A theoretical derivation of photon bremsstrahlung, induced by the interactions of an energetic quark in a hot and dense quark-gluon plasma, is given in the framework of the reaction operator approach. For the physically relevant case of hard jet production, followed by few in-medium interactions, we find that the Landau-Pomeranchuk-Migdal suppression of the bremsstrahlung photon intensity is much stronger than in the previously discussed limit of on-shell quarks and a large number of soft scatterings. This result is incorporated in the first systematic study of direct photon production in minimum bias d+Cu and d+Au and central Cu+Cu and Au+Au heavy ion collisions at the Relativistic Heavy Ion Collider at center of mass energies √s = 62.4 GeV and 200 GeV. We find that the contribution of the photons created via final-state interactions is limited to 35% for 2 GeV < pT < 5 GeV and at high transverse momenta the modification of the direct photon cross section is dominated by initial-state cold nuclear matter effects.

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I. INTRODUCTION

In the highly successful hard probes program at the Relativistic Heavy Ion Collider (RHIC), the interplay of nuclear effects that alter the cross section for direct photon production is not yet well understood. Direct γ measurements [1, 2] have provided an important baseline to help establish the dominance of final-state effects for the γ production and the subsequent quark scattering, which is given in the framework of the reaction operator approach. For the physically relevant case of hard jet production, followed by few in-medium interactions, we find that the Landau-Pomeranchuk-Migdal suppression of the bremsstrahlung photon intensity is much stronger than in the previously discussed limit of on-shell quarks and a large number of soft scatterings. This result is incorporated in the first systematic study of direct photon production in minimum bias d+Cu and d+Au and central Cu+Cu and Au+Au heavy ion collisions at the Relativistic Heavy Ion Collider at center of mass energies √s = 62.4 GeV and 200 GeV. We find that the contribution of the photons created via final-state interactions is limited to 35% for 2 GeV < pT < 5 GeV and at high transverse momenta the modification of the direct photon cross section is dominated by initial-state cold nuclear matter effects.

Recent phenomenological refinements [14, 17] suggest that the QGP enhancement of direct γ production may be smaller than previously expected and partly cancelled by the quenching of fragmentation photons. Still, there is no calculation to date that consistently includes known nuclear matter effects, such as the Cronin effect [18, 19, 20], shadowing [21, 22], and cold nuclear matter energy loss [23], to provide quantitative guidance for the relative strength of initial- and final-state modifications in the observed γ cross sections. An additional serious deficiency in the theory and phenomenology of direct photon production in heavy ion collisions is the absence of systematic studies in proton-nucleus (p+A) and nucleus-nucleus (A+A) collisions for different system sizes and center of mass energies. This is especially true now, when new experimental results from RHIC are soon expected to become available [14, 21]. Last but not least, only through extensive detailed comparison between theory and data [5] can one test the model validity and gain confidence in the extracted quantitative properties of the dense matter created in heavy ion reactions.

With this motivation, we derive the QGP-induced γ spectrum for hard quark production in finite size plasmas. The same model of jet-medium interactions is used to calculate the quark conversion cross section and the suppression of fragmentation photons. These theoretical results, when applicable, are combined in a numerical simulation with cold nuclear matter effects to provide model predictions for d+Cu, d+Au, Cu+Cu and Au+Au reactions at center of mass energies of 62.4 GeV and 200 GeV per nucleon pair at RHIC. This article is organized as follows: in section II we highlight the differences between gluon and photon bremsstrahlung and identify the theoretical approach that reproduces the known Bethe-Heitler spectrum. The derivation of the final-state medium-induced photon radiation in the reaction operator approach is given in section III. Numerical results, relevant to the phenomenology of direct γ production are also shown. In section IV we carry out a systematic investigation of cold and hot nuclear matter effects that alter the mid-rapidity photon cross section in ultra-relativistic collisions of heavy nuclei at RHIC.
summary and conclusions are presented in section V.

II. PHOTON VERSUS GLUON BREMSSTRAHLUNG

The computation of photon bremsstrahlung is usually considered to be easier than that of gluon bremsstrahlung due to the absence of self-interactions of the gauge boson. However, it is not well appreciated that the two physics processes are quite different. To illustrate this, we first examine the radiative amplitude that corresponds to the case of single scattering of a fast on-shell quark 11, 20:

\[ \mathcal{M}_{\text{rad}}(k) \propto 2ig_s\epsilon_\perp \cdot \left( \frac{k_\perp}{k^2} - \frac{(k-q_\perp)}{(k-q)^2} \right) e^{i\frac{q_\perp \cdot z^+}{2}} T^c, T^a \]. \hspace{1cm} (1)

In Eq. (1) \( g_s \) is the strong coupling constant, \( k^\mu = [k^+, k^-, k_\perp] \) is the momentum of the radiated gluon in light-cone coordinates and \( \epsilon \) is its polarization vector. We denote by \( q^\mu = [q^+, 0, q_\perp] \) the momentum exchange with the medium at position \( z \) and by \( T^c, T^a \in SU(3) \) the color matrices at the emission and interaction vertices. Evidently, it is the color rotation of the parent parton and the re-interaction of the bremsstrahlung gluon in nuclear matter that determine the gluon emission intensity and allow neglect of the deflection of the jet. If we, however, take the Quantum Electro-Dynamics (QED) limit \( g_s \rightarrow \epsilon_q = (\pm 1/3, \pm 2/3)\epsilon, \ T_e \rightarrow 1 \) in Eq. (1) we find \( \mathcal{M}_{\text{rad}}(k) \rightarrow 0 \). Therefore, a theoretical approach developed to describe gluon emission cannot be directly generalized to photon emission and vice versa 27. Our conclusion is independent of the specific example of incoherent parton scattering. All regimes of coherent inelastic scattering can be treated in the unified framework of the reaction operator approach 22. The reaction operator \( \hat{R}_n \) describes the effect of one additional correlated in-medium scattering at position \( z_n \) at the cross section level and is process dependent. Taking the above mentioned QED limit of \( \hat{R}_n \) derived for gluon bremsstrahlung 11, one finds:

\[ \hat{R}_n = T_a T_\alpha - (C_F/2)1 - (C_F/2)1 \equiv 0. \hspace{1cm} (2) \]

In Eq. (2) \( C_F \) is the quadratic Casimir in the fundamental representation of \( SU(3) \). Thus, better treatment of jet-medium interactions is needed for both incoherent and coherent photon emission calculations.

With these results in mind, we first identify the refinement of the kinematic approximations necessary to derive the induced \( \gamma \) spectrum. The scattering of a fast quark in nuclear matter is modeled via interactions with an external non-Abelian field \( V^{\mu, c}(q) \) 18:

\[ V^{\mu, c}(q) = -n^\mu 2\pi \delta(q^+) T^c(q) e^{iq_\perp \cdot z^+} \]

\[ g_s V^c(q) \equiv v(q) T^c(t) \]. \hspace{1cm} (3)

Here, the four-vector \( n^\mu = \delta^{\mu, -} = [0, 1, 0_\perp] \) and the color matrix \( T^c(t) \in SU_c(3) \) in Eq. (3) represents the target charge that creates the non-Abelian field. We take the Fourier transform \( v(q) \) to be of color-screened Yukawa type but with Lorentz boost invariance:

\[ v(q) \equiv \frac{4\pi\alpha_s}{-q^2 + \mu^2} = \frac{4\pi\alpha_s}{q_\perp^2 + \mu^2} = v(q_\perp) \], \hspace{1cm} (4)

where we have used the \( q^+) = 0 \) choice of frame. This specific form of \( v(q) = v(q_\perp) \) is particularly useful since in-medium interactions in both hot and cold nuclear matter are of finite range \( r_{\text{int.}} = \mu^{-1} \) and we shall assume that \( \lambda_\mu \gg 1 \), where \( \lambda \) is the quark mean free path.

The differential photon bremsstrahlung spectrum arises from single-Born scattering diagrams shown in the top panel of Fig. 1. Using a high energy approximation for the quark to simplify the interaction and emission vertices we obtain:

\[ iM_{\text{RHS}}^D(k) = \int \frac{d^4q}{(2\pi)^4} \left( -ie \right) \frac{i\epsilon^\mu(2p_f + k)_\mu}{(p_f + k)^2} \frac{v^{\nu, c}(q)(2p_f + 2k - q)_\nu}{(p_f + k - q)^2 + i\epsilon} \]

\[ \approx -i\int \frac{d^2q_\perp}{(2\pi)^2} e^{i\frac{q_\perp \cdot z^+}{2}} T^c(p) T^c(t) \]

\[ \times e^{ie \cdot p_f \cdot k / k \cdot p_f} e^{iz^+ k^-}. \hspace{1cm} (5) \]

For the second \( \gamma \) emission diagram similar considerations lead to:

\[ iM_{\text{LHS}}^D(k) \approx \left[ -i\int \frac{d^2q_\perp}{(2\pi)^2} e^{i\frac{q_\perp \cdot z^+}{2}} T^c(p) T^c(t) \right] \]

\[ \times e^{ie \cdot p_f \cdot k / k \cdot p_f} e^{iz^+ k^-}, \hspace{1cm} (6) \]

and the radiative matrix element at position \( z_i \) reads:

\[ \mathcal{M}_{\text{rad}}(k, \{i\}) = e^{ie \cdot p_f \cdot k / k \cdot p_f} e^{iz_i^+ k^-}. \hspace{1cm} (7) \]

![FIG. 1: Top panel: single-Born diagrams for medium-induced \( \gamma \) emission. Bottom panel: the corresponding double-Born diagrams.](image-url)
In Eq. (7) the collisional amplitude is not shown. Let the initial- and final-state momenta of a fast on-shell quark be $p_i = [E^+, Q_2, \gamma_{i-1}/(2E^+), Q_\perp i]$, $p_f = [E^+, Q_2', \gamma_{i+1}/(2E^+), Q_\perp i]$, such that $Q_\perp i - Q_\perp i - 1 = Q_\perp i$. The double differential medium-induced photon distribution is then given by:

\[
k^+ \frac{dN^\gamma(k; \{i\})}{dk^+ d^2k_\perp} = \frac{\alpha_{em}}{\pi^2} \left( \frac{k^+}{\left( k^+ + \gamma_{i-1}/2 \right)^2} \right)^2 \left( \frac{Q_\perp i}{Q_\perp i - 1} \right)^2,
\]

and is dominated by emission coincident with the directions of the incoming and the outgoing quarks. Changing variables $k = k^+ - k_\perp^{(pole)}$, $k_\perp^{(pole)} = Q_\perp i, k^+/E^+, Q_\perp i - 1/k^+/E^+$, respectively, we obtain the QED double logarithmic result:

\[
N^\gamma(\{i\}) \approx 2 \frac{\alpha_{em}}{\pi} \ln \frac{k^+}{k^+_{\min}} \ln \frac{q^2_{\max}}{m^2}, \tag{9}
\]

where $m^2$ regulates the collinear divergence.

In the case of coherent gluon emission in finite media with few subsequent scatterings, the interference between the hard (vacuum) and soft (medium-induced) bremsstrahlung largely determines the LPM cancellation pattern \[3, 10\]. Double-Born diagrams, with two momentum exchanges at the same position $z_i = z_i'$, can contribute at any fixed order in opacity. The relevant double-Born diagrams for photon emission are shown in the bottom panel of Fig. 1. We find:

\[
iM^V_{RHS}(k) = \int \frac{d^3q_1}{(2\pi)^3} \frac{d^3q_2}{(2\pi)^3} \frac{(-ie)^2(2p_1 + k\gamma)_{\mu}}{(p_f + k)^2 + ie} \times (-ig_s T^\mu(\gamma)) \frac{iV_{\nu,c}(q_2)(2p_f + k^\mu - q_{\nu})_{\nu}}{(p_f + k - q_{\nu})^2 + ie} \times (-ig_s T^\nu_d(q_2)(2p_f + k^\nu - q_{\nu})_{\nu})\frac{iV_{\nu,c}(q)(2p_f + k^\nu - q_{\nu})_{\nu}}{(p_f + k - q^\nu)^2 + ie} \approx \frac{1}{A_\perp} \left[ -\frac{1}{2} \int \frac{d^2q_\perp}{(2\pi)^2} |v(q_\perp)|^2 C_i C_{\mu}\frac{m^2}{d_A} \right] \times e \left( \frac{e}{k \cdot p_f} \right) \delta^2(q_\perp - q_\perp') \times e^{i\varepsilon^+ k^-}. \tag{10}
\]

To obtain the result in Eq. (10) we averaged over the initial and summed over the final parton colors. We also carried out the average over the position of the scattering center in the transverse plane: $A_{\perp}^{-1} \int d^2z_\perp \exp[-iz_\perp \cdot (q_\perp + q_\perp')] = A_{\perp}^{-1} (2\pi)^2 \delta^2(q_\perp + q_\perp')$. For single-Born interactions, see Eqs. \[5\] and \[6\], such averages are possible only after squaring the amplitudes. The result differs from the one for virtual interactions in yielding $\delta^2(q_\perp - q_\perp')$, and in the absence of the factor $-1/2$, which accompanies the collision term.

Similarly, for the second virtual diagram in Fig. 1 we obtain:

\[
iM^V_{RHS}(k) \approx \frac{1}{A_\perp} \left[ \frac{1}{2} \int \frac{d^2q_\perp}{(2\pi)^2} |v(q_\perp)|^2 C_i C_{\mu}\frac{m^2}{d_A} \right] \times e \left( \frac{e}{k \cdot p_f} \right) e^{i\varepsilon^+ k^-}. \tag{11}
\]

Adding Eqs. (10) and (11) we see that the same general radiation matrix element, Eq. (7), can be factorized for double-Born interactions. In Eqs. (10) and (11) $d\sigma/d^2q_\perp = (C_i C_{\mu}/d_A) |v(q_\perp)|^2/(2\pi)^2$ is the differential scattering cross section, calculated in the Born approximation. However, $q_\perp + q_\perp' = 0$, and $p_i = p_f$ implies that the double-Born interaction does not contribute a new photon bremsstrahlung amplitude.

### III. Differential Photon Bremsstrahlung Spectrum to All Orders in Opacity

To define the iterative procedure of computing the photon bremsstrahlung contribution from multiple scattering, we first consider the action of the direct operator $\mathcal{D}_n$ at position $z_n$ on a radiative amplitude with $n - 1$ correlated scatterings. In what follows we have dropped the collisional amplitudes since they were shown to yield an elastic scattering cross section per order in opacity for both diffusion \[18, 28\] and radiative \[10\] processes. The result of such action can be represented as:

\[
\mathcal{D}_n \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k) \equiv \hat{(1 + \mathcal{B}_n)} \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k) = \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k) + \left(-\frac{1}{2}\right)^{N_v} \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}} \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k, \{n\}). \tag{12}
\]

Here, the factor $\left(-\frac{1}{2}\right)^{N_v}$ arises because every virtual contact interaction in the amplitude gives a factor $-\frac{1}{2}$, and $N_v(\mathcal{M}^{rad}_{i_1 \cdots i_{n-1}})$ is their number. The first term in Eq. (12) corresponds to a momentum exchange with the energetic jet. In the high energy limit we do not keep track of the transverse modification of the parent parton except for the contribution to the soft photon bremsstrahlung. This approximation does not affect the calculation of the intensity spectrum $d\Gamma/dk^+$ but the angular distribution must be convoluted with the medium-induced jet acoplanarity, which we here neglect. The second term in Eq. (12) is the new radiative contribution, Eq. (7), at position $z_n$, and the prefactor accounts for the number of preceding virtual interactions in the amplitude. Next, we consider the double-Born interaction of the quark at position $z_n$. From section \[1\] we know that there is no new $\gamma$ bremsstrahlung contribution since $p_f = p_i$. However, a factor $-1/2$ arises for the forward elastic scattering. The modification of the amplitude $\mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k)$ is found to be:

\[
\mathcal{V}_n \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k) = -\frac{1}{2} \mathcal{M}^{rad}_{i_1 \cdots i_{n-1}}(k). \tag{13}
\]
In both Eq. (12) and Eq. (13) we have omitted the color factors since for the simple case of individual parton propagation these are trivially absorbed in the elastic scattering cross section or inverse mean free path per order in opacity 18.

Let \( i_n = 0, 1, 2 \) indicate for each \( n \): no interaction, direct interaction and virtual interaction with the medium, respectively 10. Expansion in powers of \( \sigma^2 \) or, equivalently, \( 1/\lambda_q \) requires that in the conjugate amplitude \( \tilde{\mathcal{M}}_{rad}^{i_1 \cdots i_n}(k) \) we have \( i_n = 2 - i_n \). Consequently, the contribution to the radiation pattern at \( n \)-th order in opacity is:

\[
\begin{align*}
&dN^\gamma(k, n) \propto 2 \sum_{i_1 \cdots i_n=0} \tilde{\mathcal{M}}_{rad}^{i_1 \cdots i_n}(k) \mathcal{M}_{rad}^{i_1 \cdots i_n}(k) \\
&= \sum_{i_1 \cdots i_n=0} 2 \tilde{\mathcal{M}}_{rad}^{i_1 \cdots i_n-1}(k)(\hat{D}^\dagger D + \hat{V}^\dagger + \hat{V}) \mathcal{M}_{rad}^{i_1 \cdots i_n-1}(k) \\
&= \sum_{i_1 \cdots i_n=0} 2 \tilde{\mathcal{M}}_{rad}^{i_1 \cdots i_n-1}(k)(\hat{B}^\dagger_n \hat{B}_n + \hat{B}^\dagger_n + \hat{B}_n) \mathcal{M}_{rad}^{i_1 \cdots i_n-1}(k) .
\end{align*}
\]

(14)

For medium induced photon emission, the reaction operator \( \hat{R}_n = \hat{D}^\dagger D_n + \hat{V}^\dagger + \hat{V}_n = \hat{B}^\dagger_n \hat{B}_n + \hat{B}^\dagger_n + \hat{B}_n \) has a particularly simple form. The first term in Eq. (14) vanishes beyond first order \( (n = 1) \) in opacity 10:

\[
\begin{align*}
&\sum_{i_1 \cdots i_n=0} 2 \tilde{\mathcal{M}}_{rad}^{i_1 \cdots i_n-1}(k) \hat{B}^\dagger_n \hat{B}_n \mathcal{M}_{rad}^{i_1 \cdots i_n-1}(k) \\
&= |\mathcal{M}_{rad}(k, \{n\})|^2 \sum_{i_1 \cdots i_n=0} 2 (-\frac{1}{2}) \tilde{N}_e (-\frac{1}{2}) N_e = 0 ,
\end{align*}
\]

(15)

where we have used \( (-\frac{1}{2} - \frac{1}{2} + 1)^{n-1} = 0 \) for \( n \geq 2 \). The 2nd and 3rd terms in Eq. (14) yield:

\[
\begin{align*}
2 \text{Re} \mathcal{M}_{rad}^\gamma(k, \{n\}) &\sum_{i_1 \cdots i_n=0} 2 (-\frac{1}{2}) \tilde{N}_e \mathcal{M}_{rad}^{i_1 \cdots i_n-1}(k) \\
&= 2 \text{Re} \mathcal{M}_{rad}^\gamma(k, \{n\}) |\mathcal{M}_{rad}(k, \{n-1\})|^2 \\
&\times \sum_{i_1 \cdots i_n=0} 2 (-\frac{1}{2}) \tilde{N}_e (-\frac{1}{2}) N_e = 0 ,
\end{align*}
\]

(16)

if \( n \geq 3 \). Therefore, unlike the case of gluon bremsstrahlung, photon bremsstrahlung contributions vanish beyond second order in opacity.

Taking into account the interactions of the parent quark along its trajectory through the QGP and the average over the transverse momentum transfers, the main theoretical result derived in this letter reads:

\[
k^+ \frac{dN^\gamma(k)}{dk^+ d^2k_\perp} = \alpha_{em} \frac{\pi^2}{\lambda_q \sigma_{el}^2} \int d^2q_\perp \int \frac{1}{d^2\sigma^\gamma} \frac{d^2\sigma^\gamma}{d^2q_\perp} \\
\times \left[ |\mathcal{M}_{rad}(\{1\})|^2 + 2 \mathcal{M}_{rad}^\gamma(\{1\}) \mathcal{M}_{rad}(\{0\}) \right] \cos(k^- \Delta z^\gamma) \\
+ \prod_{i=1}^2 \int d^2q_\perp \frac{1}{\lambda_q \sigma_{el}^2} \\
\times 2 \mathcal{M}_{rad}^\gamma(\{2\}) \mathcal{M}_{rad}(\{1\}) \cos(k^- \Delta z^\gamma) .
\]

(17)

In Eq. (17) \( \Delta z^\gamma = z^\gamma - z^\gamma_{i-1} \), the \( \Delta z \) integrals are nested, and \( \tau_f^{-1} = k^2/(2\omega) \approx \sqrt{2} k^- \) is the inverse photon formation time. When \( \tau_f^{-1} \lambda_q \gg 1 \) the photons decohere early from the parent quark and our result reduces to incoherent emission from individual scattering centers. Our result differs from previous findings 7, 8 in several important ways: first, it treats the case of finite \( L/\lambda_q \sim \text{few} \), relevant to heavy ion physics. Second, we find that the most significant contribution to the LPM effect for photons comes from the interference of the medium-induced photon radiation with the hard emission from the large \( Q^2 \) scattering of the parent quark. Finally, we find that there can be non-linear corrections \( \sim L^2 \) to the dominant linear in \( L \) behavior of the photon spectrum. Note that calculations of final-state \( \gamma \) emission have also been carried out in deep inelastic scattering on nuclei 22, 29 using the high twist approach 11.

We are now ready to study numerically the final-state QGP-induced photon spectrum. Our results are limited to first order in opacity. The question of whether destructive interference may cancel part of the hard photon bremsstrahlung is deferred for future studies. With Eq. (7) representing the amplitude of both light and heavy fermions, the intensity spectrum is easily generalized to quarks with physical and thermal mass, \( m^2 = M_q^2 + C_F g^2 T^2 \). Notable differences from the \( m = 0 \) case include regulation of the poles in Eq. (5) via \( (k^+/E^+) \)2\( m^2 \), appearance of new terms \( \propto m^2 \), and finite quark velocity, \( \beta < 1 \). We show an example of a quark propagating outward from the center of the medium created in \( b = 3 \) fm Au+Au collisions at RHIC. For a typical integrated gluon rapidity density, \( dN^g/dy \approx 1150 \), distributed proportional to the 2D participant number density \( dN_{\text{part}}/d^2x_\perp \), we calculate the temperature \( T(x_\perp, t) \) as a function of time and position in the transverse plane, assuming approximate Bjorken expansion of the QGP near midrapidity \( y = 0 \). The necessary Debye screening scale \( \mu(T) = m_D(T) = g_s T \), elastic scattering cross section \( \sigma^g(T) \), and quark mean free path \( \lambda_q(T) = 1/\sigma^g(T) \rho(T) \) are evaluated as in 3 with \( g_s = 2.5 \). The top panel of Fig. 2 shows the medium-induced photon number spectrum \( dN^\gamma/dx = \langle e/\epsilon_q \rangle^2 dN^\gamma/dx \), \( x = k^+/E^+ \), normalized by the squared fractional quark electric charge (bar will denote such scaling for any physics quantity). We considered light, \( M_{u,d} = 0 \) GeV, and heavy, \( M_c = 1.5 \) GeV, \( M_b = 4.5 \) GeV, quarks of energy \( E_q = 100 \) GeV in the Bethe-Heitler limit and to first
order in opacity. In the absence of coherence, the induced \( \gamma \) spectrum is dominated by photon energies \( \omega \sim \text{a few } xT \). The contributions of the heavy quark sector to the medium-induced photon multiplicity and the energy loss due to photon emission are strongly suppressed. When the interference between the vacuum and the medium-induced photon radiation is taken into account, we find that the spectrum \( dN_\gamma/dx \) is suppressed much more effectively than in the limit of very large number of soft scatterings for on-shell quarks \( [4] \). The bottom panel of Fig. 2 shows the dependence of the partial fractional energy loss due to \( \gamma \) emission versus the mean number of quark interactions \( \langle n \rangle = L/\lambda_q \). Insert illustrates the energy dependence of \( \Delta E_\gamma^{(7)}/E_q \). It is correlated with deviations from the naive expectation, \( \Delta E_\gamma^{(7)} \propto E_q \), as illustrated in the insert in Fig. 2. Medium-induced photon emission in both the coherent and incoherent limits is too small, \( \Delta E_\gamma^{(7)}/E_q < 1\% \), to contribute to the quenching of quark jets. Last but not least, we emphasize that the number of interactions in the QGP, even for jets emerging from the center of the heavy ion collision region, is small, \( \langle n \rangle = 3.1 \). As in the case of gluon emission, this is a clear indication that, even in the most central Au+Au reactions at RHIC, we are not in the limit of large number of scatterings. For consistency with the calculation of light hadron attenuation, in our numerical estimates we use an effective geometry of \( L = \text{6 fm} \) and uniform distribution of partons \( [4] \). In this case \( \langle n \rangle \) is reduced to 2.4, accounting for jets close to the periphery of the interaction region.

IV. PHENOMENOLOGY OF HARD PHOTON PRODUCTION IN P+A AND A+A COLLISIONS

With the results from the previous section at hand, we now turn to the question of hard, \( p_T > \text{a few GeV} \), photon production in heavy-ion collisions. QGP and cold nuclear matter effects are identified through the nuclear modification ratio:

\[
R_{AB}(p_T, b) = \frac{d\sigma_{AB}/dyd^2p_T}{N^{coll}_{AB}(b) d\sigma_{pp}/dyd^2p_T}, \tag{19}
\]

where the number of binary collisions, \( N^{coll}_{AB}(b) \) in Eq. (19), is computed in an optical Glauber model. The baseline \( p+p \) cross section is evaluated in factorized perturbative QCD to lowest order and leading twist as fol-

FIG. 2: Top panel: medium-induced photon number spectrum versus \( x = k^+ / E^+ \) for \( E_q = 100 \text{ GeV} \) light, charm and bottom quarks in central Au+Au collisions at \( \sqrt{s} = 200 \text{ GeV} \). Both the incoherent Bethe-Heitler (solid) and the coherent final-state (dashed) bremsstrahlung cases are shown. Bottom panel: partial fractional quark energy loss \( \Delta E_\gamma^{(7)}/E_q \) due to \( \gamma \) emission versus the mean number of quark interactions \( \langle n \rangle = L/\lambda_q \). Insert illustrates the energy dependence of \( \Delta E_\gamma^{(7)}/E_q \).

FIG. 3: Direct photon production cross section in \( p+p \) collisions at \( \sqrt{s} = 62.4 \text{ GeV} \) and 200 GeV. Data at the higher RHIC energy is from PHENIX \( [4] \). Insert shows the fraction of fragmentation to all direct photons.
Figure 3 shows the cross sections for direct $\gamma \bar{\gamma}$ production in the nuclear medium is controlled by the state multiple scattering \[20\]. In our calculations the theoretical approaches to the Cronin effect are well documented \[31\]. We note that in semi-inclusive deep inelastic scattering (SIDIS) the final-state broadening is neglected in direct photon phenomenology, is a systematic correction to the initial-state energy loss. The complementary $p_T > 5 - 7$ GeV part of phase space is characterized by $R_{dA}(p_T) < 1$ with the isospin effect, included in all calculations, being one of the major contributors. The EMC effect only becomes noticeable at the lower C.M. energy and at the highest transverse momenta. Initial-state, cold nuclear matter energy loss have not yet been studied, in our calculations we make a similar approximation. We note that the weaker falloff with energy of the incoming parton flux, when compared to the hard scattered (additional $\sim 1/p_T^2$) final-state quark and gluon distributions, implies values of $\kappa$ closer to unity. The specific choice that we make, $\kappa = 0.7$, was constrained through the evaluation of the $p_T^0$ production cross section in $d+Au$ collisions at RHIC \[3\] which provided a good description of the experimental data.

Our results for minimum-bias $d+Au$ (solid lines) and $d+Cu$ (dashed lines) collisions are shown in Fig. 4. The calculated $R_{dA}(p_T)$ at $\sqrt{s} = 200$ GeV and 62.4 GeV are presented in the top and bottom panels, respectively. In all calculations isospin effect at high $p_T$. We emphasize that, at transverse momenta $\sim 15$ GeV, nuclear effects on the direct photon cross section can be as large as 20% at $\sqrt{s} = 200$ GeV and 40% at $\sqrt{s} = 62.4$ GeV in minimum
approximation $[12]$ that stems from the limit of small $t$-channel momentum transfers to energetic jets is implicit. For a parent quark propagating through the QGP:

$$N_{\text{conv}}(c) = \int_{t_0}^{L} dt \rho(T)\sigma_{\text{tot}}^{\gamma\gamma\rightarrow\gamma\gamma}(T), \quad (25)$$

where time and position dependence is taken via $T(x,t)$. The cross section in Eq. (24), with $s \approx 2m_D E$ and $t \in (m^2_D, s/4)$, consistent with the forward scattering approximation, reads:

$$\sigma_{\gamma\gamma\rightarrow\gamma\gamma} = \frac{\pi \alpha_s \alpha_{\text{em}}}{6m_D E} \ln \frac{E}{2m_D}. \quad (26)$$

Numerical results in central Au+Au and Cu+Cu collisions at the intermediate and top RHIC energies are shown in Fig. 5. The final-state gluon rapidity densities $dN^\gamma/dy = 1150$ (Au, 200 GeV), $800$ (Au, 62.4 GeV), $370$ (Cu, 200 GeV), $255$ (Cu, 62.4 GeV) are constrained by measured and extrapolated particle multiplicities at RHIC $[4]$. Of the final-state effects in the QGP, the left panels only include the quenching of fragmentation photons. Since a small fraction of the direct $\gamma$ come from fragmentation processes and quark attenuation is significantly smaller, $\sim C_F/C_A$, when compared to gluon attenuation, the observable QGP modification is also very small. In fact, $R_{AA}(p_T)$ is dominated by cold nuclear matter effects, such as the ones shown in Fig. 3, amplified by the presence of two large nuclei. $R_{AA}(p_T)$ from published $[1]$ and preliminary $[2, 4]$ data are shown for comparison. Given the error bars, only large Cronin enhancement at $\sqrt{s} = 62$ GeV is excluded, suggestive of the role of initial-state inelastic jet scattering in controlling the magnitude of the Cronin effect.

The right panels in Fig. 5 include the medium-induced photons and the jet conversion contribution. For completeness, we have also shown results for the incoherent Bethe-Heitler radiation spectrum and absence of initial state energy loss. This scenario gives direct $\gamma$ enhancement as large as factors of 3 and 4 over the p+p baseline at $\sqrt{s} = 200$ GeV and 62 GeV, respectively. Not only is this case not supported theoretically by the results derived in this paper and in Ref. $[22]$, but when compared to data, even with the present large experimental error bars, it is clearly excluded. In calculating the coherent final state photon emission rate, Eq. (17), and the jet conversion rate, Eq. (25), we also account for the time-dependent quenching of the quark as it propagates through the QGP:

$$R_{AA}(p_T, t) = (1 - f(t)) + R_{AA}(p_T)f(t). \quad (27)$$

Here, $R_{AA}(p_T)$ is the full final-state quark quenching, evaluated as in $[4]$ for different system sizes and center of mass energies, and $f(t)$ interpolates between 0 and 1 to give the time dependence of radiative energy loss $[22]$. We find that this effect reduces $dN^\gamma_{\text{med}}/dz$ and $N^\gamma_{\text{conv}}$ by $\sim 30\%$. Our results show that at transverse momenta

$\begin{align*}
\end{align*}$

bias $d+A$ collisions. Preliminary experimental data $[24]$ is consistent with the theoretical expectation. However, given the large error bars, it cannot discriminate between different cold nuclear matter effects or constrain their magnitudes. Careful experimental investigation is needed to pinpoint these effects and, since they will be enhanced in $A+A$ collisions, caution should be exercised in the interpretation of the $R^\gamma_{AA}(p_T)$ findings.

In $A+A$ collisions, QGP-induced modification of direct photon production cross section includes competing effects: the quenching of fragmentation photons and the tree level (jet conversion) and bremsstrahlung photon enhancement. The quenching of fragmentation photons is modelled in the same way as the quenching of hadrons $[21]$ and can be combined with the medium-induced $\gamma$ contributions as follows:

$$D_{\gamma/c}(z) = \int_0^{1-z} d\epsilon P(\epsilon) \frac{1}{1 - \epsilon} D_{\gamma/c} \left( \frac{z}{1 - \epsilon} \right)
+ \frac{dN^\gamma_{\text{med}}(c)}{dz} + N^\gamma_{\text{conv}}(c)\delta(1-z). \quad (24)$$

Here, $P(\epsilon)$ is the probability distribution of the fractional jet energy loss $\epsilon = \Delta E/E \quad [4, 32]$, $dN^\gamma_{\text{med}}(c)/dz$ is the QGP-induced bremsstrahlung that we calculated in section $[11]$ and $N^\gamma_{\text{conv}}(c)$ is the number of jets converted to photons. For such conversions, in Eq. (24) the $p_\gamma \approx p_c$
$p_T < 5 \text{ GeV}$ the contribution from jet conversion to the total photon cross section is limited to $\sim 25\%$ and the contribution of medium-induced $\gamma$ is limited to $\sim 10\%$. In the high $p_T$ range the total enhancement contribution is found to be $\sim 5\%$.

V. CONCLUSIONS

In this paper, we provided a theoretical derivation of the final-state QGP-induced photon bremsstrahlung for the experimentally relevant case of hard jet production. We demonstrated that while the physics processes that control photon and gluon emission differ, the common Landau-Pomeranchuk-Migdal interference between the radiation from the hard scattering and the radiation from the subsequent soft quark interactions in the plasma leads to a significant suppression of the $\gamma$ intensity. The photon spectrum was found to be attenuated at least by a factor of several for jet energies relevant for RHIC phenomenology, in contrast with the estimated modest $30\%$ attenuation for asymptotic $t = -\infty$ on-shell jets in the limit of a very large number of soft interactions, where the interference with the bremsstrahlung that accompanies hard jet production has not been taken into account.

We found that the suppression of $dI^\gamma/dk^+$ also leads to non-linear dependence of $\Delta E^\gamma_q(E/\lambda_q, E_q)$ on the system size and sub-linear dependence on the parent quark energy that have previously been associated primarily with gluon emission.

To help identify the significance of both cold and hot nuclear matter effects on direct photon production we carried out the first systematic phenomenological study of $R_{AA}^\gamma(p_T)$ in midrapidity $d+Cu$, $d+Au$, $Cu+Cu$ and $Au+Au$ reactions at RHIC energies of $\sqrt{s} = 20.4 \text{ GeV}$ and $200 \text{ GeV}$. As expected, in all cases we found that in the absence of QGP formation the nuclear modification factor at intermediate transverse momenta is dominated by the Cronin effect and at high transverse momenta by isospin effects and initial-state parton energy loss. Surprisingly, however, the contribution of final-state QGP effects to the direct photon cross section in nucleus-nucleus collisions was small: less than $-25\%$ from the quenching of fragmentation photons, less than $+25\%$ from jet conversion and less than $+10\%$ from medium-induced $\gamma$ bremsstrahlung. While experimental measurements are not yet precise enough to disentangle such modest effects, they have already put severe constraints on theoretical models that suggest a dominant role of jet-plasma interactions in the $p_T \geq 2 \text{ GeV}$ part of phase space. In particular, PHENIX data is compatible with strong coherent suppression of the bremsstrahlung $\gamma$, which is the main theoretical result of this paper.

In conclusion, we suggest that only a systematic study of direct photons in heavy ion reactions for various system sizes and center of mass energies will help uncover...
their full potential both as a baseline for jet tomography and as an independent probe of nuclear effects. In this exploration, precision d+A data is critical, since our theoretical results support the possibility that, in both proton-nucleus and nucleus-nucleus collisions, cold nuclear matter effects play a dominant role in altering the cross section for direct photon production.

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