High-energy photons from expanding quark-gluon plasma and hot hadronic matter

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The emission of high-energy photons from an expanding quark-gluon plasma (QGP), which undergoes a first-order phase transition to a hot hadronic gas before freeze-out, is evaluated. Compton and annihilation processes in the quark-gluon plasma and an exhaustive array of $hh \to h\gamma$ reactions as well as decays in hot hadronic matter are considered. We find that, if the initial temperature of the QGP is more than twice the critical temperature, then, beyond a transverse momentum of about 2–3 GeV, the photons from the QGP outshine those having their origin in hadronic matter. We further note that the increase in degrees of freedom in the hot hadronic matter reduces the lifetime of the mixed phase.

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One of the most spectacular predictions of QCD is that with increasing density and/or temperature hadronic matter is dissolved into a quark-gluon plasma (QGP). The hope that such temperatures and densities may be realized in ultrarelativistic heavy-ion collisions has led to extensive experimental and theoretical efforts to identify possible signatures of the quark-gluon plasma. It is generally accepted [1–7] that high-energy photons originating from the annihilation process $q\bar{q} \to g\gamma$ and QCD Compton process $qg \to q\gamma$ and $\bar{q}g \to \bar{q}\gamma$ can be an excellent probe of the plasma, as they interact only electromagnetically with the medium and consequently have a very long mean free path. Thus the photons retain the memory of the early times of the QGP rather effectively.

It is important, however, to study the characteristics of these photons in order to distinguish them from the photons which have their origin in direct QCD processes, in hadronic matter, created during the cooling and hadronization of the plasma and in the hadronic decay after the freeze-out. Thus it is generally felt [8] that the photons originating from hadronic matter dominate in the region of small transverse momenta ($p_T < 2$ GeV) and the direct QCD photons do so at large p_T (~ 4 GeV or higher), leaving open a window around $p_T \sim 2$–4 GeV, where the photons radiating from the QGP dominate. Obviously, the precise location of this window is determined by the initial temperature of the plasma and the dynamics of the space-time evolution.

These aspects of the photons as a signature of the QGP have been brought into sharp focus by the recent work of Kapusta, Lichard, and Seibert [9], who have evaluated the rates for emission of photons from QGP and hadrons to the same order and found them to be nearly identical at a given temperature. We find in the present work that by taking into account the hydrodynamic expansion of the plasma in space and time, the photons from the QGP dominate over those from hadronic matter for $p_T > 2$–3 GeV if the initial temperature of the plasma is twice the critical temperature (T_c) or more.

We would like to add a word of caution at this stage. The results of Kapusta, Lichard, and Seibert are often misinterpreted to imply the inadequacy of photons as a signature of the QGP. In the context of ultrarelativistic heavy-ion collisions, however, the plasma formed has to undergo a dynamical evolution with time governed by an isentropic and *not* isothermal expansion. We show in this work that taking account of the entropy conservation during the evolution leads to a separation of the kinematical region where contributions from the hadronic phase are overshadowed by those from the QGP part.

We start by noting that the net rate [9] from hard and soft momentum transfers in $q\bar{q} \to g\gamma$ and $q(\bar{q})g \to q(\bar{q})\gamma$ processes is given by

$$E\frac{dR}{d^3p} = \frac{5}{9}\frac{\alpha\alpha_s}{2\pi^2}T^2 e^{-E/T}\ln\left|\frac{2.912}{g^2}\frac{E}{T}\right| , \tag{1}$$

where we have taken $g^2=5$ with $\alpha_s=0.4$ and considered only u and d quarks.

The rate for the hadronic reaction $hh \to h\gamma$ is given by

$$E\frac{dR_i}{d^3p} = \frac{1}{2}\frac{\text{deg}}{(2\pi)^3}\int \frac{d^3p_1}{(2\pi)^3 2E_1}\frac{d^3p_2}{(2\pi)^3 2E_2}\frac{d^3p_3}{(2\pi)^3 2E_3} \times f_1(E_1)f_2(E_2)[1+f_3(E_3)] \times (2\pi)^4\delta(p_1^\mu+p_2^\mu-p_3^\mu-p^\mu)|M_i|^2 \tag{2}$$

("deg" denotes degeneracy), where M_i, the amplitude, is obtained from

$$\frac{d\sigma_i}{dt} = \frac{|M_i|^2}{16\pi\lambda(s,m_1^2,m_2^2)} , \tag{3}$$

with

$$\lambda(a,b,c) = a^2+b^2+c^2-2(ab+bc+ca) \tag{4}$$

and the differential cross sections for the reactions $\pi\pi \to \rho\gamma$, $\pi\rho \to \pi\gamma$, $\pi\pi \to \eta\gamma$, and $\pi\eta \to \pi\gamma$ are taken from Kapusta, Lichard, and Seibert [9]. The purely electromagnetic annihilation $\pi^+\pi^- \to \gamma\gamma$ is considered for the sake of completeness. We have also considered the vector decay $\omega^0 \to \pi^0\gamma$, $\rho^0 \to \pi^+\pi^-\gamma$, as well as $\pi^0 \to \gamma\gamma$.

In Fig. 1 we have plotted our results for the comparison of rates for photons from QGP at $T=160$ MeV and those from the reactions $hh \to h\gamma$, as well as vector decays. We find that the rates are quite similar at $T=160$ MeV as well, as was found at $T=200$ MeV earlier [9].

We also show the emission rate of photons from hadronic matter consisting of an ideal gas of *pions*, only, when the $\pi^+\pi^- \to \gamma\gamma$ and the $\pi^0 \to \gamma\gamma$ processes contribute. This result provides a justification for the neglect of the hadronic contribution to the photon count rate before freeze-out for hadronic matter consisting of pions only [10]. Interestingly enough, we find that the contribution of the $\pi^0 \to \gamma\gamma$ channel is about 5 times that for the $\pi^+\pi^- \to \gamma\gamma$ before freeze-out.

The results of Fig. 1 have relevance for the mixed phase, when the quark matter and hadronic matter coexist at the same temperature. In favorable cases in heavy-ion collisions at ultrarelativistic energies, the QGP will be formed at an initial temperature T_i (larger than the critical temperature T_c). The QGP will expand and cool until it undergoes a phase transition to a hadronic gas, which will further expand until it freezes out. Thus the dynamics of the expansion will play a crucial role in providing the window for the photons from the QGP.

We take the bag-model equation of state for quark matter for simplicity and write

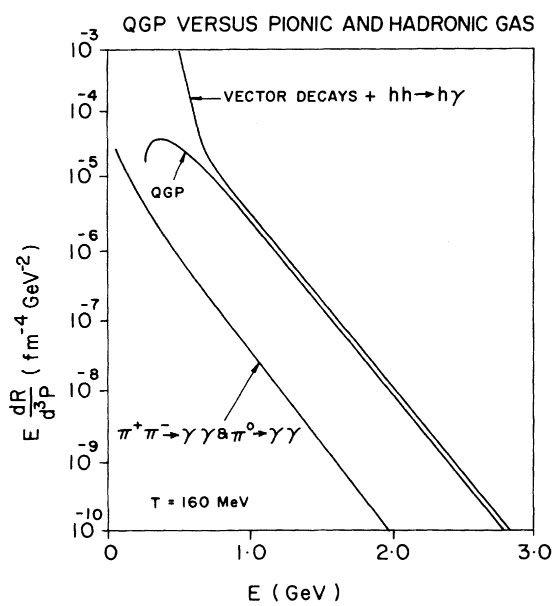

FIG. 1. Emission rate per unit volume for photons from a quark-gluon plasma and a hot hadronic gas at $T=160$ MeV. For the QGP, Compton ($qg \to q\gamma$) and annihilation ($q\bar{q} \to \gamma g$) processes are considered. The reactions $\pi\pi \to \rho\gamma$, $\pi\rho \to \pi\gamma$, $\pi\pi \to \eta\gamma$, and $\pi\eta \to \pi\gamma$ and decays $\omega^0 \to \pi^0\gamma$ and $\rho^0 \to \pi^+\pi^-\gamma$ have been considered for the case of the hadronic gas. In addition, the emission from a pionic gas are also shown for comparison.

$$P_Q = aT^4 - B, \quad \epsilon_Q = 3aT^4 + B, \quad s_Q = 4aT^3, \tag{5}$$

with

$$a = \frac{\pi^2}{90}(2\times 8 + \tfrac{7}{8}\times 2\times 2\times 3\times 2), \tag{6}$$

for an ideal gas of (u and d) quarks and gluons. The hadronic matter is considered to be a mixture of noninteracting π, ρ, ω, and η mesons, such that

$$\epsilon_H = \sum_i \frac{g_i}{(2\pi)^3} \int \frac{d^3p_i (p_i^2+m_i^2)^{1/2}}{\exp\{(p_i^2+m_i^2)^{1/2}/T\}-1}, \tag{7a}$$

$$P_H = \sum_i \frac{g_i}{(2\pi)^3} \int \frac{p_i^2 d^3p_i}{3(p_i^2+m_i^2)^{1/2}} \frac{1}{\exp\{(p_i^2+m_i^2)^{1/2}/T\}-1}, \tag{7b}$$

and

$$s_H = \frac{\epsilon_H + P_H}{T} \equiv 4a_{\text{eff}}(T)T^3 = 4\pi^2 g_{\text{eff}}(T)\frac{T^3}{90}, \tag{8}$$

where g_i is equal to 3, 9, 3, and 1 for π, ρ, ω, and η, respectively.

Thus we can visualize the finite-mass hadronic matter as considered here as consisting of massless hadrons having an effective degeneracy $g_{\text{eff}}(T)$ such that the expression for s_H is satisfied and use it later to obtain lifetimes, etc., of the mixed and hadronic phases. We note that $g_{\text{eff}}=3$ for an *ideal gas of (massless) pions* usually taken to describe the hot hadronic gas. We insist that the high temperature of the hadronic matter created during the hadronization does not allow us to treat the hadronic matter as an ideal gas of (noninteracting) pions. We further recall that an interacting pion gas can be shown to have properties similar to a noninteracting ideal gas consisting of pions and resonances [11]. The velocity of sound (c_s) is then obtained as

$$c_s^{-2} = \frac{T}{s}\frac{ds}{dT} = \left[\frac{1}{a_{\text{eff}}(T)}\frac{da_{\text{eff}}(T)}{dT}T + 3\right]. \tag{9}$$

We find that $c_s^2 \approx 0.238$ over the relevant temperature range for hadronic matter (see also [12]). This provides sufficient support for the methodology adopted here for handling the hadronic phase.

Realizing that, we can write [6]

$$T_c = \left[\frac{B}{a - a_{\text{eff}}}\right]^{1/4}.$$

We find that the change in the value of the critical temperature due to the change in the composition of the hadronic matter envisaged here is only a few percent for a given bag constant B. Thus the richness of the hadronic matter considered here does not change the lifetime of the QGP to any significant extent.

We use Bjorken hydrodynamics [13] with only longitudinal isentropic expansion so that $s\tau=$const through the QGP and the mixed and hadronic phases.

Now the result for emission of photons from an expanding QGP up to the freeze-out point is given by

$$\frac{dN}{d^2p_T dy} = \pi R_A^2 \int \tau\, d\tau\, d\eta \left\{ \left[E\frac{dR}{d^3p}\right]_{\mathrm{QGP}} \Theta(s-s_Q) + \left[\left[E\frac{dR}{d^3p}\right]_{\mathrm{QGP}}\left[\frac{s-s_H}{s_Q-s_H}\right] + \left[\frac{s_Q-s}{s_Q-s_H}\right]\left[E\frac{dR}{d^3p}\right]_{\mathrm{had}}\right]\Theta(M) + \left[E\frac{dR}{d^3p}\right]_{\mathrm{had}}\Theta(s_H-s)\right\}, \tag{10a}$$

where

$$\Theta(M)=\Theta(s_Q-s)\Theta(s-s_H) . \tag{10b}$$

Here s_Q and s_H are the entropy densities at the critical temperature.

The basic rates for the emission of photons from the QGP and hadronic matter are taken from the relations given above. A Monte Carlo program has been written to perform the entire integration numerically. The photons radiating from quark matter evidently originate from a pure QGP phase as well as the QGP part of the mixed phase. Similarly, the sum of the contributions of the hadronic phase and the hadronic part of the mixed phase would be called the contribution of hadronic matter.

In Fig. 2 we show our results for two typical initial temperatures $T_i=1.5T_c$ and $2T_c$ for a $^{208}\mathrm{Pb}+^{208}\mathrm{Pb}$ system. The initial formation time of the QGP is taken as $\tau_i=1$ fm/c. These along with a value for $T_c=160$ MeV here uniquely determine the particle rapidity density through its relation with the entropy density [6,14] (which is also shown on the figure for easy reference). Results for $T_c=200$ MeV are given in Fig. 3.

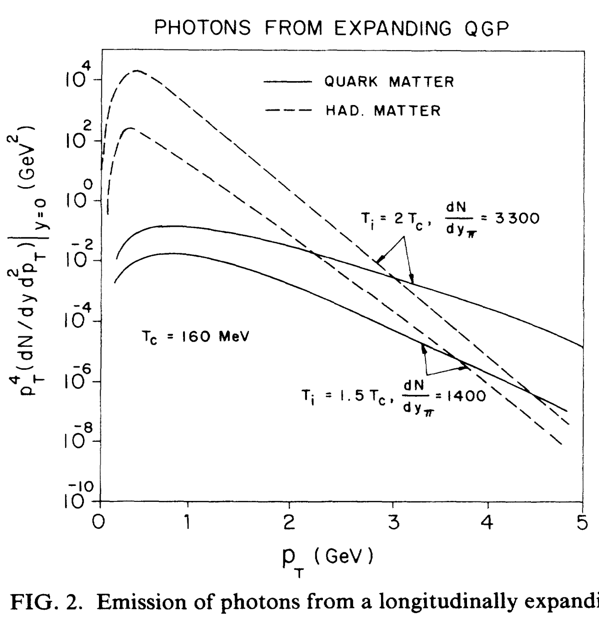

FIG. 2. Emission of photons from a longitudinally expanding QGP up to the time of freeze-out, considering a first-order phase transition at $T_c=160$ MeV. The formation time of the plasma, τ_i, is taken as 1 fm/c (see text). The solid curves give the contribution of quark matter, and the dashed curves give the results for hadronic matter. Two typical initial temperatures are chosen as $1.5T_c$ and $2T_c$, which along with τ_i uniquely determine the particle rapidity density (dN/dy_π) for the $^{208}\mathrm{Pb}+^{208}\mathrm{Pb}$ system considered here. The freeze-out temperature $T_f=140$ MeV.

We see that when the initial temperature is $1.5T_c$ the photons from quark matter outshine the photons from hadronic matter for p_T larger than 3–3.5 GeV. The situation improves substantially for the larger initial temperature considered. For the most optimistic cases considered likely to be realized at CERN Large Hadron Collider (LHC) and BNL Relativistic Heavy Ion Collider [14], the photon signal from quark matter considerably outshines the background from hadronic matter beyond $p_T \sim 2.5$–3.5 GeV. This is a very convenient window as the direct QCD photons do not contribute significantly here [15]. In fact, one can argue that for higher energies at RHIC and/or LHC the hadronic background at larger p_T could perhaps be reliably evaluated using perturbative estimates of minijets [16–18], making the photon signals cleaner indeed.

It may be further argued that $\tau_i=1$ fm/c assumed here is definitely not the only acceptable [19–23] choice. In fact, the uncertainty principle provides that if the initial temperature of the plasma is T_i, then $\tau_i \sim 1/3T_i$ as the corresponding energy scale is $\sim 3T_i$ [16,24]. We also realize that, because of the onset of minijet phenomenon at larger incident energies, the average transverse momentum of the particle increases [25], which should again imply a smaller τ_i [24]. This along with

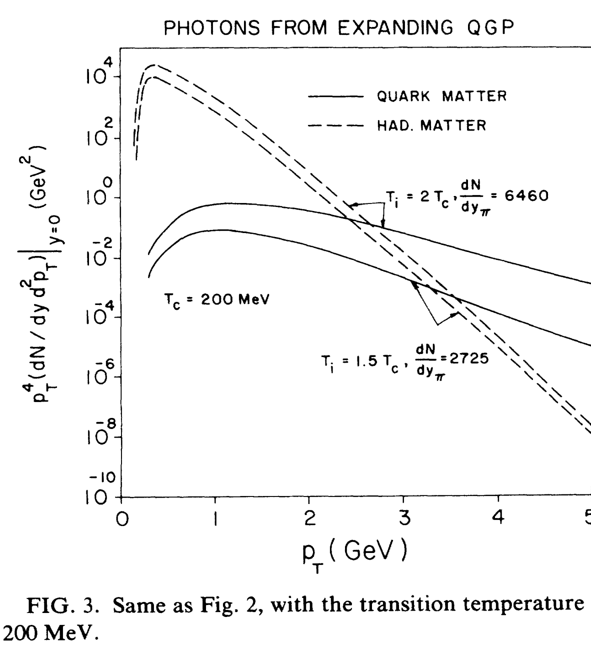

FIG. 3. Same as Fig. 2, with the transition temperature T_c as 200 MeV.

$dN/dy_\pi \sim T_i^3 \tau_i$ can give much higher initial temperatures for the same multiplicity densities as considered here [19]. We note that the contributions of the mixed and hadronic phases remain unaltered when T_c and (dN/dy_π) are fixed. The results for higher initial temperatures are then easily visualized from Figs. 2 and 3 by having an increased QGP contribution. We would like to remark here that this favorable shift in the window for QGP photons to a lower p_T has been possible because the reduced lifetime of the mixed phase as envisaged here.

In order to understand this, we note that during the mixed phase the quark matter and hadronic matter are at the same temperature T_c, and hence the contribution from the mixed phase is proportional to $(\tau_H^2 - \tau_Q^2)$, where τ_H gives the end of mixed phase and τ_Q the end of the QGP phase in the rest frame of the plasma. Thus, if the hadronic matter consists of *only pions,*

$$\tau_H = \frac{a}{a_\pi} \tau_Q \ , \tag{11}$$

whereas for the present case of hadronic gas consisting of π ρ, ω, and η mesons,

$$\tau_H = \frac{a}{a_{\text{eff}}(T_c)} \tau_Q \ . \tag{12}$$

Referring to Eq. (9), we note that g_{eff} is 4.9 at 160 MeV and 6.3 at 200 MeV. This leads to a decrease by a factor of 1.7 and 2.3, respectively, in the lifetimes of the mixed phases consisting of heavier mesons compared with the case of a pionic gas. The freeze-out time for our case is also reduced correspondingly [19]. Noting that the "richness" of hadronic matter does not affect the lifetime of the pure QGP phase, we see immediately that a realistic description of hadronic matter considerably reduces the hadronic background to the QGP signal. These arguments also hold for the dilepton spectrum.

It could be of interest to consider a scenario when there is no phase transition. We start by finding the temperature of a "hot hadronic gas" which has *the same initial entropy density* as the QGP considered here upon formation, i.e., at τ_i, for a meaningful comparison. In Table I we give the result for such an exercise for the four cases considered here.

We see that the hadronic gas scenario without a phase transition has to have very large initial temperatures indeed in order to provide the reasonable particle rapidity densities considered here. Even though we find it very difficult to imagine a hot hadronic gas at such large temperatures, we have made such a "mindless" comparison as a purely academic exercise of the count rate for photons for the two scenarios. In Fig. 4 we consider the case when the QGP is formed at an initial temperature of $T_i = 1.5T_c$ with $T_c = 160$ MeV. We see that, because of the larger initial temperature, the slope of the spectrum in the hot hadronic gas scenario is smaller. The difference would be even larger for the other cases in Table I. However, we strongly feel that too much importance should not be attached to this result, as we have used ground-state hadronic properties (obtained at $T=0$) even at such higher temperatures. The temperature dependence of these quantities will only introduce a further uncertainty.

The conclusions of the present work should be qualified with the following observations. We have not included the bremsstrahlung processes leading to the production of the photons in the QGP and hadronic matter. However, hadronic bremsstrahlung may give photons having a $p_T \approx 100$ MeV [15]. The photons from bremsstrahlung in quark matter, on the other hand, make only a negligible contribution beyond $p_T \approx 1.5T_i$ [6,19], where the photons from Compton and annihilation processes dominate.

One might question the neglect of the contribution of strange quarks. It is felt [26] that if we include the contribution of s quarks in quark matter, we should also include the contribution of strange mesons in hadronic matter. It is quite likely then that the two contributions to the basic rates are rather similar and then our basic findings will not be altered.

In all the results reported here, the temperature dependence [27] of the masses and coupling constants in the hadronic phase has been ignored. It is generally felt [26], though, that the present estimates of these dependencies have their largest uncertainty around the critical temperature. Extracting reliable estimates in this region, which is of interest in the present context, is a nontrivial task,

TABLE I. Comparison of the initial temperature in a QGP scenario with that for a hadronic gas scenario, corresponding to identical entropy densities at the initial time $\tau_i = 1$ fm/c.

	T_i in the QGP scenario	T_i in the hot hadronic gas scenario (MeV)
$T_c = 160$ MeV	$T_i = 1.5T_c$	362
	$T_i = 2T_c$	435
$T_c = 200$ MeV	$T_i = 1.5T_c$	415
	$T_i = 2T_c$	513

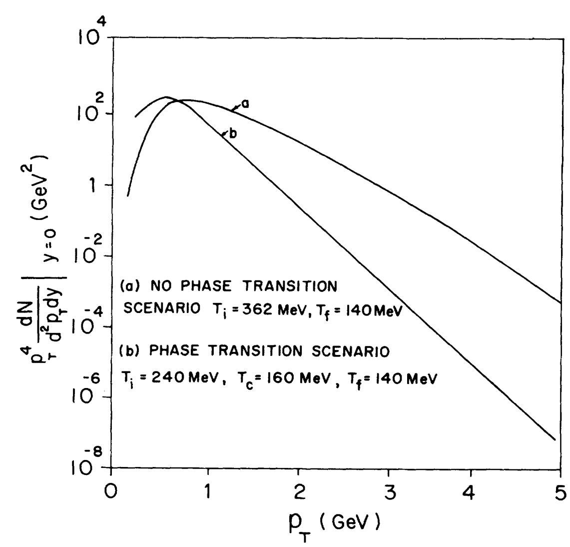

FIG. 4. Comparison of the count rate of the photons from an expanding plasma of (a) hot hadrons and (b) the QGP at the same initial density. We have used $\tau_i = 1$ fm/c, which then corresponds to $dN/dy_\pi = 1400$.

but work along these lines is in progress.

The assumption of only a longitudinal expansion up to the time of freeze-out is not quite justifiable. However, it is known [28,29] that transverse expansion starts playing a significant role only after the rarefaction wave hits the center of the plasma, i.e., after a proper time $\sim R_A/c_s$. Thus transverse expansion is not significant for emissions from the QGP phase. The reduction of the lifetime of the mixed phase envisaged here, because of the increased hadronic degrees of freedom, reduces its importance for the duration of the mixed phase as well. The photons from the hadronic phase may "experience" a transverse kick, though, and appear in the p_T window identified here, thus polluting the signals of the QGP to some extent. On the other hand, the lifetimes of the mixed and hadronic phases will be reduced because of transverse expansion, thus reducing their contribution. These competing aspects along with the possible inclusion of strange and heavier mesons are under investigation [30].

In brief, we have obtained cross sections for the emission of high-energy photons from an expanding quark-gluon plasma by convoluting their rates, as a result of the Compton and annihilation processes in quark matter and an exhaustive array of hadronic reactions and decays in hadronic matter with the space-time history of the plasma. We find that if the initial temperature of the QGP is more than $2T_c$, then for $p_T > 2-3$ GeV the photons having their origin in quark matter dominate over those originating from hadronic matter. Such a scenario is crucially dependent on the reduction of the lifetime of the mixed phase containing heavy mesons.

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[1] E. L. Feinberg, Nuovo Cimento A **34**, 39 (1976).
[2] E. V. Shuryak, Yad. Fiz. **28**, 796 (1978) [Sov. J. Nucl. Phys. **28**, 408 (1978)].
[3] K. Kajantie and H. I. Miettinen, Z. Phys. C **9**, 341 (1981).
[4] F. Halzen and H. C. Liu, Phys. Rev. D **25**, 1842 (1982).
[5] Bikash Sinha, Phys. Lett. **128B**, 91 (1983).
[6] R. C. Hwa and K. Kajantie, Phys. Rev. D **32**, 1109 (1985).
[7] G. Staadt, W. Greiner, and J. Rafelski, Phys. Rev. D **33**, 66 (1986).
[8] Sibaji Raha and Bikash Sinha, Int. J. Mod. Phys. A **6**, 517 (1991), and references therein.
[9] J. Kapusta, P. Lichard, and D. Seibert, Phys. Rev. D **44**, 2774 (1991).
[10] D. K. Srivastava and Bikash Sinha, Phys. Lett. B **261**, 1 (1991).
[11] G. M. Welke, R. Venugopalan, and M. Prakash, Phys. Lett. B **245**, 137 (1990).
[12] E. Shuryak, Yad. Fiz. **16**, 395 (1973) [Sov. J. Nucl. Phys. **16**, 220 (1973)].
[13] J. D. Bjorken, Phys. Rev. D **27**, 140 (1983).
[14] U. Heinz, P. Koch, and B. Friman, in *Proceedings of the ECFA Large Hadron Collider Workshop,* Aachen, Germany, 1990, edited by G. Jarlskog and D. Rein (CERN Report No. 90-10, Geneva, 1990), Vol. II, p. 1079.
[15] M. Neubert, Z. Phys. C **42**, 231 (1989).
[16] D. K. Srivastava, Bikash Sinha, M. Gyulassy, and X. N. Wang, Phys. Lett. B **276**, 285 (1992).
[17] K. Kajantie, P. V. Landshoff, and J. Lindfors, Phys. Rev. Lett. **59**, 2527 (1987).
[18] M. Jacob and P. V. Landshoff, Mod. Phys. Lett. **1**, 657 (1986).
[19] D. K. Srivastava and Bikash Sinha, J. Phys. G (to be published).
[20] E. V. Shuryak, *The QCD Vacuum, Hadrons and the Superdense Matter* (World Scientific, Singapore, 1988), p. 323.
[21] M. Gyulassy and T. Matsui, Phys. Rev. D **29**, 419 (1984).
[22] R. C. Hwa and K. Kajantie, Phys. Rev. Lett. **56**, 696 (1986).
[23] L. McLerran and T. Toimela, Phys. Rev. D **31**, 541 (1985).
[24] J. Kapusta, L. McLerran, and D. K. Srivastava, Phys. Lett. B **283**, 145 (1992); Theoretical Physics Institute, University of Minnesota, Report No. TPI-MINN/3-T (unpublished).
[25] X. N. Wang and M. Gyulassy, Phys. Rev. D **44**, 3501 (1991).
[26] J. Kapusta (private communication).
[27] G. E. Brown and Mannque Rho, Phys. Rev. Lett. **66**, 2720 (1991).
[28] K. Kajantie, M. Kataja, L. McLerran, and P. V. Ruuskanen, Phys. Rev. D **34**, 811 (1986).
[29] Bhaskar Datta, Sibaji Raha, and Bikash Sinha, Nucl. Phys. **A490**, 733 (1988).
[30] J. Alam, S. Chakrabarty, D. K. Srivastava, Bikash Sinha, and Sibaji Raha (unpublished).