Final-state interactions in the decays $B^0 \rightarrow \chi_{c0} K^{*0}$ and $B^+ \rightarrow \chi_{c0} K^{*+}$

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Abstract

In this article, we study the final-state rescattering effects in the decays $B^0 \rightarrow \chi_{c0} K^{*0}$ and $B^+ \rightarrow \chi_{c0} K^{*+}$, and observe the corrections are zero in the $SU(3)$ limit, which is warranted by the heavy quark symmetry. It is difficult to accommodate the experimental data without fine-tuning.

PACS numbers: 13.20.He; 14.40.Lb

Key Words: Final-state interactions, $B$-decay

1 Introduction

Recently, the Babar collaboration reported the observation of the decay $B^0 \rightarrow \chi_{c0} K^{*0}$ as well as evidence of the decay $B^+ \rightarrow \chi_{c0} K^{*+}$ with an 8.9 and a 3.6 standard deviation significance respectively [1]. The measured branching fractions are $Br(B^0 \rightarrow \chi_{c0} K^{*0}) = (1.7 \pm 0.3 \pm 0.2) \times 10^{-4}$ and $Br(B^+ \rightarrow \chi_{c0} K^{*+}) = (1.4 \pm 0.5 \pm 0.2) \times 10^{-4}$. They also obtained the upper limit $Br(B^+ \rightarrow \chi_{c0} K^{*+}) < 2.1 \times 10^{-4}$ at the 90% confidence level. The decays take place through the process $b \rightarrow sc\bar{c}$ (or more precise $\bar{b} \rightarrow s\bar{c}\bar{c}$, they relate with each other by charge conjugation, in this article, we calculate the amplitudes for the process $b \rightarrow sc\bar{c}$, then take charge conjunction to obtain the branching fractions.) at the quark-level. The quantitative understanding of those decays depends on our knowledge about the nonperturbative hadronic matrix elements of the operators entering the effective weak Hamiltonian. The factorizable contributions in the decays $B^0 \rightarrow \chi_{c0} K^{*0}$ and $B^+ \rightarrow \chi_{c0} K^{*+}$ are zero, we have to resort to special mechanism to overcome the difficulty.

Final-state interactions play an important role in the hadronic $B$-decays, the color-suppressed neutral modes such as $B^0 \rightarrow D^0 \pi^0, \pi^0 \pi^0, \rho^0 \pi^0, K^0 \pi^0$ are enhanced substantially by the long-distance rescattering effects [2]. In Refs.[3, 4], the authors study the rescattering effects of the intermediate charmed mesons for the decays $B^- \rightarrow \chi_{c0} K^-, h_c K^-$, and observe the final-state interactions can lead to larger branching fractions to account the experimental data. The factorizable amplitude in the decay $B^0 \rightarrow \eta_c K^*$ is too small to accommodate the experimental data [5], and the effects of the final-state interactions can smear the discrepancy [6]. It is intersecting to study the effects of the final-state interactions in the decays $B^0 \rightarrow \chi_{c0} K^{*0}$ and $B^+ \rightarrow \chi_{c0} K^{*+}$.

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The article is arranged as: in Section 2, we study the final-state rescattering effects in the decays $B^0 \rightarrow \chi_{c0} K^{*0}$ and $B^+ \rightarrow \chi_{c0} K^{*+}$; in Section 3, the numerical result and discussion; and Section 4 is reserved for conclusion.

2 Final-state rescattering effects in the decays $B^0 \rightarrow \chi_{c0} K^{*0}$ and $B^+ \rightarrow \chi_{c0} K^{*+}$

The effective weak Hamiltonian for the decay modes $b \rightarrow s c \bar{c}$ can be written as (for detailed discussion of the effective weak Hamiltonian, one can consult Ref. [7])

$$H_w = \frac{G_F}{\sqrt{2}} \left\{ V_{cb} V_{cs}^* \left[ C_1(\mu) O_1(\mu) + C_2(\mu) O_2(\mu) \right] + h.c. \right\} ,$$

where the $V_{ij}$'s are the CKM matrix elements, the $C_i$'s are the Wilson coefficients calculated at the renormalization scale $\mu \sim O(m_b)$ and the relevant operators $O_i$ are given by

$$O_1 = (\bar{c} \sigma_{\alpha} b_\alpha) V_{-A} (\bar{c} \beta c_\beta) V_{-A},$$

$$O_2 = (\bar{c} \sigma_{\alpha} b_\alpha) V_{-A} (\bar{c} \beta c_\alpha) V_{-A},$$

(2)

here $\alpha$ and $\beta$ are color indexes. From the effective weak Hamiltonian $H_w$, we can see that the factorizable amplitudes are zero.

The decays $B^0 \rightarrow D D_s, D D_s^*, D^* D_s, D^* D_s^*$ are color enhanced due to the large Wilson coefficient $C_2$.

$$\langle D_s(q) D_s^*(p) | H_w | B(P) \rangle = 2 \tilde{a}_2 P \cdot \epsilon^*(p) f_{D_s} M_{D^*} A_0(q^2),$$

$$\langle D_s(q) D_s^*(p) | H_w | B(P) \rangle = 2 \tilde{a}_2 P \cdot \epsilon^*(q) f_{D_s} M_{D^*_s} F_1(q^2),$$

$$\langle D_s(q) D_s^*(p) | H_w | B(P) \rangle = \tilde{a}_2 (M_D^2 - M_{D^*_s}^2) f_{D_s} F_0(q^2),$$

$$\langle D_s^*(q) D_s^*(p) | H_w | B(P) \rangle = \tilde{a}_2 f_{D_s} M_{D_s^*} \left[ \frac{2 \epsilon^{\mu \nu \alpha \beta} \epsilon^*_\nu(q) \epsilon^*_\mu(p) P_\alpha P_\beta V(q^2)}{M_B + M_{D^*}} - \epsilon^*(q) \cdot \epsilon^*(p) \right]$$

$$\left( (M_B + M_{D^*}) A_1(q^2) + \frac{2 P \cdot \epsilon^*(q) \epsilon^*(p) \cdot q A_2(q^2)}{M_B + M_{D^*}} \right),$$

(3)

where $\tilde{a}_2 = \frac{G_F}{\sqrt{2}} V_{cb} V_{cs}^* (C_2 + \frac{C_1}{3})$, the $f_{D_s}$, $f_{D_s^*}$, $f_{D_s}$, $f_{D_s^*}$ are the weak decay constants, and the $A_0(q^2)$, $A_1(q^2)$, $A_2(q^2)$, $V(q^2)$, $F_0(q^2)$, $F_1(q^2)$ are the weak form-factors.

2 The corresponding decays for the $B^+$ can be studied in the same way, in the following, we only present the technical details for the $B^0$ decays.
defined by [8, 9],

\[
\langle D(p)|\bar{\sigma}_\mu (1-\gamma_5)b|B(P) \rangle = (P+p)_\mu F_1(q^2) - \frac{M_B^2 - M_D^2}{q^2} q_\mu [F_1(q^2) - F_0(q^2)],
\]

\[
\langle D^*(p)|\bar{\sigma}_\mu (1-\gamma_5)b|B(P) \rangle = \epsilon_\mu^{\nu \alpha \beta} e^p_\alpha p_\beta 2V(q^2) \frac{2M_V q \cdot \epsilon^*}{M_B + M_V} q_\mu A_0(q^2) - \left( \epsilon^* \frac{q \cdot \epsilon^*}{q^2} q_\mu \right) (M_B + M_V) A_1(q^2) + \left[ (P+p)_\mu - \frac{M_B^2 - M_D^2}{q^2} q_\mu \right] q \cdot \epsilon^* \frac{A_2(q^2)}{M_B + M_V},
\]

and the \( \epsilon_\mu \) is the polarization vector of the vector meson, \( q_\mu = P_\mu - p_\mu \).

The decays \( B^0 \to \chi_{c0} K^0 \) and \( B^+ \to \chi_{c0} K^+ \) can take place through the decay cascades \( B \to DD_s, DD_s^*, D^* D_s, D^* D_s^* \to \chi_{c0} K^* \), the rescattering amplitudes of \( DD_s, DD_s^*, D^* D_s, D^* D_s^* \to \chi_{c0} K^* \) may play an important role.

The final-state interactions can be described by the following effective lagrangians,

\[
\mathcal{L}_{\chi_{c0} DD} = g_{\chi_{c0} DD} \chi_{c0} DD^\dagger, \tag{5}
\]

\[
\mathcal{L}_{\chi_{c0} D^* D^*} = g_{\chi_{c0} D^* D^*} \chi_{c0} D^* D^\dagger, \tag{6}
\]

\[
\mathcal{L}_{DDV} = -ig_{DDV} D_1^{\mu \nu} \bar{D}_\mu D^\dagger_j (V^\mu)^j, \tag{7}
\]

\[
\mathcal{L}_{D^* DV} = -2f_{D^* DV} \epsilon^{\mu \nu \alpha \beta} (\partial_\mu V^\nu) [D^\dagger_i \partial_\alpha D^1_j - D^*_1 \partial_\alpha D^i_j], \tag{8}
\]

\[
\mathcal{L}_{D^* DV} = ig_{D^* DV} D_1^{\mu \nu} \partial_\mu D^{* \nu} (V^\mu)^j + 4if_{D^* DV} D_1^{\mu \nu} \partial_\mu (\partial^\mu V^\nu - \partial^\nu V^\mu)^j D^{* \nu} j, \tag{9}
\]

where the indexes \( i, j \) stand for the flavors of the light quarks, \( D^{(s)} = (\bar{D}^{0, s}) \), \( (D^{(s)} = (\bar{D}^{0, s})^T \), and \( V \) is the 3 \times 3 matrix for the nonet vector mesons,

\[
V = \begin{pmatrix}
\rho^0 \sqrt{2} + \omega \sqrt{2} & \rho^+ \sqrt{2} & K^{*+} \\
\rho^- \sqrt{2} & -\rho^0 \sqrt{2} + \omega \sqrt{2} & K^{*0} \\
K^{*-} & \bar{K}^{*0} & \phi
\end{pmatrix}. \tag{10}
\]

The lagrangians \( \mathcal{L}_{DDV}, \mathcal{L}_{D^* DV} \) and \( \mathcal{L}_{D^* DV} \) are taken from Ref.[2], and the \( \mathcal{L}_{\chi_{c0} DD^*} \) and \( \mathcal{L}_{\chi_{c0} D^* D^*} \) are constructed from the heavy quark theory in this article.

The rescattering effects can be taken into account by eight Feynman diagrams, see Fig.1. We calculate the absorptive parts (or imaginary parts) of the rescattering amplitudes \( \text{Abs}(i) \) by the Cutkosky rule,

\[
\text{Abs}(i) = \frac{1}{2} \int \frac{d^3 \vec{p}_1}{(2\pi)^3 2E_1} \int \frac{d^3 \vec{p}_2}{(2\pi)^3 2E_2} (2\pi)^4 \delta^4 (P - p_1 - p_2) T^i_{B \to \chi_{c0} K^*} T^i_{\chi_{c0} \to B}, \tag{11}
\]

where the amplitudes \( T^i_{B \to \chi_{c0} K^*} \) stand for the corresponding factorizable contributions.
presented in Eq.(3), and the rescattering amplitudes $T_{i\to\chi_{c0}K^*}^i$ are given by

$$T_{DD_s^*\to\chi_{c0}K^*}^a = -2ig_{DDV}p_2 \cdot \epsilon^*(p_4) \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}DD},$$

$$T_{DD_s^*\to\chi_{c0}K^*}^b = 2ig_{DDV}p_2 \cdot \epsilon^*(p_4) \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}DD},$$

$$T_{DD_s^*\to\chi_{c0}K^*}^c = 4if_{D^*D^*V}\epsilon^{\mu\alpha\beta} p_{4e\epsilon}(p_4)p_{2e\epsilon}(p_2) \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}D^*D^*}\epsilon^*(q) \cdot \epsilon(p_1),$$

$$T_{DD_s^*\to\chi_{c0}K^*}^d = 4if_{D^*D^*V}\epsilon^{\mu\alpha\beta} p_{4e\epsilon}(p_4)p_{2e\epsilon}(p_2) \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}D^*D^*}\epsilon^*(q) \cdot \epsilon(p_1),$$

$$T_{DD_s^*\to\chi_{c0}K^*}^e = -4if_{D^*D^*V}\epsilon^{\mu\alpha\beta} p_{4e\epsilon}(p_4)p_{2e\epsilon}(p_2) \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}D^*D^*}\epsilon^*(q) \cdot \epsilon(p_1),$$

$$T_{DD_s^*\to\chi_{c0}K^*}^f = -4if_{D^*D^*V}\epsilon^{\mu\alpha\beta} p_{4e\epsilon}(p_4)p_{2e\epsilon}(p_2) \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}D^*D^*}\epsilon^*(q) \cdot \epsilon(p_1),$$

$$T_{DD_s^*\to\chi_{c0}K^*}^g = \{4if_{D^*D^*V} [p_4 \cdot \epsilon(p_2)\epsilon^*(p_4) \cdot \epsilon(q) - p_4 \cdot \epsilon(q)\epsilon^*(p_4) \cdot \epsilon(p_2)] + 2ig_{D^*D^*V}\epsilon^*(p_4) \cdot \epsilon(p_2) \cdot \epsilon(p_1) \} \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}D^*D^*}\epsilon^*(q) \cdot \epsilon(p_1),$$

$$T_{DD_s^*\to\chi_{c0}K^*}^h = \{-4if_{D^*D^*V} [p_4 \cdot \epsilon(p_2)\epsilon^*(p_4) \cdot \epsilon(q) - p_4 \cdot \epsilon(q)\epsilon^*(p_4) \cdot \epsilon(p_2)] - 2ig_{D^*D^*V}\epsilon^*(p_4) \cdot \epsilon(p_2) \cdot \epsilon(p_1) \} \frac{1}{t - M_D^2} \mathcal{F}^2(t)ig_{\chi_{c0}D^*D^*}\epsilon^*(q) \cdot \epsilon(p_1),$$

where $t = q^2$, $q = p_1 - p_3 = p_4 - p_2$, and the $\epsilon_\mu$ is the polarization vector of the corresponding vector meson $V$, $\epsilon_\alpha(q)\epsilon_\beta(q) = -g_{\alpha\beta} + \frac{q_{\alpha q_\beta}}{M_V^2}$. The $p_1, p_2, p_3$ and $p_4$ stand for the momenta of the mesons $D, D_s, \chi_{c0}$ and $K^*$ respectively in the amplitude $T_{DD_s^*\to\chi_{c0}K^*}^a$; the momenta in other amplitudes can be understood analogously. The off-shell effects of the $t$-channel exchanged mesons $D, D^*, D_s$ and $D_s^*$ are taken into account by introducing a monopole form-factor [2],

$$F(M_i, t) = \frac{\Lambda_i^2 - M_i^2}{\Lambda_i^2 - t},$$

and the cutoff $\Lambda_i$ are parameterized as $\Lambda_i = M_i + \alpha\Lambda_{QCD}$, where $\alpha$ is a free parameter and $\Lambda_{QCD} = 0.225\text{GeV}$. In fact, the $g_sF(M_i, t)$ are the momentum dependent strong coupling constants, we can vary the parameter $\alpha$ to change the effective strong couplings, here we use the notation $g_s$ to denote all the strong coupling constants.

The dispersive parts (or real parts) of the rescattering amplitudes can be obtained via the dispersion relation,

$$\text{Dis}(i)(M_B^2) = \frac{1}{\pi} \mathbb{P} \int_{s_{th}}^\infty \frac{\text{Abs}(i)(s')}{s' - M_B^2} ds',$$

where the thresholds $s_{th}$ are given by $s_{th} = (M_D + M_{D_s})^2, (M_{D^*} + M_{D_s})^2, (M_D + M_{D_s})^2, (M_{D^*} + M_{D_s})^2$ for any specific diagram.
3 Numerical result and discussions

In the flavor $SU(3)$ limit, there are strong cancelation among the rescattering amplitudes, $T^0_D,_{D_s\to\chi_c K^*} + T^0_D,_{D_s\to\chi_c K^*} = 0$ and $T^0_D,_{D_s\to\chi_c K^*} + T^0_D,_{D_s\to\chi_c K^*} = 0$. From the Particle Data Group, $M_D = 1.87 \text{ GeV}$, $M_{D_s} = 1.97 \text{ GeV}$, $M_{D_s} = 2.010 \text{ GeV}$ and $M_{D_s} = 2.112 \text{ GeV}$ [10], we can see that the $SU(3)$ breaking effects are small. However, the experimental data from the CLEO collaboration, $f_D = 222.6\pm16.7^{+2.8}_{-3.4} \text{ MeV}$ [11, 12] and $f_{D_s} = (0.274 \pm 0.013) \text{ GeV}$ [13] show that the $SU(3)$ breaking effects are rather large, $f_D/f_{D_s} = 1.23$, while most of theoretical calculations indicate $f_D/f_{D_s} \approx 1.1$, the discrepancy maybe indicate new physics beyond the standard model [14]. If we take into account the small $SU(3)$ breaking effects, the rescattering amplitudes $T^i$, ($i = a, b, g, h$) have contributions.

For the rescattering amplitudes $T^i$, $i = c, d, e, f$, there are no cancelation. We carry out the integrals formally,

\[ \text{Abs}(i) = \int \frac{d^3p_1}{2E_1} \frac{d^3p_2}{2E_2} \delta^4(P - p_1 - p_2) f(p_1, p_2, p_3, p_4) \epsilon^{\mu
u\alpha\beta} p_{4\mu} \epsilon^{*}_\nu(p_4) p_{2\alpha} p_{1\beta} \]

\[ = \epsilon^{\mu
u\alpha\beta} p_{4\mu} \epsilon^{*}_\nu(p_4) [A g_{\alpha\beta} + B p_{3\alpha} p_{3\beta} + C p_{3\alpha} p_{4\beta} + D p_{4\alpha} p_{4\beta}] \]

where we have introduced the formal notations $f(p_1, p_2, p_3, p_4)$ and $A, B, C, D$ (scalar coefficients), they have no contributions.

The above conclusion and the following discussion also hold for the decay $B^+ \rightarrow \chi_c K^{*+}$.

The strong coupling constants $f_{D^*DV}$, $f_{D^*D^*V}$, $g_{DDV}$, $g_{D^*D^*V}$, $g_{\chi_c DD}$ and $g_{\chi_c D^*D^*}$ are not free parameters. In the heavy quark limit, the strong coupling constants $f_{D^*DV}$, $f_{D^*D^*V}$, $g_{DDV}$ and $g_{D^*D^*V}$ can be related to the basic parameters \( \lambda \) and \( \beta \) in the heavy quark effective lagrangian and relevant parameters, we neglect them for simplicity,

\[ f_{D^*DV} = \frac{f_{D^*D^*V}}{M_{D^*}} = \frac{\lambda g_V}{\sqrt{2}}, \]

\[ g_{DDV} = \frac{g_{D^*D^*V}}{M_{D^*}} = \frac{\beta g_V}{\sqrt{2}}, \]

where $g_V = 5.8$ from the vector meson dominance theory [15]; we can also calculate them with the light-cone QCD sum rules [17, 18]. The strong coupling constants $g_{\chi_c DD}$ and $g_{\chi_c D^*D^*}$ can be estimated with the universal Isgur-Wise form-factor at zero recoil $\xi(1)$ and the assumption of dominance of the intermediate $\chi_c K$ meson for the scalar heavy quark current $\tau_c$ [19],

\[ g_{\chi_c DD} = \frac{2M_{D^*}M_{\chi_c}}{f_{\chi_c}}, \]

\[ g_{\chi_c D^*D^*} = -\frac{2M_{D^*}M_{\chi_c}}{f_{\chi_c}}, \]

(17)
where the decay constant $f_{\chi_c}$ is defined by $f_{\chi_c} M_{\chi_c} = \langle 0 | \bar{c}(0) c(0) | \chi_c \rangle$.

The only free parameter is the $\alpha$ in the off-shell form-factors $F(M_i, t)$. We may take into account the experimental data from the Babar collaboration with fine-tuning of the momentum dependent strong coupling constants $g_s F(M_i, t)$. We will not resort to the fine-tuning mechanism.

In fact, the $t$-channel rescattering amplitudes $DD_s, D^* D_s, D D_s^* \rightarrow \chi_c K^*$ can be described by the collective strong coupling constant $g_s$, $g_s^2 = g_s^2 F^2(M_i, t) \frac{1}{t-M_i^2}$.

At the level of quark-gluon degrees of freedom, the rescattering occur through $\bar{c} q_1 + \bar{c} q_2 \rightarrow \chi_c + V$. In the heavy quark limit, the heavy quarks decouple from the light degrees of freedom, the $\bar{q}_1 q_2$ rearrange to the vector meson $V$ (or pseudoscalar meson $P$). Conservation of the heavy quark spin warrants the $\bar{c} c$ pair rearrange to the $J/\psi$ or $\eta_c$, not the $\chi_c$, because additional relative $P$-wave between the $\bar{c} c$ pair is required to form the $\chi_c$, $\chi_{c1}$ and $\chi_{c2}$. It is not unexpected that the total rescattering amplitudes are nearly zero, the small discrepancy due to the $SU(3)$ breaking effects and the $c$ quark is not heavy enough. This case is contrary to the $3P_0$ model [20] [21], where the $\bar{q} q$ pairs with the quantum numbers $3P_0$ are created from the QCD vacuum, the relative $P$-wave between the $\bar{c} c$ pair is canceled out with the relative $P$-wave between the $\bar{q} q$ pair, the decays $\chi_c + \bar{q} q \rightarrow D \bar{D}, D^* \bar{D}^*$ can occur, if they are kinetically allowed, see the effective lagrangians in Eqs.(5-6). The final state interactions may play an important role in the precesses with the final states $\chi_c J + S, A$ and $J/\psi, \eta_c + P, V$, where the $S, P, V$ and $A$ stand for the scalar, pseudoscalar, vector and axial-vector mesons respectively.

4 Conclusion

In this article, we study the final-state rescattering effects in the decays $B^0 \rightarrow \chi_c K^{*0}$ and $B^+ \rightarrow \chi_c K^{*+}$, and observe the corrections are zero in the $SU(3)$ limit, which is warranted by the heavy quark symmetry. It is difficult to accommodate the experimental data without fine-tuning. The final state interactions may play an important role in the decays $B \rightarrow \chi_c J + S, A; J/\psi, \eta_c + P, V$.

Acknowledgments

This work is supported by National Natural Science Foundation, Grant Number 10775051, and Program for New Century Excellent Talents in University, Grant Number NCET-07-0282.

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Figure 1: The Feynman diagrams for the final-state interactions.