**Research Article**

**B → J/ψK(1270) and B → J/ψK(1400) Decays in QCD Factorization Approach**

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The hadronic $B \to J/\psi K(1270)$ and $B \to J/\psi K(1400)$ has been analyzed in “QCD factorization” approach and generalized factorization. The effective coefficients $a_h$ have been calculated for three helicity states $h = 0, +, -$ which give three different contribution of amplitudes. We consider that $J/\psi$ behaves as a light meson in compared to $B$ meson. For $B^+ \to J/\psi K^+(1270)$ and $B^+ \to J/\psi K^+(1400)$, experimental data of branching ratios are $(1.8 \pm 0.5) \times 10^{-3}$ and $<5 \times 10^{-4}$, respectively. Our best obtained results are $(1.79 \pm 0.01) \times 10^{-3}$ at $\theta = 58^\circ$ and $(1.42 \pm 0.01) \times 10^{-3}$ at $\theta = 32^\circ$ for $B^+ \to J/\psi K^+(1270)$ decay. And we have $4.27 \times 10^{-4}$ at $\theta = 58^\circ$ and $3.45 \times 10^{-4}$ at $\theta = 32^\circ$ for $B^+ \to J/\psi K^+(1400)$, which are in agreement with experiment.

1. **Introduction**

Recent experimental results obtained by BABAR, Belle, and CLEO have opened an interesting area of research about production of axial-vector mesons in $B$ decays. Two body $B$ decays look for CP violation and overconstrain the CKM parameters in the Standard Model. Exclusive modes containing $B \to PP, PV, and VV$, which have been extensively discussed in the literature have confirmed such expectation [1]. In this research, the two-body hadronic decays of $B$ meson into $VA$ ($V$: vector, $A$: axial-vector meson) are studied in QCD factorization method. First we introduce the structure of axial-vector meson. Then we study the hadronic decays $B \to J/\psi K(1270)$ and $B \to J/\psi K(1400)$, particularly. There are two distinct types of axial-vector meson, namely, $^3P_1$ and $^1P_1$. In the quark model, two nonets of $J^{PC} = 1^+$ axial-vector meson are expected as the orbital excitation of the $q\bar{q}$ system. In terms of the spectroscopic notation $^{2S+1}L_J$, there are two types of $p$-wave mesons ($^3P_1$ and $^1P_1$). These two nonets have distinctive $C$ quantum numbers, $C = +$ and $C = -$, respectively. Experimentally, the $J^{PC} = 1^{++}$ nonet consists of $a_1(1260), f_1(1285), f_1(1420),$ and
$K_{1A}$, while the $1^{+}$ nonet has $b_{1}(1235)$, $h_{1}(1170)$, $h_{1}(1380)$, and $K_{1B}$. There are two mixing effects for axial-vector meson: one is the mixing between $^{3}P_{1}$ and $^{1}P_{1}$ states, for example, $K_{1A}$ and $K_{1B}$, and the other is mixing among $^{3}P_{1}$ or $^{1}P_{1}$ states themselves [2]. The $^{3}P_{1}$ meson behaves similarly to the vector meson this is not the case for the $^{1}P_{1}$ meson. For the latter, its decay constant vanishes in $SU(3)$ limit. Their light-cone distribution amplitudes are given by using the QCD sum rule method, and the chiral-even two-parton light-cone distribution amplitudes of the $^{3}P_{1}$ ($^{1}P_{1}$) meson are symmetric (antisymmetric) under the exchange of quark and antiquark momentum fractions in the $SU(3)$ limit due to G-parity. The $B$ decays involving an axial-vector meson and vector meson in final state have three polarization states. We have studied the two body $B$ decays involving axial-vector meson $K(1270)$ or $K(1400)$, and a vector meson $J/ψ$ in final state. The simplest approach to obtain the hadronic matrix elements in decay amplitudes is naive factorization, where the matrix elements have been parameterized into the product of decay constant and form factors. In generalized factorization, it is well established that nonfactorizable effects such as vertex corrections, hard spectator interactions, and annihilation contributions are calculable in hard scattering approach, and they have been studied in [3, 4] for the decays $B \rightarrow a_{1}(1260)\pi, a_{1}(1260)K, B \rightarrow h_{1}(1235)K^{*}, b_{1}(1235)K^{*}$ and $b_{1}(1235)\rho$ [5].

Using the product expansion, the low-energy effective Hamiltonian for $B \rightarrow M_{1}M_{2}$ decays can be written generally as [6, 7]

$$H_{\text{eff}} = \frac{G_{F}}{\sqrt{2}} \sum_{p=d,s} \lambda_{p}^{D} \left( C_{1} O_{1}^{p} + C_{2} O_{2}^{p} + \sum_{i=3}^{10} C_{i} O_{i} \right) + \text{h.c.},$$

where $\lambda_{p}^{D} = V_{\rho p} V_{\rho D}^{*}$ $(D = d, s)$ are the product of CKM elements and $C_{i}$’s are Wilson coefficients which have been evaluated at next leading order (NLO). The expressions of local operators are [8].

**Current-current operators:**

$$O_{1}^{p} = \left( \bar{p}_{a} b_{\bar{b}} \right)_{V-A} \left( \bar{q}_{\bar{b}} p_{a} \right)_{V-A'},$$

$$O_{2}^{p} = \left( \bar{p}_{a} b_{\bar{b}} \right)_{V-A} \left( \bar{q}_{\bar{b}} p_{a} \right)_{V-A'}.$$  

**QCD penguin operators:**

$$O_{3} = \left( \bar{q}_{a} b_{\bar{a}} \right)_{V-A} \sum_{q} \left( \bar{q}_{\bar{q}} q_{a} \right)_{V-A'},$$

$$O_{4} = \left( \bar{q}_{a} b_{\bar{a}} \right)_{V-A} \sum_{q} \left( \bar{q}_{\bar{q}} q_{a} \right)_{V-A'}.$$
Electroweak penguin operators:

\[ O_5 = (\bar{q}_a b_a)_{V-A} \sum_q (\bar{q}_\rho q_\sigma)_{V+A} \]

\[ O_6 = (\bar{q}_a b_b)_{V-A} \sum_q (\bar{q}_\rho q_a)_{V+A} \]

\[ (1.3) \]

Here \((\bar{q}_1 q_2)_{V^\pm A} = \bar{q}_1 (1 \pm \gamma_5) q_2\). For hadronic weak decays, the short distance effects can be calculated by perturbative theory but the computation of nonperturbative long distance effects is difficult. The fundamental problem in computation of hadronic matrix elements is due to nonperturbative effects arising from the strong interactions. The simplest approach to hadronic matrix elements is naive factorization hypothesis, where the hadronic matrix elements have been parameterized into the product of the decay constants and the form factors. Beneke et al. suggested QCD formula to compute the hadronic matrix elements in the heavy quark limit, combining the hard scattering approach with power counting in \(1/m_b\) [9, 10]. In heavy quark limit \(m_b \gg \Lambda_{QCD}\), up to power corrections of order of \(\Lambda_{QCD}/m_b\), the QCD formula for \(B \to M_1 M_2\) can be written as

\[ \langle M_1 M_2 | O_i | B \rangle = \sum_j F_j^{B \to M_i} \left( m_2^2 \right) \int_0^1 du T_{ij}^l(u) \phi_{M_2} + (M_1 \leftrightarrow M_2) \\
+ \int_0^1 d\xi du dv T_{ij}^{II}(\xi, u, v) \phi_B(\xi) \phi_{M_1}(v) \phi_{M_2}(u). \]

\[ (1.5) \]

\(M_1\) and \(M_2\) are both light. Which \(F_j^{B \to M}\) is form factor for \(B\) to \(M\) transition and \(\phi_M\) is the light cone distribution amplitude for meson. \(\xi, u, v\) are momentum fractions of constituent quarks in \(B, M_1\) and \(M_2\) mesons, respectively. \(T_i\) and \(T_{ij}\) are hard-scattering kernels arising from hard gluon exchanges. Hence, the hadronic matrix elements can be separated into short distance part (hard-scattering kernels) and long distance part. In QCD factorization, nonfactorizable loop effects and spectator scattering contribution have been considered. At
leading order, there is a single diagram with no hard gluon interaction, and the contribution to \( T_1(u) \) in (1.5) is independent of \( u \), and \( u \)-integral reduces to the normalization condition for the \( M_2 \) wave function. Consequently, the factorization formula (1.5) reproduces naive factorization, if we neglect gluon exchange. In naive factorization, one neglects all corrections of order \( \Lambda_{\text{QCD}}/m_b \) and of order \( \alpha_s \). In QCD factorization, one computes systematically corrections to higher order in \( \alpha_s \), but still neglects power correction of order \( \Lambda_{\text{QCD}}/m_b \). In factorization, we have factorizable and nonfactorizable contributions. There are soft gluons in the factorizable contributions which are absorbed into the physical form factor. The diagrams also have hard contributions, which go into the short distance coefficient. The diagrams containing gluon exchanges which do not belong to the form factor are called nonfactorizable contributions. At order \( \alpha_s \), these contributions can be divided into four groups: vertex corrections, penguin diagrams, hard spectator interactions, and annihilation diagrams. A detailed discussion of QCD factorization approach can be found in [8–10].

This method works well for the case with two light mesons in which the final-state mesons carry large momenta. When there is a heavy quark in the final state such as \( \bar{B} \to D^+\pi^- \), this method still works when a spectator quark of \( B \) meson is absorbed by \( D \) meson. However, when the spectator quark is absorbed by a light quark, for example, in \( \bar{B} \to D\pi^0 \), nonfactorizable contributions are infrared divergent, and the factorization breaks down. When we consider \( B \to J/\psi K(1270) \) and \( B \to J/\psi K(1400) \), it looks ambiguous at first sight whether we can apply the same method used in \( B \to \pi\pi \), or \( \bar{B} \to D^+\pi^- \), since the spectator quark in the \( B \) meson goes into light \( K(1270) \) or \( K(1400) \) meson. However, what is special about \( J/\psi \) is that the size of the charmonium is so small (\(-1/\alpha_s m_{J/\psi}\)) that the charmonium has a negligible overlap with the \((B,K)\) system [11].

### 2. Input Parameters

#### 2.1. Light-Cone Distribution Amplitude (LCDA)

The QCD corrections can change the local quark-anti quark operators into a series of nonlocal operators. If we assume that the \( J/\psi \) behave as a light meson, we can describe the light cone distribution. Note that in QCDF, the LCDAs of the light vector meson are written as [12]

\[
\langle J/\psi(p,\epsilon)|\bar{c}_{1a}(y)c_{2\beta}(x)|0\rangle
= -i\int_0^1 du \exp \{i(up \cdot y + \bar{u}p \cdot x)\]
\times \left\{ f_{J/\psi} m_{J/\psi} \left( \frac{p^\mu z^\nu}{p^\mu z^\nu} \right) \frac{1}{s_{12}} + \frac{1}{2} e_{\mu\nu} \frac{p^\mu p^\nu}{(p \cdot z)^2} \frac{1}{s_{12}} \right\} - \left\{ f_{J/\psi} m_{J/\psi} \left( \frac{p^\mu z^\nu}{p^\mu z^\nu} \right) \frac{1}{s_{12}} \right\},
\]

(2.1)
where \( z = y - x \) with \( z^2 = 0 \). \( \phi_{\parallel}(u) \), \( \phi_{\perp}(u) \) are twist-2 DAs; \( g^{(a)}_{\parallel}(u) \), \( g^{(a)}_{\perp}(u) \), \( h^{(s)}_{\parallel}(u) \) and \( h^{(l)}_{\parallel}(u) \) are twist-3 ones. \( u \) is momentum fraction of \( c \) quark in \( J/\psi \), \( \bar{u} = 1 - u \). \( f_{J/\psi} \), \( f^{1}_{J/\psi} \) are vector and tensor decay constants, respectively. Similarly, for axial-vector meson \( 3^{(1)}P_1 \), we have

\[
\langle \frac{3^{(1)}P_1(p, \varepsilon)}{d_\ast(y)} \frac{u_2\delta(x)}{0} \rangle = -\frac{i}{4} \int_0^1 du \exp [i(v \cdot p + \bar{v} \cdot x)]
\]

\[
\times \left\{ f^{3(1)}_{3\|} m_{3(1)}^{3(1)}(p) + f^{3(1)}_{3\perp} m_{3(1)}^{3(1)}(v) + \frac{1}{4} \epsilon^{\nu \rho \sigma} \epsilon^{\ast \mu} p^\mu z^\nu y^\gamma g_{3\perp}^{3(1)}(v) \right\} + f^{3(1)}_{3\parallel} P_1
\]

\[
\times \left( \frac{f_{1\|} m_{1\|}^{3(1)}(v)}{(p \cdot z)^2} - \frac{m_{1\perp}^{3(1)}(v)}{(p \cdot z)^2} \right) - \frac{m_{1\perp}^{3(1)}(v)}{(p \cdot z)^2} + \frac{h_{3\perp}^{3(1)}(p)}{2} \right) \right\}_{\alpha \beta}.
\]

(2.2)

Applying equation of motion to LCDAs, one can obtain the Wandzura-Wilczek relations in which twist-3 LCDAs related to the twist-2 ones [2]. Also \( \phi_{\pm}(x) \) are defined as

\[
\phi_{\pm}(x) = \frac{g^{(a)}_{\pm}(x)}{4} + \frac{g^{(a)}_{\parallel}(x)}{4} = \int_x^1 \frac{\phi_{\parallel}(u)}{u} du,
\]

(2.3)

which are transverse components of twist-3 LCDAs. While \( \phi_{\pm}^{(a)}(x) \) are longitudinal components of twist-3 ones for vector (axial-vector) meson. We specify the light-cone distribution amplitudes (LCDAs) as

\[
\phi_V(x, \mu) = 6x(1 - x) \left[ 1 + \sum_{n=1}^{\infty} a_n^V(\mu) C_n^{3/2}(2x - 1) \right],
\]

(2.4)

\[
\phi_a(x, \mu) = 3 \left[ 2x - 1 + \sum_{n=1}^{\infty} a_n^V(\mu) P_{n+1}(2x - 1) \right],
\]

for the vector meson, where \( C_n(x) \) are Gegenbauer polynomials and \( P_n(x) \) are Legendre polynomials. The normalization of LCDAs for vector meson is

\[
\int_0^1 dx \phi_V(x) = 1,
\]

(2.5)

\[
\int_0^1 dx \phi_a(x) = 0.
\]
Also,

\[
\phi^A_\parallel(x) = 6x(1-x)\left[1 + 3a_{1}^{l=2}P_1(2x-1)\right],
\]

\[
\phi^A_\perp(x) = 6x(1-x)\left[a_{0}^{l=2}P_1 + 3a_{1}^{l=2}P_1(2x-1)\right],
\]

\[
\phi^A_a(x,\mu) = 3\left[ 2(2x-1) + \sum_{n=1}^{\infty} a_{n}^{l=2}P_1(\mu)P_{n+1}(2x-1) \right],
\]  

for $^3P_1$ axial-vector meson and

\[
\phi^A_\parallel(x) = 6x(1-x)\left[a_{0}^{l=1}P_1 + 3a_{1}^{l=1}P_1(2x-1)\right],
\]

\[
\phi^A_\perp(x) = 6x(1-x)\left[1 + 3a_{1}^{l=1}P_1(2x-1)\right],
\]

\[
\phi^A_a(x,\mu) = 3\left[ (2x-1) + \sum_{n=1}^{\infty} a_{n}^{l=1}P_1(\mu)P_{n+1}(2x-1) \right],
\]  

for $^1P_1$ axial-vector meson. The normalization conditions are

\[
\int_0^1 dx\phi_\parallel(x) = 1,
\]

\[
\int_0^1 dx\phi_\perp(x) = a_{0}^{l=2}P_1,
\]

\[
\int_0^1 dx\phi_a(x) = 0,
\]  

for $^3P_1$ axial-vector meson and

\[
\int_0^1 dx\phi_\parallel(x) = a_{0}^{l=1}P_1,
\]

\[
\int_0^1 dx\phi_\perp(x) = 1,
\]

\[
\int_0^1 dx\phi_a(x) = 0,
\]  

for $^1P_1$ axial-vector meson.

Here $\phi_V$, $\phi^A_\parallel$, and $\phi^A_\perp$ are twist-2 LCDAs; likewise $\phi_a$ and $\phi^A_a$ are twist-3 ones for vector and axial-vector mesons, respectively. The parameters $a_{1}^{l=2}$ and $a_{1}^{l=1}$ are defined by the matrix element of a twist-2 conformal operator with conformal spin 3 in [13]. These parameters are unknown for $J/\psi$; then we consider asymptotic forms of amplitudes for $J/\psi$. Also, the Gegenbauer moments for $K_{1A}$ and $K_{1B}$ are given in [14, 15].
For the $B$ meson, we can write the projection as

$$
\langle 0 | \bar{b}_a(x) q_\beta(0) | B(p) \rangle = -i \frac{f_B}{4} \phi_B(\rho) \left[ (p + m_b) \gamma_5 \right] \beta_\gamma.
$$

(2.10)

A detailed dissection of wave function of $B$ meson can be found [9, 10, 16–18]

$$
\langle 0 | \bar{b}_a(x) q_\beta(0) | B(p) \rangle \bigg|_{(x_+ = x_- = 0)} = -i \frac{f_B}{4} \left[ (p + m_b) \gamma_5 \right] \beta_\gamma 
$$

$$
\times \int_0^1 d\bar{\rho} e^{-\frac{\bar{\rho}^2}{2}} \left[ \phi_1^B(\bar{\rho}) + \frac{m_B}{\lambda_B} \phi_2^B(\bar{\rho}) \right] \bigg|_{\gamma_\alpha}
$$

(2.11)

with $n_- = (1, 0, 0, -1)$, and the normalization conditions are

$$
\int_0^1 d\bar{\rho} \phi_1^B(\bar{\rho}) = 1,
$$

$$
\int_0^1 d\bar{\rho} \phi_2^B(\bar{\rho}) = 0.
$$

(2.12)

The $B$ wave function corresponds to $\lambda_B$ defined by $\int_0^1 d\bar{\rho} \phi_1^B(\bar{\rho}) / \bar{\rho} = m_B / \lambda_B$. $\lambda_B$ is the first inverse moment of the $B$ meson’s distribution which is of order $\Lambda_{\text{QCD}}$.

### 2.2. Decay Constant

Decay constants of vector and axial-vector mesons are defined as

$$
\langle V(p, \epsilon) | \bar{q}_2 \gamma_\mu q_1 \rangle | 0 \rangle = -i f_V m_V \epsilon_\mu^*,
$$

$$
\langle A(p, \epsilon) | \bar{q}_2 \gamma_\mu \gamma_5 q_1 \rangle | 0 \rangle = i f_A m_A \epsilon_\mu^*.
$$

(2.13)

Also, the transverse decay constants are defined via the tensor current as

$$
\langle V(p, \epsilon) | \bar{q}_2 \sigma_{\mu\nu} q_1 \rangle | 0 \rangle = -f_V^\perp \left( \epsilon_\mu^* p^\nu - \epsilon_\nu^* p^\mu \right),
$$

$$
\langle A(p, \epsilon) | \bar{q}_2 \gamma_\mu \gamma_5 q_1 \rangle | 0 \rangle = i f_A m_A \epsilon_\mu^*.
$$

(2.14)

In general, the decay constants $f_{1_{\gamma_\mu}}$ and $f_{3_{\gamma_\mu}}$ are zero in the $SU(3)$ limit.
2.3. Form Factor

The form factors of \( B \rightarrow A \) transitions are defined as

\[
\langle A(p, \varepsilon) | A_\mu | B(p_B) \rangle = i \frac{2}{m_B - m_A} \epsilon_{\mu \nu a b} \varepsilon^\nu_{\eta} p^a_B p^b A^{BA}(q^2),
\]

\[
\langle A(p, \varepsilon) | V_\mu | B(p_B) \rangle = - \left\{ (m_B - m_A) \epsilon_\mu V_1^{BA}(q^2) - (\varepsilon^* \cdot p_B) (p_B + p) \frac{V_2^{BA}(q^2)}{m_B - m_A} \right. \\
-2m_A \frac{\varepsilon^* \cdot p_B}{q^2} q^\mu \right. \\
\left. \left[ V_3^{BA}(q^2) - V_0^{BA}(q^2) \right] \right\},
\]

and for \( B \rightarrow V \) transitions, we have

\[
\langle V(p, \varepsilon) | V_\mu | B(p_B) \rangle = \frac{2}{m_B + m_V} \epsilon_{\mu \nu a b} \varepsilon^\nu_{\eta} p^a_B V^{BV}(q^2),
\]

\[
\langle V(p, \varepsilon) | A_\mu | B(p_B) \rangle = - \left\{ (m_B + m_V) \epsilon_\mu A_1^{BV}(q^2) - \frac{\varepsilon^* \cdot p_B}{m_B + m_V} (p_B + p) \frac{A_2^{BV}(q^2)}{m_B + m_V} \right. \\
-2m_V \frac{\varepsilon^* \cdot p_B}{q^2} q^\mu \right. \left. \left[ A_3^{BV}(q^2) - A_0^{BV}(q^2) \right] \right\},
\]

where \( q = p_B - p \), and for \( q^2 = 0 \) we have \( A_0^{BV} = A_0^{BV} \) and \( V_3^{BA} = V_0^{BA} \). \( \epsilon_\mu \) is the polarization four-vector, and \( V, A_1, A_2, A_3, \) and \( A_0 \) are form factors and

\[
V_3^{BA}(q^2) = \frac{m_B - m_A}{2m_A} V_1^{BA}(q^2) - \frac{m_B + m_A}{2m_A} V_2^{BA}(q^2),
\]

\[
A_3^{BV}(q^2) = \frac{m_B + m_V}{2m_V} A_1^{BV}(q^2) - \frac{m_B - m_V}{2m_V} A_2^{BV}(q^2).
\]

The form factors of \( B \rightarrow K_{1A} \) and \( B \rightarrow K_{1B} \) transitions have been calculated in the light-cone sum rule (LCSR) [19]. The momentum dependence of form factors is calculated in the LCSR approach which is parameterized in three-parameter form:

\[
F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^4/m_B^4)},
\]

where \( q^2 = m_{J/\psi}^2 \), and two parameters \( a \) and \( b \) are given in Table 1.

2.4. \( B \rightarrow J/\psi K(1270) \) and \( B \rightarrow J/\psi K(1400) \) Decays in QCD Factorization

Now, we want to calculate the branching ratios of \( B \rightarrow J/\psi K(1270) \) and \( B \rightarrow J/\psi K(1400) \) decays in QCD factorization method. First, we have calculated the branching ratios in generalized factorization which the effective Wilson coefficients are not dependent on
where \( K \) is the physical decaying amplitudes as linear combinations. So, the \( d_i^{\text{eff}} \) coefficients are

\[
a_i^{\text{eff}} \equiv c_i^{\text{eff}} + \frac{c_{i+1}^{\text{eff}}}{N_c} \quad (i = \text{odd}),
\]

\[
a_i^{\text{eff}} \equiv c_i^{\text{eff}} + \frac{c_{i-1}^{\text{eff}}}{N_c} \quad (i = \text{even}).
\]

As it known the physical states \( K_{1}(1270) \) and \( K_{1}(1400) \) are mixture of states \( ^1P_1 \) and \( ^3P_1 \), whose relations could be parameterized by

\[
K_{1}(1270) = K_{1A} \sin \theta + K_{1B} \cos \theta,
\]

\[
K_{1}(1400) = K_{1A} \cos \theta - K_{1B} \sin \theta,
\]

where \( K_{1A} \) and \( K_{1B} \) are \( ^3P_1 \) and \( ^1P_1 \) axial-vector mesons, respectively, and \( \theta \) is the mixing angle. From experimental data on masses and partial ratios of \( K_{1}(1270) \) and \( K_{1}(1400) \), it is found two solutions for the mixing angle with a two-fold ambiguity, \( \theta = \pm 32^\circ \) and \( \theta = \pm 58^\circ \) [1]. Hence, the physical decaying amplitudes are given by

\[
A(B \rightarrow J/\psi K_{1}(1270))_{\rho} = A(B \rightarrow J/\psi K_{1A}) \sin \theta + A(B \rightarrow J/\psi K_{1B}) \cos \theta,
\]

\[
A(B \rightarrow J/\psi K_{1}(1400))_{\rho} = A(B \rightarrow J/\psi K_{1A}) \cos \theta - A(B \rightarrow J/\psi K_{1B}) \sin \theta.
\]
Since the final states of $B \rightarrow VA$ carry spin degrees of freedom, we have calculated the helicity amplitudes, separately. In our work, the longitudinal amplitude $h = 0$ and the transverse amplitudes $h = \pm 1$ are calculated for $B \rightarrow J/\psi K_{1A}$ and $B \rightarrow J/\psi K_{1B}$ decays [22]. If we assume that $J/\psi$ behaves as a light meson, due to its size, we can describe the light-cone distribution. For $B \rightarrow J/\psi K(1270)$ and $B \rightarrow J/\psi K(1400)$ decays, we have colour suppressed contribution $a_2$, $a_5$, $a_7$, and $a_9$. Then, the contributions of vertex corrections for these decays are only for $i = 2, 3, 5, 7, 9$. Hence, three different helicity amplitudes have been written as

\begin{align}
A^h(B \rightarrow J/\psi K_{1A}) &= \left[V_{cb} V_{cs}^* a_2^h - V_{tb} V_{ts}^* \left(a_3^h + a_5^h + a_7^h + a_9^h\right)\right] X^{(BK_{1A},J/\psi)}_h, \\
A^h(B \rightarrow J/\psi K_{1B}) &= \left[V_{cb} V_{cs}^* a_2^h - V_{tb} V_{ts}^* \left(a_3^h + a_5^h + a_7^h + a_9^h\right)\right] X^{(BK_{1B},J/\psi)}_h,
\end{align}

(2.22)

where the factorizable amplitude is written as

\begin{align}
X^{(BK_{1A,B},J/\psi)}_h &\equiv \langle J/\psi | (\bar{c}c)_{V-A} | 0 \rangle \langle K_{1A,B} | (\bar{b} s)_{V-A} | B \rangle.
\end{align}

(2.23)

The helicity dependencies read as

\begin{align}
X^0_0^{(BK,J/\psi)} &= \frac{if_{J/\psi}}{2m_K} \left[m_B^2 - m_K^2 - m_{J/\psi}^2 (m_B - m_K) V_{1}^{BK} \left(q^2\right)^2 - \frac{4m_B^2 p_c^2}{m_B - m_K} V_{2}^{BK} \left(q^2\right)^2\right], \\
X^{\pm}_{\pm}^{(BK,J/\psi)} &= -i f_{J/\psi} m_B m_{J/\psi} \left[1 - \frac{m_K^2}{m_B^2}\right] V_{1}^{BK} \left(q^2\right)^2 \mp \frac{2p_c}{m_B - m_K} A^{BK} \left(q^2\right)^2,
\end{align}

(2.24)

where $p_c$ is c.m. momentum of decays in the $B$ rest frame

\begin{align}
p_c = \sqrt{\frac{(m_B^2 - (m_1 + m_2)^2)(m_B^2 - (m_1 - m_2)^2)}{2m_B}},
\end{align}

(2.25)

$m_1$ and $m_2$ are the mass of final state mesons. The branching ratio is given by [18]

\begin{align}
BR(B \rightarrow VA) &= \frac{p_c}{8\pi m_B^3 \tau_B} \left(\frac{G_F}{\sqrt{2}} \right)^2 \left|A(B \rightarrow VA)\right|^2,
\end{align}

(2.26)

where $G_F$ is Fermi constant and $\tau_B$ is life time of $B$ meson. If the final states are two light vector mesons, it is expected that $|A_0| > |A_+| > |A_-|$ for $B$ meson decays. (For $\bar{B}$ decays, exchange $- \leftrightarrow +$.) It is because of that the amplitude $|A_+|$ is suppressed by a factor of $\sqrt{2m_2/m_B}$, while the $|A_-|$ amplitude is further suppression in $m_1/m_B$. But for $B$ to heavy-light final states, $\sqrt{2m_{heavy}/m_B}$ is of order unity [23]. In QCD factorization, the effective coefficients $a_i^h$ basically are Wilson coefficients in conjunction with short distance
nonfactorizable corrections such as vertex corrections and hard spectator interactions. Then, we have [2]

\[ a_i^h(J/\psi K_{1A}(K_{1B})) = \left( C_i + \frac{C_{i+1}}{N_c} \right) N_i(J/\psi) \int_0^1 \phi_{i/\psi,h}^h(x) dx + \frac{C_{i+1} \alpha_s}{N_c} \left( V_i^h(J/\psi) + \frac{4\pi^2}{N_c} H_i^h(J/\psi K_{1A}(K_{1B})) \right) \]  \tag{2.27}

where \( i = 1, \ldots, 10 \), the lower (upper) signs apply when \( i \) is even (odd), \( C_i \) are Wilson coefficients (we use the next to leading order coefficients, NLO, obtained in the naive dimensional regularization scheme, NDR, at the energy scale \( \mu = m_b(m_b) \)), \( C_F = (N_c^2 - 1)/(2N_c) \) with \( N_c = 3 \), \( J/\psi \) is emitted meson and \( K_{1A}(K_{1B}) \) have the spectator quark of the meson. \( V_i^h(J/\psi) \) is the vertex corrections and \( H_i^h(J/\psi K_{1A}(K_{1B})) \) is the hard-scattering interactions, obtained due to exchange soft and hard gluons. The superscript \( h \) denotes the helicity of the final-state meson, that is, \( h = 0, \pm \). The LCDA \( \phi_{i/\psi}^h \) is \( \phi_{i/\psi}^h \) for \( h = 0 \) and \( \phi_{i/\psi}^\pm \) for \( h = \pm \). And \( N_i^h(J/\psi) = 1 \). We can take the infinite mass limit of the b quark in which \( m_b \) goes to infinity while \( m_{J/\psi} \) is fixed (\( m_{J/\psi}/m_b \to 0 \)). The vertex corrections are given by

\[
V_i^0(J/\psi) = \begin{cases} 
\int_0^1 dx \phi_i^{1/\psi}(x) \left[ 12 \ln \frac{m_b}{\mu} - 18 + g(x) \right], & i = 2, 3, 9, \\
\int_0^1 dx \phi_i^{1/\psi}(x) \left[ -12 \ln \frac{m_b}{\mu} + 6 - g(1 - x) \right], & i = 5, 7, 
\end{cases} \tag{2.28}
\]

\[
V_i^h(J/\psi) = \begin{cases} 
\int_0^1 dx \phi_i^{1/\psi}(x) \left[ 12 \ln \frac{m_b}{\mu} - 18 + g_T(x) \right], & i = 2, 3, 9, \\
\int_0^1 dx \phi_i^{1/\psi}(x) \left[ -12 \ln \frac{m_b}{\mu} + 6 - g_T(1 - x) \right], & i = 5, 7. 
\end{cases} \tag{2.29}
\]

With

\[
g(x) = 3 \left( \frac{1 - 2x}{1 - x} \ln x - i\pi \right) + \left[ 2Li_2(x) - \ln^2 x + \frac{2\ln x}{1 - x} - (x \leftrightarrow 1 - x) \right], 
\]

\[
g_T(x) = g(x) + \frac{\ln x}{1 - x},
\]

we consider asymptotic function for \( \phi_i^{1/\psi}(x) = 6x(1 - x) \), and also we have \( \phi_+(x) = 3(1 - x)^2 \) and \( \phi_-(x) = 3x^2 \). Hard spectator corrections are given by

\[
H_i^0(J/\psi K_{1A(1B)}) = \int_{\mathcal{B}} \frac{f_{1/\psi}(K_{1A(1B)})}{X_0(K_{1A(1B)}, J/\psi)} \frac{m_B}{\Lambda_B} 
\times \int_0^1 dy \, dx \left( \frac{\phi_i^{K_{1A(1B)}}(y) \phi_i^{1/\psi}(x)}{(1-y)(1-x)} - r^i_x \phi_i^{K_{1A(1B)}}(y) \phi_i^{1/\psi}(x) \right), \tag{2.31}
\]
for $i = 2, 3, 9$, and

$$
H^0_i (J/\psi K_{1A(1B)}) = -\frac{if_{\bar{B}f_f} f_{K_{1A(1B)}} m_B}{X^{(BK_{1A(1B)}, J/\psi)}_0} \frac{m_B}{\lambda_B} \times \int_0^1 dy dx \left( \frac{\phi_{K_{1A(1B)}}^{K_{1A(1B)}}(y) \phi^{J/\psi}_+(x)}{(1-y)x} - r_x^{K_{1A(1B)}} \frac{\phi^{K_{1A(1B)}}(y) \phi^{J/\psi}_+(x)}{(1-y)(1-x)} \right),
$$

(2.32)

for $i = 5, 7$. $r_x^{K_{1A(1B)}} = 2(m_{K_{1A(1B)}} f_{K_{1A(1B)}})/(m_\mu f_{K_{1A(1B)}})$ are chirally enhanced parameters. The transverse hard spectator terms are

$$
H^+_i (J/\psi K_{1A(1B)}) = -\frac{2if_{\bar{B}f_f} f_{K_{1A(1B)}} f_{J/\psi} m_{K_{1A(1B)}} m_B}{m_B X^{(B K_{1A(1B)}, J/\psi)}_+} \frac{m_B}{\lambda_B} \times \int_0^1 dy dx \frac{\phi_{K_{1A(1B)}}^{K_{1A(1B)}}(y) \phi^{J/\psi}_+(x)}{(1-y)^2 x},
$$

(2.33)

$$
H^-_i (J/\psi K_{1A(1B)}) = -\frac{2if_{\bar{B}f_f} f_{K_{1A(1B)}} f_{J/\psi} m_{K_{1A(1B)}} m_B}{m_B X^{(B K_{1A(1B)}, J/\psi)}_-} \frac{m_B}{\lambda_B} \times \int_0^1 dy dx \frac{(1-y-x) \phi_{K_{1A(1B)}}^{K_{1A(1B)}}(y) \phi^{J/\psi}_+(x)}{(1-y)^2 (1-x)^2},
$$

(2.34)

for $i = 2, 3, 9$, and

$$
H^-_i (J/\psi K_{1A(1B)}) = -\frac{2if_{\bar{B}f_f} f_{K_{1A(1B)}} f_{J/\psi} m_{K_{1A(1B)}} m_B}{m_B X^{(B K_{1A(1B)}, J/\psi)}_-} \frac{m_B}{\lambda_B} \times \int_0^1 dy dx \frac{(y-x) \phi_{K_{1A(1B)}}^{K_{1A(1B)}}(y) \phi^{J/\psi}_+(x)}{(1-y)^2 x^2},
$$

(2.35)
for $i = 5, 7$. In our calculations, the transverse hard spectator term $H^s_{ij} (J/\psi K_{1A})$ is zero. There are some divergences in other hard spectator terms. We have used the unknown parameter $X_H$ to eliminate these divergences, defined as

$$X_H = \int_0^1 \frac{dx}{1-x} = \ln \left( \frac{m_B}{\Lambda_{\text{QCD}}} \right) (1 + r), \quad (2.36)$$

where $\Lambda_{\text{QCD}}$ is hadronic parameter and $r = \rho_H e^{i \varphi_H}$ which $\rho_H$ and $\varphi_H$ are phases parameters $(0 \leq \rho_H \leq 1, -180 \leq \varphi_H \leq 180)$. We consider $r = 0$ in these decays and $B \to J/\psi K$ [24]. This parameter phenomenologically is significant. Hence, the non factorizable contributions (vertex corrections and hard spectator interactions) have been calculated by (2.28)–(2.35).

In the other hand, we know that $J/\psi$ is a heavy vector meson, although in [10] has been supposed that quarkonium is light meson relative to the $B$ meson. If we consider other limit in which $m_c$ goes to infinity, $m_{J/\psi}/m_b$ is held fixed. In this scenario, the amplitude wave function of $J/\psi$ has represented like $B$ meson. Also, to the leading order in $1/m_c$, the $J/\psi$ wave function is [18]

$$\langle J/\psi (p, \epsilon) | \varepsilon_u(x) c_\beta(0) | 0 \rangle \bigg|_{(x,x_i=0)} = \frac{f_{J/\psi}}{4} \left[ e^{i \varphi (p + m_b) r_y} \right]_H \int_0^1 d\xi e^{-i p \cdot x} \left[ \phi^{1/2}_1(\xi) + n \phi^{1/2}_2(\xi) \right] \bigg|_{r_y}. \quad (2.37)$$

Using [18], the effect of higher twist for $J/\psi$ may not converge enough. Because the charmed quark carries a momentum fraction of order $-m_c/m_{J/\psi}$, the distribution amplitudes of $J/\psi$ vanish in the end point region. $\phi_1$ is adapted as the DA of nonlocal vector current of $J/\psi$ rather $g_1^{n}$ as the DA of $\varepsilon_1$ component since the latter does not vanish at the end point. Comparing (2.1) with (2.37) the leading twist DA of $J/\psi$ is

$$\phi^{1/2}_1(u) = \phi^{1/2}_1(u) - 6 u \bar{u},$$

$$f_{J/\psi} = f_{1/2}^{1/2}.$$  \quad (2.38)

Then, for vertex corrections [18], we have

$$V^0_i (J/\psi) = \left\{ \begin{array}{l} \int_0^1 dx \phi^{1/2}_1(x) \left[ 12 \text{Ln} \frac{m_c}{\mu} \cdot 18 + f(x, z) \right] + \int_0^1 dx \phi^{1/2}_1(x) g(x, z), \quad i = 2, 3, 9, \\ \int_0^1 dx \phi^{1/2}_1(x) \left[ -12 \text{Ln} \frac{m_c}{\mu} \cdot 6 + f(x, z) \right] + \int_0^1 dx \phi^{1/2}_1(x) g(x, z), \quad i = 5, 7, \end{array} \right.$$  \quad (2.39)
with

\[
f(x, z) = \frac{2zx}{1 - z(1 - x)} + (3 - 2x) \frac{\ln x}{1 - x} \\
+ \left( -\frac{3}{1 - zx} + \frac{1}{1 - z(1 - x)} - \frac{2zx}{(1 - z(1 - x))^2} \right) xz \ln zx \\
+ \left( 3(1 - z) + 2zx + \frac{2z^2x^2}{1 - z(1 - x)} \right) \ln(1 - z) - i\pi \frac{1}{1 - z(1 - x)} + (1 - z) \\
\times \left\{ -\frac{4x}{(1 - z)(1 - x)} \ln x + \frac{zx}{(1 - z(1 - x))^2} \ln(1 - z) \\
+ \left( \frac{1}{(1 - zx)^2} - \frac{1}{(1 - z(1 - x))^2} + \frac{2(1 + z - 2zx)}{(1 - z)(1 - zx)^2} \right) xz \ln zx \\
- i\pi \frac{zx}{(1 - z(1 - x))^2} \right\},
\]

\[\text{(2.40)}\]

\[
g(x, z) = -4r \frac{\ln x}{1 - x} + \frac{4zx \ln zx}{1 - z(1 - x)} - 4zr \frac{\ln(1 - z) - i\pi}{1 - z(1 - x)} \\
+ (1 - z) \left\{ \frac{4r}{(1 - z)(1 - x)} \ln x - \frac{4r zx}{(1 - z)(1 - zx)} \ln zx \right\},
\]

\[
h(x, z) = \frac{2zx}{1 - z(1 - x)} + (3 - 2x) \frac{\ln x}{1 - x} \\
+ \left( -\frac{3}{1 - zx} + \frac{1}{1 - z(1 - x)} - \frac{2zx}{(1 - z(1 - x))^2} \right) xz \ln zx \\
+ \left( 3(1 - z) + 2zx + \frac{2z^2x^2}{1 - z(1 - x)} \right) \ln(1 - z) - i\pi \frac{1}{1 - z(1 - x)} \\
\]

\[
I(x, z) = -4r \frac{\ln x}{1 - x} + \frac{4zx \ln zx}{1 - z(1 - x)} - 4zr \frac{\ln(1 - z) - i\pi}{1 - z(1 - x)},
\]

where \( z = m_{J/\psi}^2/m_b^2 \) and \( r = f_{J/\psi}^+ m_c / f_{J/\psi} m_{J/\psi} \). In [11] by contracting (2.14) for \( J/\psi \) with \( p_v \) and applying the equation of motion and (2.13) for \( J/\psi \), it finds that \( f_{J/\psi}^+ m_{J/\psi} = 2f_{J/\psi} m_c \). Therefore

\[
\frac{f_{J/\psi}^+ m_c}{f_{J/\psi} m_{J/\psi}} = 2 \left( \frac{m_c}{m_{J/\psi}} \right)^2,
\]

\[\text{(2.41)}\]
In [25], the charmed quark carries a momentum fraction of order \(-m_c/m_{J/\psi}\), which is the same fraction momentum \(x\), then \(r = 2x^2\). The contributions of hard spectator interaction have been written as

\[
H_i^0(J/\psi K_{1A(1B)}) = \frac{if_B f_{1/\psi} f_{K_{1A(1B)}}}{X_0^{(BK_{1A(1B)}J/\psi)}} \frac{m_B}{\lambda_B} \left( \phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x) \right) \times \int_0^1 dy dx \left( \frac{1-y}{(1-y)(1-x)} - r_X \frac{\phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x)}{1-y}{(1-x)} \right),
\]

for \(i = 2, 3, 9\),

\[
H_i^0(J/\psi K_{1A(1B)}) = \frac{if_B f_{1/\psi} f_{K_{1A(1B)}}}{X_0^{(BK_{1A(1B)}J/\psi)}} \frac{m_B}{\lambda_B} \left( \phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x) \right) \times \int_0^1 dy dx \left( \frac{1-y}{(1-y)(1-x)} - r_X \frac{\phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x)}{1-y}{(1-x)} \right),
\]

for \(i = 2, 3, 7\). The transverse hard spectator terms are

\[
H^\perp_i(J/\psi K_{1A(1B)}) = \frac{2if_B f_{1/\psi} f_{K_{1A(1B)}}}{m_B X_0^{(BK_{1A(1B)}J/\psi)}} \frac{m_B}{\lambda_B} \frac{1}{1-z} \left( \phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x) \right) \times \int_0^1 dy dx \frac{1-y}{(1-y)^2 x},
\]

\[
H^\perp_i(J/\psi K_{1A(1B)}) = \frac{2if_B f_{1/\psi} f_{K_{1A(1B)}}}{m_B X_0^{(BK_{1A(1B)}J/\psi)}} \frac{m_B}{\lambda_B} \frac{1}{1-z} \left( \phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x) \right) \times \int_0^1 dy dx \frac{1-y-x}{(1-y)^2 x},
\]

for \(i = 2, 3, 9\), and

\[
H^\perp_i(J/\psi K_{1A(1B)}) = \frac{2if_B f_{1/\psi} f_{K_{1A(1B)}}}{m_B X_0^{(BK_{1A(1B)}J/\psi)}} \frac{m_B}{\lambda_B} \frac{1}{1-z} \left( \phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x) \right) \times \int_0^1 dy dx \frac{1-y}{(1-y)^2 x},
\]

for \(i = 2, 3, 9\), and

\[
H^\perp_i(J/\psi K_{1A(1B)}) = \frac{2if_B f_{1/\psi} f_{K_{1A(1B)}}}{m_B X_0^{(BK_{1A(1B)}J/\psi)}} \frac{m_B}{\lambda_B} \frac{1}{1-z} \left( \phi_{1}^{K_{1A(1B)}}(y) \phi_{1/\psi}(x) \right) \times \int_0^1 dy dx \frac{1-y-x}{(1-y)^2 x},
\]
In generalized factorization, the branching ratios of $B \to J/\psi K_{1A(1B)}$ behave as

$$H_i^+(J/\psi K_{1A(1B)}) = -\frac{2if_B f_{K_{1A(1B)} f_{J/\psi}} m_{J/\psi} m_{K_{1A(1B)}}}{m_B X^{(B K_{1A(1B)}, J/\psi)}_+} \frac{m_B}{\lambda_B} \frac{1}{1-z} \times \int_0^1 dydx \frac{\phi_{K_{1A(1B)}^+}^+(y) \phi_{J/\psi}^+(x)}{(1-y)^2 x^2}, \tag{2.46}$$

for $i = 5, 7$. Hence, the nonfactorizable contributions which are the same vertex corrections and hard spectator interactions have calculated by (2.28)–(2.35) for first scenario and by (2.39)–(2.46) for second one. Finally, the branching ratios of $B \to J/\psi K(1270)$ and $B \to J/\psi K(1400)$ have calculated at $\mu = m_B$ and at mixing angles $\theta = 32^\circ, 58^\circ$ in Tables 2 and 3.

3. Input Quantities

In this section, we have been introduced the essential input quantities [26]:

$$m_B = 5.28 \text{ GeV}, \quad m_{J/\psi} = 3.1 \text{ GeV}, \quad m_{K_{1A}} = 1.367 \text{ GeV}, \quad m_{K_{1B}} = 1.310 \text{ GeV},$$

$$f_B = 210 \pm 20 \text{ MeV}, \quad f_{J/\psi} = 405 \text{ MeV}, \quad f_{K_{1A}} = 0.250 \text{ GeV},$$

$$f_{K_{1B}}^+ (\mu = 1 \text{ GeV}) = 0.065 \text{ GeV},$$

$$a_{0K_{1A}} = 0.26, \quad a_{1K_{1A}} = -1.08, \quad \lambda_B (1 \text{ GeV}) = 225 \text{ MeV},$$

$$f_{K_{1B}} = -0.028 \text{ GeV}, \quad f_{K_{1B}}^+ (\mu = 1 \text{ GeV}) = 0.190 \text{ GeV}, \quad a_{1K_{1B}} = -1.95,$$

$$C_F = \frac{N^2 - 1}{2N}, \quad G_F = 1.166 \times 10^{-5}, \quad N = 3,$$

$$|V_{cb} V_{c\bar{s}}^*| = 0.039, \quad |V_{tb} V_{t\bar{s}}^*| = 0.041, \quad \tau_B = (1.638 \pm 0.011) \times 10^{-12} \text{ s}. \tag{3.1}$$

From (2.19), the magnitudes of factorizable amplitudes are obtained as

$$X_0^{(B K_{1A}, J/\psi)} = 3.54i, \quad X_+^{(B K_{1A}, J/\psi)} = -1.25i, \quad X_-^{(B K_{1A}, J/\psi)} = -5.56i,$$

$$X_0^{(B K_{1B}, J/\psi)} = -5.36i, \quad X_+^{(B K_{1B}, J/\psi)} = 1.21i, \quad X_-^{(B K_{1B}, J/\psi)} = 5.09i. \tag{3.2}$$

4. Discussion

In generalized factorization, the branching ratios of $B \to J/\psi K(1270)$ and $B \to J/\psi K(1400)$ decays are small in order to experiment, although they are in range of experimental data for these decays in mixing angles as shown in Tables 2 and 3. In this work, we have been tried to analyze these decays within “QCD factorization” method. Hence, the factorizable amplitudes in (2.24) have been calculated by using the input data of form factors, decay constants, and the non-factorizable contributions corresponding to the vertex corrections, and hard spectator interactions in $a_i^H$ have been calculated at two scenarios. First, $f_{J/\psi}$ behaves as a light meson in compared to $B$ meson ($z \to 0$). And in second, we have considered mass
of $J/\psi$ in which $m_{J/\psi}/m_b$ is held fixed. The effective coefficients $a_i^h$ have been calculated at three helicity states $h = 0, +, -$, which give three different contributions of amplitude. The branching ratios are calculated at these conditions and for the mixing angles $\theta = 32, 58$, where given in Tables 2 and 3. As it is shown, the obtained results in QCD factorization for these decays are not suitable assumption for all of the mixing angles. For $B \to J/\psi K(1270)$ and $B \to J/\psi K(1400)$, the experimental branching ratios are $(1.8 \pm 0.5) \times 10^{-3}$ and $< 5 \times 10^{-4}$, respectively [26]. In first scenario, the best obtained results for $B \to J/\psi K(1270)$ decay are $1.79 \times 10^{-3}$ at $\theta = 58^0$ and $1.42 \times 10^{-3}$ at $\theta = 32^0$; for $B \to J/\psi K(1400)$, these are $1.79 \times 10^{-3}$ at $\theta = 58^0$ and $1.42 \times 10^{-3}$ at $\theta = 32^0$, while in second one, as shown in Tables 2 and 3.

References

and the radiative decay


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