')$ | $0.122 \pm 0.006$ | $0.076 \pm 0.004$ | $\chi_{c1} f_0$ | $0.122 \pm 0.016$ | $0.076 \pm 0.010$ |
| $\chi_{c1}\,\phi$ | | $0.0345 \pm 0.006$ | | | |

TABLE III. Branching ratios ($\times 10^4$) of the decays $B_s \to M_{c\bar{c}} L$ using the form factors in [20] (CDSS) and in [21] (BZ). The experimental results for $B_s \to J/\psi(\psi(2S))\phi$ are taken from PDG [12]; the branching fractions of $B_s \to J/\psi\eta(\eta')$, measured by the Belle Collaboration [27], are quoted combining the errors in quadrature. For $B_s \to M_{c\bar{c}} f_0$, the effective coefficient $a_2$ obtained from the $B \to K$ mode and the form factors CDSS and BZ are used; the two experimental results are due to the LHCb [28] and Belle [29] Collaborations.

| Mode | $\mathcal{B}$ (CDSS) | $\mathcal{B}$ (BZ) | Exp. | Mode | $\mathcal{B}$ (CDSS) | $\mathcal{B}$ (BZ) |
|---|---|---|---|---|---|---|
| $J/\psi\eta$ | $4.3 \pm 0.2$ | $4.2 \pm 0.2$ | $3.32 \pm 1.02$ [27] | $\eta_c\eta$ | $4.0 \pm 0.7$ | $3.9 \pm 0.6$ |
| $J/\psi\eta'$ | $4.4 \pm 0.2$ | $4.3 \pm 0.2$ | $3.1 \pm 1.39$ [27] | $\eta_c\eta'$ | $4.6 \pm 0.8$ | $4.5 \pm 0.7$ |
| $\psi(2S)\eta$ | $2.9 \pm 0.2$ | $3.0 \pm 0.2$ | | $\eta_c(2S)\eta$ | $1.5 \pm 0.8$ | $1.4 \pm 0.7$ |
| $\psi(2S)\eta'$ | $2.4 \pm 0.2$ | $2.5 \pm 0.2$ | | $\eta_c(2S)\eta'$ | $1.6 \pm 0.9$ | $1.5 \pm 0.8$ |
| $J/\psi\phi$ | — | $16.7 \pm 5.7$ | $13 \pm 4$ [12] | $\eta_c\phi$ | — | $15.0 \pm 7.8$ |
| $\psi(2S)\phi$ | — | $8.3 \pm 2.7$ | $6.8 \pm 3.0$ [12] | | | |
| $\chi_{c1}\eta$ | $2.0 \pm 0.2$ | $2.0 \pm 0.2$ | | $\chi_{c1}f_0$ | $1.88 \pm 0.77$ | $0.73 \pm 0.30$ |
| $\chi_{c1}\eta'$ | $1.9 \pm 0.2$ | $1.8 \pm 0.2$ | | $\chi_{c1}\phi$ | — | $3.3 \pm 1.3$ |
| $J/\psi f_0$ | $4.7 \pm 1.9$ | $2.0 \pm 0.8$ | $3.2 \pm 1.3$ [28] | $\eta_c f_0$ | $4.1 \pm 1.7$ | $2.0 \pm 0.9$ |
| | | | $2.32 \pm 0.96$ [29] | | | |
| $\psi(2S)f_0$ | $2.3 \pm 0.9$ | $0.89 \pm 0.36$ | | $\eta_c(2S)f_0$ | $0.58 \pm 0.38$ | $1.3 \pm 0.8$ |

TABLE IV. Branching ratios ($\times 10^4$) of $B_s$ decays into p-wave charmonia.

| Mode | $\mathcal{B}$ | Mode | $\mathcal{B}$ | Mode | $\mathcal{B}$ |
|---|---|---|---|---|---|
| $\chi_{c0}\eta$ | $0.85 \pm 0.13$ | $\chi_{c2}\eta$ | $<0.17$ | $h_c\eta$ | $<0.23$ |
| $\chi_{c0}\eta'$ | $0.87 \pm 0.13$ | $\chi_{c2}\eta'$ | $<0.17$ | $h_c\eta'$ | $<0.23$ |
| $\chi_{c0}f_0$ | $1.15 \pm 0.17$ | $\chi_{c2}f_0$ | $<0.29$ | $h_c f_0$ | $<0.30$ |
| $\chi_{c0}\phi$ | $1.59 \pm 0.38$ | $\chi_{c2}\phi$ | $<0.10(0.62 \pm 0.17)$ | $h_c\phi$ | $(<1.9)$ |

ratios in Tables III and IV. In Table III the available experimental data are also included, with a general agreement with the corresponding theoretical results.

Several remarks are in order. The values of $a_2^{\rm eff}$ derived from the form factors in Ref. [20] are larger than the ones derived from Ref. [21]; they also turn out to be channel-dependent. Their range (0.2–0.3) or (0.3–0.5) is larger than the one obtained in the QCDF and pQCD approaches which undershoot the data. As appears from Tables III and IV, all the modes have sizeable branching fractions, so that they are promising candidates for measurements of $\beta_s$. The modes involving $\eta$, $\eta'$, $f_0$ present, with respect to $B_s \to J/\psi\phi$, the advantage that the final state is a $CP$ eigenstate, not requiring any angular analysis. However, the channels with $\eta$ and $\eta'$ could be useful only after a number of events will be accumulated, since at least two photons are required for the reconstruction. On the other hand, it should be observed that a possible future Super B factory running at the $\Upsilon(5S)$ peak could promisingly study modes such as $B_s \to J/\psi\eta^{(\prime)}$ [30] with a better performance with respect to an hadronic machine [31].

As discussed in [23,32,33], the mode $B_s \to J/\psi f_0(980)$ has appealing features since, compared with the $\eta$ and $\eta'$, the $f_0$ can be easily reconstructed in the $\pi^+\pi^-$ final state, which occurs with a large rate: $\mathcal{B}(f_0(980) \to \pi^+\pi^-) = (50^{+7}_{-8})\%$ [34], so that this channel can be accessed.[2] Indeed, the decay $B_s \to J/\psi f_0$ has been very recently observed by the LHCb [28] and Belle [29] Collaborations, and a preliminary result is also reported by the CDF Collaboration [35].

The LHCb Collaboration has measured, in proton-proton collisions at c.o.m. energy of 7 TeV and using a fit to the $\pi^+\pi^-$ spectrum with interfering resonances, the ratio [28]:

$$R_{f_0/\phi} = \frac{\Gamma(B_s \to J/\psi f_0; f_0 \to \pi^+\pi^-)}{\Gamma(B_s \to J/\psi\phi; \phi \to K^+K^-)} = 0.252^{+0.046}_{-0.032}({\rm stat})^{+0.027}_{-0.033}({\rm syst}). \qquad (18)$$

[2]The quark content of $f_0(980)$ is not completely known. Under the $\bar{q}q$ assignment, this meson might be a mixture of the isosinglet $\bar{n}n$ and $\bar{s}s$ ($n = u$, $d$) components. The mixing angle can be fixed using experimental information on, for instance, the decays $J/\psi \to \phi f_0$ and $J/\psi \to \omega f_0$: $\mathcal{B}(J/\psi \to \phi f_0) = (3.2 \pm 0.9) \times 10^{-4}$, $\mathcal{B}(J/\psi \to \omega f_0) = (1.4 \pm 0.5) \times 10^{-4}$, which might signal a nonstrange component of $f_0$ and the consequent reduction of $\mathcal{B}(B_s \to J/\psi f_0)$ by about 30%.

This result, together with the experimental measurement of $\mathcal{B}(B_s \to J/\psi \phi)$ included in Table III and $\mathcal{B}(\phi \to K^+K^-) = (48.9 \pm 0.5) \times 10^{-2}$ [12], corresponds to $\mathcal{B}(B_s \to J/\psi f_0) = (3.2 \pm 1.3) \times 10^{-4}$, combining all the uncertainties in quadrature. The Belle Collaboration, analyzing 121.4 fb$^{-1}$ of data collected at the $\Upsilon(5S)$ resonance, has obtained [29]:

$$\mathcal{B}(B_s \to J/\psi f_0; f_0 \to \pi^+\pi^-) = [1.16^{+0.31}_{-0.19}(\text{stat})^{+0.15}_{-0.17}(\text{syst})^{+0.26}_{-0.18}(N_{B_s^{(*)}\bar{B}_s^{(*)}})] \times 10^{-4} \tag{19}$$

which corresponds, combining the uncertainties in quadrature, to $\mathcal{B}(B_s \to J/\psi f_0) = (2.32 \pm 0.96) \times 10^{-4}$. The preliminary result reported by the CDF Collaboration reads [35]:

$$R_{f_0/\phi} = \frac{\mathcal{B}(B_s \to J/\psi f_0)}{\mathcal{B}(B_s \to J/\psi \phi)} \frac{\mathcal{B}(f_0 \to \pi^+\pi^-)}{\mathcal{B}(\phi \to K^+K^-)} = 0.292 \pm 0.020(\text{stat}) \pm 0.017(\text{syst}), \tag{20}$$

which corresponds to $\mathcal{B}(B_s \to J/\psi f_0) = (3.7 \pm 1.3) \times 10^{-4}$. The three measurements are all compatible with our prediction.

The results for $B_s \to \eta_c L$ are also included in Table III. Although these channels have sizable branching fractions, they present the drawback of the difficult reconstruction of the $\eta_c$.

Let us now consider $B_s$ decays to $p$-wave charmonia. We have stressed that, among these decays, the only one with nonvanishing amplitude in the factorization assumption is that with $\chi_{c1}$ in the final state. In the other cases, i.e. for modes involving $\chi_{c0,2}$ and $h_c$ collected in Table IV, the results are obtained determining the decay amplitudes from the $B$ decay data by making use of the $SU(3)_F$ symmetry. In this case, the differences between the $B$ and $B_s$ decays arise from the phase space and lifetimes of the heavy mesons. As for the mechanism inducing such processes, one possibility, put forward in [36], is that rescattering can be responsible for the observed branching fractions. Among these channels, $B_s \to \chi_{c0}\phi$ is of prime interest and very promising for both hadron colliders and $B$ factories. Even though $\mathcal{B}(B_s \to \chi_{c0}\phi)$ is 1 order of magnitude smaller than $\mathcal{B}(B_s \to J/\psi \phi)$ it has appealing features, in particular, as far as the potential of the LHCb experiment is concerned. As a matter of fact, the $\chi_{c0}$ dominant decay modes are $\chi_{c0} \to \rho^+\pi^-\pi^0$, $\rho^-\pi^+\pi^0$, $\pi^+\pi^-\pi^+\pi^-$, with branching fractions of order $10^{-2}$. Therefore the final state consists of six charged hadrons, and the particle identification information from the RICH detectors could suppress the background. Furthermore, the vertex detector might be particularly efficient for these channels.

Considering finally the polarization fractions, the results for the longitudinal polarization fractions $f_L$ for the modes with two $J = 1$ mesons in the final state can be found in Table V. There is agreement with experiment for $B_s \to J/\psi \phi$, the only mode for which data on $f_L$ are available. This is at odds with the case of a few suppressed $B$ decays to two light vector mesons, in which the experimental datum is not reproduced assuming factorization.

TABLE V. Longitudinal polarization fraction $f_L$ ($\times 10^2$) for $B_s$ decays to two $J = 1$ mesons.

| Mode | Prediction | Experiment |
|---|---|---|
| $J/\psi\phi$ | $51.3 \pm 5.8$ | $54.1 \pm 1.7$ |
| $\psi(2S)\phi$ | $41.0 \pm 3.7$ | |
| $\chi_{c1}\phi$ | $43.9 \pm 4.4$ | |

Before concluding this section, we would like to comment about the accuracy of the results collected in Tables III, IV, and V. As we have already stressed, the uncertainties quoted in Tables III and IV are due to the errors affecting the parameters of the form factors used in the calculation and to the errors on the experimental branching ratios of the corresponding $B_d$ decay modes exploited to obtain $a_2^{\text{eff}}$. In principle, other uncertainties affect such predictions due to the limited accuracy of $SU(3)_F$ and to the neglect of possible nonfactorizable terms.

$SU(3)_F$ is usually considered a reliable assumption, and in the context of $B_s$ decays such an issue has been discussed in several analyses. In [37] it was shown that this symmetry allows to reliably relate $B_d$ and $B_s$ decay amplitudes not only in magnitude but also as far as their strong phases are concerned. Another argument is provided by the analysis in [38], where it has been shown that $B_s \to K^+K^-$ can be related to $B_d \to \pi^+\pi^-$ using U-spin symmetry, which amounts to replace all the $d$ quarks in the decay by $s$ quarks, an assumption similar but weaker than the $SU(3)_F$ adopted here.

Concerning the role of nonfactorizable terms, it is more difficult to assess how large such terms are. However, as pointed out in [36,39], it is likely and supported by available data that such terms are negligible in those cases in which there is a sizeable factorizable term, while they could be relevant whenever the leading factorizable term is either absent (as in the case of the modes in Table IV) or it is effectively suppressed (being, e.g., loop induced or strongly CKM suppressed). In [39] it was shown that the polarization fractions are sensible probes of such nonfactorizable contributions. Actually, naive and generalized factorization provide the same result in the case of the polarization fractions, since they differ only for the value of $a_2^{\text{eff}}$ which cancels in the ratio defining a polarization fraction. To modify the prediction for $f_L$ one should either consider approaches in which the three polarization fractions (the longitudinal and the two transverse ones) involve different Wilson coefficients, or invoke again other mechanisms such as rescattering. The first case occurs in QCD factorization and in pQCD. Among the nonfactorizable mechanisms, rescattering has been proposed as a solution to the puzzle of the polarization fractions in some $B$ decays to two light vector mesons, when the considered process is suppressed as in the penguin induced mode $B \to \phi K^*$ [39]. In the case of the modes in Table III the agreement of our prediction for the longitudinal polarization fraction $f_L$ with the experimental datum for $B_s \to J/\psi \phi$, the only $B_s$ decay mode for which $f_L$ has been measured, supports this argument.

Finally, concerning the modes in Table IV, as we have already stressed, the results do not rely on factorization, but only on the use of $SU(3)_F$ to obtain the amplitude, hence only the issue of the $SU(3)_F$ accuracy applies to these modes.

## III. NEW PHYSICS EFFECTS IN NONLEPTONIC $B_s$ DECAYS: GENERAL ANALYSIS

As mentioned in the Introduction, hints of deviations from SM predictions have recently been found in $B_s$ phenomenology, hence it is worth considering the effects of new physics in the $B_s$ sector, which may show up in mixing and/or in decay amplitudes.

New physics in $B_s - \bar{B}_s$ mixing can modify the mixing phase $\beta_s$. We refer to this phase as to $\beta_s^{\rm eff}$, which contains SM as well as NP contributions: $\beta_s^{\rm eff} = \beta_s^{\rm SM} + \beta_s^{\rm NP}$. This effect is the same for all decay modes, and simply shifts the value of $\beta_s$. On the other hand, NP in the decay amplitudes can affect various channels in different ways, even for modes induced by the same quark transition, as we specify in the following.

Let us discuss the possibility that experimental results for nonleptonic $B_s$ decays deviate from the predictions given in the previous section. Such predictions rely on $SU(3)$ flavor symmetry and on experimental data on corresponding $B_d$ decays, in which no NP effects have been detected at the present level of accuracy. Deviations in $B_s$ decay rates with respect to the predictions could be due to a violation of $SU(3)_F$ symmetry, which is generally expected at a few percent level. However, there is the more exciting possibility of deviations due to NP effects with small contributions in $B_d$ oscillations and decays and detectable contributions in $B_s$, an eventuality which is interesting to consider for the modes studied in this paper. Such modes receive contribution both from tree-level and loop diagrams, so that one would expect NP to affect them negligibly. However, there are scenarios in which the contribution of new particles in loop diagrams can be competitive with the SM tree level diagrams. This is the case, for example, of supersymmetric scenarios in which one loop gluino exchanges for $b \to s$ transition could give a sizeable contribution to the $b \to c\bar{c}s$ induced modes. This would affect the branching ratios of such modes, and the $CP$ asymmetries, since new phases could arise through the soft supersymmetry breaking terms. To avoid constraints on such phases from existing limits on dipole electric moments, one should consider flavor dependent phases.

Here we do not focus on a specific NP model, rather we parametrize the effects of new physics in a general way, i.e. in terms of an amplitude, a weak and a strong phase, and discuss how these quantities can be constrained by experimental data on the modes considered above.

In a customary notation, $\mathcal{A}_f$ is the amplitude for $B_s \to f$ decay to a generic final state $f$ ($CP$ eigenstate, common to $B_s$ and $\bar{B}_s$) which, in our case, is of the kind $M_{c\bar{c}}L$.[3] The corresponding $\bar{B}_s$ decay amplitude is denoted as $\bar{\mathcal{A}}_f$. Being interested in $CP$ asymmetries, we introduce the quantity

$$\lambda_f = e^{-2i\beta_s^{\rm eff}}\left(\frac{\bar{\mathcal{A}}_f}{\mathcal{A}_f}\right), \tag{21}$$

in terms of which one can write the mixing-induced $CP$ asymmetry $S_f$ and the direct $CP$ asymmetry $C_f$:

$$S_f = \frac{2\Im(\lambda_f)}{1+|\lambda_f|^2}, \qquad C_f = \frac{1-|\lambda_f|^2}{1+|\lambda_f|^2}. \tag{22}$$

Assuming that there is a single dominant NP amplitude (or that all NP amplitudes have the same weak and strong phases relative to the SM), we write:

$$\mathcal{A}_f = \mathcal{A}_f^{\rm SM}(1 + Re^{i\theta_{\rm NP}}e^{i\delta_{\rm NP}}), \tag{23}$$

where $R = \frac{|\mathcal{A}_f^{\rm NP}|}{|\mathcal{A}_f^{\rm SM}|}$ is the ratio of the modulus of the NP amplitude and that of the SM one, while $\theta_{\rm NP}(\delta_{\rm NP})$ is the strong (weak) NP phase with respect to the SM part. Our working hypothesis is that no NP affects $B$ decays: actually, the new weak phase $\delta_{\rm NP}$ would be the same in $B$ and $B_s$ decays, depending only on the underlying quark transition, however $R$ and $\theta_{\rm NP}$ depend on the matrix elements of the operators between initial and final states, which can be different.

Using the definition in Eq. (23) and considering that: $\bar{\mathcal{A}}_f = \bar{\mathcal{A}}_f^{\rm SM}(1 + Re^{i\theta_{\rm NP}}e^{-i\delta_{\rm NP}})$, we get for the $CP$-averaged branching fraction

$$\mathcal{B}^{\rm exp} = \mathcal{B}^{\rm SM}[1 + 2R\cos(\theta_{\rm NP})\cos(\delta_{\rm NP}) + R^2], \tag{24}$$

where from now on we shall omit the label $f$ in the quantities in Eq. (24) though they are channel-dependent (except for $\delta_{\rm NP}$ which is the same for all the modes induced by the same underlying quark transition). From Eq. (21), we obtain:

$$\lambda_f = e^{-2i\beta_s^{\rm eff}}\frac{1 + Re^{i(\theta_{\rm NP}-\delta_{\rm NP})}}{1 + Re^{i(\theta_{\rm NP}+\delta_{\rm NP})}}\frac{\bar{\mathcal{A}}_f^{\rm SM}}{\mathcal{A}_f^{\rm SM}}. \tag{25}$$

Neglecting double Cabibbo suppressed contributions to the SM amplitudes we have:

[3] In the case of two vectors in the final state $f$ denotes one of the final state components being the $CP$ eigenstate.

$$S_f = -\eta_f \frac{\sin(2\beta_s^{\rm eff}) + 2R\cos\theta_{\rm NP}\sin(2\beta_s^{\rm eff} + \delta_{\rm NP}) + R^2\sin(2\beta_s^{\rm eff} + 2\delta_{\rm NP})}{1 + 2R\cos\theta_{\rm NP}\cos\delta_{\rm NP} + R^2}, \tag{26}$$

$$C_f = -\frac{2R\sin\theta_{\rm NP}\sin\delta_{\rm NP}}{1 + 2R\cos\theta_{\rm NP}\cos\delta_{\rm NP} + R^2}, \tag{27}$$

$\eta_f$ being the $CP$ eigenvalue of the final state $f$. In absence of NP we recover the SM results $\lambda_f = e^{-2i\beta_s}$, $S_f = -\eta_f\sin(2\beta_s)$, $C_f = 0$. These results would also hold if NP contributes only to $B_s - \bar{B}_s$ mixing with $\beta_s \to \beta_s^{\rm eff}$.

The three Eqs. (24), (26), and (27) allow to determine the NP parameters $R$, $\theta_{\rm NP}$, $\delta_{\rm NP}$, once experimental data on $\mathcal{B}(B_s \to f)$, $S_f$ and $C_f$ are available for a given final state $f$. Assuming $R \ll 1$, we obtain:

$$\theta_{\rm NP} = \arctan\left(\frac{-C_f}{\tilde{S}_f}\right), \tag{28}$$

$$\delta_{\rm NP} = \arctan\left[\frac{(1+\Sigma)}{\Sigma}\tilde{S}_f\right], \tag{29}$$

$$R = \frac{\Sigma}{2\cos(\theta_{\rm NP})\cos(\delta_{\rm NP})}, \tag{30}$$

where $\Sigma$ and $\tilde{S}_f$ parametrize deviations from the SM:

$$\Sigma = \frac{\mathcal{B}^{\rm exp}}{\mathcal{B}^{\rm SM}} - 1, \tag{31}$$

$$\tilde{S}_f = \frac{-\eta_f S_f - \sin(2\beta_s^{\rm eff})}{\cos(2\beta_s^{\rm eff})}. \tag{32}$$

In Fig. 1 we plot the direct $CP$ asymmetry $C_f$ versus the mixing-induced $CP$ asymmetry for several values of the strong phase $\theta_{\rm NP}$, using the relation (28). Once for a given channel $f$ there will be data available for $(S_f, C_f)$ one could find a range for $\theta_{\rm NP}$. The two panels in Fig. 1 are obtained assuming different values for the $B_s - \bar{B}_s$ mixing phase, which is fixed to $2\beta_s^{\rm eff} = 0.77 \pm 0.37$ rad (one of the values obtained by HFAG averaging the experimental results provided by the Tevatron Collaborations CDF and D0) in the left panel, while in the right panel it is fixed to the SM value $\beta_s = 0.017$ rad (no errors are attached to the SM value, hence the various regions shrink to lines). In these figures only positive values of $\theta_{\rm NP}$ have been considered, since $C_f(-\theta) = C_f(\frac{\pi}{2} + \theta) = -C_f(\theta)$.

Equation (29) allows to determine the weak phase $\delta_{\rm NP}$ using the measured $\mathcal{B}(B_s \to f)$ and $S_f$. To appreciate how this could be obtained, we consider the final state $J/\psi\eta$, with $\eta_f = 1$, and plot in Fig. 2 $\delta_{\rm NP}$ versus $S_{J/\psi\eta}$ for $\Sigma$ corresponding to the values $\mathcal{B}^{\rm exp} = 3.32 \pm 1.02$ (see Table III) and $\mathcal{B}^{\rm SM} = 4.25 \pm 0.28$ (the average of the (CDSS) and (BZ) predictions in Table III).

Last, we consider Eq. (30). In this case, all of the three observables $\mathcal{B}(B_s \to f)$, $S_f$ and $C_f$ are required to constrain $R$. Moreover, once experimental information is available about the three observables $\mathcal{B}(B_s \to f)$, $S_f$ and $C_f$ in at least two decay modes, also $\beta_s$ can be constrained, since only the NP parameters $R$ and $\theta_{\rm NP}$ are channel-specific, while $\delta_{\rm NP}$ is the same for all the modes induced by the weak transition $b \to c\bar{c}s$. Hence, measuring in two channels the six observables $(\mathcal{B}(B_s \to f_1), S_{f_1}, C_{f_1})$ and $(\mathcal{B}(B_s \to f_2), S_{f_2}, C_{f_2})$ it would be possible to determine $R_1$, $\theta_{\rm NP,1}$, $R_2$, $\theta_{\rm NP,2}$, $\delta_{\rm NP}$ and $\beta_s$.

A final remark concerns the role of double Cabibbo suppressed SM contributions to the $B_s$ processes of the type discussed here. In SM, taking into account also such contributions, the expression for the $\bar{B}_s \to M_{c\bar{c}}L$ decay amplitude can be written as [40]:

$$A^{\rm SM}(\bar{B}_s \to M_{c\bar{c}}L) = \mathcal{A}^{\rm SM}\left(1 - \frac{\lambda^2}{2}\right)[1 + \epsilon a^{\rm SM} e^{i\theta} e^{-i\gamma}] \tag{33}$$

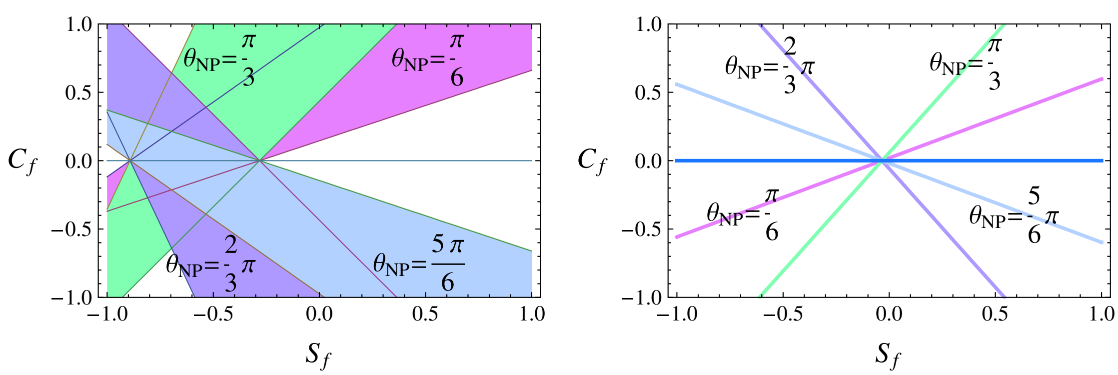


FIG. 1 (color online). Direct $CP$ asymmetry $C_f$ versus the mixing-induced $CP$ asymmetry $S_f$ for several values of the strong phase $\theta_{\rm NP}$.

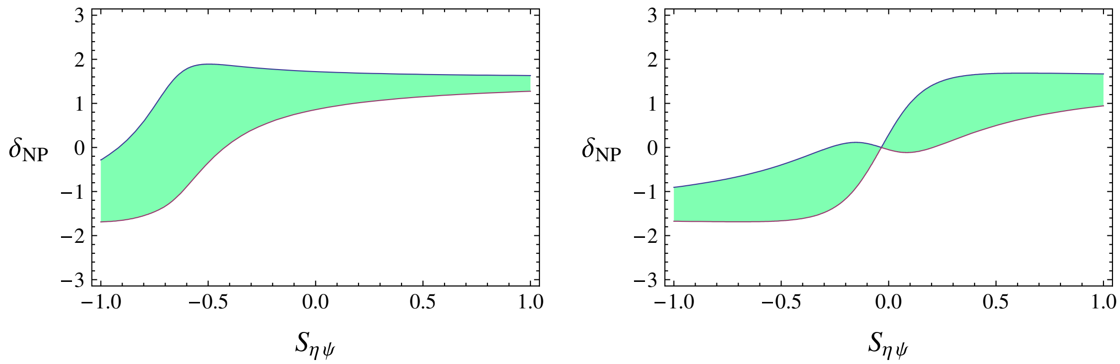


FIG. 2 (color online). Weak phase $\delta_{\rm NP}$ versus $S_{J/\psi\eta}$ in the case $\theta_{\rm NP}=0$ and for $\beta_s$ fixed to the HFAG value (left) or to the SM value (right).

where $\mathcal{A}^{\rm SM}=A_T+A_P^{(c)}-A_P^{(t)}$ and $a^{\rm SM}e^{i\theta}=\frac{A_P^{(u)}-A_P^{(t)}}{A_T+A_P^{(c)}-A_P^{(t)}}$ are combinations of the tree operator (3) and of the penguin operators with internal quark ($q$) in (4). The amplitude involves parameters of the CKM matrix, such as $\lambda$, the sine of the Cabibbo angle, and $\gamma$, the argument of $V_{ub}^*$, together with the strong phase $\theta$. $\epsilon$ represents the combination $\epsilon=\frac{\lambda^2}{1-\frac{\lambda^2}{2}}\simeq 0.053$, and its small value, together with the small ratio of the Wilson coefficients of penguin and tree operators, is the argument to justify the neglect of the term involving $a^{\rm SM}$. The reliability of this approximation has been tested several times, with the conclusion that penguin corrections to both direct and indirect $CP$ induced asymmetries are, at most, of $\mathcal{O}(10^{-3})$ in the case of the $B_d$ system [41]. The issue has been reconsidered for $B_s$ modes, such as $B_s\to J/\psi\phi$ [40], as well as in $B_d\to J/\psi K_{S,L}$ [42], and it has been suggested to use experimental data to get information on such contributions using modes in which they are not Cabibbo suppressed, namely $B_d\to J/\psi K^{*0}(K^+\pi^-)$ for $B_s\to J/\psi\phi$, and $B_d\to J/\psi\pi^0$ for $B_d\to J/\psi K_{S,L}$, employing U-spin or $SU(3)_F$ symmetries.

Double Cabibbo suppression also characterizes the SM penguin pollution in all the processes we have considered in our study, and the idea of using control modes to bound their sizes can be applied extensively. This is the case, first, of processes involving $\psi(2S)$, namely $B_s\to\psi(2S)\phi$, the control process of which can be $B_d\to\psi(2S)K^{*0}(K^+\pi^-)$ accessible at the B factories. Analogously, under $SU(3)_F$ symmetry the $B_s\to J/\psi(\psi(2S))\eta(\eta')$ modes have $B_d\to J/\psi(\psi(2S))\pi^0$ as control processes, from which some information has already been gained. As for $B_s\to\chi_{c0}\phi$, a control mode could be $B_d\to\chi_{c0}K^{*0}$. Detailed analyses would be possible with the availability of experimental measurements; the variety of possibilities makes the set of decay modes discussed in this study of great experimental and theoretical interest.

## IV. CONCLUSIONS

Recent results in the $B_s$ sector require efforts to identify promising ways to unveil new physics. We have considered decay channels induced by the $b\to c\bar{c}s$ transition, using generalized factorization together with $SU(3)_F$ symmetry to predict their branching fractions in the standard model. Modes with a charmonium state plus $\eta$, $\eta'$, $f_0(980)$ are interesting, since they are $CP$ eigenstates and do not require angular analyses. The case of $f_0$ is particularly suitable in view of its reconstruction in the $\pi^+\pi^-$ mode.

If NP affects the $B_s$ sector, it can either contribute to the $\Delta B=2$ induced mixing amplitude, in a channel independent way, modifying the value of the mixing phase with respect to its SM value, or modify the decay amplitudes, in a way that can vary from one channel to the other. In this case, the mixing-induced $CP$ asymmetry $S_f$ would no more be equal to $-\eta_f\sin 2\beta_s$ and the direct $CP$ asymmetry $C_f$ would differ from zero, conditions which instead hold in the SM or in the case that NP affects only the oscillation process. Both cases can also happen. We have considered a general parametrization of new physics in which the second condition holds, fixing $\beta_s$ either to the average of CDF and D0 results or to its SM value. With data on the branching ratio and the two $CP$ asymmetries $S_f$ and $C_f$ for a given final state $f$ it will be possible to constrain NP parameters.


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[1] Updated results can be found at: http://www.utfit.org/UTfit/ and http://ckmfitter.in2p3.fr/.
[2] A. Abulencia *et al.* (CDF Collaboration), Phys. Rev. Lett. **97**, 242003 (2006).
[3] V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. Lett. **98**, 121801 (2007).
[4] V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. Lett. **101**, 241801 (2008).
[5] V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. D **76**, 057101 (2007); T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. Lett. **100**, 121803 (2008); **100**, 161802 (2008).
[6] E. Barberio *et al.* (Heavy Flavor Averaging Group), arXiv:0808.1297.
[7] L. Oakes (on behalf of the CDF Collaboration), *FPCP 2010, Torino, Italy, 2010*.
[8] V. M. Abazov *et al.* (The D0 Collaboration), Phys. Rev. D **82**, 032001 (2010); Phys. Rev. Lett. **105**, 081801 (2010).
[9] V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. D **74**, 092001 (2006).
[10] A. Lenz *et al.*, Phys. Rev. D **83**, 036004 (2011).
[11] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[12] K. Nakamura *et al.* (Particle Data Group), J. Phys. G **37**, 075021 (2010).
[13] M. Neubert and B. Stech, Adv. Ser. Dir. High Energy Phys. **15**, 294 (1998).
[14] A. J. Buras and L. Silvestrini, Nucl. Phys. **B548**, 293 (1999).
[15] Y. Y. Keum, H. n. Li, and A. I. Sanda, Phys. Lett. B **504**, 6 (2001); Phys. Rev. D **63**, 054008 (2001); C. D. Lu, K. Ukai, and M. Z. Yang, Phys. Rev. D **63**, 074009 (2001).
[16] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Phys. Rev. Lett. **83**, 1914 (1999); Nucl. Phys. **B591**, 313 (2000).
[17] A. Khodjamirian, Nucl. Phys. **B605**, 558 (2001).
[18] B. Melic, Phys. Rev. D **68**, 034004 (2003); Z. G. Wang, L. Li, and T. Huang, Phys. Rev. D **70**, 074006 (2004).
[19] P. Colangelo and A. Khodjamirian, in *At the Frontier of Particle Physics/Handbook of QCD*, edited by M. Shifman (World Scientific, Singapore, 2001), Vol. 3, pp. 1495–1576.
[20] P. Colangelo, F. De Fazio, P. Santorelli, and E. Scrimieri, Phys. Rev. D **53**, 3672 (1996); **57**, 3186(E) (1998).
[21] P. Ball and R. Zwicky, Phys. Rev. D **71**, 014015 (2005).
[22] P. Ball and R. Zwicky, Phys. Rev. D **71**, 014029 (2005).
[23] P. Colangelo, F. De Fazio, and W. Wang, Phys. Rev. D **81**, 074001 (2010).
[24] T. Feldmann, P. Kroll, and B. Stech, Phys. Rev. D **58**, 114006 (1998); Phys. Lett. B **449**, 339 (1999); T. Feldmann, Int. J. Mod. Phys. A **15**, 159 (2000).
[25] F. Ambrosino *et al.* (KLOE Collaboration), Phys. Lett. B **648**, 267 (2007).
[26] F. De Fazio and M. R. Pennington, J. High Energy Phys. 07 (2000) 051.
[27] I. Adachi *et al.* (Belle Collaboration), arXiv:0912.1434.
[28] R. Aaij *et al.* (LHCb Collaboration), Phys. Lett. B **698**, 115 (2011).
[29] J. Li *et al.* (Belle Collaboration), Phys. Rev. Lett. **106**, 121802 (2011).
[30] M. Bona, J. Garra Tico, E. Grauges Pous *et al.*, arXiv:0709.0451.
[31] See, e.g., K. Anikeev *et al.*, arXiv:hep-ph/0201071.
[32] S. Stone and L. Zhang, Phys. Rev. D **79**, 074024 (2009); arXiv:0909.5442.
[33] O. Leitner, J. P. Dedonder, B. Loiseau, and B. El-Bennich, Phys. Rev. D **82**, 076006 (2010).
[34] M. Ablikim *et al.* (BES Collaboration), Phys. Rev. D **70**, 092002 (2004); **72**, 092002 (2005).
[35] CDF note 10404: "Measurement of branching fractions of Bs -> J/psi f0(980) decay at CDF", March 11, 2011, available at http://www-cdf.fnal.gov/physics/new/bottom/110224.blessed-B0Jpsif0/cdf10404.pdf.
[36] P. Colangelo, F. De Fazio, and T. N. Pham, Phys. Lett. B **542**, 71 (2002); Phys. Rev. D **69**, 054023 (2004).
[37] M. Gronau and J. L. Rosner, Phys. Lett. B **669**, 321 (2008).
[38] R. Fleischer and R. Knegjens, arXiv:1012.0839.
[39] P. Colangelo, F. De Fazio, and T. N. Pham, Phys. Lett. B **597**, 291 (2004).
[40] S. Faller, R. Fleischer, and T. Mannel, Phys. Rev. D **79**, 014005 (2009).
[41] See, e.g., H. Boos, T. Mannel, and J. Reuter, Phys. Rev. D **70**, 036006 (2004); M. Ciuchini, M. Pierini, and L. Silvestrini, Phys. Rev. Lett. **95**, 221804 (2005); H. n. Li and S. Mishima, J. High Energy Phys. 03 (2007) 009; M. Gronau and J. L. Rosner, Phys. Lett. B **672**, 349 (2009).
[42] S. Faller, M. Jung, R. Fleischer, and T. Mannel, Phys. Rev. D **79**, 014030 (2009).