# Intrinsic charm of light mesons and *CP* violation in heavy quark decay

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We investigate the impact of intrinsic heavy quark states on predicted values of $CP$ asymmetries in decays of heavy mesons. It is shown that the intrinsic charm contribution, although dynamically suppressed in QCD, is favored by the weak interaction, and therefore it can significantly dilute the predicted values of $CP$-violating asymmetries. This introduces an additional nonperturbative uncertainty into the estimate of direct $CP$-violating effects. We provide a phenomenological estimate of the intrinsic charm content of $\eta$ and $\eta'$ mesons by expanding various amplitudes in terms of the heavy-light quark mixing angle and discuss theoretical uncertainties in the estimates of direct $CP$-violating asymmetries in $B\to\eta^{(\prime)}K^{(*)}$. [S0556-2821(98)03815-6]

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## I. INTRODUCTION

Heavy quark decays offer a wide array of methods for testing the standard model and searches for the signatures of new physics. In particular, they proved to be a powerful tool in the studies of the weak mixing matrix and prepare a fertile ground for the exploration of the smallest Cabibbo-Kobayashi-Maskawa (CKM) matrix elements $V_{ub}$ and $V_{td}$ and, most importantly, studies of $CP$ violation. The observation of signatures of direct $CP$ violation involves comparisons of the partial decay rates of charged mesons. As a result, the necessary requirement is a possibility for two distinct pathways (i.e., amplitudes with different weak and strong phases) to reach a given final state. It is realized, for instance, in the $B$ decays to the final states not containing charmed quarks. In this example, $CP$ violation can occur either from the interference of the tree-level and penguin amplitudes, or from the interference of the penguin amplitudes with different quark flavors in the loop. The strong phases are generated by allowing internal quarks in the penguin loop to go on their mass shells by virtue of the so-called Bander-Silverman-Soni (BSS) mechanism [1]. It is expected that the resulting strong phases for different amplitudes are different, in the first case because tree-level graphs do not produce perturbative strong phases, and in the second case, because of the different mass thresholds for different quark species in the penguin loop.

Let us consider a process governed by the quark $b\to s$ transition to a given final state $f$. In general, there are two distinct ways to reach $f$: by the tree-level $b\to u\bar{u}s$ amplitude or by a penguin $b\to sq\bar{q}$ transition. It is clear that the tree-level amplitude is proportional to the CKM matrix elements $V_{ub}^{*}V_{us}$ and thus Cabibbo suppressed as it scales as $\lambda^4$ in Wolfenstein parametrization. On the other hand, the leading penguin effects, although suppressed by loop factors, scale as $V_{tb}^{*}V_{ts}$, or $\lambda^2$. This makes the tree amplitude comparable in strength with the one-loop penguin diagrams, thus enhancing the interference term and allowing for sizable $CP$-violating asymmetry

$$A_{CP}^{\text{dir}}=\frac{\Gamma_{\bar{B}\to\bar{f}}-\Gamma_{B\to f}}{\Gamma_{\bar{B}\to\bar{f}}+\Gamma_{B\to f}}. \tag{1}$$

These asymmetries have been calculated for a number of final states $f$, such as $B\to K\bar{K},K\pi$, etc.

It is plausible to assume that while the BSS mechanism does represent a way to produce a nonzero $CP$-violating asymmetry, the soft nonperturbative effects might produce larger final state interaction (FSI) phases in the exclusive transitions, thus introducing a nonperturbative uncertainty to the calculation of the asymmetry of Eq. (1). For example, the rescattering of physical *hadrons* produced in the reaction provides an additional source for the FSI phase [2–5].

Here we shall argue that soft FSI contributions do not exhaust the list of possible nonperturbative uncertainties of $A_{CP}^{\text{dir}}$. The Fock state expansion could include a nonvanishing contribution of the heavy quark (e.g., charm) states to the light meson's wave function. Although higher Fock state contributions are dynamically suppressed in QCD, weak transitions of the $b$ quark to the heavy quark states are Cabibbo favored. Therefore, the weak interaction selects intrinsic charm states of the light mesons making their contributions competitive with the direct $b\to u$ transitions to the leading Fock states of the light mesons. We admit that the intrinsic charm content of light mesons is considerably difficult to estimate. However, at least for a particular class of light vector and pseudoscalar mesons this contribution can be phenomenologically accounted for by allowing the ''mixing'' of the heavy mesons with hidden charm with the light mesons bearing the same quantum numbers. For instance, there is a nonvanishing probability for the mixing of $J/\psi$ and $\phi$, $\eta_c$, and $\eta'$, etc. In what follows we consider the possible effects of intrinsic heavy quark states of light mesons on $CP$-violating asymmetries.

The paper is organized as follows. In Sec. II we consider an upper bound on the value of the heavy-light quark mixing angle. In Sec, III we discuss how the intrinsic charm content of the light mesons affects direct $CP$-violating asymmetries concentrating on the phenomenologically interesting transitions $B\to\eta^{(\prime)}K^{(*)}$. We summarize our results in Sec. IV.

## II. HEAVY-LIGHT QUARK MIXING

There are many possible approaches in QCD that account for the intrinsic heavy quark states in light quark systems. It is therefore reasonable to employ a phenomenological de-scription of the mixing, extracting the values of the mixing angles from the experiment. These values can be later compared to the results obtained using various models. Let us parametrize a mixing of the heavy and light pseudoscalar mesons in terms of the following matrix:

$$\begin{pmatrix} \eta' \\ \eta \\ \eta_c \end{pmatrix} = \begin{pmatrix} a_P - b_P & c_P - d_P & -\sin\alpha_P \cos(\phi_P - \theta_P) \\ -c_P - d_P & a_P + b_P & -\sin\alpha_P \sin(\phi_P - \theta_P) \\ \sin\alpha_P \cos\phi_P & \sin\alpha_P \sin\phi_P & \cos\alpha_P \end{pmatrix} \begin{pmatrix} \eta_0 \\ \eta_8 \\ \eta_{c0} \end{pmatrix}, \quad (2)$$

where $a_P = \cos^2(\alpha_P/2)\cos\theta_P$, $b_P = \sin^2(\alpha_P/2)\cos(2\phi_P - \theta_P)$, $c_P = \cos^2(\alpha_P/2)\sin\theta_P$, and $d_P = \sin^2(\alpha_P/2)\sin(2\phi_P - \theta_P)$. Here $\eta_0$, $\eta_8$, and $\eta_{c0}$ represent the flavor SU(3) singlet, octet, and pure $c\bar{c}$ states, respectively, and $\eta'$, $\eta$, and $\eta_c$ are physical states

$$|\eta_0\rangle = \frac{1}{\sqrt{3}}|u\bar{u} + d\bar{d} + s\bar{s}\rangle,$$

$$|\eta_8\rangle = \frac{1}{\sqrt{6}}|u\bar{u} + d\bar{d} - 2s\bar{s}\rangle, \quad (3)$$

$$|\eta_{c0}\rangle = |c\bar{c}\rangle.$$

The mixing angles $\alpha_P$, $\theta_P$, and $\phi_P$ have a simple physical meaning [6]: $\alpha_P$ represents an ‘‘admixture’’ of the heavy quarks to the light ones, $\theta_P$ represents a mixing of the light quarks among themselves, and $\phi_P$ gives a light quark SU(3) admixture to a heavy $Q\bar{Q}$ state. It is clear that one recovers a standard $\eta - \eta'$ mixing matrix as $\alpha_P \to 0$ and heavy quarks decouple from the light ones. Sometimes, it is more convenient to work in the quark basis[1]

$$|\eta'\rangle = \frac{1}{\sqrt{2}} X_{\eta'}|u\bar{u} + d\bar{d}\rangle + Y_{\eta'}|s\bar{s}\rangle + Z_{\eta'}|c\bar{c}\rangle,$$

$$|\eta\rangle = \frac{1}{\sqrt{2}} X_{\eta}|u\bar{u} + d\bar{d}\rangle + Y_{\eta}|s\bar{s}\rangle + Z_{\eta}|c\bar{c}\rangle, \quad (4)$$

$$|\eta_c\rangle = \frac{1}{\sqrt{2}} X_{\eta_c}|u\bar{u} + d\bar{d}\rangle + Y_{\eta_c}|s\bar{s}\rangle + Z_{\eta_c}|c\bar{c}\rangle,$$

where we have generalized the construction of [7] to include intrinsic charm states. Using Eq. (2) it is easy to show that

$$X_{\eta_i}^2 + Y_{\eta_i}^2 + Z_{\eta_i}^2 = 1 \quad (5)$$

for each $\eta_i = \{\eta, \eta', \eta_c\}$. This equation might be violated by the presence of some other pseudoscalars that might mix with $\eta_i$ as well. This would result in the limitation of the described method to the prediction of the upper bound on the value of $\eta_c$ contribution to the mixing angle $\alpha_P$.

The construction of Eq. (2) corresponds to the introduction of the additional, charmed singlet current in addition to the ‘‘standard’’ SU(3) singlet and octet currents

$$A_8^\mu = \frac{1}{\sqrt{6}}(\bar{u}\gamma^\mu\gamma_5 u + \bar{d}\gamma^\mu\gamma_5 d - 2\bar{s}\gamma^\mu\gamma_5 s),$$

$$A_0^\mu = \frac{1}{\sqrt{3}}(\bar{u}\gamma^\mu\gamma_5 u + \bar{d}\gamma^\mu\gamma_5 d + \bar{s}\gamma^\mu\gamma_5 s), \quad (6)$$

$$A_c^\mu = \bar{c}\gamma^\mu\gamma_5 c.$$

These induce the following matrix elements:

$$\langle\eta'|\bar{u}\gamma_\mu\gamma_5 u|0\rangle = -i\sqrt{\frac{2}{3}}\left[(a_P - b_P)F_0 + \frac{1}{\sqrt{2}}(c_P - d_P)F_8\right]p_\mu,$$

$$\langle\eta'|\bar{s}\gamma_\mu\gamma_5 s|0\rangle = -i\sqrt{\frac{2}{3}}[(a_P - b_P)F_0 - \sqrt{2}(c_P - d_P)F_8]p_\mu, \quad (7)$$

$$\langle\eta'|\bar{c}\gamma_\mu\gamma_5 c|0\rangle = i\sqrt{2}\sin\alpha_P\cos(\phi_P - \theta_P)F_{\eta_{c0}}p_\mu.$$

Similar matrix elements exist for the $\eta$. Here $F_{0,8}$ are the singlet and octet decay constants and $F_{\eta_{c0}}$ is the $\eta_{c0}$ decay constant. In the limit of SU(3) symmetry $F_8 = F_\pi$. The SU(3)-violating corrections have been calculated in [8] in the framework of chiral perturbation theory and found to modify this relation by approximately 25%, i.e., $F_8/F_\pi \equiv 1/x_8 = 1.25$. We shall use this value in the following analysis. The value of $x_0$, on the other hand, is not fixed by the SU(3) symmetry arguments, so we shall keep it as a free parameter fixing it later by fitting to the experimental data (in the limit of the ‘‘nonet’’ symmetry $F_0/F_\pi \equiv 1/x_0 = 1$).

The parameters of the mixing matrix can be obtained phenomenologically, as they contribute to the decays of charmonia to the light $\eta, \eta'$, and $\pi^0$ mesons and to radiative decays of the light mesons. Here our essential assumption is that the mixing angles $\alpha_P$ and $\phi_P$ are sufficiently small. This assumption, however, is rather loose, and is valid even for a relatively large charm content of $\eta'$ [9,10], but it allows us

[1]Note that we are *not* using any specific quark model.

to carry out a perturbative expansion in these angles. In the following analysis we will only keep terms linear in $\alpha_P$ and $\phi_P$.

The bulk of information about the relevant mixing angles comes from the radiative decays of $\eta$, $\eta'$, and $\eta_c$ mesons. Normalizing the decay widths to the width of $\pi^0\to\gamma\gamma$ and using Eq. (2) we obtain

$$\rho_{\eta'}\equiv\frac{3}{8}\left(\frac{m_\pi}{m_{\eta'}}\right)^3\frac{\Gamma(\eta'\to\gamma\gamma)}{\Gamma(\pi^0\to\gamma\gamma)}$$

$$=\left[\frac{F_\pi}{F_0}[a_P-b_P]+\frac{F_\pi}{F_8}\frac{c_P-d_P}{\sqrt{8}}-\sqrt{\frac{3}{8}}\frac{CF_\pi F_{\eta_{c0}}}{m_{\eta_c}^2}\sin\alpha_P\cos(\phi_P-\theta_P)\right]^2,$$

$$\rho_\eta\equiv 3\left(\frac{m_\pi}{m_\eta}\right)^3\frac{\Gamma(\eta\to\gamma\gamma)}{\Gamma(\pi^0\to\gamma\gamma)}$$

$$=\left[\frac{F_\pi}{F_0}[a_P+b_P]-\frac{F_\pi}{F_8}\sqrt{8}[c_P+d_P]-\frac{1}{\sqrt{3}}\frac{CF_\pi F_{\eta_{c0}}}{m_{\eta_c}^2}\sin\alpha_P\sin(\phi_P-\theta_P)\right]^2, \tag{8}$$

$$\rho_{\eta_c}\equiv\frac{3}{8}\left(\frac{m_\pi}{m_{\eta_c}}\right)^3\frac{\Gamma(\eta_c\to\gamma\gamma)}{\Gamma(\pi^0\to\gamma\gamma)}$$

$$=\left[\frac{F_\pi}{F_0}\sin\alpha_P\cos\phi_P+\frac{F_\pi}{F_8}\frac{\sin\alpha_P\sin\phi_P}{\sqrt{8}}+\sqrt{\frac{3}{8}}\frac{CF_\pi F_{\eta_{c0}}}{m_{\eta_c}^2}\cos\alpha_P\right]^2,$$

where $a_P$, $b_P$, and $c_P$ were defined in Eq. (2) and

$$C=\sqrt{2}\left(\frac{8\pi}{3}\right)^2. \tag{9}$$

Expanding in terms of $\alpha_P$ and $\phi_P$ and keeping only the linear part we arrive at

$$\rho_{\eta'}\simeq\left[\frac{x_8}{\sqrt{8}}\sin\theta_P+\cos\theta_P\left(x_0-\sqrt{\frac{3}{8}}\frac{CF_\pi F_{\eta_{c0}}}{m_{\eta_c}^2}\alpha_P\right)\right]^2,$$

$$\rho_\eta\simeq\left[x_8\cos\theta_P-\sin\theta_P\left(\sqrt{8}x_0-\sqrt{\frac{1}{3}}\frac{CF_\pi F_{\eta_{c0}}}{m_{\eta_c}^2}\alpha_P\right)\right]^2, \tag{10}$$

$$\rho_{\eta_c}\simeq\left[x_0\alpha_P+\sqrt{\frac{3}{8}}\frac{CF_\pi F_{\eta_{c0}}}{m_{\eta_c}^2}\right]^2.$$

As it is seen from Eq. (10), the dependence on $\phi_P$ drops out at this order, so we set $\phi_P=0$ in what follows. In time, when the accuracy of experimental measurements improves, the second order in the ''angle expansion'' should constrain the value of $\phi_P$ as well.

It is clear from Eq. (10) that three equations do not constrain four parameters: $\theta_P$, $\alpha_P$, $x_0$, and $F_{\eta_{c0}}$. There are several ways to proceed at this point. For instance, one can fix $x_0=1$ by assuming a nonet symmetry and fit the rest of the parameters from the Eq. (10). Here we shall take a different approach. The ''angle expansion'' can be carried out for other processes involving $\eta$ and $\eta'$, such as $J/\psi\to\eta^{(\prime)}\gamma$ and $J/\psi\to\eta_c\gamma$, or $\eta'\to\rho\gamma$. In the limit of SU(3) invariance, only the singlet state $\eta_0$ or $\eta_{c0}$ can be coupled to $J/\psi$. The amplitudes of the radiative decay of a vector charmonium state into a final state containing the SU(3) singlet light quark state $|\eta_0\rangle$ or $|\eta_{c0}\rangle$ can be written as

$$A(J/\psi\to\eta_0\gamma)=A\epsilon_{\mu\nu\alpha\beta}\epsilon_\psi^\mu\epsilon_\gamma^\nu p_\psi^\alpha k_\gamma^\beta,$$

$$A(J/\psi\to\eta_{c0}\gamma)=B\epsilon_{\mu\nu\alpha\beta}\epsilon_\psi^\mu\epsilon_\gamma^\nu p_\psi^\alpha k_\gamma^\beta. \tag{11}$$

Performing the ''angle expansion'' we find that[2]

$$\frac{\Gamma(J/\psi\to\eta_c\gamma)}{\Gamma(J/\psi\to\eta'\gamma)}=\frac{1}{\cos^2\theta_P}\left[\frac{p_{\eta_c}}{p_{\eta'}}\right]^3\left[\frac{1+\alpha_P(A/B)}{\alpha_P-(A/B)}\right]^2,$$

$$\frac{\Gamma(J/\psi\to\eta_c\gamma)}{\Gamma(J/\psi\to\eta\gamma)}=\frac{1}{\sin^2\theta_P}\left[\frac{p_{\eta_c}}{p_\eta}\right]^3\left[\frac{1+\alpha_P(A/B)}{\alpha_P-(A/B)}\right]^2, \tag{12}$$

$$\frac{\Gamma(J/\psi\to\eta'\gamma)}{\Gamma(J/\psi\to\eta\gamma)}=\frac{1}{\tan^2\theta_P}\left[\frac{p_{\eta'}}{p_\eta}\right]^3.$$

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[2]This result agrees with the result of [11] in the limit $A/B\ll\alpha_P$. Please note that it is not possible to extract $\alpha_P$ from this decay mode *alone* without invoking additional dynamical arguments about the size of $A/B$ (cf. [11]).

As seen from Eq. (12), the *ratio* of radiative decay widths of charmonia into $\eta'$ and $\eta$ is independent of the mixing angle $\alpha_P$ *at this order* and can be used to extract the value of $\theta_P$. Alternatively, the ratio

$$\frac{\Gamma(\eta'\to\rho\gamma)}{3\Gamma(\omega\to\pi\gamma)}\left[\frac{p_\pi}{p_\rho}\right]^3 \simeq X_{\eta'}^2 \sim \frac{1}{3}[\sqrt{2}\cos\theta_P+\sin\theta_P]^2 \tag{13}$$

does not depend on the heavy-light mixing angle and can be used to extract $\theta_P$. A similar analysis is possible for $\rho\to\eta\gamma$, etc. Extracting $\theta_P$ from either decay mode and feeding it into Eq. (10) leads to the constraints on all of the mixing parameters. This calculation finds $\theta_P\approx-20°$, $x_0\approx 0.92$ (which corresponds to $F_0/F_\pi\approx 1.05$, in full accord with [7,8]). The situation is more complicated with respect to the heavy-light mixing angles. Uncertainties in the extractions of the light-quark mixing angle $\theta_P$ and decay rates complicate the extraction of the value of the mixing angle $\alpha_P$. In fact, all of the experimental results can be successfully fit assuming a zero value for the mixing angle $\alpha_P$. In order to overcome this problem, additional constraints on the parameters in Eq. (10) must be imposed. For instance, the value of $F_{\eta_{c0}}$ can be extracted from other decay modes that are less sensitive to the heavy-light mixing, e.g., $B\to\eta_c X_s$, where the $c\bar{c}$ component is enhanced by the CKM matrix element $V_{cb}$. Using this method, $F_{\eta_{c0}}$ can be fixed to $F_{\eta_{c0}}\simeq 0.29$ GeV. This and the third line of Eq. (10) implies that $|\sin\alpha_P|\leqslant 0.03$ $(1.7°)$, $\phi_P=0$. It is interesting to compare this bound to what already exists in the literature. It appears that all the models of intrinsic charm based on heavy-light meson mixing [11,12] (see also [13]) satisfy this bound. On the other hand, there exists a class of operator-product-expansion-(OPE)-based calculations [9,10] that clearly violates this bound, predicting $\alpha_P\simeq 7°$ if the expansion is terminated at the level $1/m_c^2$. As we shall see later, these values of $\alpha_P$ significantly dilute the direct $CP$ asymmetries in $B$ decays.[3]

A similar construction is available for the vector mesons, where we use the same notations with the obvious replacement of the subscript $P$ by $V$,

$$\begin{pmatrix}\phi\\ \omega\\ J/\psi\end{pmatrix}=\begin{pmatrix} a_V-b_V & c_V-d_V & -\sin\alpha_V\cos(\phi_V-\theta_V)\\ -c_V-d_V & a_V+b_V & -\sin\alpha_V\sin(\phi_V-\theta_V)\\ \sin\alpha_V\cos\phi_V & \sin\alpha_V\sin\phi_V & \cos\alpha_V\end{pmatrix}\begin{pmatrix}\omega_0\\ \omega_8\\ \psi_0\end{pmatrix}. \tag{14}$$

Numerically, $\sin\alpha_V\cos(\phi_V-\theta_V)\simeq 1.2\times 10^{-4}$ [6].

## III. AMPLITUDES AND *CP*-VIOLATING ASYMMETRIES

In $B$ decays, the tree-level amplitude is suppressed by a small $V_{ub}$, which makes it comparable to the one-loop penguin amplitude. On the other hand, the transition $B\to[Q\bar{Q}]s\to Ms$ with $Q$ being a heavy (charm) quark and $M$ a light final state meson is not CKM suppressed. Therefore the latter mechanism becomes competitive with the direct $b\to u$ transition to the light quarks constituting the light mesons. Since the $b\to c\bar{c}s$ amplitude does not contain a $CP$-violating weak phase, it may significantly reduce the predicted value of $CP$ asymmetry. In what follows we consider two cases for $M$ being a pseudoscalar and a vector meson.

The generic amplitude for the decays of a $B$ meson to the charmless final state $f$ can be written as

$$\begin{aligned}A_{B\to f}&=\xi_u A_T+\xi_c A_M+\sum_{i=u,c,t}\xi_i A_P^i\\ &=\xi_u[A_T+A_P^{ut}]+\xi_c[A_P^{ct}+A_M].\end{aligned} \tag{15}$$

Here $A_T$ is a tree-level amplitude, $A_M$ is a ‘‘mixing’’ amplitude, and $A_P$ is a penguin amplitude. As usual, unitarity of the CKM matrix was used to write $\xi_t=-\xi_c-\xi_u$, $A_P^{ct}\equiv A_P^c-A_P^t$, $A_P^{ut}\equiv A_P^u-A_P^t$, and $\xi_i=V_{ib}^*V_{is}$.

In order to form a $CP$-violating asymmetry, we also must consider the corresponding amplitude for the decay of the $\bar{B}$:

$$\bar{A}_{\bar{B}\to\bar{f}}=\xi_u^*[\bar{A}_T+\bar{A}_P^{ut}]+\xi_c^*[\bar{A}_P^{ct}+\bar{A}_M]. \tag{16}$$

Using the fact that $\bar{A}_T=A_T$ and $\bar{A}_M=A_M$, the asymmetry of Eq. (1) can be formed as

$$\begin{aligned}\Delta_f&=\Gamma_{\bar{B}\to\bar{f}}-\Gamma_{B\to f},\\ \Delta_f&=\lambda_f\,\mathrm{Im}\,\xi_u^*\xi_c\,\mathrm{Im}[A_T+A_P^{ut*}][A_P^{ct}+A_M],\end{aligned} \tag{17}$$

where $\lambda_f=\sqrt{[1-(x_{f_1}+x_{f_2})^2][1-(x_{f_1}-x_{f_2})^2]}/(4\pi m_B)$ is a phase space factor, divided by the sum of the corresponding decay rates. As seen from Eq. (17), direct $CP$ violation in the $B\to\eta_i K$ mode arises not only from the interference of

[3]The OPE-based calculations successfully explain the unexpectedly large branching ratio for $B\to\eta' K$. They however have certain phenomenological difficulties in the case of the inclusive $\eta'$ production in $B$ decays, as well as describing ratios of branching fractions of $B$ mesons decaying to $PP$ and $PV$ final states [11,26,27]. The fact that the mass of the charmed quark is not sufficiently large for the fast convergence of the OPE might explain the variance in the results of these calculations [14].

the penguin diagrams with the different internal quark flavors, but also, because of the complicated quark content of the $\eta_i$, from the interference of the Cabibbo-suppressed tree-level amplitudes with the penguin amplitudes. As we shall see later, the intrinsic charm contribution affects $CP$-violating asymmetries calculated using perturbative BSS phases. It is therefore instructive to study the dependence of the asymmetry on the parameter $\hat{q}=q^2/m_B^2$, which parametrizes the momentum flowing through the penguin vertex [16]. The final answer, of course, is independent of $\hat{q}$ and should be obtained by either integrating the asymmetry with respect to $\hat{q}$ smeared by some function defined by the momentum distribution of quarks in the final state mesons, or by fixing $\hat{q}$ using quark model arguments. We shall use the second method for our predictions, presenting the graphs asymmetry vs $\hat{q}$ to show the threshold structure of $CP$-violating asymmetries. We comment on the effects of soft FSI phases in the Conclusion.

In the following discussion we shall first use the effective Hamiltonian calculated at the leading order in QCD, i.e., with no QCD corrections associated with the penguin part. Next, the full next-to-leading order effective Hamiltonian is employed.

It is well known that the calculation of the two-body nonleptonic decays of heavy mesons cannot be performed without invoking a particular model. This model dependence is partially cancelled in the asymmetry Eq. (1). In our calculation we choose the factorization approximation [17–21] and the Bauer, Stech, and Wirbel (BSW) model [15] to estimate relevant form factors.

### A. Leading order calculations

The ‘‘no QCD Hamiltonian’’ reads

$$\mathcal{H}_{\rm eff}^{\rm LO}=\frac{4G_F}{\sqrt{2}}\left\{\xi_Q\sum_{i=1}^{2} C_i O_i^Q-\frac{\alpha_s}{8\pi}\left[\sum_{i=u,c,t}\xi_i F_i\right]\times\left[-\frac{O_3}{N_c}+O_4-\frac{O_5}{N_c}+O_6\right]\right\}, \quad (18)$$

$$O_1^q=\bar{s}_\alpha\gamma_\mu L Q_\beta \bar{Q}_\beta\gamma^\mu L b_\alpha, \quad O_2^q=\bar{s}\gamma_\mu L Q\bar{Q}\gamma^\mu L b,$$

$$O_{3(5)}=\bar{s}\gamma_\mu L b\sum_q \bar{q}\gamma^\mu L(R)q,$$

$$O_{4(6)}=\bar{s}_\alpha\gamma_\mu L b_\beta\sum_q \bar{q}_\beta\gamma^\mu L(R)q_\alpha .$$

Here, $Q=\{u,c\}$, $q=\{u,d,s\}$, and $F_i$ are the Inami-Lim functions for the flavor $i$. A similar construction is available for the effective $b\to d$ transitions (although the effects of intrinsic charm states are largely suppressed in these modes since $b\to u\bar{u}d$ decays are not CKM suppressed compared to $b\to c\bar{c}d$). In what follows we drop the contributions from the electroweak penguin operators and dipole operators for the sake of simplicity.

The decay amplitude $A_{\eta^{(\prime)}K}=-i\langle\eta^{(\prime)}K|\mathcal{H}_{\rm eff}|B\rangle$ can be written as

$$A_{B\to\eta^{(\prime)}K}=\xi_u[A_T^{\eta^{(\prime)}K}+F_{ut}A_P^{\eta^{(\prime)}K}]+\xi_c[F_{ct}A_P^{\eta^{(\prime)}K}+A_M^{\eta^{(\prime)}K}],$$

$$A_T^{\eta^{(\prime)}K}=-G_F m_B^2[a_2F^{uu}_{\eta^{(\prime)}}f_+^K(m^2_{\eta^{(\prime)}})L_k(\mu_i)+a_1F_Kf_+^{\eta^{(\prime)}}(0)L_\eta(\mu_i)], \quad (19)$$

$$A_P^{\eta^{(\prime)}K}=G_F m_B^2\frac{\alpha_s}{8\pi}\left(1-\frac{1}{N_c^2}\right)a_P,$$

$$A_M^{\eta^{(\prime)}K}=G_F m_B^2\sin\alpha_P\times\cos(\phi_P-\theta_P)F_{\eta_c}f_+^K(m^2_{\eta^{(\prime)}})L_k(\mu_i),$$

where $a_1=C_2+\chi C_1$, $a_2=C_1+\chi C_2$, $\mu_i=m_i^2/m_B^2$, and the following notations are used:

$$a_P=F_Kf_+^{\eta^{(\prime)}}(0)L_\eta(\mu_i)+F^{ss}_{\eta^{(\prime)}}f_+^K(m^2_{\eta^{(\prime)}})L_k(\mu_i)+2F_Kf_+^{\eta^{(\prime)}}(0)\frac{m_K^2}{m_sm_b}M_\eta(x_i^2)+\overline{F}^{ss}_{\eta^{(\prime)}}f_+^K(m^2_{\eta^{(\prime)}})\frac{m^2_{\eta^{(\prime)}}}{m_sm_b}M_k(\mu_i). \quad (20)$$

Our choice of the form factors and kinematic parameters is explained in the Appendix. In the standard factorization approach $\chi=1/N_c$ and all the octet-octet nonfactorizable corrections are neglected. These corrections can be accounted for phenomenologically by treating $\chi$ as a free parameter and fitting it to the available data assuming universality of these corrections for different final states [22]. In this approach, $a_2\simeq 0.25$ [23,24].

In order to maintain gauge invariance and unitarity, the calculation must be performed to order $\alpha_s^2$ in perturbative QCD (PQCD) expansion [25]. The $CP$-violating asymmetry reads

$$\Delta_{\eta^{(\prime)}K}=\lambda_{\eta^{(\prime)}K}{\rm Im}\,\xi_u^*\xi_c\Bigg((A_P^{\eta^{(\prime)}K})^2\times\left\{{\rm Im}\,F_{ut}^*{\rm Re}\,F_{ct}+{\rm Re}\,F_{ut}^*\left[{\rm Im}\,F_{ct}-\frac{n_F}{6}{\rm Re}\,F_{ct}\right]\right\}+A_M^{\eta^{(\prime)}K}A_P^{\eta^{(\prime)}K}{\rm Im}\,F_{ut}^*+A_P^{\eta^{(\prime)}K}A_T^{\eta^{(\prime)}K}\times\left[{\rm Im}\,F_{ct}-\frac{n_F}{6}{\rm Re}\,F_{ct}\right]\Bigg). \quad (21)$$

Dividing $\Delta_{\eta' K}$ by the sum of the decay rates, the asymmetry Eq. (1) can be calculated. As it is seen from Fig. (1), the intrinsic charm reduces the $CP$ asymmetry by approximately 30–50 %, if our estimate of $\alpha_P$ is used thus complicating the extraction of $CP$-violating parameters of the CKM matrix

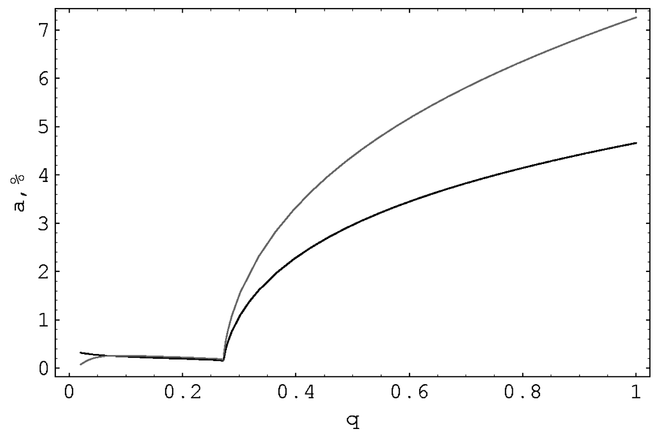


FIG. 1. $CP$ asymmetry calculated at the leading order in QCD as a function of the parameter $q^2/m_B^2$ for $B\to\eta' K$, without intrinsic charm (gray) and with intrinsic charm $\alpha_P=0.03$ (black).

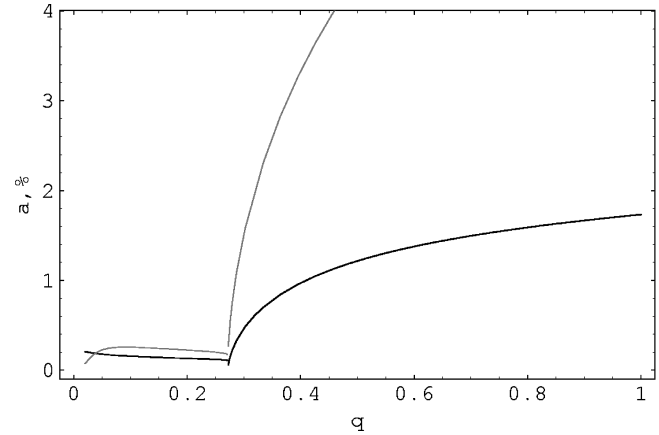


FIG. 2. Same as Fig. 1 but for $\alpha_P\simeq 0.12$.

from the decay modes of this type. Please note, that if by some reason, the intrinsic charm content of the $\eta'$ is increased [9,10], this mode becomes practically useless for the observation of the direct $CP$-violating effects. It is clear from Fig. (2) that $CP$-violating asymmetry is significantly diluted even for $\alpha_P\simeq 7°$, which corresponds to the lower bound of the prediction in OPE-based calculations. This however, simplifies the time-dependent analysis of the decay $B_d^0\to\eta' K_s$ as it suppresses direct $CP$-violating amplitudes. The behavior of $CP$-violating asymmetries for the case of $\eta K$ final states is similar to the case described above. We therefore refrain from displaying the shape of the $A_{CP}^{\rm dir}(\hat{q})$ function, but rather show the numerical value of the $A_{CP}^{\rm dir}$ in Table I with $\hat{q}$ fixed by the quark model arguments. The analysis of the $\eta^{(')}V$ final state is completely similar to the calculation described above. The results are presented in Table I as well.

### B. Next-to-leading order calculations

The next-to-leading order (NLO) QCD effective Hamiltonian reads

$$\mathcal{H}_{\rm eff}^{\rm NLO}=\frac{4G_F}{\sqrt{2}}\left\{\sum_{Q=u,c}\xi_Q\left[\sum_{i=1}^{2}C_iO_i^Q+\sum_{k=3}^{6}C_kO_k\right]\right\},\tag{22}$$

$$O_1^q=\overline{s}_\alpha\gamma_\mu LQ_\beta\overline{Q}_\beta\gamma^\mu Lb_\alpha,\quad O_2^q=\overline{s}\gamma_\mu LQ\overline{Q}\gamma^\mu Lb,$$

$$O_{3(5)}=\overline{s}\gamma_\mu Lb\sum_q\overline{q}\gamma^\mu L(R)q,$$

$$O_{4(6)}=\overline{s}_\alpha\gamma_\mu Lb_\beta\sum_q\overline{q}_\beta\gamma^\mu L(R)q_\alpha.$$

In our calculation we used renormalization scheme independent effective Wilson's coefficients defined as

$$C_{2i+1}^{\text{eff }q}(m_b)=\overline{C}_{2i+1}(m_b)+\frac{\alpha_s(m_b)}{8\pi N_c}\left[G(m_q,q^2,m_b)-\frac{10}{9}\right]\overline{C}_2,$$

$$C_{2i}^{\text{eff }q}(m_b)=\overline{C}_{2i}(m_b)-\frac{\alpha_s(m_b)}{8\pi}\left[G(m_q,q^2,m_b)-\frac{10}{9}\right]\overline{C}_2\tag{23}$$

with $i=1,2$ and Wilson's coefficients $\overline{C}_k$ at $\mu=m_b$ with $\alpha_s(m_Z)=0.118$ are given by

$$\overline{C}_1(m_b)=-0.313,\quad \overline{C}_2(m_b)=1.150,$$
$$\overline{C}_3(m_b)=0.017,\quad \overline{C}_4(m_b)=-0.037,\tag{24}$$
$$\overline{C}_5(m_b)=0.010,\quad \overline{C}_6(m_b)=-0.046.$$

TABLE I. Absolute values of $CP$-violating asymmetries $A_{CP}^{\rm dir}$ for $\hat{q}=q^2/m_b^2\simeq 0.5$, $\rho=0.05$, $\eta=0.36$, and $m_c=1.6-1.3$ GeV.

| Mode | $A_{CP}^{\rm dir}(\alpha_P=0.0)$,% | $A_{CP}^{\rm dir}(\alpha_P=0.03)$,% | $A_{CP}^{\rm dir}(\alpha_P=0.12)$,% |
|---|---|---|---|
| $B^-\to\eta' K^-$, LO | 1.9–3.5 | 1.6–2.7 | 1.0–1.5 |
| $B^-\to\eta' K^-$, NLO | 2.5–4.4 | 1.7–2.9 | 0.7–1.2 |
| $B^-\to\eta K^-$, LO | 3.1–7.3 | 2.7–6.1 | 1.9–3.9 |
| $B^-\to\eta K^-$, NLO | 3.0–7.1 | 2.2–5.4 | 0.8–2.6 |
| $B^-\to\eta' K^{*-}$, LO | 2.9–6.7 | 1.9–4.1 | 0.9–1.6 |
| $B^-\to\eta' K^{*-}$, NLO | 7.6–16.5 | 3.3–7.5 | 0.3–1.3 |
| $B^-\to\eta K^{*-}$, LO | 6.7–19.8 | 4.5–12.5 | 1.9–4.7 |
| $B^-\to\eta K^{*-}$, NLO | 15.4–38.3 | 6.3–18.3 | 1.4–1.9 |

The function $G(m_q, q^2, \mu)$ is given by

$$G(m_q,q^2,\mu)=-4\int dxx(1-x)\log\frac{m_q^2-x(1-x)q^2}{\mu^2} \tag{25}$$

where $q$ is a momentum of the gluon in the penguin diagram. It is clear that the strong phase is generated every time $q^2>4m_q^2$ for each quark species in the loop.

The factorization calculation is completely similar to the one performed in the previous section. The only difference arises from the fact that the effective constants $a_i=C_{2i-1}+\chi C_{2i}$ are not vanishing and have to be taken into account along with the corresponding form factors. Those are the combinations of the form factors and the decay constants defined in the previous section. The results are presented in Table I. In our calculation we fixed the value of $a_2\simeq 0.25$ fixed by the experimental data while dropping the nonfactorizable contributions to the matrix elements of penguin operators which are difficult to estimate reliably. The asymmetry is seen not to change significantly in going to NLO approximation with the intrinsic charm contribution still diluting it at the level of 30–50% for the estimated value of $\alpha_P\simeq 0.03$. The effect becomes stronger for higher values of mixing angles. As it is seen, even the perturbative result is very uncertain. The major source of uncertainty is by far dominated by the value of the charmed quark mass that affects the position of the charmed quark threshold. Because of the SU(3) symmetry relations, the values of the hadronic form factors, although important for the decay width predictions, are not seen to significantly affect the predicted values of $CP$ asymmetries. An additional source of uncertainty is the value of nonfactorizable corrections usually summarized in the effective parameter $\chi$. In the case of the asymmetries induced by the penguin-tree interference, it can change the balance of these contributions, thus shifting the value of $A_{CP}^{\rm dir}$. It is usually assumed that the nonfactorizable corrections can be taken into account by replacing the factors of $1/N_c$ in $a_i$ throughout the calculation with the *single* parameter $\chi$. This fact also induces the uncertainty into the estimate of both decay width and $CP$-violating asymmetry.[4] It is interesting to note that in the limit $\chi\to 0$ the contribution from the intrinsic charm amplitude becomes *negative* actually reducing the predicted values of branching fractions: for instance, $B(B^-\to\eta' K^-)=2.8\times 10^{-5}$ for moderate values of the form factors and $\alpha_P\simeq 0.3$ drops to $B(B^-\to\eta' K^-)=1.6\times 10^{-5}$ if $\chi=0$. Of course, higher values for the branching fraction are still possible considering large uncertainties in the values of hadronic form factors [26].

[4]The contribution of the octet matrix element is usually associated with the quark final state rescatterings into the colorless hadron states. It is clear that the interactions of the rescattered and spectator quarks might introduce different $\chi$'s for different amplitudes. For instance, there would be different contributions from $b\to u\overline{u}s$ and $b\to d\overline{d}s$ octet amplitudes to $B^-\to\eta^{(\prime)}K^{(*)}$. The effect is similar to the effect of Pauli interference in nuclear physics.

Similar calculations can be performed for the decays involving vector particles in the final state, such as $B\to\phi K$. The calculation is simplified considerably since the decay is dominated by a single penguin amplitude. The $CP$ asymmetry for the leading order case is

$$\Delta_{\phi K}=\lambda_{\phi K}{\rm Im}\ \xi_u^*\xi_c[A_P^{\phi K2}\{{\rm Im}\ F_{ut}^*{\rm Re}\ F_{ct}+{\rm Re}\ F_{ut}^*{\rm Im}\ F_{ct}\} +A_M^{\phi K}A_P^{\phi K}{\rm Im}\ F_{ut}^*]. \tag{26}$$

Here, the following notations are used:

$$A_P^{\phi K}=-G_F m_\phi F_\phi(p_B+p_K)\cdot\epsilon f_+^K(m_\phi^2)\frac{\alpha_s}{8\pi}$$

$$A_M^{\phi K}=-G_F m_\phi F_\psi(p_B+p_K)\cdot\epsilon f_+^K(m_\phi^2)a_2 \times\sin\alpha_V\cos(\phi_V-\theta_V). \tag{27}$$

As usual, one expects partial cancellations among the first and the second term in Eq. (26) because of the Glashow-Iliopoulos-Maiani (GIM) mechanism. Thus, the intrinsic charm amplitude is potentially important as it does not suffer from this cancellation. Fortunately, the small value of the heavy-light mixing angle for the vector mesons makes the intrinsic charm contribution extremely small. The correction to the asymmetry $A_{CP}^{\rm dir}\simeq 0.3\%$ is less than 1% for the same choice of CKM parameters as before. We therefore see no point in performing the NLO studies of this class of decay modes. Clearly, final states containing vector particles are much less affected by the higher Fock state "pollution." The resulting $CP$-violating asymmetries, however, are significantly smaller.

## IV. CONCLUSIONS AND OUTLOOK

We have investigated the impact of the intrinsic heavy quark states on the $CP$-violating asymmetries in $B$ decays. It arises because of the fact that the intrinsic charm quark states, although suppressed by QCD dynamics of the process, are less affected by weak Cabibbo suppression. The most dramatic effect occurs in the case of the $\eta' V$ final states. Unfortunately, the impact of the intrinsic charm states on the $CP$ asymmetry is very difficult to test experimentally since the direct $CP$ asymmetry explicitly depends on the strong phase of $S$ matrix elements relating two states produced by the weak interaction. This quantity is notoriously difficult to estimate theoretically since it comprises not only perturbative BSS phases but also soft nonperturbative phases generated by the rescattering of physical hadrons [2]. Soft FSI contributions, although important, are not seen to change the shape of the graphs asymmetry vs parameter $q$ but rather move them up or down, whereas the described mechanism certainly affects the shape. Of course, in the final result the function described by $A_{CP}^{\rm dir}(\hat{q})$ has to be integrated over $\hat{q}$ with a suitable choice of a "smearing function" representing momentum distributions of quarks inside the final hadrons, or $\hat{q}$ has to be fixed by the quark model arguments. In either case, the intrinsic charm contribution is seen to reduce the value for $CP$-violating asymmetry sizably. This effect is not universal, but rather specific to the final states containing $\eta^{(\prime)}$ mesons, so it cannot be attributed to the long-distance part of the penguin diagram.

The mechanism described in this paper, along with the uncertainty in the determination of FSI phases complicates the use of the direct $CP$-violating asymmetries as a so-called ''consistency check'' in the determination of the angles of the CKM triangle. For instance, the numerical values of the angle $\gamma$ determined from various decay modes have to be consistent with each other unless there exists a new physics contribution that violates this requirement. Thus, the possible inconsistency should manifest a nonstandard model mechanism affecting the decay processes. As one can see from the discussion above, there exist possible sources of violation of this consistency check *within* the standard model.

Finally, we would like to note that this mechanism does not markedly affect the decay modes where electroweak (EW) penguin diagrams are manifested (e.g., $B\to\eta'\pi$) since there the tree-level and mixing amplitudes contribute at the same order in Wolfenstein parameter λ with mixing amplitudes additionally suppressed by the small values of heavy-light mixing angles.

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## APPENDIX: HADRONIC MATRIX ELEMENTS AND KINEMATICAL FACTORS

We use the following definitions for the pseudoscalar hadronic form factors:

$$\langle M|\overline{q}\gamma_\mu b|B\rangle = f_+^M(q^2)(p_B+p_M)_\mu + f_-^M(q^2)(p_B-p_M)_\mu ,$$

$$\langle M|\overline{q_1}\gamma_\mu\gamma_5 q_2|0\rangle = -i\sqrt{2}F_M^{q_1q_2}p_{M\mu}. \tag{A1}$$

We take $F_K=0.12$ GeV. Usual relations among decay constants (and transition form factors) of different mesons in the pseudoscalar octet imposed by SU(3)-symmetry are modified in the presence of $\eta-\eta'$ mixing and intrinsic charm components of $\eta$ and $\eta'$, e.g.,

$$F_{\eta'}^{uu} = \frac{1}{\sqrt{3}}\left\{(a_P-b_P)F_0+\frac{1}{\sqrt{2}}(c_P-d_P)F_8\right\},$$

$$F_{\eta'}^{ss} = \frac{1}{\sqrt{3}}\{(a_P-b_P)F_0-\sqrt{2}(c_P-d_P)F_8\}, \tag{A2}$$

$$f_+^{\eta'} = \frac{1}{\sqrt{3}}f_+^0\left\{(a_P-b_P)+\frac{1}{\sqrt{2}}(c_P-d_P)\right\},$$

where $F_0\simeq 1.05F_\pi$, $F_8\simeq 1.25F_\pi$ as explained in the text, and $F_\pi=0.093$ GeV. Also, $f_+^0=0.33$. Similar expressions exist for the $\eta$ meson as well. Vector and axial vector form factors are defined as

$$\begin{aligned}\langle K^*|\overline{s}\gamma_\mu(1+\gamma_5)b|B\rangle &= ig(q^2)\epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}(p_B+p_K)^\alpha(p_B-p_K)^\beta \\ &\quad +(m_B^2-m_K^2)f_1(q^2)\epsilon_\mu^* + (\epsilon^*\cdot q) \\ &\quad \times[f_2(q^2)(p_B+p_K)_\mu + f_3(q^2)(p_B-p_K)_\mu],\end{aligned} \tag{A3}$$

$$\langle M^*|\overline{q_1}\gamma_\mu q_2|0\rangle = -i\sqrt{2}F_{M*}m_{M*}\epsilon_\mu^* .$$

The relationship of the form factors used in this paper to those of [15] is

$$g(q^2)=\frac{iV(q^2)}{m_B+m_{K*}},\quad f_1(q^2)=\frac{iA_1(q^2)}{m_B-m_{K*}},$$

$$f_2(q^2)=\frac{iA_2(q^2)}{m_B+m_{K*}},$$

$$f_3(q^2)=\frac{i}{q^2}[2m_{K*}A_0-(m_B+m_{K*})A_1+(m_B-m_{K*})A_2], \tag{A4}$$

where we take $A_1=A_2=A_3=0.33$ GeV for the sake of simplicity. $L(\mu_i)$ and $M$ are the kinematical parameters

$$L_k(\mu_i)=1-\mu_K+\frac{f_-^K(q^2)}{f_+^K(q^2)}\mu_\eta,$$

$$L_\eta(\mu_i)=1-\mu_\eta+\frac{f_-^\eta(q^2)}{f_+^\eta(q^2)}\mu_K,$$

$$M_k(\mu_i)=\frac{1}{2}\left[[3-y+(1-3y)\mu_K-(1-y)\mu_\eta] + \frac{f_-^K(q^2)}{f_+^K(q^2)}[1-y+(1-y)\mu_K+(1+y)\mu_\eta]\right], \tag{A5}$$

and $y\simeq 1/2$–1 is related to the momentum distribution of quarks inside of the mesons.

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