Features of photons radiated off quarks escaping from a quark-gluon plasma

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(Received 24 February 1994; revised manuscript received 6 January 1995)

Radiation of photons off quarks escaping from a quark-gluon plasma and thereby accelerated by the confining color field is studied. The rate of emission of these photons is found to increase with their rapidity for a given energy. The p_T spectrum of the photons thus radiated is compared to those due to the Compton and annihilation processes in the quark-gluon plasma (QGP) after accounting for the space-time evolution of the plasma. The largest contributions are seen if the current mass for the quarks is used and then the photons produced due to the color-confining mechanism operating on the quarks are found to give a contribution of about 100 % at very low p_T and about 10–20 % at the transverse momenta of 2–3 GeV as compared to the photons due to the Compton and annihilation processes in the plasma. However, if a larger value is associated with the mass of the quark, this radiation is very strongly suppressed. This uncertainty raises doubts about the usefulness of such radiations as a signature of a QGP, which is rather unfortunate as these photons are known to be polarized.

PACS number(s): 12.38.Mh, 13.85.Qk, 25.75.+r

I. INTRODUCTION

One of the most important challenges in high energy nuclear physics at the moment is the creation, detection, and study of the quark-gluon plasma — a deconfined state of strongly interacting matter. It is likely to be produced in collisions involving heavy nuclei at the Relativistic Heavy Ion Collider (RHIC) at Brookhaven National Laboratory and at the Large Hadron Collider (LHC) and the Super Proton Synchrotron (SPS) at CERN. It does not seem to be unlikely that the quark-hadron phase transition is already taking place in S+Au collisions at the SPS. This has led to one of the best coordinated international efforts both theoretically and experimentally in recent times, which has paid rich dividends. Thus by now a series of very reasonable signatures of the formation of a quark-gluon plasma (QGP) have been identified, and detailed experiments have been planned to verify them.

High energy photons and dileptons provide one of the most reliable messengers of the advent of the QGP as they have a very large mean free path which allows them to escape the system without any further (final state) interaction. Their production cross section is also a very rapidly increasing function of the temperature, and thus they are produced most copiously just when the system is at its hottest (and densest) which allows them to probe the nascent QGP.

The experiments planning to use high energy photons to probe the QGP have to confront two rather daunting problems. The first one involves a very large background of photons produced by the decay of π^0 and η mesons. Fortunately, the problem gets easier with the increase in transverse momentum of the detected photons, and for $\gamma/\pi^0 \geq 0.05$ and $p_T \geq 1$ GeV, it may become tractable [1, 2].

The other problem has its origin in the finding of Kapusta *et al.* [3] which simply states that a hot hadronic gas "shines just as brightly" as the QGP at the same temperature. However [4], a hadronic gas will have a smaller entropy density as compared to the QGP at the same temperature. Additionally, it will appear only during the mixed phase and the hadronic phase of the evolution of the system produced in relativistic nuclear collisions when the temperatures are lower than during the (early) QGP phase. In spite of this, as the system will undergo a rapid transverse expansion during the latter stages of the evolution [5], the thermal photons from hot hadronic matter get "blueshifted" into the (transverse momentum) window available for thermal photons originating from the QGP. This makes the identification of "golden photons" produced from QGP somewhat involved.

In view of the above, the study of mechanisms which could lead to enhanced production of high energy photons from the QGP *but not from* hot hadronic matter can be very rewarding indeed. Two such efforts immediately come to mind. First, the fact that the quarks and the

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gluons are confined within the boundaries of the QGP leads to an enhancement of their distribution function by many orders of magnitude above the Boltzmann distribution normally assumed for this purpose, for larger transverse momenta [6]. This effect is obviously absent for the hadronic gas and may lead to a greatly enhanced production of high energy photons due to the Compton and annihilation processes in the QGP [7].

In the present work, we concentrate on yet another *unique* mechanism which could provide high energy photons from the QGP *but is absent for a hot hadronic gas.* Consider the quarks confined in the QGP. The quarks near the surface having momenta pointing outwards will try to escape the plasma. The collective color field of the plasma will get into action, masquerading as a chromoelectric flux tube. Now one of the following two things can happen. Either the flux tube will fragment, releasing a color singlet $q\bar{q}$ pair, or the flux tube will pull the quark back into the plasma, ensuring confinement. However, in this latter process, the quark will suffer a rapid acceleration which will lead to a radiation of photons similar to the magnetic bremsstrahlung [8] in ordinary electromagnetic plasma. It has been suggested [8] that this radiation may be as strong, *or even an order of magnitude stronger than the ("standard") radiation due to the Compton and annihilation processes.* In the following we reexamine the importance of this signature of the QGP in the light of recent (and more accurate) determination of the rates for photon emission from QGP due to the Compton and annihilation processes, and study its features in detail, specifically at SPS, RHIC, and LHC energies.

II. FORMULATION

A. General considerations

Let us briefly recall the basic formalism [8] for the radiation of photons off the accelerating quarks near the surface of the QGP. We use the chromoelectric flux tube model [9] to visualize the confinement. This model approximates the interaction between the plasma volume and any colored object crossing the boundary by a constant force $\sigma \approx 0.2$ GeV^2 which pulls the object back to the plasma. The force is assumed to act along the normal to the plasma surface (see Fig. 1). It is expected [10] that the temperature of the QGP at the RHIC (LHC) could be as high as 500 (900) MeV. Thus the de Broglie wavelength of the quarks is much smaller than the transverse dimensions of the QGP, making a semiclassical description of their motion meaningful. At the same time the energy of the u and d quarks $\langle E\rangle$ ($\approx 3T$) during the QGP phase is many orders of magnitude larger than their (current) mass $\approx$ 5–10 MeV. (We shall return to the consequences of taking a larger mass accounting for the many body effects on the quarks, or the constituent mass of the quark, later.) Therefore, as pointed out by the authors of Ref. [8], the so-called "coherent radiation length" [11]

$$l_{\rm coh} \;\approx\; 4\pi\, E\,(E-\omega)\,\omega^{-1}\,m^{-2} \tag{1}$$

becomes very large compared with the radius of curvature

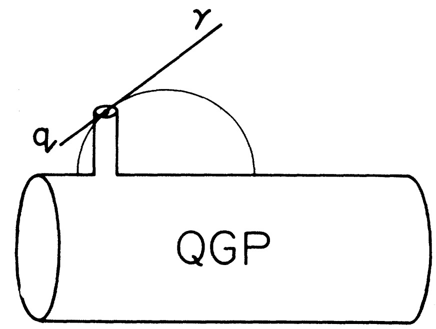

FIG. 1. A quark near the surface tries to escape the quark-gluon plasma. A chromoelectric flux tube is formed. It may either fragment, releasing a color singlet $q\bar{q}$ pair, or it may pull the quark back into the plasma. In the latter case the quark will experience strong acceleration and radiate photons.

of the trajectory of the quarks, $\rho \approx E/\sigma$, even when ω is quite close to the total energy of the radiating particle. This provides that the radiated power is independent of the mass of the particle [11], and, therefore,

$$\frac{dI}{dt} = 0.37\, e_q^2\, \alpha\, (E\,\sigma\, \sin\phi)^{2/3} \;, \tag{2}$$

where e_q is the charge of the quark in units of charge of the electron, α is the fine structure constant, and ϕ is the angle between the momentum of the quark and the normal to the surface of the plasma. The spectral radiation density has been shown [8] to be given by

$$\frac{dI}{d\omega\, dt} = 0.52\, e_q^2\, \alpha\, \omega^{1/3}\, (\sigma\, \sin\phi/E)^{2/3}, \quad 0 \le \omega < E \;. \tag{3}$$

This relation is invalid only when ω is very close to E. The radiated photons are confined to a narrow angle ($\approx m/E$) and we assume them to be collinear with the quarks. Using the laws of motion for the quarks moving under the influence of the confining force near the boundary and the distribution of the quarks in the plasma, the spectral distribution of the photons emitted per unit time per unit surface area of the plasma has been obtained [8] as

$$\frac{dN_\gamma}{dS\, dt\, \omega^2\, d\omega\, d\Omega} = \frac{0.52\, g\, \alpha}{(2\pi)^3\, \sigma^{1/3}} \times \frac{3}{7}\omega^{2/3} \left[\sum_{i=q,\bar{q}} e_i^2\right] \sin^{2/3}\phi_0 \times \int_1^\infty d\xi\, \exp\left(-\frac{\omega}{T}\xi\right)(\xi^{7/3}-1). \tag{4}$$

In the above, ϕ_0 is the angle between the normal to the surface of the plasma and the direction of the emitted photon, g is the number of the quark degrees of freedom due to spin and color and is equal to $2\times 3 = 6$, $\sum_{i=q,\bar{q}} e_i^2$ is sum of the squares of the charge fractions of the quarks and the antiquarks given by $2 \times (e_u^2 + e_d^2)$, and T is the temperature of the plasma. The number of photons per unit time per unit surface area is obtained by performing the integration over the direction Ω and the energy ω as

$$\frac{dN_\gamma}{dS\,dt} = A \left[\sum_{i=q,\bar{q}} e_i^2\right] \alpha\, T^{11/3}\, \sigma^{-1/3} , \tag{5}$$

$$A = \frac{1.56\, g\, 2^{5/3}\, \Gamma^2(4/3)}{(2\pi)^2} \approx 0.60 . \tag{6}$$

In order to have an estimate of the importance of this process, we apply it to the case of a longitudinally expanding plasma according to the Bjorken hydrodynamics [12] which is expected to be valid during the QGP phase. This provides that $dS\,dt = 2\,\pi\, R_T\, dz\, dt$ where R_T is the transverse dimension of the system, and we obtain the number of photons radiated per unit length per unit time as

$$\frac{dN_\gamma^{\rm rad}}{dz\,dt} = 2\pi\, R_T\, A \left[\sum_{i=q,\bar{q}} e_i^2\right] \alpha\, T^{11/3}\, \sigma^{-1/3}. \tag{7}$$

This may be compared to radiation of photons due to Compton and annihilation processes in the plasma, which are volume processes, and for which one of the best estimates to date [3] is given by

$$E_\gamma \frac{dN}{d^4x d^3k_\gamma} = \frac{5}{9} \times \frac{\alpha\, \alpha_s}{2\pi^2}\, T^2 \ln\left(\frac{2.912\, E_\gamma}{4\pi\, \alpha_s\, T} + 1\right) \times \exp(-E_\gamma/T) , \tag{8}$$

where E_γ is the photon energy and α_s is the QCD coupling constant. The integration over the photon momentum can be performed to get the number of photons emitted per unit time per unit volume. We have found that to a great degree of accuracy and for not too large values of α_s this can be written as

$$\frac{dN_\gamma^{C+{\rm ann}}}{d^4x} = \frac{5}{9} \times \frac{4\pi\, \alpha\, \alpha_s}{2\pi^2} \times 0.7\, T^4 \ln(1/\alpha_s) = \frac{5\,\pi\, \alpha\, \alpha_s}{63.5} \ln(1/\alpha_s)\, T^4 . \tag{9}$$

It is worthwhile to recall that this estimate for photons is more than twice the value obtained in early works [13]. This certainly reduces the relative strength of the radiated photons studied here by a factor of 2. Again, we write for the plasma undergoing a longitudinal expansion according to Bjorken hydrodynamics $d^4x = \pi\, R_T^2\, dz\, dt$ and thus the number of photons radiated per unit time per unit length of the plasma due to the Compton and annihilation processes is given by

$$\frac{dN_\gamma^{C+{\rm ann}}}{dz\,dt} = \frac{5}{9} \times 1.4\, R_T^2\, \alpha\, \alpha_s\, T^4 \ln(1/\alpha_s) . \tag{10}$$

This provides for an easy reckoning of the importance of the mechanism leading to photons radiated off quarks due to the action of the confining color field at the boundary, in situations likely to be attained during the early life of the QGP as created in relativistic collisions of nuclei. Thus we find that the ratio ($\mathcal{R}$) of the number of photons radiated per unit time per unit length of the cylindrical plasma as compared to the thermal photons emitted due to Compton and annihilation processes is given by

$$\mathcal{R} = \frac{dN_\gamma^{\rm rad}/\,dz\,dt}{dN_\gamma^{C+{\rm ann}}/\,dz\,dt} \tag{11}$$

$$= 1.063\,[\alpha_s \ln(1/\alpha_s)\, R_T\, T^{1/3}\, \sigma^{1/3}]^{-1} \tag{12}$$

$$= 5.0325\, [T({\rm GeV})]^{-1/3}\, [R_T\, ({\rm fm})]^{-1} \tag{13}$$

for $\alpha_s = 0.3$ and $\sigma = 0.2$ GeV^2. This expression clearly demonstrates the increasing importance of the radiation process considered here as the size of the system decreases and the temperature becomes smaller. Thus, for a collision involving two lead nuclei, this ratio is about 0.9 at T=500 MeV and increases to 1.2 at T=200 MeV. On the other hand, the corresponding values for collisions involving two oxygen nuclei are as high as 2.14 and 2.90, respectively. Unfortunately, the rate of emission of the photons decreases very rapidly with the decrease in temperature. Otherwise this could have led to the exciting possibility that the experiments looking for electromagnetic signatures of QGP could have become feasible and viable even for collisions involving smaller nuclei, which would have the additional advantage of providing a less messy background.

Thus we see that for situations likely to be attained at the SPS, RHIC, and LHC the radiation of photons due to the color-confining mechanism considered here could be as intense as that due to the Compton and the annihilation processes in the plasma, *if we take the current mass of the quark for these evaluations.* However, as remarked earlier, whereas the rates for the latter are similar to that for radiation from hot hadronic matter due to hadronic reactions and decay of vector mesons [3], there is no analogue of the radiation process considered here [8]. In fact, it can be argued that, even though we still have quark matter during the mixed phase, it is quite likely to be found in much smaller bubbles, where the conditions necessary for radiation of photons due to the color-confining mechanism operating on the quarks are no longer satisfied.

B. The transverse momentum spectrum

After these general considerations let us evaluate the transverse momentum spectrum of the photons radiated off the quarks near the boundary when the dynamic evolution of the plasma is accounted for. As a first step we find the rate per unit time per unit surface area for the photons.

Equation (4) can be manipulated after some algebra to give

$$\frac{dN_\gamma^{\rm rad}}{dS\,dt\,d^2k_T\,dy} = \frac{0.52\,g\,\alpha}{2\,(2\pi)^4\,\sigma^{1/3}} \times \frac{6}{7}\left[\sum_{i=q,\bar q} e_i^2\right] E_\gamma^{5/3} \times \int_0^{2\pi} d\alpha\,\left(1-\cos^2\alpha/\cosh^2 y\right)^{1/3} \times \int_1^\infty d\xi\, e^{-E_\gamma\xi/T}(\xi^{7/3}-1)\ . \tag{14}$$

A very interesting feature of this expression is that the rate of emission of a photon with energy E_γ is a function of not only the energy of the photon but also its rapidity (the polar angle), unlike for the case of the Compton plus annihilation processes above [Eq. (8)]. This has its origin in the dependence of the rate of radiation on the angle between the direction of motion of the quark and the confining force which acts along the normal to the surface of the plasma [see Eq. (2)]. Recalling that $E_\gamma = k_T \cosh y$, this implies an enhanced production of photons having smaller transverse momenta, for a given photon energy. This is best understood in terms of a universal function $\mathcal{Y}$ defined below (see Fig. 2) which gives the ratio of these rates for any rapidity to that for zero rapidity (largest k_T) for any given energy:

$$\mathcal{Y} = \frac{\int_0^{2\pi} d\alpha\,\left(1-\cos^2\alpha/\cosh^2 y\right)^{1/3}}{\int_0^{2\pi} d\alpha\,\left(1-\cos^2\alpha\right)^{1/3}}\ . \tag{15}$$

We shall see later that the use of the boost-invariant hydrodynamics of Bjorken while integrating over the space-time history of the plasma "robs" us of this unique rapidity dependence of the photon spectrum. It remains to be seen if a more general (boost-noninvariant) hydrodynamics [14] can retain and perhaps magnify the rapidity dependence of the spectra for the radiated photons compared to those due to Compton and annihilation processes and also the ones due to decay of π^0 mesons.

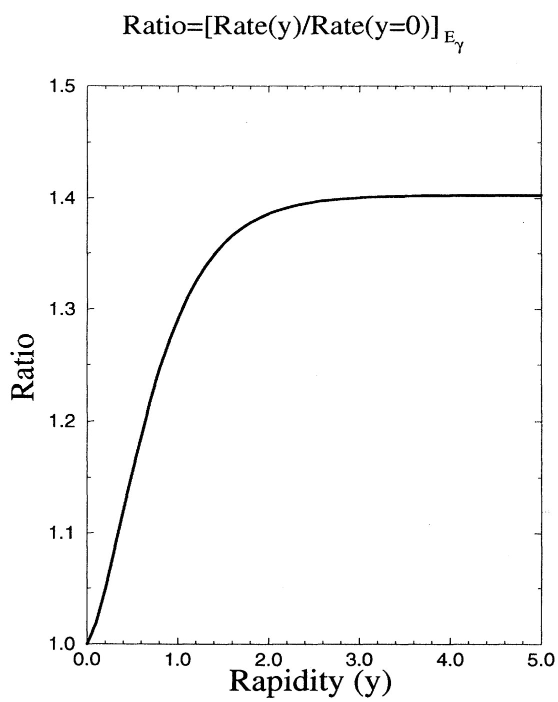

FIG. 2. The universal function $\mathcal{Y}$ giving the ratio of radiation of photons of a given energy per unit time per unit surface area and having a rapidity y to those with $y = 0$.

In the following we shall consider the emission of photons only during the QGP phase, and thus using Bjorken's hydrodynamics we write

$$dS\,dt = 2\,\pi\,R_T\,dz\,dt = 2\,\pi\,R_T\,\tau\,d\tau\,d\eta\ , \tag{16}$$

where τ is the proper time and η is the space-time rapidity of the fluid. Similarly, we write, for the four-volume element,

$$d^4x = \pi\,R_T^2\,dz\,dt = \pi\,R_T^2\,\tau\,d\tau\,d\eta. \tag{17}$$

The cooling law for Bjorken hydrodynamics is simply given by $T^3\tau = T_0^3\tau_0$, where T_0 is the temperature of the plasma [10] at the initial (formation) time τ_0. We shall consider emission during the time when the plasma cools from the initial temperature T_0 to the transition temperature T_c, which will be attained at the proper time $\tau_c = [T_0/T_c]^3\,\tau_0$.

Now the number of photons due to the radiation process considered here becomes

$$\frac{dN_\gamma^{\rm rad}}{d^2k_T\,dy} = \frac{1}{2} \times \frac{0.52\;g\,\alpha}{(2\pi)^4\,\sigma^{1/3}} \times \frac{6}{7}\left[\sum_{i=q,\bar q} e_i^2\right] k_T^{5/3}\,2\,\pi\,R_T \times \int_{\tau_0}^{\tau_c} \tau\,d\tau \int_{-\infty}^{\infty} d\eta\,\cosh(y-\eta) \times \int_0^{2\pi} d\alpha\,\left[\cosh^2(y-\eta) - \cos^2\alpha\right]^{1/3} \times \int_1^\infty d\xi\, e^{-k_T\,\xi\,\cosh(y-\eta)/T}(\xi^{7/3}-1)\ . \tag{18}$$

The corresponding expression for photons emitted due to the Compton plus annihilation processes is given by

$$\frac{dN_\gamma^{C+\rm ann}}{d^2k_T\,dy} = \frac{5}{9} \times \frac{\alpha\,\alpha_s}{2\,\pi^2}\pi\,R_T^2 \int_{-\infty}^{\infty} d\eta \int_{\tau_0}^{\tau_c} d\tau\,\tau\,T^2 \times \ln\left[\frac{2.912\,k_T\,\cosh(y-\eta)}{4\pi\,\alpha_s\,T} + 1\right] \times \exp\left[-k_T\,\cosh(y-\eta)/T\right]\ . \tag{19}$$

The total number of photons produced in any rapidity window $Y_1 \le y \le Y_2$ is then easily obtained by integrating the above expressions over d^2k_T. It may again be noted that the use of Bjorken's boost-invariant hydrodynamics renders the distribution of photons independent of the rapidity. We have evaluated these expressions numerically, without any further approximation.

C. Initial conditions

We shall relate the pion multiplicity to the initial time and initial temperature appropriate for Bjorken hydrodynamics as

$$T_0^3 \tau_0 = \frac{2\,\pi^4}{45\,\zeta(3)\,\pi\,R_T^2\,4\,a_Q}\;\frac{dN_\pi}{dy}\;, \tag{20}$$

where dN_π/dy is the pion multiplicity and $a_Q = 37\pi^2/90$ for a system consisting of u and d quarks. Further assuming that the initial time and the initial energy of the particles should be related via the uncertainty principle, we get $\tau_0 \approx 1/(3T_0)$ (see [5, 10]). In the following we shall take the pion multiplicity dN_π/dy as 624 at SPS energies, 1735 at RHIC energies, and 5624 at LHC energies for central collisions involving two lead nuclei. This leads to T_0=320 MeV and τ_0=0.206 fm/c at the SPS, T_0=532 MeV and τ_0=0.124 fm/c at the RHIC, and T_0=958 MeV and τ_0= 0.069 fm/c at the LHC. The transition temperature has been taken as 160 MeV.

D. Results at the SPS, RHIC, and LHC

In Fig. 3 we have shown our results for the photons radiated due to the action of the color-confining mechanism on the quarks. The photons due to the Compton plus annihilation processes during the QGP phase of the interacting system are also shown. A central collision of two lead nuclei at SPS energies is considered.

We see that at the lowest k_T the photons radiated due to the color-confining mechanism operating on the quarks are comparable to those due to the Compton plus annihilation processes. However, at larger k_T they are only of the order of 10–20 % of the photons emitted due to the latter process. Similar results are seen at RHIC and LHC energies (see Figs. 4 and 5). The different transverse momentum dependence of the spectra is remarkable.

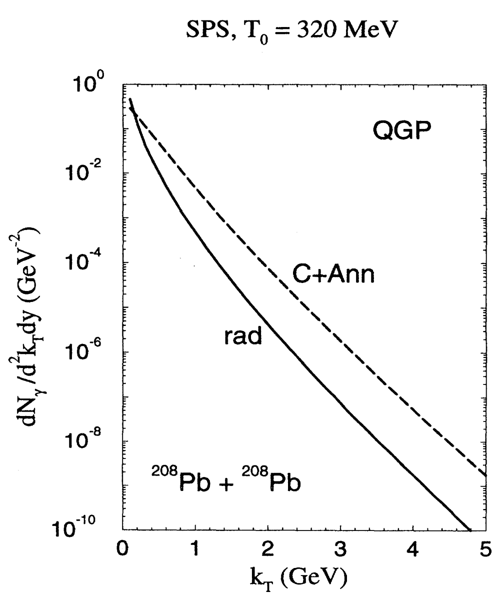

FIG. 3. Photon spectrum at SPS energies for central collisions involving two lead nuclei. The solid curve gives the results for the photons radiated due to the color-confining process considered here, and the dashed curve gives the photons due to the Compton plus annihilation processes.

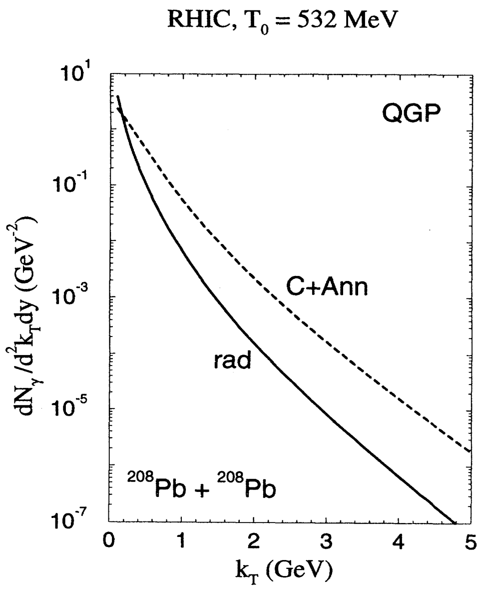

FIG. 4. Same as Fig. 3, at RHIC energies.

E. Consequences of a larger mass for the quarks

In order to reach Eq. (2), the current quark mass and *not* the constituent quark mass has been used. Thus the results given so far provide the most optimistic estimate for the radiation considered here.

However, if we choose the thermal mass ($\approx \sqrt{\frac{2\pi}{3}}\,\alpha_s\,T$) or the constituent mass ($\approx$ 310 MeV) for the quarks, the situation will alter drastically. Thus, for an arbitrary value of quark mass, the radiated power is given by [11]

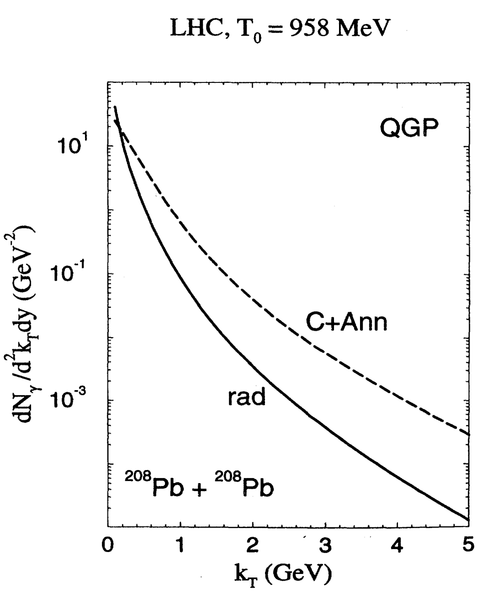

FIG. 5. Same as Fig. 3, at LHC energies.

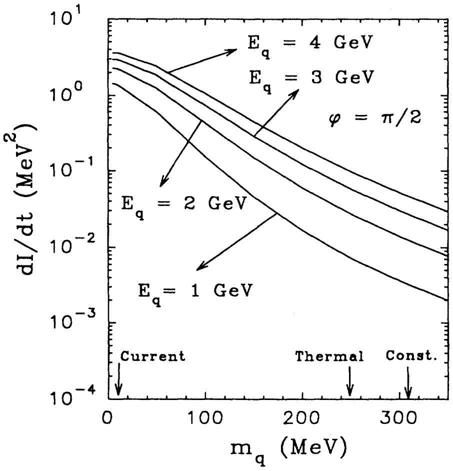

FIG. 6. The radiated power for some typical quark energies, as a function of the quark mass. Note the strong suppression as the mass changes from the current to the thermal ($T = 160$ MeV) and the constituent quark masses.

$$\frac{dI}{dt} = -\frac{e^2\, m^2\, \chi^2}{2\sqrt{\pi}} \int_0^\infty x\, dx\, \frac{4 + 5\chi\, x^{3/2} + 4\chi^2\, x^3}{(1 + \chi\, x^{3/2})^4}\, \mathrm{Ai}'(x)\ , \tag{21}$$

where

$$\chi = \frac{\sigma\, E_q\, \sin\phi}{m_q^3} \tag{22}$$

and $\mathrm{Ai}'(x)$ is the derivative of the Airy function. We have plotted the above in Fig. 6 for $\phi = \pi/2$ and some representative values of quark energies, as a function of the mass of the quark. The radiated power is seen to decrease drastically with increase in mass. For the thermal mass of the quark (simulating the many body effects) or the constituent quark mass, it is reduced by two orders of magnitude below the value for the current mass of the quarks.

III. DISCUSSION AND CONCLUSIONS

Before concluding, let us briefly examine some of the inputs into the treatment followed here. To begin with, one may perhaps argue that any transverse expansion of the system will destroy the sharp boundary considered here. Fortunately, this expansion is only minimal during the QGP phase [5].

It is conceivable that a system of hot hadrons produces a sort of attractive "mean field" which holds the hadrons [15]. Then hadrons trying to leave the system will experience a force similar in spirit to the case discussed here, and can radiate photons. However, the attractive potential in the case of hadrons is only of the order of a few tens of MeV [15], and it will also fall off rapidly as we move away from the surface. It is also necessary to realize that for the case of hadrons the temperature is not likely to be much more than 200 MeV, which limits their energy to the order of 600 MeV, which is only a factor of 4 more than the mass of pions, say. Thus it is felt that a vastly reduced force and an enhanced mass in the case of hadrons will largely suppress this radiation at larger k_T.

Finally, we have not included the consequences of the Landau-Pomeranchuk effect on the radiation of photons along the trajectory of the quark. This effect is known to be negligible [16] for photons having large momenta, and thus we feel that our results for k_T larger than 1 GeV will be unaltered.

We conclude that the radiation of photons due to the color-confining mechanism operating on the quarks near the surface of the quark-gluon plasma could be an interesting signature of the formation of the QGP, if we were justified in using the current mass of the quark, in which case they are seen to contribute at a level of 10–100 % of the photons due to the Compton plus annihilation processes at SPS, RHIC, and LHC energies in central collisions involving heavy nuclei. However, if we are to take the mass of the quark as the thermal mass or as constituent mass, this radiation is suppressed strongly, which is rather unfortunate as these photons are polarized [8], and have other interesting properties.

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