Puzzles in $B \rightarrow h_c(\chi_{c2})K$ Decays and QCD Factorization

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We study the factorization-forbidden decays $B \rightarrow h_c K$ and $B \rightarrow \chi_{c2} K$ in the QCD factorization approach. If neglecting the vertex corrections and regularizing the end-point singularities in spectator corrections properly, we get small branching ratios for both the two decay modes, which are roughly consistent with the experimental upper limits. This is in contrast to another factorization-forbidden decay $B \rightarrow \chi_{c0} K$, for which a large decay rate is obtained in the same approach.

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1. Introduction. The branching ratios (BRs) of $B \rightarrow h_c(\chi_{c0}, \chi_{c2})K$ decays are expected to be small since they are naively factorization-forbidden. But non-factorizable contributions to these modes could be large and enhance their BRs. This is the case for $B \rightarrow \chi_{c0} K$, which was observed by Belle with surprisingly large BR of $(6.0^{+2.1}_{-1.8} \pm 1.1) \times 10^{-4}$ and was soon confirmed by BaBar with BR of $(2.7 \pm 0.7) \times 10^{-4}$, being comparable to those of factorization-allowed but color-suppressed decays, such as $B \rightarrow J/\psi(\eta_c, \chi_{c1-3})K$.

One may then expect similarly large BRs for $B \rightarrow h_c(\chi_{c2})K$ decays. Accordingly, a number of authors have suggested that $h_c$ could be easily observed in B-factories. However, recent measurements set the upper limit BRs of $B \rightarrow \chi_{c2} K$ and $B \rightarrow h_c K$ at 90% C.L.: \[
\begin{align*}
\text{Br}(B^+ \rightarrow \chi_{c2} K^+) &< 3.0 \times 10^{-5}, \\
\text{Br}(B^0 \rightarrow \chi_{c2} K^0) &< 4.1 \times 10^{-5}, \\
\text{Br}(B^+ \rightarrow \chi_{c2} K^+) &< 3.8 \times 10^{-5},
\end{align*}
\]

which are about an order of magnitude smaller than that of $B \rightarrow \chi_{c0} K$. The puzzling hierarchies challenge our understanding of the properties of the non-factorizable mechanism for these decay modes.

The non-factorizable contributions to these modes have been studied in different approaches. The final-state soft re-scattering effects were suggested and large and comparable rates for all these three modes were predicted, which are evidently inconsistent with new experimental data. On the other hand, with the Light-Cone Sum Rules, the non-factorizable contributions were found to be too small to account for the large $B \rightarrow \chi_{c0} K$ decay rate. Recently, $B \rightarrow \chi_{c0} K^{(*)}$ were also studied in the PQCD approach, where the $k_T$ factorization was adopted. They predicted a large rate of $B \rightarrow \chi_{c0} K$ by including the spectator contributions only, but no predictions for rates of $B \rightarrow h_c(\chi_{c2}) K$ were given.

Within the framework of QCD factorization, $B \rightarrow \chi_{c0,2} K$ decays were studied earlier, and we found that there exist infrared divergences in the vertex corrections and end-point singularities in the leading twist spectator corrections, which implies that soft contributions may be large in these decays. These contributions come from the soft gluon exchange between the $K$ or $B$ meson and the $c\bar{c}$ pair which is emitted as a color-octet at the short-distance weak interaction vertex. This may imply a connection of the infrared behavior in the exclusive decays with the color-octet contributions in the inclusive $B$ decays to charmium, which are found to be the dominant mechanism. The results of Ref. suggest that a sizable fraction of the large color-octet partial rate of inclusive $B$ decay into charmonium does in fact end up in some two-body decay modes. This is even true for the factorization-allowed decays, such as $B \rightarrow J/\psi K$, for which the infrared safe leading-twist contributions only result in a BR about $1 \times 10^{-4}$, which is an order of magnitude smaller than the experimental data. However, after considering the soft contributions arising from higher-twist spectator interactions, the prediction can account for the data quite well. The prominent effects of soft spectator interactions are also used in Ref. to explain the large BR of $B \rightarrow \psi(3770)K$ decay.

Similarly, these soft contributions to $B \rightarrow \chi_{c0} K$ were also estimated by the authors of Ref. and by us in QCD factorization. Both they and we found that there were linear singularities in chirally enhanced twist-3 spectator interactions, which are numerically large and dominant in $B \rightarrow \chi_{c0} K$ decay, but the results of theirs and ours were quite different because different treatments of the end-point singularities in the spectator interactions. They also got a small BR for $B \rightarrow \chi_{c2} K$ by introducing undetermined imaginary parts for soft spectator interactions.

Since the $s(\bar{s})$ quark emitted from the weak interaction vertex moves fast in the $B$ meson rest frame, we may expect that the soft gluon exchange is dominated by that between the $c\bar{c}$ pair and the spectator quark that goes into the kaon. For $B \rightarrow \chi_{c0} K$ decay, this fact has been confirmed in QCD factorization and PQCD factorization approaches. In particular, in we find that the vertex corrections could be small, and the spectator corrections are large and dominant. In this letter,
we will evaluate the BRs of $B \to h_c(\chi_{c2})K$ by using the same schemes given in \cite{17}. We will show that if the contributions from vertex corrections are neglected, we can fit the experimental data of $B \to h_c(\chi_{c2})K$ quite well.

2. Amplitudes in QCD factorization. We treat the charmonium as a non-relativistic $c\bar{c}$ bound state. Let $p$ be the total momentum of the charmonium and $2q$ be the relative momentum between $c$ and $\bar{c}$ quarks, then $v^2 \sim 4q^2/p^2 \sim 0.25$ can be treated as a small expansion parameter. For the P-wave charmonium, because the wave function at the origin $\mathcal{R}_1(0) = 0$, the amplitudes need to be expanded to the first order in the relative momentum $q$ (see, e.g., \cite{13}).

$$\mathcal{M}(B \to 2S+1P_j(c\bar{c})) = \sum_{LS} \int \frac{d^4q}{(2\pi)^3} \langle m_c | P_{SS}(p, q) \rangle \mathcal{O}(q) \mathcal{O}(0),$$

where $\mathcal{O}(q)$ represents the rest of the decay matrix elements and are expected to be further factorized as product of $B \to K$ form factors and hard kernel or as the convolution of a hard kernel with light-cone wave functions of $B$ meson and $K$ meson within QCD factorization approach. The spin projection operators $P_{SS}(p, q)$ are constructed in terms of quark and anti-quark spinors as

$$P_{00}(p, q) = \frac{3}{m_c} \sum_{s_1, s_2} \langle \frac{p}{2} - q, s_2 | \bar{u}(\frac{p}{2} + q, s_1) | 0 \rangle$$

$$= -\frac{3}{4m_c^2} \langle \frac{p}{2} - q, s_2 | \bar{u}(\frac{p}{2} + q, s_1) | 0 \rangle$$

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and

$$\mathcal{O}(0) = \mathcal{O}(q)|_{q=0},$$

$$\mathcal{O}(q) = \mathcal{O}(0) + \mathcal{O}(1).$$

In Eq. (2) we take charmonium mass $M \simeq 2m_c$. Here $m_c$ is the charm quark mass.

The integral in Eq. (2) is proportional to the derivative of the P-wave wave function at the origin by

$$\int \frac{d^4q}{(2\pi)^3} \langle \frac{p}{2} - q, s_2 | \bar{u}(\frac{p}{2} + q, s_1) | 0 \rangle = \frac{3}{4m_c^2} \langle \frac{p}{2} - q, s_2 | \bar{u}(\frac{p}{2} + q, s_1) | 0 \rangle.$$

We use the following polarization relations respectively for $h_c(J = 1)$ and $\chi_{c2}(J = 2)$:

$$\sum_{LS} \mathcal{O}(q) = \mathcal{O}(0) + \mathcal{O}(1).$$

$$\mathcal{O}(0) = \mathcal{O}(q)|_{q=0}.$$

(7)

where $\mathcal{O}(q)$ is the polarization vector for $h_c$, and the polarization tensor $\mathcal{O}(q)$, which is symmetric under the exchange $\alpha \leftrightarrow \beta$, is the one appropriate for $\chi_{c2}$.

The effective Hamiltonian relevant for $B \to h_c(\chi_{c2})K$ is written as \cite{24}:

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} \left[ V_{cb} V^*_{cs} C_1 \mathcal{O}^I_1 + C_2 \mathcal{O}^I_2 \right] \times A_j. \tag{8}$$

where in the parentheses, "$-"$ is corresponding to $h_c$ and "$+"$ to $\chi_{c2}$, respectively. Here and afterwards, we will use the subindex $j$ to denote different charmonium states, and $j = 1, 2$, represent $h_c$ and $\chi_{c2}$ respectively. Thus the coefficients $A_j$ are given by

$$A_j = \frac{i6R_1(0)}{\sqrt{2 \pi M_j^2}} \frac{\alpha_c C_F}{4\pi N_c} \left[ F_1(M_j^2) \cdot n_j (fI_j + fIII_j) \right] \tag{9}$$

Here $F_I$ is calculated from the four vertex diagrams (a, b, c, d) and $fIII_j$ is calculated from the two spectator diagrams (e, f) in \cite{10}. The function $fIII_j$ receives contributions from both twist-2 and twist-3 light-cone distribution amplitudes of the K meson, and we will simply symbolize them as $fII_j$ and $fIII_j$, respectively.

The factors $n_j$ in (10) are given by

$$n_1 = 2e^* \cdot p_b, n_2 = -\frac{4\sqrt{2}e^* \cdot p_b}{(1 - z_2) m_B} \tag{11}$$

where $z_j = M_j^2 / m_B^2 \simeq 4m^2 / m_B^2$.

For $B \to h_c(\chi_{c2})K$, as just like $B \to \chi_{c0}K$, the soft gluons coupled to the charm quark pair through color...
dipole interactions, which are proportional to the relative momentum \( q \) and give the leading order contributions both in \( 1/m_b \) and NR expansion (see Eq. (2)). As a result, there are generally infrared divergences in \( f_{Ij} \), while \( f_{II}^{2,3} \) suffer from logarithmic and linear end-point singularities, respectively. The infrared divergences in \( f_{Ij} \) can be regularized by a gluon mass \( m_g \) or by the binding energies \( b_j = M_j - 2m_c \) of charmonia,

\[
f_{Ij} = -\frac{8z(1-z + \ln z)}{(1-z)^2} \ln \left( \frac{m_g^2}{m_b^2} \right) + \text{finite terms},
\]

\[
= -\frac{4z((1-z)(1+3z)+4z\ln z)}{(1-z)^2} \ln \frac{b_j}{M_j} + \text{finite terms}, \tag{12}
\]

which are consistent with Ref. [12, 17] and Ref. [16], respectively. Here and afterwards we will omit the subscript of \( z_j \) unless it is necessary.

To derive the functions \( f_{II} \), we use the light-cone projector of \( K \) meson up to twist-3,

\[
M^K_{\alpha\beta}(p') = \frac{if_K}{4} \left\{ p' \gamma_5 \phi_K(y) - i\sigma_{\mu\nu} p' \frac{\partial}{\partial k_{2\mu}} \frac{\phi_K(y)}{6} \right\}_{\alpha\beta}, \tag{13}
\]

where \( k_{2(1)} \) is the momentum of the anti-quark (quark) in \( K \) meson, and the derivative acts on the hard-scattering amplitudes in the momentum space only. The chirally enhanced mass scale \( \mu_K = m_K^2/(m_c(\mu) + m_d(\mu)) \) is comparable to \( m_b \), which ensures that the twist-3 spectator interactions are numerically large, though they are suppressed by \( 1/m_b \).\(^1\) The twist-2 light-cone distribution amplitude (LCDA) \( \phi_K(y) \) and the twist-3 ones \( \phi^K_1(y) \) and \( \phi^K_2(y) \) are symmetric under exchange \( y \leftrightarrow (1-y) \) in the limit of SU(3) symmetry. In practice, we choose the asymptotic form \( \phi_K(y) = \phi^K_1(y) = 6y(1-y) \) and \( \phi^K_2(y) = 1 \) for these LCDAs.

The projectors in Eq. (13) have been used by us in Ref. [16] and there exist logarithmic and linear end-point singularities in twist-2 and twist-3 spectator interactions, respectively. Similar singularities are also found by the authors of Ref. [16]. But they use a different projector, which can be derived from Eq. (13) by adopting an integration by parts on \( y \) and dropping the boundary terms. As having mentioned in Ref. [16], the two projectors can be consistent with each other only when the singularities in all the regions with small-virtualities are carefully regularized. That is, to introduce a relative off-shellness \( \lambda \) to regularize every factor "\( y \)" in the denominator. For linear singularities, we have

\[
\int \frac{y^n \, dy}{(y + \lambda)^{n+2}} = \frac{1}{(n+1)\lambda} - 1 + O(\lambda), \quad n = 0, 1, 2, ... \tag{14}
\]

The difference between Eq. (14) in our scheme for small-virtuality and that in Ref. [16] is quite evident. If simply parameterizing the linear singularities on the left hand side of Eq. (14) as \( \int dy / y^2 = m_B / \Lambda_B \) following Ref. [16], one would get the same results for all \( n \geq 0 \). We will see that the non-trivial factors \((n+1)^{-1}\) in Eq. (14) play an important role in evaluating the BRs of \( B \to h_c(\chi_{c2})K \) decays.

Under this scheme, as we expected, the projectors in Eq. (13) and in Ref. [16] give the same forms of \( f_{II}^{2,3} \):

\[
\begin{align*}
    f_{II}^1 &= a_1 \frac{m_B}{\Lambda_B} [3z + O(\lambda)], \\
    f_{II}^2 &= a_2 \frac{m_B}{\Lambda_B} [6z \ln \lambda + 15z + O(\lambda)], \tag{15}
\end{align*}
\]

\[
\begin{align*}
    f_{II}^1 &= a_1 r_K \frac{m_B}{\Lambda_B} \frac{3z - (1+z) \ln \lambda - \frac{3}{2} (1-z)}{2 \lambda}, \\
    f_{II}^2 &= a_2 r_K \frac{m_B}{\Lambda_B} \frac{z - (1+2z) \ln \lambda - \frac{1}{2} (3+2z)}{\lambda} + O(\lambda)[(16)]
\end{align*}
\]

where \( r_K = 2\mu_K/m_b \) is of \( O(1) \) and the factors \( a_j \) are defined as

\[
a_j = \frac{8\pi^2 f_K f_B}{N_c (1-z)^2 m_B^2 F_1 (m_B^2 / \Lambda_B^2)} \tag{17}
\]

In Eq. (15) and Eq. (16), the integral with LCD of B meson is conventionally parameterized as \( \int_0^1 \frac{d\xi}{\xi} \phi_{B(\xi)} = \frac{m_B}{\Lambda_B} \) with \( \Lambda_B \approx 300 \text{ MeV} \).[16]

\(^1\) In fact, the twist-3 contributions to these decay modes are not power suppressed because of the linear singularities contained in them.
It is worth emphasizing that this scheme is physical since the off-shellness of quarks and gluons are naturally serve as infrared cutoffs when $y \to 0$. Following Ref. [17], we use the binding energy $b_j$ to determine the relative off-shellness $\lambda_j$:

$$
\lambda_j \equiv \frac{z_j}{1-z_j} \cdot \frac{b_j}{M_j}.
$$

Before explicitly evaluating these amplitudes, we should emphasize the importance of the twist-3 spectator contributions again. Normalizing all the amplitudes by the finite part of $f_{II}$, we can see that both $f_{III}$'s and $f_{II}^2$'s are of the order of $\ln \sim \ln(m_B/\Lambda)$, but $f_{II}^3$'s are of order $\frac{m_B}{\Lambda} \sim \frac{m_B}{\Lambda}$ because of the chiral enhancement and the linear singularities contained in them. These unusual power counting rules give support to the naive expectation that the soft spectator interactions would be dominant in these three decay modes. At the qualitative level, this statement has been validated in Ref. [16, 17] for $B \to \chi_{c0} K$.

3. Numerical results and discussions. For numerical analysis, we use the following input parameters with two values for the QCD renormalization scale $\mu = 1.45$ GeV and $\mu = 4.4$ GeV:

$$
M_1 = 3.524 \text{ GeV}, \quad M_2 = 3.556 \text{ GeV}, \quad m_c = 1.5 \text{ GeV},
$$

$$
m_B = 5.28 \text{ GeV}, \quad \Lambda_B = 300 \text{ MeV}, \quad R_{\chi_c}(0) = 0.375 \text{ GeV}^{5/2},
$$

$$
f_B = 216 \text{ MeV}^{21}, \quad f_K = 160 \text{ MeV},
$$

$$
F_1(M_1^2) = 0.80, \quad F_1(M_2^2) = 0.82^{22},
$$

$$
r_K(\mu) = 0.85(1.3), \quad \alpha_s(\mu) = 0.34(0.22).
$$

In Eq. (19) the $\mu$-dependent quantities at $\mu = 1.45$ GeV ($\mu = 4.4$ GeV) are shown without (with) parentheses. The Wilson coefficients $c_i$'s are evaluated at leading order by renormalization group approach [20], since the amplitudes in Eq. (10) are only of the leading order in $\alpha_s$.

We use LoopTools [23] for numerical calculations of $f_{II}$ with a gluon mass $m_g$, and the obtained BRS $Br_j$ are listed in Tab. I with $m_g$ varying from 200 Mev to 500 Mev. The values of $f_{II}$ are also shown in Tab. I. We see the $\mu$-dependence of $f_{II}$ arising from $r_K(\mu)$ is largely cancelled by $\alpha_s(\mu)$ [10]. In Tab. I, for comparison, we also list the results of $B \to \chi_{c0} K$ decay, which can be found in Ref. [17]. Here and afterward, we will use subscript "hc, $\chi_{c2}, \chi_{c0}" to symbolize the quantities for $h_c, \chi_{c2}$ and $\chi_{c0}$ respectively.

We predict a relative small BR for $B \to \chi_{c2} K$ with $R_{\chi_{c2}} \equiv Br_{\chi_{c2}}/Br_{h_c} = 3 - 6$ in Tab. I, which gives a signal of hierarchy without fine tuning of parameters. However, the BR of $B \to h_c K$ is large and the ratio $R_1 \equiv Br_{\chi_{c0}}/Br_{h_c} \simeq 2$, which is not large enough to account for experimental data.

Further, we regularize the infrared divergences in $f_{I}(j = h_c, \chi_{c2}; \chi_{c0})$ using the binding energy. The results for $f_{I}(\chi_{c2})$ and $f_{I}(\chi_{c0})$ are available in Ref. [16]. We find both $f_{I}(\chi_{c2})$ and $f_{I}(\chi_{c0})$ in the binding energy scheme are smaller than those in the gluon mass scheme (with the gluon mass values chosen above). Since the large $Br_{\chi_{c0}}$ are mainly due to the large twist-3 spectator contribution $f_{II}^3$, it is not very sensitive to the value of $f_I$, but Br$_{\chi_{c2}}$ is. As a result, we get $R_{\chi_{c2}} > 20$ no matter which $\mu$ is chosen. Numerical calculations indicate that the value of $f_{I} h_c$ in the binding energy scheme is also smaller than that in the gluon mass scheme.

The largeness of $R_{\chi_{c2}}$ is mainly due to the difference of $f_{II}^3$ and $f_{III}^3$, and can be traced to the infrared behavior of the spectator interactions. This is the reason why our results are different from those of Ref. [16], in which the coefficient of the linear pole in $f_{II}^3$ is two times larger than ours in [16]. As a result, the ratio $R_{\chi_{c2}}$ can be large only by adding adjusted imaginary parts to spectator functions $f_{II}$ there.

As we have mentioned, since the $s(\bar{s})$ quark going out from the weak interaction vertex moves fast in the $B$ meson rest frame, the soft gluon emitted from the $c\bar{c}$ quark pair is favored to be absorbed by the spectator quark. This picture has been confirmed by the power counting rules deduced above and by our analysis for $B \to \chi_{c0} K$ decay [17]. On the other hand, the vertex functions $f_{I}(j)$ in general are numerically small but sensitive to regularization schemes. So it might be reasonable to evaluate amplitudes with only spectator contributions $f_{II}^3$ and $f_{III}^3$, and treat vertex corrections as uncertainties, of which the contributions to $Br_j$ are expected to be smaller than $1 \times 10^{-4}$, as in the case of factorizable decay $B \to J/\psi K$. And it could be even smaller for nonfactorizable decays.

\[2\] Here, the functions $f_{II}^{2,3}$ are extracted from Eq. (10) with $n_{\chi_{c0}} = -(1-z)m_B/\sqrt{z^2}$, so their definitions differ from those in Ref. [17] just by a minus sign.
With the approximation of neglecting contributions from vertex corrections, the results are listed in Tab.II.

In this approximation, the soft spectator interactions dominate, and there is no surprise that we get a larger $\text{Br}_{\chi_{c0}}$ and a much smaller $\text{Br}_{\chi_{c2}}$ simultaneously. The hierarchy is more evident and $R_{\chi_{c}} \sim 4$. Furthermore, the rate of $h_c$ is smaller than that in Tab.I, and $R_{h_c} \sim 4$, which is roughly consistent with the experimental data if we take into account the new PDG average value $\text{Br}_{\chi_{c0}} = (1.6 \pm 0.4) \times 10^{-4}$.

Since soft scattering can introduce a strong-interaction phase, the functions $f_{II}$ could be complex. In our scheme, the phases may emerge with logarithmic singularities in $f_{II}$ if $\lambda$ is negative. To account for this effect, here every logarithm in (15) and (16) can be multiplied by a universal factor $(1 + e^{i\theta})$ with $0 < \rho < 1.5$. The phase $\theta$ is expected to be small, say, varying from 0 to $\pi/2$. We find that the corrections to those $BR$’s do not change our conclusions drastically. For example, choosing $\rho = 1.3$ with $\theta$ varying from 0 to $\pi/2$, we get $\text{Br}_{\chi_{c0}} \simeq (0.4 \pm 4.0) \times 10^{-5}$, $R_{\chi_{c0}} \simeq 100$ and $R_{h_c} \simeq 4$, which are roughly consistent with experimental data, respectively.

In summary, since the non-factorizable contributions arising from soft spectator interactions could be large for $B$ to charmonium exclusive decays, the smallness of experimental upper limits of $BR$s of $B \to h_c(\chi_{c2})K$ are surprising. We study these two modes within the framework of QCD factorization. If neglecting the vertex corrections and regularizing the end-point singularities in spectator corrections properly, we get small $BR$s of $0.27 \times 10^{-4}$ and $0.03 \times 10^{-4}$ for $B \to h_cK$ and $B \to \chi_{c2}K$ respectively, while the predicted $BR$ for $B \to \chi_{c0}K$ is large, as shown in Table II. These are roughly consistent with experiments. On the other hand, at present we do not have a rigorous treatment for the vertex corrections because of the appearance of infrared divergences. Nevertheless, if we use the gluon mass (or binding energy) to regularize the infrared divergences (with reasonable values for the gluon mass or binding energy), we find their effects are possibly small, and may therefore be viewed as theoretical uncertainties. If we take these uncertainties into account, the predictions for $BR$s of $B \to h_c(\chi_{c2})K$ become $\text{Br}_{h_c} = (3 \pm 10) \times 10^{-5}$ and $\text{Br}_{\chi_{c2}} = (0 \pm 10) \times 10^{-5}$. Here the large errors in predictions are mainly attributed to uncertain but sub-important (even power-suppressed) contributions arising from those vertex corrections.

Note. Very recently, the decay mode $B \to h_cK$ has also been studied in PQCD factorization approach [25]. As in the case of $B \to \chi_{c0}K$ [4], the authors of [25] neglected the vertex corrections, and get a small rate $\text{Br}(B \to h_cK) = 4.5 \times 10^{-5}$, which is roughly consistent with ours.

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