The transverse momentum ($p_t$) dependence of $J/\psi$ production in heavy ion collisions is investigated in a transport model with both initial production and continuous regeneration of charmonia. The competition between the two production mechanisms results in a $p_t$ suppression in central collisions, the gluon multi-scattering in the initial stage leads to a high $p_t$ enhancement, and the regeneration populates $J/\psi$s at low $p_t$ region and induces a minimum in $R_{AA}(p_t)$. These three phenomena are indeed observed in both 200 GeV Cu+Cu and Au+Au collisions at RHIC energy. 

a) The averaged transverse momentum square ($\langle p_T^2 \rangle$) is strongly suppressed in central collisions, which is very different from the SPS data \cite{4} where ($\langle p_T^2 \rangle$) gets saturated for central Pb+Pb collisions at energy $\sqrt{s_{NN}} \approx 17.3$ GeV;

b) The nuclear modification factor $R_{AA}$ is about 1/2 at low $p_t$ but becomes around unity at high $p_t$;

c) There exists a minimum in $R_{AA}$ in low $p_t$ region.

Can these phenomena tell us something new about the nature of $J/\psi$ production at RHIC energy? The $p_t$ broadening at SPS energy is generally attributed to the gluon multi-scattering in the initial state \cite{6} and the leakage effect in the final state \cite{7}. In a pA collision, a gluon of the proton scatters from target nucleons before it fuses with a gluon from the target to form a $J/\psi$. The gluon rescattering in the initial state is treated as a random walk in transverse momentum and the observed ($\langle p_T^2 \rangle$) is predicted to increase linearly with the mean length of the path of the incident gluon. In an AB collision, both gluons which fuse to a $J/\psi$ are affected by the rescattering. The leakage effect on $J/\psi$ production has already been considered 20 years ago \cite{1,8,9}. The anomalous suppression inside the QGP is not an instantaneous process, but takes a certain time. During this time the $J/\psi$s with high transverse momenta may leak out of the source of the anomalous suppression. As a consequence, low $p_t$ $J/\psi$s are absorbed preferentially but high $p_t$ $J/\psi$s can survive.

While charm quark production at SPS energy is expected to be small, there are more than 10 $c \bar{c}$ pairs produced in a central Au+Au collision at RHIC energy and probably more than 200 pairs at LHC energy \cite{10}. The uncorrelated charm quarks in the QGP can be recombined to form $J/\psi$s. Obviously, the regeneration will enhance the $J/\psi$ yield and alter its momentum spectrum. The regeneration approach for $J/\psi$ production at RHIC has been widely discussed with different models, such as thermal production on the hadronizaton hypersurface according to statistic law \cite{11}, the coalescence mechanism \cite{12}, and the kinetic model \cite{13} which considers continuous regeneration in a QGP. Recently, the $J/\psi$ transverse momentum distribution at RHIC energy was discussed \cite{14} in the frame of a two-component model \cite{13} which includes both initial $J/\psi$ production through nucleon-nucleon (NN) interaction and regeneration.

The medium created in high-energy nuclear collisions evolves dynamically. In order to extract information about the medium by analyzing the $J/\psi$ distribution, both the hot and dense medium and the $J/\psi$ production processes must be treated dynamically. Since the massive $J/\psi$s are unlikely fully thermalized with the medium, their phase space distribution is governed by a transport equation. In the transport approach \cite{16} which includes dissociation and regeneration processes in a QGP, the $p_t$ broadening due to initial gluon rescattering and the leakage effect in the final state are respectively taken into account through the initial condition and the free streaming term of the transport equation. To comprehensively treat the $J/\psi$ distribution in the transport approach, the $J/\psi$ transport equation is solved together with hydrodynamic equations which characterize the space-time evolution of the QGP. In this Letter, we investigate in the transport model the $J/\psi$ transverse momentum distributions in heavy ion collisions at RHIC and LHC energies, including the ($\langle p_T^2 \rangle$) as a function of centrality and the $R_{AA}$ as a function of $p_t$.

In pp collisions, the $J/\psi$s from the feed-down of $\psi'$ and $X_c$ are respectively about 10% and 30% of the total final $J/\psi$s \cite{17}. Since $\Psi(= J/\psi, \psi', X_c)$ is heavy, the distribution function $f_\Psi(p_t, x_t, \tau | b)$ in central rapidity region and transverse phase space ($p_t, x_t$) at time $\tau$ and fixed impact parameter $b$ is controlled by a classical Boltzmann-type transport equation \cite{16}.

\[ \partial f_\Psi / \partial \tau + v_\Psi \cdot \nabla f_\Psi = -\alpha_\Psi f_\Psi + \beta_\Psi. \] (1)
The second term on the left-hand side arises from free-streaming of Ψ with transverse velocity \( \mathbf{v}_\Psi = \mathbf{p}_\tau / \sqrt{p_\tau^2 + m_\Psi^2} \), which leads to the leakage effect and is important for those high momentum charmonia. The suppression and regeneration in hot medium are reflected in the loss term \( \alpha_{\Psi}(p_t, x_t, \tau|b) \) and gain term \( \beta_{\Psi}(p_t, x_t, \tau|b) \). Considering only the gluon dissociation process \( g + \Psi \rightarrow c + \bar{c} \), \( \alpha_{\Psi} \) can be expressed as

\[
\alpha_{\Psi}(p_t, x_t, \tau|b) = \frac{1}{2E_{\Psi}} \int \frac{d^3p_g}{(2\pi)^3 2E_g} W_{g\Psi}^\infty(s) f_g(p_g, T, u) \Theta(T - T_c) / \Theta(T_d^g - T),
\]

where \( E_g \) and \( E_{\Psi} \) are the gluon and charmonium energies, \( f_g = 1/((E_{\Psi}^g u/\epsilon_{\Psi}^g - T) - 1) \) is the gluon thermal distribution, \( T(x_t, \tau|b) \) and \( u(x_t, \tau|b) \) are the local temperature and velocity of the hot medium, and \( W_{g\Psi}^\infty \) is the transition probability of the gluon dissociation as a function of \( s = (p + p_g)^2 \) calculated with perturbative Coulomb potential [18].

For \( J/\psi \), the dissociation cross section reads

\[
\sigma_{J/\psi}(\omega) = A_0 \frac{(\omega/\epsilon_{J/\psi} - 1)^3/2}{(\omega/\epsilon_{J/\psi})^5},
\]

with \( A_0 = (2^{11/27} \pi^5 \epsilon_{J/\psi})^{-1/2} \), where \( \omega \) is the gluon energy in the rest frame of \( J/\psi \), \( m_c \) is the charm quark mass, and \( \epsilon_{J/\psi} \) is the binding energy of \( J/\psi \). From what discussed in [19], to take the relativistic effect into consideration and to avoid non-physics divergence in the regeneration cross section, we replace the \( J/\psi \) binding energy by the gluon threshold energy.

The step function in the numerator of the loss term \( \alpha \), controlled by the critical temperature \( T_c \), deconfinement phase transition, means that the dissociation happens only in the QGP phase, namely \( \alpha = 0 \) for \( T < T_c \). Since the hadronic phase occurs later in the evolution of heavy ion collisions when the density of the system is lower compared to the early hot and dense period, we have neglected the hadronic dissociation. The step function in the denominator of \( \alpha \), characterized by the dissociation temperature \( T_d^\Psi \), indicates that any \( \Psi \) cannot survive when the temperature of the QGP is higher than \( T_d^\Psi \), namely \( \alpha = \infty \) for \( T > T_d^\Psi \). From the lattice QCD simulation [20], the \( J/\psi \) spectral function remains almost the same when the temperature is between \( T_c \) and \( T_d \), but the sharp peak of \( J/\psi \) vanishes suddenly when the temperature reaches \( T_d \). The introduction of the dissociation temperature \( T_c \) is also consistent with the idea of sequential charmonium suppression [19, 21].

The gain term \( \beta \) in the transport equation (1), which is a function of \( c \) and \( \bar{c} \) recombination transition probability \( W_{c\bar{c}}^\infty \), can be obtained from the lose term \( \alpha \) using detailed balance [13, 16]. Instead of the gluon distribution \( f_g \) in \( \alpha \), the charm quark distribution \( f_c \) in \( \beta \) is assumed to be in the form of

\[
f_c(p_c, x_t, \tau|b) = \frac{d\sigma_{NN}^{c\bar{c}}}{dy} \bigg|_{y=0} \frac{T_A(x_t) T_B(x_t - b)}{\tau} f_{cp}(p_c|b),
\]

where \( d\sigma_{NN}^{c\bar{c}}/dy|_{y=0} \) is the production cross section of charm quark pairs at central rapidity region in NN collisions, and \( T_A \) and \( T_B \) are the thickness functions of the two colliding nuclei defined as \( T(x_t) = \lim_{z_1 \to -\infty, z_2 \to -\infty} T(x_t, z_1, z_2) \) with \( T(x_t, z_1, z_2) = \int_{z_1}^{z_2} d\rho(x_t, z) \) and \( \rho(r) \) being the Woods-Saxon nuclear density profile. From the PHENIX data, open charmars carries an elliptic flow \( v_2 \sim 0.1 \) [22] which means that charm quarks are likely thermalized with the medium, we take the charm quark momentum distribution \( f_{cp} \) as the Fermi-Dirac function \( f_{cp}(p_c, T, u|b) \sim 1/((E_{c\bar{c}}^{p_d}/u_{c\bar{c}} + T) + 1) \).

With the known suppression and regeneration terms \( \alpha \) and \( \beta \), the transport equation can be solved analytically with the result [16]

\[
f_{\Psi}(p_t, x_t, \tau|b) = f_{\Psi}(p_t, x_t - \mathbf{v}_\Psi(\tau - \tau_0), \tau_0|b) e^{-\int_{\tau_0}^{\tau} d\tau' \alpha(p_t, x_t - \mathbf{v}_\Psi(\tau - \tau'), \tau'|b)} + \int_{\tau_0}^{\tau} d\tau' \beta(p_t, x_t - \mathbf{v}_\Psi(\tau - \tau'), \tau'|b) e^{-\int_{\tau_0}^{\tau} d\tau'' \alpha(p_t, x_t - \mathbf{v}_\Psi(\tau - \tau''), \tau'|b)}.
\]

The first and second terms on the right-hand side indicate the contributions from the initial production and continuous regeneration, respectively. Both suffer anomalous suppression. The coordinate shift \( x_t \to x_t - \mathbf{v}_\Psi \Delta \tau \) reflects the leakage effect during the time period \( \Delta \tau \).

Since the collision time for NN interactions at RHIC energy is about 0.1 fm/c and less than the starting time \( \tau_0 \) of the medium evolution which is about 0.5 fm/c, the nuclear absorption for the initially produced charmonia has ceased before the QGP evolution and is reflected in the transport process as the initial distribution \( f_{\Psi}(p_t, x_t, \tau_0|b) \) of the solution [16]. Considering a finite formation time of charmonia which is about 0.5 fm/c and larger than the collision time, the nuclear absorption can be safely neglected for the calculations in heavy ion collisions at high energy [23]. In this case, the initial distribution can be written as

\[
f_{\Psi}(p_t, x_t, \tau_0|b) = \frac{5}{4\pi} \frac{d\sigma_{NN}^{c\bar{c}}}{dy} \bigg|_{y=0} \int dz_A dz_B \rho_A(x_t, z_A) \rho_B(x_t - b, z_B) \frac{1}{(p_t^2)_N} \left( 1 + \frac{p_t^2}{4(p_t^2)_N} \right)^{-6}
\]

with a normalized power-law momentum distribution extracted from the PHENIX data for pp collisions [24], where \( d\sigma_{NN}^{c\bar{c}}/dy|_{y=0} \) is the charmonium production cross section at central rapidity region in NN collisions, and \( (p_t^2)_N \) is
the averaged transverse momentum square after NN collisions \[^6\]

\[
\langle p_T^2 \rangle_N(x_t, z_A, z_B|b) = \langle p_T^2 \rangle_{NN} + a_{NN} \rho_0^{-1} (T_A(x_t, -\infty, z_A) + T_B(x_t - b, z_B, \infty))
\]

with \(\rho_0\) being normal nuclear density. For given values \(b\) and \(x_t\) in the transverse plane, suppose a \(\Psi\) is produced at longitudinal coordinates \(z_A\) and \(z_B\) in nuclei A and B, respectively. The two gluons which fuse to form the \(\Psi\) carry transverse momentum from two sources: 1) Intrinsic \(p_t\), because they had been confined to nucleons. The intrinsic part is observable via \(NN \to \Psi\) process and leads to \(\langle p_T^2 \rangle_{NN}\) in Eq.\(^7\). 2) Collisional contribution to \(p_t\), because in a nuclear matter, the gluons traverse thickness \(T_A(x_t, -\infty, z_A)\) and \(T_B(x_t - b, z_B, \infty)\) of nuclear matter in A and B, respectively, and acquire additional transverse momentum via \(gN\) collisions. This is the origin of the second term in Eq.\(^7\). The constant \(a_{NN}\) is usually adjusted to the data for pA collisions.

The local temperature \(T(x_t, \tau|b)\) and fluid velocity \(u_\mu(x_t, \tau|b)\), which appear in the thermal gluon and charm quark distribution functions and control the suppression and regeneration region via the two step functions in \(\alpha\) and \(\beta\), are determined by the \((2+1)\) dimensional Bjorken’s hydrodynamic equations for the medium evolution,

\[
\begin{align*}
\partial_\tau E + \nabla \cdot \mathbf{M} &= -(E + p)/\tau, \\
\partial_\tau M_x + \nabla \cdot (M_x \mathbf{v}) &= -M_x/\tau - \partial_x p, \\
\partial_\tau M_y + \nabla \cdot (M_y \mathbf{v}) &= -M_y/\tau - \partial_y p, \\
\partial_\tau R + \nabla \cdot (R \mathbf{v}) &= -R/\tau
\end{align*}
\]

with the definitions \(E = (\epsilon + p)\gamma^2 - p, \mathbf{M} = (\epsilon + p)\gamma \mathbf{v}\) and \(R = \gamma\eta\), where \(\gamma\) is the Lorentz factor, and \(\epsilon, p\) and \(\mathbf{v}\) are the local energy density, pressure and transverse velocity of QGP.

To close the hydrodynamical equations we need to know the equation of state of the medium. We follow Ref.\(^{24}\), where the deconfined phase at high temperature is an ideal gas of massless u, d quarks, 150 MeV massed s quarks and gluons, and the hadron phase at low temperature is an ideal gas of all known hadrons and resonances with mass up to 2 GeV.\(^{25}\) There is a first order phase transition between these two phases. In the mixed phase, the Maxwell construction is used. The mean field repulsion parameter and the bag parameter are chosen as \(K = 450\) MeV fm\(^3\)\(^{25}\) and \(B^{1/3} = 236\) MeV to obtain the critical temperature \(T_c = 165\) MeV at vanishing baryon number density. The initial condition of the hydrodynamic equation at RHIC energy is the same as in Ref.\(^{16}\).

Solving the local temperature and fluid velocity from the hydrodynamic equations and then substituting them into the transport solution \(^{10}\) for \(J/\psi, \psi'\) and \(\chi_c\), we obtain the distribution function \(f_{J/\psi}(p_t, x_t, \tau_f|b)\) for the final state \(J/\psi s\) at the freeze-out time \(\tau_f\). By employing the Cooper-Frye formula \(^{27}\) with a longitudinal Hubble-like fluid \(^{28}\), any physical observable \(A\) for \(J/\psi\) can be estimated by integrating \(A \cdot f_{J/\psi}\) over the freeze-out surface.

We now calculate the nuclear modification factor \(R_{AA}\) and averaged transverse momentum square \(\langle p_T^2 \rangle\) as functions of centrality decided by the number of participant nucleons \(N_p\) for \(J/\psi s\) produced in Cu+Cu and Au+Au collisions at RHIC energy. The following parameters are used in our calculations: the charm quark mass \(m_c = 1.87\) GeV\(^{12}\), the charmonium mass \(m_{J/\psi} = 3.6\) GeV and \(m_{\psi'} = m_{\chi_c} = 3.7\) GeV, the starting time of hydrodynamic evolution \(\tau_0 = 0.6\) fm, the charm quark and charmonium production cross section at \(s = 200\) GeV as \(d^2 \sigma^c_{NN}/dy_{|y| = 0} = 120\) mb\(^{29}\) and \(d^2 \sigma^\psi_{NN}/dy_{|y| = 0} = 26.4, 4.4,\) and 13.2 nb for \(\Psi = J/\psi, \psi'\) and \(\chi_c\).\(^{24}\)

The \(\Psi\) binding energy or the gluon threshold energy in hot and dense medium should be smaller than its value in vacuum. It is estimated to be less than 220 MeV in the QGP phase\(^{30,31}\). We will take \(\epsilon_{J/\psi} = 150\) MeV. From our numerical results, a not very large deviation from this value does not lead to a sizeable change in \(R_{AA}\) and \(\langle p_T^2 \rangle\). The quasi-free cross section has a larger width than that of the gluon dissociation with binding energy.

The results of \(R_{AA}(N_p)\) and \(\langle p_T^2 \rangle(N_p)\) at central rapidity for Au+Au collisions are shown in Fig.\(^{14}\) and compared with the RHIC data.\(^{14}\) The \(R_{AA}\) is defined as the ratio of \(J/\psi s\) produced in a nuclear collision to that in a corresponding pp collision, normalized by the number of binary collisions. When the centrality increases, the contribution from the initial production drops down monotonously, due to the suppression in the QGP, while the number of regenerated \(J/\psi s\) goes up monotonously, because of the number of charm quarks increases with the collision centrality. From the experimental data, there exists a flat region at \(50 < N_p < 170\) with \(R_{AA} \simeq 0.6\), and at \(N_p > 170\) the \(R_{AA}\) continues to decrease with increasing centrality. This suppression structure can be reproduced in our transport approach by choosing the dissociation temperature \(T_d^{J/\psi} = T_d^{\chi_c} = T_c\) and \(T_d^{J/\psi} = 1.92T_c\). After the plateau ends, the space-time region with temperature \(T > T_d^{J/\psi}\) increases with centrality, and more \(J/\psi s\) are fully eaten up by the extremely hot QGP. This is the reason of the further decrease of \(R_{AA}\) at \(N_p > 170\). Since the QGP with \(T > T_d\) exists only in the early stage of the fireball and lasts for a short time, the initially produced charmonia experience this stage, while most of the regenerated charmonia that are created later do not. This is why the \(R_{AA}\) for regenerated \(J/\psi s\) is much less influenced by the dissociation temperature. The participant number fluctuations becomes important to \(R_{AA}\) at extremely large \(N_p\), as discussed at SPS energy\(^{32}\), and may explain the deviation of our result from the data in very central collisions.

We also calculated the \(R_{AA}\) for Cu+Cu collisions and compared it with the RHIC data.\(^{33}\) When the colliding energy is fixed, the temperature of the fireball is mainly controlled by the participant number, and therefore the result for Cu+Cu is almost the same as that for semi-central Au+Au collisions with \(N_p < 110\), where 110 is the
maximum participant number for a Cu+Cu collision. Since Cu is much lighter than Au, the fireball formed in Cu+Cu collisions cannot reach the temperature for full $J/\psi$ dissociation.

Most of the produced particles are with low momentum, they dominate the centrality dependence of the yield, and the information carried by high $p_t$ charmonia is therefore screened in $R_{AA}(N_p)$. In order to understand the behavior of those high $p_t$ $J/\psi$s which are more sensitive to the production and suppression mechanisms, we consider the averaged transverse momentum square $\langle p_t^2 \rangle$ in the bottom panel of Fig.1. In our calculation, the values of $\langle p_t^2 \rangle_{NN} = 4.14$ (GeV/c)$^2$ [24] and $a_{gN} = 0.1$ (GeV/c)$^2$/fm [14] are used for NN collisions, see eq. (7). At SPS, the gluon rescattering parameter was taken as $a_{gN} = 0.077$ (GeV/c)$^2$/fm [5, 34]. As one can see, the initial contribution to $\langle p_t^2 \rangle$ increases smoothly at $N_p \leq 170$ and shows a saturation when $N_p > 170$. This behavior is very similar to the case at SPS energy [2] where there is almost no regeneration and the initial production can be considered as the total result. The regenerated $J/\psi$s are from the thermalized charm quarks, therefore, their transverse momentum is rather small and the averaged value is almost centrality independent, in comparison with the initially produced charmonia, see the dotted line. Since the regeneration becomes more important in central collisions, the total $\langle p_t^2 \rangle$ is strongly suppressed at large $N_p$ as a consequence of the competition between the initial production and regeneration. While almost all the models with and without regeneration mechanism can describe the $J/\psi$ yield after at least one parameter is adjusted, the $p_t$ suppression in central collisions seems to be a signature of charmonium regeneration in Au+Au collisions at RHIC energy. Again, the calculation for Cu+Cu collisions is very similar to the result for Au+Au in the region of $N_p < 110$.

To look into the details of $J/\psi$ production and suppression mechanisms, we now turn to the calculation of $R_{AA}(p_t)$ as a function of $p_t$ which is discussed in different models [14, 35]. The result for central Cu+Cu and Au+Au collisions at midrapidity is shown and compared with the RHIC data in Fig.2. The monotonous increase of $R_{AA}$ with only initial production comes from three aspects. One is the $p_t$ dependence of the gluon dissociation cross section. The gluon energy relative to $J/\psi$ in the cross section is scaled by $J/\psi$ binding energy which is much smaller than 1 GeV. Therefore, gluons with small energy are more likely to dissociate a $J/\psi$, or in other

---

FIG. 1: The nuclear modification factor $R_{AA}$ and averaged transverse momentum square $\langle p_t^2 \rangle$ as functions of participant nucleon number $N_p$ in central rapidity region for Au+Au collisions at RHIC energy. Dot-dashed, dotted and solid lines represent respectively the calculations with only initial production, only regeneration and both contributions. The data are from PHENIX collaboration [2, 3].
words, \(J/\psi\)s with low momentum are easy to be eaten up by the hot medium. The second reason is the leakage effect with which the high momentum charmonia can escape the anomalous suppression region. The last one is the \(p_t\) broadening happened in the initial state. Gluons get additional transverse momentum via collisions with participant nucleons before they fuse to \(J/\psi\)s. At SPS energy the \(J/\psi R_{AA}\) has been observed to exceed unity at high \(p_t\) which can be well explained by the gluon multi-scattering [7]. At RHIC energy, from our numerical calculation, the gluon rescattering is essential for the increasing of \(R_{AA}\) in high \(p_t\) region.

In comparison with the initially produced \(J/\psi\)s which can carry high momentum from the hard process, the regenerated \(J/\psi\)s from thermalized charm quarks are distributed at low momentum region. In Fig.2 these effects are clearly shown by dot-dashed and dotted lines for initial production and regeneration, respectively. In the low \(p_t\) region, the competition between the initial production which increases with \(p_t\) and the regeneration which decreases with \(p_t\) leads to a relatively flat structure for central Cu+Cu collisions.

For the \(p_t\) dependence of \(R_{AA}\) for central Au+Au collisions, shown in the bottom panel of Fig.2 while the trend of the curves are similar to that for Cu+Cu collisions, there are obvious features arisen from the stronger suppression and stronger regeneration in Au+Au collisions. In comparison with Cu+Cu collisions, the fireball formed in Au+Au collisions is much hotter and larger and lasts much longer, and the initially produced \(J/\psi\)s are strongly suppressed on their way out of the fireball. Since not all the regenerated \(J/\psi\)s are from the extremely hot medium, the effect of the stronger suppression in Au+Au collisions on the regeneration is not so important as for the initial production. Considering the fact that the regeneration is proportional to the square of the number of binary collisions, it is more important in central Au+Au collisions than in Cu+Cu collisions. At low \(p_t\) the regeneration contribution even exceeds the initial production. It is the competition between the stronger suppression of the initially produced \(J/\psi\)s and the stronger regeneration at low \(p_t\) that explains the overall reduced \(J/\psi R_{AA}\) and its minimum at \(p_t \sim 2.5\) GeV/c in Au+Au collisions. For central Cu+Cu collisions, there seems to have a minimum too, see the top panel of Fig.2 However, due to the relatively small contribution from the regeneration, our calculations do not show such a clear minimum.

As a prediction, we calculated the \(J/\psi R_{AA}\) and \(\langle p_t^2 \rangle\) at LHC energy, the results are shown in Fig.3. The input to the transport equation, namely the charmonia and charm quark production cross sections in NN collisions, is taken as \(d\sigma^{J/\psi}_{NN}/dy = 2\) \(\mu\)b and \(d\sigma^{c\bar{c}}_{NN}/dy = 0.7\) mb, estimated by the CEM model and PYTHIA simulation [36]. In comparison with the heavy ion collisions at RHIC, the formed fireball at LHC is hotter, bigger and longer lived, and almost all the initially produced charmonia are eaten up by the medium. On the other hand, there are much

![Graph showing nuclear modification factor \(R_{AA}\) as a function of transverse momentum \(p_t\) in central rapidity region for Cu+Cu and Au+Au collisions at RHIC energy. Dot-dashed, dotted and solid lines represent respectively the calculations with only initial production, only regeneration and both contributions. The data are from STAR [4] for Cu+Cu and PHENIX [3] for Au+Au.](image)
more charm quarks generated at LHC energy, the regeneration controls the population in central and semi-central collisions. This regeneration dominance leads to an increasing $R_{AA}$ and a much stronger $p_t$ suppression at LHC for large enough $N_p$. The rapid change of $R_{AA}$ and $\langle p_t^2 \rangle$ in the small $N_p$ region is due to the strong competition between the initial production and regeneration, and the degree of this competition decreases with increasing centrality and ends when $N_p$ is large enough. The increasing $R_{AA}$ and especially the small and saturated $\langle p_t^2 \rangle$ for central collisions can be regarded as signatures of regeneration dominance at LHC, which are not expected in a model with initial production only.

In summary, we investigated the transverse momentum dependence of $J/\psi$ production in heavy ion collisions at RHIC and LHC energies. By considering both initial production and regeneration and solving the coupled hydrodynamic equations for medium evolution and the transport equation for $J/\psi$ motion, we focused on the centrality dependence of $\langle p_t^2 \rangle(N_p)$ and $p_t$ dependence of $R_{AA}(p_t)$ in Au+Au and Cu+Cu collisions. While the high $p_t$ behavior is characterized by the initial production, the regeneration and initial production become equally important at low $p_t$. At RHIC, the competition between the two production mechanisms leads to the decrease of $J/\psi$ $\langle p_t^2 \rangle(N_p)$ in central collisions and the minimum of $R_{AA}(p_t)$ at low $p_t$. It is necessary to emphasize that the initial production alone cannot reproduce the $p_t$ suppression and the minimum of $R_{AA}$ found at RHIC. At LHC, almost all the initially produced charmonia are eaten up by the hotter, larger and longer lived fireball, the $J/\psi$ behavior is dominated by the regeneration when $N_p$ is large enough. As a result, the $J/\psi$ $R_{AA}$ increases with centrality and decrease with $p_t$ in central collisions. On the other hand, there is almost no regeneration at FAIR energy, and the $J/\psi$ production is governed by the initial production. In this case, both the $p_t$ suppression and the minimum of $R_{AA}$ will disappear in heavy ion collisions at FAIR energy.

Acknowledgments: We are grateful to Xianglei Zhu and Li Yan for their help in numerical calculations. The work is supported by the NSFC grant No. 10735040, the 973-project No. 2007CB815000, and the U.S. Department
of Energy under Contract No. DE-AC03-76SF00098.

[1] T.Matsui and H.Satz, Phys. Lett. **B178**, 416(1986).
[2] J.Lajoie [PHENIX Collaboration], J. Phys. **G34**, S191(2007).
[3] A.Adare *et al.*, [PHENIX Collaboration], Phys. Rev. Lett. 98 (2007) 232301.
[4] B.I.Abelev *et al.*, [STAR Collaboration], arXiv:0904.0439
[5] A.Adare *et al.*, [PHENIX Collaboration], Phys. Rev. Lett. **B215**, 218(1988); J.P.Blaizot and J.Y.Ollitrault, Phys. Lett. **B217**, 392(1989).
[6] J.Huefner and P.Zhuang, Phys. Lett. **B559**, 193(2003); X.Zhu and P.Zhuang, Phys. Rev. **C67**, 067901(2003).
[7] J.P.Blaizot and J.Y. Ollitrault, Phys. Lett. **B 199**, 499(1987).
[8] B.I.Abelev *et al.*, [STAR Collaboration], arXiv:0904.0439.
[9] M.C.Abreu *et al.*, [NA50 Collaboration], Phys. Lett. **B499**, 85(2001).
[10] S.Gavin and M.Gyulassy, Phys. Lett. **B214**, 241(1988); J.Hüfner, Y.Kurihara and H.J.Pirner, Phys. Lett. **B215**, 218(1988); J.P.Blaizot and J.Y.Ollitrault, Phys. Lett. **B217**, 392(1989).
[11] J.Huefner and P.Zhuang, Phys. Lett. **B559**, 193(2003); X.Zhu and P.Zhuang, Phys. Rev. **C67**, 067901(2003).
[12] J.P.Blaizot and J.Y. Ollitrault, Phys. Lett. **B 199**, 499(1987).
[13] F. Karsch and R. Petronzio, Z. Phys. **C37**, 627(1988).
[14] R.V.Gavai *et al.*, Int. J. Mod. Phys. **A10**, 2999(1995).
[15] P.Braun-Munzinger and J. Stachel, Phys. Lett. **B490**, 196(2000) and Nucl. Phys. **A690**, 119(2001); M.I.Gorenstein, A.P.Kostyuk, H.Stöcker and W. Greiner, Phys. Lett. **B509**, 277(2001); L.Grandchamp and R.Rapp, Phys. Lett. **B523**, 60(2001) and Nucl. Phys. **A709**, 415(2002).
[16] K.Hagiwara *et al.*, Particle Data Group, Phys. Rev. **D66**, 010001(2002).
[17] C.-Y.Wong, *Introduction to High-Energy Heavy-Ion Collisions*, World Scientific, 1994, p.27-34.
[18] Y.Zhang, J. Phys. **G35**, 104022(2008); A.Adare *et al.*, Phys. Rev. Lett. **97**, 252002(2006); D.Hornback, [PHENIX Collaboration], J. Phys. **G35**, 104113(2008).
[19] F.Karsch, M.T.Mehr and H.Satz, Z. Phys. **A783**, 249(2007).
[20] L.Grandchamp, R.Rapp, Phys. Lett. **B523**, 60(2001).
[21] C.-Y.Wong, *Introduction to High-Energy Heavy-Ion Collisions*, World Scientific, 1994, p.27-34.
[22] Y.Zhang, J. Phys. **G35**, 104022(2008); A.Adare *et al.*, Phys. Rev. Lett. **97**, 252002(2006); D.Hornback, [PHENIX Collaboration], J. Phys. **G35**, 104113(2008).
[23] P.Karsch, M.T.Mehr and H.Satz, Z. Phys. **C37**, 617(1998).
[24] L.Grandchamp, R.Rapp, Phys. Lett. **B523**, 60(2001).
[25] J.-P.Blaizot, P.M.Dinh, J.-Y.Ollitrault, Phys. Rev. Lett. **85**, 4012(2000).
[26] S.X.Oda *et al.*, [PHENIX Collaboration], J. Phys. **G35**, 104134(2008).
[27] T.Gunji, H.Hamagaki, T.Hatsuda, T.Hirano and Y.Akamatsu, J. Phys. **G35**, 104137(2008).