Open charm and dileptons from relativistic heavy-ion collisions

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Dileptons are considered as one of the cleanest signals of the quark-gluon plasma (QGP); however, the QGP radiation is masked by many background sources from either hadronic decays or semileptonic decays from correlated charm pairs. In this study, we investigate the relative contribution of these channels in heavy-ion collisions from $\sqrt{s_{NN}} = 8$ GeV to 5 TeV with a focus on the competition between the thermal QGP radiation and the semileptonic decays from correlated D-meson pairs. As a tool, we employ the parton-hadron-string dynamics (PHSD) transport approach to study dilepton spectra in Pb + Pb (Au + Au) collisions in a wide energy range, incorporating for the first time a fully microscopic treatment of the charm dynamics and their semileptonic decays. We find that the dileptons from correlated D-meson decays dominate the thermal radiation from the QGP in central Pb + Pb collisions at the intermediate masses (1.2 GeV < M < 3 GeV) for $\sqrt{s_{NN}} > 40$ GeV, while for $\sqrt{s_{NN}} = 8$ to 20 GeV the contribution from D, \bar{D} decays to the intermediate mass dilepton spectra is subleading such that one should observe a rather clear signal from the QGP radiation. We furthermore study the p_T spectra and the $R_{AA}(p_T)$ of single electrons at different energies as well as the excitation function of the inverse slope of the m_T spectra for intermediate-mass dileptons from the QGP and from charm decays. We find moderate but characteristic changes in the inverse slope parameter for $\sqrt{s_{NN}} > 20$ GeV which can be observed experimentally in high statistics data. Additionally, we provide detailed predictions for dilepton spectra from Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV.

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I. INTRODUCTION

Relativistic heavy-ion collisions are well suited to generate hot and dense matter in the laboratory, although the matter is produced within small space-time regimes. Whereas in low-energy collisions one produces dense nuclear matter with moderate temperature T and large baryon chemical potential μ_B , ultrarelativistic collisions at Relativistic Heavy Ion Collider (RHIC) or Large Hadron Collider (LHC) energies produce extremely hot matter at small baryon chemical potential. In order to explore the phase diagram of strongly interacting matter as a function of T and μ_B , both types of collisions are mandatory. According to lattice calculations of quantum chromodynamics (IQCD) [1–3], the phase transition from hadronic to partonic degrees of freedom (at vanishing baryon chemical potential $\mu_B = 0$ is a crossover. This phase transition is expected to turn into a first-order transition at a critical point (T_r, μ_r) in the phase diagram with increasing baryon chemical potential μ_B . Since this critical point cannot be determined theoretically in a reliable way, the beam energy scan (BES) program at RHIC aims to find the critical point and the phase boundary by gradually decreasing the collision energy [4,5]. Furthermore, new facilities such as the Facility for Antiproton and Ion Research (FAIR) and Nuclotron-Based Ion Collider Facility (NICA) are under construction to explore in particular the intermediate energy range where one might study also the competition between chiral symmetry restoration and deconfinement as suggested in Refs. [6,7].

Since the partonic phase in relativistic heavy-ion collisions appears only for a couple of fm/c, it is quite a challenge for experiments to investigate its properties. The heavy-flavor mesons are considered to be promising probes in this search since the production of heavy flavor requires a large energymomentum transfer. Thus, it takes place early in the heavy-ion collisions, and-due to the large energy-momentum transfershould be described by perturbative quantum chromodynamics (pOCD). The produced heavy flavor then interacts with the hot dense matter (of partonic or hadronic nature) by exchanging energy and momentum. As a result, the ratio of the measured number of heavy flavors in heavy-ion collisions to the expected number in the absence of nuclear or partonic matter is suppressed at high transverse momentum, and the elliptic flow of heavy flavor is generated by the interactions in noncentral heavy-ion collisions. The experimental data at RHIC and LHC show that the suppression of heavy-flavor hadrons at high transverse momentum and its elliptic flow v_2 are comparable to those of light hadrons [8,9]. This is a puzzle for heavy-flavor production and dynamics in relativistic heavy-ion collisions, as pointed out by many groups [10-23], and is a subject of intense studies both theoretically and experimentally. For recent reviews, we refer the reader to Refs. [24,25].

Furthermore, the electromagnetic emissivity of strongly interacting matter is a subject of long-standing interest [26–29]

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and is explored also in relativistic nucleus-nucleus collisions, where the photons (and dileptons) measured experimentally provide a time-integrated picture of the collision dynamics. Dileptons are of particular interest since their invariant mass provides an additional scale compared to photons and allows us to partly separate the production channels from the early (possibly partonic) phase with those from the late hadronic phase. After decades of experimental and theoretical studies, it has become clear that dileptons with invariant masses below about 1.2 GeV preferentially stem from hadronic decays, providing some glimpse of the modification of hadron properties in the dense and hot hadronic medium (cf. Refs. [29,30] and references therein) while the intermediate mass regime 1.2 GeV < M < 3 GeV should provide information about thermal dileptons from the QGP $(q + \bar{q} \rightarrow e^+e^-, q + \bar{q} \rightarrow e^+e^-)$ $g + \gamma^*, g + q(\bar{q}) \rightarrow, q(\bar{q}) + e^+e^-)$ as well as the amount of correlated open charm (semileptonic) decays from early production of $c\bar{c}$ pairs. Whereas at RHIC and LHC energies the background from $D\bar{D}$ pairs overshines the contribution from the QGP in the intermediate-mass regime [30], one might expect to find some window in bombarding energy where the partonic sources dominate since the charm production drops rapidly with decreasing bombarding energy. In this work, we intend to quantify this expectation and to identify optimal systems for future measurements at FAIR/NICA or at the RHIC Beam-Energy-Scan (BES) as well as those at the Super Proton Synchrotron (SPS) by the NA61 Collaboration.

We recall that previously we have studied the contribution of semileptonic decays from D mesons to the dilepton spectra at RHIC ($\sqrt{s_{NN}} = 200 \text{ GeV}$) and LHC ($\sqrt{s_{NN}} = 2.76 \text{ TeV}$) energies based on an extended statistical hadronization model [31,32]. The charm production in AA collisions was accounted for by scaling the contribution from p + p collisions with the number of binary NN collisions. However, in these studies only the semileptonic decays of correlated (and unscattered) $D\bar{D}$ pairs were considered whereas the contribution from rescattered D and \overline{D} mesons had been neglected. Also, only hadronic rescattering has been incorporated for the decorrelation of the produced DD pair. Since these assumptions are too crude to correctly reflect the actual experimental measurements with their detailed acceptance cuts, a fully microscopic reanalysis of the charm dynamics and charm pair angular correlation is mandatory.

We here employ the microscopic parton-hadron-string dynamics (PHSD) approach, which differs from the conventional Boltzmann-type models in the aspect [33] that the degrees of freedom for the QGP phase are off-shell massive strongly interacting quasiparticles that generate their own mean-field potential. The masses of the dynamical quarks and gluons in the QGP are distributed according to spectral functions whose pole positions and widths, respectively, are defined by the real and imaginary parts of their self-energies [30]. The partonic propagators and self-energies, furthermore, are defined in the dynamical quasiparticle model (DQPM) in which the strong coupling and the self-energies are fitted to lattice QCD results [34].

We recall that the PHSD approach has successfully described numerous experimental data in relativistic heavy-ion collisions from the Alternating Gradient Synchrotron (AGS), SPS, RHIC, and LHC energies [30,33,35–37]. More recently, the charm production and propagation has been explicitly implemented in the PHSD and detailed studies on the charm dynamics and hadronization-fragmention have been performed at top RHIC and LHC energies in comparison to the available data [38–40]. In the PHSD approach, the initial charm and anticharm guarks are produced by using the PYTHIA event generator [41], which is tuned to the transverse momentum and rapidity distributions of charm and anticharm quarks from the fixed-order next-to-leading logarithm (FONLL) calculations [42]. The produced charm and anticharm quarks interact in the QGP with off-shell partons and are hadronized into D mesons close to the critical energy density ($\sim 0.5 \,\text{GeV}/\text{fm}^3$) for the crossover transition either through fragmentation or coalescence. We stress that the coalescence is a genuine feature of heavy-ion collisions and does not show up in p + p interactions. The hadronized D mesons then interact with light hadrons in the hadronic phase until freeze-out and final semileptonic decay. We have found that the PHSD approach, which has been applied for charm production in Au + Au collisions at $\sqrt{s_{NN}} = 200 \text{ GeV} [38]$ and in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [39], describes the R_{AA} of D mesons in reasonable agreement with the experimental data from the STAR Collaboration [43,44] and from the ALICE Collaboration [45,46] when including the initial shadowing effect in the latter case. In this work, we furthermore apply the PHSD approach to charm and dilepton production in relativistic heavy-ion collisions from $\sqrt{s_{NN}} = 8$ GeV to 2.76 TeV, analyze the angular correlation between the charm quarks or D mesons, respectively, and evaluate the contribution to the dilepton spectra from their semileptonic decays. Furthermore, we will give predictions for dilepton mass spectra from Pb + Pb collisions at the top LHC energy of $\sqrt{s_{NN}} = 5.02$ TeV for low and intermediate invariant masses.

This paper is organized as follows: The production of heavy mesons in p + p collisions is described in Sec. II and $c\bar{c}$ pair multiplicities in central Pb+Pb collisions are evaluated within PHSD as a function of invariant energy. We then present the heavy quark interactions in the QGP, their hadronization and hadronic interactions, respectively, in Sec. III as well as the semileptonic decays of the charm hadrons. Section IV is devoted to the description of the dilepton sources incorporated in the actual PHSD calculations while in Sec. V we calculate the R_{AA} of single electrons from open charm mesons at midrapidity as a function of transverse momentum and the modification of the $c\bar{c}$ correlation angle due to the partonic and hadronic interactions in central Pb+Pb collisions from $\sqrt{s_{NN}} = 8$ to 200 GeV. We continue with excitation functions for dilepton spectra in these collisions and investigate separately the contributions from hadronic and partonic sources as well as semileptonic decays from open charm. In Sec. VI, we will compare the PHSD calculations for dilepton spectra with experimental data from $\sqrt{s_{NN}} =$ 19.6 GeV to 2.76 TeV and present predictions for dilepton mass spectra from Pb+Pb collisions at the top LHC energy of $\sqrt{s_{NN}} = 5.02$ TeV. Section VII closes our study with a summary, while Appendixes A and B include the details of the partonic production channels for lepton pairs as well as an examination of the uncertainties in the charm cross



FIG. 1. (a) The $c\bar{c}$ pair cross section in p + p reactions as a function of the invariant energy $\sqrt{s_{NN}}$ as implemented in PHSD. The symbols denote experimental data from Refs. [47–49]. (b) The number of primary $c\bar{c}$ pairs in Pb+Pb collisions at b = 2 fm as a function of $\sqrt{s_{NN}}$. The shaded area in panel (b) shows the uncertainty in the number of $c\bar{c}$ pairs due to the uncertainty in the charm production cross section in p + p collisions.

section and the effects of experimental cuts on the dilepton spectra.

II. CHARM PAIRS FROM p + p COLLISIONS

As pointed out in the introduction, the charm quark ($c\bar{c}$) pairs are produced through initial hard nucleon-nucleon scattering in relativistic heavy-ion collisions. We employ the PYTHIA event generator to produce the heavy-quark pairs and modify their transverse momentum and rapidity such that they are similar to those from the FONLL calculations at RHIC and LHC energies (cf. Ref. [40]). At SPS and lower energies, we do not employ any modification of the PYTHIA results. Figure 1(a) shows the charm production cross section for p + p collisions (as implemented in PHSD) as a function of the invariant energy $\sqrt{s_{NN}}$ which is fitted to a wide range of experimental data. We can see a rather fast drop of the $c\bar{c}$ cross section with decreasing energy, especially close to the threshold energy for charm-pair production. Note, however, that the data show an uncertainty of about a factor of two, which implies a corresponding uncertainty in the following PHSD calculations.

A. Multiplicities for $c\bar{c}$ pairs in central Pb + Pb reactions

We recall that in heavy-ion reactions the number of $c\bar{c}$ pairs produced is approximately given by the number of binary nucleon-nucleon collisions $N_{bin}(b)$ (at given impact parameter b) times the probability to produce a $c\bar{c}$ pair in an inelastic nucleon-nucleon collision at given $\sqrt{s_{NN}}$, which is the ratio of the $c\bar{c}$ cross section to the inelastic N + N cross section. The scaling of the $c\bar{c}$ multiplicity with the number of binary N + N collisions is rather well reproduced in actual PHSD calculations where additionally the smearing of $\sqrt{s_{NN}}$ by Fermi motion is taken into account as well as fluctuations in the number of binary nucleon-nucleon collisions $N_{bin}(b)$ on an event-by-event basis. The corresponding PHSD results for Pb+Pb collisions at b=2 fm are displayed in Fig. 1(b) as a function of $\sqrt{s_{NN}}$ and demonstrate that the average $c\bar{c}$ pair multiplicity in central collisions is far below unity at SPS and FAIR-NICA energies. In this case, we may gate in the PHSD calculations on events with a single $c\bar{c}$ pair selected by Monte Carlo from the number of possible binary N + N reactions—and follow the dynamics of the charm quarks throughout the time evolution in PHSD, i.e., partonic scattering, hadronization by coalescence or fragmentation, and final hadronic rescattering of charmed mesons and baryons (see below). At the end, all observables have to be multiplied by the probability for the charm event as illustrated in Fig. 1(b). The shaded area in Fig. 1(b) shows the uncertainty in the number of $c\bar{c}$ pairs due to the uncertainty of the charm cross section in p + p collisions [cf. Fig. 1(a)]. Note that for $\sqrt{s_{NN}} < 20$ GeV no data are available and the number of $c\bar{c}$ pairs entirely stem from a parameterized function which takes into account the phase space of final states.

B. Fragmentation of charm and bottom in p+p collisions

The produced charm and bottom quarks in hard nucleonnucleon collisions are hadronized by emitting soft gluons, which is denoted by "fragmentation." As in Ref. [38], we use the fragmentation function of Peterson, which reads as [50]

$$D_Q^H(z) \sim \frac{1}{z[1 - 1/z - \epsilon_Q/(1 - z)]^2},$$
 (1)

where z is the momentum fraction of the hadron H fragmented from the heavy quark Q while ϵ_Q is a fitting parameter which is taken to be $\epsilon_Q = 0.01$ for charm [38] and 0.004 for bottom [39]. We note that the fragmentation function is applied only for the transverse momentum of the hadron while the rapidity is assumed to be the same as before the fragmentation. The chemical fractions of the charm quark decay into $D^+, D^0, D^{*+}, D^{*0}, D_s$, and Λ_c are taken to be 14.9, 15.3, 23.8, 24.3, 10.1, and 8.7% [39,51–53], respectively, and those of the bottom quark decay into $B^-, \overline{B}^0, \overline{B}^0_s$, and Λ_b are 39.9, 39.9, 11, and 9.2% [54]. After the momentum and the species of the fragmented particle are decided by Monte Carlo, the energy of the fragmented particle is adjusted to be on shell. Furthermore, the D^* mesons first decay into $D + \pi$ or $D + \gamma$, and then the D- mesons produce single electrons through the semileptonic decay [55], which is evaluated within PYTHIA.

III. HEAVY QUARK DYNAMICS IN A+A COLLISIONS

We here briefly recall the various interactions of charm quarks (or charm hadrons) in the partonic (hadronic) medium as introduced in Ref. [40].

A. Heavy-quark interactions in the QGP

In PHSD, the baryon-baryon and baryon-meson collisions at high-energy produce strings. If the local energy density is above the critical energy density ($\sim 0.5 \, \text{GeV}/\text{fm}^3$), the strings melt into quarks and antiquarks with masses determined by the temperature-dependent spectral functions from the DQPM [34]. Massive gluons are formed through flavor-neutral quark and antiquark fusion, in line with the DQPM. In contrast to normal elastic scattering, off-shell partons may change their mass after the elastic scattering according to the local temperature T in the cell (or local space-time volume) where the scattering happens. This automatically updates the parton masses as the hot and dense matter expands, i.e., the local temperature decreases with time. The same holds true for the reaction chain from gluon decay to quark + antiquark $(g \rightarrow q + \bar{q})$ and the inverse reaction $(q + \bar{q} \rightarrow g)$ following detailed balance. The local temperature is determined from the local energy density in the rest frame of the cell by employing the lattice QCD equation of state from Ref. [56].

Because of the finite spectral width of the partonic degrees of freedom, the parton spectral function has timelike as well as spacelike parts. The timelike partons propagate in space-time within the light cone while the spacelike components are attributed to a scalar potential energy density [33]. The gradient of the potential energy density with respect to the scalar density generates a repulsive force in relativistic heavy-ion collisions and plays an essential role in reproducing experimental flow data and transverse momentum spectra of hadrons with light quarks (see Ref. [30] for a review). For charm quarks, we assume in this study that the heavy quark has a constant (on-shell) mass: The charm quark mass is taken to be 1.5 GeV; however, the light quarks and antiquarks as well as gluons are treated fully off shell.

The heavy quarks and antiquarks produced in early hard collisions-as described above-interact with the dressed lighter off-shell partons in the QGP. The cross sections for the heavy-quark scattering with massive off-shell partons have been calculated by considering explicitly the mass spectra of the final-state particles in Refs. [57,58]. The elastic scattering of heavy quarks in the QGP is treated by including the nonperturbative effects of the strongly interacting quark-gluon plasma (sQGP) constituents, i.e., the temperature-dependent coupling $g(T/T_c)$, which rises close to T_c , the multiple scattering, etc. The multiple strong interactions of quarks and gluons in the sQGP are encoded in their effective propagators with broad spectral functions (imaginary parts). As pointed out above, the effective propagators, which can be interpreted as resummed propagators in a hot and dense QCD environment, have been extracted from lattice data in the scope of the DQPM [34].

We recall that the divergence encountered in the *t*-channel scattering is cured self-consistently, since the infrared regulator is given by the finite DQPM gluon mass and width. For further details, we refer the reader to Refs. [57,58].

We recall that charm interactions in the QGP—as described by the DQPM charm scattering cross sections—differ substantially form the pQCD scenario; however, the spacial diffusion constant for charm quarks $D_s(T)$ is consistent with the lQCD data [39,59] within error bars.

B. Heavy-quark hadronization

The heavy-quark hadronization in nucleus-nucleus collisions is realized via "dynamical coalescence" and fragmentation. Here "dynamical coalescence" means that the probability of finding a coalescence partner is calculated from the Wigner density in coordinate and momentum space and the coalescence is realized by Monte Carlo in the vicinity of the critical energy density $0.4 \le \epsilon \le 0.75 \,\text{GeV}/\text{fm}^3$ as described in Ref. [40]. We note that such a dynamical realization of heavyquark coalescence is in line with the dynamical hadronization of light quarks in the PHSD. Summing up the coalescence probabilities from all candidates, whether the heavy quark or heavy antiquark hadronizes by coalescence or not, and which quark or antiquark among the candidates will be the coalescence partner, is decided by Monte Carlo. If a random number is above the sum of the coalescence probabilities, it is tried again in the next time step until the local energy density is lower than $0.4 \text{ GeV}/\text{fm}^3$. The heavy quark or heavy antiquark, which does not succeed to hadronize by coalescence throughout the expansion phase of the partonic subsystem, then hadronizes through fragmentation as in p + p collisions. We recall that charm quarks with low transverse momenta p_T dominantly hadronize by coalescence while those with large p_T undergo fragmentation [40].

C. Interactions of charm mesons with the hadronic medium

After the hadronization of heavy quarks and their subsequent decay into D, D^* mesons, the final stage of the evolution concerns the interaction of these states with the hadrons forming the expanding bulk medium. A realistic description of the hadron-hadron scattering—potentially affected by resonant interactions—includes collisions with the states $\pi, K, \bar{K}, \eta, N, \bar{N}, \Delta, \bar{\Delta}$. A description of their interactions has been developed in Refs. [60–68] using effective field theory. Moreover, after the application of an effective theory, one should implement a unitarization method to the scattering amplitudes to better control the behavior of the cross sections at moderate energies.

The details of the interaction for the four heavy states follows quite in parallel by virtue of the "heavy-quark spinflavor symmetry." It accounts for the fact that if the heavy masses are much larger than any other typical scale in the system, like Λ_{QCD} , temperature, and the light hadron masses, then the physics of the heavy subsystem is decoupled from the light sector, and the former is not dependent on the mass nor on the spin of the heavy particle. This symmetry is exact in the ideal limit $m_Q \rightarrow \infty$, with m_Q being the mass of the heavy quark confined in the heavy hadron. In the opposite limit $m_Q \rightarrow 0$, one can exploit the chiral symmetry of the QCD Lagrangian to develop an effective realization for the light particles. This applies to the pseudoscalar meson octet (π, K, \bar{K}, η) . Although both symmetries are broken in nature (as in our approach, when implementing physical masses), the construction of the effective field theories incorporates the breaking of these symmetries in a controlled way. In particular, it provides a systematic expansion in powers of $1/m_H$ (inverse heavy-meson mass) and powers of p, m_l (typical momentum and mass of the light meson). Following these ideas, we use two effective Lagrangians for the interaction of a heavy meson with light mesons and with baryons, respectively.

In the scattering with light mesons, the scalar (D) and vector (D^*) mesons are much heavier than the pseudoscalar meson octet $(\pi, K, \overline{K}, \eta)$. The latter have, in addition, masses smaller than the chiral scale $\Lambda_{\chi} \simeq 4\pi f_{\pi}$, where f_{π} is the pion decay constant. In this case, one can exploit standard chiral perturbation theory for the dynamics of the (pseudo-) Goldstone bosons, and add the heavy-quark mass expansion up to the desired order to account for the interactions with heavy mesons. In our case, the effective Lagrangian is kept to next to leading order in the chiral expansion, but to leading order in the heavy-quark expansion [61,63]. From this effective Lagrangian, one can compute the tree-level amplitude (or potential), which describes the scattering of a heavy meson off a light meson as worked out in Refs. [66,67].

For the heavy meson-baryon interaction, we use an effective Lagrangian based on a low-energy realization of a *t*-channel vector meson exchange between mesons and baryons. In the low-energy limit, the interaction provides a generalized Weinberg-Tomozawa contact interaction as worked out in Refs. [60,62,64,65]. The effective Lagrangian obeys SU(6) spin-flavor symmetry in the light sector, plus heavy-quark spin symmetry (HQSS) in the heavy sector (which is preserved either the heavy quark is contained in the meson or in the baryon).

The tree-level amplitudes for meson-meson and mesonbaryon scattering have strong limitations in the energy range in which they should be applied. It is limited for those processes in which the typical momentum transfer is low and below any possible resonance. To increase the applicability of the treelevel scattering amplitudes and restore exact unitarity for the scattering-matrix elements, we apply a unitarization method, which consists in solving a coupled-channel Bethe-Salpeter equation for the unitarized scattering amplitude T_{ij} using the potential V_{ij} as a kernel,

$$T_{ij} = V_{ij} + V_{ik}G_kT_{kj},\tag{2}$$

where G_k is the diagonal meson-meson (or meson-baryon) propagator which is regularized by dimensional regularization in the meson-meson (or meson-baryon) channel. We adopt the "on-shell" approximation to the kernel of the Bethe-Salpeter equation to reduce it into a set of algebraic equations. We refer the reader to Refs. [60,62,64–67] for technical details and individual results.

The unitarization procedure allows for the possibility of generating resonant states as poles of the scattering amplitude T_{ij} in the complex plane. Even when these resonances are not explicit degrees of freedom, and we do not propagate them

The resulting (unitarized) cross sections for the binary scattering of D, D^* (with any possible charged states) with $\pi, K, \bar{K}, \eta, N, \bar{N}, \Delta, \bar{\Delta}$ are implemented in the PHSD code considering both elastic and inelastic channels. About 200 different channels are taken into account. Although the unitarization method helps to extend the validity of the tree-level amplitudes into the resonant region, one cannot trust the final cross sections for higher energies. Beyond the resonant region, the transition between the high- and low-energy regimes is interpolated such that the cross sections are continuous.

IV. DILEPTON PRODUCTION CHANNELS

We recall that in the hadronic sector PHSD is equivalent to the hadron-string-dynamics (HSD) transport approach [69] that has been used for the description of pA and AA collisions from SIS to SPS energies and has lead to a fair reproduction of hadron abundances, rapidity distributions, and transverse momentum spectra as well as dilepton spectra. In particular, HSD incorporates off-shell dynamics for vector mesons and a set of vector-meson spectral functions [70] that covers possible scenarios for their in-medium modification, i.e., in particular a collisional broadening of the vector resonances. Note that in the off-shell transport description, the hadron spectral functions change dynamically during the propagation through the medium and evolve toward the on-shell spectral function in the vacuum. The dilepton production by a (baryonic or mesonic) resonance R decay can be schematically presented in the following way:

$$BB \to RX,$$
 (3)

$$mB \to RX,$$
 (4)

$$R \to e^+ e^- X, \tag{5}$$

$$R \to mX, \quad m \to e^+ e^- X,$$
 (6)

$$R \to R'X, \quad R' \to e^+e^-X;$$
 (7)

i.e., in a first step, a resonance *R* might be produced in baryon-baryon (*BB*) or meson-baryon (*mB*) collisions. Then this resonance can couple to dileptons directly (5) (e.g., Dalitz decay of the Δ resonance: $\Delta \rightarrow e^+e^-N$) or decays to a meson *m* (+baryon) (6), which produces dileptons via direct decays (ρ, ω, ϕ) or Dalitz decays (π^0, η, ω). The resonance *R* might also decay into another resonance *R'* (7) which later produces dileptons via Dalitz decay. Note that in the combined model the final particles—which couple to dileptons—can be produced also via nonresonant mechanisms, i.e., "background" channels at low and intermediate energies or string decay at high energies. In addition to the hadronic channels above, we account for the 4π channels, i.e., the dilepton production in the two-body reactions $\pi + \rho$, $\pi + \omega$, $\rho + \rho$, and $\pi + a_1$ as described in detail in Ref. [71]. The latter provide the background from hadronic channels in the intermediate-mass regime 1.2 GeV < M < 3 GeV [71], which is not shown explicitly in this study since the contribution of 4π channels is much smaller than the contribution from open charm decays and the QGP radiation.

We recall that the influence of in-medium effects on the vector mesons (ρ, ω, ϕ) has been extensively studied within the PHSD approach in the past (cf. Refs. [30,70,71]) and it has been shown that the collisional broadening scenario for the in-medium vector-meson spectral functions is consistent with experimental dilepton data from SPS to LHC energies in line with the findings by other groups [29]. Accordingly, in the present study we will adopt the collisional broadening scenario for the vector-meson spectral functions as the default scenario.

In order to address the electromagnetic radiation of the partonic phase, off-shell cross sections of $q\bar{q} \rightarrow \gamma^*$, $q\bar{q} \rightarrow \gamma^* g$, and $qg \rightarrow \gamma^* q$ ($\bar{q}g \rightarrow \gamma^* \bar{q}$) reactions taking into account the effective propagators for quarks and gluons from the DQPM have been calculated in Ref. [72]. Here γ^* stands for the e^+e^- or $\mu^+\mu^-$ pair. Dilepton production in the QGP—as created in early stages of heavy-ion collisions—is calculated by implementing these off-shell processes into the PHSD transport approach on the basis of the same partonic propagators as used for the time evolution of the partonic system. For a review on electromagnetic production channels within PHSD, we refer the reader to Ref. [30] and for the details of the dilepton cross sections from off-shell partonic channels to Appendix A.

V. RESULTS FOR HEAVY-ION REACTIONS

So far we have described the interactions of the heavy flavor—produced in relativistic heavy-ion collisions—with partonic and hadronic degrees of freedom. Since the matter produced in heavy-ion collisions is extremely dense, the interactions with the bulk matter suppresses heavy flavors at high p_T . On the other hand, the partonic or nuclear matter is accelerated outward (exploding), and a strong flow is generated via the interactions of the bulk particles and the repulsive scalar interaction for partons. Since the heavy flavor strongly interacts with the expanding matter, it is also accelerated outward. Such effects of the medium on the heavy-flavor dynamics are expressed in terms of the nuclear modification factor defined as

$$R_{AA}(p_T) \equiv \frac{dN_{AA}/dp_T}{N_{\text{binary}}^{AA} \times dN_{pp}/dp_T},$$
(8)

where N_{AA} and N_{pp} are, respectively, the number of particles produced in heavy-ion collisions and that in p + p collisions, and N_{binary}^{AA} is the number of binary nucleon-nucleon collisions in the heavy-ion collision for the centrality class considered. Note that if the heavy flavor does not interact with the medium in heavy-ion collisions, the numerator of Eq. (8) will be similar to the denominator. For the same reason, a R_{AA} smaller (larger) than one in a specific p_T region implies that the nuclear matter suppresses (enhances) the production of heavy flavors in that transverse momentum region.

In noncentral heavy-ion collisions, the produced matter expands anisotropically due to the different pressure gradients between in plane and out of plane. If the heavy flavor interacts strongly with the nuclear matter, then it also follows this



FIG. 2. The transverse momentum spectra of *D* mesons (a) and the $R_{AA}(p_T)$ of single electrons from semileptonic decay of *D* mesons (b) as a function of the transverse momentum p_T in central Pb + Pb collisions from PHSD at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV at midrapidity.

anisotropic motion to some extent. The anisotropic flow is expressed in terms of the elliptic flow v_2 which reads

$$v_2(p_T) \equiv \frac{\int d\phi \cos 2\phi (dN_{AA}/dp_T d\phi)}{2\pi dN_{AA}/dp_T},$$
(9)

where ϕ is the azimuthal angle of a particle in momentum space.

A. Nuclear modification of dielectrons from heavy flavor

In this subsection, we focus on the $c\bar{c}$ dynamics, the dielectrons produced from heavy-flavor pairs, and their modification in relativistic heavy-ion collisions.

Figure 2(a) shows the transverse momentum spectra of *D* mesons in central Pb+Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV for |y| < 1. Since the cross section for charm production increases with collision energy as shown in Fig. 1(a), the transverse momentum spectrum of *D* meson enhances strongly with increasing collision energy and also becomes harder.

Figure 2(b) displays the nuclear modification factor of single electrons from *D* meson semileptonic decays at midrapidity (|y| < 1) for the same set of central Pb + Pb collisions. We

mention that for the semileptonic decays of heavy flavors we use the subroutine "pydecay" of the PYTHIA event generator [41]. Contrary to the R_{AA} at RHIC and LHC energies, we find ratios well above unity at $\sqrt{s_{NN}} = 8$ and 11.5 GeV, which implies an enhancement of the yield (at higher momenta) rather than the familiar suppression at RHIC and LHC. The enhanced R_{AA} at low energies (8 and 11.5 GeV) may be dominantly attributed to the Fermi motion of nucleons in the colliding nuclei, which does not exist in p + p collisions and slightly increases the collision energy in binary nucleon-nucleon scattering. Since the collision energies are close to the threshold energy for charm-pair production, where the production cross section increases rapidity as shown in Fig. 1(a), a small enhancement of the collision energy gives a sizeable increase of the charm production and subsequently the decay products. We note in passing that the R_{AA} of single electrons at $\sqrt{s_{NN}} = 39$ and 200 GeV is consistent with our recent results in Ref. [40], where the R_{AA} is shown also for higher transverse momenta.

Since heavy flavor is always produced by pairs, there is an angular correlation between the heavy quark and heavy antiquark. If the heavy quark and antiquark from the same pair (through semileptonic decays) produce a positron and an electron, respectively, the produced dielectron also has an angular correlation. On the other hand, the matter produced in heavy-ion collisions changes the transverse momentum of each heavy flavor and consequently also the angular correlation of the heavy flavor pair. It has been suggested that the analysis of the azimuthal angular correlation might provide information on the energy loss mechanism of heavy quarks in the QGP [73], because stronger interactions should result in less pronounced angular correlations. Since in the PHSD we can follow up the fate of an initial heavy quark-antiquark pair throughout the partonic scatterings, the hadronization and final hadronic rescatterings, the microscopic calculations allow us to shed some light on the correlation between the in-medium interactions and the final angular correlations.

Figure 3 shows the azimuthal angular distribution between the transverse momentum of charm (D) and that of anticharm (D) for each charm pair at midrapidity (|y| < 1) before (dashed lines) and after the interactions with the medium (solid lines) in central Pb+Pb collisions at $\sqrt{s_{NN}} = 17.3$ and 200 GeV. The azimuthal angle between the initial charm and anticharm quarks is provided by the PYTHIA event generator and peaks around $\phi = 0$ for $\sqrt{s_{NN}} = 17.3$ GeV, while we find a maximum around $\phi = \pi$ for $\sqrt{s_{NN}} = 200$ GeV. After the interaction with the hadronic and partonic matter, however, the azimuthal angle between the D and \overline{D} has a maximum near $\phi = 0$ at both collision energies. In other words, the azimuthal angle changes little in low-energy collisions, but changes considerably in high-energy collisions. As shown in our previous study [39], the shift of the maximum in the azimuthal angle from π to 0 at $\sqrt{s_{NN}} = 200 \text{ GeV}$ can be attributed to the strong interaction of charm with radial flow.

Figure 4 shows the invariant mass spectra of dielectrons from charm pairs with (red lines) and without the interactions with hot medium (blue lines) in central Pb + Pb collisions at $\sqrt{s_{NN}} = 17.3$ [Fig. 4(a)] and 200 GeV [Fig. 4(b)]. We can see that the invariant mass spectrum of dielectrons changes little for $\sqrt{s_{NN}} = 17.3$ GeV, while it is considerably suppressed at



FIG. 3. Azimuthal angular distribution between the transverse momentum of a heavy-flavor meson and that of an antiheavy-flavor meson for each heavy flavor pair at midrapidity (|y| < 1) before (dashed lines) and after the interactions with the medium (solid lines) in central Pb + Pb collisions at $\sqrt{s_{NN}} = 17.3$ (a) and 200 GeV (b).

large invariant mass at $\sqrt{s_{NN}} = 200$ GeV. This suppression can be understood from Figs. 2 and 3, considering that the invariant mass of the dielectron depends on the momenta of electron and positron and also on the angle between them. Figures 2 and 3 clearly show that the momenta of electron and positron are suppressed and the azimuthal angle between them decreases at $\sqrt{s_{NN}} = 200$ GeV; both effects decrease the invariant mass of the dielectron. On the other hand, the momenta of electron and positron and the azimuthal angle do not change much at $\sqrt{s_{NN}} = 17.3$ GeV such that the dielectron spectrum stays approximately unchanged.

B. Excitation function of dielectron production in Pb + Pb collisions from $\sqrt{s_{NN}} = 8$ to 200 GeV

As mentioned in the previous sections, the dileptons produced in relativistic heavy-ion collisions can be classified into three parts: (i) dileptons from heavy-flavor pairs, (ii) from partonic scatterings in the QGP phase, and (iii) from hadronic interactions in the hadronic (HG) phase. In this subsection, we compare the separate contributions in central Pb+Pbcollisions at various energies from 8 to 200 GeV.



FIG. 4. Invariant mass spectra of dielectrons from charm pairs with (red lines) and without the interactions with the hot medium (blue lines) in central Pb+Pb collisions at $\sqrt{s_{NN}} = 17.3$ (a) and 200 GeV (b).

Figure 5 shows the dielectron mass spectra from hadronic channels (a), from partonic interactions in the QGP (b), and from the semileptonic decays of $D\bar{D}$ pairs (c) in central Pb + Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV at midpseudorapidity $|\eta_e| < 1$ for the leptons. We find that the contribution from the hadronic channels increases moderately with collision energy (in line with the hadron abundances), the contribution from the QGP raises more steeply (in line with the enhanced space-time volume of the QGP phase), and that from $D\bar{D}$ pairs is most dramatically increasing [in line with the number of $c\bar{c}$ pairs, cf. Fig. 1(b)]. Accordingly, the contribution from heavy flavor is small at low-energy collisions, but becomes more important with increasing collision energy in competition with the production from the QGP channels.

In order to show the separate contributions explicitly, we compare in Fig. 6 the contributions from the QGP (red lines) and from $D\bar{D}$ pairs (green lines) with the total dielectron spectrum (blue lines) at different collision energies for central Pb+Pb collisions. In low-energy collisions, the dielectrons from hadronic channels dominate in the low-mass region and those from partonic interactions dominate in the intermediate-mass range while the contribution from $D\bar{D}$ pairs is negligible. With increasing collision energy, the contribution from $D\bar{D}$



10

10

10⁻

10⁻²

10⁻³

10⁻⁴ 10^{-€}

dN^{ee}/dM (GeV⁻¹



8 GeV

FIG. 5. The invariant mass spectra of dileptons from the hadronic sources (HG) (a), the QGP (b), and $D\bar{D}$ pairs (c) in central Pb + Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV from the PHSD.

pairs becomes more significant and comparable to that from partonic interactions at $\sqrt{s_{NN}} \approx 39$ GeV in the intermediatemass range. Finally, it overshines the partonic contribution at $\sqrt{s_{NN}} = 200$ GeV (and above).

Figure 7 compares the contributions from $D\bar{D}$ pairs (green lines) to three partonic channels, i.e., $q + \bar{q} \rightarrow e^+ + e^-$, $q + \bar{q} \rightarrow g + e^+ + e^-$, and $q(\bar{q}) + g \rightarrow q(\bar{q}) + e^+ + e^-$, for



FIG. 6. The invariant mass spectra of dileptons from partonic interactions (red lines) and from $D\bar{D}$ pairs (green lines) together with total dielectron spectrum including hadronic contributions (blue lines) in central Pb + Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV from the PHSD at midpseudorapidity for the leptons.

intermediate mass dileptons (1.2 GeV < M < 3 GeV) as a function of collision energy $\sqrt{s_{NN}}$ for Pb+Pb collisions at b = 2 fm. The figure clearly shows that the contribution from partonic interactions, especially from $q + \bar{q} \rightarrow e^+ + e^-$, dominates the intermediate-mass range in low-energy collisions. However, the contribution from $D\bar{D}$ pairs rapidly increases

with increasing collision energy, because the scattering cross section for charm production grows fast above the threshold energy as shown in Fig. 1(a). It overshines the contribution from partonic interactions around $\sqrt{s_{NN}} \approx 40$ GeV and dominates at higher energies. Since the detectors of different collaborations have a different acceptance, we show in Fig. 7(b) the results



FIG. 7. The contributions to intermediate-mass dielectrons (1.2 GeV < M < 3 GeV) from $D\bar{D}$ pairs (green lines), different channels of partonic interactions, $q + \bar{q} \rightarrow e^+ + e^-$, $q + \bar{q} \rightarrow g + e^+ + e^-$, $q(\bar{q}) + g \rightarrow q(\bar{q}) + e^+ + e^-$ (see legend) as a function of $\sqrt{s_{NN}}$ for Pb+Pb collisions at b = 2 fm (for midrapidity leptons). The red solid line displays the sum of the partonic contributions.

without any acceptance cuts, while Fig. 7(a) shows the results for a midpseudorapidity cut on leptons of $|\eta^e| < 1$. However, the contributions from the partonic interactions and from $D\bar{D}$ pairs show similar behavior in both cases.

One of most important issues in heavy-ion physics is to find and study the properties of partonic nuclear matter which is created in a small space-time volume in relativistic heavy-ion collisions. To this end, one needs observables that are not blurred by hadronic interactions. Our results in Figs. 6 and 7 clearly demonstrate that the window to study partonic matter by dielectrons at intermediate masses without substantial background from heavy flavor decays opens for collision energies $\sqrt{s_{NN}} < 40$ GeV.

C. Transverse mass spectra at midrapidity

In this subsection, we explore central Pb + Pb collisions at various energies with a focus on the transverse mass spectra of dileptons with intermediate mass at midrapidity. To this end, we show in Fig. 8 the Lorentz-invariant transverse mass spectra for (b = 2 fm) Pb + Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5,





FIG. 8. The transverse mass spectra of dileptons with the invariant mass between 1.2 and 3 GeV from the QGP (a), $D\overline{D}$ pairs (b), and all sources (c) in central Pb + Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV from the PHSD. The fat solid lines show exponential fits to the PHSD results in the transverse mass range [1.75, 2.95] GeV.

17.3, 39, and 200 GeV for the dielectrons with the invariant mass between 1.2 and 3 GeV from the QGP [Fig. 8(a)], from D mesons [Fig. 8(b)], and the dileptons from all channels (including especially D, \overline{D} decay) [Fig. 8(c)]. All spectra show an approximately exponential decay (fat solid lines) in the transverse mass m_T for 1.75 GeV $< m_T < 2.95$ GeV, which



FIG. 9. The inverse slope parameters of intermediate-mass dielectrons from the QGP (red line with dots), $D\bar{D}$ pairs (green line with squares), and all sources (blue line with triangles) in central Pb+Pb collisions at $\sqrt{s_{NN}} = 8$, 11.5, 17.3, 39, and 200 GeV from the PHSD.

can be characterized by an inverse slope parameter β which is different for dileptons from open charm and those from the QGP at all bombarding energies.

The excitation function in the inverse slope parameters β is shown in Fig. 9 for the three cases of Fig. 8, i.e., dileptons with the invariant mass between 1.2 and 3 GeV from the QGP (red line with dots), $D\overline{D}$ pairs (green line with squares), and all dilepton sources (blue line with triangles) in central Pb+Pb collisions at $\sqrt{s_{NN}} = 8, 11.5, 17.3, 39$, and 200 GeV. We find that the inverse slope parameter from the QGP contribution (red line with dots) is larger than the contribution from D decays (green line with squares) at all energies and almost identical to the inverse slope for the total dilepton spectra (blue line with triangles) in the transverse mass range [1.75,2.95] GeV at SPS energies. Since the contribution from the D decays increases with bombarding energy, a small wiggle in $\sqrt{s_{NN}}$ can be found in the inverse slope for the total dilepton spectra (blue line with triangles) in the lower RHIC energy regime. This wiggle should be seen in experiments provided that high statistics data become available for intermediate-mass dileptons.

VI. PHSD VERSUS EXPERIMENTAL DATA AND PREDICTIONS FOR THE TOP LHC ENERGY

A. Au + Au and Pb + Pb collisions from 19.6 GeV to 2.76 TeV

In this section, we compare the invariant mass spectra of dielectrons from the PHSD to the experimental data in Au + Au collisions from $\sqrt{s_{NN}} = 19.6$ to 200 GeV from the STAR collaboration and those in Pb + Pb collisions from the ALICE collaboration at $\sqrt{s_{NN}} = 2.76$ TeV. We note that the experimental data from the STAR Collaboration and those from the ALICE Collaboration have different centralites and different acceptance cuts. The STAR data are obtained for minimumbias Au + Au collisions and electrons and positrons with transverse momenta $p_T \ge 0.2$ GeV and pseudorapidities $|\eta^e| < 1.0$. On other hand, the ALICE data are obtained for 0–10% central Pb + Pb collisions and the electrons and positrons with transverse momenta $p_T \ge 0.4$ GeV and pseudorapidities $|\eta^e| < 0.8$. The sensitivity of the invariant mass spectra of dielectrons to the cross section for charm production and cuts in p_T and pseudorapidity η^e is discussed in more detail in Appendix B.

The first five panels of Fig. 10 show the invariant mass spectra of dielectrons from the beam-energy scan (BES) at $\sqrt{s_{NN}} = 19.6, 27, 39$, and 62.4 GeV and from the top RHIC energy. As discussed in the previous subsection, the contribution from hadrons is dominant in the low-mass region and signals a broadening of the ρ -meson spectral function in dense nuclear matter (cf. Ref. [30]. On the other hand, the intermediate-mass range originates predominantly by dielectrons from partonic interactions and those from heavy flavor decays. Similar to the Pb + Pb collisions in Fig. 6, the contribution from heavy flavor becomes more important with increasing collision energy. The contribution from heavy flavors and from partonic interactions cross around invariant masses $M \approx 1$ GeV in Au+Au collisions at $\sqrt{s_{NN}} = 19.6$ GeV. However, the crossing point shifts to 1.6 GeV at $\sqrt{s_{NN}} = 27$ GeV and to ~ 2.0 GeV at $\sqrt{s_{NN}} = 39$ and 62.4 GeV. At the top RHIC energy they cross at \sim 2.4 GeV.

The last panel of Fig. 10 is the invariant mass spectrum of dielectrons in central Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV. As in Au+Au collisions at the RHIC energies, the lowmass range is dominated by the dielecrons from hadronic channels and the intermediate-mass region by those from partonic interactions and heavy-flavor decays. However, the crossing point of the contribution from partonic interactions and that from heavy flavor is lower than at the top RHIC energy, which is due to a couple of effects: (i) The cross section for charm production no longer increases rapidly at the LHC energies as shown in Fig. 1(a). It is also seen in Fig. 1(b), which shows the number of produced charm pairs as a function of collision energy. As a result, the growth in the number of produced charm pairs is not faster than the growth of dielectrons from partonic interactions. Additionally, the shadowing effect, which is the modification of the parton distribution function in nuclei [74], considerably suppresses charm production at the LHC energies [39] (see below). (ii) Another reason is the stronger suppression of the charm four-momentum by partonic scattering at the LHC energies. As already discussed in the context of Fig. 4, the strong interaction of heavy flavor with the medium reduces the invariant mass of dielectrons. Since the interaction is stronger at the LHC energies, we can expect a larger suppression of the dielectron spectrum at larger invariant masses. (iii) Furthermore, at the LHC energies the contribution from semileptonic $B\bar{B}$ decays becomes important. Comparing the lower two panels of Fig. 10, the contribution from $B\bar{B}$ decays is found to be larger than that from $D\bar{D}$ decays above $M \approx 2.2$ GeV in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV, while the contribution from $B\bar{B}$ decays is larger only above $M \approx 2.8 \,\text{GeV}$ in Au + Au collisions $\sqrt{s_{NN}} =$ 200 GeV. Since the contribution from $B\bar{B}$ decays amounts to about 50% of the contribution from partonic interactions at the LHC energies, it will distort the information on partonic matter in the intermediate-mass range of the dielectron spectrum.

Besides the interesting points mentioned above, we close this subsection with the comment that the dilepton invariant



FIG. 10. The invariant mass spectra of dielectrons from the PHSD in comparison to the STAR data in Au + Au collisions from $\sqrt{s_{NN}} = 19.6$ to 200 GeV [75,76] and to the ALICE data in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV [77]. The total yield is displayed in terms of the blue lines while the different contributions are specified in the legends. Note that the contribution from J/Ψ and Ψ' decays are not included in the PHSD calculations.

mass spectra from the PHSD describe reasonably well the available experimental data for collision energies from 19.6 GeV to 2.76 TeV, although the experimental data at $\sqrt{s_{NN}} = 2.76$ TeV are available only for invariant masses $M \leq 1$ GeV.

B. Predictions for central Pb + Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV

Based on the successful description of experimental data from the BES for $\sqrt{s_{NN}} = 19.6$ GeV to the LHC energy at $\sqrt{s_{NN}} = 2.76$ TeV, we here provide predictions for dielectron production in central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV.



FIG. 11. The $R_{AA}(p_T)$ (a) and the elliptic flow $v_2(p_T)$ (b) for *D* meson at 0–10% central Pb + Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV (from PHSD) as a function of the transverse momentum with (solid line) and without shadowing effects (dashed line). Experimental data are from the CMS collaboration [78,79].

As mentioned above, a proper description of heavy-flavor production and interactions in heavy-ion collisions is necessary to allow for reliable predictions.

Figure 11 shows the R_{AA} [Fig. 11(a)] and the elliptic flow v_2 [Fig. 11(b)] of D mesons as functions of transverse momentum in 0–10% central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV. In both panels, the dashed lines are the results without the shadowing effect and the solid lines with EPS09 shadowing [74] included. The upper panel shows that shadowing reduces the R_{AA} considerably at low transverse momentum, which can be explained as follows: If the collision energy is very large, charm quark pairs with small transverse momentum are dominantly produced by partons with a small energy-momentum fraction x of the nucleon. On the other hand, the parton distribution function of a nucleon in a heavy nucleus is considerably suppressed at small x in such high-energy collisions [74]. which leads to a suppression of charm production at low transverse momentum. Figure 11(a) clearly shows that the shadowing effect is necessary to explain the experimental data from the ALICE Collaboration. We note that the PHSD results





FIG. 12. The invariant mass spectra of dielectrons for 0–10% central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV from the PHSD for $|p_T^e| > 0.4$ GeV and $|\eta^e| < 0.8$.

are presently available only up to $p_T = 20 \text{ GeV/c}$ due to the limited statistics and huge CPU time required. In the case of the open charm elliptic flow $v_2(p_T)$, the statistics do no allow for robust results for $p_T > 6 \text{ GeV/c}$. On the other hand, the shadowing effect is seen to have no substantial effect on the elliptic flow of *D* mesons up to $p_T \approx 6 \text{ GeV/c}$ since shadowing changes the production of charm from initial hard collisions but does not change the interactions of produced charm in the partonic medium. Figure 11 demonstrates that both the R_{AA} and the elliptic flow v_2 of *D* mesons are approximately described at $\sqrt{s_{NN}} = 5.02$ TeV by the PHSD. Although the v_2 of *D* mesons is slightly underestimated, this will have practically no effect on the dielectron spectrum.

Figure 12 shows the prediction from PHSD for the invariant mass spectra of dielectrons in 0–10% central Pb+Pb collisions at $\sqrt{s_{NN}} = 5.02$ TeV within the acceptance ($p_T > 0.4$ GeV, $|\eta^e| < 0.8$) as in Fig. 10(f). Comparing with the results at 2.76 TeV, we find no dramatic change in the shape of the spectrum except for an overall enhancement of the dielectron yield. The yields of dielectrons from hadronic channels, from partonic interactions, and from heavy-flavor decays are, respectively, enhanced by 55%, 54%, and 36% at $\sqrt{s_{NN}} = 5.02$ TeV. We note that the dielectron yield from hadronic channels and that from partonic interactions increase by a similar amount, suggesting that both dielectron yields are produced from bulk matter whereas the dielectron yield from heavy-flavor decays is less enhanced due to a lower increase in the charm production cross section.

VII. SUMMARY

We have studied correlated electron (e^+e^-) production through the semileptonic decay of charm hadrons in relativistic heavy-ion collisions from $\sqrt{s_{NN}} = 8$ GeV to 5 TeV within the PHSD transport approach in extension of our work on *D*-meson production in relativistic heavy-ion collisions at RHIC and LHC energies [38–40] and low-mass dilepton production from SIS to RHIC energies [30].

In the PHSD, the charm partons—produced by the initial hard nucleon-nucleon scattering—interact with the massive

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quarks and gluons in the QGP by using the scattering cross sections calculated in the dynamical quasiparticle model (DQPM), which reproduces heavy-quark diffusion coefficients from lattice QCD calculations at temperatures above the deconfinement transition. When approaching the critical energy density for the phase transition from above, the charm (anti)quarks are hadronized into D mesons through the coalescence with light (anti)quarks. Those heavy quarks, which fail in coalescence until the local energy density is below $0.4 \,\text{GeV}/\text{fm}^3$, hadronize by fragmentation as in p + p collisions. The hadronized D mesons then interact with light hadrons in the hadronic phase with cross sections that have been calculated in an effective Lagrangian approach with heavy-quark spin symmetry. Finally, after freeze-out of the D mesons, they produce single electrons through semileptonic decays with the branching ratios given by the PYTHIA event generator.

The dilepton production from hadronic and partonic channels in central Pb + Pb (or Au + Au) collisions has been calculated, including also the contribution from the semileptonic decays of heavy flavors in PHSD for the first time on a fully microscopic level. We recall that also the cross sections for dilepton production have been calculated by employing the same propagators and couplings as incorporated in the partonic dynamics in PHSD (cf. Appendix A). We find that even in central Pb+Pb collisions at $\sqrt{s_{NN}} = 8$ to 20 GeV, the contribution from D, \overline{D} mesons to the intermediate mass dilepton spectra is subleading and one should have a rather clear signal from the QGP radiation whereas at the top RHIC energy this contribution overshines the intermediate mass dileptons from the QGP. It is interesting to note that the dielectrons from D, \bar{D} mesons do not increase any more relative to partonic interactions at the LHC energies for a couple of reasons: (i) the cross section for charm production does not grow as fast as at low energies; (ii) the shadowing effects, which suppress charm production at low transverse momentum, are stronger at LHC than at RHIC energies (cf. Fig. 11); (iii) the charm quark pair loses more four-momentum in the partonic medium produced at the LHC, which suppresses the invariant mass of the dielectrons from the semileptonic decays. Furthermore, the contribution from B, \overline{B} meson decays becomes more important and supersedes the contribution from D,D meson decays above M = 2.2-2.3 GeV at the LHC energies and amounts to about half the contribution from partonic interactions. All these effects strongly distort the information about partonic matter from intermediate-mass dielectrons at the LHC energies. The dilepton spectra at lower masses ($0.2 \text{ GeV} \leq M \leq 0.7 \text{ GeV}$) at SPS, FAIR-NICA, and BES RHIC energies show some sensitivity to the medium modification of the ρ meson, where the data favor an inmedium broadening as pointed out in the earlier studies on dilepton production reviewed in Refs. [29,30].

Additionally, we have explored the transverse mass spectra of dileptons in the invariant mass range from 1.2 to 3 GeV in central Pb + Pb collisions for $\sqrt{s_{NN}} = 8$ to 200 GeV and find approximately exponential spectra for transverse masses in the energy range [1.75,2.95] GeV (cf. Fig. 8). Since the inverse slope parameters differ for the contributions from the QGP

and are higher than that from D decays, we expect a wiggle in the excitation function of the inverse slope parameter for these intermediate mass dileptions (cf. Fig. 9), which should be seen experimentally in high statistics data.

In general, the PHSD calculations compare well with the available dilepton data from the BES program at RHIC as well as the LHC energy of $\sqrt{s_{NN}} = 2.76$ TeV where, unfortunately, only low-mass dilepton data are available so far. Explicit predictions for central Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV have been provided (cf. Fig. 12); however, the partonic contribution in the intermediate-mass range has a large background from D, \bar{D} as well as B, \bar{B} correlated semileptonic decays. As noted above, this background—in the intermediate-mass range—is by far subleading at lower SPS and FAIR-NICA energies, which provides promising perspectives for the future dilepton measurements at these facilities and allows for a fresh look at the electromagnetic radiation from the QGP at finite baryon chemical potential.

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APPENDIX A

1. Leading-order contribution

In this Appendix, we provide the details on the cross sections for the processes $q + \bar{q} \rightarrow \gamma^*(e^+e^-)$, $q + g \rightarrow q + \gamma^*(e^+e^-)$, and $q + \bar{q} \rightarrow g + \gamma^*(e^+e^-)$, considering the offshell nature of the interacting partons in line with the DQPM.

The invariant matrix element for the process $q + \bar{q} \rightarrow l^+ + l^-$, as shown in Fig. 13, is given by

$$M = \bar{u}(p_3)ie\gamma^{\nu}v(p_4)\frac{-ig_{\mu\nu}}{q^2}\bar{v}^a(p_2)ig\gamma^{\mu}\delta_{ab}u^b(p_1), \quad (A1)$$



FIG. 13. Feynman diagram for $q + \bar{q} \rightarrow l^+ + l^-$.

where *a* and *b* are the color indices of the incoming quark and antiquark. The matrix element squared then reads as

$$|\overline{M}|^{2} = \frac{8(4\pi\alpha)^{2}}{N_{c}} \frac{1}{s^{2}} \{ (p_{1} \cdot p_{3})(p_{2} \cdot p_{4}) + (p_{1} \cdot p_{4})(p_{2} \cdot p_{3}) + m_{1}m_{2}(p_{3} \cdot p_{4}) + m_{l}^{2}(p_{1} \cdot p_{2}) + 2m_{l}^{2}m_{1}m_{2} \},$$
(A2)

with $N_c = 3$ for the number of colors. Here, m_1 , m_2 , and m_l are, respectively, quark and antiquark masses and lepton mass, and the following color algebra is used:

$$|\overline{M}|^{2} \sim \frac{1}{N_{c}^{2}} \bar{v}^{a}(p_{2}) \delta_{ab} u^{b}(p_{1}) \bar{u}^{b'}(p_{1}) \delta_{a'b'} v^{a'}(p_{2})$$
$$\sim \frac{1}{N_{c}^{2}} \delta_{aa'} \delta_{ab} \delta_{a'b'} \delta_{bb'} = \frac{1}{N_{c}^{2}} \delta_{aa} = \frac{1}{N_{c}}.$$
 (A3)



FIG. 14. Feynman diagrams for the process $q + g \rightarrow q + \gamma^*$.

We note that m_1 is not necessarily equal to m_2 since the masses of the incoming quark and antiquark have spectral distributions which depend on the local temperature as defined by the DQPM. The phase-space integration for the scattering cross section is straightforward (cf. Ref. [55]).

2. Next-to-leading-order contributions

The invariant matrix element for the process $q + g \rightarrow q + \gamma^*(e^+e^-)$ in Fig. 14 reads

$$M = \bar{u}_{b}(p_{5}) \bigg[T^{a}_{bc} \frac{ig\gamma^{\alpha}i(p_{5} - p_{2} + m_{q})\delta_{cd}ie_{q}\gamma^{\beta}}{(p_{2} - p_{5})^{2} - m_{q}^{2} + i2(p_{20} - p_{50})\Gamma_{q}} + \frac{\delta_{bc}ie_{q}\gamma^{\beta}i(p_{1} + p_{2} + m_{q})ig\gamma^{\alpha}}{(p_{1} + p_{2})^{2} - m_{q}^{2} + i2(p_{10} + p_{20})\Gamma_{q}} T^{a}_{cd} \bigg] u_{d}(p_{1})\varepsilon^{a*}_{\alpha}(p_{2})L_{\beta}, \quad (A4)$$

where p_1 , p_2 , and p_5 are the momenta of incoming quark, incoming gluon, and outgoing quark with p_{10} , p_{20} , and p_{50} denoting their zeroth components; d, a, and b are their color indices; α and β are spin indices of the incoming gluon and outgoing photon; e_q is the electric charge of the quark; m_q and Γ_q are, respectively, the mass and the spectral width of the intermediate quark; and L_β is defined as

$$L_{\beta} = \bar{u}_b(p_3)ie\gamma_{\beta}v(p_4)/(iq^2), \tag{A5}$$

with p_3 and p_4 denoting the momenta of electron and positron and $q = p_3 + p_4$. In the case of real photon production, L_β is replaced by $\varepsilon_\alpha(q)$. The first term in Eq. (A4) is the *u* channel and the second term is the *s* channel. We note that the imaginary part of the quark propagator in the *u* channel is taken proportional to $p_{20} - p_{50}$ (downward) rather than to $p_{50} - p_{20}$ (upward), because the gluon mass is roughly twice the quark mass.

The invariant matrix element squared is written as follows:

$$\begin{split} |M|^{2} &= 8(4\pi)^{2} \alpha \alpha_{s} \left(\frac{e_{q}}{e}\right)^{2} \frac{N_{c}^{2} - 1}{2} L_{\beta} L_{\beta'}^{*} \left[\left\{ (4m_{5}m_{q} - 2p_{5} \cdot (p_{5} - p_{2})) \right\} \left\{ p_{1}^{\beta} (p_{5} - p_{2})^{\beta'} + p_{1}^{\beta'} (p_{5} - p_{2})^{\beta} - p_{1} \cdot (p_{5} - p_{2})g^{\beta\beta'} \right\} \\ &+ \left\{ (p_{5} - p_{2})^{2} - m_{q}^{2} \right\} \left\{ p_{5}^{\beta} p_{1}^{\beta'} + p_{5}^{\beta'} p_{1}^{\beta} - p_{5} \cdot p_{1} g^{\beta\beta'} \right\} + 2m_{1}g^{\beta\beta'} \left\{ m_{5}(p_{5} - p_{2})^{2} + m_{5}m_{q}^{2} - m_{q} p_{5} \cdot (p_{5} - p_{2}) \right\} \right) \\ &\times \frac{-1}{\left\{ (p_{2} - p_{5})^{2} - m_{q}^{2} \right\}^{2} + 4(p_{20} - p_{50})^{2}\Gamma_{q}^{2}} + \left\{ (4m_{1}m_{q} - 2p_{1} \cdot (p_{1} + p_{2})) \right\} \left\{ p_{5}^{\beta} (p_{1} + p_{2})^{\beta'} + p_{5}^{\beta'} (p_{1} + p_{2})^{\beta'} \right. \\ &- p_{5} \cdot (p_{1} + p_{2})g^{\beta\beta'} \right\} + \left\{ (p_{1} + p_{2})^{2} - m_{q}^{2} \right\} \left\{ p_{5}^{\beta} p_{1}^{\beta'} + p_{5}^{\beta'} p_{1}^{\beta} - p_{5} \cdot p_{1}g^{\beta\beta'} \right\} \\ &+ 2m_{5}g^{\beta\beta'} \left\{ m_{1}(p_{1} + p_{2})^{2} + m_{1}m_{q}^{2} - m_{q}p_{1} \cdot (p_{1} + p_{2}) \right\} \frac{-1}{\left\{ (p_{1} + p_{2})^{2} - m_{q}^{2} \right\}^{2} + 4(p_{10} + p_{20})^{2}\Gamma_{q}^{2}} \\ &- 2\left(-(p_{5} \cdot p_{1})\left\{ (p_{5} - p_{2})^{\beta} (p_{1} + p_{2})^{\beta'} + (p_{5} - p_{2}) \cdot (p_{1} + p_{2})g^{\beta\beta'} - (p_{5} - p_{2})^{\beta'} (p_{1} + p_{2})^{\beta} \right\} \\ &+ p_{5}^{\beta} \left\{ p_{1} \cdot (p_{5} - p_{2})(p_{1} + p_{2})^{\beta'} + (p_{5} - p_{2}) \cdot (p_{1} + p_{2})g^{\beta\beta'} - (p_{5} - p_{2})^{\beta'} (p_{1} + p_{2})^{\beta} \right\} \\ &- (m_{3}^{2} - p_{5} \cdot p_{2}) \left\{ p_{1}^{\beta} (p_{1} + p_{2})^{\beta'} + p_{1}^{\beta'} (p_{1} + p_{2})^{\beta} - (m_{1}^{2} + p_{1} \cdot p_{2})(p_{5} - p_{2})^{\beta'} \right\} \\ &- (m_{3}^{2} - p_{5}) \cdot (p_{1} + p_{2})p_{1}^{\beta} + (m_{1}^{2} + p_{1} \cdot p_{2})(p_{5} - p_{2})^{\beta} - p_{1} \cdot (p_{5} - p_{2})(p_{1} + p_{2})p_{1}^{\beta'} + (m_{1}^{2} + p_{1} \cdot p_{2})(p_{5} - p_{2})^{\beta} - p_{1} \cdot (p_{5} - p_{2})(p_{1} + p_{2})^{\beta} \right\} \\ &- m_{3}^{\beta'} \left\{ (p_{5} - p_{2}) \cdot (p_{1} + p_{2})p_{1}^{\beta} + (m_{1}^{2} + p_{1} \cdot p_{2})(p_{5} - p_{2})^{\beta} - p_{1} \cdot (p_{5} - p_{2})(p_{1} + p_{2})^{\beta'} \right\} \\ &- m_{3}^{\beta'} \left\{ p_{1}^{\beta} (p_{1} + p_{2})^{\beta'} + p_{5}^{\beta'} (p_{1} + p_{2})^{\beta} - p_{5} \cdot (p_{1} + p_{2})g^{\beta\beta'} \right\} + 2m_{5}m_{q}p_{1}^{\beta} (p_{1} + p_{2})^{\beta'} + 2m_{5}m_{q$$

$$+2m_{q}^{2}p_{1}^{\beta}p_{5}^{\beta'}+2m_{1}m_{q}(p_{5}-p_{2})^{\beta}p_{5}^{\beta'}-m_{5}m_{q}^{2}m_{1}g^{\beta\beta'})\times \frac{\{(p_{2}-p_{5})^{2}-m_{q}^{2}\}\{(p_{1}+p_{2})^{2}-m_{q}^{2}\}+4(p_{10}+p_{20})(p_{20}-p_{50})\Gamma_{q}^{2}}{\left(\{(p_{2}-p_{5})^{2}-m_{q}^{2}\}^{2}+4(p_{20}-p_{50})^{2}\Gamma_{q}^{2}\right)\left(\{(p_{1}+p_{2})^{2}-m_{q}^{2}\}^{2}+4(p_{10}+p_{20})^{2}\Gamma_{q}^{2}\right)}\right],$$
(A6)

where it is assumed that $\varepsilon_{\alpha}^{a*}(p_2)\varepsilon_{\alpha'}^{a'}(p_2) = -g_{\alpha\alpha'}\delta_{aa'}$ according to the Lorentz gauge, and

$$L^{\beta}L^{\beta'*} = 16\pi\alpha \frac{p_3^{\beta}p_4^{\beta'} + p_4^{\beta}p_3^{\beta'} - g^{\beta\beta'}(q^2/2)}{q^4}.$$
 (A7)

We note that the nonvanishing width of the quark spectral function removes divergences which appear in some kinetic regions. The first three lines in the bracket of Eq. (A6) are the squared *u* channel and the next three lines are the squared *s* channel and the rest the mixed term of *u* and *t* channels. We note that the squared *u* channel and the squared *s* channel are equivalent to each other, if p_1 and $-p_5$ and m_1 and m_5 are exchanged.

The invariant matrix element for the process $q + \bar{q} \rightarrow g + \gamma^*$ is calculated from Fig. 15 as

$$M = \bar{v}_d(p_2) \left[\frac{\delta_{db} i e_q \gamma^{\beta} i(\not p_1 - \not p_5 + m_q) i g \gamma^{\alpha}}{(p_1 - p_5)^2 - m_q^2 + i 2(p_{10} - p_{50})\Gamma_q} T_{bc}^a + T_{db}^a \frac{i g \gamma^{\alpha} i(-\not p_2 + \not p_5 + m_q) \delta_{bc} i e_q \gamma^{\beta}}{(-p_2 + p_5)^2 - m_q^2 + i 2(-p_{20} + p_{50})\Gamma_q} \right] u_c(p_1) \varepsilon_{\alpha}^a(p_5) L_{\beta}.$$
(A8)

Compared to the $|M|^2$ from Eq. (A4), the invariant matrix element squared for $q + \bar{q} \rightarrow g + \gamma^*$ is equivalent to that for $q + g \rightarrow q + \gamma^*$ with p_2 and $-p_5$, and m_2 and m_5 being exchanged and an additional overall minus sign. The integration over phase space for the next-to-leading-order (NLO) processes is given by

$$\int \frac{d^3 p_5}{(2\pi)^3 2E_5} \int \frac{d^3 p_3}{(2\pi)^3 2E_3} \int \frac{d^3 p_4}{(2\pi)^3 2E_4} (2\pi)^4 \delta^{(4)}(p_1 + p_2 - p_3 - p_4 - p_5).$$
(A9)

Introducing $q = p_3 + p_4$, the phase space is factorized as following:

$$\int \frac{d^3 p_5}{(2\pi)^3 2E_5} \int d^4 q \,\delta^{(4)}(p_1 + p_2 - q - p_5) \int \frac{d^3 p_3}{(2\pi)^3 2E_3} \int \frac{d^3 p_4}{(2\pi)^3 2E_4} (2\pi)^4 \delta^{(4)}(q - p_3 - p_4), \tag{A10}$$

where the second line is nothing but the phase space for a two-body decay, which affects only the dilepton part of Eq. (A7):

$$\int \frac{d^{3}p_{3}}{(2\pi)^{3}2E_{3}} \int \frac{d^{3}p_{4}}{(2\pi)^{3}2E_{4}} (2\pi)^{4} \delta^{(4)}(q-p_{3}-p_{4}) L^{\beta} L^{\beta'*}$$

$$= 16\pi\alpha \int \frac{d^{3}p_{3}}{(2\pi)^{3}2E_{3}} \int \frac{d^{3}p_{4}}{(2\pi)^{3}2E_{4}} (2\pi)^{4} \delta^{(4)}(q-p_{3}-p_{4}) \frac{p_{3}^{\beta}p_{4}^{\beta'}+p_{4}^{\beta}p_{3}^{\beta'}-g^{\beta\beta'}q^{2}/2}{q^{4}}$$

$$= 16\pi\alpha \int \frac{d^{3}p_{3}}{(2\pi)^{3}2E_{3}} \int \frac{d^{3}p_{4}}{(2\pi)^{3}2E_{4}} (2\pi)^{4} \delta^{(4)}(q-p_{3}-p_{4}) \frac{q^{\beta}q^{\beta'}-2p_{3}^{\beta}p_{3}^{\beta'}-g^{\beta\beta'}q^{2}/2}{q^{4}}$$

$$= 16\pi\alpha \int \frac{d^{3}p_{3}}{(2\pi)^{3}2E_{3}} \int \frac{d^{3}p_{4}}{(2\pi)^{3}2E_{4}} (2\pi)^{4} \delta^{(4)}(q-p_{3}-p_{4}) \frac{q^{\beta}q^{\beta'}-2E_{3}^{2}g^{\beta0}g^{\beta'0}-(2/3)|\vec{p}_{3}|^{2}g^{\beta i}g^{\beta' i}-g^{\beta\beta'}q^{2}/2}{q^{4}}, \quad (A11)$$

considering

$$\int d^3 p_3 p_3^i p_3^j = \frac{1}{3} \int d^3 p_3 |\vec{p}_3|^2 \delta^{ij}.$$
(A12)

In the q rest frame, Eq. (A11) reduces to

$$\frac{16\pi\alpha}{q^4} \left\{ q^\beta q^{\beta'} - \frac{q^2}{2} (g^{\beta\beta'} + g^{\beta0} g^{\beta'0}) - \frac{2|\vec{p_3}|^2}{3} g^{\beta i} g^{\beta' i} \right\} \frac{|\vec{p_3}|}{4\pi q} \equiv \overline{L^\beta L^{\beta'*}} \frac{|\vec{p_3}|}{4\pi q} = \overline{L^\beta L^{\beta'*}} \frac{1}{8\pi} \sqrt{1 - \frac{4m_l^2}{q^2}}, \tag{A13}$$

where $\overline{L^{\beta}L^{\beta'*}}$ is the lepton pair tensor averaged over phase space. The rest part of phase space, the first line in Eq. (A10), can be simplified as follows:

$$\int \frac{d^3 p_5}{(2\pi)^3 2E_5} \int d^4 q \delta^{(4)}(p_1 + p_2 - q - p_5) = \frac{1}{(2\pi)^2} \int \frac{d p_5 p_5^2 d \cos \theta}{2E_5} = \frac{1}{16\pi^2 \sqrt{s}} \int dq^2 d \cos \theta p_5$$
$$= \frac{1}{16\pi^2 \sqrt{s}} \int dq^2 d \cos \theta \sqrt{\frac{\left(s + m_5^2 - q^2\right)^2}{4s} - m_5^2}, \tag{A14}$$

with $p_1 + p_2 = p_5 + q$ and $dq^2 = d(p_1 + p_2 - p_5)^2 = -2\sqrt{s}dE_5$ in the center-of-mass frame.

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FIG. 15. Feynman diagrams for $q + \bar{q} \rightarrow g + \gamma^*$.

Combining Eqs. (A13) and (A14), the differential cross section is given by

$$\frac{d\sigma}{dq^2 d\cos\theta} = \frac{1}{8(4\pi)^3 p_i s} \sqrt{\frac{\left(s + m_5^2 - q^2\right)^2}{4s} - m_5^2} \\ \times \sqrt{1 - \frac{4m_l^2}{q^2}} \overline{|M|}^2, \tag{A15}$$

with $L^{\beta}L^{\beta'*}$ being substituted with $\overline{L^{\beta}L^{\beta'*}}$ and $\overline{|M|}^2 = |M|^2/96$ for $q + g \to q + \gamma^*$ and $\overline{|M|}^2 = |M|^2/36$ for $q + \bar{q} \to g + \gamma^*$ from spin + color degeneracies; p_i is the momentum of the initial particle in the center-of-mass frame:

$$p_i = \sqrt{\frac{\{s - (m_1 + m_2)^2\}\{s - (m_1 - m_2)^2\}}{4s}}.$$
 (A16)

The numerical calculations are carried out in the *q*-rest frame:

$$p_{1}^{\mu} = (E_{1}, 0, 0, |\vec{p}_{1}|),$$

$$p_{2}^{\mu} = (E_{2}, 0, |\vec{p}_{5}| \sin \psi, |\vec{p}_{5}| \cos \psi - p_{1}),$$

$$p_{5}^{\mu} = (E_{5}, 0, |\vec{p}_{5}| \sin \psi, |\vec{p}_{5}| \cos \psi),$$

$$p_{3}^{\mu} = (E_{3}, |\vec{p}_{3}| \sin \theta_{1} \sin \phi, |\vec{p}_{3}| \sin \theta_{1} \cos \phi, |\vec{p}_{3}| \cos \theta_{1}),$$

$$p_{4}^{\mu} = (E_{4}, -\vec{p}_{3}),$$
(A17)

where

$$\begin{aligned} |\vec{p}_3| &= \sqrt{\frac{q^2}{4} - m_l^2}, \\ E_3 &= E_4 = \sqrt{m_l^2 + p_3^2}, \\ E_5 &= \frac{s - q^2 - m_5^2}{2q^2} \text{ from } s = (p_5 + q)^2, \\ |\vec{p}_5| &= \sqrt{E_5 - m_5^2}, \\ E_2 &= \frac{m_2^2 + q^2 - t}{4E_3} \text{ from } t = (p_2 - q)^2, \\ E_1 &= \sqrt{s + |\vec{p}_5|^2} - E_2 \text{ from } s = (p_1 + p_2)^2 \\ \cos \psi &= \frac{m_2^2 + |\vec{p}_1|^2 + |\vec{p}_5|^2 - E_2^2}{2|\vec{p}_1||\vec{p}_5|}, \end{aligned}$$

from $E_2^2 - |\vec{p}_5|^2 \sin^2 \psi - (|\vec{p}_5| \cos \psi - p_1)^2 = m_2^2$. Independent variables are then *s*, *t*, θ_1 , and ϕ . Integrating over θ_1 and



FIG. 16. The invariant mass spectra of dielectrons from $D\bar{D}$ pairs in minimum-bias Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV without any acceptance cuts (dotted blue line), only with the p_T cut (dashed red line), only with the pseudorapidity cut $|\eta^e| < 1$ (dot-dashed orange line), and with both cuts included (solid green line). The full dots are the experimental data from the STAR Collaboration for the total dilepton spectra.

 ϕ in the q-rest frame, we get

$$\overline{L^{\beta}L^{\beta'*}} = 16\pi\alpha \frac{(q^2/2) - (2/3)|\vec{p}_3|^2}{q^4} g^{\beta i} g^{\beta' i}$$
$$= 16\pi\alpha \frac{q^2 - 2m_l^2}{3q^4} g^{\beta i} g^{\beta' i}$$
(A18)

from Eq. (A13).

APPENDIX B

In this Appendix, we study the effect of acceptance cuts on the invariant mass spectrum of dielectrons and the dependence of the dielectron spectrum on the total cross section for charm production (within the experimental uncertainties) by considering minimal-bias Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV (as an example).

Figure 16 shows the invariant mass spectra of dielectrons from $D\bar{D}$ pairs with several acceptance cuts incorporated. The dotted blue line is the dielectron spectrum without any acceptance cuts and is naturally much higher than the data. The dashed red line is the spectrum with the p_T cut for electrons as well as positrons ($|p_T^e| > 0.2$ GeV). This cut reduces the dielectron mass spectrum slightly more at low invariant mass than at large invariant masses and thus enhances the apparent slope for intermediate masses. This results from the fact that the electron and the positron with large invariant mass have large momenta such that the p_T^e cut is less effective. On the other hand, the pseudorapidity cut $(|\eta^e| < 1)$ reduces considerably the dielectron spectrum at large invariant mass for the same reason. If the momenta of electron and positron-composing a dielectron-are large due to the large invariant mass, they tend to lie outside the pseudorapidity cut. The solid green line, finally, is the dielectron spectrum after both cuts, which is essentially the same as the green line in Fig. 10(e).

According to the STAR measurements, the total cross section for charm production in p + p collisions at $\sqrt{s_{NN}} =$ 200 GeV is $797 \pm 210^{+208}_{-295} \ \mu b$ [48] and thus has a statistical and systematic error about a factor of two. Figure 17 shows the invariant mass spectrum of dielectrons with the charm cross section from the STAR Collaboration considering its statistical and systematic errors while including the cuts in p_T^e and η^{e} . Since the dielectrons from $D\bar{D}$ semileptonic decays are the most dominant contribution in the intermediate-mass range, the total dielectron spectrum is primarily sensitive to the charm cross section employed. The figure shows that the PHSD results with the mean value of charm cross section (from STAR) overestimate the dielectron data and the inclusion of both statistical and systematic errors is necessary to achieve an agreement with the experimental data. For our present study, we use the charm cross sections fitted to the experimental data within a wide range of collision energies as shown in Fig. 1(a), where the cross section at $\sqrt{s_{NN}} = 200 \text{ GeV}$ is about 400 μ b, which is still within the statistical and systematic error bars of the STAR Collaboration and which is close to the recent results from the PHENIX Collaboration [49].

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FIG. 17. The invariant mass spectra of dielectrons for minimumbias Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV with the total cross section for charm production from the STAR Collaboration with statistical and systematic errors displayed in terms of the shaded areas.

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