Event-by-event bottomonia suppression in relativistic heavy-ion collisions

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Abstract. Using the screened Cornell potential and the next-to-leading order perturbative QCD to determine, respectively, the properties of bottomonia and their dissociation cross sections in a quark-gluon plasma, we have studied in a 2+1 ideal hydrodynamics the effect of initial fluctuations on bottomonia production in relativistic heavy-ion collisions. We have found that while initial fluctuations hardly affect the yield of the 1S ground state bottomonium, their effect on that of excited bottomonium states is not small. Compared to the case with smooth initial conditions, the survival probabilities of excited bottomonium states are reduced at low transverse momentum in the case of large initial fluctuations while they are increased in the case of small initial fluctuations as a result of the smearing effect introduced in solving the hydrodynamic equations. The observed suppression of excited bottomonia relative to the ground state bottomonium of an average transverse momentum of 6.25 GeV/c by the Compact Muon Solenoid (CMS) collaborations can, however, be described at present with both smooth and fluctuating initial conditions.

1. introduction

Since the suggestion of $J/\psi$ suppression as a possible signature of the quark-gluon plasma (QGP) formed in relativistic heavy ion collisions \cite{1}, extensive studies have been carried out both theoretically and experimentally not only on charmonia production but also on bottomonia production in these collisions \cite{2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17}. Compared to $J/\psi$ production, the ground state bottomonium is, however, a more promising probe of the hot dense matter created in relativistic heavy ion collisions because of the relatively small contributions from regeneration in the QGP \cite{19, 18}. Recently, suppressions of bottomonia production in relativistic heavy-ion collisions relative to those expected from p+p collisions at same energies have been observed at both the Relativistic Heavy-Ion Collider (RHIC) \cite{20} and the Large Hadron Collider (LHC) \cite{21, 22}. To understand these experimental results, we have studied in Ref.\cite{19} bottomonia production in these collisions by including their production from both initial hard scattering of colliding nucleons and the regeneration from bottom and antibottom quarks in the produced quark-gluon plasma based on a schematic hydrodynamics for the bulk collision dynamics \cite{19}. We found that the contribution from regeneration was very small for bottomonia as expected and the modification of the thermal properties of bottomonia in hot dense matter was helpful in describing the experimental data. The above study has, however, neglected the temperature fluctuation in the produced hot dense matter since we have taken all...
thermal quantities to be uniform in the schematic hydrodynamics model. In a recent study [23], we have improved our previous results by using a 2+1 ideal hydrodynamics to include also the effect of initial fluctuations on bottomonia production in relativistic heavy ion collisions. In the following, we present results from this study.

2. The 2+1 ideal hydrodynamics for relativistic heavy ion collisions

The hydrodynamic equations are based on the conservations of energy-momentum and various charges. Because chemical potentials are negligible in heavy-ion collisions at LHC [24], only the energy-momentum conservations are used in our study. Assuming the boost invariance, energy-momentum conservations are then expressed in the \((\tau, x, y, \eta)\) coordinate system, with \(\tau = \sqrt{t^2 - z^2}\) and \(\eta = \frac{1}{2} \ln \left(\frac{z + \sqrt{z^2 + t^2}}{z - \sqrt{z^2 + t^2}}\right)\), as [25, 26]

\[
\begin{align*}
\partial_\tau (\tau T^{00}) + \partial_x (\tau T^{0x}) + \partial_y (\tau T^{0y}) & = -p, \\
\partial_\tau (\tau T^{0x}) + \partial_x (\tau T^{xx}) + \partial_y (\tau T^{xy}) & = 0, \\
\partial_\tau (\tau T^{0y}) + \partial_x (\tau T^{xy}) + \partial_y (\tau T^{yy}) & = 0,
\end{align*}
\]

where \(T^{\mu\nu}\) and \(p\) are, respectively, the energy-momentum tensor and pressure.

To solve Eq. (1) requires information on the initial conditions of a collision, particularly the initial entropy density and the equation of state of produced matter. For the initial entropy density, it is given as

\[
\frac{ds}{d\eta} = C \left\{ (1 - \alpha) \frac{n_{\text{part}}}{2} + \alpha n_{\text{coll}} \right\},
\]

where \(n_{\text{part}}\) and \(n_{\text{coll}}\) are, respectively, the number densities of participants and binary collisions. In the case of smooth initial conditions as used in Ref. [9], they are given by

\[
\begin{align*}
n_{\text{part}}(\vec{r}) & \equiv \frac{d^2 N_{\text{part}}}{\tau_0 d\tau dy} = AT_A(\vec{r}) \left[ 1 - \{ 1 - T_B(\vec{b} - \vec{r})\sigma_{in} \}^B \right] \\
& \quad + BT_B(\vec{b} - \vec{r}) \left[ 1 - \{ 1 - T_A(\vec{r})\sigma_{in} \}^A \right], \\
n_{\text{coll}}(\vec{r}) & \equiv \frac{d^2 N_{\text{coll}}}{\tau_0 d\tau dy} = \sigma_{in} ABT_A(\vec{r}) T_B(\vec{b} - \vec{r}),
\end{align*}
\]

where \(\tau_0\) is the initial thermalization time; \(A\) and \(B\) are mass numbers of colliding nuclei; \(\vec{r}\) and \(\vec{b}\) are, respectively, the transverse position and impact parameter vectors; \(\sigma_{in}\) is the nucleon-nucleon inelastic cross section and has a value of 64 mb at LHC energies [28]; \(T_{A(B)} = \int d\vec{r} A_{A(B)}(\vec{r}, z)\) is the thickness function with \(A_{A(B)}\) being the nucleon distribution function in nucleus \(A(B)\) for which the Wood-Saxon model is used.

In the case of fluctuating initial conditions, the positions of colliding nucleons are randomly distributed according to \(A_{A(B)}\), and two nucleons are considered as participants and a binary collision takes place at their middle point if the transverse distance between a nucleon from nucleus \(A\) and a nucleon from nucleus \(B\) is shorter than \(\sqrt{\frac{\sigma_{in}}{\pi}}\). In generating the initial conditions for the hydrodynamic evolution, a smearing parameter \(\sigma\) is introduced in evaluating the number densities of participants and binary collisions, i.e.,

\[
n_{\text{part(coll)}}(\vec{r}) = \frac{1}{2\pi\sigma^2\tau_0} \sum_{i=1}^{N_{\text{part(coll)}}} \exp \left( - \frac{|\vec{r}_i - \vec{r}|^2}{2\sigma^2} \right),
\]

where \(\vec{r}_i\) is the transverse position of \(i\)-th participant or binary collision. We have used the same smearing parameter \(\sigma\) for both number densities [30] and consider the two values of \(\sigma = 0.4\) and \(0.8\) fm, corresponding to large and small initial fluctuations, respectively.
For the equations of state (EoS), we have used the quasiparticle model based on early lattice QCD data for the QGP and the resonance gas model for the hadron gas as in Refs. [32, 9]. This model thus assumes the presence of a first-order phase transition and the critical temperature $T_c$ is 170 MeV. We note that compared to recent lattice data from the hotQCD [33] and the Wuppertal-Budapest collaboration [34], our EoS is similar to that from the hotQCD Collaboration and the difference is smaller than that between the two latest lattice QCD calculations [35].

Figure 1. (Color online) Transverse momentum spectra of $\pi^-$ and antiproton in 0-5 % central Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [29] compared with those from the ideal hydrodynamic model with the initial thermalization time, chemical and kinetic freeze-out temperatures of 0.6 fm/c, 160 MeV and 130 MeV, respectively.

The parameters $C$ and $\alpha$ in Eq. (2) as well as the initial thermalization time $\tau_0$ are determined from fitting the centrality dependence of the charged-particle multiplicity [31] and the transverse momentum spectra of $\pi^-$ and antiproton measured by the ALICE Collaboration in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [29], and they are shown in Fig. 1. Using the Cooper-Frye freeze-out formula and assuming that the multiplicity does not change after chemical freeze-out at temperature $T = 160$ MeV [24], we have obtained $\tau = 0.6$ fm/c, $C = 29.0$ and $\alpha = 0.15$. We note that this value of $C$ is slightly larger than that in our previous study based on a schematic viscous hydrodynamics [27] in order to reproduce the additional entropy generated in the viscous hydrodynamics.

In Fig. 2, we compare the temperature distributions in the transverse plane at the initial thermalization time $\tau_0 = 0.6$ fm/c (upper panels) and at a later time of $\tau = 1.0$ fm/c (lower panels) in Pb+Pb collisions at center of mass energy $\sqrt{s_{NN}} = 2.76$ TeV and impact parameter $b = 2.45$ fm for the smooth (left panels) and fluctuating initial conditions with $\sigma = 0.4$ (middle panels) and 0.8 fm (right panels). The latter are from one of the one hundred different events that are generated for this centrality.

3. Thermal properties of bottomonia in QGP
The potential energy between a pair of heavy quark and antiquark is modified in QGP due to the effect of color Debye screening. The free energy of this heavy quark system has been extracted from lattice QCD calculations, from which the internal energy can then be determined from the thermodynamics relation. Whether the free energy or the internal energy is more appropriate
for describing the potential energy between a heavy quark and antiquark pair in QGP is still controversial [36]. In our study, we have used instead the screened Cornell potential [37],

\[ V(r, T) = \frac{\bar{\sigma}}{\mu(T)} \left[ 1 - e^{-\mu(T)r} \right] - \frac{\alpha}{r} e^{-\mu(T)r}, \tag{5} \]

with \( \bar{\sigma} = 0.192 \text{ GeV}^2 \) and \( \alpha = 0.471 \). The screening mass \( \mu(T) \) depends on temperature, and we have used the one given in pQCD, i.e., \( \mu(T) = (N_c/3 + N_f/6)^{1/2} g T \), where \( N_c \) and \( N_f \) are numbers of colors and light quark flavors, respectively, and the coupling constant \( g \) is taken to be 1.87 [27]. Compared to the results from the lattice QCD, this potential is close to the free energy around critical temperature and becomes more similar to the internal energy with increasing temperature.

Solving the Schrödinger equation with the potential in Eq. (5) for the bottom quark mass \( m_b = 4.746 \text{ GeV} \), we have obtained the dissociation temperatures 681, 285, 190, 257, and 185 MeV for the 1S, 2S, 3S, 1P and 2P bottomonium states, respectively [19]. A quarkonium then cannot be produced in regions where the temperature is higher than its dissociation temperature.

Even though a quarkonium can be formed in less hot region, it can still be dissociated by scattering with quarks and gluons in the QGP. This effect can be quantified by the thermal decay width of a bottomonium in QGP. For a bootomonium at rest in QGP, this is given by

\[ \Gamma(T) = \sum_i \int \frac{d^3k}{(2\pi)^3} v_{\text{rel}}(k)n_i(k, T)\sigma_{i}^{\text{diss}}(k, T), \tag{6} \]

where \( i \) denotes the quarks and gluons in the QGP; \( n_i \) is the momentum distribution function of parton species \( i \) in grand canonical ensemble; and \( v_{\text{rel}} \) is the relative velocity between the
scattering bottomonium and parton. For the dissociation cross sections of bottomonia $\sigma_{i}^{\text{diss}}$, we have calculated them up to the next-to-leading order (NLO) in pQCD [38, 39]. While in the leading order (LO) a bottomonium is dissociated by absorbing a thermal gluon, in the NLO it is dissociated by the gluon emitted from a quark or gluon in the QGP. Because of the dipole nature of quarkonia, their dissociation cross sections depend on the derivative of their wavefunctions with respect to the relative momentum between the constituent heavy quark and antiquark. Since the derivative of the wavefunctions obtained from the screened Cornell potential rapidly increases near the dissociation temperature of quarkonium, the thermal decay widths of bottomonia increase with increasing temperature and diverge at their dissociation temperatures as shown in Fig. 3 (a). For a bottomonium of finite velocity in an expanding QGP, its thermal decay width is affected by the modified QGP parton distribution in its rest frame. We have generalized Eq. (6) to include this effect [40], and the results are shown in Fig. 3 (b). It is seen that although the thermal decay width of bottomonium changes significantly with temperature, it does not vary much with the velocity [38].

We note that although a heavy quark pair is instantly produced in high-energy nuclear collisions, heavy quarkonium takes time to be formed from heavy quark pair [41, 42, 43]. In our study, we have ignored the thermal dissociation of bottomonia before their formation. Also, the excited states are expected to take longer time to form than the ground state. For simplicity, we have used the same formation time for both the ground state bottomonium and its excited states.

Bottomonia are produced with probabilities proportional to $n_{\text{coll}}$ in Eq. (3) for the case of smooth initial conditions and Eq. (4) for the case of fluctuating initial conditions. With its motion isotropically distributed in the azimuthal angle $\phi$, a bottomonium produced at $\vec{r}_0$ and moving with velocity $\vec{v}$ then has the survival probability

$$S(\vec{r}_0, \vec{v}) = \exp \left[ -\frac{1}{\gamma} \int_{t_0}^{\infty} d\tau \Gamma(\vec{r}_0 + \vec{v}\tau, \tau) \right]$$

with $\gamma = 1/\sqrt{1 - v^2}$.

**Figure 3.** (Color online) (a) Thermal decay widths of 1S, 2S, 1P, 3S, and 2P state bottomonia shown, respectively, by solid, dashed, dotted, dash-dotted, and dash-dot-dotted lines. (b) Thermal decay width of 1S state bottomonium as a function of velocity in QGP at temperature $T = 2.0$, 2.6, and 3.2 $T_c$. 

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4. Results

![Figure 4](image_url)

**Figure 4.** (Color online) Nuclear modification factor $R_{AA}$ of $\Upsilon(1S)$ as a function of participant number for smooth initial conditions (solid line) and fluctuating initial conditions with smearing parameter $\sigma = 0.4$ fm (dashed line) and 0.8 fm (dotted line) in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. Experimental data are taken from Ref. [22].

Neglecting their regeneration from bottom and antibottom quarks in QGP, we have calculated the nuclear modification factors of bottomonia in Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. Fig. 4 shows the nuclear modification factor $R_{AA}$ of $\Upsilon(1S)$ as a function of participant number for smooth initial conditions (solid line) and fluctuating initial conditions with smearing parameter $\sigma = 0.4$ fm (dashed line) and 0.8 fm (dotted line) together with the experimental results from the CMS Collaboration [22]. These results are obtained with the contributions from $\chi_b(1P)$, $\chi_b(2P)$, $\Upsilon(2S)$ and $\Upsilon(3S)$ to $\Upsilon(1S)$ decays taken to be 27.1, 10.5, 10.7 and 0.8 %, respectively, [44] and the velocity of bottomonia determined from the average transverse momentum of $\Upsilon(1S)$ since only the transverse momentum of 1S state has been measured [22]. As in our previous study, most suppression of $\Upsilon(1S)$ comes from the dissociation of its excited states. We note that the $R_{AA}$ from both the smooth and fluctuating initial conditions are similar to our previous result based on a schematic hydrodynamics [19], and they are also similar to each other.

Fig. 5 shows the nuclear modification factor $R_{AA}$ of 1S (upper lines), 2S (middle lines), and 3S (lower lines) bottomonium states as functions of transverse momentum in minimum-bias Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. Solid lines are from smooth initial conditions, and dashed and dotted lines are from fluctuating initial conditions with smearing parameter $\sigma = 0.4$ and 0.8 fm, respectively. The $R_{AA}$ of 1P and 2P states are similar to those of 2S and 3S states, respectively. It is seen that the $R_{AA}$ of 1S state is not much changed by initial fluctuations as a result of its strong binding and high dissociation temperature. The initial fluctuating effect on 2S and 3S states is, however, not small. Their $R_{AA}$ in the case of large fluctuating initial conditions ($\sigma = 0.4$ fm) are smaller than those in the case of smooth initial conditions. As shown in Fig. 2, nucleon-nucleon collisions are more locally concentrated in the case of fluctuating initial conditions, resulting in the formation of hot spots at which there is a relatively larger number of binary collisions. Although more bottomonia are produced at these hot spots, their survival probability from thermal dissociation decreases unless they have enough transverse momentum to escape these regions and enhance the so-called leakage effect. As a result, the $R_{AA}$ increases more rapidly with transverse momentum in the case of large fluctuating initial conditions. On
Figure 5. (Color online) Nuclear modification factor $R_{AA}$ of 1S (upper lines), 2S (middle lines), and 3S (lower lines) bottomonium states as functions of transverse momentum for smooth initial conditions (solid lines) and fluctuating initial conditions with smearing parameter $\sigma = 0.4$ fm (dashed lines) and 0.8 fm (dotted lines) in minimum-bias Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV.

On the other hand, the $R_{AA}$ in the case of small fluctuating initial conditions ($\sigma = 0.8$ fm) are close to or larger than that in the case of smooth initial conditions in all $p_T$.

Figure 6. (Color online) (a) Nuclear modification factor $R_{AA}$ of 1S bottomonium state as a function of transverse momentum in minimum-bias Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [22]. (b) Ratio of the yield of 2S and 3S bottomonium states to that of 1S state in minimum-bias Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV divided by that in p+p collisions at same energy. Solid, dashed, and dotted lines are from smooth and fluctuating initial conditions with smearing parameter $\sigma = 0.4$ and 0.8 fm, respectively. Experimental data are taken from Ref. [21] based on the average transverse momentum of 1S bottomonium state [22].
In Fig. 6 (a), we show the nuclear modification factor $R_{AA}$ of 1S bottomonium state as a function of transverse momentum in minimum-bias Pb+Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [22]. It is seen that initial fluctuations do not change much the result as in Fig. 4. Fig. 6 (b) shows the ratio of the yield of 2S and 3S bottomonium states to that of 1S state in Pb+Pb collisions divided by that in p+p collisions at same energy. This double ratio has the advantage that the cold nuclear matter effect is canceled if they are the same for the ground and excited bottomonium states. The experimental data shown in Fig. 6 has the value $0.31 \pm 0.03$ [21] at the average transverse momentum of 1S state [22]. The relative yield of 2S and 3S states in p+p collisions are obtained from Ref. [44]. Our results show that compared to that for smooth initial conditions, this ratio is suppressed for large initial fluctuations but is enhanced for small initial fluctuations. This is due to the fact that the dissociation probabilities of excited bottomonia are smaller for $\sigma = 0.8$ fm and larger for $\sigma = 0.4$ fm than for smooth initial conditions. Because of their large errors and limited momentum range, present experimental data can, however, be described with both smooth and fluctuating initial conditions.

5. Summary

Bottomonium and its excited states are promising probes of the properties of the hot dense matter created in relativistic heavy-ion collisions. In our study, we have investigated the effect of initial fluctuations in heavy-ion collisions on bottomonia suppression. For a more realistic description of the expansion dynamics of produced hot matter, a 2+1 ideal hydrodynamic model has been used. The thermal properties of bottomonia have been obtained from the screened Cornell potential and the dissipations of bottomonia by thermal partons were calculated up to NLO in pQCD. We have neglected, however, the small cold nuclear matter effect and the regeneration effect. We have found that the initial fluctuations hardly affect the survival probability of 1S state while the effect on excited states is not small. With large initial fluctuations corresponding to the smearing parameter $\sigma = 0.4$ fm, the survival probability of excited bottomonium state are suppressed at small transverse momentum, resulting in a survival probability that increases more rapidly with transverse momentum than in the case of smooth initial conditions. On the other hand, small initial fluctuations ($\sigma = 0.8$ fm) enhance the survival probabilities of excited bottomonium states at all transverse momenta as a result of the smearing effect. The available experimental data on the double ratio of excited states to ground state of bottomonium in Pb+Pb collisions to p+p collisions at an averaged transverse momentum can, however, be described at present with both smooth and fluctuating initial conditions.

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