Eur. Phys. J. A **23**, 129–133 (2005)
DOI 10.1140/epja/i2004-10069-9

THE EUROPEAN
PHYSICAL JOURNAL A

P-wave charmonium decays into baryon and antibaryon pairs in quark pair creation model

R.G. Ping[1,2,a], B.S. Zou[1,2,3,4], and H.C. Chiang[1,2,3,4]

[1] CCAST (World Lab.), P.O. Box 8730, Beijing 100080, PRC
[2] Institute of High Energy Physics, Chinese Academy of Sciences, P.O. Box 918(4), Beijing 100039, PRC
[3] Institute of Theoretical Physics, Chinese Academy of Sciences, P.O. Box 2735, Beijing 100080, PRC
[4] Center of Theoretical Nuclear Physics, National Laboratory of Heavy Ion Accelerator, Lanzhou 730000, PRC

Received: 9 August 2004 /
Published online: 3 December 2004 –
Communicated by V.V. Anisovich

Abstract. A simple quark pair creation model is introduced to study exclusive decays of χ_{c_J} into baryon-antibaryon pairs. With this simple model, some exclusive decay processes, for example, $\chi_{c0} \to \mathrm{B\bar{B}}$ (B $= \Lambda, \Sigma^0, \Xi^-$) are investigated and their decay widths are evaluated by inclusion of the properties of outgoing baryons, and the results show that the strengthened decay channels $\chi_{c_J} \to \Lambda\bar{\Lambda}(J=0,2)$ are easily understood by considering only the color singlet contribution of P-wave charmonium.

PACS. 13.25.Gv Decays of J/ψ, Υ, and other quarkonia – 12.39.Jh Nonrelativistic quark model – 14.20.Jn Hyperons

1 Introduction

Decays of P-wave charmonia have continuously attracted interests of both theoretical and experimental experts [1]. Exclusive P-wave charmonium decays were once supposed as a good place to test their bound properties and decay mechanisms. In earlier studies on the mass spectroscopy of charmonia in the context of non-relativistic quark model, $\chi_{c_J}(J=0,1,2)$ states are described as P-wave bound states of $\mathrm{c\bar{c}}$ quarks, and their narrow decay widths are once phenomenologically understood by the OZI rule. At one time the non-relativistic treatment on charmonia decays was widely accepted because of there being two different energy scales involved in their decay processes, *i.e.*, the charm quark mass (m_c) and the bound energy ε ($m_\mathrm{c} \gg \varepsilon$). In this approximate scheme, the total hadronic decay widths of $\chi_{c_J}(J=0,2)$ are proportional to the square of derivative of the radial wave functions at origin $|R'(0)|^2$ [2], and the decay of χ_{c1} into two gluon is forbidden. This non-relativistic description of P-wave quarkonia is partly supported by experiments [3].

Exclusive decays of charmonia into hadrons are also supposed to be an efficacious laboratory for testing perturbative QCD. The earliest perturbative QCD treatment on exclusive decays of S-wave charmonium into the light hadrons was carried out by Lepage and Brodsky [4]. Their calculation was based on the assumption that the annihilation of the heavy quark and antiquark is a short-distance process which can be computed in perturbative theory due to the asymptotic freedom of QCD and the non-perturbative effects from light hadrons are factored into the wave functions at the origin. Due to the large-Q^2 value involved in these processes, the mass of the quarks emitted from a gluon can be ignored. Thus gluons couple with massless quarks and the helicity conservation leads to a simple selection rule [5], which states that a spin-0 particle cannot be allowed to decay into two fermions with opposite helicities. Thus it immediately forbids the decay $\chi_{c0} \to \mathrm{p\bar{p}}$ [6]. However, the experimental value for this channel does not vanish. The massless quark approximation seems to fail for these exclusive processes, and the effects arising from the quark mass should be taken into account in explaining the available experimental data. The mass correction to this “forbidden” charmonium decay has already been considered by Anselmino [6]. They found that, assigning to the quark a constituent mass rather than a current mass, one obtained non-zero values for this process. Other corrections from the non-relativistic approximation are also supposed to be essential in explanation of some forbidden decays. Non-perturbative corrections have also been carefully considered for charmonium decays [7]. Since P-wave charmonium lies between the boundary of perturbative and non-perturbative scale, it is not clear whether perturbative QCD alone should account for a correct description or other non-perturbative effects should

[a] e-mail: pingrg@mail.ihep.ac.en

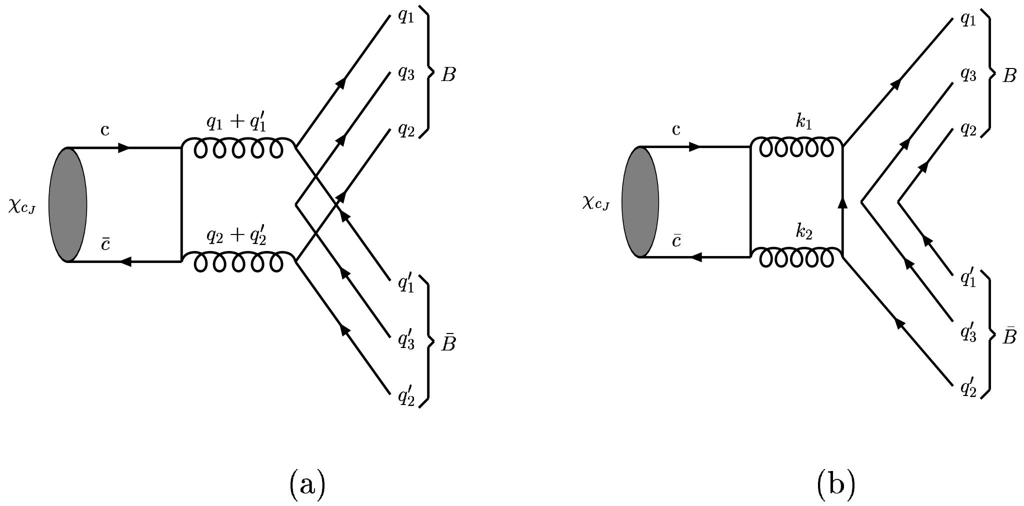

Fig. 1. The Feymman diagram to lowest order α_s for $\chi_{c_J} \to \mathrm{B\bar{B}}$ ($\mathrm{B} = \mathrm{p}, \Lambda, \Sigma^0, \Xi^-$) decays. a) One quark pair is created and b) two quark pairs are created from QCD vacuum.

still be non-negligible. Among these attempts to explain these exclusive decays, higher-order Fock states [8], gluonic contribution [9] and two-quark correlation [10] have been considered.

Nowadays, there is a renewed interest in studying decays of P-wave charmonia, not only due to the theoretical inconsistence mentioned above, but also due to the recent progress in experimental results from BES Collaboration [11]. This measurement shows that the branching ratios of $\chi_{c_J} \to \Lambda\bar{\Lambda}$ ($J = 0, 1, 2$) are almost over two times larger than $\chi_{c_J} \to \mathrm{p\bar{p}}$ ($J = 0, 1, 2$), respectively. Some theoretical studies on the χ_{c_J} decays argued that the lowest Fock state expansion, *i.e.* the color singlet mechanism, is insufficient to describe P-wave quarkonium decays, so that the next higher Fock state, *i.e.* the color octet mechanism, should be taken into account. Recently, a prediction based on this higher Fock state scheme came with a reasonable agreement with measurement on decays $\chi_{c_J} \to \mathrm{p\bar{p}}$ ($J = 1, 2$); however, the predicted width for $\chi_{c_J} \to \Lambda\bar{\Lambda}$ is about half of that for decay $\chi_{c_J} \to \mathrm{p\bar{p}}$ ($J = 1, 2$) [12].

In this paper, a quark pair creation model is presented for evaluating phenomenologically the branching ratios for decays $\chi_{c_J}(J = 0, 2) \to \mathrm{B\bar{B}}$ ($\mathrm{B} = \Lambda, \Sigma^0, \Xi^-$) by an explicit inclusion of the baryonic properties. We will show that, in this simple model, the branching ratios for decays $\chi_{c_J} \to \Lambda\bar{\Lambda}$ ($J = 0, 2$) are well reproduced, and we also compare their decay widths in terms of two different decaying mechanisms.

2 The model description and formalism

The decay widths of charmonia are assumed to be suppressed by charmonia decay via a disconnected quark line diagram. In the language of perturbative QCD, this "disconnected" decay mechanism can be described as $\mathrm{c\bar{c}}$ quarks annihilate into two gluons, which then materialize into the outgoing (anti)quarks. If the χ_{c_J} states are simply assumed as the bound states of $\mathrm{c\bar{c}}$ quarks, there are two possible decay modes which are allowed by the OZI rule as shown in fig. 1. Actually, the color factor of the mode (b) $c_\mathrm{b} = 2\sqrt{3}/9$ is twice as large as that for the mode (a) $c_\mathrm{a} = \sqrt{3}/9$. In the decay mode (a), the exclusive decay processes of $\chi_{c_J} \to \mathrm{B\bar{B}}$ are assumed via two steps. First, the $\mathrm{c\bar{c}}$ quarks annihilate into two gluons, then, the two gluons are materialized into two quark-antiquark pairs. Due to the quark-gluon coupling, another quark pair is allowed to be created from QCD vacuum with the quantum number $J^{PC} = 0^{++}$, thereafter, the three (anti)quarks hadronize into the outgoing (anti)baryon. Instead, in the decay mode (b), the two gluons only couple to light quarks to produce an intermediate state, which then decay into the outgoing baryon and antibaryon pair through combining the two created quark-antiquark pairs. In our model, we also assume that mode (a) is not allowed to mix with mode (b), nor taking the higher Fock states of P-wave charmonium states into account.

In those decay processes, the created quark pairs are all described by the quark pair creation model. Similarly to the common quark pair creation model, a phenomenological parameter, *i.e.* the strength of a created quark pair, is assumed to be equal for both light- and strange-quark pairs. Generally, the created quark pairs with any color and any flavor can be generated anywhere in space, but only those whose color-flavor wave functions and spatial wave functions overlap with those of outgoing baryons can make a contribution to the final decay width. Following the usual procedure, the Hamiltonian for the created quark pair can be defined in the modified $^3P_0$ model [13] in terms of quark and antiquark creation operators b^+ and d^+,

$$H_I = \sum_{i,j,\alpha,\beta,s,s'} \int \mathrm{d}^3k \; g_I[\bar{u}(\vec{k}, s)v(-\vec{k}, s')] \times b^+_{\alpha,i}(\vec{k}, s)d^+_{\beta,j}(-\vec{k}, s')\delta_{\alpha\beta}\hat{C}_I \,, \quad (1)$$

where $\alpha(\beta)$ and $i(j)$ are the flavor and color indices of the created quarks (antiquarks), and $u(k,s)$ and $v(k',s')$ are free Dirac spinors for quarks and antiquarks, respectively. They are normalized as $u^+(k,s)u(k,s') = v^+(k,s)v(k,s') = \delta_{ss'}$. $\hat{C}_I = \delta_{ij}$ is the color operator for $q\bar{q}$ and g_I is the strength of the decay interaction, which is assumed as a constant in these processes. In the non-relativistic limit g_I can be related to γ, the strength of the conventional $^3P_0$ model, by $g_I = 2m_q\gamma$ [13]. It is obvious that both created quarks with opposite helicities, as well as with apparel helicities can make contributions to the transition amplitude.

We attempt to evaluate the relative decay widths for the processes $\chi_{c_J} \to B\bar{B}$ ($B = \Lambda, \Sigma^0, \Xi^-$) over that of the processes $\chi_{c_J} \to p\bar{p}$, and those P-wave states are described at the level of hadron, approximately, the contribution from the $c\bar{c}$ quark dynamics, as well as from bound states are all phenomenologically parameterized into an overall constant, which only associates with χ_{c_J} quantum numbers. We feel that this treatment is reasonable at least in the limit of the non-relativistic approximation. Phenomenologically, this constant can be determined by the measurement of the decay width from one of the processes $\chi_{c_J} \to B\bar{B}$.

2.1 Mode a

The decay widths for the processes $\chi_{c_J} \to B\bar{B}$ are evaluated directly from the following expression:

$$\frac{d\Gamma(B\bar{B})}{d\Omega} = \frac{1}{32\pi^2}\overline{\sum_{s_z,s'_z}}|M_J(s_z,s'_z)|^2\frac{|\vec{P}_B|}{M^2_{\chi_{c_J}}}, \tag{2}$$

where $\vec{P}_B$ is the momentum vector for the outgoing baryon, and $M_J(s_z,s'_z)$ is the transitional amplitude of the process $\chi_{c_J} \to B\bar{B}$, which is expressed as follows:

$$\begin{aligned} M_J(s_z,s'_z) &\equiv \langle\Psi_B(q,s_z)\Psi_{\bar{B}}(q'_i,s'_z)|T_J|\chi_{c_J}\rangle \\ &= \sum_{s_i,s'_i}\int\left(\prod_{i=1}^{3}d\vec{q}_i d\vec{q}\,'_i\right) \\ &\quad\times\langle\Psi_B(q,s_z)\Psi_{\bar{B}}(q'_i,s'_z)|q_i,s_i,q'_i,s'_i\rangle \\ &\quad\times\langle q_i,s_i,q'_i,s'_i|T_J|\chi_{c_J}\rangle, \end{aligned} \tag{3}$$

where $\Psi_B(q_i,s_z)$ and $\Psi_{\bar{B}}(q'_i,s'_z)$ are the totally asymmetric wave functions for baryon and antibaryon, respectively. The hard-scattering amplitudes $\langle q_i,s_i,q'_i,s'_i|T_J|\chi_{c_J}\rangle$ for the processes $\chi_{c_J} \to qqq + \bar{q}\bar{q}\bar{q}$ are easily obtained according to standard Feynman rules.

If we neglect dynamical contributions from $c\bar{c}$ quarks, and parameterize the decay constant of χ_{c_J}, the color factor, the strong coupling constant in these decays, into an overall constant $C_J(J=0,1,2)$, one obtains:

$$\begin{aligned} &\langle q_i,s_i,q'_i,s'_i|T_0|\chi_{c0}\rangle = \\ &\frac{C_0 g_{\mu\nu}O^\mu(q_1,s_1;q'_1,s'_1)O^\nu(q_2,s_2;q'_2,s'_2)Q(q_3,s_3;q'_3,s'_3)}{(q_1+q'_1)^2(q_2+q'_2)^2} \\ &\times\delta^3(q_3-q'_3) + (q_1,q'_1 \leftrightarrow q_3,q'_3) + (q_2,q'_2 \leftrightarrow q_3,q'_3), \end{aligned} \tag{4}$$

for χ_{c0} decays, where $O^\mu(q_i,s_i;q'_i,s'_i)$, $O^\nu(q_i,s_i;q'_i,s'_i)$, $Q(q_i,s_i;q'_i,s'_i)$ are operators defined as:

$$\begin{aligned} O^\mu(q_1,s_1;q'_1,s'_1) &= \bar{u}(q_1,s_1)\gamma^\mu v(q'_1,s'_1), \\ O^\nu(q_2,s_2;q'_2,s'_2) &= \bar{u}(q_2,s_2)\gamma^\nu v(q'_2,s'_2), \\ Q(q_3,s_3;q'_3,s'_3) &= g_I\bar{u}(q_3,s_3)v(q'_3,s'_3). \end{aligned} \tag{5}$$

where g_I is the strength of the created quark pair from QCD vacuum.

As for χ_{c2} decay we have

$$\begin{aligned} &\langle q_i,s_i,q'_i,s'_i|T_2|\chi_{c2},m\rangle = \\ &\frac{C_2\Phi_{\mu\nu}(m)O^\mu(q_1,s_1;q'_1,s'_1)O^\nu(q_2,s_2;q'_2,s'_2)Q(q_3,s_3;q'_3,s'_3)}{(q_1+q'_1)^2(q_2+q'_2)^2} \\ &\times\delta^3(q_3-q'_3) + (q_1,q'_1 \leftrightarrow q_3,q'_3) + (q_2,q'_2 \leftrightarrow q_3,q'_3), \end{aligned} \tag{6}$$

where C_2 is an overall constant and $\Phi_{\mu\nu}(m)$ is the covariant spin wave function of χ_{c2} with the helicity value m, which can be built out of the polarization vector from the relation

$$\Phi_{\mu\nu}(m) = \sum_{m_1,m_2}\langle 1m_1,1m_2|2m\rangle\,\epsilon_\mu(m_1)\epsilon_\nu(m_2), \tag{7}$$

where ϵ_μ is the polarization vector for the spin-1 particles.

2.2 Mode b

The variables involved in this decay mode are defined as follows:

$$\begin{aligned} &\vec{q}_i = -\vec{q}\,'_i; \qquad q_i^0 = q_i'^0; \\ &k_1 + k_2 = (M_{\chi_c},\vec{0}); \qquad k_1 - k_2 = (0,2\vec{k}). \end{aligned} \tag{8}$$

The transitional amplitudes for $\chi_{c_J} \to B\bar{B}$ are expressed as follows:

$$\begin{aligned} M_J(s_z,s'_z) &= \sum_{s_i,s'_i}\int\left(\prod_{i=1}^{3}d\vec{q}_i\right) \\ &\quad\times\langle\Psi_B(q,s_z)\Psi_{\bar{B}}(q'_i,s'_z)|q_i,s_i,q'_i,s'_i\rangle \\ &\quad\times\langle q_i,s_i,q'_i,s'_i|T_J|\chi_{c_J}\rangle, \end{aligned} \tag{9}$$

where $\Psi_B(q_i,s_z)$ and $\Psi_{\bar{B}}(q'_i,s'_z)$ are the totally asymmetric functions for baryon and antibaryon, respectively. We define an operator $A^{\mu\nu}$

$$\begin{aligned} &A^{\mu\nu}(q_1s_1,q'_1s'_1) \\ &\quad= \bar{u}(q_1s_1)\gamma^\mu\frac{\not{q}_1-\not{k}_1+m}{(q_1-k_1)^2-m^2}\gamma^\nu v(q'_1s'_1). \end{aligned} \tag{10}$$

Then the hard-scattering amplitude for the decay $\chi_{c0} \to B\bar{B}$ reads:

$$\begin{aligned} &\langle q_i,s_i,q'_i,s'_i|T_0|\chi_{c_0}\rangle \\ &= \int\frac{d^3\vec{k}}{(2\pi)^3 2k_1^0}C_0 g_{\mu\nu}A^{\mu\nu}(q_1s_1,q'_1s'_1)\bar{u}(q_2s_2)v(q'_2s'_2) \\ &\quad\times\bar{u}(q_3s_3)v(q'_3s'_3)\frac{1}{k_1^2+i\epsilon}\frac{1}{k_2^2+i\epsilon} + (1\leftrightarrow 3) + (2\leftrightarrow 3) \\ &= -\int\frac{|\vec{k}_1|^2 d\Omega_k}{16\pi^2}C_0\bar{u}(q_1s_1)\gamma^\mu\frac{2\not{k}_1-2\not{q}_1+4m}{(q_1-k_1)^2-m^2} \\ &\quad\times\gamma^\nu v(q_1)\bar{u}(q_2s_2)v(q'_2s'_2)\bar{u}(q_3s_3)v(q'_3s'_3) \\ &\quad+(1\leftrightarrow 3)+(2\leftrightarrow 3). \end{aligned} \tag{11}$$

In the above equation, we have used the on-shell approximation for the gluonic propagator. *i.e.* $1/(k_1^2+i\epsilon)(k_2^2+i\epsilon) \to -2\pi^2\delta(k_1^2)\delta(k_2^2)$.

Similarly, the hard-scattering amplitude of the $\chi_{c2} \to \mathrm{B\bar{B}}$ decay is given by

$$\begin{aligned}
&\langle q_i, s_i, q_i', s_i'|T_2|\chi_{c_2}, m\rangle \\
&= -\int \frac{|\vec{k}_1|^2 \mathrm{d}\Omega_k}{16\pi^2} C_2 \Phi_{\mu\nu} A^{\mu\nu} \bar{u}(q_2 s_2) v(q_2' s_2') \bar{u}(q_3 s_3) \\
&\quad \times v(q_3' s_3') + (1 \leftrightarrow 3) + (2 \leftrightarrow 3), \qquad (12)
\end{aligned}$$

where $\Phi_{\mu\nu}$ is defined as in eq. (7).

3 Numerical results

For cancellation of the overall constant dependence, we evaluate the decay width ratio $\Gamma(\chi_{c_J} \to \mathrm{B\bar{B}})/\Gamma(\chi_{c_J} \to \mathrm{p\bar{p}})$. Therefore, the inclusion of the properties for the outgoing baryons is essential in our model. However, little information is known about baryon structures from perturbative QCD theory, since the hadron lies out of the asymptotic region. We simply account for the properties of bound states by explicitly including the bound-state wave functions in the naive quarks model. For example, the flavor-spin wave function of the proton may be explicitly constructed in the representation of the $SU(6)$ group if one ignores the mass difference between the u, d light quarks. The spin and flavor wave functions of the proton and antiproton are taken as:

$$\Psi_{\mathrm{SF}}^{\mathrm{p}} = \Psi_{\mathrm{SF}}^{\bar{\mathrm{p}}} = \frac{1}{\sqrt{2}}(\chi^\rho \phi^\rho + \chi^\lambda \phi^\lambda), \qquad (13)$$

where $\chi^\rho(\phi^\rho)$ and $\chi^\lambda(\phi^\lambda)$ are the mixed-symmetry pair spin(isospin) wave functions. We assume that the spatial distribution for constituent quarks (u, d) in a baryon can be described by a simple harmonic-oscillator eigenfunction in their center-of-mass (c.m.) system, *i.e.*

$$\phi_{\mathrm{p}}(\vec{k}_\rho, \vec{k}_\lambda) = \frac{1}{(\pi\beta)^{3/2}} e^{-(\vec{k}_\lambda^2 + \vec{k}_\rho^2)/2\beta}, \qquad (14)$$

where $\beta = m\omega$ is the harmonic-oscillator parameter and $\vec{k}_\rho, \vec{k}_\lambda$ are defined as

$$\begin{aligned}
\vec{k}_\rho &= \frac{1}{\sqrt{6}}(\vec{k}_1 + \vec{k}_2 - 2\vec{k}_3), \\
\vec{k}_\lambda &= \frac{1}{\sqrt{2}}(\vec{k}_1 - \vec{k}_2), \qquad (15)
\end{aligned}$$

where $\vec{k}_1$, $\vec{k}_2$ and $\vec{k}_3$ are the momenta for the three quarks in the c.m. system.

In case of hyperons, one can construct the spin-flavor wave functions analogously to that of the proton if one ignores the mass differences between light quarks and strange quarks. However, because this mass difference is substantially compared to the average quark momentum, it is advantageous to use a basis that makes explicit $SU(3)_F$ symmetry breaking under exchange of unequal-mass quarks. The flavor wave functions for strange baryons are taken as

$$\begin{aligned}
\phi_\Lambda &= \frac{1}{\sqrt{2}}(ud - du)s, \\
\phi_\Sigma &= \frac{1}{\sqrt{2}}(ud + du)s, \\
\phi_\Xi &= ssd, \qquad (16)
\end{aligned}$$

and the construction of the spin wave functions χ proceeds analogously to that of the flavor wave functions. The spatial wave function of the hyperon is chosen to be of the same form as the proton's except that the unequal-mass quark can be assigned in the c.m. system and it reads

$$\phi_{\mathrm{Y}}(\vec{k}_\rho, \vec{k}_\lambda) = \frac{1}{(\pi^2\beta_\rho\beta_\lambda)^{3/4}} e^{-\left(\frac{\vec{k}_\rho^2}{2\beta_\rho} + \frac{\vec{k}_\lambda^2}{2\beta_\lambda}\right)}, \qquad (17)$$

where $\beta_\rho = (3km)^{1/2}$ and m is the identical quark mass, and $\beta_\lambda = (3km_\lambda)^{1/2}$ with $m_\lambda = 3mm_3/(2m+m_3) > m$, and m_3 is the unequal quark mass. Adopting these relations we relate $\beta_\rho(\Lambda,\Sigma,\Xi)$ and $\beta_\lambda(\Lambda,\Sigma,\Xi)$ to β, the harmonic-oscillator parameter of proton by relations: $\beta_\rho(\Lambda,\Sigma) = \beta, \beta_\rho(\Xi) = \sqrt{m_{\mathrm{s}}/m_{\mathrm{u}}}\beta$ and $\beta_\lambda(\Lambda,\Sigma,\Xi) = \sqrt{m_\lambda/m}\beta_\rho(\Lambda,\Sigma,\Xi)$.

In the laboratory system of χ_{c_J}, the outgoing baryons move very fast. For example, the momentum of the outgoing proton for the decay $\chi_{c0} \to \mathrm{p\bar{p}}$ is about $1.427\,\mathrm{GeV}/c$. However, the wave functions of baryons in the nonrelativistic quark model are given in their c.m. systems. One has to make a Lorentz transformation of the wave functions of baryons from their c.m. system to the laboratory system. In general, this correction involves two aspects, one is the Lorentz boost of the spatial wave functions of baryons from their c.m. systems to the χ_{c_J} rest system; the other is the Melosh rotation of quark spinors. For simplicity, we feel that it may be a reasonable approximation to ignore the effects from the Melosh rotation of quark spinors and only perform the Lorentz boosts for spatial wave function [14] *i.e.*

$$\Phi_{\mathrm{B}}(\vec{q}_\rho, \vec{q}_\lambda) = \left|\frac{\partial(\vec{k}_\rho, \vec{k}_\lambda)}{\partial(\vec{q}_\rho, \vec{q}_\lambda)}\right|^{1/2} \Phi_{\mathrm{B}}(\vec{k}_\rho, \vec{k}_\lambda), \qquad (18)$$

where $\vec{q}_\rho$ and $\vec{q}_\lambda$ have a similar form as eq. (15) except that $\vec{k}_1$, $\vec{k}_2$ and $\vec{k}_3$ are replaced by three momenta of the outgoing quarks $\vec{q}_1$, $\vec{q}_2$ and $\vec{q}_3$ in laboratory system, respectively.

In our calculation, there are four parameters to be determined, *i.e.* the overall constant C_J, the light-quark mass $m_{\mathrm{u,d}}$, the strange-quark mass m_{s} and the harmonic parameter β. The constant C_J can be determined by using decays $\chi_{c_J} \to \mathrm{p\bar{p}}$ $(J = 0, 1, 2)$. In the relativized quark model [15], the light-quark mass and the harmonic parameter are, respectively, taken as $m_{\mathrm{u,d}} = 0.22\,\mathrm{GeV}$ and $\beta = 0.16\,\mathrm{GeV}^2$. We choose these parameters as preferable values in our calculation. In various quark models, the

Table 1. The ratio of decay widths $\Gamma(\chi_{cJ} \to \mathrm{B\bar{B}})/\Gamma(\chi_{cJ} \to \mathrm{p\bar{p}})$. The central values correspond to the choice for parameters $m_\mathrm{s} = 0.37\,\mathrm{GeV}$, $\beta = 0.16\,\mathrm{GeV}^2$, the uncertainty comes from the adjustment of the strange-quark mass m_s from 0.32 GeV to 0.42 GeV. Measured values are taken from [11].

$\frac{\Gamma(\chi_{cJ}\to\mathrm{B\bar{B}})}{\Gamma(\chi_{cJ}\to\mathrm{p\bar{p}})}$	Mode a	Mode b	Measured values
$\frac{\Gamma(\chi_{c0}\to\Lambda\bar{\Lambda})}{\Gamma(\chi_{c0}\to\mathrm{p\bar{p}})}$	4.1 ± 0.8	0.97 ± 0.25	2.14 ± 0.26
$\frac{\Gamma(\chi_{c0}\to\Sigma^0\bar{\Sigma}^0)}{\Gamma(\chi_{c0}\to\mathrm{p\bar{p}})}$	1.10 ± 0.10	0.82 ± 0.18	–
$\frac{\Gamma(\chi_{c0}\to\Xi^-\bar{\Xi}^+)}{\Gamma(\chi_{c0}\to\mathrm{p\bar{p}})}$	0.95 ± 0.30	0.42 ± 0.05	–
$\frac{\Gamma(\chi_{c2}\to\Lambda\bar{\Lambda})}{\Gamma(\chi_{c2}\to\mathrm{p\bar{p}})}$	2.10 ± 0.11	0.60 ± 0.05	3.37 ± 1.70
$\frac{\Gamma(\chi_{c2}\to\Sigma^0\bar{\Sigma}^0)}{\Gamma(\chi_{c2}\to\mathrm{p\bar{p}})}$	5.40 ± 1.10	0.59 ± 0.05	–
$\frac{\Gamma(\chi_{c2}\to\Xi^-\bar{\Xi}^+)}{\Gamma(\chi_{c2}\to\mathrm{p\bar{p}})}$	0.70 ± 0.30	0.22 ± 0.09	–

difference between a strange-quark mass and a light-quark mass is always chosen within the range 0.12–0.22 GeV. Using these parameters, we calculate the ratio of the decay widths for the processes $\Gamma(\chi_{cJ} \to \mathrm{B\bar{B}})$ to $\Gamma(\chi_{cJ} \to \mathrm{p\bar{p}})$, as shown in table 1. The central values are related to our preferred parameters $m_\mathrm{u,d} = 0.22\,\mathrm{GeV}$, $m_\mathrm{s} = 0.37\,\mathrm{GeV}$, $\beta = 0.16\,\mathrm{GeV}^2$. The change of the strange-quark mass m_s within the range 0.32–0.42 GeV gives the uncertainty of central values. From these results, it seems that the decay mechanism of mode (b) is unlikely possible, since in this mode, the decay widths for $\chi_{cJ} \to \mathrm{p\bar{p}}$ $(J = 0, 1, 2)$ are the largest ones among these decays. As for the decay mechanism in mode (a), we find that these decays, $\chi_{c0} \to \Lambda\bar{\Lambda}$, $\chi_{c2} \to \Lambda\bar{\Lambda}$, $\Sigma^0\bar{\Sigma}^0$, are strengthened. Although the calculated width for $\chi_{c0} \to \Lambda\bar{\Lambda}$ is still larger than the measured one, we find that in this simple model, the measured values for decays $\chi_{cJ} \to \Lambda\bar{\Lambda}$ $(J = 0, 2)$ can be understandable only by taking into account the color singlet contribution in χ_{cJ} decays. We expect that in the near future other strengthened decay channels are to be measured on BESII to test our results.

To conclude, we present a simple quark pair creation model to study the χ_{cJ} exclusive decays. With this model, the limits of the helicity conservation rule can be removed, so that some "forbidden" decay processes, *i.e.* $\chi_{c0} \to \mathrm{B\bar{B}}$ can be investigated. Phenomenologically, this model is equivalent to the quark mass correction to the Brodsky approach. The full decay widths of the processes $\chi_{cJ} \to \mathrm{B\bar{B}}$ $(J = 0, 2)(\mathrm{B} = \Lambda, \Sigma^0, \Xi^-)$ are evaluated by explicitly including of the baryonic properties, and the results show that the strengthened decay channels $\chi_{cJ} \to \Lambda\bar{\Lambda}$ $(J = 0, 2)$ are understood by including only the color singlet contribution of the P-wave charmonia.

This work is partly supported by the National Natural Science Foundation of China under grants Nos. 10225525, 10055003 and by the Chinese Academy of Sciences under project No. KJCX2-SW-N02. The author R.G. Ping thanks C.Z. Yuan, Z.Y. Zhang and P.N. Shen for useful discussions and comments.

References

1. L. Kopke, N. Wermes, Phys. Rep. **174**, 67 (1989).
2. V.A. Novirov *et al.*, Phys. Rep. C **41**, 1 (1978).
3. Particle Data Group, Phys. Rev. D **66**, 10001 (2002).
4. G.P. Lepage, S.J. Brodsky, Phys. Rev. D **22**, 2157 (1980).
5. S.J. Brodsky, G.P. Lepage, Phys. Rev. D **24**, 2848 (1981).
6. M. Anselmino, F. Caruso, F. Murgia, Phys. Rev. D **42**, 3218 (1990); M. Anselmino, F. Murgia, Phys. Rev. D **47**, 3977 (1983); **50**, 2321 (1994).
7. P. Moxhay, J.L. Rosner, Phys. Rev. D **28**, 1132 (1983); R. Mcclary, N. Byers, Phys. Rev. D **28**, 1692 (1983).
8. M. Benayoun, V.L. Chernyak, I.R. Zhitnitsky, E. Braaten, G.P. Lepage, Phys. Rev. D **46**, R1914 (1992).
9. S.J. Brodsky, G.P. Lepage, S.F. Tuan, Phys. Rev. Lett. **59**, 621 (1982); M. Anselmino, M. Genovese, E. Predazzi, Phys. Rev. D **44**, 1597 (1991).
10. M. Anselmino, F. Caruso, S. Forte, Phys. Rev. D **44**, 1438 (1991); M. Anselmino, F. Caruso, S. Joffily, J. Soares, Mod. Phys. Lett. A **6**, 1415 (1991).
11. J.Z. Bai *et al.*, Phys. Rev. D **67**, 112001 (2003).
12. S.M.H. Wong, Eur. Phys. J. C **14**, 643 (2000).
13. E.S. Ackleh, T. Barnes, E.S. Swanson, Phys. Rev. D **54**, 6811 (1996).
14. R.G. Ping, H.C. Chiang, B.S. Zou, Phys. Rev. D **66**, 054020 (2002); R.G. Ping *et al.*, Chin. Phys. Lett. **19**, 1592 (2002).
15. S. Godfrey, N. Isgur, Phys. Rev. D **32**, 189 (1985).