Physics Letters B 772 (2017) 819–824

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Bottomonium continuous production from unequilibrium bottom quarks in ultrarelativistic heavy ion collisions

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ARTICLE INFO

Article history:
Received 25 April 2017
Received in revised form 11 July 2017
Accepted 25 July 2017
Available online 27 July 2017
Editor: W. Haxton

ABSTRACT

We employ the Langevin equation and Wigner function to describe the bottom quark dynamical evolutions and their formation into a bound state in the expanding Quark Gluon Plasma (QGP). The additional suppressions from parton inelastic scatterings are supplemented in the regenerated bottomonium. Hot medium modifications on $\Upsilon(1S)$ properties are studied consistently by taking the bottomonium potential to be the color-screened potential from Lattice results, which affects both $\Upsilon(1S)$ regeneration and dissociation rates. Finally, we calculated the $\Upsilon(1S)$ nuclear modification factor $R_{AA}^{\rm rege}$ from bottom quark combination with different diffusion coefficients in Langevin equation, representing different thermalization of bottom quarks. In the central Pb–Pb collisions ($b=0$) at $\sqrt{s_{NN}}=5.02$ TeV, we find a non-negligible $\Upsilon(1S)$ regeneration, and it is small in the minimum bias centrality. The connections between bottomonium regeneration and bottom quark energy loss in the heavy ion collisions are also discussed.

1. Introduction

Since charmonium was proposed as a probe for the existence of the deconfined matter called "Quark–Gluon Plasma" (QGP) [1], its production mechanisms has been widely studied based on coalescence model [2–4] and transport models [5–7] in nucleus–nucleus collisions. The nuclear modification factor R_{AA} is a measurement of cold and hot medium suppressions on quarkonium yields. The cold nuclear matter effects include the nuclear absorption [8], Cronin and shadowing effects [9–11]. The first one is negligible at LHC colliding energies due to strong Lorentz dilation, where "spectator" nucleons already move out of the colliding region before the formation of a quarkonium eigenstate. Cronin effect will shift the momentum distribution of primordially produced hidden- and open-charm (or bottom) states [12,13]. This can be included by a proper modification of their transverse momentum distributions in pp collisions [14,15]. Shadowing effect is weak at RHIC colliding energies, but important at the LHC colliding energies. All the cold nuclear matter effects can be included in the heavy flavor initial distributions prior to the hot medium effects. In nucleus–nucleus collisions at $\sqrt{s_{NN}}=2.76$ TeV and 5.02 TeV, regeneration from charm and anti-charm quarks is widely believed to dominate the prompt charmonium yields [16,17]. This is supported by the enhancement of J/ψ R_{AA} and the suppression of J/ψ mean transverse momentum square $\langle p_T^2\rangle_{J/\psi}$ observed at 2.76 TeV and 5.02 TeV: the regenerated J/ψs from thermalized charm quarks carry small momentum compared with the primordially produced ones, this will pull down the $\langle p_T^2\rangle_{J/\psi}$ of final prompt J/ψ in nucleus–nucleus collisions [18].

However for bottomonium, the situation seems not so clear. Transport model calculations suggest a non-negligible bottomonium regeneration in $\sqrt{s_{NN}}=5.02$ TeV Pb–Pb collisions [19]. Also, experimental data hinted a stronger bottomonium regeneration at 5.02 TeV compared with 2.76 TeV, but within its large uncertainty which prevents solid conclusions [20,21]. Considering that heavy quark mass is very large, it takes some time to reach kinetic thermalization [22,23] in the fast cooling QGP with an initial temperature ~ 500 MeV in AA collisions. Non-thermalization of bottom quark momentum distribution will suppress the combination probability of b and $\bar{b}$ quarks in QGP. However, the ratio of hidden- to open-bottom states is at the order of 0.1% which is smaller than the charm flavor. This may make the yield from $(b+\bar{b}\to\Upsilon(1S)+g,\ b+\bar{b}+\zeta\to\Upsilon(1S)+\zeta$ with $\zeta=g,q,\bar{q})$ not negligible compared with primordially produced $\Upsilon(1S)$. Different from Ref. [19] which simply introduces a factor (smaller than unit) to account for the non-thermalization effect of bottom quarks on Υ regeneration, now we do the realistic evolutions of bottom

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http://dx.doi.org/10.1016/j.physletb.2017.07.054

quarks with the Langevin equation. Its parameters are fixed by the observables of open heavy flavor mesons. In this work, we develop a more realistic model to consider the dynamical evolutions of bottom quarks, bottomonium regeneration process and the following hot medium suppression after they are regenerated. We employ the "Langevin equation + Wigner function + (gluon, quasi-free) dissociations" for bottom quark and bottomonium evolutions in heavy ion collisions. Furthermore, we consider the hot medium modifications on bottomonium properties at finite temperature by taking the color-screened heavy quark potential (extracted from Lattice free energy $F(r,T)$ [24]). With color-screened heavy quark potential in time independent Schrödinger equation, we obtain the mean radius and binding energy of $\Upsilon(1S)$ at different temperatures. Each of them will be used in Wigner function (regeneration rate) and (gluon, quasi-free) dissociation rates.

Our paper is organized as follows. In Section 2, we introduce the Langevin equation and Wigner function for heavy quark dynamical evolutions and their combination. The hot medium modifications on $\Upsilon(1S)$ properties (such as the mean radius and binding energy) are also studied based on potential model. In Section 3, we introduce the hydrodynamic equations for QGP expansion in nucleus–nucleus collisions. The relevant inputs of heavy quarks are presented in Section 4. In Section 5, we give the $\Upsilon(1S)$ nuclear modification factor $R_{AA}^{\rm rege}$ from the combination of b and $\bar{b}$ quarks in QGP. Different coupling strength between heavy quarks and QGP (controlling the heavy quark thermalization) are studied in the $\Upsilon(1S)$ regeneration. We summarize the work in Section 6.

2. Dynamical evolutions of heavy quarks

2.1. Langevin equation for heavy quark diffusion

The heavy quark diffusion in Quark Gluon Plasma can be treated as a Brownian motion, which is widely studied with Langevin equation [25,26]. During the evolution of heavy quarks, they can combine into a quarkonium [2,3,27,28]. The probability of the combination of heavy quark Q and $\bar{Q}$ depends on their distributions in phase space and also the properties of the produced quarkonium at finite temperature [29,30]. Instead of dealing with heavy quark distributions, we employ the Langevin equation to simulate Q and $\bar{Q}$ evolutions in the hot medium. Combination process ($Q+\bar{Q}\to(Q\bar{Q})_{\rm bound}+g$, $Q+\bar{Q}+\zeta\to(Q\bar{Q})_{\rm bound}+\zeta$ with $\zeta=g,q,\bar{q}$) can be included through the Wigner function $W_{Q\bar{Q}\to\Upsilon}$ [3]. The Langevin equation is written as

$$\frac{d\vec{p}}{dt}=-\vec{\eta}_D(p)\vec{p}+\vec{\xi} \tag{1}$$

where $\vec{\eta}_D(p)$ and $\vec{\xi}$ are the drag force and the noise of the hot medium on heavy quarks. $\vec{\xi}$ satisfies the correlation relation

$$\langle\xi^i(t)\xi^j(t')\rangle=\kappa\delta^{ij}\delta(t-t') \tag{2}$$

κ is the diffusion coefficient of heavy quarks in momentum space, which is connected with spatial diffusion coefficient D by $\kappa=2T^2/D$. The drag force in Langevin equation can then be determined by the fluctuation–dissipation relation [31]

$$\eta_D(p)=\frac{\kappa}{2TE} \tag{3}$$

T is the temperature of the bulk medium and $E=\sqrt{m_Q^2+|\vec{p}|^2}$ is the heavy quark energy.

In order to numerically solve the Langevin equation for heavy quark diffusions in QGP, it is discretized as below

$$\vec{p}(t+\Delta t)=\vec{p}(t)-\eta_D(p)\vec{p}\Delta t+\vec{\xi}\Delta t \tag{4}$$

$$\vec{X}_Q(t+\Delta t)=\vec{X}_Q(t)+\frac{\vec{p}}{E}\Delta t \tag{5}$$

$$\langle\xi^i(t)\xi^j(t-n\Delta t)\rangle=\frac{\kappa}{\Delta t}\delta^{ij}\delta^{0n} \tag{6}$$

where $\vec{X}_Q$ is the coordinate of a heavy quark in QGP. As long as the time step Δt for numerical evolutions is small enough, one can assume free motions for heavy quarks with a constant velocity $\vec{v}_Q=\vec{p}/E$ during this time step, and update the heavy quark momentum by Eq. (4) at the end of each Δt due to hot medium effects. The medium-induced radiative energy loss [32,33] and parton elastic collisions [34] of heavy quarks can be included in the terms of the drag force $\vec{\eta}_D(p)$ and the noise $\vec{\xi}$. $\xi^i(t)_{i=1,2,3}$ in Eq. (6) is sampled randomly based on a Gaussian function with the width $\sqrt{\kappa/\Delta t}$.

The initial transverse momentum distribution as an input of Eq. (4) is obtained from PYTHIA simulations. As the initial energy density of QGP changes with coordinates, the heavy quark initial distribution in QGP affects their evolutions, which will finally affect the heavy quark thermalization degree and quarkonium regeneration. Heavy quark pairs are produced from parton hard scatterings, their density (within rapidity region Δy) is proportional to the number of binary collisions,

$$\frac{dN_{\rm PbPb}^{Q\bar{Q}}}{d\vec{x}_T}=\sigma_{pp}^{Q\bar{Q}}(\Delta y)\times T_{\rm Pb}(\vec{x}_T-\frac{\vec{b}}{2})T_{\rm Pb}(\vec{x}_T+\frac{\vec{b}}{2}) \tag{7}$$

where $\sigma_{pp}^{Q\bar{Q}}(\Delta y)$ is the heavy quark production cross section in proton–proton collisions within rapidity Δy. $T_{\rm Pb}(\vec{x}_T)=\int dz\rho_{\rm Pb}(\vec{x}_T,z)$ is the thickness function of lead. Nucleus density $\rho_{\rm Pb}(\vec{x}_T,z)$ is taken to be the Woods–Saxon distribution. When colliding energy is at the order of ∼TeV, theoretical and experimental studies indicate a strong nucleus (anti-)shadowing effect on heavy quark (quarkonium) production [11]. Furthermore, this effect depends on the nucleon density, which gives different modification on heavy quark production at different positions of the nucleus. We employ a theoretical model (EPS09 NLO) to obtain a shadowing factor $r_S(\vec{x}_T,p_T,y)$. The initial distribution of heavy quarks (quarkonium) in Pb–Pb collisions is then taken as $\frac{dN_{\rm PbPb}^{Q\bar{Q}}}{d\vec{x}_T}\times r_S(\vec{x}_T,p_T,y)$.

2.2. Heavy quark recombination process with Wigner function

In the nucleus collisions, heavy quark pairs are produced and evolve inside the QGP as a Brownian motion. During QGP evolutions, Q and $\bar{Q}$ can meet each other and combine into a bound state, which may survive from the hot medium due to its large binding energy. If Q and $\bar{Q}$ can reach kinetic equilibrium, their relative momentum ($\vec{p}_Q-\vec{p}_{\bar{Q}}$) and relative distance ($\vec{X}_Q-\vec{X}_{\bar{Q}}$) will be small, which enhances the combination probability of Q and $\bar{Q}$ [35]. The discussion about connections between heavy quark thermalization and the quarkonium regeneration is left to the next section. Wigner function is widely used for hadron production in the coalescence model. It gives the probability of Q and $\bar{Q}$ combining into a bound state with relative distance $\vec{r}=\vec{X}_Q-\vec{X}_{\bar{Q}}$, and momentum $\vec{q}=\vec{p}_Q-\vec{p}_{\bar{Q}}$, see the function below [3]

$$f(r,q)=A_0\exp(-\frac{r^2}{\sigma^2(T)})\exp(-q^2\sigma^2(T)) \tag{8}$$

A_0 is the normalization factor. Here we neglect the contributions of momentum carried by light partons in the heavy quark formation process, and use the "$b\bar{b}\to\Upsilon$" to represent both $b+\bar{b}\to\Upsilon(1S)+g$ and $b+\bar{b}+\zeta\to\Upsilon(1S)+\zeta$. In realistic simulations, only 0.01% ∼ 0.1% of bottom quarks can form a bound state in the expanding QGP with a lifetime of ∼ 10 fm/c and an initial temperature of ∼ 500 MeV [36]. Most of them become open heavy-flavor hadrons with a light (anti-)quark. Considering the large ratio of $\sigma_{pp}^{b\bar{b}}/\sigma_{pp}^{\Upsilon(1S)} \sim 1000$ (much larger than the charm flavor ∼ 200), this combination process may be important for the bottomonium nuclear modification factor R_{AA}. Bottomonium properties such as binding energy and shape of the wave function can be modified by the hot medium. This can affect the probability of heavy quark combination. We include hot medium effects on bottomonium properties by introducing the temperature dependence of Gaussian width $\sigma(T)$. It is connected with the bottomonium mean radius square at finite temperature by $\sigma(T) = \sqrt{8\langle r^2\rangle_\Upsilon(T)/3}$ [3].

In order to study the bottomonium regeneration in QGP, we put one b and $\bar{b}$ in QGP and evolve each of them with Langevin equation to obtain the probability $W^{b\bar{b}\to\Upsilon}$ of their combination to regenerate a $\Upsilon(1S)$. The total yield of regenerated bottomonium within rapidity Δy in nucleus–nucleus collisions is scaled by the number of bottom pairs as

$$N_{\rm PbPb}^{b\bar{b}\to\Upsilon}|_{\Delta y} = \sigma_{pp}^{b\bar{b}}|_{\Delta y} N_{coll} \times W_{b\bar{b}\to\Upsilon} \quad (9)$$

where N_{cill} is the number of binary collisions. The average number of $b\bar{b}$ pairs produced in central rapidity region is only ∼ 1, the regenerated bottomonium yield is proportional to the number of $b\bar{b}$ pairs. The initial momentum and position of heavy quarks are randomly generated based on the distributions from PYTHIA simulations and Eq. (7) (both with modifications of the shadowing effect). At each time step, one can obtain their relative distance r and relative momentum q, and their combination probability $P(r,q) = r^2q^2 f(r,q)$. If the probability is larger than a random number between 0 and 1, then the formation process of bottomonium happens. Otherwise, they continue the evolutions with Langevin equation independently until moving out of QGP (hadronization as open bottom hadrons). After $\Upsilon(1S)$ is regenerated inside QGP, it will decay due to parton inelastic scatterings and color screening from QGP. We supplement this part by the rate equation $dN_\Upsilon^{\rm rege}/dt = -\Gamma_\Upsilon^{\rm diss}(T)N_\Upsilon$, where $\Gamma_\Upsilon^{\rm diss}$ is the decay rate from gluon and quasi-free dissociations. It is connected with Wigner function (see Eq. (8)). $N_\Upsilon^{\rm rege}(t)$ is the number of regenerated Υ at time t.

After one $\Upsilon(1S)$ is regenerated at the time t_0, the initial condition of Eq. (10) becomes $N_\Upsilon^{\rm rege}(t_0) = 1$, and $N_\Upsilon^{\rm rege}(t)$ decreases with time based on the rate equations,

$$N_\Upsilon^{\rm rege}(t_1+\Delta t) = N_\Upsilon^{\rm rege}(t_1)e^{-\Gamma^{\rm diss}(T)\Delta t} \quad (10)$$

$$\vec{R}_\Upsilon(t_1+\Delta t) = \vec{R}_\Upsilon(t_1) + \frac{\vec{P}_\Upsilon}{E_\Upsilon}\Delta t \quad (11)$$

Note that $N_\Upsilon^{\rm rege}(t \geq t_0) \leq 1$ where t_0 is the time of $\Upsilon(1S)$ regeneration. $\vec{P}_\Upsilon \approx \vec{p}_b + \vec{p}_{\bar{b}}$ and $\vec{R}_\Upsilon$ are the momentum and coordinate of the center of the regenerated $\Upsilon(1S)$. From hydrodynamic equations, the local temperature of QGP is different at different coordinate $\vec{R}$ and time. Therefore, we update $\Upsilon(1S)$ position at each time step, and take a new local temperature T of QGP at $\vec{R}_\Upsilon(t_1+\Delta t)$ to recalculate the decay rate $\Gamma^{\rm diss}$ for $\Upsilon(1S)$ evolution at the next time step. We continue Eqs. (10)–(11) from the time t_0 of $\Upsilon(1S)$ regeneration until it moves out of QGP (where local temperature is smaller than the critical temperature T_c). Doing sufficient events, $N_{b\bar{b}}^{\rm events}$, of putting one b and $\bar{b}$ in the expanding QGP, and sum $\Upsilon(1S)$ final yields to be $N_\Upsilon^{\rm rege}(\mathrm{tot})$, one can obtain the $\Upsilon(1S)$ regeneration probability from b and $\bar{b}$ evolutions in QGP, $W_{b\bar{b}\to\Upsilon} = N_\Upsilon^{\rm rege}(\mathrm{tot})/N_{b\bar{b}}^{\rm events}$.

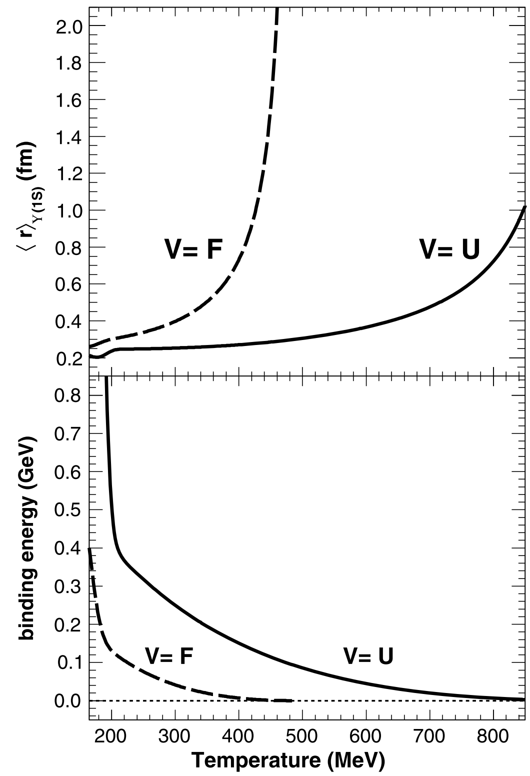

Fig. 1. Upper panel: mean radius of $\Upsilon(1S)$ at finite temperature with the heavy quark potential to be two limits: free energy $V = F$ (dashed line) and internal energy $V = U$ (solid line). Lower panel: binding energy of $\Upsilon(1S)$ with $V = F$ and $V = U$ respectively.

2.3. Bottomonium dissociation and regeneration rates at finite temperature

QGP color screening reduces the binding energy of bottomonium, and increases its decay rate at finite temperature. With heavy quark potential to be $V_{b\bar{b}} = F$ or $V_{b\bar{b}} = U$ [37], one can obtain the $\Upsilon(1S)$ binding energy from solving time-independent radial Schrödinger equation (with $\hbar = c = 1$)

$$[-\frac{1}{2m_\mu}\frac{\partial^2}{\partial r^2} + V_{b\bar{b}}(r,T)]\psi(r) = \mathcal{E}_\Upsilon(T)\psi(r) \quad (12)$$

Here $m_\mu = m_b/2$ is the reduced mass in the center of Υ mass frame. $\mathcal{E}_\Upsilon(T)$ and $\psi(r)$ are the binding energy and radial wave function of a Υ eigenstate. The mean radius and binding energy (see Fig. 1) obtained consistently from Eq. (12) will be used in the Wigner function for bottomonium regeneration and the parton dissociations, respectively.

The decay rates from gluon dissociation and quasi-free dissociation [38,39] with temperature dependent binding energy are plotted in Fig. 2. When the $\Upsilon(1S)$ is strongly bound, gluon dissociation dominates the $\Upsilon(1S)$ decay rate, such as at the temperature region of $T < 0.2$ GeV with $V_{b\bar{b}} = U$ (see red dashed and solid lines). However at a high temperature like $T = 300$ MeV, strong color screening effect reduces the $\Upsilon(1S)$ binding energy to be around 0.25 GeV with $V = U$ and 0.04 GeV with $V = F$, which are far below the vacuum value ∼ 1.1 GeV [40]. This makes quasi-free dissociation dominates the decay rate. We include both contributions on the hot medium suppression of regenerated $\Upsilon(1S)$ at the entire temperature region $T > T_c$.

The realistic heavy quark potential is between $V = F$ and $V = U$. Strong binding limit ($V = U$) corresponds to the fast dissociation of $\Upsilon(1S)$ and assumes no energy exchange between $\Upsilon(1S)$

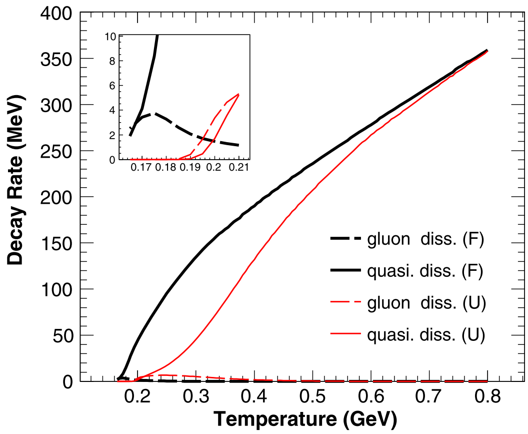

Fig. 2. (Color online.) Decay rates of $\Upsilon(1S)$ as a function of temperature. Dashed and solid lines are the decay rates from gluon and quasi-free dissociations respectively. Thick black and thin red lines are with heavy quark potential to be the free energy F and internal energy U.

and the thermal bath during the dissociation process. Compared with the weak binding limit ($V=F$) which assumes a very slow dissociation of $\Upsilon(1S)$, the limit of $V=U$ gives a larger binding energy of $\Upsilon(1S)$. In that case, it takes more energy to break $\Upsilon(1S)$ after its regeneration. Therefore, more $\Upsilon(1S)$ can be regenerated and survive from QGP with the heavy quark potential to be $V=U$, see Fig. 4.

3. Hydrodynamic model

We employ the $(2+1)$ dimensional ideal hydrodynamics to simulate strong expansion of finite sized QGP produced in ultra-relativistic heavy ion collisions,

$$\partial_\mu T^{\mu\nu}=0 \tag{13}$$

Here $T^{\mu\nu}=(e+p)u^\mu u^\nu - g^{\mu\nu}p$ is the energy–momentum tensor, and (e,p,u^μ) are the energy density, pressure and four-velocity of fluid cells. The equation of state of the deconfined medium is taken as an ideal gas of massless (u,d) quarks, 150 MeV massed s quarks and gluons [41]. Hadron phase is an ideal gas of all known hadrons and resonances with mass up to 2 GeV [42]. From the scaling of initial temperature at $\sqrt{s_{NN}}=2.76$ TeV Pb–Pb collisions which is $T_0=485$ MeV, we set the initial maximum temperature at $\sqrt{s_{NN}}=5.02$ TeV to be $T_0=510$ MeV [43]. Based on hydrodynamic model studies, light hadron spectra at RHIC 200 GeV Au–Au and LHC 2.76 TeV Pb–Pb collisions indicate a same time scale of QGP reaching local equilibrium $\tau_0^{200\ \mathrm{GeV}}\approx\tau_0^{2.76\ \mathrm{TeV}}\approx 0.6$ fm/c [44,45]. Therefore, we still take the same value of $\tau_0=0.6$ fm/c at 5.02 TeV Pb–Pb collisions due to its weak dependence on colliding energy. The transverse expansion of QGP controlled by Eq. (13) starts from τ_0.

4. Inputs of bottom flavor

The bottomonium regeneration requires the number of bottom quarks in nucleus–nucleus collisions. We determine this by using the production cross section in pp collisions and binary collision scaling in Pb–Pb collisions, $N_{\mathrm{PbPb}}^{b\bar b}=\sigma_{pp}^{b\bar b}N_{coll}(b)$. Lack of experimental data about $\sigma_{pp}^{b\bar b}$ at 5.02 TeV pp collisions, we extract its value by the linear interpolation between the cross sections at central rapidity at 1.96 TeV and 2.76 TeV collisions. At 1.96 TeV pp collisions, CDF collaboration published the cross section of b-hadrons integrated over all transverse momenta in the rapidity $|y|<0.6$ to be $17.6\pm0.4(stat)^{+2.5}_{-2.3}(syst)$ μb [46]. With this we extract the central value of the differential cross section to be $d\sigma_{pp}^{b\bar b}/dy=14.7$ μb. Combined with the $d\sigma_{pp}^{b\bar b}/dy=23.28\pm2.70(stat)^{+8.92}_{-8.70}(syst)$ μb in the central rapidity at 2.76 TeV [47], we obtain the differential cross section $d\sigma_{pp}^{b\bar b}/dy=47.5$ μb in the central rapidity at 5.02 TeV pp collisions. Our purpose is to study the contribution of regeneration component in the experimentally measured inclusive $\Upsilon(1S)$ yields, presented as the nuclear modification factor [15],

$$R_{AA}^{\mathrm{inclu}}(\Upsilon(1S))=\frac{N_{AA}^{\mathrm{prim}}(\Upsilon(1S))+N_{AA}^{b\bar b\to\Upsilon(1S)}}{\frac{d\sigma_{pp}^{\Upsilon(1S)}}{dy}\Delta y\cdot N_{coll}(b)} = R_{AA}^{\mathrm{prim}}+R_{AA}^{\mathrm{rege}} \tag{14}$$

$$N_{AA}^{b\bar b\to\Upsilon(1S)}=W_{b\bar b\to\Upsilon(1S)}^{\mathrm{Lan+Wigner}}\left(\frac{d\sigma_{pp}^{b\bar b}}{dy}\Delta y\cdot N_{coll}\right) \tag{15}$$

where N_{AA}^{prim} in the numerator of Eq. (14) represents the $\Upsilon(1S)$ primordial production including direct production and decay contributions from excited states $(1P,2P,2S,3S)$. The second term $N_{AA}^{b\bar b\to\Upsilon(1S)}$ is for the *new* bottomonium from b and $\bar b$ combination during QGP evolutions, which is closely connected with the bottom quark diffusions in the expanding QGP and so our main interest in this work. $W_{b\bar b\to\Upsilon(1S)}^{\mathrm{Lan+Wigner}}$ is the probability of one b and $\bar b$ quark combine into a $\Upsilon(1S)$. $d\sigma_{pp}^{\Upsilon(1S)}/dy$ in the denominator of Eq. (14) is the $\Upsilon(1S)$ inclusive cross section. With the differential cross sections $d\sigma_{pp}^{\Upsilon(1S)}/dy=27\pm1.5$ nb at 1.8 TeV [48] and $d\sigma_{pp}^{\Upsilon(1S)}/dy=80\pm9$ nb at 7 TeV from CMS Collaboration [49] in the central rapidity of pp collisions, we extract the central value of inclusive cross section to be $d\sigma_{pp}^{\Upsilon(1S)}/dy=59.8$ nb at 5.02 TeV, which gives the ratio of $N_{pp}^{\Upsilon(1S)}/N_{pp}^{b\bar b}$ in central rapidity to be 0.13%, close to the typical order 0.1% of hidden- to open-bottom state ratio in elementary hadronic collisions [50].

The coupling strength between bottom quarks and QGP is indicated by the drag coefficient in Langevin equation. The spatial diffusion coefficient is taken to be $D(2\pi T)=C$ [51]. Different values of C will be employed to study the effects of bottom quark thermalization on bottomonium regeneration.

5. Bottomonium continuous regeneration in Pb–Pb collisions

In the previous work [36], we employ the Langevin equation plus Wigner function to calculate J/ψ and $\psi(2S)$ regeneration from dynamical evolutions of (anti-)charm quarks, with an assumption that they are produced at each certain temperature, without the following suppression from hot medium after their regeneration. This can well describe the relationship between elliptic flows of regenerated charmonium eigenstates $(1S,2S)$, but can not give a good description of their yields, due to the lack of hot medium suppression effect. In this work, we improve our approach of "Langevin equation + Wigner function" by considering hot medium modifications on quarkonium properties, which change quarkonium regeneration and dissociation rates through the mean radius $\langle r\rangle_\Upsilon(T)$ and binding energy $\mathcal{E}_\Upsilon(T)$ of quarkonium.

In the realistic simulations, we generate one b and $\bar b$ randomly in the coordinate and momentum space based on the probability distributions given in previous sections. Then we evolve them separately with two individual Langevin equations, and check if they can form a $\Upsilon(1S)$ at each time step. In central collisions, b and $\bar b$ are easier to lose energy and meet each other to form a new $\Upsilon(1S)$. Smaller value of the parameter C indicate a stronger coupling strength between bottom quarks and QGP, which results in larger regeneration rate $W_{b\bar b\to\Upsilon}^{\mathrm{Lan+Wigner}}$.

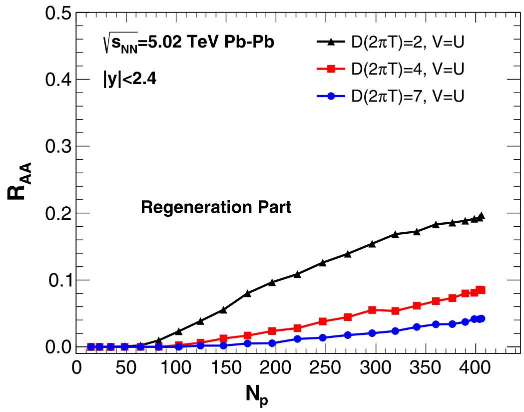

Fig. 3. (Color online.) Nuclear modification factor of regenerated $\Upsilon(1S)$ from b and $\bar{b}$ quark combination as a function of number of participants N_p in central rapidity region $|y| < 2.4$ at 5.02 TeV Pb–Pb collisions. The heavy quark potential is taken as internal energy $V = U$ (used to determine the mean radius and binding energy of $\Upsilon(1S)$). Lines with triangle, square and circle markers are for diffusion coefficient $D(2\pi T) = 2, 4, 7$ respectively. Note that these D values satisfy the relation of $1 \lesssim D(2\pi T) \lesssim 7$ from pQCD and Lattice calculations [52–54].

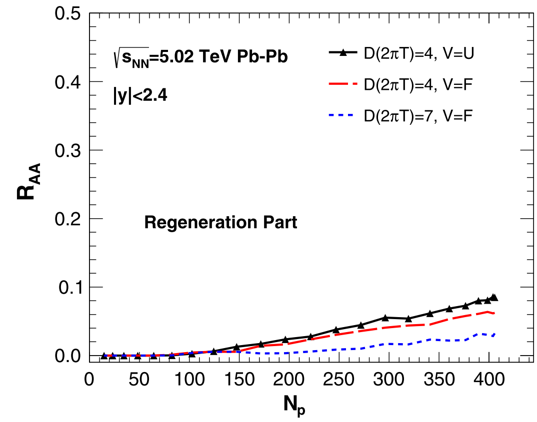

Fig. 4. (Color online.) Nuclear modification factor of regenerated $\Upsilon(1S)$ from b and $\bar{b}$ quark combination as a function of number of participants N_p in central rapidity region $|y| < 2.4$ at 5.02 TeV Pb–Pb collisions. The heavy quark potential is taken as $V = U$ and $V = F$.

Now we do the full calculations of bottomonium regeneration in heavy ion collisions, and give the regeneration part (see Eqs. (14)–(15)) of inclusive $R_{AA}^{\Upsilon(1S)}$ in Fig. 3. In central collisions where QGP temperature is high, both number of $b\bar{b}$ pairs and the combination probability of b and $\bar{b}$ quarks become large in Pb–Pb collisions. This makes $R_{AA}^{\Upsilon(1S)}$ increase with N_p. The slope of $R_{AA}^{\Upsilon(1S)}(N_p)$ is larger than the slopes of $N_{\mathrm{PbPb}}^{b\bar{b}}(N_p)$ and $W_{b\bar{b}\to\Upsilon}^{\mathrm{Lan+Wigner}}(N_p)$. From the nuclear modification factor of B-hadrons, the situation of $D(2\pi T) = 4$ seems better to describe the heavy quark energy loss in $\sqrt{s_{NN}} = 5.02$ TeV Pb–Pb collisions [51]. In the most central collisions, the regeneration contributes to the inclusive $R_{AA}^{\Upsilon(1S)}$ with around $(10 \sim 20)\%$ ($V = U$) and $(5 \sim 10)\%$ ($V = F$). But it is only $\sim 4\%$ at minimum bias centrality (with $N_p \approx 200$). Note that transport model calculations gave the regeneration R_{AA} to be $\sim 8\%$ in semi-central ($N_p \approx 200$) and most central Pb–Pb collisions in the rapidity $|y| < 2.4$ at 2.76 TeV [19, 50].

One way to testify the contribution of $\Upsilon(1S)$ regeneration in heavy ion collisions is to study the rapidity dependence of the p_T−integrated $R_{AA}^{\Upsilon(1S)}(y)$. The bottom quark differential cross section decreases with rapidity, which will suppress the $\Upsilon(1S)$ regeneration at forward rapidity. For charmonium, the decreasing tendency of $R_{AA}^{J/\psi}(y)$ with rapidity is very strong and explained well by the regeneration mechanism [15]. As charmonium regeneration mainly dominates at the low p_T bins and drops to zero at middle and high p_T bins. Therefore, $R_{AA}^{J/\psi}$ shows strong decreasing tendency with rapidity at $p_T > 0$ (where regeneration dominates) and almost no rapidity dependence at $p_T \gtrsim 4$ GeV/c. Considering the fraction of $\Upsilon(1S)$ regeneration is only $10 \sim 15\%$ in its inclusive R_{AA} at the impact parameter $b = 0$, we expect a weaker rapidity dependence of $R_{AA}^{\Upsilon(1S)}(y)$. In the minimum bias centrality, this tendency should be much weaker and $R_{AA}^{\Upsilon(1S)}(y)$ almost shows a flat feature, just as experimental data shows [20,55].

We also did the calculations in $\Upsilon(1S)$ weak binding scenario ($V = F$), see Fig. 4. With weak binding ($V = F$), $\Upsilon(1S)$ can only be regenerated at $T < 400$ MeV where its binding energy is nonzero. Also, the dissociation rate of regenerated $\Upsilon(1S)$ is larger for $V = F$ compared with $V = U$, which makes regenerated $\Upsilon(1S)$ easier to be dissociated. $\Upsilon(1S)$ regeneration with $V = F$ is smaller (see Fig. 4). This is also consistent with transport model calculations. The $\Upsilon(1S)$ R_{AA} from regeneration contribution is around 6% (with $D(2\pi T) = 4$) and 3% (with $D(2\pi T) = 7$) in the most central collisions. In all centralities, its contribution is smaller than the value of $V = U$, see Fig. 4.

Considering that $D(2\pi T) = 4 \sim 7$ can well explain the nuclear modification factor of open heavy flavor meson, we expect that regeneration contribute ~ 0.1 to the final inclusive $R_{AA}^{\Upsilon(1S)}$. This is also consistent with the results in Ref. [56] where R_{AA}^{rege} is around 0.1.

6. Summary

We employ the Langevin equation to describe the diffusion of bottom quarks, and Wigner function for the b and $\bar{b}$ quarks to regenerate $\Upsilon(1S)$s during the QGP evolutions. After the regeneration of a $\Upsilon(1S)$, it also suffers the color screening and parton inelastic scatterings from QGP. The dissociation and regeneration rates of $\Upsilon(1S)$ are connected with each other by the temperature dependent binding energy $\mathcal{E}_\Upsilon(T)$ and mean radius $\langle r\rangle_\Upsilon(T)$, which can be obtained simultaneously from Schrödinger equation with the color screened heavy quark potential extracted from Lattice calculations. Including cold nuclear matter effects, we give the full calculations of $\Upsilon(1S)$ regeneration in Pb–Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. With different drag coefficient in Langevin equation, the energy loss of b and $\bar{b}$ quarks is different. Strong coupling between bottom quarks and QGP can make b and $\bar{b}$ lose more energy, and increases the probability of their formation into a bound state. In the scenario of $V = U$ ($V = F$), we obtain the $\Upsilon(1S)$ regeneration R_{AA}^{rege} to be $0.1 \sim 0.2$ ($0.06 \sim 0.1$) in the most central collisions, but negligible in minimum bias centrality. With realistic evolutions of bottom quarks and their hadronization process, we study the regeneration contribution to the $\Upsilon(1S)$ nuclear modification factor R_{AA} measured in experiments, and also build the connection between bottom quark energy loss and bottomonium regeneration in this work.

Acknowledgements

Baoyi Chen thanks X. Du for helpful discussions. We acknowledge that TAMU group is also studying bottomonium regeneration based on Langevin equation independently. Please see their manuscript [56]. This work is supported by the NSFC under Grant No. 11547043.

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