Physics and Signatures of the Quark-Gluon Plasma*

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This is a critical review of the various observables that have been proposed to signal the change from dense hadronic matter to a quark-gluon plasma at high temperature or baryon density. I discuss current models of quark-gluon plasma formation in relativistic heavy ion collisions and analyze the virtues and ambiguities of various signatures.

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I. INTRODUCTION

The central goal of relativistic heavy ion physics is the discovery of the new state of strongly interacting matter that is called the quark-gluon plasma and the study of its properties [1]. In order to accomplish this goal, we need reliable signals for the formation of such a state. This review tries to capture the essential ideas and the current status of the theoretical studies on the most promising quark-gluon plasma signals. The reader is encouraged to also consult the earlier reviews of Kajantie and McLerran [2] and Singh [3] and the Proceedings of the Strasbourg Workshop [4], which contain detailed accounts of various signatures.

In order to shed light on the connections between the many proposed quark-gluon plasma signatures I will here group them in five categories, according to the physical properties of superdense hadronic matter to which they are sensitive. These are: thermodynamic variables measuring the equation of state; probes for chiral symmetry restoration; probes of the color response function; probes of the electromagnetic response function; and “exotic” signatures of the quark-gluon plasma. Some signals, e.g. strangeness production, may be sensitive to more than one particular aspect. In this case they have been grouped in the section which I considered to be most relevant.

The review begins with a brief survey of our current picture of the dynamical and structural properties of the quark gluon plasma, of its creation and evolution. It then proceeds to the analysis of the various proposed signals and concludes with an assessment of their merits and ambiguities.

II. BRIEF SURVEY OF QUARK-GLUON PLASMA PHYSICS

Experiments at Brookhaven and CERN have provided compelling evidence that highly energetic collisions between heavy ions proceed through an intermediate phase that is characterized by energy densities in excess of 1 GeV/fm$^3$ [5]. This can be concluded from the analysis of rapidity distributions of particle multiplicities and transverse energy, as well as from the observed depletion of charmonium production which requires the presence of a dense cloud of comovers. It is less certain at this point whether this dense hadronic phase can be considered as locally quasi-thermal, but theoretical studies indicate that local thermalization of matter in relativistic heavy ion collision actually proceeds more rapidly as the bombarding energy grows. Estimates of the thermalization time, which are based on perturbative QCD and the parton model [6–8], as well as recent nonperturbative studies of lattice gauge theory [9] yield values below 0.5 fm/c at collider energies (above 100 GeV per nucleon in the c.m. system).

The problem of the equation of state of superdense hadronic matter is a fundamental issue of quantum chromodynamics. Theory tells us that it makes little sense to base a description of strongly interacting matter at energy densities far above 1 GeV/fm$^3$ on a picture beginning with noninteracting hadrons. Rather it appears to be appropriate to start from the fundamental constituents of hadrons, i.e. quarks and gluons, because interactions between quarks and gluons become weak at short distances. This approach, anticipated in pre-QCD days by P. Carruthers [10], was pioneered by Collins and Perry [11], Baym and Chin [12], McLerran et al. [13], Shuryak [14], and others. It is complicated by the fact that the QCD plasma retains essentially nonperturbative aspects even at the highest densities because of the existence of collective modes (plasmons) and the absence of screening of long-range static color-magnetic forces in perturbation theory. Although plasma properties can be classified according to orders of the QCD coupling constant $g$ (not $\alpha_s = g^2/4\pi$), $g \ll 1$ does not hold even at the highest conceivable energies ($g \approx 0.5$ at the Planck scale).

A. Structure of the Quark-Gluon Plasma

The theory of the equation of state of the gluon plasma is conceptually quite simple, because it is directly based on the fundamental QCD Lagrangian

$$\mathcal{L}_{QCD} = \frac{1}{4} \sum_a F^{a\mu\nu} F_{a\mu\nu} + \sum_{f=1}^{N_f} \overline{\psi}(i\gamma^\mu \partial_\mu - g\gamma^\mu A^a_\mu \frac{\lambda^a}{2} - m_f)\psi$$  \hspace{1cm} (1)

where the subscript $f$ denotes the various quark flavors $u, d, s, c$, etc., and the nonlinear glue field strength is given by

$$F^{a\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + gf_{abc} A^b_\mu A^c_\nu.$$  \hspace{1cm} (2)

QCD predicts a weakening of the quark-quark interaction at short distances (or high momenta $Q^2$), because the one-loop series for the gluon propagator yields a running coupling constant
\[ g^2(Q^2) = \frac{16\pi^2}{(11 - \frac{2}{3} N_f)} \ln(Q^2/\Lambda^2) \xrightarrow{Q^2 \to \infty} 0, \]

where \( N_f \) is the number of active quark flavors. The QCD scale parameter \( \Lambda \) is now quite well determined \[15] to have the value \( \Lambda \approx 200 \text{ MeV} \).

Assuming that interactions of quark and gluons are sufficiently small at high energy, the energy density \( \varepsilon \) and pressure \( P \) of a quark-gluon plasma at temperature \( T \) and quark chemical potential \( \mu_f \) can be calculated by thermal perturbation theory. Neglecting quark masses in first-order perturbation theory, the equation of state is

\[ \varepsilon = \left( 1 - \frac{15}{4\pi \alpha_s} \right) \frac{8\pi^2}{15} T^4 + N_f \left( 1 - \frac{50}{21\pi \alpha_s} \right) \frac{7\pi^2}{10} T^4 + \sum_f \left( 1 - \frac{2}{\pi \alpha_s} \right) \frac{3}{\pi \mu_f^2} \left( \pi^2 T^2 + \frac{1}{2} \mu_f^2 \right) + B, \]

where \( B \) is the difference between the energy density of the perturbative and the nonperturbative QCD vacuum (the bag constant). The pressure is given by \( P = \frac{1}{3}(\varepsilon - 4B) \); the entropy density is \( s = (\partial P/\partial T)_\mu \). One observes that \( \langle \rangle \) is essentially the equation of state of a gas of massless particles with corrections due to the QCD trace anomaly and perturbative interactions. These are always negative, reducing the energy density at given temperature by about a factor two when \( \alpha_s = 0.5 \). The perturbative plasma phase becomes unstable at temperatures below \( T_c \approx 0.8B^{1/4} \approx 170 \text{ MeV} \) (for \( \mu = 0 \)), for standard values of the bag constant.

More reliable predictions concerning this phase transition can at present only be obtained by numerical simulations of the QCD equation of state on a finite discretized volume of space-time, usually referred to as lattice gauge theory. In this approach \[16\] one approximately calculates the partition function for a discretized version of the QCD Lagrangian \[19\] by Monte-Carlo methods. In principle, this technique should accurately describe the quark-gluon plasma as well as the hadronic phase but, in practice, its accuracy especially at low temperature is severely limited by finite size effects and other technical difficulties. Where the numerical results are most reliable, i.e. for the pure gluon theory without dynamical quarks, the calculations predict a sudden jump in the energy density at a certain temperature while the pressure rises more gradually.

When dynamical quarks are added, the role of the quarks becomes less clear for two reasons. One is that the calculations involving fermion fields on the lattice are much more time consuming, and hence the numerical results are less statistically meaningful and reliable. Moreover, the definition of quark confinement becomes rather fuzzy in the presence of light quarks, because the color flux tube between two heavy quarks can break by creation of a light quark pair: \( \bar{Q}Q \rightarrow (\bar{q}q)(\bar{q}Q) \). E.g., highly excited states of charmonium can break up into a pair of D-mesons. Thus, in the calculations the \( \bar{q}Q \) potential does not rise linearly with distance, but is effectively screened.

For massless dynamical quarks there exists a new order parameter, the quark-antiquark condensate in the vacuum \langle 0 | \bar{q}q | 0 \rangle\. When it assumes a nonzero value, chiral symmetry is spontaneously broken, as can be seen as follows: The scalar quark density has the chiral decomposition \( \bar{q}q = \bar{q}_L \gamma_5 q_L + \bar{q}_R \gamma_5 q_R \), hence the broken vacuum state contains pairs of quarks of opposite chirality. A left-handed quark, say, can therefore annihilate on a left-handed antiquark in the vacuum condensate, liberating its right-handed partner. This process is perceived as change of chirality of the free quark, which provides exactly the same effect as a nonvanishing quark mass. However, in reality the mass of the light quarks \( u, d \) is nonzero, and the chirality of a light quark is never exactly conserved, even when the quark condensate vanishes.

All one can do, therefore, is to look for sudden changes in the distance at which color forces are screened, or in the quark condensate. If these are discontinuous, one deals with a phase transition, otherwise with a possibly rapid, but continuous change of internal structure as it occurs, e.g., in the transformation of an atomic gas into an electromagnetic plasma. The identification of the nature of the phase change is complicated by finite size effects. The best published results, by the Columbia group \[17\][18\], for a \( 16^3 \times 4 \) lattice indicate a surprisingly strong dependence of the phase diagram on the magnitude of the strange quark mass. For the physical mass \( m_s \approx 150 \text{ MeV} \) the presence of a discontinuity has not been established, but a very rapid change in the energy density over a small temperature range (about 10 MeV) is clearly seen in the numerical results.

In the next order of the coupling constant \( \alpha_s \), the gluon energy density is found to be infrared divergent. The physical reason for this divergence is that gluon and quark degrees of freedom develop an effective mass, which leads to screening of long-range color-electric forces. Technically, the screening mass is obtained by summing an infinite chain of one-loop insertions in the gluon propagator, which can be summed analytically and yield a contribution to the gluon energy of order \( \alpha_s^{3/2} \) with a rather large coefficient \[13\].

One obtains more insight into the properties of the interacting quark-gluon plasma by considering the gluon propagator \( D_{\mu\nu}(k) \). Because of gauge invariance, \( k^\mu D_{\mu\nu}(k) = 0 \), it can be decomposed into longitudinal and transverse parts, which are scalar functions of the variables \( \omega = k^0 \) and \( k = |k| \). These are most conveniently written in the form
\[ D_L(\omega, k) = \frac{1}{\varepsilon_L(\omega, k)k^2}, \] (5)

\[ D_T(\omega, k) = \frac{1}{\varepsilon_T(\omega, k)\omega^2 - k^2}, \] (6)

where the color-dielectric functions are given by:

\[ \varepsilon_L(\omega, k) = 1 + \frac{g^2 T^2}{k^2} \left[ 1 - \frac{\omega}{1k} \ln \left( \frac{\omega + k}{\omega - k} \right) \right]; \] (7)

\[ \varepsilon_T(\omega, k) = 1 - \frac{g^2 T^2}{2k^2} \left[ 1 - \left( 1 - \frac{k^2}{\omega^2} \right) \frac{\omega}{2k} \ln \left( \frac{\omega + k}{\omega - k} \right) \right]. \] (8)

Several things are noteworthy about (5–8). First they imply that static longitudinal color fields are screened:

\[ D_L(0, k) = \frac{1}{\varepsilon_L(0, k)k^2} = \frac{1}{k^2 + g^2 T^2}. \] (9)

The Debye length obviously is \( \lambda_D = (gT)^{-1} \). On the other hand, (3,8) show that static transverse (magnetic) color fields remain unscreened at this level of approximation. The static magnetic screening length is of higher order in the coupling constant: lattice gauge calculations [21] as well as analytical studies [22] have shown that \( \lambda_M = Cg^2T \) with \( C \approx 0.31 \) for SU(3) gauge theory.

For a finite frequency \( \omega \) the in-medium propagators (5,6) have poles corresponding to propagating collective modes of the glue field. The dispersion relation for the longitudinal mode:

\[ \varepsilon_L(\omega, k) = 0, \] (10)

called the plasmon, has no counterpart outside the medium. The analogous relation for the transverse mode:

\[ \varepsilon_T(\omega, k) = k^2/\omega^2 \] (11)

describes the effects of the medium on the free gluon. The behavior of both modes is remarkably similar. For \( k \to 0 \) they yield an effective plasmon mass

\[ \omega_L, \omega_T \xrightarrow{k \to 0} m_g^* = \frac{1}{\sqrt{3}} gT, \] (12)

whereas for large momenta (\( k \to \infty \)) one finds

\[ \omega_L(k) \to k, \quad \omega_T(k) \to \sqrt{k^2 + \frac{1}{2} g^2 T^2}. \] (13)

For plasma conditions realistically attainable in nuclear collisions (\( T \approx 300 \) MeV, \( \alpha_s \approx 0.3 \)) the effective gluon mass \( m_g^* \) is of the order of the temperature itself. We must conclude therefore, that the notion of almost free gluons (and quarks) in the high-temperature phase of QCD is quite far from the truth.

Let us discuss some consequences of these results:

1. The potential between two static color charges, such as two heavy quarks, is screened in the quark-gluon plasma phase. The Fourier transform of eq. 8 yields the potential

\[ V_{QQ}(r) \sim \frac{1}{r} e^{-r/\lambda_D} \] (14)

with screening length \( \lambda_D \approx 0.4 \) fm at \( T \approx 250 \) MeV. The screening of long-range color forces is, of course, the origin of quark deconfinement in the high-temperature phase. An important consequence, to be discussed later, is the disappearance of the bound states of a charmed quark pair (\( c\bar{c} \)) in the quark-gluon plasma [23].

2. The color screening at large distances cures most infrared divergences in scattering processes between quarks and gluons. A self-consistent scheme implementing this mechanism has been devised by Braaten and Pisarski [24].
involves the resummation of gluon loops involving gluons with momenta of order $gT$ and has been shown to be gauge invariant when vertex corrections are also taken into account.

3. The finite effective gluon mass $m_g^*$ leads to the suppression of long-wavelength gluon modes with $k \leq gT$ in the quark-gluon plasma. As a result, the pressure is reduced. We now have two mechanisms that can be responsible for $P < \frac{4}{3} \varepsilon$: an effective gluon mass $m_g^*$ and a nonvanishing vacuum energy $B$. Fits [27] to SU(3) lattice gauge theory results [26] show that probably both mechanisms are at work.

Whereas the depletion of long-wavelength modes in the quark-gluon plasma appears to be rather well established, its dynamical origin is by no means clear. If the above picture is correct, their population is simply suppressed by the collective mass of the colored plasmons On the other hand, it has also been speculated that low-momentum modes in the QCD plasma can survive at high temperature in spite of color screening [28] (but see [29]). A phenomenological model has been constructed which ascribes the different behavior of $\varepsilon(T)$ and $P(T)$ near to the postulated cluster structure of the quark-gluon plasma [30]. Several signals are expected to be very sensitive to the microscopic structure of the QCD plasma at low momenta, e.g. production of multistrange baryons or low-mass lepton pairs. A better understanding of the long-range structure of the quark-gluon plasma, and of the interpretation of the pertinent lattice-QCD results, is therefore of paramount importance.

B. Properties of the Quark-Gluon Plasma

1. Thermalization

All models of the formation of the quark-gluon plasma in nuclear collisions require information about its rate of thermalization. Does it thermalize sufficiently fast, so that a thermodynamical description makes sense? In the picture based on quasi-free quarks and gluons moving through the plasma, thermalization proceeds mainly via two-body collisions, where the color force between the colliding particles is screened. The technically easiest way of looking at this is to consider the quark (gluon) damping rate $\Gamma$, i.e. twice the imaginary part of the quark (gluon) self energy.

For quarks and gluons of typical thermal momenta ($p \approx T$) one finds [24] with the help of the techniques discussed in the previous section:

$$\Gamma_Q = -\frac{1}{2p} \text{Im}[\text{Tr}\Sigma(p)] \approx \frac{g^2}{3\pi} T \left(1 + \ln \frac{1}{\alpha_s}\right), \quad (15)$$

$$\Gamma_G = -2 \text{Im}[\omega_T(p)] \approx 0.54g^2T. \quad (16)$$

The larger factor for gluons simply reflects the fact that gluons have a higher color charge than quarks and therefore scatter more often.

One can argue that a better way to look at thermalization is to consider the rate of momentum transfer between particles, i.e. weighting the differential scattering cross section by $\sin^2 \theta$, where $\theta$ is the scattering angle. Here one finds [31]:

$$\Gamma_Q^{(tr)} \approx 2.3\alpha_s^2T \ln \frac{1}{\alpha_s}; \quad \Gamma_G^{(tr)} = 3\Gamma_Q^{(tr)}. \quad (17)$$

The transport rate $\Gamma^{(tr)}$ is closely related to the shear viscosity of the quark gluon plasma. Both approaches yield quite similar numbers for values of the strong coupling constant in the range $\alpha_s = 0.2 - 0.5$. The characteristic equilibration time, as defined as the inverse of the rates $\Gamma_i$ or $\Gamma_i^{(tr)}$ is:

$$\tau_G \approx 1 \text{ fm}/c, \quad \tau_Q \approx 3 \text{ fm}/c, \quad (18)$$

i.e. gluons thermalize about three times faster than quarks [3]. Initially, therefore, probably a rather pure glue plasma is formed in heavy ion collisions, which then gradually evolves into a chemically equilibrated quark-gluon plasma. This picture is supported by microscopic models of relativistic nuclear collisions, but rather independent of

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1(Note the additional factor two compared to the standard definition, $\Gamma = 2\gamma$).
the details entering into these models [7,8]. It is worthwhile pointing out one consequence: during its hottest phase the QCD plasma is thought to be mainly composed of gluons, which are not accessible to electromagnetic probes, such as lepton pairs and photons. The gluon plasma can be best probed by strongly interacting signals, such as charmed quarks or jets.

2. Dynamical chaos

In the early 1980’s it was discovered by Matinyan, Savvidy and others [32] that classical Yang-Mills fields represent a chaotic dynamical system in the extreme infrared limit. The full space-time dynamics can be investigated by evolving the gauge fields in real time on a lattice. From a systematic study of the evolution of randomly chosen field configurations Trayanov and Müller [33] and Gong [34] concluded that nonabelian gauge fields are characterized by a universal maximal Lyapunov exponent $\lambda_0$, which for the SU(3) gauge theory obeys the equation $\lambda_0 \approx \frac{1}{16} g^2 \langle E_p \rangle$, where $\langle E_p \rangle$ denotes the average energy per lattice plaquette. For a thermalized system $\langle E_p \rangle = \frac{\pi^2}{3} T$ holds in SU(3), hence the Lyapunov exponent satisfies

$$\lambda_0 \approx 0.54 g^2 T.$$  \hfill (19)

This corresponds to the rate with which, in the classical limit, a small field perturbation grows. It numerically coincides with the thermal damping rate of long-wavelength plasmons [10] which describes the rate of decay of a deviation from thermal equilibrium in perturbation theory. The characteristic time $\tau_s = \lambda_0^{-1} \approx \Gamma_c(0)^{-1}$ corresponds to the thermalization of the most long-ranged modes. For the temperature range expected from a violent heavy ion collision this thermalization time is less than 0.5 fm/c.

C. Formation and Evolution of the Quark-Gluon Plasma

1. Formation of a Thermalized State

Presumably, a quark-gluon plasma state can only be created in the laboratory by collisions between two nuclei at very high energy. But how do the fully coherent parton wave functions of two nuclei in their ground states evolve into locally quasi-thermal distributions of partons as they are characteristic of the quark-gluon plasma state? There are mainly two approaches to this problem that have been extensively investigated: (a) QCD string breaking, and (b) the partonic cascade.

In the string picture developed from models of soft hadron-hadron interactions, one assumes that nuclei pass through each other at collider energies with only a small rapidity loss (on average about one unit), drawing color flux tubes, or strings, between the “wounded” nucleons. If the area density of strings is low (not much greater than 1 fm$^{-2}$) they are supposed to fragment independently by quark pair production on a proper time scale or order 1 fm/c. Most realizations of this picture are based on the Lund string model [35], e.g. Fritiof [36], Attila [37], Spacer [38], VENUS [39], QGSM [40] and RQMD [41]. When the density of strings grows further, at very high energy and for heavy nuclei, the formation of “color ropes” instead of elementary flux tubes has been postulated [42,43].

Alternatively, a continuum description based on the Schwinger model of (1+1)-dimensional QED with heuristic back-reaction—“chromohydrodynamics”—has been invoked to describe the formation of a locally equilibrated quark-gluon plasma [44,45]. One general aspect of these models is that initially part of the kinetic energy of the colliding nuclei is stored in coherent glue field configurations, which subsequently decay into quark pairs. The flux tubes carry no identifiable entropy. The entropy associated with a thermal state is produced in the course of pair creation. In particular, there is no distinction between gluon and quark thermalization.

The parton cascade approach [46,47] which has precursors in the work of Boal [48], Hwa and Kajantie [49], Blaizot and Mueller [50] and Eskola et al. [51], is founded on the parton picture and renormalization-group improved perturbative QCD. Whereas the string picture runs into conceptual difficulties at very high energy when the string density becomes too large, the parton cascade becomes invalid at lower energies, where most partonic scatterings are too soft to be described by perturbative QCD.

We distinguish three-regimes in the evolution of an ultra-relativistic heavy-ion collision. Immediately after the Lorentz contracted nuclear “pancakes” have collided, scattered partons develop an incoherent identity and evolve into a quasithermal phase space distribution by free streaming separation of the longitudinal spectrum. Rescattering of these partons finally leads to the thermalization after a time probably much less than 1 fm/c. The thermalized quark-gluon plasma then evolves according to the laws of relativistic hydrodynamics, until it has cooled to the critical temperature $T_c \approx 150 - 200$ MeV, where it begins to hadronize.
The physics governing the evolution during the three states is quite different. When the two Lorentz-contracted nuclei collide, some of their partons will scatter and then continue to evolve incoherently from the remaining partons. A parton-parton scattering can be described by perturbative QCD if the momentum transfer involved is sufficiently large. It is not entirely clear where perturbative QCD becomes invalid, but popular choices for the lower momentum cut-off are $p_T^{\min} \approx 1.7 - 2 \text{ GeV}/c$. The time it takes for the scattered parton wave functions to decohere from the initial parton cloud depends on the transverse momenta of the scattered partons. Usually, one argues that the partons must have evolved at least one-half transverse wavelength $\lambda_T \approx \pi/p_T$ away from their original position before they can be considered as independent quanta. The considerations underlying this argument are similar to those for the Landau-Pomeranchuk effect.

Complete calculations following the evolution of the parton distributions microscopically until the attainment of thermal equilibrium have been carried out by K. Geiger who found almost fully thermalized phase space distributions of gluons in Au + Au collisions at RHIC energy ($E_{cm} = 100 \text{ GeV/u}$) with $T \approx 325 \text{ MeV}$ after proper time $\tau \approx 1.8 \text{ fm}/c$ (see Figure 10). These results make it very likely that a quark-gluon plasma will be produced in experiments at RHIC.

2. Expansion and Chemical Equilibration

Once the quark-gluon plasma has reached local thermal equilibrium, its further evolution can be described in the framework of relativistic hydrodynamics. The hydrodynamic equations for an ultrarelativistic plasma with $P = \frac{1}{3} \rho$ admit a boost-invariant solution describing a longitudinally expanding fireball with constant rapidity density. When transverse expansion effects are taken into account, longitudinal boost invariance is partially destroyed, but the overall picture remains intact. This scenario has been extensively studied, and is described in several publications and reviews, to which the interested reader is referred.

Thermal equilibration does not always imply chemical equilibration at the parton level. A general result of all partonic cascade simulations is that phase space equilibration occurs much faster for gluons than for quarks. As a result, the QCD plasma in its hottest stage is predominantly a gluon plasma. In the thermalized phase, this can be described by different fugacities for gluons and quarks, $\lambda_G$ and $\lambda_Q$, which evolve with time due to chemical reactions among partons and the hydrodynamic longitudinal expansion of the quark-gluon plasma.

The chemical equilibration of heavier quarks, especially strange and charmed quarks, has been investigated by many authors, because they are thought to be excellent probes of the physical conditions in dense hadronic matter. Strange quarks are predicted to equilibrate on a time scale of about $5 \text{ fm}/c$ even charmed and bottom quarks may be abundantly produced in the plasma phase, but it is quite unlikely that a complete phase space equilibration can be reached for these heavy flavors.

3. Hadronization

An important aspect of the late evolution of the quark-gluon plasma is its hadronization. Mostly it is assumed that the plasma expands and cools until it reaches the critical temperature $T_c \approx 170 \text{ MeV}$ and then converts into a hadronic gas while maintaining thermal and chemical equilibrium. More detailed descriptions of the dynamics of hadronization have been developed in connection with the problem of strange hadron production and pion production.

A totally different approach, explored more recently, consists in following the partonic reactions at a microscopic level, until the parton density has become sufficiently low to permit the formation of individual hadrons. A great deal is known about the mechanism of final state hadron production in $e^+e^-$ and NN-scattering but, unfortunately, we do not know whether this knowledge applies to the hadronization of a quark-gluon plasma in bulk.

At present, it is hard to judge the merits of either approach on a theoretical basis. Their validity and usefulness simply depends on whether the microscopic processes during hadronization proceed approximately at thermodynamical equilibrium or not. There are indications from strange hadron yields that hadronization may occur quite suddenly. The possibility of large deviations from thermal equilibrium during the chiral phase transition has recently attracted considerable interest, since it could lead to the formation of “disoriented” chiral condensate states that would manifest themselves in unusual pion charge ratios. This phenomenon will be discussed in more detail later (section III.B.2).
III. QUARK-GLUON PLASMA SIGNATURES

Theoretical speculations about the nature of the quark-gluon plasma would remain academic if there existed no experimental tools to observe its formation and to study its properties. There are several arguments against the existence of unambiguous signatures of a quark-gluon plasma formed in heavy-ion collisions. The size of the plasma volume is expected to be small, at most a few fermis in diameter, and it does not live long, maybe between 5 and 10 fm/c. Furthermore, all signals emerging from the plasma are processed through, or receive background from, the hot hadronic gas phase that follows the hadronization of the plasma phase. Nonetheless, a remarkable wealth of ideas has been put forward in the past decade as to how the identification and investigation of the short-lived quark-gluon plasma phase could be accomplished.

It is impossible to present an exhaustive review of quark-gluon plasma signatures here, and we must concentrate on the most promising ones. Many omitted details can be found in other reviews and conference proceedings. Since signals are designed to be sensitive to certain physical properties of the quark-gluon plasma phase, they will be classified here in the following way: (1) signals sensitive to the equation of state; (2) signals of chiral symmetry restoration; (3) probes of the color response function (including deconfinement); (4) probes of the electromagnetic response function; (5) various other signatures that escape simple classification.

A. Probes of the Equation of State

The basic idea behind this class of signatures is the identification of modifications in the dependence of energy density $\epsilon$, pressure $P$, and entropy density $s$ of superdense hadronic matter on temperature $T$ and baryochemical potential $\mu_b$. One wants to search for a rapid rise in the effective number of degrees of freedom, as expressed by the ratios $\epsilon/T^4$ or $s/T^4$, over a small temperature range. These quantities would exhibit a discontinuity in the presence of a first-order phase transition, at least if we were dealing with systems of infinite extent. More realistically, we can expect a steep, step-like rise as predicted by recent lattice simulations.

Of course, one requires measurable observables that are related to the variables $T$, $s$, or $\epsilon$. It is customary to identify those with the average transverse momentum $\langle p_T \rangle$, and with the rapidity distribution of the hadron multiplicity $dN/dy$, or transverse energy $dE_t/dy$, respectively. One can then, in principle, invert the $\epsilon-T$ diagram and plot $\langle p_T \rangle$ as function of $dN/dy$ or $dE_t/dy$. If there occurs a rapid change in the effective number of degrees of freedom, one expects an S-shaped curve, whose essential characteristic feature is the saturation of $\langle p_T \rangle$ during the persistence of a mixed phase, later giving way to a second rise when the structural change from color-singlet to colored constituents has been completed. Detailed numerical studies in the context of the hydrodynamical model have shown that this characteristic feature is rather weak in realistic models. The strength of the transverse flow signal can be enhanced by studying higher moments of the momentum distribution, or heavier hadrons such as baryons. It has been shown that a pion gas is probably too dilute to accelerate baryons efficiently, hence an observed collective baryon flow must be attributed to an earlier phase of the evolution of the hadronic fireball. The transverse flow signal could also be enhanced by the existence of a sharp phase boundary or detonation wave.

Flow effects in transverse momentum spectra are notoriously difficult to detect. While model calculations and phenomenological analyses of measured transverse momentum distributions of particles at presently accessible energies (up to 200 GeV/nucleon) point to the presence of transverse flow, the phenomenon is far from established in this energy range. Most likely, a credible analysis will require the study of anisotropic flow patterns allowing for the determination of the reaction plane in noncentral collisions. The techniques that have proved very successful at lower energies may also be useful in the domain of ultrarelativistic energies.

In order to trace these effects in nuclear collisions one probably has to vary the beam energy in small steps, as will be feasible at RHIC. In nucleon-antinucleon collisions, however, one may make use of the existence of large fluctuations in the total multiplicity even for central N-N collisions. This idea has been probed by the E-735 collaboration at Fermilab, who found a continued rise of $\langle p_T \rangle$ for antiprotons and hyperons with multiplicity, exceeding 1 GeV/c, for the most violent events $(dN/dy > 20)$. An analysis in terms of a schematic, scaling flow model showed that these experimental results would be consistent with transverse flow, but this hydrodynamic model is at variance with the small source sizes deduced from two-pion correlations. It is more likely that minijets are the origin of the apparent transverse flow. The importance of the minijet component is expected to grow with increasing center-of-mass energy. Minijets are predicted to be a major mechanism of energy deposition in heavy ion reactions at collider energies, as discussed before. It is also worthwhile mentioning that an increase in the average transverse momentum of produced particles with multiplicity is observed even in $e^+e^-$ annihilations into hadrons, where it can be explained by minijet branching processes.
Models of the space-time dynamics of nuclear collisions need independent confirmation, especially concerning the correctness of their geometrical assumptions. Such a check is provided by identical particle interferometry, e.g. of $\pi\pi$, KK, or NN correlations, which yield information on the reaction geometry. By studying the two-particle correlation function $F_2(q) = 1 + C_2(q)$ in different directions of phase space, it is possible to obtain measurements of the transverse and longitudinal size, of the lifetime, and of flow patterns of the hadronic fireball at the moment when it breaks up into separate hadrons \cite{103,104}. The transverse sizes found in heavy ion collisions \cite{107,110}, as well as in N-N collisions at high multiplicity \cite{113}, are larger than the radii of the incident particles clearly exposing the fact that produced hadrons rescatter before they are finally emitted. Recent theoretical work has pointed out the importance of the finite lifetime of the fireball \cite{105} and of shadowing effects \cite{106}. Since interferometric size determinations will be possible on an event-by-event basis when Pb or Au beams become available, the correlation of global parameters like $\langle p_T \rangle$ and $dN/dy$ with the fireball geometry may be performed on individual collision events.

Vector meson decays into lepton pairs also may provide information about flow patterns and the lifetime of a mixed phase of quark-gluon plasma droplet and hadronic matter. These effects will be discussed in section III.D.

### B. Signatures of Chiral Symmetry Restoration

#### 1. Strangeness Enhancement

The most often proposed signature for a restoration of spontaneously broken chiral symmetry in dense baryon-rich hadronic matter are enhancements in strangeness and antibaryon production. The basic argument in both cases is the lowering of the threshold for production of strange hadrons and baryon-antibaryon pairs. An optimal signal is obtained by considering strange antibaryons, which combine both effects \cite{13,11}.

An enhancement of strange particle production in nuclear collisions has been observed in many experiments \cite{112–117}. However it has also been learned that such an enhancement alone does not make a reliable signature for the quark-gluon plasma. Strange particles, especially K mesons and $\Lambda$ hyperons can be copiously produced in hadronic reactions before the nuclear fireball reaches equilibrium. This mechanism seems to work very efficiently at presently accessible energies \cite{118,119}. The processes leading to enhanced K and $\Lambda$ production have been studied in detail in the framework of hadron cascade models, where it was found that most of the enhancement comes from meson-baryon reactions initiated by mesons produced in first collisions, which have a strongly nonthermal spectrum \cite{115}.

However, such calculations based on the notion of a cascade of binary, sequential hadron interactions fail to describe the full spectrum of data obtained from heavy ion collisions at 200 GeV/nucleon. The enhancement of $\Lambda$-hyperon production over a wide rapidity range \cite{117}, and especially the observed strongly enhanced yields of $\bar{\Lambda}, \Xi$, and even $\Omega$, $\bar{\Omega}$ hyperons \cite{113,116} require severe modifications of the hadronic cascades. The data can be explained, if more exotic mechanisms such as color rope formation \cite{43,120}, multiple string breaking \cite{121}, or decaying multi-quark droplets \cite{123,12} are incorporated into the cascade models. Of course, such mechanisms should be considered as phenomena associated with the onset of quark-gluon plasma formation.

A simple thermal model based on the assumption of a rapidly decaying quark-gluon plasma with strange quark fugacity $\lambda_s$ near unity ($\mu_s = 0$) and light quark fugacity $\lambda_q \approx 1.5$ ($\mu_q/T \approx 0.4$) describes the observed strange hadron multiplicities rather well \cite{123,126}. However, a thermally, but not quite chemically equilibrated hadron gas can also explain the hadron yield ratios \cite{124,122}. The equilibrated fireball model, especially for a quark-gluon plasma, provides a simple explanation for the relatively large abundance of anti-hyperons, such as $\bar{\Lambda}, \Xi$, and $\bar{\Omega}$. The value $\mu_s = 0$ is natural in an environment where quarks are deconfined and a baryon-antibaryon asymmetry does not affect strange quarks, but it is rather unnatural in a hadronic environment at nonvanishing net baryon density \cite{123}. The plasma model is favored when the entropy/baryon content of the disintegrating fireball is also considered \cite{7,128}.

However, certain inconsistencies remain. It is not completely clear how a dense, hot fireball with temperature around 200 MeV can explosively disintegrate without final state interactions that modify the hadron yields. Furthermore, the fireball assumption contradicts the conclusions about the reaction dynamics derived from more sophisticated hadronic cascade models that predict a larger entropy per baryon \cite{123}. On the other hand, one can argue that the existing data point to a sequential freeze-out scenario for strange and nonstrange hadrons \cite{130}. The hadron yields may also be sensitive to structural details of the quark-hadron phase transition \cite{131}.

It has been argued that strange particles, and especially antibaryons, could also be produced more abundantly, if their masses are modified in the hadronic phase due to medium effects \cite{132}. The mass of K-mesons can be substantially lowered at finite baryon number density (this could even lead to kaon condensation \cite{133,133}), and the effective mass of antibaryons would be substantially reduced if mean-field meson theories could be applied to the description of dense hadronic matter. It remains to be seen, though, whether models which predict an enhancement
in the $\overline{\Lambda}/\pi$ ratio by an order of magnitude $[136]$ can provide a consistent description of the many other aspects of a heavy ion collision.

Strangeness enhancement may also affect the $\phi$-meson channel $[137]$. This was indeed observed in experiments at 200 GeV/nucleon ($S + U$ and $O + U$ collisions) where a clear enhancement in the $\phi/(\rho + \omega)$ yield ratio compared with $p + U$ interactions was found $[138]$. This enhancement can be explained on the basis of quark-gluon plasma formation as well as by purely hadronic scenarios $[139,140]$.

In summary, the analysis of the implications of strange hadron abundances is complicated by the fact that strangeness carrying hadrons interact strongly. Information carried by strange hadrons about their original source may be lost in final state interactions. Although these interactions have been modeled in considerable detail $[66,67]$, many predictions are quite model dependent. The lack of experimental information about many interactions involving strange hadrons is also an impediment factor $[123]$. Furthermore, it has been recognized $[131]$ that the dynamical and structural aspects of the quark-hadron transition must be much better understood before conclusions about the existence of a quark-gluon plasma can be safely drawn from hadron yields. Despite these obstacles, however, strange particle yields provide some of the most stringent tests for dynamic models of ultrarelativistic nuclear collisions $[141]$.

2. Disoriented Chiral Condensates

A direct and virtually unambiguous signal for the restoration of chiral symmetry in nuclear collisions could come from domains of disoriented chiral condensate. These correspond to isospin singlet, coherent excitations of the pion field, and would decay into neutral and charged pions with the probability distribution

$$P (\frac{N_{\pi^0}}{N_\pi} ) \sim \frac{1}{2} \sqrt{\frac{N_{\pi}}{N_{\pi^0}}},$$

(20)

Although the average ratio $\langle N_{\pi^0}/N_\pi \rangle = \frac{1}{4}$ as required by isospin symmetry, final states with a large surplus of charged pions over neutral pions, as they were observed in Centauro events $[142]$, would occur with significant probability $[143,144]$. The most likely origin of a coherent low-energy excitation of the pion field would be a collective isosinglet excitation of the Goldstone boson field ($\sigma, \vec{\pi}$) associated with the spontaneous breaking of chiral symmetry, which can be described as a nonlinear wave in the sigma model $[145,146]$. Such a wave can be excited by the growth of local instabilities during the transition from the chirally restored high-temperature phase of QCD to the low-temperature phase, in which chiral symmetry is broken $[147,148]$. These instabilities are a direct consequence of the instability of the chirally symmetric ground state below the temperature $T_c$ of the chiral phase transition. The growth of long wavelength modes in the chiral order parameter then occurs quite naturally, if the transition proceeds out of equilibrium $[149]$.

Numerical calculations have shown that the coherence length of these collective excitations remains quite small, of order 1 fm, if the chiral field gets too far out of equilibrium $[150]$. As the coherence length is inversely proportional to the growth rate of the instabilities $[151,152]$, larger domains of coherently excited pion field may emerge if the chiral order parameter is somewhat, but not too much, away from its equilibrium $[153]$. Of course, such a scenario is a priori more likely to occur in a relativistic heavy ion collision, because the chiral transition is accompanied by a large change in the energy density. The need to absorb the latent heat reduces the speed with which the temperature falls in the vicinity of $T_c$, implying that the system cannot get too far out of equilibrium until the chiral transition is completed. Boost invariant models of the dynamics of the chiral order parameter have been studied $[154,155]$ with the conclusion that coherent domains of the order of 3 – 5 fm can be formed in the cooling process $[158,159]$. It is quite unlikely that a Bose condensed state could be formed simply by isentropic expansion of a dense pion gas, as this contradicts simple entropy considerations $[160]$. It has also been shown $[161]$ that even large deviations from isospin neutrality would not significantly destroy the signature provided by the distribution (20). The observation of pion charge ratios $N_{\pi^0}/N_\pi$ significantly different from $\frac{1}{4}$, or nonzero charge correlations $[160]$, would therefore be a very strong and direct signature of the chiral phase transition.

C. Color Response Function

The basic aim in the detection of a color deconfinement phase transition is to measure changes in the color response function

$$\Pi_{\mu\nu}(q^2) = \int d^4x d^4y \, e^{iq(x-y)} \langle j_\mu^a(x) j_\nu^b(y) \rangle,$$

(21)
where \( j_\mu^a(x) \) is the color current density. Although this correlator is not gauge invariant (except in the limit \( q \to 0 \)), its structure can be probed in two ways: (1) the screening length \( \lambda_D^a = \langle \delta^a_{\mu\nu} D^{r\mu} D^{s\nu} \rangle \) leads to dissociation of bound states of a heavy quark pair \((Q\bar{Q})\), and (2) the energy loss \( dE/dx \) of a quark jet in a dense medium is sensitive to an average of \( \Pi^{ab}_\mu(q^2) \) over a wide range of \( q \).

### 1. Quarkonium Suppression

The suppression of \( J/\psi \) production \([23]\) is based on the insight that a bound state of a \((c\bar{c})\) pair cannot exist when the color screening length \( \lambda_D \approx 1/qT \) is less than the bound state radius \( \langle r_J^2 \rangle ^{1/2} \). Thus a \((c\bar{c})\) pair formed by fusion of two gluons from the colliding nuclei cannot bind inside the quark-gluon plasma \([162]\). Lattice simulations of SU(3) gauge theory \([163,164]\) show that this condition should be satisfied slightly above the deconfinement temperature \( T/T_c > 1.2 \). The screening length appears to be even shorter when dynamical fermions are included in the lattice simulations \([165]\). In addition, the \( D \)-meson is expected to dissociate in the deconfined phase, lowering the energy threshold \( \Delta E^c \) for thermal breakup of the \( J/\psi \) into two \( D \)-mesons. The combination of these two effects leads to a rapid rise of the dissociation probability beyond \( T_c \), reaching unity at about 1.2 \( T_c \) \([170]\). Excited states of the \((c\bar{c})\) system, such as \( \psi' \) and \( \chi_{c} \), are even more rapidly dissociated and should disappear as soon as the temperature exceeds \( T_c \). For the heavier \((bb)\) system similar considerations apply, although shorter screening lengths are required than for the charmonium states \([166]\). The dissociation temperature of the \( Y \) ground state is predicted to be around 2.5 \( T_c \), that of the larger \( Y' \) state around 1.1 \( T_c \).

Owing to its finite size, the formation of a \((c\bar{c})\) bound state requires a time of the order of 1 fm/c. This formation time can be understood classically, as the time required for the \((c\bar{c})\) pair to get separated by a distance equal to the average diameter of the bound state, or quantum mechanically, as the time needed for the various eigenstates of the system to decohere \([167,169]\). The \( J/\psi \) may still survive, if it escapes from the region of high density and temperature before the \( c\bar{c} \) pair has been spatially separated by more than the size of the bound state \([23]\). This will happen either if the quark-gluon plasma cools very fast, or if the \( J/\psi \) has sufficiently high transverse momentum \( p_T \geq 3 \text{ GeV/c} \). The phenomenology of the \( J/\psi \) suppression due to quark-gluon plasma formation has been extensively modeled, and its dependence on total transverse energy and nuclear size, as well as on transverse momentum, has been studied by many authors \([171,181]\). The details of \( J/\psi \) suppression near \( T_c \) are quite complicated and could require a rather long lifetime of the quark-gluon plasma state before becoming clearly visible because there are precursor effects \([182]\).

On the other hand, the \( J/\psi \) may also be destroyed in a hadronic scenario by sufficiently energetic collisions with comoving hadrons \([183,184]\), leading to dissociation into a pair of \( D \)-mesons. This mechanism has been carefully analyzed \([185–187]\). Moreover, \( J/\psi \) production is also “suppressed” in hadron-nucleus collisions, i.e. the cross section for \( hA \to J/\psi \) grows only as \( A^\alpha \) with \( \alpha \approx 0.93 < 1 \) \([188]\), as well as in \( \mu A \) collisions \([189]\). Comprehensive analyses of the available data have led to the conclusion that several effects contribute to the observed suppression at small \( x \): nuclear shadowing of gluons, initial state scattering of partons resulting in a widened transverse momentum distribution, and final state absorption on nucleons \([181,190,192]\).

These mechanisms do not explain, however, the full extent of \( J/\psi \) suppression observed in nucleus-nucleus collisions by the NA38 experiment at CERN \([193]\). Additional absorption of \( J/\psi \) on hadronic comovers (pions and resonances) with a density around 1 fm\(^3\) and an absorption cross section \( \sigma_{ab} \approx 3 \text{ mb} \) can provide a good global description of the NA38 data \([194]\). An equally good description of the present data can be obtained by invoking the plasma dissociation mechanism \([195]\). The presence of absorption by comovers is also supported by the observed stronger suppression of the \( \psi' \) state \([193]\), but this observation could be readily explained by hadronic comovers \([194]\). Similar considerations, comparing \( J/\psi \) suppression with the observed \( \phi \)-meson enhancement \([138] \) also lead to the conclusion that hadronic comovers and a quark-gluon plasma can describe the available data equally well \([139]\). In view of this persistent ambiguity, the observed \( J/\psi \) suppression cannot at present be considered as evidence for the formation of a quark-gluon plasma in nuclear collisions. Whether this ambiguity can be relieved at higher (collider) energies where \( T, T' \) suppression can also be studied, as has been argued in \([166] \), remains to be seen. In spite of these ambiguities, heavy vector meson suppression is a clear indicator of the presence of a high-density environment in the central rapidity region formed in relativistic heavy ion collisions.

### 2. Energy Loss of a Fast Parton

Another possible way of probing the color structure of QCD matter is by the energy loss of a fast parton (quark or gluon). The mechanisms are similar to those responsible for the electromagnetic energy loss of a fast charged particle in matter, i.e. energy may be lost either by excitation of the penetrated medium or by radiation.
The energy loss of a fast quark in the quark-gluon plasma was first studied in perturbative QCD (196) and later in the framework of transport theory (197). The connection between energy loss of a quark and the color-dielectric polarizability of the medium can be established in analogy to the theory of electromagnetic energy loss (198, 200). The magnitude of the energy loss is proportional to the strong coupling constant \( \alpha_s^2 \). Different authors obtain a stopping power between 0.4 and 1 GeV/fm for a fast quark (198, 201), which is somewhat smaller than the energy loss of a fast quark in nuclear matter.

Although radiation is a very efficient energy loss mechanism for relativistic particles, it is strongly suppressed in a dense medium by the Landau-Pomeranchuk effect (22). In the case of QCD this effect has recently been analyzed comprehensively (52). The suppression of soft radiation (22) limits the radiative energy loss to about 1 GeV/fm (201). Adding the two contributions it thus appears that the stopping power of a fully established quark-gluon plasma is probably higher than that of hadronic matter.

D. Electromagnetic Response Function

Photons and lepton pairs (202–204) are in many respects the cleanest signals for the quark-gluon plasma because they probe the earliest and hottest phase of the evolution of the fireball, and are not affected by final state interactions. Their drawbacks are the rather small yields and the relatively large backgrounds from hadronic decay processes, especially electromagnetic hadron decays. Electromagnetic signals probe the structure of the electromagnetic current response function

\[
W_{\mu\nu}(q^2) = \int d^4x \, d^4y \, e^{iq(x-y)} \langle j_\mu(x)j_\nu(y) \rangle.
\]  

(22)

In the hadronic phase, \( W_{\mu\nu}(q^2) \) is dominated by the \( \rho^0 \) resonance at 770 MeV, whereas perturbative QCD predicts a broad continuous spectrum above twice the thermal quark mass \( m_q = gT/\sqrt{6} \) in the high temperature phase. At \( q^2 < 100 \) MeV collective modes are predicted to exist in both phases. The collective plasma excitation, the “plasmino” (203) has a higher effective mass than the collective \( \pi^+\pi^- \) mode (206), but both lie in a region where their signal will probably be overwhelmed by background, in particular by lepton pairs emitted from non-collective modes, and by the Dalitz pair background (207). On the other hand, lepton pair production in the quark-gluon plasma may be sensitive to preequilibrium phenomena, e.g. collective plasma oscillations of large amplitude (208). Such oscillations are known to occur in the framework of the chromo-hydrodynamic model (14, 14) where the collision energy is first stored in a coherent color field which later breaks up into \( q\bar{q} \) pairs. The ensuing collective flow effects could strongly enhance the production of lepton pairs of high invariant mass.

1. Lepton Pairs

Lepton pairs have been considered as probes of the quark-gluon plasma since the earliest days (203, 204, 210, 216). Many of these original calculations concentrated on lepton pairs emitted with invariant masses in the energy range below the \( \rho \)-meson mass. However, with an improved understanding of the collision dynamics and the hadronic backgrounds (217), it has since become clear that lepton pairs from the quark-gluon plasma can probably only be observed for invariant masses above 1 – 1.5 GeV (218, 221). Assuming thermalization of the quark-gluon plasma on a time scale of about 1 fm/c, the thermal dilepton spectrum is superseded by Drell-Yan pairs from first nucleon-nucleon collisions at invariant masses around 2 – 2.5 GeV.

The great progress in understanding the mechanisms of thermalization made recently has completely changed this outlook. It has become clear that the yield of high-mass dileptons critically depends on, and provides a measure of, the thermalization time (222). Lepton pairs from the equilibrating quark-gluon plasma may dominate over the Drell-Yan background up to invariant masses in the range 5 – 10 GeV, as predicted by the parton cascade (223), and other models of the early equilibration phase of the nuclear collision (224, 220). If lepton pairs can be measured above the Drell-Yan background up to several GeV of invariant mass, the early thermal evolution of the quark-gluon phase can be traced in a rather model independent way (227). Dileptons from charm decay are predicted to yield a substantial contribution to the total dilepton spectrum and could, because of their different kinematics, provide a useful determination of the total charm yield (228). This could serve as an indirect probe of the partonic preequilibrium phase, where the total charm yield is enhanced due to rescattering of gluonic partons (7, 13), if the direct background is sufficiently well understood (228).

Although they do not constitute direct probes of the quark-gluon plasma, lepton pairs from hadronic sources in the invariant mass range between 0.5 and 1 GeV could be valuable signals of the dense hadronic matter formed in
nuclear collisions. The suggestion \[230\] that the disappearance of the \(\rho\)-meson peak in the lepton pair mass spectrum would signal the deconfinement transition has recently been revived \[231\]. The basic idea is to utilize the fact that the quark-gluon plasma phase should exhibit the higher temperature than the hadronic phase, and therefore lepton pairs from the quark-gluon plasma should dominate at high \(p_T\) over those originating from hadronic processes. Note that this reasoning may become inconclusive when one allows for collective transverse flow. Because of its larger mass, the \(\rho\)-meson spectrum is much more sensitive to the presence of flow than the quark spectrum in the quark-gluon plasma phase. Kataja et al. \[232\] have shown that the \(\rho\)-meson peak becomes more prominent at large \(p_T\), even if a quark-gluon plasma phase exists temporarily.

Nonetheless, the lepton pairs from \(\rho\)-meson decay can be a very useful tool for probing the hadronic phase of the fireball. Heinz and Lee \[233\] have pointed out that the \(\rho\)-peak is expected to grow strongly relative to the \(\omega\) and \(\phi\) peaks in the lepton pair mass spectrum, if the fireball lives substantially longer than 2 fm/c. This occurs because of the short average lifetime of the \(\rho\) (1.3 fm/c), so that several generations of thermal \(\rho^0\) mesons would contribute to the spectrum. In the limit of a very long-lived fireball the ratio of lepton pairs from \(\rho^0\) and \(\omega\) decays would approach the ratio of their leptonic decay widths. The \(\rho/\omega\) ratio can therefore serve as a fast “clock” for the fireball lifetime.

The widths and positions of the \(\rho\), \(\omega\), and \(\phi\) peaks should also be sensitive to medium-induced changes of the hadronic mass spectrum, especially to precursor phenomena associated with chiral symmetry restoration \[234, 238\]. The general conclusion, however, is that modifications of the peak positions are probably small except in the immediate vicinity of the phase transition. Changes are predicted to occur sooner, if the hadronic phase contains an appreciable net baryon density. On the other hand, the increase in the width of the \(\phi\)-meson due to collision broadening is substantial \[239\]. This could serve as a measure of the density of the mixed phase. A change in the K-meson mass also would affect the width of the \(\phi\) meson \[240, 241\]. A double \(\phi\) peak in the lepton pair spectrum would be indicative of a long-lived mixed phase \[242\].

### 2. Direct Photons

Direct photons, the second electromagnetic probe of dense matter, compete with the formidable background from the decays of pseudoscalar mesons. Reconstruction and subtraction of these decays with sufficient accuracy is difficult, but remarkable progress in this respect has been made \[243\]. At collider energies, the background from decays is predicted to extend to large transverse momenta due to the influence of minijets \[244\], but this contribution may be strongly suppressed by minijet rescattering and thermalization.

Even if decay photons can be subtracted, there remains the competition between radiation from the quark-gluon plasma and from the hadronic phase. In contrast to the lepton-pair spectrum the hadronic radiation spectrum is not concentrated in a single narrow resonance. The dominant source of photons from the thermal hadron gas is the \(\pi\rho \rightarrow \gamma\rho\) reaction \[245\], in which the very broad \(a_1\)-meson may be an important contribution \[246\]. In the quark phase the Compton scattering process \(gg \rightarrow \gamma g\) dominates. Infrared singularities occurring in perturbation theory are cured by taking into account medium effects \[247\]. The result is that a hadron gas and a quark-gluon plasma in the vicinity of the critical temperature \(T_c\) emit photon spectra of roughly equal intensity and similar spectral shape \[248\]. However, a clear signal of photons from the quark-gluon plasma would be visible for transverse momenta \(p_T\) in the range 2 – 5 GeV/c if a very hot plasma is initially formed \[249\]. The photon spectrum in the \(p_T\) range 1 – 2 GeV/c is mostly emitted from the mixed phase \[249\]. Transverse flow effects make the separation of the contributions from the different phases more difficult \[250\], and destroy the correlation between the slope of the photon spectrum in the intermediate \(p_T\) range and the temperature of the mixed phase \[251\].

Photon pairs are of interest, because they are emitted more abundantly than lepton pairs of the same invariant mass, and because they probe components of the electromagnetic response function \[22\] with different spin \[252, 253\]. Moreover, the photon pair correlations in momentum space can—in principle—be utilized for a density interferometric analysis that retains information about the quark-gluon plasma phase \[254, 255\]. The photon pair correlations are very sensitive to transverse flow effects and could be utilized to analyze the dynamics of the hadronization transition.

### E. Metastable Quark Matter

There exist general arguments for the possible metastability of cold quark matter with a high content of strange quarks \[257, 260\], which are supported by calculations in the framework of the MIT-bag model \[261, 265\]. These models also predict that strange quark matter could remain metastable at finite temperature \[266, 267\]. Whereas there exist strong astrophysical limits on the presence of absolutely stable quark matter \[268\], the existence of metastable drops
of strange quark matter (strangelets) is not excluded experimentally. Estimate of the lifetime of such objects against weak decay yield values of the order of $10^{-7} - 10^{-6}$ s \[269\].

These predictions are of interest for relativistic heavy ion collisions, because a high density of strange quarks is produced \[270\]. There exists a mechanism for the separation of strange quarks and antiquarks due to the evaporation of K-mesons from a baryon-rich quark-gluon plasma \[271\,273\]. Since this strangeness distillation mechanism requires a nonvanishing baryon number density, it should be especially effective at presently accessible energies where baryons are scattered into the central rapidity region. So far, experimental searches have been negative \[274\,275\], but highly improved experiments are in progress. The strangeness distillation mechanism can be probed by density interferometry with neutral kaons \[279\,280\]. The observation of multiple strange hypernuclei \[281\], though interesting in itself, would provide strong evidence against the metastability of small drops of strange quark matter.

IV. SUMMARY

Experimental detection of a quark-gluon plasma in relativistic heavy ion reactions at collider energies requires a combination of signals, which probe different aspects of the high-temperature phase of QCD. It is first necessary to establish the main reaction mechanism, to obtain experimental information on the initial conditions, and to determine the lifetime of the hot, dense fireball. Electromagnetic probes, charm yield, and density interferometry are promising tools for this purpose. It would be especially interesting to demonstrate the preponderance of gluonic degrees of freedom in the initial entropy production, as predicted by parton cascade models. If the thermalization times are as short, and the initial densities as high as currently predicted, the observation of a state consisting of quasifree quarks and gluons should be possible with little ambiguity.

The existence of a rather long-lived mixed phase at the quark-hadron phase transition should be visible in flow effects and in the lepton-pair spectrum. Strange baryons are an excellent probe of dense baryon-rich matter. The observed enhancements are perhaps the best indication, so far, that hadronic reaction models are insufficient. The observation of disoriented chiral domains would clearly demonstrate the presence of the chiral phase transition. Metastable strangelets, if produced and found, would provide unassailable evidence for the existence of quark matter.

The primary goal of relativistic heavy ion physics remains the exploration of the reaction mechanism at high energies. Is there a transition from the low energy regime, where reactions can be successfully described as interacting hadronic cascades, to a high-energy regime, where quark and gluon constituents provide a much simpler description of the first few fm/c of the reaction? Do the pomeron-dominated soft strong interactions disappear at sufficiently high energy, giving way to bulk interactions that can be much more economically described in the framework of perturbative QCD? An experimental demonstration of these features would constitute, in itself, the discovery of the quark-gluon plasma. The fact that hadronic cascade models are beginning to fail and to become inconsistent at the highest currently accessible energies \[282\] provides reason for optimism that this transition may be near.

ACKNOWLEDGMENTS

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[1] For an introduction to this field, see e.g.: B. Müller, The Physics of the Quark-Gluon Plasma, Lecture Notes in Physics, Vol. 225 (Springer-Verlag, Berlin-Heidelberg 1985); L. McLerran, Rev. Mod. Phys. 58, 1021 (1986); Quark-Gluon Plasma, edited by R. C. Hwa (World Scientific, Singapore, 1991).
[2] K. Kajantie and L. McLerran, Ann. Rev. Nucl. Part. Sci. 37, 293 (1987).
[3] C. P. Singh, Phys. Rep. 236, 147 (1993).
[4] Quark-Gluon Plasma Signatures, edited by V. Bernard et al. (Editions Frontieres, Paris, 1990).
[5] J. Stachel and G. R. Young, Ann. Rev. Nucl. Part. Sci. 42, 537 (1992).
[6] E. Shuryak, Phys. Rev. Lett. 68, 3270 (1992).
[7] K. Geiger, Phys. Rev. D46, 4965 and 4986 (1992); Phys. Rep. (in print).
[8] T. S. Biró, E. van Doorn, B. Müller, M. H. Thoma, and X. N. Wang, Phys. Rev. C48, 1275 (1993).
[9] T. S. Biró, C. Gong, B. Müller, and A. Trayanov, Int. J. Mod. Phys. C5, 113 (1994).
[10] P. Carruthers, Collective Phenomena 1, 147 (1973).
[11] J. C. Collins and M. Perry, *Phys. Lett.* **34**, 1353 (1975).
[12] G. Baym and S. A. Chin, *Phys. Lett.* **62B**, 241 (1976); S. A. Chin, *Phys. Lett.* **78B**, 552 (1978).
[13] B. Friedman and L. McLerran, *Phys. Rev. D**17**, 1109 (1978); L. D. McLerran, *Phys. Rev. D**24**, 450 (1981).
[14] E. V. Shuryak, *Phys. Rep.* **61**, 71 (1980).
[15] See: *Review of Particle Properties 1994*, section 25, *Phys. Rev. D**50**, 1173 (1994).
[16] M. Creutz, *Quarks, Gluons and Lattices* (Cambridge University Press, Cambridge, 1983).
[17] F. R. Brown, et al., *Phys. Rev. Lett.* **65**, 2491 (1990).
[18] N. H. Christ, *Nucl. Phys.* **A544**, 81c (1992).
[19] J. I. Kapusta, *Nucl. Phys.* **B148**, 461 (1979); T. Toimela, *Z. Phys.* **C17**, 365 (1983).
[20] V. P. Silin, *Sov. Phys. JETP* **11** 1136 (1960); V. V. Klimov *Sov. Phys. JETP* **55**, 199 (1982); H. A. Weldon, *Phys. Rev. D**26**, 1394 (1982).
[21] A. Billoire, G. Lazarides, and Q. Shafi, *Phys. Lett.* **103B**, 450 (1981); T. A. DeGrand and D. Toussaint, *Phys. Rev. D**25**, 526 (1982); G. Lazarides and S. Sarantakos, *Phys. Rev. D**31**, 389 (1985).
[22] T. Matsui and H. Satz, *Phys. Lett. B**178B**, 416 (1986).
[23] T. S. Biró and B. Müller, *Nucl. Phys. A**561**, 477 (1993).
[24] T. Biró and H. Satz, *Phys. Lett. B**178B**, 416 (1986).
[25] R. D. Pisarski, *Phys. Rev. Lett.* **63**, 1120 (1989); E. Braaten and R. D. Pisarski, *Phys. Rev. D**42**, 2156 (1990).
[26] T. Biró, P. Lévai, and B. Müller, *Phys. Rev. D**42**, 3078 (1990).
[27] D. H. Rischke, preprint CU-TP-649, Columbia University (1994).
[28] A. Ukawa, *Nucl. Phys.* **A498**, 227c (1989).
[29] C. DeTar and J. Kogut, *Phys. Rev. Lett.* **59**, 399 (1987); C. Bernard et al., *Phys. Rev. Lett.* **68**, 2125 (1992); S. Schramm and M. C. Chu, *Phys. Rev. D**48**, 2279 (1993).
[30] C. Bernard et al. (MILC collaboration), preprint UUHEP-92-10, [hep-lat/9211031](http://arxiv.org/abs/hep-lat/9211031).
[31] D. Rischke, J. Schaffner, M. Gorenstein, A. Schäfer, H. Stöcker, and W. Greiner, *Z. Phys.* **C56**, 325 (1992); *Phys. Lett. B**278**, 19 (1992).
[32] M. H. Thoma, *Phys. Lett. B**269**, 144 (1991); G. Baym, H. Monien, C. J. Pethick, and D. G. Ravenhall, *Phys. Rev. Lett.* **64**, 1867 (1990).
[33] S. G. Matinyan, G. K. Savvidy, and N. G. Ter-Arutyunyan-Sahvidy, *Sov. Phys. JETP* **53**, 421 (1981); *JETP Lett.* **34**, 590 (1981); B. V. Chirikov and D. L. Shepelyanski, *JETP Lett.* **34**, 163 (1981); *Sov. J. Nucl. Phys.* **36**, 908 (1982); G. K. Savvidy, *Nucl. Phys.* **B246**, 302 (1984). See also: T. S. Biró, S. G. Matinyan, and B. Müller, *Chaos and Gauge Theories* (World Scientific, Singapore, in print) and references therein.
[34] B. Müller and A. Trayanov, *Phys. Rev. Lett.* **68**, 3387 (1992).
[35] C. Gong, *Phys. Lett.* **298**, 257 (1993).
[36] B. Andersson, G. Gustafsson, G. Ingelman, and T. Sjöstrand, *Phys. Rep.* **97**, 31 (1983).
[37] B. Nilsson-Almqvist and E. Stenlund, *Comp. Phys. Comm.* **43**, 387 (1987).
[38] M. Gyulassy, preprint CERN-TH-4784 (1987, unpublished).
[39] T. Csörgő, J. Zimányi, J. Bondorf, and H. Heiselberg, *Phys. Lett. B**222**, 115 (1989).
[40] K. Werner, *Z. Phys.* **C42**, 85 (1989).
[41] N. S. Amelin, K. K. Gudima, and V. D. Toneev, * Yad. Fiz.* **51**, 512 (1990).
[42] M. Sorge, H. Stöcker, and W. Greiner, *Nucl. Phys.* **A498**, 567c (1989); *Ann. Phys.* **192**, 266 (1989).
[43] T. S. Biró, H. B. Nielsen, and J. Knoll, *Nucl. Phys.* **B245**, 449 (1984).
[44] H. Sorge, M. Berenguer, H. Stöcker, and W. Greiner, *Phys. Lett. B**289**, 6 (1992).
[45] A. Bialas and W. Czyż, *Phys. Rev. D**31**, 198 (1985); *Nucl. Phys. B**267**, 242 (1986).
[46] S. Kagiyaama, A. Nakamura, and A. Minaka, *Prog. Theor. Phys. Phys. B**36**, 583 (1986).
[47] K. Kajantie and T. Matsu, *Phys. Lett. B**164**, 373 (1985); G. Gatoff, A. K. Kerman, and T. Matsui, *Phys. Rev. D**36**, 114 (1986); M. Asakawa and T. Matsui, *Phys. Rev. D**43**, 2871 (1991); G. Gatoff, preprint ORNL/CCIP/91/24, Oak Ridge, 1991.
[48] K. Geiger and B. Müller, *Nucl. Phys.* **B369**, 600 (1992).
[49] D. Boal, *Phys. Rev. C**33**, 2206 (1986).
[50] R. C. Hwa and K. Kajantie, *Phys. Rev. Lett.* 56, 696 (1986).
[51] J. P. Blaizot and A. H. Mueller, *Nucl. Phys.* **B289**, 847 (1987).
[52] K. J. Eskola, K. Kajantie, and J. Lindfors, *Phys. Lett. B**214**, 613 (1988).
[53] A. B. Migdal, *Sov. Phys. JETP* **5**, 527 (1957).
[54] A. H. Sørenson, *Z. Phys.* **C53**, 595 (1992).
[55] M. Gyulassy and X. N. Wang, *Nucl. Phys.* **B420**, 583 (1994).
[56] J. D. Bjorken, *Phys. Rev. D**27**, 140 (1983).
[57] H. von Gersdorff, L. McLerran, M. Kataja, and P. V. Ruuskanen, *Phys. Rev. D**34**, 794 (1986); M. Kataja, P. V. Ruuskanen, L. McLerran, and H. von Gersdorff, *Phys. Rev. D**34**, 3755 (1986).
[58] J. P. Blaizot and J. Y. Ollitrault, in *Quark-Gluon Plasma* (World Scientific, Singapore, 1991) p. 393.
[59] K. Geiger and J. I. Kapusta, *Phys. Rev. D**47**, 4905 (1993).
[165] F. Karsch and H. W. Wyld, Phys. Lett. 213B, 505 (1988).
[166] F. Karsch and H. Satz, Z. Phys. C51, 209 (1991).
[167] V. Černý, I. Horváth, R. Lietava, A. Nogova, and J. Pišút, Z. Phys. C46, 481 (1990).
[168] J. Hufner, B. Povh, and S. Gardner, Phys. Lett. B238, 103 (1990).
[169] J. Cleymans and R. L. Thews, Z. Phys. C45, 391 (1990); R. L. Thews, Nucl. Phys. A525, 685c (1991).
[170] D. Blaschke, Nucl. Phys. A525, 269c (1991).
[171] J. P. Blaizot and J. Y. Ollitrault, Phys. Lett. 199B, 499 (1987).
[172] F. Karsch and R. Petronzio, Phys. Lett. B212B, 255 (1988).
[173] P. V. Ruuskanen and H. Satz, Z. Phys. C37, 623 (1988).
[174] M. C. Chu and T. Matsui, Phys. Rev. D37, 1851 (1988).
[175] H. Satz, Nucl. Phys. A488, 511 (1988).
[176] T. Matsui, Z. Phys. C38, 245 (1988).
[177] J. Ftačnik, P. Lichard, N. Pišutova, and J. Pišút, Z. Phys. C42, 139 (1989).
[178] G. Röpke, D. Blaschke, and H. Schultz, Phys. Rev. D38, 3589 (1988).
[179] S. Raha and B. Sinha, Phys. Lett. B318, 413 (1989).
[180] R. Lietava, Z. Phys. C50, 107 (1991).
[181] M. Gaździcki and S. Mrówczyński, Z. Phys. C49, 546 (1991).
[182] S. Hioki, T. Kanki and O. Miyamura, Prog. Theor. Phys. 84, 317 (1990); 85, 603 (1991).
[183] S. Gavin, M. Gyulassy, and A. Jackson, Phys. Lett. 207B, 257 (1988).
[184] R. Vogt, M. Prakash, P. Koch, and T. H. Hansson, Phys. Lett. 208B, 263 (1988).
[185] S. Gavin and M. Gyulassy, Phys. Lett. B214, 241 (1988).
[186] S. Gavin, R. Vogt, Nucl. Phys. B345, 104 (1990).
[187] R. Vogt, S. J. Brodsky, and R. Hoyer, Nucl. Phys. B360, 67 (1991).
[188] M. Alde et al., Phys. Rev. Lett. 66, 133 (1991).
[189] P. Amandruz et al., preprint CERN-PPE-91-198 (1991).
[190] S. Gupta and H. Satz, Phys. Lett. B283, 439 (1992).
[191] R. Vogt, S. J. Brodsky, and P. Hoyer, Nucl. Phys. B383, 643 (1992).
[192] M. A. Doncheski, M. B. Gay Ducati, and F. Halzen, preprint MAD-PH-742, Univ. Wisconsin-Madison (1993), hep-ph/9302262.
[193] C. Baglin et al. [NA38 collaboration], Phys. Lett. B220, 471 (1989); B251, 465 and 471 (1990); B262, 362 (1991); B268, 453 (1991); B270, 105 (1991); M. C. Abreu et al. [NA38 collaboration], Nucl. Phys. A544, 209c (1992).
[194] S. Gavin, H. Satz, R. L. Thews, and R. Vogt, Z. Phys. C61, 351 (1994); S. Gavin, preprint BNL 49421, Brookhaven (1993).
[195] J. G. Shen and X. J. Ziu, Nucl. Phys. A55a, 708 (1993); Z. Phys. C50, 85 (1991).
[196] J. S. Bjorken, Fermilab publication 82/59, Batavia (1982), unpublished.
[197] B. Svetitsky, Phys. Rev. D37, 2484 (1988).
[198] M. H. Thoma and M. Gyulassy, Nucl. Phys. B351, 491 (1991); E. Braaten and M. H. Thoma, Phys. Rev. D44, R2625 (1991).
[199] S. Mrówczyński, Phys. Lett. B269, 383 (1991).
[200] Y. Koike and T. Matsui, Phys. Rev. D45, 3237 (1992).
[201] M. Gyulassy, M. Plümer, M. H. Thoma, and X. N. Wang, Nucl. Phys. A538, 37c (1992).
[202] E. V. Shuryak, Phys. Lett. 78B, 150 (1978).
[203] L. D. McLerran and T. Toimela, Phys. Rev. D31, 545 (1985).
[204] S. Raha and B. Sinha, Int. J. Mod. Phys. A6, 517 (1991).
[205] H. A. Weldon, Phys. Rev. Lett. 66, 283 (1991).
[206] C. Gale and J. Kapusta, Phys. Rev. D43, 3080 (1991).
[207] E. Braaten, R. D. Pisarski, and T. C. Yuan, Phys. Rev. Lett. 64, 2242 (1990).
[208] M. Asakawa and T. Matsui, Phys. Rev. D43, 2871 (1991).
[209] E. L. Feinberg, Nuovo Cim. 34A, 391 (1976).
[210] G. Domokos and J. I. Goldman, Phys. Rev. D23, 203 (1981).
[211] K. Kajantie and H. I. Miettinen, Z. Phys. C9, 341 (1981); C14, 357 (1982).
[212] S. A. Chin, Phys. Lett. 119B, 51 (1982).
[213] G. Domokos, Phys. Rev. D28, 123 (1983).
[214] L. McLerran and T. Toimela, Phys. Rev. D31, 545 (1985).
[215] J. Cleymans and J. Fingberg, Phys. Lett. 168B, 405 (1986); J. Cleymans, J. Fingberg and K. Redlich, Phys. Rev. D35, 2153 (1987).
[216] R. Hwa and K. Kajantie, Phys. D32, 1109 (1985).
[217] J. Cleymans, K. Redlich, and H. Satz, Z. Phys. C52, 517 (1991).
[218] K. Kajantie, M. Kataja, L. McLerran, and P. V. Ruuskanen, Phys. Rev. D34, 811 (1986); K. Kajantie, J. Kapusta, L. McLerran, and A. Mekjian, Phys. Rev. D34, 2746 (1986).
[275] M. Aoki et al. [E-858 collaboration], Phys. Rev. Lett. 69, 2345 (1992).
[276] K. Borer et al. [Newmass collaboration], Phys. Rev. Lett. 72, 1415 (1994).
[277] R. S. Longacre et al. [E-810 collaboration], Nucl. Phys. A566, 167c (1994).
[278] G. Appelquist et al. [NA52 collaboration], Nucl. Phys. A566, 507c (1994).
[279] C. Greiner and B. Müller, Phys. Lett. B219, 199 (1989).
[280] M. Gyulassy, Phys. Lett. B286, 211 (1992).
[281] J. Schaffner, C. B. Dover, A. Gal, C. Greiner, and H. Stöcker, Phys. Rev. Lett. 71, 1328 (1993).
[282] K. Werner, Phys. Rev. Lett. 73, 1594 (1994).