Centrality-dependent Direct Photon $p_t$ spectra in Au+Au Collisions at RHIC

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We calculate the centrality-dependence of transverse momentum ($p_t$) spectra for direct photons in Au+Au collisions at the RHIC energy, based on a realistic data-constrained (3+1) dimensional hydrodynamic description of the expanding hot and dense matter, a reasonable treatment of the propagation of partons and their energy loss in the fluid, and a systematic study of the main sources of direct photons. The resultant $p_t$ spectra agree with recent PHENIX data in a broad $p_t$ range. The competition among the different direct photon sources is investigated at various centralities. Parton energy loss in the plasma is considered for photons from fragmentation and jet photon conversion, which causes about 40% decrease in the total contribution. In the high $p_t$ region, the observed $R_{AA}$ of photons is centrality independent at the accuracy of 5% based on a realistic treatment of energy loss. We also link the different behavior of $R_{AA}$ for central and peripheral collisions, in the low $p_t$ region, to the fact that the plasma in central collisions is hotter compared to peripheral ones.

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I. INTRODUCTION

The formation and observation of a quark-gluon plasma (QGP) in heavy ion collisions are important goals of modern nuclear physics [1,2]. Suppression of high $p_t$ hadron yields [3] is one of the most important features observed at the Relativistic Heavy Ion Collider (RHIC). Theoretically this is attributed to the interaction between jets (hard partons) and the bulk matter [4,5,6,7]. Experimentally, absence of the suppression in $d+Au$ collisions [3] reveals that the suppression results from a final-state effect and, in turn, that the hot and dense matter is created in Au+Au collisions. The amount of suppression depends significantly on the centrality of the collision [3], which implies various sizes of hot dense matter are formed in heavy ion collisions at various centralities. This offers us an excellent opportunity to study the interaction of partons inside the system and, consequently, properties of the matter under extreme conditions.

Hadron production in heavy ion collisions involves bulk hadronization of the thermal partons at low $p_t$, the fragmentation of quenched hard partons at high $p_t$ and the hadronization contributed from both thermal and hard partons at intermediate $p_t$. However, it is quite difficult to systematically describe all these hadronization processes since some of them are beyond the perturbative treatment and usually contain many parameters without full understanding. Low $p_t$ hadrons also strongly interact with each other after hadronization and cannot carry direct information from inside the hot matter. Under this situation, a systematic study of direct photons in a wide range of transverse momentum and centrality can serve as a guide to understand the whole reaction processes of heavy ion collisions, since we do not need to treat hadronization itself nor interaction between produced direct photons and the bulk matter thanks to the large mean free path of direct photons compared to the typical size of the system in heavy ion collisions. Competition among different sources at various centralities may be also useful in understanding of production mechanism of direct photons.

In this paper, we first study the role of jet quenching on the centrality dependence of direct photon production. For this purpose, a reliable treatment of hard parton energy loss is needed. This is formulated via the BDMPS framework [8] and tested on pion suppression at various centralities. Since neutral pions and other mesons are significantly suppressed in central Au+Au collisions, and since the suppression has an evident centrality dependence, the following question arises naturally: What is the role of hard parton energy loss on direct photon production? Main purpose of this paper is to answer this question.

We also investigate the interplay among the various sources of direct photons. Similar to hadron suppression, photons from parton fragmentation are expected to offer information on the interaction between hard partons and the bulk via jet quenching. Thermal photons and photons from parton-bulk interactions are penetrating probes of the hot matter, respectively, through interaction of partons inside bulk matter and interaction between primary partons and the bulk matter. It is interesting to see whether one reproduces the observed photon spectra by considering all photon sources simultaneously and consistently at different collision centralities. By identifying the dominant sources of direct photons at given values of $p_t$, we will be able to discuss in which way the different $p_t$-regions of the photon spectra provide information about the different production processes.

We need a realistic description of the hot and dense
matter to investigate the effect of bulk matter on photon emission. This is achieved by using three dimensional (3D) hydrodynamic simulations of bulk matter which have already been tested against a vast body of low $p_t$ hadron data at RHIC.

To perform a systematic study of direct photon production (from sources other than neutral meson decays) in relativistic heavy ion collisions, we shortly review the possible sources in the following.

**Primordial NN scattering.** The direct photon production via Compton scattering and quark-antiquark annihilation can be calculated in perturbation theory using the conventional parton distribution functions and the factorization hypothesis. In principle one should consider at this stage also higher order contributions, like bremsstrahlung of photons accompanying for example two-jet production in hard parton-parton scattering. However, we consider this component as a part of the so-called jet fragmentation (or bremsstrahlung) contribution to be affected by the thermalized matter, which we will discuss separately.

**Thermal photons.** In high energy nuclear collisions, the density of secondary partons is so high that the quarks and gluons re-scatter and eventually thermalize to form a bubble of hot QGP. The plasma expands, cools down, and goes through a phase transition to hadronic gas (HG) phase. Thermal photons can be produced during the whole history of the evolution of hot matter from the QGP phase to the HG phase through the mixed phase due to collisions of or radiations from thermalized particles. Yields of photons from a thermal source are exponentially damped so that contribution to very high $p_t$ region is negligible. However, contribution to low $p_t$ is expected to be dominant in central collisions in which the size and the temperature of a hot and dense matter are large enough.

**Jet-photon conversion.** When hard partons pass through thermalized matter, they may interact. Collisions between jets and deconfined partons via quark-antiquark annihilation and quark-gluon Compton scattering can produce direct photons. This is often called jet-photon conversion.

**Jet fragmentation.** Photon production also occurs as a higher order effect in purely partonic initial hard scatterings: at any stages of the evolution of a jet (final state parton emission), there is a possibility of emitting photons. Existence of a QGP again affects the results of fragmented photons since energetic partons lose their energy prior to fragmentation. In this work, we assume fragmentation of partons only outside the plasma, which is similar to high $p_t$ hadron production from jet fragmentation.

There are possible contributions to photon production which are not included in the present study: The medium-induced radiation is supposed to be a higher order contribution. At the RHIC energy this contributes much less than jet photon conversion at low and intermediate $p_t$ and much less than fragmentation at high $p_t$. So the contribution from medium-induced radiation is ignored this paper. In the time interval between the primordial collisions at $\tau = 0$ and the thermalization of the hot matter at $\tau_0$, the interaction between non-equilibrated soft partons and hard partons may also produce direct photons. We neglect the contribution in the preequilibrium stage since the time interval is much shorter than the life time of the equilibrated matter ($\sim 20$ fm/$c$).

The paper is organized as follows: In Sec. [II] we first give a brief review on the space-time evolution of the hot matter created in Au+Au collisions at different centralities based on a (3+1)-dimensional ideal hydrodynamical calculation. In Sec. [III] we discuss parton energy loss in the QGP. We investigate neutral pion production in the high $p_t$ region in order to fix the parameters of the energy loss scheme. We discuss sequentially the contributions from various sources to direct photon $p_t$ spectra in Sec. [IV]. We show our results and compare them with recent experimental data in Sec. [V]. Section [VI] is devoted to conclusion of the present study.

## II. SPACE-TIME EVOLUTION OF THE HOT AND DENSE MATTER

Several sources of direct photon production in heavy ion collisions depend on the bulk dynamics of hot and/or dense matter and the matter along trajectories of energetic partons. So a realistic description of reaction dynamics is indispensable for the quantitative analysis of photon production. In our calculation, fully three-dimensional (3D) ideal hydrodynamics is employed to describe the space-time evolution of the hot and dense matter created in Au+Au collisions at RHIC energy at various centralities. We solve the equations of energy-momentum conservation

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

in full 3D space ($\tau, x, y, \eta$) under the assumption that the local thermal equilibrium is reached (maintained) at (after) an initial time $\tau_0 = 0.6$ fm/$c$. Here $\tau$ and $\eta$ are the proper time and the space-time rapidity, respectively. $x$ and $y$ are transverse coordinates. In the transverse plane, the centers of two colliding nuclei are located at $(x, y) = (b/2, 0)$ and $(-b/2, 0)$ before the collision at an impact parameter $b$. Ideal hydrodynamics is characterized by the energy-momentum tensor

$$T^{\mu \nu} = (e + P)u^\mu u^\nu - P g^{\mu \nu},$$

where $e$, $P$, and $u^\mu$ are energy density, pressure, and local four velocity, respectively. We neglect the finite net-baryon density which is small near the mid-rapidity at RHIC. For the high temperature ($T > T_c = 170$ MeV) QGP phase we use the equation of state (EOS) of massless non-interacting parton gas ($u, d, s$ quarks and glu-
The bag constant is tuned to be $B^\frac{4}{3} = 247.19$ MeV to match pressure of the QGP phase to that of a hadron resonance gas at critical temperature $T_c = 170$ MeV. A hadron resonance gas model at $T < T_c$ includes all hadrons up to the mass of the $\Delta(1232)$ resonance. Our hadron resonance gas EOS implements chemical freeze-out at RHIC [13]. This is achieved by introducing appropriate temperature-dependent chemical potentials for all hadronic species $i$ in a way that their numbers $\tilde{N}_i$ including all decay contributions from higher-lying resonances, $\tilde{N}_i = N_i + \sum_b b_{R\rightarrow iX} N_{R^b}$, are conserved during the evolution [11, 14, 15, 16, 17, 18]. Here $N_i$ is the average multiplicity of the $i$-th hadron species, and $b_{R\rightarrow iX}$ is the effective branching ratio (a product of branching ratio and degeneracy) of a decay process $R \rightarrow i + X$. In this partial chemical equilibrium (PCE) model [11] only strongly interacting resonances with large decay widths (whose decays do not alter $\tilde{N}_i$) remain chemically equilibrated below $T_{ch}$. It should be noted that the hadronic chemical composition described by hydrodynamics using the PCE model is roughly consistent with that of the hadronic cascade models [19], as long as the latter are initialized at $T_{sw} = 169$ MeV with thermal and chemical equilibrium distributions.

We assume that, at $\tau_0 = 0.6$ fm/c, the initial entropy distributions is proportional to a linear combination of the number density of participants (85%) and binary collisions (15%) [10]. Centrality dependence of charged particle multiplicity observed by PHOBOS [20] has been well reproduced by full 3D hydrodynamics simulations with the above setups [19]. In the following calculations, hydrodynamic outputs at representative impact parameters $b = 3.2, 5.5, 7.2, 8.5, 9.7,$ and $10.8$ fm are chosen for 0-10%, 10-20%, ..., and 50-60% centrality, respectively.

So far, the space-time evolution of the QGP fluid obtained as above has been also exploited for a quantitative study of hard and rare probes such as azimuthal jet anisotropy, nuclear modification factor of identified hadrons, disappearance of back-to-back jet correlation, and $J/\psi$ suppression [21].

In Table I, initial temperatures at the plasma center, $T_0 = T(\tau_0, 0, 0, 0)$, are shown for various centralities. These temperature values will be important to interpret the centrality dependence of the slope of pt spectra from thermal radiation, which will be discussed later. Figure 1 shows the time evolution of energy density at the center of fluids $(x, y, \eta) = (0, 0, 0)$ for various centralities. Clearly for any given proper time $\tau$, the more central collisions one obtains higher energy densities at the plasma center.

For convenience of the following calculations, we introduce $f_{QGP}(\tau, x, y, \eta)$ as the fraction of the QGP phase in a fluid element. It is obvious that $f_{QGP} = 1$ (0) in the QGP (hadronic) phase. In the mixed phase, the fraction of the QGP is calculated via

$$f_{QGP} e_{QGP} + (1 - f_{QGP}) e_{had} = e(\tau, x, y, \eta)$$

with $e_{QGP}$ and $e_{had}$ being the energy densities of the QGP phase and the hadron phase at $T = T_c$, respectively.

### III. Parton Energy Loss in a Plasma

Energy loss of hard partons in a plasma affects both jet photon conversion and jet fragmentation. The momentum distribution of jets (energetic gluons or quarks with different flavors) from primordial nucleus-nucleus scattering is calculated as [22]

$$\frac{dN^{AB\rightarrow jet}}{dy d^2p_t} = KT_{AB}(b) \sum_{abcd} \int dx_a dx_b G_{a/A}(x_a, M^2)$$

$$\times G_{b/B}(x_b, M^2) \frac{\hat{s} d\sigma}{\hat{s}} (ab \rightarrow cd) \delta(\hat{s} + \hat{t} + \hat{u})$$

where $T_{AB}(b)$ is the nuclear overlapping function at an impact parameter $b$ for each centrality, $G_{a/A}(x_a, M^2)$ and $G_{b/B}(x_b, M^2)$ are parton distribution functions in nuclei $A$ and $B$. We take MRST 2001 LO parton distributions in proton [23]. The elementary cross sections for $ab \rightarrow cd$ can be found in Ref. [22]. We set the factorization scale...
M and renormalization scale \( Q \) to be \( M = Q = p_t \). \( K = 2 \) is chosen to take into account higher order contributions. These parameters are chosen as to reproduce high \( p_t \) pion data in \( pp \) collisions at RHIC, which will be discussed later. The above formula for \( p_8 \) spectra was extensively tested in \( pp \) (\( pp \)) collisions in an energy range from \( \sqrt{s} = 27.4 \) GeV to 630 GeV. Nuclear shadowing effect and EMC effect are taken into account through EKS98 scale dependent nuclear ratios \( R_{a/P}^{EKS}(x, A) \) [24]. Isospin of a nucleus with mass \( A \), neutron number \( N \), and proton number \( Z \) is corrected as follows:

\[
G_{a/A}(x) = \left[ \frac{N}{A} G_{a/N}(x) + \frac{Z}{A} G_{a/P}(x) \right] R_{a/P}^{EKS}(x, A). \tag{5}
\]

The isospin mixture and nuclear shadowing eventually cause a decrease of nuclear modification at high \( p_t \) region, which will be shown in Sec. V.

We assume that all jets are produced at \( \tau = 1/Q \approx 0 \) with the phase space distribution

\[
f_0(\vec{p}, \vec{r}) \propto \frac{dN}{dp^2} T_A \left( x - \frac{b}{2} y \right) T_B \left( x + \frac{b}{2} y \right) \delta(z) \tag{6}
\]

where \( \vec{r} = (x, y, z) \) is the coordinate of a jet, \( b \) is the impact parameter, and \( T_A \) and \( T_B \) are thickness functions of nuclei \( A \) and \( B \). The \( \delta \)-function reflects the highly Lorentz-contracted colliding nuclei \( A \) and \( B \). The phase space distribution of hard partons is normalized as

\[
\int f_0(\vec{p}, \vec{r})d^3r = (2\pi)^3 \frac{dN}{dp^2} \tag{7}
\]

Energetic partons can suffer interactions with the fluid and lose their energies. We employ the BDMS formula [3] to calculate parton energy loss in a plasma created in heavy ion collisions. For a parton of type \( i = g, q \) with initial momentum \( \vec{p}_0 \) formed at \( \vec{r}_0 \), the whole path length of a parton traversing the expanding QGP (including the mixed phase) is

\[
L(\vec{p}_0, \vec{r}_0) = \int_{\tau_0}^{\infty} d\tau \theta(\mathcal{F}_{QGP}(\tau, \vec{x}(\tau))). \tag{8}
\]

Here \( \vec{x}(\tau) \) is a trajectory of a parton, \( \mathcal{F}_{QGP}(\tau, \vec{x}(\tau)) \) is the fraction of the QGP phase at a position \( (\tau, \vec{x}(\tau)) \), and \( \theta \) is a step function, which gives \( \theta(\mathcal{F}_{QGP}) \) equal unity in the QGP and the mixed phases and zero in the hadron phase.

The total energy loss along this path is calculated as

\[
\Delta E(i, \vec{p}_0, \vec{r}_0) = D \int_{\tau_0}^{\infty} d\tau \epsilon(i, \tau, \vec{x}(\tau)) \theta(\mathcal{F}_{QGP}(\tau, \vec{x}(\tau))). \tag{9}
\]

Here \( D \) is an adjustable parameter, \( \epsilon(i, \tau, \vec{x}(\tau)) \) is the energy loss per unit distance for a parton \( i \) at a position \( (\tau, \vec{x}(\tau)) \), given as [3]

\[
\epsilon(i, \tau, \vec{x}(\tau)) = \alpha_s \sqrt{\mu^2 E^*/\lambda_i}.
\]

Here, the temperature-dependent running coupling constant – assuming a similar formula as the lowest order one in perturbation theory – can be obtained by fitting the numerical results from lattice quantum chromodynamics (QCD) simulations [25] as

\[
\alpha_s(T) = \frac{6\pi}{(33 - 2N_f) \ln(8T/T_c)}. \tag{10}
\]

The Debye screening mass is given as \( \mu = gT \), with \( g^2/4\pi = \alpha_s(T) \). The energy of a hard parton in the local rest frame is \( E^* = p^* \mu \), where \( p^* \) is the four momentum of the hard parton in the laboratory frame and \( \mu \) is a local fluid velocity. All hard partons are treated as on-shell massless particles. The mean free path, \( \lambda_i \), of a parton \( i \), is given as

\[
\lambda_i^{-1} = \sigma_{gq} f_{QGP} + \sigma_{qq} f_{QGP}, \tag{11}
\]

\[
\lambda_q^{-1} = \sigma_{qg} f_{QGP} + \sigma_{qq} f_{QGP}, \tag{12}
\]

with cross sections \( \sigma_i = C_i \alpha_s \pi/T^2 \) [26]. The color factors \( C_i \) are 4/9, 1, and 9/4 for \( qq \), \( gg \), and \( gg \) scattering, respectively. The parton densities \( \rho_g \) and \( \rho_q \) are obtained from the EOS of the massless relativistic gas. The fraction of the QGP phase \( f_{QGP} \) is considered in the mixed phase to ensure a smooth transition from the QGP phase to the HG phase. Note that the above quantities, i.e., temperature \( T \), fluid velocity \( u_\mu \), parton densities \( \rho_i \), and, in turn, mean free path \( \lambda_i \), depend on the location of the parton \( \vec{x}(\tau) \) and can be obtained from full 3D hydrodynamics simulations discussed in the previous section.

Various sizes of the plasma are formed in heavy ion collisions at different centralities. We use the common energy loss formula Eq. (9) for all of these media. The main purpose in the present paper is a systematic study of direct photon production rather than a detailed analysis of parton energy loss. So we admit ourselves to introduce an adjustable parameter \( D \) in Eq. (9) to fit simultaneously neutral \( \pi \)-meson data in Au+Au collisions at different centralities [9].

We first discuss pion production in proton-proton collisions. We calculate neutral \( \pi \)-meson production assuming pQCD factorization, Eq. (11).

\[
\frac{dN^{\pi^0}_{pp}}{dy dp_t^2} = \sum_{\epsilon = g, q} \int d\vec{z} \frac{dN^{pp-c}_{\epsilon}}{dy dp_t^2} \frac{1}{z^2} D^0_{\epsilon/c}(z_c, Q^2), \tag{13}
\]

where \( D^0_{\epsilon/c}(z_c, Q^2) \) is \( \pi^0 \) fragmentation functions parameterized by Kniehl et al. [27]. In Fig. 2 \( p_t \) spectra for neutral pions in \( pp \) collisions at \( \sqrt{s} = 200 \) GeV calculated with \( M = Q = 2p_t \), \( p_t \), and \( p_t/2 \) are compared to PHENIX data [28]. In the high \( p_t \) region where the pQCD is expected to work, we reasonably reproduce the experimental data with the above setup with \( M = 2 = 2 \) and \( M = Q = p_t \). We use the \( p_t \) spectrum as a reference spectrum in the following calculations.

The effect of parton energy loss is taken into account through the medium modified fragmentation function
leading order contribution to direct photon production

Figure 2: (Color Online) Neutral pion production in pp collisions at $\sqrt{s}=200$ GeV is compared to PHENIX data. Three lines from top to bottom correspond to $Q = p_t/2$, $p_t$, and $2p_t$ respectively.

Figure 3 shows the nuclear modification factors for neutral pions in Au+Au collision at $\sqrt{s_{NN}}=200$ GeV for different centralities. Solid lines are calculated with the BDMPS energy loss formula (9) with $\lambda_c = 0.5$, $s=200$ GeV is compared to PHENIX data [28]. With a common value of the parameter $D=1.5$, our results undershoot the experimental data due to absence of neutral pion production from bulk components in this calculation. Notice that low $p_t$ pion data have already been described well [19] by using hydrodynamic simulations employed in the present study. In the following photon calculations, we always use the BDMPS energy loss formula (9) with $D=1.5$.

IV. THE DIFFERENT SOURCES OF DIRECT PHOTON PRODUCTION

Leading order contribution. Similar to Eq. (16), the leading order contribution to direct photon production in nucleus-nucleus collisions reads

$$\frac{dN_{AB}^{AB-\gamma}}{dyd^2p_t} = T_{AB}(b) \sum_{ab} \int dx_a dx_b G_{a/A}(x_a, M^2)$$

$$\times G_{b/B}(x_b, M^2) \frac{s}{\pi} \frac{d\sigma}{dt}(ab \rightarrow \gamma + X) \delta(s + t + u) \delta \tilde{s}$$ (16)

where the elementary processes $ab \rightarrow \gamma + X$ are Compton scattering $qg \rightarrow \gamma q$ and annihilation $q\bar{q} \rightarrow \gamma\gamma$.

Fragmentation contribution. Higher order contributions in pp collisions are due to jet fragmentation. We
can calculate them as
\[ \frac{dN_{\text{frag}}^{PP}}{dyd^2p_t} = \sum_{c=g,q_i} \int dz_c \frac{dN_{\text{PP}}^{c}}{dyd^2p_t^c} \frac{1}{z_c^2} D_{\gamma/c}(z_c, Q^2), \]
(17)

with photon fragmentation functions \( D_{\gamma/c}(z_c, Q^2) \) being the probability for obtaining a photon from a parton \( c \) which carries a fraction \( z \) of the parton’s momentum. So \( z'_{c} = z_{c} p_{c} / p_{c} \) is the transverse momentum carried by the parton \( c \) before fragmentation and \( d^{3}p / E = z_c^2 d^3p / E^' \). The effective fragmentation functions for obtaining photons from partons can be calculated perturbatively. We use the parameterized solutions by Owens \[22\].

In case of heavy ion collisions, parton energy loss in a plasma should be taken into account. This can be done via modified fragmentation functions \[28\]
\[ D_{\gamma/c}(z_c, Q^2, \Delta E_c) = \left( 1 - e^{-\frac{\Delta E_c}{T}} \right) \left[ \frac{z'}{z_c} D_{\gamma/c}(z_c', Q^2) + \frac{L}{\lambda_c} \frac{z'}{z_c} D_{\gamma/g}(z_c', g^2, Q^2) \right] + e^{-\frac{\Delta E_c}{T}} D_{\gamma/c}(z_c, Q^2), \]
(18)

with \( z'_{c} = z_{c} (p_{c} - \Delta E_c) / \lambda_{c} \) and \( \phi'_{c} = (L / \lambda_{c}) p_{c} / \Delta E_c \) being the rescaled momentum fractions carried by the parton \( c \) and the emitted gluons before fragmentation. \( \lambda_{c} \) is mean free path of the parton \( c \) in the plasma, \( L \) is the path length of each parton traversing the plasma defined in Eq. (5). Thus, in heavy ion collisions, contributions from fragmentation become
\[ \frac{dN_{\text{frag}}^{AB}}{dyd^2p_t} = \sum_{c=g,q_i} \int dz_c \frac{dN_{\text{AB}}^{c}}{dyd^2p_t^c} \frac{1}{z_c^2} D_{\gamma/c}(z_c, Q^2, \Delta E_c). \]
(19)

The above formula counts only fragmented photons outside the plasma. In principle, when fragmentation into photons happens inside the plasma photon can escape the plasma due to the long mean free path. However it is not evident when and where fragmentation happens.

**Thermal production.** The emission rate of photons is \( \Gamma = E d^{3}R / d^{3}p \), where \( R \) is the number of photons emitted from a medium per unit space-time volume with temperature \( T \). Total yields of thermal photons can be obtained by summing the emission rate over the space-time volume as
\[ \frac{dN_{\text{thermal}}}{dyd^2p_t} = \int d^{4}x \Gamma(E^*, T) \]
(20)

with \( d^{4}x = \tau d\tau d^{3}x dy \) and \( E^* = p^\mu u_{\mu} \), being the photon energy in the local rest frame. Here, \( p^\mu = (p_{t} \cos y, p_{t} \sin y, p_{t} \sin \phi, p_{t} \cos \phi) \) is the photon’s four momentum in the laboratory frame and \( u_{\mu} \) is a local fluid velocity. In our calculations, the thermal photon emission rate covers both contributions from the QGP phase and the hadronic phase
\[ \Gamma(E^*, T) = f_{\text{QGP}} \Gamma_{\text{QGP}}^{\gamma \to \gamma}(E^*, T) + (1 - f_{\text{QGP}}) \Gamma_{\text{HG}}^{\gamma \to \gamma}(E^*, T), \]
(21)

where \( f_{\text{QGP}} \) and \( T \) are the fraction of the QGP phase and temperature of the fluid, respectively, both being obtained in the hydrodynamic simulations. In the above formula, we calculate thermal photon production above the thermal freeze-out temperature \( T_{\text{th}} = 100 \) MeV. The photon emission rate from 2 \to 2 processes between thermal partons, i.e., the QCD Compton process \( qg \to \gamma g \) and annihilation \( q\bar{q} \to g\gamma \), was first calculated with the hard thermal loop resummation technique \[34\] [31]. Later, Landau-Pomeranchuk-Migdal (LPM) interference effect for emitted photons turned out to be important \[32\], leading to
\[ \Gamma_{\text{QGP}}^{\gamma \to \gamma}(E^*, T) = \sum_{i=1}^{N_{f}} \left( \frac{\alpha}{e} \right)^{2} \frac{2 \alpha_{s} T^{2}}{2 \pi^{2}} \frac{1}{e^{\gamma}} + 1 \]
\times \left[ \ln \left( \frac{T}{\sqrt{2}} \right) + \frac{1}{2} \ln(2x) + C_{22}(x) \right] + C_{\text{brems}}(x) + C_{\text{ann}}(x), \]
(22)

with \( x = E^* / T \) and
\[ C_{22}(x) = \frac{0.041}{x} - 0.3615 + 1.01 e^{-1.35x}, \]
(23)

\[ C_{\text{brems}}(x) + C_{\text{ann}}(x) = 0.633 x^{-1.5} \ln (12.28 + 1 / x) + \frac{0.154 x}{(1 + x / 16.27)^{0.5}}. \]

In the calculation we take \( N_{f} = 3 \) and a temperature dependent running coupling as in Eq. (10).

Thermal photon emission in the hadronic phase results from interactions such as \( \pi\pi \to \gamma\gamma, \pi\rho \to \pi\gamma \) and \( \rho \to \pi\pi\gamma \), etc. Interactions of mesons or baryons with strangeness can also produce photons, but the contributions are relatively small due to the phase space suppression resulting from their heavier masses. In our work, photon emission rate from the hadronic phase is based on massive Yang-Mills (MYM) calculation \[33\], where photon production from mesons with strangeness has been included as well as the axial meson \( a_{1} \) as an exchanging particle for non-strange initial states. Hadrons are compositive objects, so form factors are considered to simulate finite hadronic size effects \[34\].

**Jet photon conversion with jet energy loss.** When hard partons propagate in a plasma, they also collide with thermal partons and produce direct photons via Compton process and the quark-antiquark annihilation. We call this process jet-photon conversion, since it is a conversion of a jet into a photon with almost the same momentum as the one of originating jet. Contribution from the jet-photon conversion is calculated by integration of conversion rate over the space-time evolution of the hot and dense matter in the QGP phase
\[ \frac{dN_{\text{jpc}}}{dyd^2p_t} = \int \Gamma_{\text{jpc}}(E^*, T) f_{\text{QGP}}(x, y, \eta, \tau) d^{4}x. \]
(24)
The photon production rate by annihilation and Compton scattering of hard partons in the medium can be approximated as

$$\Gamma^{\text{pp}}(E^*, T) = \frac{\alpha s}{4\pi^2} \sum_q e_q^2 f_q(x) T^2 \ln \left[ \frac{4E^*T}{m_{th}^2} - C \right]$$

where $E^*$ is the photon energy in the local rest frame, $C = 2,323, m_{th}^2 = g^2 T^2/6$, and the strong coupling $\alpha_s = g^2/4\pi$ being temperature dependent as in Eq. (10). $\alpha$ is the electromagnetic couplings, $e_q$ and $f_q(\bar{p}, x)$ are the electric charge and the phase-space density of a hard parton of flavor $q$. The phase space distribution of hard partons at $\tau$ is obtained by considering parton energy loss as

$$f(\bar{p}, x) = f(\bar{p}, \bar{r}, \tau) = \int d^2 p_0 f_0(\bar{p}_0, \bar{r} - \bar{v} t) \delta(\bar{p}_0 - \bar{p} - \bar{v} \Delta E)$$

where $f_0(\bar{p}, \bar{r})$ is the phase space distribution at $\tau = 0$ described in Eqs. (6) and (7). The $\delta$-function expression reflects the energy loss of a parton moving along a straight line trajectory with $\bar{v} \equiv \bar{p}/E = \bar{p}_0/E_0$. $\Delta E$ is the energy loss from $\tau_0$ to $\tau$ and calculated similar to Eq. (6) but replacing the upper limit of integral $\infty$ with $\tau$.

**V. RESULTS AND DISCUSSION**

In Fig. 4, the calculated $p_t$ spectra of direct photons in Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV at centrality 0-20% and 20-40% are compared to PHENIX data [37, 38]. Here we sum over all contributions discussed in the previous section. The theoretical results for 0-20% centrality are obtained as a mixture of the calculations for 0-10% and 10-20% centrality with a weight of 50% each; a corresponding procedure applies for the 20-40% centrality results. The PHENIX data are reproduced within our multi-component model remarkably well.

In Fig. 5 we show a detailed comparison of the calculated $p_t$ spectra of direct photons with PHENIX data [39] for the centralities 0-10%, 10-20%, ..., 50-60%. Again, our results agree with data very well in a broad range of $p_t$ and centrality.

Since all the curves are almost parallel to each other, one gets more insight by using the nuclear modification factor $R_{AA}$, obtained by dividing a $p_t$ spectrum in nucleus-nucleus collisions by the $N_{coll}$-scaled $p_t$ spectrum in $pp$ collisions. In Fig. 5, we show the invariant differential cross section of direct photons in $pp$ collisions. The calculation includes the leading order contribution plus fragmentation contribution, using a scale $Q = p_t$. PHENIX data are shown as open circles [37] and filled circles [38]. In high $p_t$ regions, our result agrees with the data reasonably well: So we use it to calculate nuclear modification in Fig. 7 and Fig. 10. It also provides a baseline calculation with the LO contribution and fragmentation contribution in Au+Au collisions. Whereas, in low $p_t$ regions where pQCD is not expected to work, our result undershoots the data slightly although the error bars are large in data. The dashed line is a fit to the measured differential cross section of direct photons in $pp$ collisions at the RHIC energy

$$\frac{d\sigma}{dyd^2p_t} = 0.01834 \left( 1 + \frac{p_t^2}{1.432} \right)^{-3.27} \text{mb/GeV}^2,$$

which is employed to calculate the nuclear modification factor from thermal contribution in Fig. 7(a).

Figure 7 shows how the nuclear modification factor for direct photons, $R_{AA}$, depends on centrality and on energy loss. Data for 0-10% centrality are taken from Refs. 37 and 40. Figure 7(a) shows centrality dependence of $R_{AA}$ compared to the PHENIX data. The three curves are respectively 0-10% (dotted line), 20-30% (solid line), and 40-50% (dashed-dotted line). $R_{AA}$ has a weak centrality dependence at high $p_t$ region. This result is consistent with the observed phenomenon [39] that the $p_t$-integrated (for $p_t > 6$ GeV/c) $R_{AA}$ of direct photons is almost independent of collision centrality. Does this imply a very weak effect from jet quenching? Figure 7(b) answers this question (here for the most central collisions): Comparing calculations with (dotted line) and without energy loss (solid line), one finds a difference of up to 40%. So the effect of parton energy loss is quite visible in the $p_t$ range between 4 GeV/c and more than 20 GeV/c. If we would do the $R_{AA}$ calculations without energy loss, the difference between central and semiperipheral collisions would be about 20%, whereas the complete calculation gives the same result for all centralities, within 5%. We have to admit that we talk about small effects, requiring experimental data with relative errors of less than 5% to observe the effects.

To understand the above results, we look more closely into the different contributions. Parton energy loss in the plasma suppresses the fragmentation contributions and jet-photon conversion. So we study the ratios of the contribution with energy loss to the one without energy loss, as shown in Fig. 8. (a) for fragmentation and (b) for jet-photon conversion. Energy loss in the plasma depends on the path length of the hard parton inside the plasma, which turns out to depend on the collision centrality. We do see a similar centrality dependence of the suppression for $\pi^0$ (jet quenching effect) in fragmentation contributions and jet-photon conversion.

To understand how these energy loss features affect the total contribution, we investigate the competition from different sources in Fig. 8 for the three centralities 0-10%, 20-30%, and 40-50%. The leading order contribution (LO) from primordial elementary scatterings is plotted as dotted lines, thermal contribution as dash-dotted lines, fragmentation contribution as dashed lines, and jet photon conversion as solid lines. The latter two are calculated with parton energy loss in the plasma (left plots) and without (right). For all centralities, thermal
photons dominate at low transverse momenta and they are insignificant in the high \( p_t \) region. The leading order contribution from primordial elementary scatterings dominates in the high \( p_t \) region. This contribution is independent of bulk volume.

Let us first discuss the central collisions. Here fragmentation and conversion are an order of magnitude smaller than the LO contribution. But from Fig. 9 we know that is due to the strong energy loss effect. Without this energy loss, these two contributions would be much bigger, and this is why we find a 40% difference between the total contribution with and without energy loss.

For peripheral collisions, without energy loss, the relative contribution from conversion is smaller compared to central scatterings, since the the plasma regions is smaller. But then also the suppression from energy loss is smaller for the peripheral compared to central collisions. So at the end, including a proper energy loss treatment, for both central and peripheral collisions, conversion is roughly an order of magnitude smaller than the LO contribution, see Fig. 9(e). The relative contribution from fragmentation, without energy loss, is comparable in central and peripheral collisions, however, the energy loss is smaller in the latter ones. So fragmentation is somewhat more important in peripheral compared to central, when energy loss is considered, as can be also seen from Fig. 10 where the contribution to \( R_{AA} \) from fragmentation and conversion is shown for the different centralities: Conver-
sion contributes roughly 4%, for all centralities, fragmentation between 5% (central) and 10% (peripheral).

At the end, \( R_{AA} \) is nearly centrality independent as shown in Fig. 7(a), but a realistic (and strong) partonic energy loss is needed in order to get this scaling behavior.

In the low \( p_t \) region, contribution from thermal radiation is of significant importance. We check the centrality-dependence of the thermal contribution to \( R_{AA} \) in Fig. 11(a). At very low \( p_t \), i.e., \( p_t < 1 \text{ GeV}/c \), the thermal contributions to \( R_{AA} \) at different centralities coincide with each other. However the slope of \( R_{AA}^{\text{thermal}} \) changes and the dominant \( p_t \) region of thermal photons becomes smaller when one moves from central to peripheral collisions. This reflects the fact that the temperature in the core region depends on the collision centrality as shown in Table I in Sec. III.

So from the thermal source, the \( R_{AA} \) for central collisions exceed more and more the \( R_{AA} \) for peripheral collisions, which translates into a slight overshooting of the central total \( R_{AA} \) compared to the peripheral one, as seen in Fig. 5 in the region \( p_t < 4 \text{ GeV}/c \).

The fractions of thermal contribution as a function of \( p_t \) from different phases are shown in Fig. 11(b). Partial chemical equilibrium in the hadronic phase is used in this hydrodynamic simulation to keep the number of hadrons fixed below \( T_{ch} \). If we ignore the contribution from particle decays and use a full chemical equilibrium (FCE), the photon emission rate from hadronic phase \( \gamma_{\text{had}} \) can be used in this case. In case of the PCE, the chemical potential \( \mu_i \) for all hadronic species \( i \) will modify the photon emission rate from hadronic gas, roughly estimated by a factor of \( \exp([\mu_1 + \mu_2]/T) \) for a subprocess of \( 1 + 2 \rightarrow 3 + \gamma \) according to kinetic theory with a Maxwell-Boltzmann statistics for all particles. This factor finally increases contribution from the HG with a factor about 2. Nevertheless, in the total thermal contribution, PCE or FCE consideration does not make a visible difference. For all centralities from 0-10% to 50-60%, the QGP phase emits most of direct photons above \( p_t \sim 1 \text{ GeV}/c \). Although the volume of hadronic phase is much bigger than the one of the QGP phase due to expansion, the photon emission rate from the hadronic phase decreases even faster with temperature. The competition between volume and emission rate results in the biggest
Figure 9: (Color Online) Competition among different sources for direct photon production in Au+Au collisions at $\sqrt{s_{NN}} = 200$ GeV for different centralities. The leading order contribution (LO) from primordial elementary scatterings is plotted as dotted lines, thermal contribution as dash-dotted lines, fragmentation contribution as dashed lines and jet-photon conversion as solid lines.

contribution from the QGP phase at $p_t > 1$ GeV/$c$. In the current setting of hydrodynamic simulations at the RHIC energy, the mixed phase exists for a very long time ($\sim 8$ fm/$c$). This contributes mostly at lower $p_t$ values. By combining the results shown in Figs. 9 and 11, contribution of thermal radiation from the QGP phase is dominant in the region $1 \lesssim p_t \lesssim 4$ GeV/$c$. This momentum window may provide information inside the hot and dense matter, e.g., the initial temperature at the center, which may not be reached directly by hadron spectra.

Figure 10: (Color Online) The contribution to $R_{AA}$ from fragmentation and conversion, for different centralities.
VI. CONCLUSION

We calculated the centrality-dependence of $p_t$ spectra for direct photons in Au+Au collisions at the RHIC energy, based on a realistic data-constrained $(3+1)$-D hydrodynamic description of the expanding hot and dense matter, a reasonable treatment of propagation of partons and their energy loss, and a systematic consideration of main sources of direct photons. In this study, four main sources are considered, namely, leading order (LO) contribution from primordial elementary scatterings, thermal radiation from the fluids, fragmentation from hard partons, and jet photon conversion (JPC). Similar work [41] has been done before the appearance of the most recent data [38]. Our results agree nicely with the recent low $p_t$ data.

The role of jet quenching in the high $p_t$ region of direct photons production has been checked via fragmentation photons and jet photon conversion sources. For these two sources, the suppression of the photon rate due to parton energy loss is significant in central Au+Au collisions, and becomes less important towards peripheral collisions, similar to the suppression for meson production. Since experimentally one may separate isolated photons (LO+JPC) and associate photons (fragmentation photons), our prediction may be tested in the future.

Considering the total yields of direct photons, the contribution from fragmentation and conversion are small, contributing between 5% and 10%. However, parton energy loss plays nevertheless an important role: Without it, these second order effects would contribute significantly. Without jet quenching, the nuclear modification factors $R_{AA}$ would depend visibly on the centrality of the collisions. A strong energy loss is actually necessary to get the centrality scaling of $R_{AA}$ in our calculation – a centrality scaling which has been observed by the PHENIX collaboration. In this sense, properties of the bulk matter affect the photon yields at intermediate values of $p_t$, via the parton energy loss.

The low $p_t$ region is totally dominated by thermal radiation, providing direct information about the bulk matter. We find that $R_{AA}$ of photons at $p_t$ values below 1 GeV/$c$ is centrality independent. With increasing $p_t$, the $R_{AA}$ for peripheral collisions drops much faster than the one for central scatterings. On the other hand, thermal photons are mainly emitted from the QGP phase at $p_t>1$ GeV/$c$ even though the mixed phase and the HG phase occupy bigger space and longer time. So the different behavior of $R_{AA}$ for central and peripheral collisions, in the range 1 GeV/$c<p_t<3$ GeV/$c$, manifests the fact that the plasma in central collisions is hotter compared to peripheral collisions.

Still more investigation are needed for a precise characterization of the properties of the plasma via thermal photons. Besides, the elliptic flow of direct photons (especially thermal photons) should provide more information of the plasma, which will be discussed elsewhere.

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