B_c meson enhancement and the momentum dependence in Pb + Pb collisions at energies available at the CERN Large Hadron Collider

Yunpeng Liu,[*] Carsten Greiner, and Andriy Kostyuk

Institut für Theoretische Physik, Johann Wolfgang Goethe-Universität Frankfurt, Max-von-Laue-Strasse 1,
D-60438 Frankfurt am Main, Germany

B_c meson production in Pb + Pb collisions at $\sqrt{s} = 2.76\,A$ TeV is surveyed in both a statistical coalescence model and a transport model. The nuclear modification factor R_{AA} is predicted to be between 2 and 18 in the most central collisions, which can help to confirm the regeneration mechanism. In addition, the momentum dependence is also investigated as given by the transport model. A strong suppression of the transverse momentum is found in central collisions accompanying the enhancement in yield. The spectrum and elliptic flow of B_c are also discussed.

I. INTRODUCTION

Several kinds of signals of the hot quark matter created in relativistic heavy-ion collisions are widely discussed in the literature. One of them is based on the hard probes, for example, quenching [1] of the jets which pass through the hot medium. Another one is focused on the low-energy particles, for example, the collective flow of low-p_T particles [2], which carries the information on hydrodynamical properties of the hot matter at the initial stage of the fireball evolution.

One more significant idea utilizes heavy quarkonia as hard probes of the fireball. The anomalous J/ψ suppression in quark-gluon plasma (QGP) was predicted theoretically [3] a quarter of a century ago. Such a suppression was indeed observed in heavy-ion collisions [4]. It was however suggested that this suppression can be as well explained within the models that do not assume quark deconfinement [5–9]. Moreover, it was suggested that not only dissociation of quarkonia but also the opposite process, recombination of heavy quark-antiquark pairs, can take place in the deconfined medium [10–19]. If the regeneration of quarkonia indeed takes place, it becomes especially significant with increasing collision energy since the number of heavy quarks per collision becomes larger. When the regeneration is dominant, quarkonia can be used to detect the early fireball through its spectrum and flow like the soft probes. Even enhancement of quarkonia production has been expected [14], which could be regarded as direct evidence of the regeneration mechanism. However, such an enhancement of heavy quarkonia is not supported at Relativistic Heavy Ion Collider (RHIC) energy by the statistical model [11,20] and has never been observed at RHIC or the Large Hadron Collider (LHC) [21–26].

How can we verify the regeneration mechanism firmly?

The B_c meson was found by the Collider Detector at Fermilab (CDF) in 1998 [27]. Similar to heavy quarkonia it consists of a heavy quark and a heavy antiquark. But unlike in J/ψ or Υ mesons, the quark and antiquark have different flavors. This has a drastic effect on the production

cross section of the B_c meson in elementary hadron-hadron collisions. Indeed, creation of one $c\bar{c}$ or $b\bar{b}$ in an elementary collision is sufficient for production of, respectively, J/ψ or Υ. In contrast, production of a B_c meson in hadron-hadron collisions requires creation of at least two heavy quark pairs, $c\bar{c}$ and $b\bar{b}$, in the same collision.

Unlike in elementary hadron-hadron reactions, B_c meson production in nucleus-nucleus collisions does not require production of two heavy quark pairs in the same elementary collision. Heavy quarks and antiquarks originating from different nucleon-nucleon collisions can recombine and form a B_c meson. If such regeneration is possible, it leads to significant enhancement of B_c production in nucleus-nucleus collisions relative to proton-proton ones.

It was suggested that, due to the regeneration, the B_c meson can be observed in Au + Au collisions at RHIC, although its observation in proton-proton collisions at the same energy is hardly possible [28]. This prediction has yet to be verified in experiments.

The high beam energy of LHC makes it possible to measure B_c mesons in both $p + p$ and Pb + Pb collisions. The cross section for bottom quark production in proton-proton collisions is about one order of magnitude smaller than that for charm quarks, while the cross sections of B_c and Υ are about two orders of magnitude less than that of J/ψ [29–32]. For a very rough estimation, the yield of a meson in the regeneration process can be assumed to be proportional to the yield of each of its constituent quarks. Therefore if regeneration is present in nucleus-nucleus collisions, the yield of Υ mesons would remain two order of magnitude smaller than that of J/ψ, while the multiplicity of B_c mesons should be only one order of magnitude smaller in comparison to J/ψ. Thus the nuclear modification factor R_{AA} for B_c can be roughly one order of magnitude larger than that of J/ψ and Υ, that is, $R_{AA} > 1$, which implies enhancement of B_c instead of suppression. If such an effect is found, it will be a firm confirmation of the regeneration.

In this paper, we will calculate the yield of B_c mesons in central Pb + Pb collisions in the statistical coalescence model, and then discuss more properties including the momentum dependence in a detailed transport model. We take $\hbar = c = k_B = 1$ in the following.

————
[*]liuyp06@mails.tsinghua.edu.cn

II. BASELINE

B_c mesons are similar to quarkonia in the sense that they are composed of heavy quarks, so that the nonrelativistic approximation can be applied to describe the interaction between them. On the other hand, the decay modes of B_c are totally different from those of quarkonia. For excited B_c states with mass below the $B + D$ threshold, the only decay mode, except weak decay, is feeding down to the lower states until the ground state is reached, due to the conservation of charm and bottom. Thus the contribution from excited states is even more significant than in the case of quarkonia. The properties of ground and excited B_c mesons are discussed in different potential models [33]. In order to be simple and generalize it to finite temperature later in the transport model, we calculate the spectrum of B_c by solving the Schrödinger equation with a Cornell potential $V(r) = -\alpha/r + \sigma r$. By using this method, the mass spectra of charmonia and bottomonia can be well reproduced with the parameters $\alpha = \pi/12$, $\sigma = 0.2$ GeV^2, $m_c = 1.25$ GeV, and $m_b = 4.25$ GeV [34]. With exactly the same parameters, we can obtain the mass spectrum of B_c below the threshold as $m_{B_c}(1S) = 6.36$ GeV, $m_{B_c}(1P) = 6.72$ GeV, $m_{B_c}(2S) = 6.90$ GeV, and $m_{B_c}(1D) = 6.98$ GeV. Since the binding energy of the one-dimensional state in vacuum is only about 170 MeV $\sim T_c$, we neglect the contribution of this state in the following for simplicity. The contribution of this state in the statistical coalescence model is only a few percent. Note that there are already experimental data for B_c^+ with quantum number $J^P = 0^-$ and mass $m = 6.28$ GeV [35]; we use the experimental value instead in the following for the $1S$ state.

The cross section of B_c^+ at $\sqrt{s} = 1.96$ TeV as measured by the CDF Collaboration is $d\sigma/dy$ ($p_T > 6$ GeV) $= 15.5 \pm 5.0$ nb with rapidity $|y| < 1$ [32]. According to PYTHIA [36], the cross section for $p_T > 0$ at LHC energy $\sqrt{s} = 2.76$ TeV is larger by the factor of about 4. Thus we obtain the inclusive cross section of B_c as $d\sigma/dy(p_T > 0) = 62$ nb. The cross section of charm and bottom are estimated as $d\sigma_{c\bar{c}}/dy|_{y=0} = 620$ μb [29,30] and $d\sigma_{b\bar{b}}/dy|_{y=0} = 20$ μb [31], respectively. The elliptic flow and high-p_T suppression of open heavy flavors imply strong interaction between heavy quarks and the fireball [37]; therefore we simply take the thermal momentum distribution for the heavy quarks.

In the transport model, the momentum distribution of the initially produced B_c's is required. It is parametrized in the power-law form

$$\left.\frac{d^2\sigma}{dy\,dp_T}\right|_{y=0}(p_T)$$
$$= \frac{2(n-1)}{(n-2)\langle p_T^2\rangle}p_T\left(1 + \frac{p_T^2}{(n-2)\langle p_T^2\rangle}\right)^{-n}\left.\frac{d\sigma}{dy}\right|_{y=0}, \quad (1)$$

with $n = 4.16$ and $\langle p_T^2\rangle = 25.1$ GeV^2 estimated from PYTHIA [36] simulation. To consider the reaction rates of different B_c states, the cross section and the branch ratios of each of them are necessary. Since all the excited states feed down to the ground state without weak decays considered, we take the branch ratio as 100%. The ratio of direct production cross section of ψ' to J/ψ is around 0.23, while the same ratios for $\Upsilon(2S)$ and $\Upsilon(3S)$ to $\Upsilon(1S)$ are both above 0.35 [38].

As a rough estimation, we take the cross section of each excited component of B_c as 0.3 times that of the ground state; that is, $\sigma^{\rm dir}(1S) : \sigma^{\rm dir}(1P) : \sigma^{\rm dir}(2S) = 10:9:3$, where we have counted the degeneracy. To include the Cronin effect, a Gauss smearing is used to modify the initial momentum dependence so that the mean p_T^2 of the J/ψ in $A + A$ collisions is larger than that in $p + p$ collisions, that is, $\langle p_T^2\rangle_{AA} = \langle p_T^2\rangle_{pp} + a_{gN}l$, where l is the total path length the gluons pass through the nuclei before merging into a B_c, with the broadening factor a_{gN} taken as 0.2 GeV^2/fm. Since the initially produced B_c's suffer strong suppression, the dependence of the final observations on the parameters discussed in this paragraph for the initial production is very weak.

III. STATISTICAL COALESCENCE MODEL

The statistical hadronization model (SHM) and statistical coalescence model (SCM) have been greatly successful is describing light and heavy hadrons, respectively, with only a few parameters [11,39]. This suggested the idea that hadronization of heavy hadrons can also be described within the statistical approach. In contrast to light quarks, the masses of c and b quarks are much larger than the typical temperature of the fireball. Therefore, production of $c\bar{c}$ and $b\bar{b}$ pairs at the thermal stage of the reaction can be neglected, though pre-equilibration can slightly alter the picture for $c\bar{c}$ production at LHC energy [40]. Practically all heavy flavor pairs are produced at the initial stage of the nucleus-nucleus reaction in hard parton collisions. It is assumed that their numbers remains approximately unchanged during the fireball evolution.

The basic idea of the SCM for heavy flavors is expressed in the following balance equation [11]:

$$N_{Q\bar{Q}}^{\rm dir} = \tfrac{1}{2}g_Q N_{oQ}^{\rm th} + g_Q^2 N_{Q\bar{Q}}^{\rm th}, \quad (2)$$

where $N_{Q\bar{Q}}^{\rm dir}$ is the number of directly produced heavy quarks, and $N_{oQ}^{\rm th}$ and $N_{Q\bar{Q}}^{\rm th}$ are numbers of open and hidden heavy flavors for the hadron gas in complete (including heavy flavors) chemical equilibrium and zero value of the corresponding heavy flavor chemical potential. A fugacity g_Q is introduced to describe both the conservation of Q and that of $\bar{Q}$. When $N_{Q\bar{Q}}$ is small, the event-by-event fluctuation becomes important. This effect can be included by a modification factor $(1 + 1/N_{Q\bar{Q}})$ in the hidden part, that is, $N_{Q\bar{Q}} = (1 + 1/N_{Q\bar{Q}}^{\rm dir})g_Q^2 N_{Q\bar{Q}}^{\rm th}$, as long as the heavy flavor is mainly open heavy hadrons [12]. This modification leads to results similar to those of the widely used canonical ensemble modification [41] with a deviation within several percent. In order to take the B_c meson into consideration, we generalize the above equation into the following form:

$$N_{c\bar{c}}^{\rm dir} = \frac{1}{2}g_c\big(N_{\rm oc}^{\rm th} + g_b N_{B_c}^{\rm th}\big) + \left(1 + \frac{1}{N_{c\bar{c}}^{\rm dir}}\right)g_c^2 N_{c\bar{c}}^{\rm th}, \quad (3)$$

$$N_{b\bar{b}}^{\rm dir} = \frac{1}{2}g_b\big(N_{\rm ob}^{\rm th} + g_c N_{B_c}^{\rm th}\big) + \left(1 + \frac{1}{N_{b\bar{b}}^{\rm dir}}\right)g_b^2 N_{b\bar{b}}^{\rm th}, \quad (4)$$

where $N_{\rm oc}$ and $N_{\rm ob}$ are the yield of hadrons with charm and bottom except B_c, respectively. This formula can easily be

generalized for hadrons with even more heavy quarks. The leading order approximation for B_c is simply

$$N_{B_c} = \frac{(2N_{c\bar{c}}^{\mathrm{dir}})(2N_{b\bar{b}}^{\mathrm{dir}})}{N_{\mathrm{ob}}^{\mathrm{th}}N_{\mathrm{oc}}^{\mathrm{th}}} N_{B_c}^{\mathrm{th}}. \quad (5)$$

Higher orders account for the strict conservation of heavy quarks and the balance between open and hidden heavy flavors. In the following calculation, the full form [Eqs. (3) and (4)] is solved.

The temperature and volume of the fireball are taken as $T = 164$ MeV and $V_{\Delta y=1} = 4160$ fm^3, respectively [42], for central Pb + Pb collisions at $\sqrt{s} = 2.76\,A$ TeV, and $N_{c\bar{c}}$ and $N_{b\bar{b}}$ are derived from the Glauber model with the cross sections discussed previously. All the listed charm hadron and bottom hadron with a mass m and spin J in the particle list from the Particle Data Group [35] are included. The results are as follows: $g_c = 31.1$, $g_b = 2.39 \times 10^8$, $N_{\mathrm{ob}}/(2N_{b\bar{b}}^{\mathrm{dir}}) = 98.3\%$, $N_{B_c^+}/N_{b\bar{b}}^{\mathrm{dir}} = 1.08\%$, and $N_{\Upsilon s}/N_{b\bar{b}}^{\mathrm{dir}} = 0.60\%$.[1] Note that

$$\frac{N_{B_c}^{\mathrm{dir}}}{N_{b\bar{b}}^{\mathrm{dir}}} = \frac{\sigma_{B_c}^{\mathrm{dir}}}{\sigma_{b\bar{b}}^{\mathrm{dir}}} = \frac{62 \text{ nb}}{20\,\mu\text{b}} = 0.31\%. \quad (6)$$

The nuclear modification factor can be calculated as

$$R_{AA} = \frac{N_{B_c}}{N_{B_c}^{\mathrm{dir}}} = \frac{N_{B_c}/N_{b\bar{b}}^{\mathrm{dir}}}{N_{B_c}^{\mathrm{dir}}/N_{b\bar{b}}^{\mathrm{dir}}} = \frac{1.08\%}{0.31\%} = 3.5. \quad (7)$$

Only the scalar component of B_c mesons has been observed in experiments. If we take the probably existent vector state of B_c and the isospin partner of B^* into consideration, the results are modified as follows: $g_c = 31.1$, $g_b = 1.83 \times 10^8$, $N_{\mathrm{ob}}/(2N_{b\bar{b}}^{\mathrm{dir}}) = 96.4\%$, $N_{B_c^+}/N_{b\bar{b}}^{\mathrm{dir}} = 3.30\%$, and $N_{\Upsilon s}/N_{b\bar{b}}^{\mathrm{dir}} = 0.35\%$, which results in $R_{AA} = 3.30/0.31 = 11$. Furthermore, if $B_c(1P)$ and $B_c(2S)$ states are also considered, then R_{AA} can be as large as 13. Thus from the SCM, we do expect an enhancement of B_c mesons in central Pb + Pb collisions at LHC energy with $R_{AA} = 3.5$–13.

IV. TRANSPORT MODEL

In contrast to the statistical coalescence model, in which statistical equilibration of the heavy quark distribution among the hadrons at chemical freeze-out at hadronization is assumed, the transport model takes into account the deviation from the statistical equilibrium and traces the whole evolution of the fireball, and thus it can give more detail on the various processes in the QGP. In this model, the distribution of B_c in phase space is described by the function $f_{B_c}(\mathbf{x}, \mathbf{p}, t)$ satisfying the transport equation

$$(\partial_t + \mathbf{v} \cdot \nabla) f_{B_c} = -\alpha f_{B_c} + \beta, \quad (8)$$

where $\mathbf{v} = \mathbf{p}/E_{B_c}$ is the velocity of B_c. The dissociation rate can be expressed as

$$\alpha(\mathbf{x}, \mathbf{p}, t) = \frac{1}{E_{B_c}} \int \frac{d\mathbf{k}}{(2\pi)^3 E_k} k_\mu p^\mu f_g^{\mathrm{th}} \sigma \frac{\theta(T - T_c)}{\theta(T_d - T)}, \quad (9)$$

where $f_g^{\mathrm{th}}(k, u, T)$ is the thermal distribution of gluons and $\sigma(k, p)$ is the cross section of the gluon dissociation process $B_c^+ + g \to \bar{b} + c$, which is obtained by replacing the heavy quark mass m_Q with twice the reduced mass, 2μ, and substituting the proper binding energy in the cross sections for quarkonia obtained by the Operator Production Expansion (OPE) method [28,44–46]. In the B_c rest frame, these can be expressed as [28,44–46]

$$\sigma_{1S}(\omega) = A_0 \frac{(\omega/\epsilon_{1S} - 1)^{3/2}}{(\omega/\epsilon_{1S})^5},$$

$$\sigma_{1P}(\omega) = 4A_0 \frac{(\omega/\epsilon_{1P} - 1)^{1/2}[9(\omega/\epsilon_{1P})^2 - 20(\omega/\epsilon_{1P}) + 12]}{(\omega/\epsilon_{1P})^7},$$

$$\sigma_{2S}(\omega) = 16A_0 \frac{(\omega/\epsilon_{2S} - 1)^{3/2}(\omega/\epsilon_{2S} - 3)^2}{(\omega/\epsilon_{2S})^7},$$

where ω is the gluon energy, $A_0 = 2^{11}\pi 3^{-3}(2\mu)^{-3/2}\epsilon_{1S}^{-1/2}$, $\mu = (m_b m_c)/(m_b + m_c)$ is the reduced mass, and ϵ_J is the binding energy of the B_c state J. Here we use vacuum values; that is, the binding energy is found as a difference between the sum of the masses of the ground-state B and D mesons and the mass of the corresponding B_c state.[2] To take into account the recoil effect due to the finite mass of B_c, we further replace the binding energy ϵ by the threshold $\omega_0 = [1 + \epsilon/(2m_{B_c})]\epsilon$ as in Ref. [47].

The lifetime of B_c with the gluon dissociation taken into account is shown in Fig. 1. The binding energy of the ground state is much larger than that of the excited states, and therefore the ground state lives longer. Comparing the lifetime of the B_c meson to the radius of the nucleus $R(\mathrm{Pb}) \sim 6$ fm, one finds that the ground state suffers strong suppression at a temperature above 400 MeV and little suppression below 300 MeV, while the suppression to the excited states is already very strong at 200 MeV.

The theta functions in the numerator and denominator in Eq. (9) restrict the above process between the critical temperature T_c of light hadrons and the dissociation temperature T_d of B_c mesons. Below T_c, there are no gluons in the fireball. Above T_d, the bottom and charm quarks are strongly screened, and can no longer form a bound state.

With the heavy quark potential at finite temperature obtained from lattice QCD [48], the dissociation temperature T_d of B_c mesons can be calculated by solving the Schrödinger

---

[1]The value of the ratio $N_{\Upsilon s}/N_{b\bar{b}}^{\mathrm{dir}}$ predicted by the SCM gives the nuclear modification factor of hidden bottom mesons $R_{AA}(\Upsilon s) = 0.5$–1. For the ground state, $R_{AA}(\Upsilon 1S) = 1.2$–2.5 is about 3–6 times larger than the experimental result [43]. This suggests that the hidden bottom mesons might not be completely thermalized.

[2]In reality, the binding energies of the B_c state in a deconfined medium are modified due to the screening effect. However, in a recent study [31], the difference in Upsilon suppression between different binding energies—strong binding as in our calculation and weak binding in which heavy quarks are totally thermalized—is discussed and compared; it is found that the strong-binding calculation better explains the experimental data from RHIC and LHC. Based on these considerations, we take the binding energy in the cross section from the value in vacuum as an approximation.
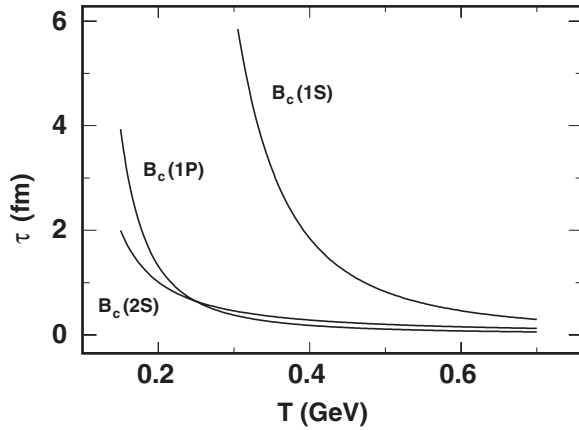
FIG. 1. Lifetime of the B_c meson with the gluon dissociation in QGP taken into account. The velocity of the QGP is taken as $v = 0$ in this figure.

equation as in Ref. [34]. The results are listed in Table I. The dissociation temperature is calculated for two extreme cases. In the first one, the potential is assumed to be equal to the internal energy: $V = U$; in the other one it is equal to the free energy: $V = F$. Note that the difference between them is the entropy term. The former corresponds to a quick adiabatic screening, while in the latter a strong heat exchange between the heavy quark system and the medium is assumed [49]. Since the difference is significant, we will do calculations in both limits in parallel.

Besides the dissociation of B_c, the inverse process $\bar{b} + c \rightarrow B_c + g$ is also considered in Eq. (8), which is described by the regeneration rate

$$
\beta = \frac{1}{2E_{B_c}} \int \frac{d\vec{k}}{(2\pi)^3 2E_g} \frac{d\vec{q}_c}{(2\pi)^3 2E_c} \frac{d\vec{q}_{\bar{b}}}{(2\pi)^3 2E_{\bar{b}}}
$$
$$
\times W(s) f_c f_{\bar{b}} (1 + f_g)(2\pi)^4 \delta^{(4)}(p + k - q_c - q_{\bar{b}}),
$$

where p, k, q_c, and $q_{\bar{b}}$ are the momenta of the B_c, the gluon, the charm quark, and the bottom quark, respectively. $W(s)$ is the transition probability, which is related to that of the dissociation process by detailed balance of the cross section entering in Eq. (9). f_c, $f_{\bar{b}}$, and f_g are distribution functions of c and b quarks and gluons. f_g is taken as a thermal distribution as in the dissociation rate α. For simplicity, we also assume kinetic thermalization for the heavy quarks, and we neglect the Pauli blocking effect. That is, the distribution function $f_Q = \rho_Q f_Q^{\text{th}}$, where the momentum part f_Q^{th} is the normalized Boltzmann distribution, and the density of heavy quarks, ρ_Q,

TABLE I. The dissociation temperature T_d of B_c mesons scaled by the critical temperature T_c. The heavy quark potential V is considered as the internal energy U and the free energy F at finite temperature, respectively. The dash means the meson does not survive above T_c.

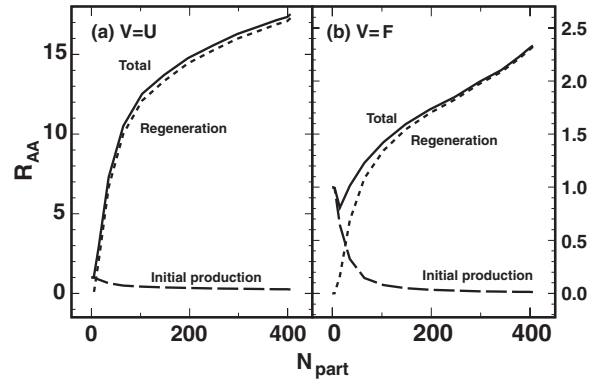| States of B_c	$1S$	$1P$	$2S$
T_d/T_c ($V = U$)	3.27	1.59	1.41
T_d/T_c ($V = F$)	1.51	–	–

FIG. 2. The nuclear modification factor R_{AA} of the B_c meson as a function of the participant nucleon number N_{part} in $V = U$ (a) and $V = F$ (b) limits. The long-dashed, dashed, and solid curves are the initial production, regeneration, and total contributions, respectively.

is determined by the conservation law of heavy quarks:

$$
\partial_t \rho_Q + \nabla \cdot (\rho_Q \vec{v}) = 0, \tag{10}
$$

where $\vec{v}$ is taken as the velocity of the medium. The initial condition is determined by the Glauber model. The regeneration process takes place in the temperature window $T_c < T < T_d$. When $T < T_c$, there are no partons, and therefore the previous process does not exist. When $T > T_d$, the dissociation rate α is infinity, and the B_c gets dissociated as soon as it forms; therefore regeneration does not occur at high temperature, either. Since the dissociation temperature T_d depends strongly on the form of the heavy quark potential, one would expect that the yield of B_c from the regeneration process also depends strongly on the potential.

As a background of the B_c suppression and regeneration, the fireball is described by $(2 + 1)$-dimensional ideal hydrodynamics with the assumption of boost invariance [50]. The equations of state (EoS) are taken as a massive ideal gas of hadrons and partons in confined and deconfined phases, respectively, with a first-order phase transition at $T_c = 165$ MeV as was done for J/ψ and Υ [51]. Multiplying this critical temperature by the ratio T_d/T_c in Table I, one obtains the dissociation temperature T_d, regardless of the absolute value of T_c in the lattice simulation. Otherwise, there could be even more enhancement of B_c. The initial condition is decided by the Glauber model, and the maximum temperature of the fireball in the most central collision is 485 MeV at a thermalization time $\tau_0 = 0.6$ fm.

The centrality dependence of the nuclear modification factor R_{AA} in both limits of $V = U$ and $V = F$ are shown in Fig. 2. If regeneration is ignored, R_{AA} of the initially produced B_c is 0.25 for $V = U$ and 0.01 for $V = F$, which are reasonable when compared to the experimental results for Υ. The R_{AA} values of $\Upsilon(1S + 2S + 3S)$ at RHIC [25] and $\Upsilon(1S)$ at LHC [52] in central collisions are consistent with the simple assumption that the ground state survives while all the excited states melt in the fireball. One can expect that the suppression of B_c is stronger than that of Υ. In the $V = U$ limit, the maximum temperature is still lower than T_d of $B_c(1S)$, and about half of the $B_c(1S)$ mesons survive the gluon dissociation, which is less than that of $\Upsilon(1S)$. In the

$V = F$ limit, the dissociation temperature is much lower, and the ground state suffers even stronger suppression.

In both cases, the contribution from the initial production is two orders of magnitude smaller than that from the regeneration, and thus the population of B_c is dominated by the latter. Note that a larger R_{AA} value for B_c than that for Υ already implies the regeneration mechanism in the fireball. In both limits of our calculations, $R_{AA} = 2.3$–17.5, exceeding unity. The temperature for regeneration in the transport model is higher than that in the SCM; thus one cannot ascertain whether the yield of B_c is smaller in the transport model or in the SCM. Actually, this interval covers the SCM results and also implies strong enhancement. The remarkable difference between the two limits comes from two factors. First, the excited states does not survive at the free energy limit, while they play an important role in the internal energy limit. In the most central collisions, the contribution from the ground state and the exited states are almost the same for $V = U$. Second, the dissociation temperature of the ground state is much higher in the internal energy limit, which allows more regeneration in the hot fireball.

When the regeneration inside the fireball is confirmed, the nuclear modification factor loses its meaning as a survival probability, because the observed B_c mesons are mostly regenerated instead of surviving. Thus R_{AA} depends not only on how the B_c mesons interact with the medium but also on the production cross section of B_c and heavy quarks in hard nucleon-nucleon collisions.

To characterize the producing and/or surviving ability of B_c from the regeneration, we take a ratio of the final yield of B_c mesons over those of the heavy quarks in the unit rapidity region [53] as

$$F \equiv \frac{dN_{B_c}/dy}{(dN_c/dy)(dN_b/dy)}. \tag{11}$$

While R_{AA} compares the final yield of B_c to the initial yield, F compares it to the source of regeneration, and thus it becomes independent of the initial cross sections of B_c and heavy quarks when the regeneration dominates. If there is no nuclear matter effect, the fraction is just

$$f_{\text{Glauber}} = \frac{(d\sigma_{B_c}/dy)\sigma_{\text{in}}}{(d\sigma_c/dy)(d\sigma_b/dy)} \frac{1}{N_{\text{coll}}}, \tag{12}$$

where σ_{B_c}, σ_c, σ_b, and σ_{in} are the cross sections of B_c, c and b quarks, and the inelastic collisions in pp collisions, respectively, and N_{coll} is the number of binary collisions. Under our assumption that the heavy quarks are conserved in the fireball, the denominators in F and F_{Glauber} are the same in Eq. (11), and thus $R_{AA} = F/F_{\text{Glauber}}$. The comparison of F between different beam energies for regeneration-dominant processes makes more sense than R_{AA}, since the initial cross sections in F_{Glauber} is dropped.

The final ratio F at LHC energy is shown in Fig. 3 compared with that of the regenerated B_c at RHIC. The final ratio is larger in the $V = U$ case, as one would expect from the dissociation temperatures. Since the temperature at LHC is much higher, the final ratio is smaller than the corresponding case at RHIC, which is consistent with the picture of color screening. The smaller F in central collisions also implies smaller producing
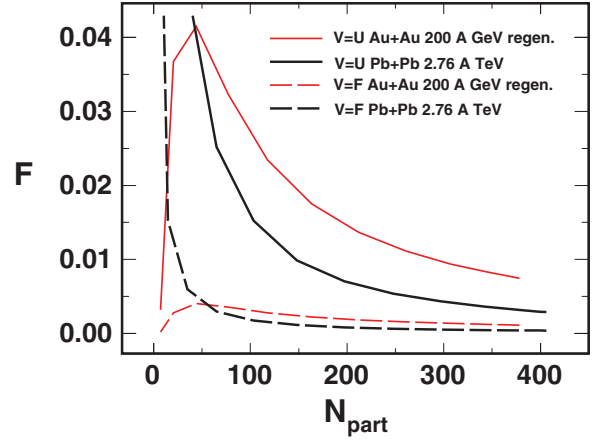

FIG. 3. (Color online) The final ratio F [defined in Eq. (11)] of the B_c meson as a function of the participant nucleon number N_{part}. The long-dashed and solid lines are $V = F$ and $V = U$ limits, respectively. The black thick lines are for LHC energy, while the red thin lines are for the regenerated B_c mesons at RHIC energy as a comparison.

and/or surviving ability of B_c in a hot, large medium, and the growth of R_{AA} with N_{part} is mainly due to the increasing multiplicity of heavy quarks. Meanwhile, the yield is very sensitive to the potential, as we found in R_{AA}, so that F with $V = F$ at RHIC is even smaller than that with $V = U$ at LHC. The difference in behavior between RHIC and LHC at small N_{part} is artificial. That is, because the initial production is not included in calculations for RHIC in Fig. 3, F as a measurement of producing and/or surviving ability only makes sense when the regeneration dominates.

The centrality dependence of F in the SCM is similar, as shown in Fig. 4. In order to apply the SCM to diffferent centralities, we assumed that the volume of the fireball is proportional to the multiplicity of charged particles, which is measured by experiments [54]. The yield of heavy quarks is from a Glauber model as before. The result can be understood
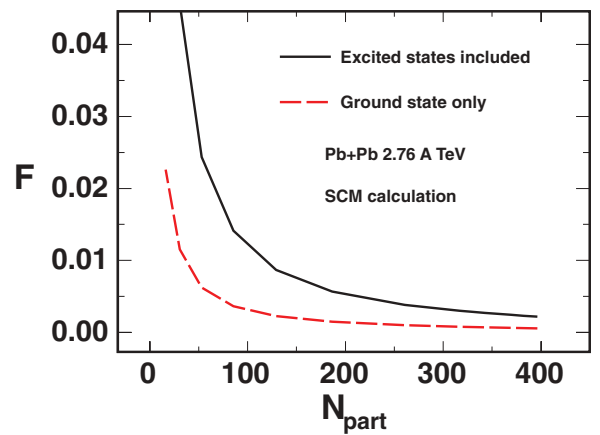

FIG. 4. (Color online) The final ratio F [defined in Eq. (11)] of the B_c meson as a function of the participant nucleon number N_{part} from a SCM calculation. The red long-dashed and black solid lines are calculations of the observed scalar B_c state only and that including excited B_c states (as listed in Table I), respectively.
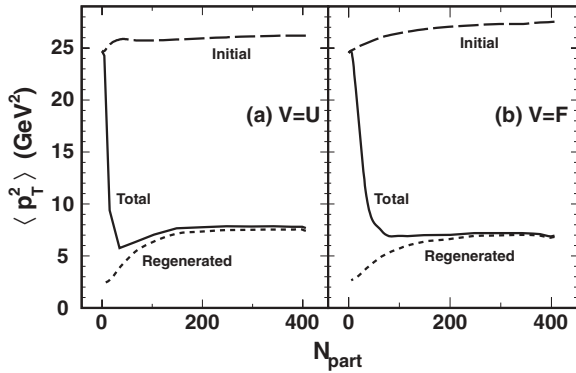
FIG. 5. The average transverse momentum squared, $\langle p_T^2 \rangle$, as a function of the participant nucleon number N_{part} in $V = U$ (a) and $V = F$ (b) limits. The long-dashed, dashed, and solid curves are for initial production, regenerated, and total contributions, respectively.

from (5), which implies in unit rapidity

$$F = \frac{4N_{B_c}^{\text{th}}}{N_{\text{oc}}^{\text{th}} N_{\text{ob}}^{\text{th}}} = \frac{4n_{B_c}^{\text{th}}}{n_{\text{oc}}^{\text{th}} n_{\text{ob}}^{\text{th}}} \frac{1}{V} \propto \frac{1}{V}. \tag{13}$$

The volume increases with N_{part}, and thus the final ratio F decreases, which implies that it is more difficult for a given pair of heavy quarks to meet and combine with each other. When considering FV, there is almost N_p scaling, with a deviation within a few percent in peripheral collisions. The deviation comes from the change in heavy quark density and the event-by-event fluctuation. In the most central collisions, F is between 6×10^{-4} and 2×10^{-3}, which is consistent with the results $[(7 \times 10^{-5})$–$(3 \times 10^{-3})]$ in previous literature [55] at volume $dV/dy = 600$–800 fm and temperature $T = 140$–260 MeV.

One of the advantages of the transport model is that it allows one to investigate the momentum of particles, which is sensitive to the dynamics. The initially produced B_c mesons come from the hard collisions, while the regenerated B_c's are merged from the heavy quarks that are softened by the medium. The typical energy scale of the medium is the temperature, which is much smaller than the typical energies of hard protons in the initial beams. Thus the transverse momentum carried by the initial production is obviously larger than that of the regeneration. When the regeneration becomes important, there is a suppression in $\langle p_T^2 \rangle$, as shown in Fig. 5. The same phenomenon is also expected for J/ψ at LHC [56].

Since thermalized heavy particles from a pointlike thermal source follow a Boltzmann distribution, that is, $dN/d\mathbf{p} \propto \exp(-E_T/T)$ at mid-rapidity, we plot $\ln(dN/d\mathbf{p})$ at mid-rapidity as a function of E_T in Fig. 6 for central Pb + Pb collisions. The plot would be a straight line for a Boltzmann distribution. The softening of the spectrum in Pb + Pb collisions is obvious through the change of the slope. If a superposition of Boltzmann distributions at different temperatures is considered, only a concave curve is expected. However, in the internal energy limit, it is a convex curve at $E_T \sim 9$ GeV. Such a behavior is mainly attributed to the suppression of excited states. If we switch off the gluon dissociation process of the excited states, the obvious bend at low p_T disappears. The
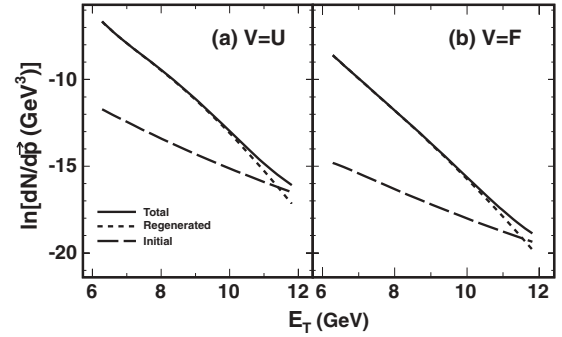

FIG. 6. B_c meson spectrum with respect to the transverse energy $E_T = \sqrt{p_T^2 + m^2}$ in the $V = U$ (a) and $V = F$ (b) limits. The long-dashed, dashed, and solid curves are for initial production, regenerated, and total contributions, respectively.

effective temperature extracted from the spectrum at low p_T according to the Boltzmann distribution is about 550 MeV, which is above the maximum temperature input of the fireball. This mainly results from the radial flow. When the velocity of the fireball is switched off in the regeneration process, the effective temperatures are around 250 MeV, while a calculation of average temperature directly from the transport model is about 230 MeV. Since the radial flow seems large, one would also expect large elliptic flow.

The elliptic flow at $b = 8.4$ fm is shown in Fig. 7. The initially produced B_cs do not thermalize with the medium, and thus they carry a small elliptic flow. The nonzero flow comes from the suppression process similar to that for jets. The regenerated B_cs are born inside the fireball, and thus they inherit a relatively large flow of the medium through heavy quarks. The total flow is dominated by the regeneration at low p_T, and it decreases at high p_T when the initial production becomes important. Since ideal hydrodynamics and the kinetic thermalization of the heavy quarks are assumed, the flow of regenerated B_c at high p_T is not reliable. Note that the assumption of the thermalization of the momenta of bottom quarks is essential in this v_2 calculation. If the bottom quarks are not thermalized, the flow of B_c can be lower than our results in the whole range of p_T, which is to be measured in future experiments.
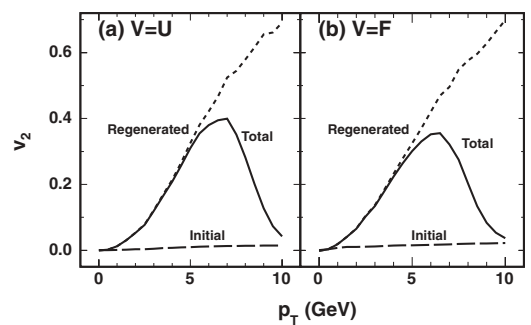

FIG. 7. The elliptic flow v_2 of B_c mesons as a function of the transverse momentum p_T at impact parameter $b = 8.4$ fm in Pb + Pb collisions. Two limits of the heavy quark potential $V = U$ (a) and $V = F$ (b) are shown. The long-dashed, dashed, and solid curves are for initial production, regenerated, and total contributions, respectively.

V. CONCLUSION

Based on the results of both the statistical coalescence model and the transport model, an enhancement of the B_c meson yield in Pb + Pb collisions relative to $p + p$ collisions at $\sqrt{s} = 2.76\,A$ TeV is predicted. If such an enhancement is observed, we can conclude firmly that the regeneration of B_c mesons in the fireball occurs. Thus the J/ψ production in Pb + Pb collisions at LHC is most likely dominated by the regeneration mechanism, and the final fraction F in Eq. (11) can be used to better characterize the hot nuclear matter effect on quarkonia in the fireball. In the opposite case, i.e., if no B_c enhancement is observed, one has to conclude that the observed Υ yield is dominated by the initial production. According to our transport results, the transverse momentum of B_c is suppressed accompanying the enhancement in yield,

since the regenerated B_cs are produced at a lower energy scale than that of the initial production. The p_T spectrum is much softer than that in pp collisions, and in the $V = U$ limit, there is a suppression at low p_T compared to the thermal distribution, which results from the suppression of the excited states. Hopefully, the study of B_c mesons will shed more light on the production and suppression of heavy quarkonia and properties of hot quark matter.

[1] K. Adcox *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **88**, 022301 (2002).

[2] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **98**, 162301 (2007).

[3] T. Matsui and H. Satz, Phys. Lett. B **178**, 416 (1986).

[4] M. Gonin *et al.* (NA50 Collaboration), Nucl. Phys. A **610**, 404c (1996).

[5] C. Spieles, R. Vogt, L. Gerland, S. Bass, M. Bleicher, H. Stocker, and W. Greiner, Phys. Rev. C **60**, 054901 (1999).

[6] J. Geiss, C. Greiner, E. Bratkovskaya, W. Cassing, and U. Mosel, Phys. Lett. B **447**, 31 (1999).

[7] N. Armesto, A. Capella, and E. G. Ferreiro, Phys. Rev. C **59**, 395 (1999).

[8] D. Kahana and S. Kahana, Prog. Part. Nucl. Phys. **42**, 269 (1999)

[9] J.-W. Qiu, J. P. Vary, and X.-F. Zhang, Nucl. Phys. A **698**, 571 (2002).

[10] S. Kabana, New J. Phys. **3**, 16 (2001).

[11] P. Braun-Munzinger and J. Stachel, Phys. Lett. B **490**, 196 (2000).

[12] M. I. Gorenstein, A. P. Kostyuk, H. Stoecker, and W. Greiner, Phys. Lett. B **509**, 277 (2001).

[13] M. I. Gorenstein, A. Kostyuk, H. Stoecker, and W. Greiner, J. Phys. G **27**, L47 (2001).

[14] R. L. Thews, M. Schroedter, and J. Rafelski, Phys. Rev. C **63**, 054905 (2001).

[15] L. Grandchamp and R. Rapp, Phys. Lett. B **523**, 60 (2001).

[16] B. Zhang, C. M. Ko, B.-A. Li, Z.-W. Lin, and S. Pal, Phys. Rev. C **65**, 054909 (2002).

[17] A. Kostyuk, arXiv:nucl-th/0502005 [nucl-th].

[18] L. Yan, P. Zhuang, and N. Xu, Int. J. Mod. Phys. E **16**, 2048 (2007).

[19] X. Zhao and R. Rapp, Phys. Lett. B **664**, 253 (2008).

[20] A. P. Kostyuk, M. I. Gorenstein, H. Stoecker, and W. Greiner, Phys. Rev. C **68**, 041902 (2003).

[21] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **98**, 232301 (2007).

[22] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **101**, 122301 (2008).

[23] B. Abelev *et al.* (STAR Collaboration), Phys. Rev. C **80**, 041902 (2009).

[24] B. Abelev *et al.* (ALICE Collaboration), Phys. Rev. Lett. **109**, 072301 (2012).

[25] R. Reed, J. Phys. G **38**, 124185 (2011).

[26] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Rev. Lett. **107**, 052302 (2011).

[27] F. Abe *et al.* (CDF Collaboration), Phys. Rev. D **58**, 112004 (1998).

[28] M. Schroedter, R. L. Thews, and J. Rafelski, Phys. Rev. C **62**, 024905 (2000).

[29] R. Averbeck *et al.* arXiv:1107.3243v2.

[30] B. Abelev *et al.* (ALICE Collaboration), J. High Energy Physics 01 (2012) 128.

[31] A. Emerick, X. Zhao, and R. Rapp, Eur. Phys. J. A **48**, 72 (2012).

[32] T. Nigmanov, K. Gibson, M. Hartz, and P. Shepard (CDF Collaboration), arXiv:0910.3013 [hep-ex].

[33] Y.-Q. Chen and Y.-P. Kuang, Phys. Rev. D **46**, 1165 (1992); **47**, 350(E) (1993).

[34] H. Satz, J. Phys. G **32**, R25 (2006).

[35] K. Nakamura *et al.* (Particle Data Group), J. Phys. G **37**, 075021 (2010)

[36] T. Sjostrand *et al.*, Comput. Phys. Commun. **135**, 238 (2001).

[37] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. C **84**, 044905 (2011).

[38] R. Vogt, Phys. Rev. C **81**, 044903 (2010).

[39] P. Braun-Munzinger, D. Magestro, K. Redlich, and J. Stachel, Phys. Lett. B **518**, 41 (2001).

[40] J. Uphoff, O. Fochler, Z. Xu, and C. Greiner, Phys. Rev. C **82**, 044906 (2010).

[41] J. Rafelski and M. Danos, Phys. Lett. B **97**, 279 (1980).

[42] A. Andronic, P. Braun-Munzinger, and J. Stachel, Phys. Lett. B **673**, 142 (2009); **678**, E516 (2009).

[43] S. Chatrchyan *et al.* (CMS Collaboration), arXiv:1208.2826 [nucl-ex].

[44] M. E. Peskin, Nucl. Phys. B **156**, 365 (1979).

[45] G. Bhanot and M. E. Peskin, Nucl. Phys. B **156**, 391 (1979).

[46] F. Arleo, P. B. Gossiaux, T. Gousset, and J. Aichelin, Phys. Rev. D **65**, 014005 (2001).

[47] A. Polleri, T. Renk, R. Schneider, and W. Weise, Phys. Rev. C **70**, 044906 (2004).

[48] O. Kaczmarek and F. Zantow, Phys. Rev. D **71**, 114510 (2005).

[49] E. V. Shuryak and I. Zahed, Phys. Rev. D **70**, 054507 (2004).

[50] P. Kolb, P. Huovinen, U. W. Heinz, and H. Heiselberg, Phys. Lett. B **500**, 232 (2001).

[51] X. Zhu, P. Zhuang, and N. Xu, Phys. Lett. B **607**, 107 (2005).
[52] S. Chatrchyan *et al.* (CMS Collaboration), arXiv:1201.5069 [nucl-ex].
[53] Y. Liu, Z. Qu, N. Xu, and P. Zhuang, J. Phys. G **37**, 075110 (2010).

[54] K. Aamodt *et al.* (ALICE Collaboration), Phys. Rev. Lett. **106**, 032301 (2011).
[55] I. Kuznetsova and J. Rafelski, Eur. Phys. J. C **51**, 113 (2007).
[56] Y. Liu, Z. Qu, N. Xu, and P. Zhuang, Phys. Lett. B **678**, 72 (2009).