Large CP Asymmetries in $B^{\pm} \rightarrow \eta_c(\chi_{c0})\pi^{\pm}$ from the $\eta_c(\chi_{c0})$ Width

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We study CP asymmetries in $B^{\pm} \rightarrow h\pi^{\pm}$ decays, where the hadronic states $h = \rho\rho, K\overline{K}\pi, \pi^+\pi^-K^+K^-$, etc., and $h = \pi^+\pi^-, K^+K^-, 2(\pi^+\pi^-)$, etc., are taken on the resonances η_c and χ_{c0}, respectively. The relatively large η_c and χ_{c0} decay widths, of about 10–15 MeV, provide the necessary absorptive phase in the interference between the resonance (going through $b \rightarrow c\overline{c}d$) and the background (through $b \rightarrow u\overline{u}d$) contributions to the amplitude. Large asymmetries of order 10% or more are likely in some modes.

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In the standard model one expects CP violation to show up in B decays in a variety of ways [1]. CP asymmetries are in principle easier to measure in charged B decays than in neutral B decays. They are, however, harder to calculate, since they usually depend on hadronic matrix elements of quark operators and on unknown final-state strong interaction phases. Aside from upper limits [2], there is no evidence for final-state phases in B decays, and it is often argued that these phases are small because of the high B mass. As an example, for the Cabibbo-suppressed process $B^- \rightarrow \psi\pi^-$ [3] (going dominantly through $b \rightarrow c\overline{c}d$), a model was used [4] to describe the rescattering from intermediate hadronic states formed by $c\overline{c}d\overline{u}$ and $u\overline{u}d\overline{u}$ to the final $\psi\pi^-$ state. A small rescattering phase was found, leading to an asymmetry at a percent level.

In the present Letter we will show that one way to overcome small final-state interaction phases is to look for B decays that go through wide $c\overline{c}$ resonances, where the resonance width provides the necessary phase. We will study charged B decay processes dominated by $b \rightarrow c\overline{c}d$, in which the $c\overline{c}$ pair forms one of the two known wide spin-zero states, η_c and χ_{c0}. These states may be identified by their hadronic decay modes h, e.g., $h = \rho\rho, K\overline{K}\pi, \pi^+\pi^-K^+K^-$ and $h = 2(\pi^+\pi^-), \pi^+\pi^-K^+K^-$, respectively, for which typical branching ratios of a few percent have been measured, or by $h = \pi^+\pi^-, K^+K^-$ into which χ_{c0} decays at a percent level [5]. The relatively large η_c and χ_{c0} decay widths, of about 10–15 MeV, provide a CP conserving phase which is effectively maximal ($\pi/2$). Interference of the resonating amplitude with a direct B decay amplitude going through $b \rightarrow u\overline{u}d$, carrying a different Cabibbo-Kobayashi-Maskawa (CKM) phase and leading to the same hadronic final states, creates a large CP asymmetry. We will show that due to this unique mechanism of CP violation, CP asymmetries in these charged B decay processes are much larger than in $B^{\pm} \rightarrow \psi\pi^{\pm}$, and are likely to reach a level of 10% or more. Furthermore, high statistics data may allow separate measurements of the resonance amplitude and of the direct amplitude which acts as a background. This would allow determining the CKM phase $\gamma = \arg(V_{ub}^*)$. (We use the standard convention [5].)

Resonance width effects in charged B decay asymmetries were studied recently [6]. The leading effect was that of the interference of two (or more) intermediate kaon resonances decaying to the same final states. Interference between a resonant and a nonresonant amplitude was already used as a CP violating mechanism in top quark decays, where the W width is the source of a CP-even phase [7].

Let us describe in quite general terms the mechanism of CP violation in $B^+ \rightarrow X_c\pi^+$, in which the charmonium state $X_c = \eta_c$ or χ_{c0} decays to one of the above hadronic final states h. For simplicity, we will consider only two-body and quasi-two-body X_c decays. In this case the B^+ decay distribution can be described in terms of s, the center-of-mass energy squared of the hadrons h, and θ, the angle between the B momentum and the momentum of one of the two X_c decay products in the h center-of-mass frame. In multibody decays, θ would be replaced by several kinematical variables. We denote by a_1 the weak decay amplitude of B^+ to $X_c\pi^+$ and by a_2 the X_c decay amplitude to h. The resonance amplitude, which also includes a Breit-Wigner form for the intermediate X_c state, is given by

$$R(s) \equiv A(B^+ \rightarrow X_c\pi^+ \rightarrow h\pi^+) = a_1a_2\frac{\sqrt{\Gamma m}}{(s - m^2) + i\Gamma m}. \quad (1)$$

m is the X_c mass and Γ is its width. To calculate the contribution of the interference of this amplitude with another amplitude to a CP asymmetry, subtraction of the partial decay width into h is required by CPT [8]. As discussed for nonperturbative cases in the first paper of Ref. [6], this procedure is equivalent to using the full width, together with some compensating processes, an absorptive part of which cancels the rescattering contribution in the original process. The compensating processes are obtained by interchanging the cut with the final state [8]. Since the partial width into h is very small relative to Γ, ~1% in our cases of interest, it

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will be neglected. Note that R does not depend on θ since X_c is spinless. The CKM phase of a_1, given by $\arg(V_{cb}^* V_{cd})$, vanishes in the standard convention which we use. Hence $a_1 a_2$ is taken to be real. Final state phases due to rescattering from other intermediate states will be absorbed into the strong phase of the direct amplitude to be discussed below.

For convenience, we will normalize the decay rate of $B^+ \to X_c \pi^+ \to h\pi^+$ by the total B decay rate

$$\frac{1}{\Gamma_B}\frac{d^2\Gamma(\text{resonance})}{dsdz} = |R(s)|^2, \qquad z \equiv \cos\theta, \quad (2)$$

such that $a_1 a_2$ is given by the product of the corresponding decay branching ratios:

$$2\pi a_1^2 a_2^2 = B(B^+ \to X_c \pi^+) B(X_c \to h). \quad (3)$$

As mentioned earlier, the B^+ decay amplitude into the final state $h\pi^+$, at $s = m^2$, consists also of a direct decay term (D) induced by $\bar{b} \to \bar{u}ud$ carrying the CKM phase γ. We neglect a small contribution from penguin amplitudes [1]. The possible slight s dependence of D around $s = m^2$ will be neglected. In general this amplitude depends on the angle θ. The direct amplitude, for $s \approx m^2$, is given by

$$D(s \approx m^2, z) \equiv A(B^+ \to h\pi^+) = \frac{d(z)}{m_B} e^{i\gamma} e^{i\delta}, \quad (4)$$

where $d(z)$ is real and δ is a final-state interaction phase. Assuming that δ is small, which motivated our search for large resonance width effects in the first place, we take hereafter $\delta = 0$. $d(z)$ can be decomposed into contributions from different spin-parity h states,

$$d(z) = \sum_{J^P} d_{(J^P)}(z), \quad (5)$$

in which an S wave ($J = 0$), for instance, corresponds to a constant term. The direct decay rate will also be normalized by the total B decay rate such that $d(z)$ becomes dimensionless

$$\frac{1}{\Gamma_B}\frac{d^2\Gamma(\text{direct}, s \approx m^2)}{dsdz} = |D(s \approx m^2, z)|^2 = \frac{d^2(z)}{m_B^2}. \quad (6)$$

The $B^+ \to h\pi^+$ amplitude at $s \approx m^2$ is given by a coherent sum of the resonance amplitude R and the direct amplitude D,

$$\frac{1}{\Gamma_B}\frac{d^2\Gamma^{(+)}}{dsdz} = |R(s) + D(s \approx m^2, z)|^2. \quad (7)$$

The corresponding amplitude for $B^- \to h\pi^-$ is obtained simply by changing the sign of the weak phase γ in $D(z)$. The difference and the sum of B^+ and B^- differential decay rates, *integrated symmetrically around* $s = m^2$, say from $s = (m - 2\Gamma)^2$ to $s = (m + 2\Gamma)^2$, are given by

$$\frac{1}{\Gamma_B}\left(\frac{d\Gamma^{(+)}}{dz} - \frac{d\Gamma^{(-)}}{dz}\right) \approx -12 a_1 a_2 d(z) \frac{\sqrt{\Gamma m}}{m_B} \sin\gamma,$$

$$\frac{1}{\Gamma_B}\left(\frac{d\Gamma^{(+)}}{dz} + \frac{d\Gamma^{(-)}}{dz}\right) \approx 6 a_1^2 a_2^2 + 16 d^2(z) \frac{\Gamma m}{m_B^2}. \quad (8)$$

The partial rate asymmetry,

$$A \equiv \frac{\Gamma^{(+)} - \Gamma^{(-)}}{\Gamma^{(+)} + \Gamma^{(-)}}, \quad (9)$$

requires an integration of (8) over z. In the numerator the resonance amplitude interferes only with a component of the direct amplitude corresponding to the hadronic system h with the charmonium J^P quantum numbers. Therefore, only one J^P term of $d(z)$ contributes, 0^- for η_c, and 0^+ for χ_{c0}. In the denominator we will assume for now that the resonance contribution to the decay rate is much larger than the direct contribution over the resonance region, $a_1^2 a_2^2/\Gamma m \gg d^2(z)/m_B^2$. We will comment on corrections to this approximation when estimating the two contributions for the relevant decays. We find

$$A \approx -2\left(\frac{d_{(0^P)}}{a_1 a_2}\right)\frac{\sqrt{\Gamma m}}{m_B}\sin\gamma. \quad (10)$$

Equation (10) is our central result. The asymmetry is given in terms of twice the ratio of magnitudes of the direct and resonance amplitudes at $s = m^2$ times $\sin\gamma$. Usually, an asymmetry contains also a sine of a CP-conserving phase. *In our case, in which interference occurs between the resonance amplitude and the direct decay amplitude corresponding to the background process at* $s = m^2$, *the strong phase difference is maximal, i.e.,* $\pi/2$.

The resonance amplitude $a_1 a_2$ is given in (3) in terms of a product of measurable branching ratios. The 0^P direct amplitude $d_{(0^P)}$ can be obtained by a partial wave analysis of the z distribution slightly off the resonance. Thus, we find

$$|A| \approx 2\sqrt{\pi\Gamma m}\sqrt{\frac{(1/\Gamma_B)\, d\Gamma(\text{direct}, s \approx m^2, 0^P)/ds}{B(B^+ \to X_c\pi^+)B(X_c \to h)}}\sin\gamma, \quad (11)$$

where $d\Gamma(\text{direct}, s \approx m^2, 0^P)/ds$ is the 0^P contribution to the differential decay rate slightly off the resonance. This expression of the asymmetry can be used to determine the weak phase γ from measurable quantities.

To estimate the asymmetry, let us relate $d_{(0^P)}$, the 0^P direct amplitude, to a measurable integrated quantity. $d(z)$, the total direct amplitude at $s \approx m^2$, may be estimated from the s-integrated $B^+ \to h\pi^+$ nonresonance branching ratio, $B(B^+ \to h\pi^+)_{\text{nr}}$. This branching ratio corresponds to h systems which do not originate in other s-channel resonances.

In order to integrate over s and z the nonresonance differential decay rate $|D(s,z)|^2$, which acts as a background, we will have to make an assumption about its s dependence. Using the variable z introduces an extra s-dependent factor into $D(s,z)$ relative to $D_{\text{inv}}(s,z)$, which is up to a constant factor the commonly used invariant amplitude [5]:

$$|D|^2 = \Phi(s)|D_{\text{inv}}|^2, \qquad \Phi(s) \equiv \sqrt{1 - \frac{4m_0^2}{s}}\left(1 - \frac{s}{m_B^2}\right). \quad (12)$$

We consider the case in which the two hadrons in h have equal masses m_0, and we set $m_\pi \approx 0$. We will assume that the nonresonance invariant amplitude, $D_{\rm inv}$, is approximately independent of s over the entire range $4m_0^2 \leq s \leq m_B^2$: $|D(m^2,z)|^2 \approx \Phi(s)d^2(z)/\Phi(m^2)m_B^2$. Integrating (6) over s and z and using (5) we find

$$B(B^+ \to h\pi^+)_{\rm nr} \approx \frac{I_\Phi}{2\Phi(m^2)} \sum_{J^P} \int_{-1}^{1} d^2_{(J^P)}(z)\, dz\,, \qquad (13)$$

where

$$I_\Phi\left(\frac{m_0^2}{m_B^2}\right) \equiv \frac{2}{m_B^2} \int_{4m_0^2}^{m_B^2} \Phi(s)\, ds \qquad (14)$$

is a standard three-body phase space factor, which is very close to 1 for $m_0 = m_\pi$ and has a value of 0.68 for $m_0 = m_\rho$.

To estimate the relative contribution of $d_{(0^P)}$ to the right-hand side of (13), one must apply model-dependent considerations. We use qualitative arguments, based on spin counting and on a partial wave analysis for the pion in $B^+ \to h\pi^+$, which may be emitted from a very close distance to the b quark up to a typical hadronic distance away from it. We use a hard sphere approximation. This leads to a suppression factor $f(0^P)$ of the 0^P decay rate relative to the total direct rate. This factor depends on the case under consideration and involves large uncertainties. For instance, for $h = \pi^+\pi^-, K^+K^-$ in which $J = 0$ requires $P = +1$, we find $f(0^P) = 0.07$–0.7, whereas for the cases $h = \rho\rho, K\overline{K}\pi$ where both parities are allowed in a $J = 0$ state, the suppression may be stronger. For definiteness, we use the above range for $f(0^P)$. Equation (13) then leads to

$$d(0^P) \approx \sqrt{f(0^P)\frac{\Phi(m^2)}{I_\Phi} B(B^+ \to h\pi^+)_{\rm nr}}\,. \qquad (15)$$

Since the above limits on $f(0^P)$ correspond to extreme assumptions, it seems to us that central values are more likely. Using (3), (10), and (15) we find

$$|A| \approx \sqrt{f(0^P)\,\frac{\Phi(m^2)}{I_\Phi}}\,\frac{\sqrt{8\pi\Gamma m}}{m_B} \times \sqrt{\frac{B(B^+ \to h\pi^+)_{\rm nr}}{B(B^+ \to X_c\pi^+)B(X_c \to h)}}\sin\gamma\,. \qquad (16)$$

Let us estimate the asymmetry under typical relevant circumstances. We will use η_c and χ_{c0} decay modes with branching ratios at a level of 1% [5]:

$$B(X_c \to h) \sim 10^{-2}. \qquad (17)$$

The two B^+ decay branching ratios in (16) will evidently be known before an asymmetry can be measured. We use the following value for the decay branching ratio into $X_c\pi^+$:

$$\begin{aligned} B(B^+ \to \eta_c(\chi_{c0})\pi^+) &\approx \left|\frac{V_{cd}}{V_{cs}}\right|^2 B(B^+ \to \eta_c(\chi_{c0})K^+) \\ &\sim \left|\frac{V_{cd}}{V_{cs}}\right|^2 B(B^+ \to \psi(\chi_{c1})K^+) \\ &\sim 5\times 10^{-5}. \end{aligned} \qquad (18)$$

The branching ratios of $B^+ \to \psi(\chi_{c1})K^+$ have been measured [5]. The common replacement $\pi^+ \leftrightarrow K^+$ with corresponding CKM factors is also justified by the recent measurement of $B^+ \to \psi\pi^+$ [3]. Recent theoretical estimates of $\Gamma(B^+ \to \eta_c K^+)/\Gamma(B^+ \to \psi K^+)$ [9] seem to indicate that $B(B^+ \to \eta_c\pi^+)$ may be larger than (18) by about a factor of 1.6 or more. On the other hand, estimates that rely on the factorization approximation indicate that we may have overestimated $B(B^+ \to \chi_{c0}\pi^+)$ [10].

The branching ratios of decays into nonresonant $h\pi^+$ states may vary somewhat from case to case. We use as a characteristic value

$$B(B^+ \to h\pi^+)_{\rm nr} \sim 10^{-5}. \qquad (19)$$

This represents typical branching ratios of low multiplicity processes of the type $b \to u\bar{u}d$, such as $B \to \pi\pi$ for which some evidence already exists [11], $B \to \pi\pi\pi, B \to \pi K\overline{K}$, etc. Branching ratios at this level were calculated for two-body and quasi-two-body decays of this type by assuming factorization [12]. Similar or even larger branching ratios are expected when a nonresonating pion is added to the final state, since at the high B mass it is easy to fragment an additional pion. There is supporting evidence for this behavior in D decays, where $B(D^+ \to \pi^+\pi^+\pi-)_{\rm nr} \approx B(D^+ \to \pi^+\pi^0)$ [5]. A statistical model for the pion multiplicity as a function of the available energy [13] predicts that in B decays the decay rate into three nonresonating pions should be larger than for two pions.

Using the above values of branching ratios and central experimental values for the η_c and χ_{c0} masses and widths [5], we find from (16) similar asymmetries for two representative processes, $B^\pm \to (\rho^+\rho^-)_{\eta_c}\pi^\pm$ and $B^\pm \to (\pi^+\pi^-)_{\chi_{c0}}\pi^\pm$:

$$|A| \sim 0.7\sqrt{f(0^P)}\sin\gamma\,. \qquad (20)$$

This is a large asymmetry for the presently allowed values of γ, $0.3 \leq \sin\gamma \leq 1$ [1] and $f(0^P) = 0.07$–0.7. Larger asymmetries are obtained for larger values of $B(B^+ \to h\pi^+)_{\rm nr}$ and for smaller values of $B(B^+ \to X_c\pi^+)B(X_c \to h)$.

We remind the reader that when deriving (10) we neglected the second term on the right-hand side of the lower Eq. (8). The large asymmetry (20) indicates that this direct contribution over the resonance region cannot be neglected. In fact, for the interval of s used to obtain (8), it reduces the numerical coefficient in (20) to about 0.5. This also affects (10) and (11) in a similar manner.

This correction can be made smaller by integrating over a narrower range around the resonance.

The above-estimated CP asymmetry applies to rather rare decay processes, $B^+ \to (h)_{s\approx m^2}\pi^+$, which have branching ratios $\mathcal{B}$ of about

$$\mathcal{B} \equiv B(B^+ \to X_c\pi^+)B(X_c \to h) \sim 5 \times 10^{-7}. \tag{21}$$

The number of charged B's required to observe such an asymmetry at a 3σ level is $N \approx 10(\mathcal{B}A^2)^{-1}$, which depends only on $B(B^+ \to h\pi^+)_{\rm nr}$ and not on $\mathcal{B}$:

$$N \approx \frac{10}{\sin^2\gamma}\,\frac{m_B^2 I_\Phi}{8\pi\Gamma m\Phi(m^2)f(0^P)B(B^+ \to h\pi^+)_{\rm nr}}\left(\frac{0.7}{0.5}\right)^2. \tag{22}$$

$1/B_{\rm nr}$ is likely to be smaller than 10^5 for favorable decay modes, such as $h = K\overline{K}\pi, \pi^+\pi^-K^+K^-$ (for η_c) and $h = \pi^+\pi^-, K^+K^-$ (for χ_{c0}). The 0^P suppression factor, $f(0^P) = 0.07$–0.7, is the most uncertain one and depends on the decay mode under consideration. Putting all numbers together, we see that typically, for $\sin\gamma \sim 1$, about 10^8–10^9 B's are needed to observe an asymmetry. For favorable cases, in which the 0^P suppression is weak (corresponding to a background that is flat in the variable z) and in which the nonresonance branching ratio is large, fewer B mesons may be required.

It is also possible to define another measurable CP violating quantity, which does not involve the 0^P suppression factor, requiring, however, measurement of the angular distributions $d\Gamma^{(\pm)}/dz$. Denoting the differential asymmetry by $a(z)$,

$$a(z) \equiv \frac{d\Gamma^{(+)}/dz - d\Gamma^{(-)}/dz}{d\Gamma^{(+)}/dz + d\Gamma^{(-)}/dz} \approx -2\left(\frac{d(z)}{a_1a_2}\right)\sin\gamma, \tag{23}$$

we define

$$\mathcal{A} \equiv \sqrt{\frac{1}{2}\int_{-1}^{1} a^2(z)\,dz}. \tag{24}$$

Measurement of $\mathcal{A}$ can be used to determine the weak phase γ from an expression identical to (11), except that here there is no restriction on J^P. Using the approximation (13), one obtains for $\mathcal{A}$ an expression as in (16), but without $f(0^P)$. Very large values, $\mathcal{A} \sim 0.5\sin\gamma$ [see (20) and the discussion below it], are expected to be measured for this quantity.

We note that in the analogous Cabibbo-allowed decays $B^\pm \to (h)_{s\approx m^2}K^\pm$, though the rates are larger than in $B^\pm \to (h)_{s\approx m^2}\pi^\pm$ by a factor $|V_{cs}/V_{cd}|^2$, the asymmetries are correspondingly smaller and harder to observe. Finally, our entire analysis applies generally to $B^+ \to \eta_c(\chi_{c0})X^+$, where X^+ is any hadronic state made from $u\overline{d}$, such as $\rho^+, \pi^+\pi^0$, etc. One may also consider the semi-inclusive processes $b \to d\eta_c(\chi_{c0})$ (similar to $b \to d\psi$ [4]) and $\eta_c(\chi_{c0}) \to h$. Their product branching ratios are expected to be about 5×10^{-6}, an order of magnitude larger than the exclusive branching ratios. Their asymmetries are as large as in the exclusive decays. The observation of such asymmetries would be easier if measurement of $b \to d\eta_c(\chi_{c0})$ were possible, in spite of the about 20 times larger background from $b \to s\eta_c(\chi_{c0})$.

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