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J/ψ Photoproduction in pp and $p\bar{p}$ Collisions

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Abstract: The production of J/ψ originating from photoproduction processes in pp and $p\bar{p}$ collisions at leading order is calculated. The color singlet and color octet mechanisms for heavy quarkonium production are reviewed within nonrelativistic QCD, and be used to deal with the direct photon and resolved photon processes respectively. Comparing with the leading order results of J/ψ production, the numerical results show that the modification of photoproduction processes for J/ψ production become obvious at large p_T region.

Key word: photoproduction; large transverse momentum region; heavy quarkonium
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1 Introduction

Clarifying the production of heavy quarkonium in various processes is interesting particularly for the test on QCD and identifying the QGP sources. Bodwin, Braaten, and Lepage[1] constructed a rigorous progress on factorization between perturbative on-shell parton level process and the quarkonium bound state dynamics. The earlier color singlet model which could not explain all the results observed at Fermilab Tevatron energies has been superseded by the color octet mechanism based on nonrelativistic QCD (NRQCD). Based on perturbative QCD (pQCD) and NRQCD, the yields of J/ψ in pp, $p\bar{p}$, and ep collisions have been investigated by many authors[2−12].

The photoproduction processes play a fundamental role in the ep deep inelastic scattering at Hadron Electron Ring Accelerator (HERA). In photoproduction processes of the ep deep inelastic scattering, a high energy photon emitted by the incident electron interacts directly with the proton via the interaction of $\gamma p \to X$. Besides, the uncertainty principle allows the high energy photon to fluctuate into a quark-antiquark pair for a short time, and then interacts with the partons of the proton. In such interactions the resolved photon can be interpreted as an extended object consisting of quarks and gluons. The interactions are the so-called resolved photoproduction processes.

We generalize the photoproduction mechanism to the heavy quarkonium production in Nucleus-Nucleus collisions. Charged partons of the incident nucleon also can emit high energy photons (and resolved photons) in relativistic Nucleus-Nucleus collisions. In relativistic heavy ion collisions, the intense heavy-ion beams represent a prolific source of quasi-real photons. Hence it is a wonderful laboratory for the extensive studies of the photoproduction physics.

In this paper, we study the direct and resolved photoproduction processes for J/ψ production in pp and $p\bar{p}$ collisions. In Sec. 2 we briefly review the color singlet mechanism and the color octet mechanism of J/ψ production. The direct and resolved

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photoproduction processes are presented in Sec. 3. We discuss the p_T distribution of the J/ψ production at $\sqrt{s}=7$ TeV pp collisions and $\sqrt{s}=1.8$ TeV and 1.96 TeV p$\bar{\text{p}}$ collisions, and discuss the comparison between photoproduction data and the leading order (LO) results. Sec. 4 is the summary.

2 Formalism

Although a quark-antiquark pair in a color octet configuration repel each other, a quark-gluon pair attract each other. A octet state of quarks can bind to some gluons whose attraction to both quarks overcomes the repulsion between quarks[13]. This picture can be derived clearly from the NRQCD. Let us begin from the QCD Lagrangian

$$\mathcal{L}=\mathcal{L}_{\text{light}}+\mathcal{L}_{\text{heavy}}, \tag{1}$$

where $\mathcal{L}_{\text{light}}$ is the Lagrangian that describe gluons and light quarks,

$$\mathcal{L}_{\text{light}}=-\frac{1}{2}\text{tr}G_{\mu\nu}G^{\mu\nu}+\sum_{n_f}\bar{q}\text{i}\not{D}q, \tag{2}$$

and $\mathcal{L}_{\text{heavy}}$ describes the heavy quarks,

$$\mathcal{L}_{\text{heavy}}=\bar{\Psi}(\text{i}\gamma^\mu D_\mu-M)\Psi, \tag{3}$$

where $G^{\mu\nu}$ is the gluon field strength tensor, q is the Dirac spinor field for a light quark, n_f is the flavors of light quarks, the gauge-convariant derivative $D^\mu=\partial^\mu+\text{i}gA^\mu$, $A^\mu=(A_0, \boldsymbol{A})$, M is the heavy quark mass. Shifting the $\mathcal{L}_{\text{heavy}}$ to a frame in which the heavy quarks are almost rest via the transformation

$$\Psi=\text{e}^{-\text{i}Mt}\widetilde{\Psi}, \tag{4}$$

we have

$$\mathcal{L}_{\text{heavy}}=\bar{\widetilde{\Psi}}(\text{i}\gamma^\mu D_\mu-M(\gamma^0-1))\widetilde{\Psi}. \tag{5}$$

After integrating out modes of momentum greater than cutoff Λ that is much less than M and using Foldy-Wouthuysen-Tani transformation[14], we get the NRQCD Lagrangian

$$\mathcal{L}_{\text{heavy}}=\mathcal{L}_0+\delta\mathcal{L}, \tag{6}$$

where

$$\mathcal{L}_0=\psi^\dagger(\text{i}D_t+\frac{\boldsymbol{D}^2}{2M})\psi+\chi^\dagger(\text{i}D_t-\frac{\boldsymbol{D}^2}{2M})\chi, \tag{7}$$

and the bilinear correction terms is

$$\begin{aligned}\delta\mathcal{L}_{\text{bilinear}}=&\frac{c_1}{8M^2}[\psi^\dagger(\boldsymbol{D}^2)^2\psi-\chi^\dagger(\boldsymbol{D}^2)^2\chi]+\\&\frac{c_2}{8M^2}[\psi^\dagger(\boldsymbol{D}\cdot g\boldsymbol{E}-g\boldsymbol{E}\cdot\boldsymbol{D})]\psi+\\&\chi^\dagger(\boldsymbol{D}\cdot g\boldsymbol{E}-g\boldsymbol{E}\cdot\boldsymbol{D})\chi]+\\&\frac{c_3}{8M^2}[\psi^\dagger(\text{i}\boldsymbol{D}\times g\boldsymbol{E}-g\boldsymbol{E}\times\text{i}\boldsymbol{D})\cdot\sigma\psi+\\&\chi^\dagger(\text{i}\boldsymbol{D}\times g\boldsymbol{E}-g\boldsymbol{E}\times\text{i}\boldsymbol{D})\cdot\sigma\chi]+\\&\frac{c_4}{2M}[\psi^\dagger(g\boldsymbol{B}\cdot\sigma)\psi-\chi^\dagger(g\boldsymbol{B}\cdot\sigma)\chi]+\\&\cdots,\end{aligned} \tag{8}$$

where E, B are respectively the electric and magnetic components of $G^{\mu\nu}$, $E^i=G^{0i}$, $B^i=\frac{1}{2}\varepsilon^{ijk}G^{jk}$. g is the QCD coupling constant, ψ is the Pauli spinor field that annihilates a heavy quark, χ is the Pauli spinor field that creates a heavy antiquark, D_t and $\boldsymbol{D}$ are the time and space components of D^μ, respectively. The effects of the high-energy modes which have been integrated out were encoded into the coefficients c_i and new local interactions of the NRQCD Lagrangian $\delta\mathcal{L}$[1]

$$\delta\mathcal{L}_{\text{4-fermion}}=\sum_n\frac{f_n(\Lambda)}{M^{d_n-4}}\mathcal{O}_n(\Lambda), \tag{9}$$

where $\mathcal{O}_n$ are local four-fermion operators.

In Coulomb gauge ($\nabla\cdot\boldsymbol{A}=0$), dynamical gluons are described by the vector potential $\boldsymbol{A}$, the Lagrangian $\mathcal{L}_{\text{heavy}}$ can be reorganized as an expansion in powers of v by using the rescaling velocity-scaling rules[15-16]. The terms in $\mathcal{L}_0+\delta\mathcal{L}_{\text{4-bilinear}}$ which are of order v^0 are

$$\begin{aligned}\mathcal{L}'_0=&\psi^\dagger(\text{i}\partial_t-gA_0+\frac{\nabla^2}{2M})\psi+\\&\chi^\dagger(\text{i}\partial_t-gA_0-\frac{\nabla^2}{2M})\chi,\end{aligned} \tag{10}$$

the terms which are of order v^1 are

$$\mathcal{L}_1=-\frac{1}{M}\psi^\dagger(\text{i}g\boldsymbol{A}\cdot\nabla)\psi+\frac{c_4}{2M}\psi^\dagger(\nabla\times g\boldsymbol{A})\cdot\sigma\psi$$

$$+\text{charge conjugate terms}\ , \tag{11}$$

and the terms which are of order v^2 are

$$\mathscr{L}_2 = -\frac{1}{2M}\psi^\dagger(g\boldsymbol{A})^2\psi + \frac{c_1}{8M^3}\psi^\dagger(\nabla^2)^2\psi - \frac{c_2}{8M^2}\psi^\dagger(\nabla^2 gA_0)\psi - \frac{c_3}{4M^2}\psi^\dagger(\nabla gA_0)\times\nabla\cdot\boldsymbol{\sigma}\psi + \frac{c_4}{2M}\psi^\dagger(\mathrm{i}g\boldsymbol{A}\times g\boldsymbol{A})\cdot\sigma\psi + \text{charge conjugate terms}\ . \tag{12}$$

We can see that, the dynamical gluons enter into the Lagrangian that are of order v^1 and higher orders. The Fock state expansion of the S-wave ortho-charmonium state ψ_c in Coulomb gauge[3] is:

$$|\psi_c\rangle = O(1)|c\bar{c}[^3S_1^{(1)}]\rangle + O(v)|c\bar{c}[^3P_J^{(8)}]g\rangle + O(v^2)|c\bar{c}[^3S_1^{(1,8)}]gg\rangle + O(v^2)|c\bar{c}[^1S_0^{(8)}]g\rangle + O(v^2)|c\bar{c}[^3D_J^{(1,8)}]gg\rangle + \cdots. \tag{13}$$

The superscripts (1) and (8) represent the singlet and octet state of the $c\bar{c}$ pair, respectively.

If we neglect the relativistic corrections in which v is taken into account, we get the color-singlet mechanism[17−18]. The quark and antiquark fly out from the initial collision point in nearly parallel direction and eventually hadronize into a J/ψ bound state by exchanging many soft gluons at a long distance. On the other hand, the relativistic corrections are considered in the color-octet mechanism[1−3, 19−21]. The high energy collision creates a $c\bar{c}$ pair in a color-octet configuration by coupling the dynamical gluons. Far away from the collision point, the color octet $|c\bar{c}g\cdots\rangle$ state transforms into a color singlet J/ψ bound state by emitting some long wavelength gluons which bleed off the color of the state but carry away virtually no energy or momentum when the $c\bar{c}$ pair already expanded to the charmonium size. Both of the color singlet mechanism and color octet mechanism can be factorized as proceeding in two parts. The first is the short-distance part which produces a $c\bar{c}$ pair and can be described by pQCD. The second is the binding of the $c\bar{c}$ pair into the J/ψ state which is called the long-distance part and the effects are consolidated into nonperturbative NRQCD matrix elements $\langle 0|\mathscr{O}_{1,8}^{J/\psi}(^{2S+1}L_J)|0\rangle$ whose values must be determined from experiment or the lattice computation.

In general, the on-shell amplitude can be expressed as

$$\sum_{ij}\langle 3i\bar{3}\bar{j}|1,8a\rangle\mathscr{A}\left(a+b\to c^i\left(\frac{P}{2}+q\right)+\bar{c}^j\left(\frac{P}{2}-q\right)+d\right) = \bar{u}\left(\frac{P}{2}+q, s_1\right)\mathscr{O}(P,q)v\left(\frac{P}{2}-q, s_2\right). \tag{14}$$

Therefore, the scattering amplitude of the processes $a+b\to c\bar{c}[^{2S+1}L_J^{(1,8a)}]+d$ are given by

$$\begin{aligned}&\mathscr{M}(a+b\to c\bar{c}[^{2S+1}L^{(1,8a)}{}_J](P)+d) = \\ &\sum_{L_ZS_Z}\sum_{s_1s_2}\sum_{ij}\int\frac{d^3\boldsymbol{q}}{(2\pi)^3 2q^0}\delta\left(q^0-\frac{|\boldsymbol{q}^2|}{2M}\right)\times \\ &\psi_{LL_Z}(\boldsymbol{q})\times\langle s_1;s_2|SS_Z\rangle\times \\ &\langle LL_Z;SS_Z|JJ_Z\rangle\langle 3i\bar{3}\bar{j}|1,8a\rangle\times \\ &\mathscr{A}\left(a+b\to c^i\left(\frac{P}{2}+q\right)+\bar{c}^j\left(\frac{P}{2}-q\right)+d\right) \\ &=\sum_{L_ZS_Z}\int\frac{d^3\boldsymbol{q}}{(2\pi)^3 2q^0}\delta\left(q^0-\frac{|\boldsymbol{q}^2|}{2M}\right)\psi_{LL_Z}(\boldsymbol{q})\times \\ &\langle LL_Z;SS_Z|JJ_Z\rangle\mathrm{tr}[\mathscr{O}(P,q)P_{SS_Z}(P,q)],\end{aligned} \tag{15}$$

where P is the momentum of $Q\bar{Q}$, q is the relative momentum of quark and antiquark, $\psi_{LL_Z}(\boldsymbol{q})$ is the non-relativistic bound-state wave function in a given angular momentum state. $\langle s_1;s_2|SS_Z\rangle$, $\langle LL_Z;SS_Z|JJ_Z\rangle$ and $\langle 3i\bar{3}\bar{j}|1,8a\rangle$ are the Clebsch-Gordan coefficients for spin angular momentum, total angular momentum and color $SU(3)$ respectively. $P_{SS_Z}(P,q)$ is the spin projection operator[22−23]

$$P_{SS_Z}(P,q)_{ij} = \sum_{s_1s_2}\langle s_1;s_2|SS_Z\rangle\, v_i\left(\frac{P}{2}+q\right)\bar{u}_J\left(\frac{P}{2}-q\right)\ , \tag{16}$$

Expanding $\mathrm{tr}[\mathscr{O}(P,q)P_{SS_Z}(P,q)]$ to zeroth and first order in q respectively, and treating q to be small, we get the S-wave and P-wave amplitudes[22, 24]. After integrating over q, we have

$$\mathscr{M}(\mathrm{a}+\mathrm{b}\to c\bar{c}[^{2S+1}S_J^{(1,8a)}]+\mathrm{d}) = \sqrt{C_S}\ \mathrm{tr}[\mathscr{O}(P,0)P_{SS_Z}(P,0)], \quad (17)$$

and

$$\mathscr{M}(\mathrm{a}+\mathrm{b}\to c\bar{c}[^{2S+1}P_J{}^{(1,8a)}]+\mathrm{d}) = \sqrt{C_P}\sum_{L_Z S_Z}\varepsilon_\alpha^*(L_Z)\langle LL_Z;SS_Z|JJ_Z\rangle \times \mathrm{tr}[\mathscr{O}^\alpha P_{SS_Z}+\mathscr{O}P^\alpha_{SS_Z}]_{q=0}, \quad (18)$$

where

$$\mathscr{O}^\alpha=\frac{\partial}{\partial q^\alpha}\mathscr{O}(P,q),\ P^\alpha_{SS_Z}=\frac{\partial}{\partial q^\alpha}P_{SS_Z}(P,q). \quad (19)$$

$C_L(L=S, P)$ is the remaining factor after the integration over q:

$$C_L=\frac{\langle 0|\mathscr{O}_n{}^{J/\psi}(^{2S+1}L_J)|0\rangle}{C_n(2J+1)}, \quad (20)$$

where $C_n=2N_c$ for color singlet and N_c^2-1 for color octet.

For 2→2 subprocesses,

$$\frac{d\hat{\sigma}}{d\hat{t}}=\frac{1}{16\pi\hat{s}^2}\sum|\mathscr{M}(\mathrm{a}+\mathrm{b}\to c\bar{c}(^{2S+1}L_J^{(1,8a)})+\mathrm{d})|^2. \quad (21)$$

3　Calculation and discussions

In direct single photoproduction processes, the parton a of the incident nucleon A can emit a high energy photon, then the high energy photon interacts with the parton b of another incident nucleon B via the interactions[25−26].

$$\begin{aligned}&\gamma+\mathrm{g_b}\to c\bar{c}[^3S_1^{(1)},\ ^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{g},\\&\gamma+\mathrm{q_b}(\bar{\mathrm{q}}_\mathrm{b})\to c\bar{c}[^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{q}(\bar{\mathrm{q}}),\end{aligned} \quad (22)$$

where $J=0, 1, 2$. The p_T distribution for J/ψ photoproduction of direct single photon subprocesses is given by:

$$E\frac{d\sigma}{dp^3}(\mathrm{A}+\mathrm{B}\to J/\psi+X)= \frac{2}{\pi}\int dx_a dx_b J(x_a, x_b, z_a, x_1, x_2)G_{a/A}(x_a, Q^2)\times f_{\gamma/q_a}(z_a)G_{b/B}(x_b, Q^2)\times \frac{d\hat{\sigma}}{d\hat{t}}(\gamma+\mathrm{b}\to c\bar{c}^{(1,8)}+X)\langle\mathscr{O}_{1,8}{}^{J/\psi}\rangle, \quad (23)$$

where $d\hat{\sigma}/d\hat{t}(\gamma+\mathrm{b}\to c\bar{c}^{(1,8)}+X)$ are the partonic cross sections, $Q^2=p_T^2$, $G_{a/A}(x_a, Q^2)$ and $G_{b/B}(x_b, Q^2)$ are the parton distributions of nucleons, $f_{\gamma/q_a}(z_a)$ is the photon spectrum from the charged parton. $J(x_a, x_b, z_a, x_1, x_2)$ is the Jacobi factor,

$$J(x_a, x_b, z_a, x_1, x_2)=\frac{z_a x_a x_b}{x_a x_b - x_a x_2}, \quad (24)$$

x_a and x_b are the momentum fractions of the partons, z_a is the momentum ratio of the photon energy and the energy of the quark,

$$z_a=\frac{x_b x_1-\tau}{x_a x_b-x_a x_2}, \quad (25)$$

where $\tau=M^2/s$, $x_1=(x_T^2+4\tau)^{1/2}e^y$, $x_2=(x_T^2+4\tau)^{1/2}e^{-y}$, $x_T=2p_T/\sqrt{s}$, M is the mass of $c\bar{c}$, and $\langle\mathscr{O}^{J/\psi}_{1,8}\rangle$ are the long-distance matrix elements.

In resolved single photoproduction processes, the parton a of the incident nucleon A can emit a resolved photon, then the parton a′ of the resolved photon interacts with the parton b of another incident nucleon B via the interactions[3, 26].

$$\begin{aligned}&\mathrm{g_{a'}}+\mathrm{g_b}\to c\bar{c}[^3S_1^{(1)},\ ^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{g},\\&\mathrm{g_{a'}}+\mathrm{q_b}(\bar{\mathrm{q}}_\mathrm{b})\to c\bar{c}[^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{q}(\bar{\mathrm{q}}),\\&\mathrm{q_{a'}}(\bar{\mathrm{q}}_{\mathrm{a'}})+\mathrm{g_b}\to c\bar{c}[^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{q}(\bar{\mathrm{q}}),\\&\mathrm{q_{a'}}+\bar{\mathrm{q}}_\mathrm{b}\to c\bar{c}[^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{g},\\&\bar{\mathrm{q}}_{\mathrm{a'}}+\mathrm{q_b}\to c\bar{c}[^1S_0^{(8)},\ ^3S_1^{(8)},\ ^3P_J^{(8)}]+\mathrm{g}.\end{aligned} \quad (26)$$

The p_T distribution for J/ψ photoproduction of resolved single photon subprocesses is given by:

$$E\frac{d\sigma}{dp^3}(\mathrm{A}+\mathrm{B}\to J/\psi+X)= \frac{2}{\pi}\int dx_a dx_b dz_{a'} J(x_a, x_b, z_a, z_{a'}, x_1, x_2) G_{a/A}(x_a, Q^2) f_{\gamma/q_a}(z_a)\times G_{b/B}(x_b, Q^2)G_{p_{a'}/\gamma}(z_{a'}, Q^2)\times \frac{d\hat{\sigma}}{d\hat{t}}(\mathrm{p_{a'}}+\mathrm{b}\to c\bar{c}^{(1,8)}+X)\langle\mathscr{O}^{J/\psi}_{1,8}\rangle, \quad (27)$$

where $d\hat{\sigma}/d\hat{t}(\mathrm{p_{a'}}+\mathrm{b}\to c\bar{c}^{(1,8)}+X)$ are the partonic cross sections, $\mathrm{p_{a'}}=\mathrm{q_{a'}}, \mathrm{g_{a'}}$. $G_{p_{a'}/\gamma}(z_{a'}, Q^2)$ is the parton distribution of theresolved photon. The Jacobi factor is

$$J(x_a, x_b, z_a, z_{a'}, x_1, x_2) = \frac{z_a z_{a'} x_a x_b}{x_a x_b z_{a'} - x_a x_2 z_{a'}}. \tag{28}$$

The momentum ratio of the photon energy and the energy of the quark,

$$z_a = \frac{x_b x_1 - \tau}{x_a z_{a'} x_b - x_a z_{a'} x_2}, \tag{29}$$

and $z_{a'}$ is the momentum ratio of the parton energy and the energy of the resolved photon.

In order to make a comparison, we calculate the LO processes of J/ψ production in pp and $p\bar{p}$ collisions via the interactions

$$g + g \to c\bar{c}[{}^3S_1^{(1)}, {}^1S_0^{(8)}, {}^3S_1^{(8)}, {}^3P_J^{(8)}] + g,$$
$$g + q(\bar{q}) \to c\bar{c}[{}^1S_0^{(8)}, {}^3S_1^{(8)}, {}^3P_J^{(8)}] + q(\bar{q}),$$
$$q + \bar{q} \to c\bar{c}[{}^1S_0^{(8)}, {}^3S_1^{(8)}, {}^3P_J^{(8)}] + g. \tag{30}$$

The p_T distribution for LO subprocesses of J/ψ production is given by:

$$E\frac{d\sigma}{dp^3}(A + B \to J/\psi + X) = \frac{2}{\pi}\int dx_a J(x_a, x_b, x_1, x_2) G_{a/A}(x_a, Q^2) \times G_{b/B}(x_b, Q^2)\frac{d\hat{\sigma}}{d\hat{t}}(a + b \to c\bar{c}^{(1,8)} + X)\langle \mathcal{O}_{1,8}^{J/\psi}\rangle, \tag{31}$$

where $d\hat{\sigma}/d\hat{t}(a + b \to c\bar{c}^{(1,8)} + X)$ are the partonic cross sections, a, b=g, q, $\bar{q}$. The Jacobi factor is

$$J(x_a, x_b, x_1, x_2) = \frac{x_a x_b}{x_a - x_1}. \tag{32}$$

We now describe our theoretical input and the kinematic conditions for our numerical analysis. The values of m_q, α, the branching ratios $B(J/\psi \to ee)$ and $B(J/\psi \to \mu\mu)$ are adopted from the Particle Data Group[27]. We set the strong coupling constant

$$\alpha_s = \frac{12\pi}{(33 - 2n_f)\ln\left(\frac{Q^2}{\Lambda^2}\right)} \tag{33}$$

with the number of quark flavors $n_f = 5$ and $\Lambda = 0.2$ GeV.

In this paper, we use the long-distance NRQCD matrix elements:

$$\langle 0|\mathcal{O}_1^{J/\psi}({}^3S_1)|0\rangle = 1.16\ \text{GeV}^3,$$
$$\langle 0|\mathcal{O}_8^{J/\psi}({}^3S_1)|0\rangle = 1.06\times 10^{-2}\ \text{GeV}^3.$$

The dominated processes of heavy quarkonium photoproduction are $a + b \to c\bar{c}({}^1S_0^{(8)}, {}^3P_J^{(8)}) + d \to J/\psi + X$, where $J = 0, 1, 2$. Therefore, the values of the color octet matrix elements $\langle 0|\mathcal{O}_8^{J/\psi}({}^1S_0)|0\rangle$ and $\langle 0|\mathcal{O}_8^{J/\psi}({}^3P_0)|0\rangle$ are important in inelastic J/ψ photoproduction. However, their values are not well determined from the experiment. The authors of Ref. [3] determined.

$$\frac{\langle 0|\mathcal{O}_8^{J/\psi}({}^1S_0)|0\rangle}{3} + \frac{\langle 0|\mathcal{O}_8^{J/\psi}({}^3P_0)|0\rangle}{m_c^2} = (2.2\pm 0.5)\times 10^{-2}\ \text{GeV}^3,$$

and we have

$$0 < \langle 0|\mathcal{O}_8^{J/\psi}({}^1S_0)|0\rangle < (6.6\pm 1.5)\times 10^{-2}\ \text{GeV}^3$$
$$0 < \frac{\langle 0|\mathcal{O}_8^{J/\psi}({}^3P_0)|0\rangle}{m_c^2} < (2.2\pm 0.5)\times 10^{-2}\ \text{GeV}^3.$$

In this paper, we follow Refs. [28, 29] and set

$$\langle 0|\mathcal{O}_8^{J/\psi}({}^1S_0)|0\rangle = 3.0\times 10^{-2}\ \text{GeV}^3$$
$$\frac{\langle 0|\mathcal{O}_8^{J/\psi}({}^3P_0)|0\rangle}{m_c^2} = 1.2\times 10^{-2}\ \text{GeV}^3.$$

At leading order, the matrix elements are related by heavy-quark spin symmetry as

$$\langle 0|\mathcal{O}_8^{J/\psi}({}^3P_J)|0\rangle = (2J+1)\langle 0|\mathcal{O}_8^{J/\psi}({}^3P_0)|0\rangle.$$

We use the photon spectrum from the charged parton[30]

$$f_{\gamma/q} = e_q^2\frac{\alpha}{2\pi}\frac{1+(1-z)^2}{z}\ln\left(\frac{Q_1^2}{Q_2^2}\right), \tag{34}$$

where e_q is the charge of the quark, α is the electromagnetic coupling parameter, z is the momentum ratio of the photon energy and the energy of the quark, Q_1^2 and Q_2^2 stand for the maximum and minimum value of the momentum transfer respectively. Q_1^2 is given by $(\hat{s}/4 - m_c^2)$[30–31], Q_2^2 is determined by $(m_q z)^2/(1-z)$[32–35] so that the photon is close to being in its mass shell, m_q is the mass of the quark. We use the parton distribution of protons[36]

$$G(x, Q^2) = R(x, Q^2)p(x, Q^2), \tag{35}$$

where $R(x, Q^2)$ is the nuclear modification of the structure function[37], $p(x, Q^2)$ is the parton distribution of proton. We use the parton distribution of the resolved photon $G_{p_{a'}/\gamma}(z_{a'}, Q^2)$

$$G_{p_{a'}/\gamma}(z_{a'}, Q^2) = \frac{9\alpha}{4\pi z_{a'}} \ln \frac{Q^2}{(0.204\ \text{GeV})^2} \times \left[t^j z_{a'}^{l} (A + B\sqrt{z_{a'}} + C z_{a'}^{m}) + t^{j'} \exp\left(-E + \sqrt{E' t^k \ln \frac{1}{z_{a'}}}\right)\right](1 - z_{a'})^D, \quad (36)$$

where factors t, j, j', k, l, m, A, B, C, D, E, and E' can be found in Ref. [33].

Figs. 1～2 give the contrast between photoproduction results (dir. + res.) and LO data (LO). Fig. 1(a) and 1(b) are the results in p$\bar{\text{p}}$ collisions, others are for pp collisions. The dash line shows the LO result, the dot line represents the photoproduction data which is the sum of the direct and resolved subprocesses, and the solid line gives out the sum of them.

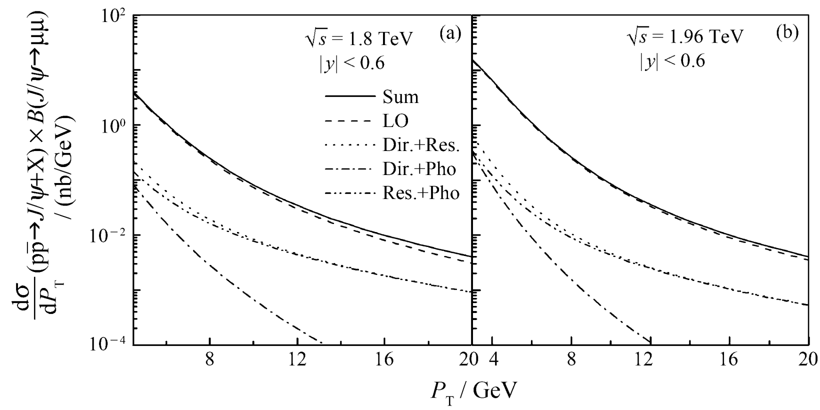

Fig. 1　The numerical results in p$\bar{\text{p}}$ collisions at $\sqrt{s}$=1.8 TeV and $\sqrt{s}$=1.96 TeV.

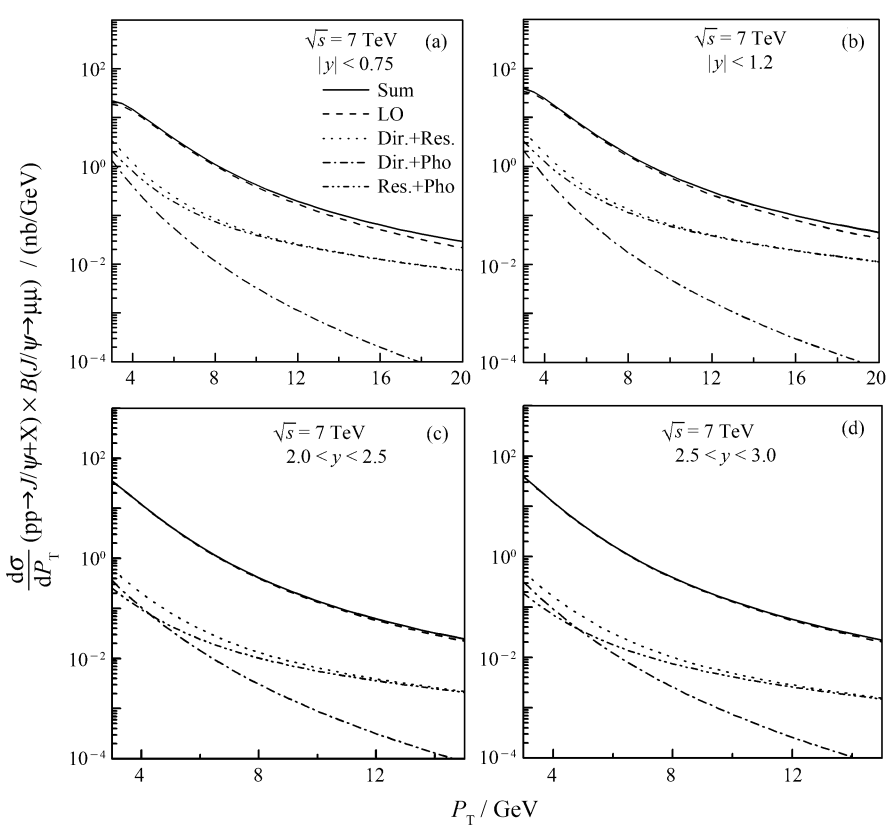

Fig. 2　The numerical results in pp collisions at $\sqrt{s}$=7 TeV.

The photoproduction corrections become obvious along with the p_T increases, this is because that the QCD coupling constant α_s reduces along with the increase of p_T such as $\alpha_s \approx 0.303$ at $p_T = 3$ GeV and $\alpha_s \approx 0.178$ at $p_T = 20$ GeV, so the total yield of J/ψ reduces along with the increase of p_T. But photoproduction is not just a pure QCD process, the production process of the high energy photons which participate in reaction are QED processes, and QED coupling constant $\alpha \approx 1/137$ keeps invariant at $p_T = 20$ GeV. Hence, the decrease amplitude of the photoproduction processes contribution along with the increase of p_T is less than the pure QCD processes such as $g + g \to J/\psi + X$, $q(\bar{q}) + g \to J/\psi + X$, and $q + \bar{q} \to J/\psi + X$.

On the other hand, the resolved singlet photon contribution dominates the J/ψ photoproduction yield and becomes more obvious in the large p_T region. The reason is that, although the number of the inelastic scattering events reduce along with the increase of p_T and so the cross section of the direct photoproduction reduces seriously in the large p_T region, the parton momentum distributions of the resolved photon are not mainly concentrated in a particular direction on p_T interval, and this furnishes a measured amount of inelastic scattering events in the large p_T region. It also can be found in the parton distribution function of the resolved photon which is directly proportional to Q^2 $(= p_T^2)$. Some authors set $Q = 2p_T$[38], but it exerts very little influence on the results, because the strong coupling constant reduces along with the increase of Q, though the photon spectrum is directly proportional to the value of Q.

In addition, from Fig. 2 we can see the photoproduction corrections to LO results with the different rapidities at $\sqrt{s_{NN}} = 7$ TeV. We can find that the correction becomes less evident when the rapidity becomes larger at the same $\sqrt{s_{NN}}$ and has the maximum in central rapidity region. The definition of rapidity means that longitudinal momentum can be neglected comparing with particle energy in central rapidity region, and the momentum transfer Q attains the maximum value. Along with the increase of rapidity, the momentum transfer becomes smaller, as the previous discussion, the correction becomes less evident.

Finally, we have not considered the fragmentation processes such as $g \to J/\psi$ and $b(\bar{b}) \to J/\psi$, moreover, we have not considered the direct double photon processes, the resolved double photon processes, and the direct and resolved mixture processes.

4 Summary

In this paper, we investigate the photoproduction of J/ψ. After reviewing the color singlet mechanism and color octet mechanism within NRQCD, the singlet-direct and singlet-resolved subprocesses have been calculated. We compare the results between photoproduction and LO production processes, the corrections become obvious at the large p_T region, on the other hand, exert less influence when the rapidity becomes larger at the same $\sqrt{s_{NN}}$. Both of them are just because that in the high momentum transfer region, the contribution of pure QCD processes reduce seriously as the decrease of the QCD coupling constant, but the decrease amplitude of the photoproduction is not very obvious. This may imply that photoproduction physics which contains QED composition have relatively obvious effect in the high p_T region.

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pp 和 $p\bar{p}$ 碰撞中 J/ψ 的光生过程

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摘要：计算了 pp 和 $p\bar{p}$ 碰撞中 J/ψ 的领头阶光生过程的产额。运用非相对论量子色动力学回顾了重夸克偶素产生的色单态机制与色八重态机制，并将它们分别用于处理直接光子过程和分解光子过程。通过与 J/ψ 产生的领头阶结果的对比可以看出，光生过程对 J/ψ 产额的修正在大横动量区域变得明显。

关键词：光生过程；大横动量区域；重夸克偶素

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