Bottomonia suppression in relativistic heavy-ion collisions

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Using the two-component model that includes both initial production from nucleon-nucleon hard scattering and regeneration from produced quark-gluon plasma, we study the effect of medium modifications of the binding energies and radii of bottomonia on their production in heavy-ion collisions. We find that the contribution to bottomonia production from regeneration is small and the inclusion of medium effects is helpful for understanding the observed suppression of bottomonia production in experiments carried out at both the BNL Relativistic Heavy Ion Collider and the CERN Large Hadron Collider.

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I. INTRODUCTION

Since the suggestion by Matsui and Satz [1] that suppressed production of J/ψ in relativistic heavy-ion collisions could be a signature for produced quark-gluon plasma (QGP), there have been many experimental [2,3] and theoretical studies [4–11] on this very interesting phenomenon; see, e.g., Refs. [12,13] for a recent review. Although their original idea was that the screening of color charges in the QGP would prohibit charm and anticharm quarks from forming the J/ψ and thus suppress its production, lattice QCD calculations of the J/ψ spectral function showed, on the other hand, that the J/ψ could survive inside QGP up to the so-called dissociation temperature [14–16]. However, the dissociation temperature depends on how the J/ψ spectral function is extracted from lattice data and it is not yet clear if it is far from the critical temperature [17]. Studying J/ψ suppression in relativistic heavy-ion collisions can thus provide not only a signature for the QGP but also a probe of its properties [2,3,18,19]. A quantitative study of J/ψ production in heavy-ion collisions is, however, complicated by their absorption in the initial cold nuclear matter and regeneration in the QGP. Since the Υ is a more strongly bound state of bottom and antibottom quarks than the J/ψ, its production in heavy-ion collisions is expected to be less affected by initial cold nuclear matter effects. Furthermore, the much smaller number of bottom quarks than charm quarks that is produced in heavy-ion collisions makes the contribution of regeneration from the QGP to Υ production also less important. Therefore, studying Υ production in heavy-ion collisions would provide a cleaner probe of the properties of QGP and also the in-medium properties of bottomonia [20,21]. Recently, the nuclear modification factor R_{AA} of the sum of bottomonia $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, defined by their yields relative to those from $p+p$ collisions multiplied by the number of initial binary collisions, in Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV was measured by the STAR Collaboration at the BNL Relativistic Heavy Ion Collider (RHIC) [22], while that of $\Upsilon(1S)$ [23] and the relative suppression of $\Upsilon(2S)$ and $\Upsilon(3S)$ to $\Upsilon(1S)$ [24] in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV were measured by the CMS Collaboration at the CERN Large Hadron Collider (LHC). In both experiments, the measured R_{AA} was seen to decrease with increasing centrality of collisions. In the present study, we use the two-component model [9–11], which was previously used for studying J/ψ production in heavy-ion collisions, to show that these experimental results allow us to obtain useful information on the properties of bottomonia in QGP.

The paper is organized as follows. We first briefly review in Sec. II the two-component model for bottomonia production in heavy-ion collisions. We then compare in Sec. III the calculated nuclear modification factors of bottomonia from the model with experimental data. Finally, a conclusion is given in Sec. IV.

II. TWO-COMPONENT MODEL

In the two-component model, bottomonia are produced from both initial nucleon-nucleon hard scattering and regeneration in produced QGP. The numbers of initially produced bottomonia are proportional to the number of binary collisions among nucleons in the two colliding nuclei. Whether these bottomonia can survive after collisions depends on many effects from both the initial cold nuclear matter and the final hot partonic and hadronic matters. The cold nuclear matter effect includes the nuclear shadowing of bottomonia production from gluon fusions as a result of the modification of the parton distribution in a nucleus from that in a nucleon as well as the absorption of initially produced bottomonia by nucleons in the colliding nuclei. Because of their small sizes, the absorption of bottomonia in cold nuclear matter is expected to be small and is neglected in the present study. Although the recently measured nuclear modification factor of bottomonia in d + Au collisions at RHIC shows a suppression [25], this has been attributed to the energy loss of gluons in nuclear matter before they fuse to produce the bottomonia [26]. For the shadowing effect, we use the EPS09 package [27] and assume that the effect is proportional to the path length

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of a parton in a nucleus [11,28]. Using the experimental data on the average transverse momentum of Υ [23,29], the ratios of gluon distribution function in a nucleus to that in a single nucleon at $x = m_T/\sqrt{s_{NN}}$ are 1.12 and 0.92 for collisions at RHIC and LHC energies, respectively, with the ratio larger than unity called the antishadowing. For thermal dissociation of bottomonia, we include the initial dissociation in the produced QGP of temperatures above their dissociation temperatures as well as the subsequent dissociation by quarks and gluons in the QGP and hadrons in the hadronic matter (HG).

For the local temperature of initially produced matter, we estimate it from the local entropy density $ds/d\eta = C[(1-\alpha)n_{\rm part}/2 + \alpha\, n_{\rm coll}]$ [30] in terms of the number densities of participants and binary collisions $n_{\rm part(coll)} = \Delta N_{\rm part(coll)}/(\tau_0 \Delta x \Delta y)$ with $\Delta N_{\rm part(coll)}$ being the number of participants (binary collisions) in the volume $\tau_0 \Delta x \Delta y$ of the transverse area $\Delta x \Delta y$ and τ_0 being the initial thermalization time. For the value of τ_0, we use 0.9 and 1.05 fm/c from the viscous hydrodynamics [31] for Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV at RHIC and Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at LHC, respectively. The parameters C and α are, respectively, 0.11 and 18.7 for RHIC and 0.15 and 27.0 for LHC from fitting measured charged particle multiplicities [11]. With the quasiparticle model of three flavors for the equation of state of QGP and the resonance gas model for that of HG [9,32], the initial maximum temperatures are 324 and 390 MeV and mean temperatures are 269 and 311 MeV for central collisions at RHIC and LHC, respectively. After initial thermalization, a schematic viscous hydrodynamics [11,33] is then used to describe the expansion of the hot dense matter with the specific shear viscosity η/s taken to be 0.16 and 0.2 for the QGP at RHIC and LHC [31], respectively, and 0.8 for the HG [34].

The thermal dissociation by partons in the QGP or hadrons in the HG of bottomonia that have survived from initial dissociations is described by the rate equation for the number N_i of bottomonia of type i

$$\frac{dN_i}{d\tau} = -\Gamma_i\left(N_i - N_i^{\rm eq}\right), \qquad (1)$$

where τ is the longitudinal proper time and Γ_i is the thermal decay width of bottomonia. The number of equilibrated bottomonia of type i is given by $N_i^{\rm eq} = \gamma^2 R n_i f V \theta(T_i - T)$, where n_i is its number density in the grandcanonical ensemble; f is the fraction of QGP in the mixed phase and is 1 in the QGP; and $\theta(T_i - T)$ is the step function with T_i being the dissociation temperature of bottomonia of type i. The chemical and kinetic off-equilibrium of bottom quarks in the QGP is included through their fugacity γ and relaxation factor R, respectively, with the former obtained from the conservation of bottom flavor [11] and the latter defined as $R(\tau) = 1 - \exp[-(\tau - \tau_0)/\tau_{\rm eq}]$ where $\tau_{\rm eq}$ is the relaxation time of bottom quarks. With $\tau_{\rm eq} \sim m_Q$ [35] and $\tau_{\rm eq} = 4$ fm/c for charm quarks [11], the relaxation time for bottom quarks is about 14 fm/c. We note that the second term on the right-hand side of Eq. (1) takes into account the regeneration of bottomonia from bottom and antibottom quarks in the QGP.

For the thermal decay widths of bottomonia in QGP, we calculate them up to the next-to-leading order (NLO) in perturbative QCD (pQCD) [36]. 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. The squared invariant amplitudes for these processes are the same as those given in Ref. [11] for charmonia except for the heavy quark mass. In terms of the resulting bottomonium dissociation cross section $\sigma^{\rm diss}$, the thermal decay width of a bottomonium is given by

$$\Gamma(T) = \sum_i \int \frac{d^3k}{(2\pi)^3} v_{\rm rel}(k) n_i(k,T) \sigma_i^{\rm diss}(k,T), \qquad (2)$$

where i denotes the quarks and gluons in the QGP; n_i is the number density of parton species i in the grand-canonical ensemble; and $v_{\rm rel}$ is the relative velocity between the scattering bottomonium and parton. The thermal width in the mixed phase is taken to be a linear combination of those in the QGP and the HG, i.e., $\Gamma(T_c) = f\Gamma^{\rm QGP}(T_c) + (1-f)\Gamma^{\rm HG}(T_c)$, where f is the fraction of QGP in the mixed phase. For the dissociation cross section of $\Upsilon(1S)$ in the hadron gas (HG), we use the factorization formula [36] and assume that those of excited bottomonia are proportional to their squared radii. In the upper panel of Fig. 1, we show the thermal decay widths of bottomonia in QGP calculated with their masses and radii in free space, i.e., 9.5 GeV and 0.23 fm, 9.97 GeV and 0.59 fm, 10.19 GeV and 0.95 fm, 9.9 GeV and 0.46 fm, and 10.25 GeV and 0.82 fm for $\Upsilon(1S)$, $\Upsilon(2S)$, $\Upsilon(3S)$, $\chi_b(1P)$, and $\chi_b(2P)$, respectively, obtained from the Cornell potential with the vacuum screening mass $\mu = 0.18$ GeV [37]. These thermal decay widths of bottomonia are appreciable and increase with increasing temperature. We note that the thermal decay widths of bottomonia in HG are significantly smaller than those in QGP.

To include medium effects on the properties of bottomonia in QGP, we consider the modification of the potential energy between bottom and antibottom quarks due to the Debye

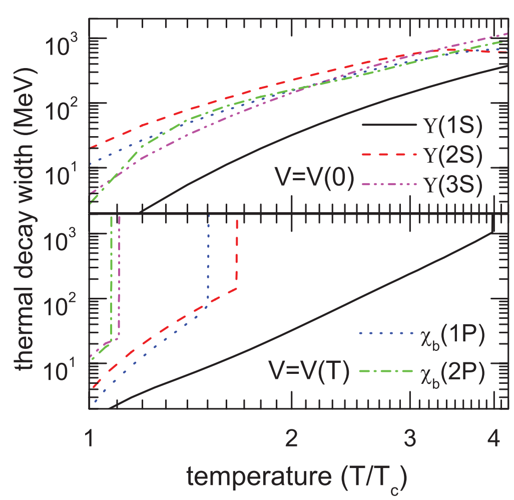

FIG. 1. (Color online) Thermal decay widths of bottomonia as functions of the temperature of QGP without (upper panel) and with (lower panel) medium effects.

screening of color charges. This is achieved by using the screened Cornell potential [37]:

$$V(r,T) = \frac{\sigma}{\mu(T)}[1 - e^{-\mu(T)r}] - \frac{\alpha}{r}e^{-\mu(T)r}, \qquad (3)$$

with $\sigma = 0.192$ GeV^2 and $\alpha = 0.471$. The screening mass $\mu(T)$ depends on temperature, and we use the one given in pQCD, i.e., $\mu(T) = \sqrt{N_c/3 + N_f/6}gT$, where N_c and N_f are numbers of colors and light quark flavors, respectively. We note that compared to results from the lattice QCD calculations for a heavy quark and antiquark pair in the QGP [38], this potential is closer to their internal energy at high temperature but to their free energy at low temperature. The screened Cornell potential thus interpolates smoothly the expected temperature-dependent potential between heavy quark and antiquark [39]. The binding energies and radii of bottomonia in the QGP can be obtained by solving the resulting Schrödinger equation for the bottom and antibottom quark pair. Taking their masses to be $m_b = 4.746$ GeV [37] and the QCD coupling constant $g = 1.87$, as in our previous study of J/ψ suppression in heavy-ion collisions [11], the results are shown in the upper and lower panels of Fig. 2. It is seen that the dissociation temperatures of $\Upsilon(1S)$, $\Upsilon(2S)$, $\Upsilon(3S)$, $\chi_b(1P)$, and $\chi_b(2P)$ are 4.0, 1.67, 1.12, 1.51, and 1.09 T_c, respectively, and their radii increase with increasing temperature. These results are similar to those obtained from the lattice nonrelativistic QCD calculations [40,41]. The thermal decay widths of bottomonia obtained with their in-medium binding energies and radii are shown in the lower panel of Fig. 1, and they increase with temperature and diverge at their dissociation temperatures. Compared to those obtained with the binding energies and radii in free space, medium effects enhance the widths significantly. We note that the inclusion of thermal decay widths of heavy quarkonia effectively takes into account the imaginary part of the potential energy between heavy quark and antiquark at finite temperature [42–44].

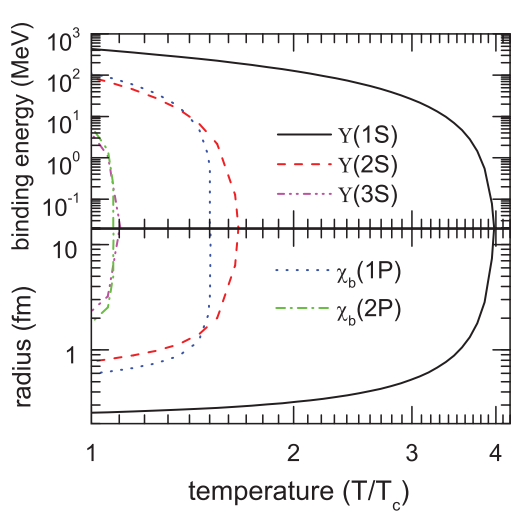

FIG. 2. (Color online) Binding energies (upper panel) and radii (lower panel) of bottomonia as functions of the temperature of QGP.

III. RESULTS

To calculate the nuclear modification factor of bottomonia in heavy-ion collisions requires information on their numbers produced in $p + p$ collisions at the same energy. Since this information is not available at LHC, we use in the present study those of $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ from the experimental data in $p+\bar{p}$ collisions at $\sqrt{s_{NN}} = 1.8$ TeV measured by the CDF Collaboration at the Fermilab [45], and those of $\chi_b(1P)$ and $\chi_b(2P)$ from their contributions to $\Upsilon(1S)$ [46] based on the branching ratios of about 0.24 and 0.13, respectively. For $p + p$ collisions at RHIC, the numbers of bottomonia were not individually measured, so we use $\sum_{n=1\sim3} B(nS) \times d\sigma/dy|_{y=0}(nS) = 114$ pb [25], where $B(nS)$ and $d\sigma/dy|_{y=0}(nS)$ are, respectively, the branching ratio and differential cross section in rapidity for $\Upsilon(nS)$, and assume their relative abundances are the same as those at the LHC. For the initial number of bottom quark pairs in determining the fugacity of bottom flavor, it is obtained from $d\sigma^{pp}_{b\bar{b}}/dy|_{y=0} = 1.34$ μb for RHIC [47] and $\sigma^{pp}_{b\bar{b}} = 17.6$ μb in the rapidity range $|y| < 0.6$ from the CDF Collaboration [48] for LHC. Also, we need the contribution of excited bottomonia to $\Upsilon(1S)$ in $p + p$ collisions at the same energy. Since this has not been measured, we use the information obtained from $p + p$ collisions at $\sqrt{s_{NN}} = 1.97$ TeV by the CDF Collaboration at the Fermilab [46], because they are known to be essentially independent of the collision energy [45]; i.e., the contributions from $\chi_b(1P)$, $\chi_b(2P)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ to $\Upsilon(1s)$ are taken to be 27.1, 10.5, 10.7, and 0.8%, respectively. The resulting nuclear modification factors (R_{AA}) of bottomonia obtained without the cold nuclear effect in Au + Au

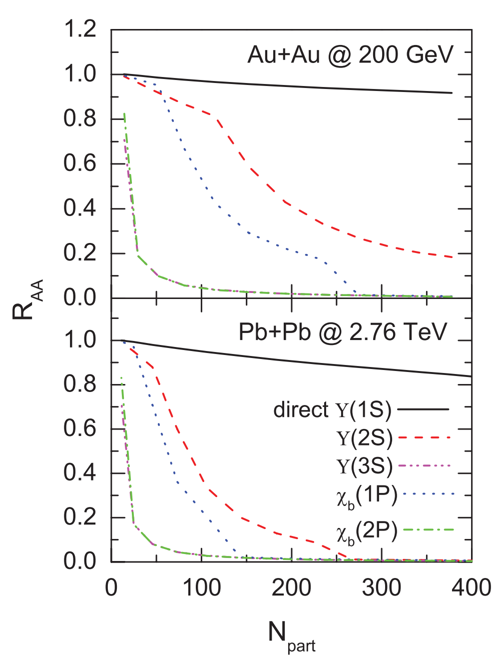

FIG. 3. (Color online) Nuclear modification factor R_{AA} of bottomonia in Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV at RHIC (upper panel) and in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at LHC (lower panel) without including the cold nuclear matter effect.

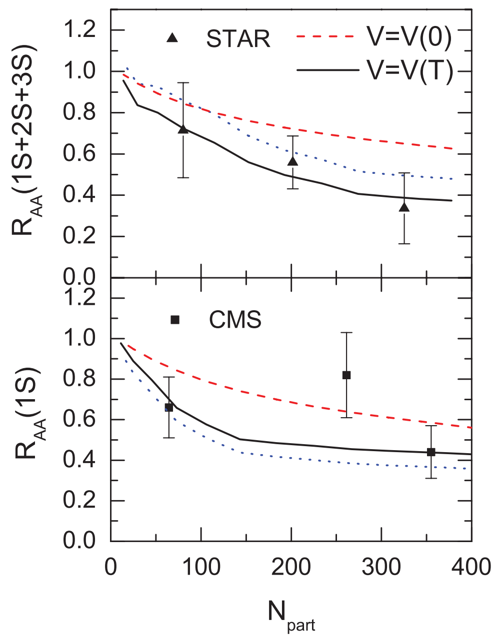

FIG. 4. (Color online) Nuclear modification factor R_{AA} of the sum of $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ in Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV at RHIC (upper panel) and that of $\Upsilon(1S)$ in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at LHC (lower panel) as functions of the participant number. Solid and dashed lines are, respectively, results with and without medium effects on bottomonia. Dotted lines are results including also the shadowing effect. Experimental data are from Refs. [22,23].

collisions at $\sqrt{s_{NN}} = 200$ GeV at RHIC and in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at LHC are shown in Fig. 3. It is seen that the R_{AA} of directly produced $\Upsilon(1S)$ is close to unity even in central collisions, while those of excited bottomonia are small.

In the upper panel of Fig. 4, the calculated R_{AA} of the sum of $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$ in Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV at RHIC as a function of the participant number is shown and compared with the experimental data from the STAR Collaboration [22]. The solid and dotted lines are results obtained without and with the shadowing effect in cold nuclear matter. It is seen that because of the large experimental errors, both can describe the data from RHIC for all centralities.

For Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV at LHC, only the R_{AA} of $\Upsilon(1S)$ has been measured by the CMS Collaboration [23]. Our results for the R_{AA} of $\Upsilon(1S)$ without the cold nuclear matter effect is shown by the solid line in the lower panel of Fig. 4. Compared with that measured by the CMS Collaboration [23], these results agree with the data for both peripheral (20–100%) and most central (0–10%) collisions. For midcentral (10–20%) collisions, our model significantly underestimates the measured R_{AA}. Similar results for LHC were also obtained in Ref. [49] based on the bottom and antibottom quark potential that was taken to be their internal energy from the lattice QCD. Also shown in the lower panel of Fig. 4 are results obtained by including the shadowing effect (dotted line), which are only slightly smaller than those obtained without the shadowing effect. We note that most of the suppression of $\Upsilon(1S)$ comes from those of its excited states, as seen from the results shown in Fig. 3. Also, the contribution from the regeneration to the R_{AA} of bottomonia is less than 1% at both RHIC and LHC and for all centralities as a result of the small number of bottom quarks and the much longer bottom quark relaxation time than the lifetime of produced QGP. For comparison, results for the R_{AA} of bottomonia without medium effects are shown by dashed lines in Fig. 4, and they are larger than those with medium effects as expected. In this case, the calculated R_{AA} of $\Upsilon(1S)$ is large compared to the experimental data from RHIC, particularly for more central collisions. This is also the case for heavy-ion collisions at the LHC except for midcentral collisions where the result obtained without medium effects can better describe the experimental data. It is not clear if this indicates a change of the bottomonia suppression mechanism in midcentral collisions. Improved experimental data are essential for resolving this puzzle.

IV. CONCLUSION

In conclusion, using the two-component model that includes both initial production from nucleon-nucleon hard scattering and regeneration from produced quark-gluon plasma, we have studied bottomonia production in heavy-ion collisions at RHIC and LHC by including the medium effects on the thermal properties of bottomonia and their dissociation cross sections. With the expansion dynamics of produced hot dense matter described by a schematic viscous hydrodynamics and including the thermal dissociation of bottomonia as well as the regeneration of bottomonia by using a rate equation, our model describes successfully the experimental data from RHIC and reasonably those from LHC on bottomonia suppression. Our results indicate that the contribution of regenerated bottomonia is small. We have also studied the cold nuclear matter effect due to the shadowing at LHC or antishadowing at RHIC of the parton distribution function in the nucleus. This was found to increase the R_{AA} of bottomonia at RHIC but decrease it at LHC. Our results with and without the cold nuclear matter effect are, however, both consistent with the experimental data because of their large errors. Furthermore, our study shows that the inclusion of medium effects on bottomonia is essential for describing the experimental observations at RHIC as well as at LHC except for midcentral collisions where results without medium effects can better describe the data. More accurate data from future experiments are needed to obtain more definitive information on the properties of bottomonia in the QGP.

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[1] T. Matsui and H. Satz, Phys. Lett. B **178**, 416 (1986).
[2] B. Alessandro *et al.* (NA50 Collaboration), Eur. Phys. J. C **39**, 335 (2005).
[3] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **98**, 232301 (2007).
[4] R. Vogt, Phys. Rep. **310**, 197 (1999).
[5] B. Zhang, C. M. Ko, B. A. Li, Z. W. Lin, and B. H. Sa, Phys. Rev. C **62**, 054905 (2000).
[6] B. Zhang, C. M. Ko, B. A. Li, Z. W. Lin, and S. Pal, Phys. Rev. C **65**, 054909 (2002).
[7] X. Zhao and R. Rapp, Phys. Lett. B **664**, 253 (2008).
[8] L. Yan, P. Zhuang, and N. Xu, Phys. Rev. Lett. **97**, 232301 (2006).
[9] T. Song, W. Park, and S. H. Lee, Phys. Rev. C **81**, 034914 (2010).
[10] T. Song, C. M. Ko, S. H. Lee, and J. Xu, Phys. Rev. C **83**, 014914 (2011).
[11] T. Song, K. C. Han, and C. M. Ko, Phys. Rev. C **84**, 034907 (2011).
[12] R. Rapp, D. Blaschke, and P. Crochet, Prog. Part. Nucl. Phys. **65**, 209 (2010).
[13] A. Andronic, P. Braun-Munzinger, K. Redlich, and J. Stachel, Nucl. Phys. A **789**, 334 (2007).
[14] M. Asakawa and T. Hatsuda, Phys. Rev. Lett. **92**, 012001 (2004).
[15] S. Datta, F. Karsch, P. Petreczky, and I. Wetzorke, Phys. Rev. D **69**, 094507 (2004).
[16] G. Aarts, C. Allton, M. B. Oktay, M. Peardon, and J.-I. Skullerud, Phys. Rev. D **76**, 094513 (2007).
[17] A. Mocsy and P. Petreczky, Phys. Rev. D **77**, 014501 (2008).
[18] G. Aad *et al.* (Atlas Collaboration), Phys. Lett. B **697**, 294 (2011).
[19] T. Dahms (f. t. C. Collaboration), J. Phys. G **38**, 124105 (2011).
[20] L. Grandchamp, S. Lumpkins, D. Sun, H. van Hees, and R. Rapp, Phys. Rev. C **73**, 064906 (2006).
[21] Y. Liu, B. Chen, N. Xu, and P. Zhuang, Phys. Lett. B **697**, 32 (2011).
[22] R. Reed (STAR Collaboration), J. Phys. G **38**, 124185 (2011).
[23] CMS Collaboration, CMS-PAS-HIN-10-006, 2011 (unpublished).
[24] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Rev. Lett. **107**, 052302 (2011).
[25] R. Reed (STAR Collaboration), J. Phys. Conf. Ser. **270**, 012026 (2011).
[26] E. G. Ferreiro, F. Fleuret, J. P. Lansberg, N. Matagne, and A. Rakotozafindrabe, arXiv:1110.5047.
[27] K. J. Eskola, H. Paukkunen, and C. A. Salgado, J. High Energy Phys. 04 (2009) 065.
[28] R. Vogt, Phys. Rev. C **71**, 054902 (2005).
[29] Matthew Cervantes (private communication).
[30] D. Kharzeev and M. Nardi, Phys. Lett. B **507**, 121 (2001)
[31] H. Song, S. A. Bass, and U. W. Heinz, Phys. Rev. C **83**, 054912 (2011).
[32] P. Levai and U. W. Heinz, Phys. Rev. C **57**, 1879 (1998).
[33] T. Song, K. C. Han, and C. M. Ko, Phys. Rev. C **83**, 024904 (2011).
[34] N. Demir and S. A. Bass, Phys. Rev. Lett. **102**, 172302 (2009)..
[35] G. D. Moore and D. Teaney, Phys. Rev. C **71**, 064904 (2005).
[36] T. Song and S. H. Lee, Phys. Rev. D **72**, 034002 (2005).
[37] F. Karsch, M. T. Mehr, and H. Satz, Z. Phys. C **37**, 617 (1988).
[38] O. Kaczmarek, Eur. Phys. J. C **61**, 811 (2009).
[39] C.-Y. Wong, Phys. Rev. C **72**, 034906 (2005).
[40] G. Aarts, S. Kim, M. P. Lombardo, M. B. Oktay, S. M. Ryan, D. K. Sinclair, and J.-I. Skullerud, Phys. Rev. Lett. **106**, 061602 (2011).
[41] G. Aarts, C. Allton, S. Kim, M. P. Lombardo, M. B. Oktay, S. M. Ryan, D. K. Sinclair, and J.-I. Skullerud, JHEP **1111**, 103 (2011).
[42] M. Laine, O. Philipsen, P. Romatschke, and M. Tassler, J. High Energy Phys. 03 (2007) 054.
[43] A. Beraudo, J.-P. Blaizot, C. Ratti, Nucl. Phys. A **806**, 312 (2008).
[44] N. Brambilla, J. Ghiglieri, A. Vairo, and P. Petreczky, Phys. Rev. D **78**, 014017 (2008).
[45] F. Abe *et al.* (CDF Collaboration), Phys. Rev. Lett. **75**, 4358 (1995).
[46] A. A. Affolder *et al.* (CDF Collaboration), Phys. Rev. Lett. **84**, 2094 (2000).
[47] Y. Morino (PHENIX Collaboration), J. Phys. G **35**, 104116 (2008).
[48] D. Acosta *et al.* (CDF Collaboration), Phys. Rev. D **71**, 032001 (2005).
[49] M. Strickland, Phys. Rev. Lett. **107**, 132301 (2011).